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Peter W. Murphy

Laser Program Annual Report
Lawrence Livermore National Laboratory

85

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Preface

This volume presents the unclassified activities and accomplishments of the Inertial Confinement Fusion and Advanced Laser Development elements of the Laser Program at the Lawrence Livermore National Laboratory for the calendar year 1985. The classified aspects of our work are reported in a companion volume. The LLNL Laser Program is supported by the United States Department of Energy. Some of the activities in the Advanced Laser Development area are supported by agencies of the United States Department of Defense.

Our objective in preparing these Laser Annuals is to present an in-depth reporting of our work for the benefit of the inertial confinement fusion and laser communities. While trying to maintain conciseness, we include sufficient detail to allow full understanding, to permit critical technical evaluation, and to make it possible for others to reproduce or extend our work. Thus, we present here new concepts, theoretical analyses, computational results, the design and configuration of experiments, and analytical data—the results of our year's research. We realized significant scientific and engineering accomplishments in the Program in 1985; these are summarized in Section 1.

We have organized this report into major sections that correspond to our principal technical activities. Section 1 provides an overview: highlights of the technical accomplishments for the year and a brief summary of Program resources and facilities. Section 2 comprises our work in target theory, design, and code development. Target development and fabrication and the related topics in materials science are contained in Section 3. Section 4 presents our work in experiments and diagnostics and includes developments in data acquisition and management capabilities. In Section 5 we discuss laser system (Nova) operation and maintenance; we have collected our activities related to supporting laser and optical technologies in Section 6. Our basic laser research and development is reported in Section 7. Section 8 contains the results of our studies in ICF applications where the work reported deals principally with the production of electric power with ICF. Finally, Section 9, a new section this year, is a comprehensive discussion of our work to date on solid state lasers for average power applications. The section is relatively large because we chose to include the detailed conceptual and background work that has led us to the specific programs and activities also reported.

L. W. Coleman

Acknowledgments

The task of writing, assembling, and publishing the LLNL Laser Program Annual Report is a large and complex one. Without the full cooperation and support of the Laser Program personnel and the Laboratory's Technical Information Department such a task would not be possible. This year's report is no exception and the essential contributions of all these people are gratefully acknowledged.

Within the Laser Program, the authors accomplished the work presented in each article; their research makes this book possible. The scientific section editors for each element of the Program worked with the authors in reviewing and assembling the articles for each section; Jerry Auerbach, Jack Campbell, Mike Campbell, Dick Foley, Bill Gagnon, Andy Hawryluk, John Hunt, Ed McCauley, Howard Powell, Charles Orth, and Ed Williams served as scientific editors this year. The administrative and clerical staff within the Program typed drafts and final articles for submission into the publication process. Finally, Jenny Masquelier, who served as the Program coordinator in the preparation of the report, deserves special thanks.

From the Technical Information Department, Louise Rufer and Peter Murphy deserve special credit for their patient and untiring role as Publication Editors. They were ably assisted in the editorial process by Pat Boyd, Donna Carnes, Kellie Essary, Bob Kirvel, Gloria Lawler, and Barbara Strack. Jane Winter served as the editorial assistant. Marla Fogle coordinated composition services and typeset most of the manuscript, and Susan Gray did the important proofreading job. Wilma Leon once again coordinated the art and layout activities. Duane Hoepker, Penny Keenan, George Kitrinis, Mark McDaniel, and John Zych were the contributing artists to this year's report. Ray Marazzi and Ron Strawser did the design. Printing coordinators Mel Moura and Margaret Sylvester made sure that the book was printed and bound according to our specifications. Finally, Joseph Martinez, Steve Skolnick, and George Watchmaker contributed their talents for some of the special photography that appears here.

We thank and acknowledge all these people, and to those whose names do not appear here but who also contributed their talents and dedication to the preparation of this report.

L. W. Coleman

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Nova's ten-beam target chamber, shown below, is an aluminum sphere 4.6 m in diameter; its walls are 13 cm thick to support atmospheric pressure, the weight of the crystal arrays, and the focusing-lens drive mechanisms while maintaining accurate alignment. The chamber is designed to accommodate a wide variety of experimental programs. One target is shown, right, in position at the center of the chamber.

In inertial confinement fusion, tiny but powerful thermonuclear explosions are produced by focusing laser beams on small DT-filled targets. The beams' intensity compresses the fuel to 10 000 times its liquid density, igniting it. In the photo below, Nova's ten beams converge on a spot 250 μm in diameter at the center of the target chamber to ignite a self-sustaining thermonuclear reaction. Such experiments contribute to the technology base supporting inertial confinement fusion as a safe and economical source of electrical power for the future.

Laser Program Overview

Laser Program Overview

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Laser Program Overview

Introduction

E. Storm and L. W. Coleman

The objectives of the Lawrence Livermore National Laboratory (LLNL) Inertial Confinement Fusion (ICF) Program continue to be to understand and develop the science and technology of ICF for the maximum benefit to the country. A complete realization of ICF's full potential will require the attainment of well-diagnosed, high-yield (~ 1000 MJ) laboratory fusion target performance with an affordable multimegajoule driver. The large, tailored fluxes of x rays and neutrons from such a target would have wide applications in support of the defense nuclear weapons program. In addition, this accomplishment would experimentally prove the scientific feasibility of economic ICF power production. In working toward our goal of demonstrating high-gain laboratory microfusion, we continued during 1985 our technical program focused on two complementary areas: experiments and advanced laser studies.

Two significant events for ICF occurred in 1985: the National Academy of Sciences (NAS) completed its review of the ICF Program (and issued its preliminary report) and, with Nova, the world's most powerful laser, we began an experimental program that will lead to imploding scaled versions of high-gain ICF targets. As of this writing (Spring 1986), Nova has already achieved the highest fusion conditions of any fusion machine, and the final NAS report reiterated and amplified the basic finding of the preliminary report, namely that, "... the ICF Program has made striking progress in the last five years and the outlook for success is more optimistic today than at the time of the last review."

In order to put this comment firmly in perspective, it is fitting to recall the conclusion of that earlier review—the 1979 ICF/DOE Experts Review Committee: "... We can see no insurmountable roadblocks to the practical achievement of electrical power generated by ICF; ... many problems and unknowns remain, but because of the wealth of promising approaches, it seems plausible that with a continued national effort a prototype experimental ICF reactor will be in operation by the turn of the century." Unfortunately, the level of funding required to ensure a prototype ICF reactor by the year 2000 has not been available. However, significant research and development progress occurred toward building a strong physics and technology base, one that supports confidence in the scientific feasibility of high-gain ICF. The chairman of the 1979 ICF/DOE Experts Review Committee was a member of a subsequent 1984 review of the LLNL ICF Program and in assessing the progress made over the last five years concluded, "I am impressed by the effectiveness of your approach. You are working on everything and have produced much more progress than I would have expected. In fact, you have made so much progress, one wonders whether there is any credible alternative approach."

Indeed, a perusal of the last five unclassified and classified *Laser Program Annual Reports* (including the 1985 volumes) shows the progress and

accomplishments that prompted those comments as well as the very positive NAS review. The highlights of these accomplishments are:

- The very rapid advances in understanding and demonstrating the significant benefit of short ($\leq 0.5\text{-}\mu\text{m}$) laser wavelength for target-interaction physics.
- The striking improvements in implosion results when using short-wavelength light as compared with the older results with $1.06\text{-}\mu\text{m}$ light.
- The remarkable progress in the Halite/Centurion Program.
- The advances in materials science and technology and in fabrication techniques for high-gain targets—specifically the low-Z, low-density, liquid-DT, saturable-foam concepts.
- The completion of the Nova, 100-kJ/100-TW-class solid state laser facility. This accomplishment also assures the technical reality of a 5- to 10-MJ, 500-TW solid state laser.
- The many innovative concepts for high-efficiency, laser-driven ICF reactor systems.

Since its completion in December 1984, the Nova laser has become an unparalleled experimental facility for high-energy-density plasma physics research. Several extensive tasks have been completed, contributing to Nova's flexibility and sophistication. In addition, the Nova two-beam target area, completed on schedule and within budget, effectively doubled our experimental capacity.

We have come from the "stone age" to the "computer age" in our experimental capability with Nova's target diagnostics. During 1985, over 30 advanced diagnostics were completed, installed, and debugged. These include unique time-resolved x-ray spectral and imaging diagnostics, optical diagnostics, and advanced neutron diagnostics.

Nova has been operated at its full-rated 100-kJ of $1\text{-}\mu\text{m}$ light; however, early in Nova's activation, we discovered that the amplifier disks contained minute (typically 10- to $50\text{-}\mu\text{m}$) platinum particles, that came from the platinum crucibles used in the production of the glass disks. At low fluences, these platinum inclusions were tolerable. However, when exposed to moderate fluence levels (~ 3 to 5 J/cm^2), these particles "exploded" and damaged the disks even though the actual platinum concentration was only about the equivalent of five parts per billion. Our Advanced Laser and Nova Groups worked with the glass companies and within six months successfully solved the platinum damage problem. However, to avoid more disk damage until replacement glass was available, we chose to limit Nova's routine operation to about 50% of its peak-rated power level, sufficient for our near-term experimental campaigns. Therefore, until the damage-prone glass has been replaced, routine operation in the ten-beam chamber is currently at 15 to 20 kJ of $0.35\text{-}\mu\text{m}$ light in a 1-ns pulse on target for a variety of laser-target interaction and implosion experiments.

By January 1986, Nova achievements included:

- Demonstrated drive conditions required for high gain with scaled, indirect-drive, high-gain targets.
- Propagated complex temporal pulse shapes of the type required for high compression through the laser chains at full power.
- Generated 10^{13} neutrons, or a fusion yield of nearly 0.2% of the incident 17 kJ of $0.35\text{-}\mu\text{m}$ light with a directly driven target.

• Started experiments for x-ray conversion efficiency, hot-electron production, and exploding foils at $0.35\text{ }\mu\text{m}$ with target plasma scale sizes greater than 5×10^3 laser wavelengths.

In the two-beam chamber, routine operation is with $\sim 6\text{ TW}$ of $0.53\text{-}\mu\text{m}$ light cylindrically focused for x-ray laser experiments. We are able to conduct experiments in both the ten-beam and two-beam target chambers on the same day: up to three shots per day and up to four shots per day, respectively.

With the recent development of a 50-mg/cm^3 CH foam that wicks up liquid DT, we now have the capability to produce the very thick ($\approx 100\text{-}\mu\text{m}$) liquid-DT shells required for efficient high-gain targets.

The activation of the Nova facility, the initial experimental results, and the target-fabrication developments are further summarized in the next article and discussed in detail in Secs. 5, 4, and 3, respectively, of this report.

Solving the Nova platinum problem as well as demonstrating $>70\%$ conversion of $1.06\text{-}\mu\text{m}$ light to $0.35\text{ }\mu\text{m}$ at the full Nova aperture once again demonstrates that the Nd:glass laser continues to be the most effective and versatile tool for ICF and laboratory weapons physics because of its scalability in energy, its ability to provide flexible, controlled pulse shaping, and its adaptability to a variety of irradiation and focusing geometries. The only remaining issue is cost.

In advanced laser research during 1985, a novel amplifier architecture for a solid state laser, and continued progress in key technology areas, give us high confidence that our goal of $\$25$ to $\$50/\text{J}$ for a multimegajoule solid state laser is indeed plausible.

Also, research has begun on a master oscillator for a free-electron laser, research that will benefit the Department of Defense (DOD). The possibility of very efficient ($>40\%$ projected) solid state laser technology, small-scale experiments with high-average-power solid state lasers, and the completed engineering study for the high-efficiency Cascade ICF reactor system, strongly suggest that, contrary to previous perception, a solid state laser driver for an ICF power plant may well be possible.

In summary, the science and technology of ICF have never looked more promising. In our opinion, the only thing missing for a high-gain laboratory facility for weapons research and military applications in 5 to 10 years and for a prototype ICF reactor system within the next 20 years is the national belief that they are needed and a firm commitment to attaining them.

Progress and Accomplishments

Experiments

Introduction. During 1985, we concentrated on preparing the Nova laser system for initial target campaigns and on understanding the performance parameters of both the laser and the frequency-conversion systems. We demonstrated that

a single Nova amplifier beam can produce 15 kJ of $1.05\text{-}\mu\text{m}$ light in a 2.5-ns pulse and that we can obtain 70 to 75% conversion of the $1.05\text{-}\mu\text{m}$ light to 0.53 or $0.35\text{ }\mu\text{m}$. We routinely conducted three highly diagnosed target experiments in a two-shift day, experiments that were performed with $0.35\text{-}\mu\text{m}$ light in a 1-ns pulse at energies of up to 20 kJ on target. We also identified system modifications that will enable us to increase Nova's routine output energy to at least its full design capability of 80 to

Fig. 1-1. The target is readied for an experiment. Note the diagnostics instruments.

Fig. 1-2. A laser target being irradiated with $0.35\text{-}\mu\text{m}$ light in the ten-beam target chamber.

120 kJ with $1.05\text{-}\mu\text{m}$ light, 50 to 80 kJ with $0.53\text{-}\mu\text{m}$ light, or 40 to 70 kJ with $0.35\text{-}\mu\text{m}$ light, each in 3- to 5-ns pulses. Figures 1-1 and 1-2 are photos taken inside the target chamber: Fig. 1-1 during a target preparation and Fig. 1-2 during a target implosion.

We expanded Nova's capability by adding a two-beam target chamber: Nova's 6th and 7th beams can now be routed to irradiate either the two- or ten-beam chambers. On July 31, 1985, we delivered

approximately 4 kJ of $0.53\text{-}\mu\text{m}$ light in a 400-ps pulse in a line focus to the center of the newly installed two-beam target chamber. In this two-beam chamber we are continuing the successful x-ray laser experiments done on Novette in 1983 and 1984 to expand the power regimes and to explore shorter wavelengths. Figure 1-3 shows the two-beam target chamber.

Because of the presence of platinum inclusions in the laser amplifier disks, we decided to limit, for the present, Nova's routine operation to modest energy and power levels ($\sim 50\%$ of its design limit). To fully understand and quantify the factors that led us to this decision, we utilized experimental stations at several positions along the laser to conduct damage and propagation experiments with full- and partial-aperture laser beams. We also used this capability to test the newly developed platinum-free, neodymium-doped phosphate laser glass that is required in the amplifiers for damage-free Nova operation at its full output capability.

We also conducted experiments to determine the reason for the lower-than-expected second- and third-harmonic frequency conversion efficiency and to evaluate stimulated Raman scattering effects created by high laser-beam fluences in long air paths. Detailed

Fig. 1-3. The Nova two-beam target chamber.

experimental results are presented in Sec. 4 of this report.

We concentrated heavily on the design, fabrication, installation, and physics activation of numerous diagnostics on both the ten-beam and two-beam target chambers; the extent and precision of these diagnostics make Nova the most sophisticated and versatile experimental facility in the world for the study of high-energy-density plasma physics.

Nova Experiments. We conducted experiments that initially served to activate the laser facility and diagnostic systems and also addressed target issues, particularly those pertaining to the physics of high-Z plasmas irradiated with 0.53- and 0.35- μm -wavelength light, at higher energy and with larger irradiation spots than had been possible with previous facilities. The multiterawatt capability of Nova at both wavelengths has allowed us to begin studying such issues as absorption, x-ray conversion efficiency, suprathermal electron production, and parametric instabilities in large plasmas (scalelengths of 10^3 to 10^4 laser wavelengths) at intensities relevant to ICF. These experiments have

expanded the results obtained in 1983 and 1984 with Novette, which we reported in our previous Laser Annuals.

We continued exploding-foil experiments to study parametric instabilities such as stimulated Raman scattering in long-scalelength, fully underdense plasmas. The goal of these experiments is to understand collisional damping of stimulated scattering with 0.35- μm irradiation.

Early experiments on the two-beam target chamber have concentrated on soft x-ray laser studies and are a continuation of the very successful Novette experiments performed in 1984. Concurrently, we are extending our modeling of laboratory soft x-ray laser schemes and are designing Nova experiments in which gain saturation for the neon-like selenium scheme, which requires several centimeters of gain length, can be achieved. We are choosing designs that minimize the refraction effects that reduce the gain achievable with these targets.

In some experiments, the length of the selenium exploding-foil targets has been extended to 2.5 cm; with these targets we continue to observe exponentiation of the

intensities of the 206- and 209-Å lines from neon-like selenium. In addition, due to the longer amplification distance in the plasma and with improved detector sensitivity, we have also observed gain at three additional lines from the same 3p manifold at 180, 220, and 260 Å. The line at 180 Å is of particular interest because its upper state at $J = 0$ had been predicted to have the highest gain of any of the 3p-3s transitions in neon-like selenium, and had not been clearly observed in the earlier experiments. With yttrium and molybdenum, we observed a total of eight x-ray lasing lines at wavelengths as short as 105 Å.

We are examining collisional excitation of excited states in nickel-like ions as a scheme to extend lasing to shorter, more penetrating wavelengths. Computer simulations indicate that there are numerous transitions in nickel-like gadolinium that have significant gain when irradiated with moderate intensity ($I \sim 0.5 - 1 \times 10^{14}$ W/cm²) 0.53-μm light and that appear in the 80- to 160-Å range. We are also analyzing photoresonant pumping schemes and recombination schemes.

We conducted joint experiments with KMS Fusion, Inc. using their four-frame holographic interferometry to measure the plasma electron-density profiles with several exploding-foil target designs and irradiation conditions. These experiments serve as important checks on the hydrodynamic modeling of soft x-ray laser design, and they permit quantitative analyses of the effects of refraction on radiation propagation through the plasma.

Experiments to characterize sub-100-ps laser-driven x-ray sources for backlighting diagnostics have also been performed in a joint effort with KMS Fusion, Inc. A wide variety of future Nova hydrodynamic and spectroscopic experiments require intense sub-100-ps x-ray diagnostic sources at a number of photon energies, the short source duration being required to minimize motion blurring of the experimental data. So far, these experiments have demonstrated bright, sub-50-ps x-ray bursts with photon energies as high as 4.7 keV.

We obtained results from a set of laser-driven impulse experiments conducted on the two-beam Nova target chamber. In these experiments, we obtained 100- to 150-kbar impulses in aluminum targets

irradiated with intensities in the range of 10^{10} to 10^{12} W/cm² in a 1-ns pulse of 0.53-μm light.

Diagnostics. We have made advances in developing new x-ray and neutron diagnostics. Highlights in this area include the fabrication and testing of multilayer mirrors for diagnostic applications and soft x-ray laser cavities, the development of the small, bright, laser-driven x-ray sources just mentioned, and the operation of a direct-recording subkilovolt x-ray charge-coupled-device camera; the latter two technologies will be extremely useful in characterizing new x-ray optical components. High-performance capsules on Nova will require precise and complex neutron diagnostics to characterize the imploded fuel state. Therefore, we are developing diagnostics that can measure the spectra, time history, and spatial extent of penetrating fusion products.

We are also continuing to develop our time-resolved measurements capability. Of particular note are the advances made in x-ray framing cameras and on an optical fiducial system for x-ray and optical streak cameras.

Data Management. The experimental programs on Nova demand a system for rapid processing and archiving of data, a service provided by the Data Management and Analysis Group (DMAAG). A highlight of their work is the activation and on-line operation of a comprehensive computer-based electronic shotbook, which enables scientists to have rapid user-friendly access to experimental results from Nova in a standardized presentation format.

Target Design. Our activity in laser-target theory, design, and analysis during 1985 was directed primarily toward preparing for the series of experimental campaigns planned for 1986 on Nova. These campaigns will address the fundamentals of hohlraum behavior using 0.35-μm light, as well as issues of symmetry and hydrodynamic stability and implosion. Much of our target design effort is reported in the classified companion to this Laser Annual.

Numerous improvements have been implemented into the LASNEX code system. These improvements include a finite-element-based diffusion model that is accurate for very badly skewed meshes. We have also implemented a detailed

configuration atomic model into LASNEX. This is our first step toward making LASNEX capable of calculating line transport for soft x-ray laser applications and for spectroscopic diagnostic modeling.

Target Fabrication. In target fabrication and materials development, our work this past year generally fell into two broad categories: (1) ICF and weapons physics targets and (2) soft x-ray laser targets. We are also developing an advanced target assembly and pre-alignment system, which will aid in the final assembly of the target components and expedite positioning and aligning a target in the Nova target chamber.

An exciting development in the past year has been the emergence of cryogenic wetted-foam targets (see Fig. 1-4). Low-Z (CH) foams with extremely small pores are now available from which direct-drive targets that are very close to the original concept for ICF can be fabricated. The foams act as a matrix to hold a thick, liquid layer of DT just as a sponge holds water. This makes it possible to fabricate very thick (greater than several hundred microns) liquid-DT shells that are >80% DT and <20% CH foam. The foams are of low enough density and Z that it is still possible to achieve a propagating burn front ignited from a DT hot spot.

The requirements on the foam material are quite demanding. Because foam materials that met all the requirements were not available commercially, we began a development effort to produce the desired material.

We are also continuing to study methods for forming homogeneous liquid-fuel layers in capsules. Currently, the most promising method makes use of a thermal gradient across the capsule. We also began an effort to produce spin-polarized DT mixtures for possible fusion applications.

Advanced Laser Studies

Introduction. The objectives of the Advanced Laser Development portion of our program are to develop advanced lasers for ICF, to develop medium-average-power solid state laser technology supported by the Defense Advanced Research Project Agency (DARPA), and to contribute to the emerging free-electron-laser program. We also use our special expertise in lasers and materials for basic research for the

Office of Basic Energy Sciences and for important special developments such as UV optical coatings for the Air Force Weapons Laboratory.

The additional laser development activities that we have undertaken this year are highly synergistic with ICF and are the outgrowth of the intense research and development efforts in solid state laser technology conducted since 1971 when the Laser Program was formalized at LLNL. These efforts, together with contributions from programs sponsored by DOD, have moved solid state laser technology into new operating regimes. We project that solid state lasers will have pulse repetition frequencies (PRF) approaching 1 kHz, and efficiencies exceeding 10%, with beam quality approaching 1/10-wave rms, which result in compact, high-quality laser systems with high flexibility of operation.

A major laser activity is directed toward meeting objectives for ICF drivers. The technology developments described in this Laser Annual are based on the continuing need to support Nova's operation as the world's most powerful and flexible target facility, to develop the basis for a high-energy, low-cost driver for high-gain ICF demonstrations, and ultimately for an efficient, economic laser driver for ICF reactor applications.

Nova Activities. Nova support activities have been aimed at raising the output power limitations of the system by raising the damage threshold of the laser glass and at accurately quantifying the details of rotational Raman scattering in the air path at the highest power levels. We are also developing a fiber optic system for diagnostic applications at 0.35 μm .

A very important materials project is the formulation and production of

Fig. 1-4. A directly-driven low-density, low-Z, DT-wetted foam capsule.

Progress and Accomplishments

platinum-free laser glass. A benefit of this development is that the production of glass with a high damage threshold appears to be compatible with low-cost, continuous-melting technology; this strongly supports the performance projections of our megajoule laser concepts.

Lasers for High-Gain ICF. Our primary, near-term laser design objective for high-gain ICF is to provide technical and cost bases for the construction of a 5- to 10-MJ, flexible, single-shot facility for high-gain ICF targets. This year, our studies of megajoule-sized solid state laser architectures have been based on target-plane holographic injection and phase conjugation. Measurements on phase conjugation have begun, and a variety of systems using phase conjugation have been conceived. The conclusions of several system studies indicate that a 5- to 10-MJ single-shot laser system could cost as little as two to three times the cost of Nova.

Notable progress has been made in the development of laser materials—crystals, glass, and coatings; in components, especially amplifiers and optical switches; and in a very thorough understanding of nonlinear crystals and crystalline laser media for advanced applications. Figure 1-5 shows a YAG crystal. A significant fraction of our advanced laser development resources has in fact been directed toward optical materials studies. Good progress has been made in acquiring a detailed understanding of damage in laser glasses,

crystals, and optical coatings; on rapid growth of nonlinear crystals; on developing a basic understanding of nonlinear optical crystals; and on developing a methodology for the development of laser ions and hosts capable of efficient ($>10\%$), moderate-PRF (1 to 10 Hz) laser performance for fusion power plant applications. Progress has continued this year with the completion of the analysis of a high-efficiency, glass-amplifier module (the single-segment amplifier) with which we demonstrated a storage efficiency of over 7%.

High-Average-Power Lasers. Our two reimbursable programs (one from DARPA begun in late FY 84 and the other from the Strategic Defense Initiative Organization begun in FY 86) for a free-electron-laser master oscillator, have as their objectives the demonstration of high-average-power solid state laser operating conditions for DOD applications. In a separate section of this Laser Annual (Sec. 9), we discuss these activities in detail.

ICF Applications Studies

The Cascade ICF reactor is proving to be the first design that can be economically competitive with fission and coal plants while maintaining the desirable safety and environmental impact features characteristic of ICF. The Cascade design achieves a very high energy-conversion efficiency (about 55%) at a reasonable cost. Its reaction chamber consists of a rotating double cone at the center of which the fusion reactions occur. Ceramic granules, poured into the ends, are held against the wall and flow through the chamber, protecting the wall, breeding tritium, and acting as the heat-transfer fluid. They give up their extremely high temperature to high-pressure helium that then goes through a Brayton cycle. The Cascade design has been under study for almost three years; design details have evolved, but the basic concept has remained the same. It has withstood a great deal of scrutiny as we and others have looked, unsuccessfully, for fatal flaws in the concept. GA Technologies has just completed a major study of the reaction chamber and the balance of

Fig. 1-5. A large YAG crystal.

plant; they concluded that the design is feasible and, if built, the cost of the electricity it would produce would be comparable to or below that from future fission and coal-powered plants.

Another reactor concept, Pulse*Star, was designed with the idea of improving on an earlier concept, the high-yield lithium-injected fusion energy (HYLIFE) reactor. With Pulse*Star studies we attempted to reduce the costs of a liquid-metal-based fusion reactor by increasing the pulse rate, by eliminating the fire hazard of liquid lithium, by eliminating the intermediate heat-exchange loop, and by adopting a pool-type design to reduce the number of pipes and connections required. During 1985, we addressed issues pertaining to the support of the lithium-lead pool, the design of the inner screen that defines the lithium-lead droplet region, and the containment of tritium. There appear to be design solutions to these issues. It is likely that the Pulse*Star design would produce electricity more cheaply than a plant based on HYLIFE would. However, since Pulse*Star is still based on the steam cycle, it would probably be less efficient than Cascade. Therefore, even though the Pulse*Star design study is still not complete, we have decided to stop working on it.

In our development scenario for ICF, we identify a number of possible applications

that include weapons physics, energy and materials production, and space propulsion. In such a scenario, it will be important to design the development facilities (especially a laboratory microfusion facility) to have multiple test chambers and to have a very flexible driver.

Author: L. W. Coleman

Program Resources

Financial resources for the ICF Program in FY 85 totaled \$68.3 million in new appropriated funding and \$82.1 million in expenditures. Expenditures included \$67.2 million in operating funds and capital equipment, and \$9.9 million in previously appropriated Nova funds. Operating costs, nearly a 12% increase over those of FY 84, represented the increased expense of operating the Nova laser.

Work force levels averaged 301 direct employees for the operating program and an additional 54 for Nova construction. The completion of Nova, however, necessitated a substantial decrease in total work force over the year, which began with 420 direct employees and ended with 252. These were primarily planned reductions associated with Nova project completion,

Fig. 1-6. Ten-year history of ICF Program resources.

Program Facilities

but funding limitations in the operating program required reductions beyond that plan.

Capital equipment funding of \$5.0 million was available to support the operating program. This level of funding, representing a 6% increase over FY 84, supported our need for increased capabilities in the diagnostic and target fabrication areas. These are directly related to the upcoming experiment campaigns on Nova. Resources for the last ten years are shown in Fig. 1-6.

Author: L. P. Altbaum

Program Facilities

The overall Laser Program has grown steadily and dramatically, both in size and in mission, during the past six years. The overall Laser Program (Inertial Confinement Fusion, Advanced Lasers, and Laser Isotope Separation) comprises a total current work force of more than 1400 people. With the successful execution of a coordinated and innovative long-range development plan, the Program today has approximately 850 000 square feet of office and laboratory floor space located in 29 permanent buildings and nearly 150 000 square feet of space in 45 temporary facilities.

Following the completion of major, long-range construction projects late in 1984 and early in 1985, the efforts of the Laser Program Facilities Group during the past year have concentrated on building activation, on utilities installation, and on laboratory conversion and renovation work. Among these 1985 projects have been the coordination of the Nova Laser System dedication, the installation of utility piping, power, and air control systems in the Nova building, and the remodeling and renovation of some laboratories.

Major work completed for the ICF and Advanced Laser Programs in 1985 has included the following:

1. The design and installation of laser laboratories and support facilities for the UC-Davis extension program.
2. The coordination of efforts for the dedication of the Nova Laser.
3. Development of a Conceptual Design Report for the Laser Optics Research Facility.

4. Modifications for the Nova two-beam target chamber area (former Shiva laser bay) and installation of the two-beam target chamber.

5. Installation of the humidity control system for the Nova laser, which included installation of a dry air system in the ten-beam target chamber area, and humidification of the Nova laser bay and switchyard areas.

6. The design of the Acid Waste Retention System.

7. The design, planning, and coordination of remodeling efforts and the completion of interior modifications for several buildings.

8. Completion of basement modifications for the main office buildings and the neutron laboratories.

9. Modifications to the high-average-power laboratory, which included the design and installation of utilities to the building, the installation of a second power supply, and the design and installation of optical water windows.

10. The design and installation of optical water windows in the Nova switchyard.

11. Continuing additions to several laser laboratories.

12. Continuing site improvements.

With the completion of large construction projects, the operational emphasis for the facilities group is now on maintenance, remodeling, and renovation work. The high volume and broad variety of current construction projects have presented a major challenge and a unique opportunity for exercising the multidimensional talents of the facilities group. Smaller construction projects, routine and emergency repairs, and the improvement of permanent buildings and temporary structures have all been developed by the facilities group, which provides the Laser Program with services that include site planning, project design, planning and estimating, material procurement, and construction supervision. In 1985, those activities contributed to the initiation, development, and eventual completion of 895 separate job orders at a value of \$6.1 million.

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The diagram, right, illustrates how a small mass perturbation in a spherical shell amplifies on implosion. The large illustration below depicts a LASNEX simulation of a shell where the initial 5% mass perturbation was magnified in the later stages of implosion thus demonstrating that good surface finish is required for adequate implosion symmetry for successful ICF.

Through the analysis of experiments and by theoretical research, target designers develop an understanding of the key physics issues pertaining to laser-matter interactions and to thermonuclear microexplosions. In turn, improved theoretical models are incorporated into the computer codes that are used to numerically simulate experiments and future designs.

The small photo below, right, shows a holographic interferogram obtained in experiments performed at KMS Fusion, Inc. in support of our laboratory x-ray laser program. On analysis, researchers at KMS provided snapshots of the density profile of our selenium foil x-ray laser targets. These compare closely with those obtained in LASNEX simulations. The diagram below, left, shows a corresponding LASNEX mesh plot.

Target Design

Target Design

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Section Editor: E. A. Williams

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Target Design

Introduction

J. D. Lindl

The Target Design Program combines the efforts of the Plasma, Code Development, and Design Groups in three basic areas:

- Developing theories of beam-plasma interaction, implosions, and thermonuclear microexplosions.
- Building plasma and implosion-burn computer codes.
- Using the theoretical and computational tools to design targets and simulate experiments.

Our effort in laser-target theory and analysis during 1985 was directed primarily toward preparing for a series of experimental campaigns planned for FY 86 on Nova. The campaigns will address the fundamentals of hohlraum behavior using 0.35- μm light, issues of symmetry and hydrodynamic stability, and a preliminary series of implosions.

One exciting development in the past year has been the emergence of cryogenic wetted-foam targets. Foams with extremely small pore size and low Z are now available, making it possible to fabricate direct-drive targets that are close to our original concept for inertial confinement fusion (ICF). The foams act as a matrix to hold a thick liquid layer of deuterium-tritium (DT). The foams are of such low density and Z that it remains possible to achieve a propagating burn front ignited from a DT hot spot.

Electron transport in the steep-gradient, high-flux regime important for ICF continues to be an important topic. We have developed a fluid description of electron heat flow, which self-consistently accounts for departures of the isotropic part of the electron-velocity distribution from Maxwell-Boltzmann. The resulting heat flow is both nonlocal and naturally flux-limited and agrees well with detailed Fokker-Planck calculations.

The 2.5-D electromagnetic direct-implicit particle code, developed to enlarge the range of problems accessible to particle simulation, is now operational and running a variety of problems. The code is capable of time steps greater by one order of magnitude than the corresponding explicit code. Hence, the direct-implicit code can simulate much larger or longer problems for which the electron-plasma time scale is unimportant.

Our efforts at modeling laboratory x-ray laser schemes have expanded substantially during 1985 with support from both the Physics Department and R Program. Following last year's successful experiments on Novette, we are designing experiments to achieve gain saturation on Nova for the neon-like selenium scheme. This goal requires propagation lengths of several centimeters, for which we have investigated several geometries. The designs are chosen to minimize refraction effects that could reduce the achievable gain from these targets. In addition, we are exploring a variety of approaches that could result in lasing at higher energies. The approaches include nickel-like collisional excitation schemes, photoresonant pumping schemes, and recombination schemes.

A wide variety of improvements have been implemented on the LASNEX code system. The improvements include a finite-element-based diffusion model that is accurate for highly skewed meshes. The model has been successfully tested in static geometry but must be further improved to maintain numerical stability in a dynamic mesh.

We have implemented a detailed-configuration atomic model into LASNEX. This is our first step toward making LASNEX capable of calculating line transport for x-ray laser applications and for spectroscopic diagnostic modeling. The S_n discrete-ray approach to radiation transport has been further improved by the addition of relativistic Compton scattering and a new numerical scheme that behaves correctly in the thermal wave limit of radiation transport.

The following articles report our progress on laser targets, ion-beam targets, plasma physics, and code development. Our work on the design and interpretation of experiments is reported in Section 4.

ICF and X-Ray Laser Target Design

Introduction

In 1985 the Laser Target Design Group continued to make significant progress in many areas of research. Specific areas of achievement include capsule design, planning of experiments on symmetry and stability, analysis of existing data, and invention and design of new schemes and targets for laboratory x-ray lasers.

We continued to study a long-range objective of achieving a hydrodynamically equivalent target (HET) on the Nova system. The HET is a target with hydrodynamic behavior much like that of a high-gain target but scaled down in size to match the energy available from Nova. The HET is too small to achieve enough hot-spot *pr* to ignite and propagate the burn into the cold, near-Fermi-degenerate fuel around it.

This year we report a major breakthrough in target design and fabrication. In the article "Direct-Drive Cryogenic Capsule Design Employing DT-Wetted Foam," we describe a method of achieving a capsule that is similar to the original laser fusion pellet designs—namely thick, frozen shells of deuterium-tritium (DT). By using the new technology of ultralight foams that are wetted with DT, we can closely approximate a thick, pure-DT shell, thus making the concept a reality.

We continued our investigations into the Rayleigh-Taylor (R-T) instability on both the experimental and theoretical fronts. In "Signal-to-Noise Ratio in Backlit Nova Experiments," we investigate some practical constraints on an experiment designed to measure instability growth rates. In "Turbulent Mix Generated by Rayleigh-Taylor and Richtmyer-Meshkov Instabilities," we use some of the latest theories of nonlinear turbulence in the exploration of problems relevant to ICF.

The target design effort on laboratory x-ray lasers, with support from both the Physics Department and R Program, has doubled in the wake of last year's success. As a result, we are pursuing a wide range of lasing scenarios. In "Progress in the Analysis of Selenium X-Ray Laser Targets," we describe the important effect of properly modeling dielectronic recombination on the analysis of our previous data. In addition we have invented several new target geometries to ensure beam propagation down lengths greater than 1 cm. Our more complete understanding of the now-conventional x-ray laser target is described in "Hydrodynamics of Exploding-Foil X-Ray Laser Targets."

Other lasing schemes with short-wavelength potential continue to occupy our interest. Improvements in our understanding of one system are described in "Investigation of the Manganese Pump Line for the Photoresonant X-Ray Laser." An entirely new scheme, analogous to our neon-like system, is discussed in

"Estimates of Gain for Nickel-Like Exploding-Foil Targets." Finally, we analyze previous experimental results of a hydrogen-like magnesium scheme in "Recombination/Rapid-Cooling X-Ray Lasers" and discuss ways to improve that general approach with respect to soft x-ray lasing.

Author: M. D. Rosen

Signal-to-Noise Ratio in Backlit Nova Experiments

Two-dimensional or streaked x-ray images of targets illuminated by a backlighter are an important part of many Nova experiments. Even when the x-ray optics and camera perform perfectly, the signal-to-noise ratio for such images is limited by shot noise caused by the finite number of photons collected for each resolution element.

The Nova Rayleigh-Taylor series will be our first set of experiments using the $22\times$ imaging system. Performance approaching the ultimate shot-noise limit is required for these experiments to be a complete success. In this article, we provide a general description of how to calculate the ultimate signal-to-noise ratio for

any imaging system involving an x-ray optic. We also give a specific estimate of the ultimate signal-to-noise ratio for the $22\times$ Wölter microscope.

The components of a backlit x-ray imaging system are shown in Fig. 2-1. One beam of the Nova laser shines on a plate made of material selected as the backlighter, producing an extended area of relatively bright illumination with the desired x-ray spectrum. The object plane of the imaging system is at the target (that is, the Rayleigh-Taylor sample, an implosion capsule, or some other target). The remaining nine Nova beams are generally used to drive the target. Optically thick regions of the target attenuate the backlighter more than do optically thin regions. Assuming that the target can be designed to be much less bright than the backlighter at the x-ray energies of interest, a shadow of the target is visible from the side opposite the backlighter. The x-ray optic (in this case, a Wölter ellipsoid-hyperboloid, grazing-incidence mirror) images the backlit target onto a photocathode or directly onto film.

If a streaked image is desired, the Nova backlighting beam remains on for the duration of the experiment. A slit of width W is placed over the photocathode,

Fig. 2-1. One Nova beam creates an extended source of backlighting radiation that illuminates a target and is imaged by an x-ray optic.

restricting the view of the target to a strip of width W/M , where M is the magnification of the x-ray optic ($M = 22$ for the $22\times$ microscope). The image of the slit is moved across the film in a direction perpendicular to the slit. The outcome is a streaked image, with position along the slit on the horizontal axis of the image and time along the vertical axis, as shown in Fig. 2-2(a).

Because such images are subsequently magnified for display to many different sizes, we may think of the dimensions on the image in terms of lengths Δx measured along the slit in the object plane (i.e., at the target) and times Δt . The ratio of the slit width to the streak speed sets a lower limit on the temporal resolution that can be achieved; however, for the purposes of this analysis, the streak speed itself is irrelevant. Our assumption is that the streak speed can be made so fast that shot noise, rather than a slow streak, limits the temporal resolution of the image. As a practical matter, the only limit on the streak speed is the ratio of the maximum distance the slit image can be moved (before it moves off the film or hits some other part of the camera) to the width of the slit. This ratio limits the total number of temporal resolution elements available on an image for the entire experiment.

If a 2-D snapshot image of the target is desired, then the Nova backlighting beam

remains on only for a brief time to freeze the motion of the target. Because there is no other shuttering, the integrated backlighter intensity during this brief flash must exceed the surface brightness of the target integrated over time. An image of this type is shown in Fig. 2-2(b). Once again, it is best to think of the dimensions in such an image as lengths Δx and Δy at the target, independent of the magnification M of the optic. In general, both resolution requirements for the experiment and performance of the optic are different along the two axes of such a snapshot. However, shot noise places a limit only on the area $\Delta x \Delta y$ of a resolution element, not on its shape.

Radiation transport codes, such as the LASNEX postprocessor code TDG, compute the specific intensity of the radiation field emerging from the object plane. This computation can be convolved with the filter response function of the imaging system to obtain the number of photons I emerging per unit time per unit surface area (of the object plane) and per unit solid angle (viewed from the object plane). The most important characteristic of any optic is the solid angle Ω that it subtends as seen from the target. Allowing for the factor Ω , the total number of photons $I\Omega$ that pass through the imaging optic per unit area of the target per unit time is obtained. Because x-ray optics operate

Fig. 2-2. (a) Long-duration backlighter flash and streak camera produce a streaked image. (b) Short backlighter flash and ordinary camera produce a 2-D image.

on the principle of grazing-incidence reflection, Ω is, unfortunately, small. For the $22 \times$ Wöhlert optic, stepped down to give a reasonable depth of focus, Ω is about 10^{-5} sr.

Our previous experience with back-lighters has been that a certain fraction of incident laser energy is converted into the desired backlighter line radiation. The fraction depends on the backlighter material, the x-ray emission lines of interest, the laser intensity, and the pulse duration. The backlighter radiation is assumed here to be hemi-isotropic. Because the angle subtended by the optic viewed from the target is small and because the backlighter is much closer to the target than to the optic, the magnification of the out-of-focus backlighter is not significantly different from the magnification of the target. Thus, the terms "per unit area in the object plane" and "per unit area of the backlighter" have essentially the same meaning.

If the laser intensity is I_L and a fraction f is converted to backlighter x-ray lines with mean energy $h\nu$, and if the transmission fraction of the target is χ , then the number of backlighter photons per unit time per unit area per unit solid angle is

$$I = \chi I_0 = \frac{\chi I_L}{2\pi h\nu} \quad (1)$$

In Eq. (1), I_L is the laser power per unit area of the backlighter as projected on the object plane, which is somewhat higher than the actual intensity on the inclined surface of the backlighter.

For the 1.5-keV aluminum K-line backlighter proposed for use in the first Rayleigh-Taylor series, the conversion fraction f is about 1% for blue light.¹ A typical transmission fraction χ for a Rayleigh-Taylor sample is 0.1, to within about one order of magnitude.

Consider first a streaked image. A single Nova beam should provide a 2-kJ, 3-ns blue backlighting pulse, which will illuminate an $800\text{-}\mu\text{m}$, 0.5-mm^2 sample disk. Hence, I_0 is about 10^{16} photons/(ns·mm²·sr) for this experiment. (Higher-energy backlighters produce fewer photons because the photon energy is higher and the conversion fraction f is lower.) $I_0\Omega = 10^{11}$ photons/(ns·mm²) could

therefore reach the image plane. If 200-ps time resolution is required for the duration of a 3-ns experiment, there are a total of 15 time resolution elements. The LLNL streak cameras have a useful streak distance of perhaps 30 mm, so that the slit width W could be no larger than 2 mm to perform the experiment with the stated resolution. This value corresponds to a strip of width $W/M = 100\text{ }\mu\text{m}$ in the object plane, which may prove to be intolerably wide and an upper limit on the usable slit width for the experiment. The maximum number of photons available for the image is thus $I_0\Omega W/M = 10^{10}$ photons/(ns·mm).

Next, consider a 2-D snapshot image. For a 200-ps flash backlighter, a Nova beam can deliver only about 300 J to the backlighter. If the illuminated area is the same $800\text{-}\mu\text{m}$ spot as before, then $E_0 = 4 \times 10^{15}$ photons/(mm²·sr) would be available for a snapshot, where E_0 is I_0 integrated over the duration of the flash. (Shorter-duration backlighter flashes produce fewer photons owing to the limitation on peak laser power that can be delivered by Nova.) Hence, $E_0\Omega = 4 \times 10^{10}$ photons/mm² could reach the image plane.

In addition to the attenuation χ of the backlighter radiation by the target, the x-ray detector in the image plane has an imperfect quantum efficiency η . For photocathodes, the relevant quantum efficiency is not the number of photoelectrons per incident photon because each photon that stops in the photocathode gives rise to a number of secondary electrons, which are emitted from the photocathode surface. Photocathode quantum efficiencies are limited by the fact that the range of the primary photoelectron is generally much less than the mean free path of the incident photon, so that the thickest useful photocathode presents a rather small optical depth to the incident x rays. For 1.5-keV photons, an optimized CsI photocathode has a quantum efficiency not exceeding $\eta = 0.1$. This is the best photocathode known.²

The quantity to be measured on the image is the optical depth difference $\Delta\tau$ between adjacent resolution elements. If N_1 photons are detected at one resolution element and N_2 at the next, then $\Delta\tau = \ln(N_1/N_2)$. The specific intensity can be measured with a relative error of $N^{-1/2}$

due to shot noise in an experiment in which N photons are detected. Hence, shot noise causes a random absolute error $\delta\Delta\tau$ in the measurement of an optical depth difference $\Delta\tau$ of

$$\delta\Delta\tau = \langle N \rangle^{-1/2}, \quad (2)$$

where $\langle N \rangle$ is the harmonic mean of N_1 and N_2 , the numbers of photons detected at the two resolution elements being compared.

The uncertainty in the measurement of an optical depth ratio for a streaked image with spatial resolution Δx and temporal resolution Δt is thus

$$\delta\Delta\tau = \left(\eta \langle \chi \rangle I_0 \Omega \frac{W}{M} \Delta x \Delta t \right)^{-1/2}, \quad (3)$$

where η is the quantum efficiency of the detector, $\langle \chi \rangle$ is the harmonic mean of the sample attenuations at the centers of the two resolution elements, I_0 is the backlighter power per unit area, Ω is the solid angle of the optic, and W/M is the slit width in the object plane.

Similarly, the uncertainty of an optical depth ratio for a 2-D image with spatial resolutions Δx and Δy along the two axes of the image is

$$\delta\Delta\tau = (\eta \langle \chi \rangle E_0 \Omega \Delta x \Delta y)^{-1/2}, \quad (4)$$

where E_0 is the backlighter energy per unit area. The uncertainties in Eqs. (3) and (4) reflect only shot noise; actual experimental uncertainties are larger owing to imperfections in the optics and in the cameras used to record the image.

The signal-to-noise ratio for the experiment is $\Delta\tau/\delta\Delta\tau$, which depends on the value of the optical depth difference to be measured as well as the required resolution in space or time for streaked images. Equation (3) or (4) can be used to determine this signal-to-noise ratio when the required resolution (Δx , Δt , and Δy) is given, or to determine the spatial and temporal resolution at which a given signal-to-noise ratio is available. These equations are also useful for determining the required slit width for a streaked image or the backlighter duration for a snapshot.

For the example of the $22 \times$ backlighting system to be used in the first Rayleigh-Taylor series, the uncertainty of the optical depth differences is $(10^{-8} \text{ mm} \cdot \text{ns}/\Delta x \Delta t)^{1/2}$ for a streaked image and $(0.25 \times 10^{-8} \text{ mm}^2/\Delta x \Delta y)^{1/2}$ for a 2-D flash image. These figures assume a mean sample attenuation $\chi = 0.1$, a detector quantum efficiency $\eta = 0.1$, a slit width $W = 2 \text{ mm}$ in the image plane for the streaked image, and a backlighter flash duration of 200 ps for the 2-D image. For this experiment, a spatial resolution of $\Delta x < 25 \mu\text{m}$ and a temporal resolution of $\Delta t < 200 \text{ ps}$ are required. Hence, for the streaked image, $\delta\Delta\tau > 1.4 \times 10^{-3}$ is the limiting possible uncertainty in $\Delta\tau$.

For the 2-D image, if a resolution of $\Delta y < 50 \mu\text{m}$ is required along the second axis, then the limiting uncertainty in $\Delta\tau$ is $\delta\Delta\tau > 1.4 \times 10^{-3}$, just as for the streaked image. Because the minimum optical-depth difference to be measured is on the order of $\Delta\tau = 0.1$, our estimates predict a comfortable signal-to-noise ratio of about 70 for these Rayleigh-Taylor experiments. The signal-to-noise ratio erodes rapidly if more stringent requirements are imposed on spatial and temporal resolution. In any case, the actual performance of the $22 \times$ imaging system cannot be allowed to fall too far below the ultimate performance limit described for the currently planned Nova experiments.

Author: D. H. Munro

Direct-Drive Cryogenic Capsule Design Employing DT-Wetted Foam

Introduction. When driven at low input energy and power, high-gain single-shell ICF designs require that the fuel be levitated as a cryogenic shell. A rough estimate of the shell thickness required can be easily derived. A target with 1-MJ input and a gain of 100 must burn 300 μg of DT. Assuming a burnup fraction on the order of 0.5 and an initial aspect ratio (shell inner radius/shell thickness) on the order of 2, we find that the shells must be a few hundred microns thick.

Fabrication of such thick shells of cryogenic DT to the high degree of uniformity believed necessary for successful implosion and efficient burn, presents a formidable technological challenge. A solid layer must pass through the liquid state as it cools, and the time for gravitational slump is short compared to the time to conductively cool past the triple point.

As discussed in "A New Direct-Drive ICF Capsule," in Sec. 3, the use of a low-Z, low- ρ , rigid foam matrix to define and stabilize the layer appears to be a promising fabrication technique. In addition to making near-term target experiments possible, this technique offers encouraging prospects for scaling to reactor target production rates (about 10^6 per day) at an acceptable unit cost. The effect of the presence of the foam material on the hydrodynamics, ignition, and burn propagation of the target must also be addressed. As a first step in studying these issues, we report on a series of 1-D LASNEX simulations of targets directly driven with 0.265- μm laser light.

LASNEX Model. We studied simple single-shell designs, such as those shown in Fig. 2-3. Neither the thin, full-density CH overcoat nor the gas bubble in the center affect target performance; both could be eliminated if fabrication or fielding were simplified by doing so. The foam/DT was treated as atomically mixed; thus, problems associated with density and absorption inhomogeneities were not addressed. The 1-D laser was assumed to be radially incident.

Targets were driven with simple "stair-step" laser pulse shapes, as shown in Fig. 2-4. Step-to-step power ratios were adjusted to maintain approximately 3:1 pressure steps between the sequential shock waves that were launched. The steps were carefully timed to minimize the shells' inflight aspect ratios (IFAR), subject to the constraint of maintaining the fuel on a nearly degenerate adiabat. Radiation from the hot corona was transported by multigroup diffusion, while the thermal electron diffusion was flux limited at 0.03 times the free-streaming velocity; both are conservative approximations with regard to target performance.

Fig. 2-3. Generic welled-foam design.

Fig. 2-4. Simple stair-step laser power pulse shape used to achieve degenerate compression and implosion.

Target Gain. Figure 2-5 shows target gain as a function of radially incident laser energy. Effects of beam diffraction, Brillouin scattering, and nonradially incident beams missing the imploding shell are not included. The result is virtually independent of initial target aspect ratio so long as the gas pressure in the center is not too high. When the unablated portion of the cryogenic shell is replaced by pure DT, the yield increase is less than 5%.

Hydrodynamic Stability. The IFAR of the imploding shell is one crude measure of a target's robustness against the deleterious effects of Rayleigh-Taylor and

Richtmeyer-Meshkov instabilities. Arguments have been presented in Refs. 3 and 4 that successful implosions must have an IFAR < 50 . Directly driven targets are known to require larger values of IFAR than do comparable indirectly driven designs—a fact frequently cited in favor of the latter type of design.⁵ Figure 2-6 shows IFAR as a function of initial aspect ratio S according to three ways of defining shell thickness. Figure 2-6(a) shows IFAR as a maximum over time; Fig. 2-6(b) shows IFAR when the shell inner surface has imploded to 75% of its initial value. With the pulse-shaping approach employed here, a $1/e$ decrease in density is frequently resolving the shock separation and, hence, may not be a meaningful measure.

Suprathermal Electrons. High-intensity lasers impinging on plasma are known to

generate nonMaxwellian electron distributions with a significant high-energy or suprathermal tail. These energetic electrons can penetrate the target, preheat the fuel, and prevent proper compression and adequate burn. Figure 2-7 shows the late-time portion of the laser intensity for two different driving scenarios at the position where electron density n is equal to 0.1, 0.25, and 1 times the critical density n_{cr} . In Fig. 2-7(a) the final step in the pulse shape is maintained at constant power; in Fig. 2-7(b) an initially slightly higher peak power is decreased smoothly as the target implodes to maintain approximately constant intensity at the critical surface. In both cases, the intensity is that needed to drive a 350-kJ design with $S = 2$. The pulse-shape modification in Fig. 2-7(b) does not affect IFAR.

Experimentally, the generation of suprathermal electrons seems to correlate with the production of Raman scattered light.⁶ For a plane, coherent light wave incident on a planar plasma slab, density gradient stabilization sets threshold intensities for Raman sidescatter (I_{TS}) and backscatter (I_{TB}) at electron density between 0.1 and 0.25 times n_{cr} at $I_{TS}\lambda^2 = 5 \times 10^{16} (L/\lambda)^{-4/3}$ at $n = 0.1 n_{cr}$ and $I_{TB}\lambda^2 = 4 \times 10^{17} (L/\lambda)^{-1} v_g/c$, where I is in W/cm^2 , λ is the laser wavelength in μm , v_g is the laser light group velocity, and c is the

Fig. 2-5. Calculated target gain vs radiantly incident laser energy.

Fig. 2-6. Shell IFAR vs initial aspect ratio S .

speed of light in vacuum. For this target $L/\lambda \approx 1000$; thus, the peak intensity is about five times the density gradient threshold for sidescatter and is below threshold for backscatter.

The effect of suprathermal generation on this target's performance is strongly dependent on the suprathermal spectrum, parameterized by T_{hot} , and on the suprathermal transport model. Figure 2-8 shows yield degradation as a function of suprathermal energy fraction f_h , temperature, T_{hot} , and flux limiter F_{aff} . This design can tolerate f_h between 1 and 40%. Designs with larger S , and hence larger IFAR, can

be driven with laser intensity that is always below Raman threshold.

Conclusion. Within the limited scope of our study for a 1-D, radially incident laser and uniformly mixed foam/DT, we see no apparent fatal flaw to the wetted-foam approach to ICF target fabrication. Target gain is not seriously degraded by the presence of the plastic contaminant. IFAR, laser intensity, hydrodynamic efficiency, and hot electron tolerance are at least as good as previously reported for directly driven targets composed of full-density CH ablators and pure cryogenic DT fuel. The good prospects for near-term, thick-shell

Fig. 2-7. Incident laser intensity to drive a 350-kJ target with $S = 2$.

Fig. 2-8. Target sensitivity to suprathermal electron production. (a) Yield degradation vs f_h for various values of flux limiter F_{aff} with an assumed 30-keV suprathermal spectrum. (b) The f_h to degrade yield by 50% vs T_{hot} using classical transport.

cryogenic experiments and for eventual large-scale economic target production make this technology an exciting one that is deserving of vigorous theoretical and experimental study.

Author: R. A. Sacks

Turbulent Mix Generated by Rayleigh-Taylor and Richtmyer-Meshkov Instabilities

Introduction. Rayleigh-Taylor (R-T) and Richtmyer-Meshkov (R-M) instabilities refer to the growth of perturbations at accelerating and shocked interfaces, respectively. The implosion of ICF capsules is vulnerable to both instabilities, which can lead to shell breakup and mixing of the shell material into the fuel. Such breakup and mixing can, in turn, reduce the thermonuclear yield of the capsule.

In the past⁷⁻⁹ our efforts have focused on single-wavelength perturbations of small amplitude (linear regime) to avoid the mathematical complexity of nonlinear waves. However, actual perturbations on an ICF capsule contain a wide range of perturbation wavelengths. Given enough e-foldings, some waves enter the nonlinear regime as the implosion proceeds (see Fig. 2-9). A different mathematical tool is needed to treat the multiwavelength problem. We have begun to solve the problem this year by using turbulence theory. In this article, we outline our results; details are given elsewhere.¹⁰

Fig. 2-9. Mixing of low- and high-density fluids after implosion is caused by random initial perturbations at the interface.

Experiments on mix generated by the R-T instability at a nonablating interface have been performed by Read and Youngs.¹¹ These authors indicate that h , the mixing length into the heavier fluid, is given approximately by

$$h = 0.07 Agt^2, \quad (5)$$

where $A = (\rho_+ - \rho_-)/(\rho_+ + \rho_-)$ is the Atwood number, $g = \text{acceleration}$, and $t = \text{time}$. The densities of the light and heavy fluids are ρ_- and ρ_+ , respectively.

From Eq. (5) we have developed a model for R-M generated mix.¹⁰ The model predicts that turbulent mixing occurs after a shock passes through an interface with random initial perturbations and that the mixing length for the case of a shock is given by

$$h = 0.14 A \Delta v t, \quad (6)$$

where Δv is the jump velocity induced by the shock. Unlike Eq. (5), which has been experimentally verified, Eq. (6) remains a prediction that will be tested by shock-tube experiments.¹²

Turbulent Energy. As energy is supplied to implode a capsule, most is converted into useful kinetic energy, $1/2 v^2$, where v is the velocity directed towards the center. A fraction of the energy, however, may be converted to turbulent energy, i.e., the swirling of eddies in the mixed region. If too much energy is lost to turbulence, then it becomes a heavy drain on the energy budget, and the capsule may fail because the shell cannot achieve the velocity required to ignite the capsule.

Our models predict¹⁰ that the energy lost to turbulence will be a small fraction of the directed energy: about 2% for R-T generated turbulence and 9% for R-M generated turbulence. When the acceleration is constant, we obtain turbulent energy E_{turb} by taking the difference in potential energy between unmixed and mixed density profiles (see Fig. 2-10). Then

$$\frac{E_{\text{turb}}}{E_{\text{dir}}} = \frac{\frac{g}{2}(\rho_+ - \rho_-)h^2 - g \int_{-h}^{+h} \rho y dy}{\frac{1}{2}v^2 \int_{-h}^{+h} \rho dy}. \quad (7)$$

The first term in the numerator is the potential energy (PE) of the unmixed fluids from which we subtract the PE of the mixed profile. If we assume, as indicated in Fig. 2-10, that the density profile $\rho(y)$ is linear, then the second term gives $(g/3)(\rho_+ - \rho_-)h^2$, so that $E_{\text{turb}} = (g/6)(\rho_+ - \rho_-)h^2$. The denominator is directed energy E_{dir} . Setting $v = gt$ and $\int \rho dy = (\rho_+ + \rho_-)h$, we obtain $E_{\text{dir}} = 1/2(gt)^2(\rho_+ + \rho_-)h$. Hence,

$$\frac{E_{\text{turb}}}{E_{\text{dir}}} = \frac{1}{3} \frac{Ah}{gt^2} = \frac{0.07}{3} A^2, \quad (8)$$

which can be, at most, 2.3% ($A < 1$).

In our model, we represent a shock as an impulsive acceleration. It is then straightforward to show (see Ref. 10) that the ratio $E_{\text{turb}}/E_{\text{dir}}$ is four times that of constant acceleration; hence, it cannot be more than 9.3%.

Canuto-Goldman Model and Eddy Sizes. We have also used the analytical model of Canuto and Goldman¹³ to calculate the turbulent energy for a class of instabilities having a power-law growth rate: $n(k) = Ck^r$, where C is a constant that may depend on r . The R-T and R-M instabilities fall in this class, with $n(k) = (gAk)^{1/2}$ and $n(k) = \Delta v Ak$, respectively. We discuss only these two cases, which are relevant to ICF capsules.

In the Canuto-Goldman model we need to specify, in addition to the growth rate, the size of the largest (λ_0) and smallest (λ_{min}) eddies, respectively. They arise as integration limits such that

$$E_{\text{turb}} = \frac{-1}{2\gamma} \int_{k_0}^{k_{\text{min}}} \frac{1}{k^2} \times \frac{d}{dk} \left[kn^{1/2} \int_{k_0}^k kn^{1/2} \frac{d}{dk} \left(\frac{n}{k^2} \right) dk \right] dk, \quad (9)$$

where $k_0 = 2\pi/\lambda_0$, $k_{\text{min}} = 2\pi/\lambda_{\text{min}}$, and γ is a constant that can be calculated within the model. Our approach¹⁰ has been to calculate E_{turb} from Eq. (9) as a function of λ_0 and λ_{min} , equate it to E_{turb} as calculated above assuming a linear density profile, and thereby obtain certain constraints on λ_0 and λ_{min} .

Figure 2-11 shows the result for the case of the R-T instability. We plot λ_0/h as a

Fig. 2-10. Density profiles for unmixed and mixed fluids.

Fig. 2-11. As $\lambda_0/\lambda_{\text{min}} \rightarrow \infty$, the ratio $\lambda_0/h \rightarrow 40\%$. Both λ_0 and h grow quadratically with time.

function of $\lambda_0/\lambda_{\text{min}}$. Two conclusions can be drawn. First, if $\lambda_0/\lambda_{\text{min}} \rightarrow \infty$, i.e., the smallest eddies λ_{min} are much smaller than the largest eddies λ_0 , then λ_0/h approaches a limit, which is approximately 40%. More explicitly,

$$\lambda_0/h \rightarrow \frac{\pi}{2} \left(\frac{9\pi}{56} \right)^2 \approx 40\%. \quad (10)$$

Second, the ratio $\lambda_0/\lambda_{\text{min}}$ must be greater than about 10. This result is obtained by requiring that the largest eddy size λ_0 be no larger than the mixing length h , i.e., $\lambda_0 < h$, which from Fig. 2-11 occurs only if $\lambda_0/\lambda_{\text{min}} > 10$.

Equation (10) implies that the largest eddies λ_0 are proportional to h . As h grows quadratically with time [see Eq. (5)], so does λ_0 , such that the ratio λ_0/h remains constant at about 40%. This observation is consistent with the snapshots presented by Read and Youngs.¹¹

No such relation is obtained when we apply the Canuto-Goldman model to turbulence generated by R-M instabilities. Instead, only one relation is obtained:

$\lambda_0 \approx 1.6 \lambda_{\min}$, i.e., the largest eddies must be about 1.6 times larger than the smallest. This outcome is in strong contrast to the R-T case, where λ_0 must be at least 10 times $\lambda_{\min}$. One possible interpretation is that R-T instabilities generate mostly chunk mix, while R-M instabilities generate mostly atomic mix if $\lambda_{\min}$ is very small.¹⁰

Finally, consider two other systems: one that is shocked at $t = 0$ and one that begins a constant acceleration at $t = 0$. In Fig. 2-12 we plot the turbulent energy per unit mass for both systems as a function of scaled time $gt/2\Delta v$, which is also the ratio of the distance $1/2gt^2$ covered by the accelerating or R-T system to the distance $\Delta v t$ covered by the shocked or R-M system. The ratio of the velocities of the two systems is also shown at the top of Fig. 2-12. By the time the R-T system has covered half the distance of the R-M system ($X_{R-T} = 1/2 X_{R-M}$), it has the same velocity Δv as the R-M system, but its E_{turb} is only 1/4 that of the R-M system. By the time the R-T system catches up ($X_{R-T} = X_{R-M}$), it is

moving at double the R-M system speed and has the same E_{turb} as the R-M system. After this time, E_{turb} continues to increase for the R-T system.

The ratio of the two turbulent energies is given by

$$\frac{E_{\text{turb}}^{\text{R-T}}}{E_{\text{turb}}^{\text{R-M}}} = \left(\frac{X_{\text{R-T}}}{X_{\text{R-M}}} \right)^2. \quad (11)$$

One implication of our outcome is that, if the shell of an ICF capsule must move a certain distance before ignition, then E_{turb} will be the same whether the shell covers that distance by a shock or a constant acceleration.

Summary. Our models predict a relatively small loss of energy to turbulence. Ablation effects, which are not included in our models, may further reduce E_{turb} because perturbations tend to grow slower at an ablating surface. Compressibility is also not included in our models, and its effect on turbulence is not well understood.

We are continuing to study what other models predict about turbulence: in particular, if the assumption of linear density gradients is crucial. Preliminary results indicate that it is not. While mix remains a possible source of yield degradation because it can cool the fuel gas, it is comforting that no more than a few percent of the laser energy can be lost to turbulence during capsule implosions.

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Fig. 2-12. Turbulent energy per unit mass in a shocked (R-M) system and an accelerating (R-T) system vs scaled time. The R-T system has less turbulent energy until it catches up ($X_{R-T} = X_{R-M}$) with the R-M system.

Progress in the Analysis of Selenium X-Ray Laser Targets

A previous experimental series on the Novette laser, using a specially designed exploding selenium foil,¹⁴ succeeded in demonstrating¹⁵ laser amplification at 206 and 209 Å. The amplification was attributed to $J = 2$ to $J = 1, 3p$ to $3s$ transitions in neon-like selenium. However, the theoretical foundation for predicted gain resulted in fundamental differences between data and theory.^{14,16}

In this article, we review progress during 1985 that has advanced our understanding of the selenium x-ray lasing system. We

also predict some potential successes and problems as we pursue the goal of demonstrating x-ray lasing at shorter wavelengths.

The major discrepancy between theory and data is the low gain observed on the $J = 0$ to $J = 1$, 3p to 3s transition compared to the predicted gain of about 10 cm^{-1} . In addition, the observed $J = 2$ to $J = 1$ gain was about 6 cm^{-1} , while the predicted value was 3 to 4 cm^{-1} .

The theory and predictions made last year contained at least one significant shortcoming. While it was widely recognized that dielectronic recombination could play an important role in atomic-level population kinetics, such recombination was only modeled in a crude manner. We used a reasonable guess as to the total rate, but connected that total rate from ground state to ground state of isoelectronic sequences. This rate was apparently correct because it reproduced gross features of $n = 3$ to $n = 2$ x-ray spectra that gave us an indication of ionization balance (approximately 40% fluorine-like and 20% neon-like selenium). However, in dielectronic recombination (e.g., from a fluorine-like ion), a ground-state electron is elevated to an $n = 3$ level or higher, while a free electron decreases in energy by an equal amount, from the continuum into a high-lying level. While autoionization is the inverse process and can undo the doubly excited neon-like state, the ion can stabilize in other ways: in particular, by radiative stabilization. That process can leave the neon-like ion in a singly excited state. Cascade processes that follow favor populating the $J = 2$ levels over the $J = 0$ levels because of their higher (by a factor of 5) statistical weight. In addition, the 3s levels can be populated in this manner, thus lowering the gain on the $J = 0$ to $J = 1$, 3p to 3s transition. When we crudely modeled dielectronic processes by proposing that the flow from fluorine-like to neon-like is directly through their ground states, rather than (correctly) through excited states, we may have correctly predicted ionization balance, but eliminated the possibility of lowering $J = 0$ gain and raising $J = 2$ gains.

As a result of such deficiencies in previous theory, we have begun to include all the doubly excited states resulting from

dielectronic recombination in the atomic physics input-data files. These files are read (along with hydrodynamic quantities from a 2-D LASNEX simulation) into the XRASER computer code. While the process of inclusion is not yet complete, it is of interest to compare our new results for gain prediction with those of last year.

The 2-D LASNEX simulation models our nominal target, which consists of a 750-Å coating of selenium on 1500 Å of a plastic called Formvar. A 2ω 2-beam Novette laser impinges on the foil with a line focus that is 200 μm in height at an intensity of $4 \times 10^{13} \text{ W/cm}^2$ in a 450-ps, full-width-at-half-maximum (FWHM) Gaussian pulse.

The 2-D results do not differ fundamentally from those obtained by 1-D simulations, discussed in the next article, "Hydrodynamics of Exploding-Foil X-Ray Laser Targets." However, the peak electron temperature for the 1-D simulation is somewhat higher than that predicted by the 2-D simulation. The temperature, density, and velocity history is fed into the XRASER calculation of gain, along with the detailed atomic physics model that includes most of the relevant doubly excited states formed by dielectronic recombination. For simplicity, we consider the conditions for a central computational zone that describes a typical part of the main lasing region.

The first result of interest from our calculation is that the ionization balance of 40% fluorine-like and 20% neon-like selenium is reproduced by the XRASER calculation, thus lending credence to our previous estimates. In Figs. 2-13 and 2-14, we plot the prediction for gain vs time for the $J = 0$ to $J = 1$ and $J = 2$ to $J = 1$ lasers, respectively. Times early with respect to the peak of the driving laser pulse are largely irrelevant because the foil has not fully exploded. Thus, the initial high density and steep density gradients of the foil deny the laser system full access to gain down its full length, and refraction bends the laser out of the high-gain medium.^{14,16} Times much later than the laser peak are uninteresting because the rapidly falling density leads to falling gain. The time of greatest interest is roughly the 100 ps immediately after the driving laser peak, which optimizes gain vs a nonrefractory density gradient.

Fig. 2-13. Predicted gain vs time for $J = 0$ to $J = 1$, $3p$ to $3s$ (183-Å) transitions in neon-like selenium.

At +100 ps in Fig. 2-13, the gain on $J = 0$ is about 7 cm^{-1} ; the gain on the upper 209-Å $J = 2$ line shown in Fig. 2-14 is about 6.5 cm^{-1} . (The 206-Å $J = 2$ line has a similar gain.) Compared with our previous model, the $J = 0$ gain is slightly reduced due, in part, to a slightly higher $3s$ population. The $J = 2$ gain is nearly doubled, due to our directly populating the upper $J = 2$, $3p$ levels via the decay of the doubly excited states. Nevertheless, the predicted $J = 0$ gain remains too high with respect to the data. Although we have, as yet, no complete explanation for the discrepancy, we now have obtained a predicted gain for $J = 0$ that is comparable to that for $J = 2$, and the $J = 2$ predicted gain is now quite close to the observed value.

Another issue arising from our data is the apparent lack of lasing lines from fluorine-like states. Even with our recent and more complete modeling, we continue

to predict a gain of about 2.5 cm^{-1} for fluorine-like lasing at 204 Å. While the value is significantly less than the gain of 6 cm^{-1} observed for $J = 2$ neon-like lines, it is curious that the lines remain unobserved given the large fraction of ions in the fluorine-like state.

As another topic, we have performed studies of the laser medium under conditions that can be characterized as non-optimal, i.e., arising from a subtle hydrodynamic effect. The nominal line focus of 1 cm is not uniform along the full 1-cm length. The laser intensity at the extreme edges of the 1-cm length is lower by a factor of 2 than that in the central part of the line of focus. This observation is based on both equivalent-plane optical measurements and on x-ray pinhole pictures. We have processed the LASNEX simulations at various intensities and created an x-ray pinhole calibration curve. From that curve and the experimental pinhole picture, we

Fig. 2-14. Predicted gain vs time for $J = 2$ to $J = 1$, $3p$ to $3s$ (209-Å) transitions in neon-like selenium.

Fig. 2-15. Electron density vs position perpendicular to the lasing-tube axis for three incident laser intensities.

have confirmed that the laser intensity at the edge is one-half that at the center.

Figure 2-15 shows the density profile of an exploding foil 100 ps after the peak of a 220-ps, FWHM Gaussian driving laser pulse. The nominal intensity I_0 is $7 \times 10^{13} \text{ W/cm}^2$. The profile for the intensity $I_0/2$ represents the density profile seen by the laser as it tries to exit the 1-cm-long tube (the intensity at the edge is reduced by a factor of 2). The outcome is not highly deleterious with respect to refraction. However, for an off-optimal shot, the central intensity may be $I_0/2$. In this case, the intensity at the edge is $I_0/4$, and the density profile at that edge shown in Fig. 2-15 is much steeper (more density) than the profile in the center of the lasing tube. The result can lead to deleterious refraction effects as the laser beam leaves the laser tube. This is a specific example of a general phenomenon: namely, off-optimal targets are increasingly sensitive to the $1/2$ edge effect.

We have made significant progress in developing better techniques for measuring density profiles. The work continues to be performed by KMS Fusion Inc., Ann Arbor, Mich. Last year we presented holographic interferograms that served as a significant check on our hydrodynamic-modeling capabilities.^{14,16} Some previous limitations to that data have now been overcome by the recent 2-beam capability at KMS Fusion. Compared to the single 20-ps snapshot per shot, KMS Fusion can now produce four exposures per shot, separated by about 120 ps. Last year the target holder obscured the critical central region of the foil; this year a special target holder was designed to allow diagnostic access to the central region of the foil.

We can now measure the electron-density profile (including the electron density n_e at the central lasing region) resulting from the two-beam irradiation of a nominal, thin selenium foil. These measurements are at the nominal intensity per

beam and pulse width described earlier in this article. Moreover, we can observe the evolution of the profile via multiple snapshots taken at successive times t_1 to t_4 from the same shot. Figures 2-16 and 2-17 show the holograms and data for a target of 750-Å selenium on 500-Å Formvar. In Fig. 2-16, a target is irradiated by a 2ω laser at 9×10^{13} W/cm² for 470 ps. In Fig. 2-17, a target is irradiated by a 2ω laser at 1.6×10^{14} W/cm² for 350 ps. The excellent agreement between our observations (data points) and the LASNEX simulations (solid lines) increases our confidence in modeling these systems.

The Nova two-beam target chamber will allow for two-sided illumination with a line focus of up to 5 cm in length. As we move to shorter-wavelength x-ray lasing schemes, our first priority is to reach saturation on our selenium scheme. Figure 2-18(a) shows the trajectories of x-ray laser beams as they traverse the length of the exploding foil for distances

Fig. 2-16. Holograms, taken at four early times, of 750-Å selenium foil illuminated from both sides. Observed electron densities closely match LASNEX simulations.

greater than 1 cm. The onset of severe degradation beyond 2 cm is quantified in Fig. 2-18(b).

We are attempting to circumvent refraction limitations to reaching saturation by inventing new target geometries for the x-ray laser. One scheme involves exploding foils that face one another and explode toward each other. This scheme creates a temporary trench that traps the x-ray laser beam in a region of high gain and allows propagation over distances much longer than 1 cm. At 3 cm, we predict a degradation of only 10% compared to the 90% degradation for the single foil shown in Fig. 2-18(b). Preliminary hydrodynamics experiments at KMS Fusion on this scheme support our optimism. However, the new target forms a confining trench in only one dimension. We are therefore pursuing cylindrical exploding foils as a 2-D analog of the trench targets. Other targets with tamping from thick Formvar substrates also show promise as nonrefractory lasing media.

Shorter-wavelength schemes using nickel-like analogs to our neon-like system or rapid-cooling recombination schemes are presented later in this section. See "Estimates of Gain for Nickel-Like Exploding-Foil Targets" and "Recombination/Rapid-Cooling X-Ray Lasers." Here we present calculations for the gain of a higher-Z, neon-like system: namely molybdenum ($Z = 42$).

The nominal target is 500 Å of molybdenum on 1000 Å of Formvar. The target is irradiated by a 2ω laser at about 2.5×10^{14} W/cm² for 500 ps. At the peak of the driving laser, electron temperatures reach 2 keV, densities reach 3×10^{20} cm⁻³, and the scalelength is about 100 μm. Lasing with a gain of about 3 cm^{-1} is predicted for both the $J = 2$ and $J = 0$ transitions.

An interesting change takes place in proceeding to higher Z. For selenium, the $J = 0$ was at 182 Å, shorter than the two $J = 2$ transitions at 206 and 209 Å. For molybdenum, the $J = 0$ at 139 Å is a longer

Fig. 2-17. Holograms, taken at four late times, of 750-Å selenium foil illuminated from both sides. Observed electron densities closely match LASNEX simulations.

Fig. 2-18. (a) Trajectories of x-ray laser beam viewed from above an exploding foil and propagating over a distance of 3 cm. (b) Radiated laser power vs rod length calculated with and without refraction effects.

wavelength than either of the two $J \approx 2$ transitions at 131 and 133 Å. The cross-over occurs near yttrium ($Z = 39$).

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Hydrodynamics of Exploding-Foil X-Ray Laser Targets

In this article, we present a similarity model for the hydrodynamics of laser-heated exploding foils. Such foils have recently been used to create soft x-ray lasers in the laboratory.^{17,18} The foils are also useful for experiments in spectroscopy and plasma physics. To date, the design of exploding-foil x-ray lasers has relied heavily on complex numerical hydrodynamics simulations.¹⁷ Our simple model appears to be a powerful design tool that reproduces the essential results of the more expensive simulations over a large and important range of parameters.

Models for laser-heated exploding foils, wires, and spheres have been discussed in previous publications.¹⁹⁻²⁰ Our approach is similar to the previous work in that we make a few assumptions that allow for a reduction of the hydrodynamic, partial differential equations to a set of coupled, ordinary differential equations. Our extensions beyond previous work include incorporation of the specific laser heating term, derivation of analytical solutions for special cases, and extensive checks of the model against LASNEX numerical hydrodynamic

simulations. In addition, we discuss applications of the model to the design of soft x-ray lasers.

Although the similarity method is easily applied in several spatial dimensions (see Refs. 21 and 26), we discuss only applications in 1-D, plane-parallel geometry. Our experience with numerical simulations shows that many essential results can be extracted from 1-D considerations.

Details of the Similarity Equation: To obtain the necessary similarity equations, we make two simplifications to the ideal, single-fluid hydrodynamic equations. First, we assume the gas is isothermal in space. Second, we assume the expansion is homogeneous—in other words, the velocity V is always a linear function of position.

By defining a time-dependent scale factor L , the condition for V is written:

$$V = \left(\frac{dL/dt}{L} \right) x. \quad (12)$$

With our two assumptions, we reduce the hydrodynamic equations to one algebraic and two ordinary differential equations

$$\text{continuity: } \rho = \frac{m}{\sqrt{2\pi}L} \exp\left(\frac{-x^2}{2L^2}\right), \quad (13)$$

$$\text{momentum: } L \frac{d^2 L}{dt^2} = \frac{p}{\rho}, \quad (14)$$

and

$$\text{energy: } c_v \frac{dT}{dt} = H = \frac{p}{\rho} \frac{dL}{dt} L^{-1}. \quad (15)$$

The term x is the spatial coordinate measured perpendicular to the foil surface; t is time; ρ , v , p , and T are mass density, velocity, pressure, and temperature, respectively; c_v and H are heat capacity and net heating, respectively, per unit mass; and m is the (constant) mass column density of the foil.

The equation of state and heat capacity are for an ideal gas with contributions of electrons to the pressure and energy and contributions of ions to the mass. In addition, we assume that the ionization state is constant over time. Although a more general treatment is allowed in the context of the similarity solution, the assumption of constant ionization state is sufficiently accurate to predict the hydrodynamic behavior of the foil, as shown by comparisons to LASNEX simulations given later. We choose the ionization state based on experience with the simulations.

The laser deposition is due to inverse bremsstrahlung. We distribute the energy equally per unit mass in the foil, consistent with the isothermal assumption. We ignore reflection of laser light at the critical surface. Because we are mainly interested in conditions after the density drops below critical, the two preceding assumptions are not highly limiting (see discussion of limitations below).

The heating rate is found from the absorbed fraction of laser flux by:

$$H = \frac{I}{m} [1 - \exp(-\tau_{ib})], \quad (16)$$

where I is total laser intensity on the foil, $\tau_{ib} = \int \kappa_{ib} dx$ is inverse-bremsstrahlung optical depth through the foil, and κ_{ib} is opacity.²⁷ Equation (16) states that the heating rate is constant for large values of τ_{ib} ($H \sim I/m$), while for small values the heating rate decreases [$H \sim I/m(\tau_{ib})$]. In the evolution of an exploding foil, τ_{ib} generally decreases with time as the density drops. As a result, the heating rate is constant at early times and decreases at later times. The time at which τ_{ib} becomes unity is denoted as the transparency time t_{trans} and is a characteristic fiducial in analyzing foil behavior.

Analytic Solutions of Similarity Equations. Useful scaling laws for the plasma parameters have been derived from analytic

solutions for flat-top laser pulses. The input parameters of the solutions are laser intensity I , laser wavelength λ , foil mass m , ion charge Z , and atomic number of the foil element A . The independent output parameters are T and L .

Figure 2-19 shows the temperature and Mach number histories from a numerical solution of Eqs. (13) to (15). The input parameters (those of a typical selenium x-ray laser target¹⁷) and main results of this run, which we refer to as the standard case, are given in Table 2-1. We discuss the three distinct time periods shown in Fig. 2-19 in order.

After the order of the sound crossing time at the initial temperature, an early asymptotic solution is approached for which the heating rate is constant. The following power-law functions satisfy the similarity equations:

$$T = \frac{1/3 HM_1}{Z} t = 11.1 \bar{I} \bar{m}^{-1} \bar{Z}^{-1} \bar{A} t, \quad (17)$$

$$L = \frac{2}{3} H^{1/2} t^{3/2} = 6.67 \times 10^{-2} \times \bar{I}^{1/2} \bar{m}^{-1/2} \bar{Z} \bar{A}^{-1} t^{3/2}, \quad (18)$$

and

$$n_e = 1.12 \times 10^{20} \bar{I}^{-1/2} \bar{m}^{3/2} \bar{Z} \bar{A}^{-1} t^{-3/2}. \quad (19)$$

Fig. 2-10. Three characteristic time periods exhibited in the numerical solution of exploding-foil similarity equations. Histories of temperature and Mach number are shown for the standard model (see Table 2-1).

Input parameter			Value		
	I		$6.5 \times 10^{13} \text{ W} \cdot \text{cm}^{-2}$		
	λ		$0.53 \mu\text{m}$		
	t_L		0.45 ns		
	m		$36 \mu\text{g} \cdot \text{cm}^{-2}$		
Output quantity					
	T_e (0.45 ns)	n_e (0.45 ns)	Fraction absorbed	Radiated/ absorbed	t_{trans}^a (ns) t_{th}^b (ns)
Model	1.38	1.17×10^{20}	0.173	0	0.043
LASNEX	1.34	1.29×10^{20}	0.171	0.09	0.083
^a For model.					
^b For LASNEX					

Table 2-1. Input parameters and output quantities for standard foil model.

Physical variable	Symbol	Scaling value
Laser intensity	I	$10^{14} \text{ W} \cdot \text{cm}^{-2}$
Laser wavelength	λ	$0.53 \mu\text{m}$
Foil-column density	m	$10^{-4} \text{ g} \cdot \text{cm}^{-2}$
Ion charge	Z	25
Atomic mass	A	80
Coulomb logarithm	Λ	5
Time	t	1 ns
Scalelength	L	1 cm
Electron temperature	T_e	1 keV
Laser spot width	W	200 μm

Table 2-2. Units for scaling formulas.

Units for L , T , and n_e are cm, keV, and cm^{-3} , respectively. Variables marked with a tilde are scaled to a set of convenient values listed in Table 2-2. During this period, the foil expands at $\sqrt{3}$ times the instantaneous isothermal sound speed.

The constant-heating period ends at t_{trans} when the foil becomes optically thin to laser light. The optical depth, obtained by evaluating Eq. (16) with the Gaussian density profile [Eq. (13)] is given by

$$\tau_{\text{th}} = 9.52 \times 10^{-3} \tilde{I}^2 \tilde{m}^2 \tilde{Z}^3 \times \tilde{A}^{-2} \tilde{\Lambda} T^{-3/2} L^{-1}, \quad (20)$$

where $\tilde{\Lambda}$ is the coulomb logarithm, evaluated at the foil center. Using Eqs. (17) to (19) in Eq. (20), we find that

$$t_{\text{trans}} = 0.157 \text{ ns } \tilde{I}^{-2/3} \tilde{\Lambda}^{2/3} \tilde{m}^{4/3} \times \tilde{Z}^{3/2} \tilde{A}^{-7/6} \tilde{\Lambda}^{1/3}, \quad (21)$$

The temperature at this time is

$$T_{\text{trans}} = 1.74 \tilde{I}^{1/3} \tilde{\Lambda}^{2/3} \tilde{m}^{1/3} \times \tilde{Z}^{1/2} \tilde{A}^{-1/6} \tilde{\Lambda}^{1/3}. \quad (22)$$

For $t_{\text{trans}} < t < t_L$, the temperature is kept nearly constant by a balance of laser heating and expansion cooling [see Fig. 2-19 and Eq. (15)]. We find power-law solutions modified by slowly varying logarithmic functions of time. Adequate scaling laws are found by fitting the logarithmic factors to weak power laws and matching to the pre-transparency solutions [Eqs. (17) to (19)] at t_{trans} . The resulting temperature is approximately constant at T_{trans} , while the scalelength and density are given by

$$L = 4.20 \times 10^{-2} \tilde{I}^{1/3} \tilde{\Lambda}^{1/6} \tilde{m}^{-1/6} \times \tilde{A}^{-7/24} \tilde{Z}^{3/8} \tilde{\Lambda}^{1/12} \tilde{I}^{5/4}, \quad (23)$$

and

$$n_e = 1.78 \times 10^{20} \tilde{I}^{-1/3} \tilde{\Lambda}^{-1/6} \tilde{m}^{7/6} \times \tilde{A}^{-17/24} \tilde{Z}^{5/8} \tilde{\Lambda}^{-1/12} \tilde{I}^{-5/4}. \quad (24)$$

The numerical coefficients in Eqs. (22) to (24) are only approximate because of the nature of the matching procedure. When compared to the numerical solution at the end of the laser pulse, for example, Eq. (23) gives a scalelength that is 50% too low, while Eq. (22) gives a temperature that is 10% too low. The relative scaling with input parameters is expected to be quite good. This period is of interest for x-ray lasers because the conditions are most uniform. It is particularly interesting that a constant temperature is maintained. During a Gaussian laser pulse, Eqs. (22) to (24) give reasonable estimates of conditions at the peak of the pulse.

After the laser is turned off, at time t_L , the foil continues to expand adiabatically. The scalelength now increases strictly linearly with time, while the temperature drops with the $2/3$ power of time, such that

$$T = T_{\text{trans}} \left(\frac{t}{t_L} \right)^{2/3}. \quad (25)$$

The scalelength and density are given by Eqs. (23) and (24), with the factor $t^{5/4}$ replaced by $t_L^{5/4} t$.

Limitations of Similarity Model. We estimate the range of validity of the similarity model by examining the approximations

made in the derivation. The isothermal and linear-velocity approximations will usually not be valid until laser burn-through. We estimate the burn-through time from the mass-ablation rate of a laser heat front^{28,29} such that

$$t_{bt} = 0.32 \tilde{I}^{-2/3} \tilde{\lambda}^2 \tilde{m}^{-4/3} \times \tilde{A}^{-7/6} \tilde{Z}^{2/3} \tilde{\Lambda}^{1/3}, \quad (26)$$

The derivation assumes that laser light is absorbed by inverse bremsstrahlung before reaching the critical surface as expected, except for extremely high intensities. The similarity model cannot be applied to laser pulses shorter than t_{bt} .

After laser burn-through, a problem with isothermality may occur if conduction is not efficient. To study this problem, we have compared the time scale for thermal equilibration via classical conduction³⁰ to the expansion time scale [$\equiv L/(dL/dt)$] during the post-transparency period. We find that isothermality should be maintained for laser fluxes larger than a certain critical value I_{cond} such that

$$I_{cond} = 5.1 \times 10^{11} \tilde{A}^{-3} \tilde{\lambda}^{-7/4} \times \tilde{Z}^{9/4} \tilde{\Lambda}^{1/2} \tilde{t}^{1/2}. \quad (27)$$

The condition $I > I_{cond}$ is not too restrictive; most of the designs of practical interest satisfy this constraint.

The assumption of 1-D, plane-parallel flow breaks down after a time t_{2d} when the transverse scalelength L becomes equal to the width of the laser spot W . Assuming that this occurs after t_{trans} , we find

$$t_{2d} = 5.52 \times 10^{-1} \text{ ns } \tilde{I}^{-4/15} \tilde{\lambda}^{-2/15} \tilde{m}^{2/15} \times \tilde{A}^{7/30} \tilde{Z}^{-3/10} \tilde{\Lambda}^{1/15} \tilde{W}. \quad (28)$$

At high laser intensities, a host of laser-plasma instabilities may set in, some of which reflect laser light, thus compromising our absorption model. A discussion of these effects is found in Ref. 31. The most critical process in our parameter range appears to be stimulated Brillouin scattering

(SBS). The intensity threshold during the transparent phase is given by

$$I_{SBS} = 8.95 \times 10^{14} \text{ W} \cdot \text{cm}^{-2} \tilde{\lambda}^{-7/2} \tilde{m}^{-1} \times \tilde{A}^{5/4} \tilde{Z}^{-3/4} \tilde{\Lambda}^{1/2} Q^{3/2}. \quad (29)$$

The threshold for stimulated Raman sidescatter (SRS) during the transparent phase can be written

$$I_{SRS} = 3.65 \times 10^{13} \text{ W} \cdot \text{cm}^{-2} \tilde{\lambda}^{-8/13} \tilde{m}^{2/13} \times \tilde{A}^{7/26} \tilde{Z}^{-9/26} \tilde{\Lambda}^{-1/13} \tilde{t}^{-15/13}. \quad (30)$$

This process is not really a limitation on the validity of the model because the main energy transfer is from the laser to hot electrons, which couple well into thermal energy of the exploding foil. Some amount of SRS is actually desirable for use as a temperature diagnostic, as described in Ref. 17. Within the bounds discussed above, we expect the similarity model to accurately reflect the physics of exploding foils.

Numerical Solutions of the Similarity Equations. We have performed extensive comparisons of numerical solutions of the similarity equations to LASNEX simulations. For the LASNEX runs, we used the multigroup, diffusion-radiation package. The electron transport flux-limiting factor was usually taken to be 0.05, and other options were standard for 1-D slab geometry. The only free parameter in the similarity solution was the ion charge, which was chosen to agree with the LASNEX result.

We first discuss the results for our standard x-ray laser target (see Table 2-1) with a 450-ps, flat-top pulse. Figure 2-20 shows that the plasma becomes nearly isothermal after the initial burn-through of the foil. The assumption of isothermality in the similarity solution is therefore accurate for a well-designed x-ray laser foil.

The electron-density profile after burn-through is shown in Fig. 2-21. The density profile is nearly Gaussian, as predicted by the similarity model. With respect to middle and late time histories of both electron density and temperature in the center of the foil (see Fig. 2-22), the agreement between the LASNEX curves and the curves for the similarity model is good. It is of

Fig. 2-20. Foil becomes isothermal early in the laser pulse as shown by the series of temperature snapshots from LASNEX. The foil is symmetric about the vertical axis.

Fig. 2-21. LASNEX density profile is approximately Gaussian. Density snapshot is taken at 0.45 ns in the standard model.

Fig. 2-22. After laser burn-through of foil, density and temperature histories of LASNEX and the similarity model are in good agreement.

particular interest that the constant temperature reached after transparency is confirmed by LASNEX even though the solutions yield different early-time behavior ($t < 100$ ps). As shown in Table 2-1, agreement to within 10% is obtained for the important results of the standard run. This level of agreement is certainly adequate for design calculations because it is within the range of reasonable uncertainties in the basic physics in the LASNEX code as well as experimental uncertainties, such as those involving laser flux and target thickness.

We also compared several runs with variations from the standard set of parameters. Figure 2-23 shows that the simple model does a good job of calculating temperature and density histories for a Gaussian and a flat-top pulse. Other parameter studies resulted in good agreement between the similarity model and LASNEX when input parameters were varied about the standard model by a factor of four.

Applications. We have considered several applications of the similarity model to the design of x-ray laser targets. Laser pulse shaping can be studied, noting the $I^{1/3}$ scaling of temperature in Eqs. (22) to (24). It appears to be possible, for example, to raise the temperature at a particular time in the pulse by raising the intensity. Figure 2-24 shows a temperature history from a LASNEX simulation to illustrate this

effect. The temperature follows closely the rise in intensity applied at 0.3 ns. Such a temperature profile may be useful to increase collisional excitation rates at a time when density drops to a value such that ionization rates are small.

The scaling laws for the constant-temperature period have also been used to map the range of temperatures and densities obtainable for a series of survey experiments on selenium laser gain. By varying the input parameters, we predict that a significant range of plasma conditions can be sampled. We intend to study experimentally how these different conditions influence gain. We are currently studying the scaling of neon-like schemes with atomic number, using both the simple hydrodynamic model and an atomic kinetics model. We have also studied the influence of radiative cooling on hydrodynamics by adding a loss term to the energy equation [Eq. (15)]. With reasonable cooling rates, we find little effect on the standard selenium foil.

Future applications include the study of x-ray laser propagation using the simple, Gaussian density profile of Eq. (13). Extensions of the similarity model to include 2-D effects and time-dependent ionization are under investigation.

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Investigation of the Manganese Pump Line for the Photoresonant X-Ray Laser

In the photoresonant scheme,³² a 1s-3p transition in hydrogen-like fluorine is excited by a 2p-3d line from beryllium-like manganese at 12.643 Å. The mismatch between these two transitions has been found to be less than 1.5 mÅ in spectroscopic studies at KMS Fusion, Inc. Lasing is expected to occur from two 2p-3d transitions in hydrogen-like fluorine at 81 Å.

In this article we report calculations carried out at LLNL and experiments conducted by KMS Fusion, Inc. The 1-D LASNEX^{33,34} and XRASER³⁵ calculations were run to obtain the time-dependent shape and position of the 12.643-Å line emitted by the manganese flashlamp, including the Doppler shift introduced by ion streaming motion. The importance

of the Doppler shift can be understood when a typical shift of 8 mÅ (0.6 eV), corresponding to a velocity of 2×10^7 cm/s, is compared to the narrow width (0.04 to 0.06 eV) of the absorption line in the fluorine.

We find that the contribution to the pump line comes from a region in the flashlamp having a large gradient in the ion streaming velocity v . Therefore, with the present flashlamp design, it becomes inherently difficult to produce a pump line with satisfactory intensity at precisely the absorption frequency.

Theory. As the laser pulse deposits energy, density and temperature profiles are established, which propagate through the flashlamp. Manganese with $n < n_a$, where n_a is the density at the ablation surface (comparable to the critical density n_c), gives rise to a red shift seen by a detector on the side of the flashlamp

Fig. 2-23. The similarity model agrees with LASNEX on the temperature and density for a Gaussian laser pulse. The pulse has a 0.45-ns FWHM centered at 0.6 ns.

Fig. 2-24. Effect of pulse shaping shown by the temperature history from LASNEX. Tripling of laser intensity at 0.3 ns causes a temperature rise by a factor of 1.4, in agreement with Eq. (22).

Fig. 2-25. Flashlamp configuration from LASNEX run for
(a) $t = 0$ and
(b) $t = 150$ ps.

Fig. 2-26. Profile of the 12.643-Å beryllium-like manganese pump line calculated by XRASER for $t = 150$ ps.

opposite from the incoming laser pulse (see Fig. 2-25). On the other hand, manganese with $n > n_c$ undergoes recoil and gives a contribution to the emitted line with a blue shift.

The initial conditions corresponding to the KMS Fusion experiments are shown in Fig. 2-25(a). The laser pulse ($\lambda = 0.53 \mu\text{m}$) has its peak at 200 ps with FWHM = 130 ps.

The flashlamp configuration at 150 ps, taken from a 1-D LASNEX calculation, is shown in Fig. 2-25(b).

The magnitude of the Doppler shift is given, in lowest order, by the relation $\Delta\lambda/\lambda = -\Delta v/v = v/c$, so that $\Delta\lambda = 4 \text{ mÅ}$ (0.3 eV) for $v = 10^7 \text{ cm/s}$ and $\lambda = 12 \text{ Å}$. As can be seen from Fig. 2-25(b), velocities in the abating material range from 0.8 to $2.5 \times 10^7 \text{ cm/s}$ (going outward to a minimum density $= n_c/3$), corresponding to red shifts of 3.2 to 10 mÅ. At the same time, recoil velocities range from 0.1 to $0.6 \times 10^7 \text{ cm/s}$, corresponding to blue shifts of 0.4 to 2.4 mÅ. The Doppler width at this time is 0.16 eV.

The 12.643-Å line calculated by XRASER at 150 ps (using input from the 1-D LASNEX run) is shown in Fig. 2-26. The line center ($h\nu = 0$) is located at 979.37 eV (12.66 Å). Negative values of $h\nu$ correspond to red shifts. There are two components to the line: red and blue. Furthermore,

the shifts away from line center, taking into account the Doppler width, are compatible with the simple estimate made above for the Doppler shift based on the LASNEX velocities.

The 1-D calculation yields a red and blue component until $t = 240$ ps, at which time the critical density surface has moved into the Si_3N_4 and only ablation velocities are found in the manganese.

Experiment. Data were recorded with an x-ray streak camera fitted with a beryl crystal ($2d = 15.95 \text{ \AA}$). The spectral resolution at $12.64 \text{ \AA}$ was calculated to be $6 \text{ m\AA}$. The output from the streak camera was amplified using a Mullard electro-optic intensifier and then recorded on Royal-X-Pan film. The streaked data had no absolute time or wavelength reference. Shifts of the peak of a line are approximately $5 \text{ m\AA}$ and are much greater than uncertainties in the measurement/analysis technique.

Intensity vs wavelength for 12.475 to $12.7 \text{ \AA}$ at $t = 150$ ps is shown in Fig. 2-27. The line of interest was found to have a $4.7\text{-\AA}$ red shift and a 1.14-eV FWHM at this time. Referring to Fig. 2-26, XRASER calculated a $4\text{-m\AA}$ red shift and approximately 1.0-eV FWHM. As time proceeds from 150 to 240 ps, the measured FWHM changes by

less than $2 \text{ m\AA}$. The red shift of the peak of the $12.64\text{-\AA}$ line is plotted in Fig. 2-28. The experimental data has been normalized to the calculations at $t = 200$ ps.

Summary. We can understand the XRASER results for the line width by taking into account the Doppler shift of the pump line emitted by the beryllium-like manganese ions moving in the direction of our line of sight. 1-D calculations yield a red shift for the peak of the pump line that is in reasonable agreement with experiment. 1-D calculations yield two distinct

Fig. 2-27. Measured spectrum from the manganese flashlamp from 12.475 to $12.70 \text{ \AA}$. The line of interest is located at $\lambda > 12.64 \text{ \AA}$.

Fig. 2-28. Theoretical and experimental red shift of the peak of the $12.64\text{-\AA}$ line.

components to the pump line for a wide range of times. Experimental measurements do not indicate this splitting; however, the instrumental broadening is 6 mÅ. A detailed comparison with experiment awaits 2-D calculations.

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Estimates of Gain for Nickel-Like Exploding-Foil Targets

Nickel-like ions have been suggested as candidates for amplifiers in the extreme ultraviolet (XUV) and soft x-ray spectral regions.^{36,37} The nickel-like configuration has a closed 3d shell in the ground state that is analogous to the closed 2p shell of the

Table 2-3. XRASER gains for nickel-like europium.

Gain (cm ⁻¹)	$J_u - J_l$	Configuration	Wavelength (Å)
6.9	0-1	4p-4d	68.06
4.4	2-1	4p-4d	100.14
4.3	0-1	4p-4d	63.32
2.9	0-1	3p4s-3p4p	94.85
2.1	1-1	4p-4d	104.20
1.5	0-1	4p-4d	81.04
1.05	2-1	3p4d-3p4f	97.23
1.04	0-1	3p4s-3p4p	151.00

Fig. 2-29. Energy levels and laser transitions for the 8-level model of nickel-like europium. Energies are in eV, and values for strength of transition gf are in brackets. ▼

neon-like configuration. Because the neon-like system has shown significant amplification in selenium,^{38,39} the analogous nickel-like system can be expected to show interesting gains.

In this article we show that an exploding-foil europium target (675 Å of europium on 1000 Å of Formvar) irradiated with 7.5×10^{13} W/cm² (400-ps FWHM pulse, $\lambda = 0.53 \mu\text{m}$) produces plasma parameters appropriate for lasing.³⁷ Representative values from a 1-D LASNEX run^{40,41} are:

- $n_e = 5 \times 10^{20} \text{ cm}^{-3}$,
- $T_e = 800 \text{ eV}$,
- $T_i = 450 \text{ eV}$,
- $n_i = 10^{19} \text{ cm}^{-3}$.

The gains calculated by a zero-dimensional XRASER code⁴² for these parameters and a 25% abundance of nickel-like europium are shown in Table 2-3.

Phenomenological 8-Level Model. We incorporate the Novette results³⁹ in the 8-level model: namely, the absence of the 0-1 transition (largest predicted gain) and a measured gain on the 2-1 lines that is somewhat larger than predicted. Anticipating a similar result for nickel-like targets, we distribute electrons in the upper states (4d) according to statistical weights. The energy-level diagram and possible laser transitions for nickel-like europium, including only the 4p and 4d excited states, are shown in Fig. 2-29.

Solving for the steady-state 4p populations, we obtain the gains listed in Table 2-4. Although the wavelengths in Tables 2-3 and 2-4 are slightly different (two different atomic codes have been used), three common lines at 68, 100, and 104 Å can be seen.

An interesting possibility arises when inner-shell excitations are considered. The $J = 1$, $3p_{3/2}4p_{3/2}$ level at 1432.9 eV has a 335.5-eV transition to the $J = 1$, $3d_{5/2}4p_{3/2}$ state at 1097.4 eV. This is a 36.95-Å transition with $(gf) = 0.217$. The upper state will certainly be populated in any experiment designed to give a reasonable nickel-like abundance. Although the gain will be decreased due to the short wavelength, mirrors could be used to increase the path length. Further work is planned along these lines.

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Recombination/Rapid-Cooling X-Ray Lasers

Introduction. During 1985 we analyzed the experiment on the Novette exploding-foil recombination laser and made comparisons with 2-D modeling calculations. Our analysis provides an explanation for the lack of observed gain and suggests several improvements in the design of future recombination experiments.

In this article, we briefly review the principles behind recombination x-ray laser schemes. We then describe experiments on Novette using an exploding foil made of magnesium. We present recent modeling results showing that regions of the magnesium foil outside of the line focus, which are heated indirectly by conduction, produce significant spontaneous emission. The spontaneous emission makes the observation of induced emission and gain difficult. Modeling calculations also show the degrading effect of intensity variation along the ~ 1 -cm line focus on observed gain. Modifications of the current x-ray laser design using a strip of magnesium or aluminum are presented. Finally, we provide schemes for photoionized recombination x-ray lasers.

Recombination/Rapid-Cooling

Schemes. The basic principles behind recombination/rapid-cooling schemes are quite simple and have been known for some time.^{43,44} We restrict our discussion to lasing in hydrogenic ions; however, most of the conclusions are valid for other ions—in particular, lithium-like ions.

The experimental approach is to fully strip a given atom and then allow the system to cool rapidly. The densities are such that the dominant recombination mechanism is three-body recombination into highly excited levels followed by cascade down to the ground-state level. During the nonequilibrium cascade, population inversions can occur between the lower levels. Population inversions occur because the radiative rates out of a hydrogenic state increase as the energy of the level decreases, allowing the lower state of a given pair to empty more rapidly. This explanation is valid provided that the lower states decay by radiative transitions rather than by collisional transitions. The ratio of radiative to collisional rates is proportional to

Gain (cm^{-1})	$J_u - J_l$	Configuration	Wavelength ($\text{\AA}$)
7.0	2-1	4p-4d	100.38
4.4	1-1	4p-4d	104.53
4.0	1-1	4p-4d	80.40
3.7	2-1	4p-4d	77.97
1.0	0-1	4p-4d	68.68

Z^7 , and for the 4-3 transitions in hydrogen-like magnesium, the electron density must be less than $4 \times 10^{20} \text{ cm}^{-3}$ for radiative transitions to be dominant. The electron density during peak recombination is expected to be $\sim 10^{20} \text{ cm}^{-3}$, therefore, the condition is satisfied. See Ref. 45 for a more detailed explanation of lasing in hydrogenic ions.

The target design for the 12 shots performed on Novette consisted of a 1100- or 2200- $\text{\AA}$ foil of magnesium placed between two 1500- $\text{\AA}$ sheets of Formvar. The Formvar absorbs the rising portion of the optical-laser pulse and helps maintain a more uniform density profile through the magnesium. The reasons for the sandwich design are discussed in more detail in Ref. 46. We suggest replacing the magnesium foil by a strip while maintaining the Formvar layers.

Double-sided illumination is provided by cylindrical lenses focusing the 0.53- μm laser to a region 0.5 to 1.0 cm by 0.01 cm. The intensity is approximately $2 \times 10^{14} \text{ W/cm}^2$ with a pulse duration of 100 to 200 ps.

The potential lasing line with the largest predicted gain, $\sim 1 \text{ cm}^{-1}$, is the 130- $\text{\AA}$, 4-3 line of hydrogen-like magnesium. This line is observed by two types of instruments. High spectral resolution, but limited time resolution, is obtained from two microchannel-plate, grazing-incidence spectrometers, referred to as McPIGS.⁴⁷ These devices view the exploding foil on- and off-axis. We present only on-axis results because a direct comparison is not possible due to lack of calibration. The three channels in the two McPIGS begin time integration of the data before the Novette laser pulse. The individual channels can be turned off at different times. The second device is a transmission-grating streak spectrograph, referred to as the TGSS-EM.⁴⁸ The TGSS-EM trades spectral resolution ($\lambda/\Delta\lambda \approx 200$) for temporal resolution (20 to 30 ps). We present selected results for each of the two complementary devices.

Table 2-4. Gains obtained from the 8-level model.

Fig. 2-30. Early-time data from the on-axis McPIGS, including tentative line identification for hydrogen-like magnesium, helium-like magnesium, and oxygen.

Figure 2-30 gives the result of the on-axis McPIGS for a channel that is shut off approximately 600 ps after the Novette pulse is completed. At this time, a significant amount of recombination is beginning to occur. We have made tentative line identification for hydrogen-like magnesium, helium-like magnesium, and oxygen. The identification of the Mg XII 4-3, O VIII 3-2, O VII 3-2, and Mg XII 5-3 lines is fairly reliable. Lines identified as Mg XII 7-3 and Mg XII 6-3 may be due to O VIII 5-2 and O VIII 4-2, respectively. Identification of the helium-like lines, Mg XI 5-3 and Mg XI 6-3, is tentative.

Results are similar for other channels of the McPIGS, which give longer time integration. The time gating of the three channels provides only limited information on the time history of line emission. However, our data clearly show a significant amount of spontaneous emission that persists for longer than 500 ps after the Novette pulse is completed. A tentative explanation for this unexpected result is given later.

The TGSS-EM allows for temporal resolution, but the loss of spectral resolution presents a problem. Oxygen in the Formvar, $C_{11}H_{18}O_5$, is initially stripped but then recombines, resulting in a broad and strong 3-2 line of helium-like oxygen

at 128 Å. This line overlaps with the potential lasing line of hydrogen-like magnesium at 130 Å. The overlap becomes significant approximately 600 ps after the Novette pulse. The time coincides with the predicted occurrence of gain in the 4-3 line. The width of the oxygen line is enhanced by the large size of our Formvar sheet. The TGSS-EM device convolves perpendicular extent with wavelength. Late in the expansion, the large physical separation of the expanding Formvar, originating from the two sides of the magnesium foil, results in an observed splitting of the 3-2 line.

The time history of the intensity of the 4-3 line of Mg XII observed by the TGSS-EM device is as follows. Intensity initially increases while the Novette laser is on and then remains approximately constant for ~350 ps after the pulse is completed. The intensity then decreases monotonically, reaching a value of one-half the maximum intensity ~700 ps after the Novette pulse is completed. Finally, the O VII 3-2 line begins to obscure the Mg XII 4-3 line.

Figure 2-31 shows TGSS-EM data for five different times, starting at the latter portion of the 200-ps Novette pulse. Figure 2-31(a) is taken at 350 ps (~50 ps

Fig. 2-31. TGSS-EM data for five different times. Times refer to the interval commencing approximately 200 ps prior to the Novette pulse.

before the pulse is completed) and shows the Mg XII 4-3 line emerging from the background. The wavelength scale was not adjusted, therefore the line appears at ~ 132 Å rather than at 130 Å. The time shown in Fig. 2-31(b) is 400 ps later, or approximately 350 ps after the Novette pulse is completed. The Mg XII 4-3 line is dominant, with the O VII 3-2 line just visible on the shoulder of the 4-3 line. Figure 2-31(c) is taken 350 ps later and shows that the Mg XII 4-3 line and the O VII 3-2 line are of comparable strength. Figure 2-31(d) is 400 ps later, when the Mg XII 4-3 line is just discernible on the shoulder of the 3-2 line. At approximately 1.5 ns after the pulse, Fig. 2-31(e) shows the splitting of the O VII 3-2 line, discussed above.

Modeling Calculations and Interpretation. We have recently completed 2-D modeling calculations that provide a tentative explanation of the Novette experimental results—in particular, the long duration of spontaneous emission in the Mg XII 4-3 line. Absorption of $0.53\text{-}\mu\text{m}$ laser light and the resulting expansion of the magnesium foil was modeled using the LASNEX code.⁴⁹ Time histories of temperature and density were used as input to the XRASER post-processor atomic physics code.⁵⁰ Another atomic physics code, RATIME,⁵¹ was used as a check. These codes calculate ionization balance and level population in the magnesium. The hydrogenic approximation is used to calculate the relevant rates and cross sections. Finally, the results from the XRASER code are used as

Fig. 2-32. On-axis view of an expanding magnesium-foil target. Region 1 is heated by the optical laser; adjacent regions are heated indirectly by electron conduction. Formvar is not shown.

Fig. 2-33. Synthetic spectra generated from the SPECTRE code showing hydrogen-like and helium-like lines of magnesium.

input to the XRASER post-processor code SPECTRE, which produces a synthetic spectra of what would be observed by various experimental devices.

The important questions resulting from the experiments were why the spontaneous emission in the Mg XII 4-3 line persisted for longer than 500 ps after the Novette pulse and whether or not such emission affected the ability to detect gain in the line. The observed time history of

spontaneous emission was not predicted by a pre-experimental analysis of the exploding-foil target, which concentrated on the region of the foil illuminated by the Novette laser. This region is heated to a temperature of 1 to 2 keV, rapidly expands, and therefore produces a relatively small amount of spontaneous emission a few hundred ps after the pulse. The high temperature implies that the magnesium is almost fully stripped and that only a small fraction (less than 20%) of hydrogen-like magnesium remains. The rapid expansion and decreasing density reduces the rate of collisional excitation of the lower hydrogen-like levels and, hence, the amount of spontaneous emission. This hot, rapidly expanding region is labeled Region 1 in Fig. 2-32.

We have identified a region of the magnesium foil that can produce significant amounts of spontaneous emission. This portion of the foil, labeled Region 2 in Fig. 2-32, is not illuminated by the Novette pulse but is heated indirectly by electron conduction from Region 1. The temperature in Region 2 ranges from 100 eV to 1 keV. Modeling calculations show that 50% of the magnesium is hydrogen-like at a temperature of ~ 500 eV. Because Region 2 is not heated as much as Region 1, it expands much more slowly. At 500 ps after the Novette pulse, Region 2 remains at a density high enough to produce a significant amount of spontaneous

emission in the 4-3 line of Mg XII. Region 3 does not reach a temperature high enough to contribute to the spontaneous emission in the line.

The SPECTRE code allows for a comparison of the modeling results with data taken by McPIGS and TGSS-EM. Figure 2-33 is the calculated synthetic spectra showing hydrogen-like and helium-like lines of magnesium at a time ~ 400 ps after the Novette pulse. A comparison with the McPIGS data in Fig. 2-30 shows that relative strengths of the Mg XI and Mg XII lines are in reasonable agreement. Formvar is included in the hydrodynamic calculations but has not yet been post-processed by XRASER and SPECTRE; thus, the O VII and O VIII lines are not present in the synthetic spectra.

The calculated intensity of the 4-3 line of Mg XII vs time, as would be observed by TGSS-EM, is shown in Fig. 2-34 for three target lengths L . Intensity is normalized to the value after the Novette pulse is completed at 400 ps. The time history of intensity is similar for the two shorter lengths except for the increase at ~ 1.1 ns associated with induced emission in the line. Induced emission is barely discernable above the spontaneous emission for the 0.5-cm target but is clearly seen for the 1-cm target. The induced emission is dominant for the 3-cm target. This result is expected because induced emission associated with gain scales exponentially with target length.

The time history of gain for the center of Region 1 (see Fig. 2-32) is shown in Fig. 2-35. The onset of gain corresponds to the increase in calculated intensity of the 4-3 line as a result of induced emission.

The fact that the Novette laser intensity is not constant along the line focus adds another complication. The intensity at the edges of the target is approximately one-half that at the center. The dominant effect of the decrease in intensity is that the predicted gain occurs at an earlier time. Therefore, at a given time, the effective length of the lasing medium is less than the actual target length. The situation could be improved in future experiments by allowing the length of the cylindrical focus to be wider than the magnesium so

Fig. 2-34. Normalized intensity of the 4-3 line of Mg XII vs time, as would be observed by TGSS-EM.

Fig. 2-35. Calculated time history of gain in the 4-3 line of Mg XII for the center of Region 1, shown in Fig. 2-32.

that the intensity along the foil or strip is more uniform.

With respect to interpretation of the Novette experiments, the decrease in effective length of the lasing medium may help explain our results. The TGSS-EM data for a 1-cm target shows a monotonic decrease in the 4-3 line of Mg XII. (The 3-2 line of O VII obscures the 4-3 line at late times.) As shown in Fig. 2-34, this result is more consistent with a shorter-length lasing medium.

Future Designs for the Recombination X-Ray Laser. We suggest that the previous magnesium foil be replaced by a narrow strip of magnesium. The width of the strip should be slightly less than 0.01 cm, the width of the Nova two-beam line focus. The strip would remove the major late-time source of spontaneous emission (Region 2 in Fig. 2-33). The reduction of spontaneous emission should greatly improve the capability of both the McPIGS and the TGSS-EM devices to measure amplification in a low-gain system.

We also suggest that the intensity along the length of strip be kept as uniform as possible. The major effect of a change in

intensity is to change the time of gain rather than the magnitude of gain. Therefore, to achieve the maximum effective length of the lasing medium, a uniform intensity is required. Uniformity can be obtained by using the center portion of the cylindrical line focus. The loss of energy is compensated by having a more uniform intensity along the strip.

The problem on TGSS-EM of the 3-2 line of O VII obscuring the 4-3 line of Mg XII at late times cannot be corrected by modification of the target. (Removing all oxygen is not practical.) Changing the Z of the lasing medium by 1, i.e., replacing magnesium with aluminum, shifts the lasing 4-3 line to 111 Å and away from the oxygen line. Preliminary calculations on an aluminum strip predicted a 50% increase in gain compared to that for magnesium.

All schemes discussed above use collisional ionization by electrons to fully strip the lasing atom. This is achieved by illuminating the target by 0.53- μ m laser light, which deposits energy into electrons that reach a thermal temperature of 1 to 2 keV. An alternative approach is to use x rays from flashlamps to strip the lasing atom. After photoionization, compared to collisional ionization, the thermal temperature of electrons is much lower, and, in principle, recombination can occur at a higher density and lead to higher gain. A gold flashlamp has a high conversion efficiency ($\sim 50\%$) for converting 0.53- μ m light to soft x rays. The majority of soft x rays have energy < 1 keV and can ionize all but the K-shell electrons in magnesium and aluminum. Stripping to only helium-like configurations allows for the possibility of lasing in a lithium-like system. The low conversion efficiency for 0.53- μ m light to ~ 2 -keV photons may require the use of the ten-beam Nova chamber for such experiments. We have begun calculations on design of x-ray laser targets for these photoionizing schemes.

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Charged-Particle Fusion

Introduction

In 1985 we concentrated on efforts that could lead to greater cost-effectiveness of charged particles for ICF applications. We have:

- Studied target gain, as reported in the following article, "Ion-Beam Targets Using Spin-Polarized DT." Our best estimate shows that substantial reduction of input energy is possible.

- Progressed in studying the use of heavy-ion fusion beam compression⁵² to perform the additional duty of pulse shaping. See "Final Beam Compression of Heavy-Ion Fusion to Shape Beam Pulse."

- Continued our studies on the symmetry issues of direct drive. We have assessed the drive asymmetries of judicious beam placement⁵³ and beam-pulse spatial profile.⁵⁴

- Studied target implosion concepts to make use of the symmetry characteristics of temporally short, direct-drive pulses.

In studying the effects of polarized fuel on ion-beam targets, we were guided by the potential economic advantages of:

- Considering usable target gain (~ 100 to 200) at up to one-half the previously estimated beam-pulse energy.

- Relaxing the peak power requirements of driver performance.

- Lowering beam-pulse energy to allow for a possible lower-cost, heavy-ion fusion (HIF) accelerator through either lower total particle charge or kinetic energy.

- Using targets with lower yield at fixed target gain to allow for smaller fusion chambers.

In terms of the target-deposition asymmetries of our scheme for beam placement,^{53,54} numerical results have shown that a residual asymmetry of 1 to 2% can be achieved on a spherical target with 32 beamlets. Initial results on target concepts using such effects might also reduce input energy by a factor of about 2.

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Ion-Beam Targets Using Spin-Polarized DT

In this article, we discuss the use of spin-polarized DT fuel for ion-beam targets. The consequence of using such fuel for single-shell targets is a reduction in energy input by roughly a factor of 2. A large polarization fraction is needed however, and it remains to be proven that the polarized fuel can be produced in quantities sufficient for commercial ICF reactors.

The reduction of input-energy requirements of an ICF target by use of spin-polarized DT fuel follows from the improvement of nuclear reaction rates when such fuel is used.⁵⁵⁻⁵⁸ As a consequence, for the same density of DT fuel and inertial confinement time, a higher fraction of fuel undergoes thermonuclear burn. Thus, the target gain with polarized fuel at nearly-fixed energy input is increased. Alternatively, for fixed target gain, the energy input is reduced when polarized fuel is used.

To achieve an appreciable increase in fusion rate, a large fraction of fuel polarization is required. In the reaction $D + T \rightarrow \text{He}^5(3/2) \rightarrow \alpha + n$, the ratio

$$\delta = \frac{\sigma v(\text{DT})}{\sigma v(\text{unpolarized})} \quad (31)$$

is determined by

$$\delta = 1.0 + 0.5 (F_D F_T), \quad (32)$$

where F_D and F_T are polarization fractions for deuterium and tritium, respectively. Even for 50% polarization, δ is only 1.125.

For a given gain, simple analytical models^{56,59} predict a driver energy reduction between $\delta^{-2.5}$ and $\delta^{-3.0}$. For $\delta = 1.5$, energy reductions between factors of 2.8 and 3.4 are indicated. The models assume that the fuel has been assembled in the desired manner and that ignition has been achieved. Our current numerical modeling gives somewhat less improvement.

There is a trade off between driver output energy and the duration of the

temporal power profile. Reduction in the driver energy implies using a smaller target; however, the duration of the power profile is roughly proportional to the target size. Smaller targets require faster rise times and more stringent timing requirements. Smaller targets may also require more preheat shielding and may be more difficult to ignite.

Figure 2-36 accounts for these provisions and shows our best estimate of gain for single-shell targets. The parameters are similar to those in previous discussions⁶⁰ of ion-beam target gains—that is, the curves are parameterized by input energy and by $r^{3/2}R$, where r is target radius in cm and R is beam ion range in g/cm^2 . The top curve in Fig. 2-36 is reproduced from Ref. 61. Because of some differences in constituent materials, the curves for unpolarized DT do not correspond exactly to the curves given in Ref. 61. More studies are needed to guarantee sufficient quantities of such polarized fuel for commercial ICF applications.

Fig. 2-36. Gain of single-shell targets with polarized DT fuel compared to gain of targets of similar materials and masses but with unpolarized DT fuel.

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Fig. 2-37. (a) Initial velocity tilt $v_z(z)$ of a 2-1/2-D run with the LLNL CONDOR particle code using a near-square initial longitudinal pulse shape. (b) Line-charge density $\lambda(z)$ when compression by a factor of 5 is attained.

Fig. 2-38. One-half the initial velocity tilt $v_z(z)$ shown in Fig. 2-37(h) results in different pulse steepness.

Final Beam Compression of Heavy-Ion Fusion to Shape Beam Pulse

In an earlier report⁶² we showed that for 1-D and 2.5-D simulations, a heavy-ion fusion beam pulse can be significantly compressed in length either spatially or temporally (as seen by target). In this article, we summarize the results of simulations showing that the final beam shape can also be significantly varied. We provide samples of shapes that are available and suggest that a wide variety of final pulse shapes can be generated from the perspective of physics, although the shapes may be subject to engineering constraints.

As discussed in the 1984 Laser Annual, the beam compresses in length because of variation in velocity along the beam bunch. Accelerator physicists refer to the process by which the beam tail catches up

with the beam head as velocity tilt in the (z, v_z) phase space.⁶²

In our earlier case, the initial longitudinal shape of the bunch is a parabola in line-charge density $\lambda(z)$. The outcome is convenient because an accurate analytic solution exists for 1-D problems with no longitudinal temperature variation. The initial velocity tilt is such that the velocity of particles in the ion beam bunch is approximately a decreasing linear function of spatial distance z from beam tail to beam head.

There is no reason to restrict the initial shape of the charge bunch to a parabola. In fact, its shape after leaving the accelerator and before bunch compression is more likely to be a square pulse (perhaps with slightly rounded ends).

Figures 2-37 and 2-38 show the time evolution in beam compression of two pulses with nearly square initial shape, but with different amounts of initial velocity tilts. The amount of tilt or dv_z/dz is directly related to the steepness of the final pulse shape $\lambda(z)$. The gridding and beam parameters are similar to those in Ref. 62.

The most obvious variation of Figs. 2-37 and 2-38 is to give the beam head a ramp that is less steep than that for the beam tail. This is done by imposing an initial tilt, shown in Fig. 2-39(a), with two different slopes but with continuous $v_z(z)$. The initial square pulse evolves to the pulse shape of Fig. 2-39(b) at the time that line-charge density (or power) is increased by a factor of 5. Similar techniques might reproduce two ramped pulses, perhaps even approaching the pulse shapes needed by some targets, if we used multiple pulses of

Fig. 2-39. Combining the effects in Figs. 2-37 and 2-38 results in different ramps in the front and rear of the pulse.

Fig. 2-40. Pulse shape at a 5-fold increase of peak power when 3-component initial tilt is chosen.

Fig. 2-41. Additional example of a shaped pulse when 3-component initial tilt is chosen.

this type. (See the illustration in Ref. 63, noting in particular how the target sees the pulse reversed because the beam head reaches the target before the beam tail.)

We can obtain more interesting pulse shapes with three or more linear (or

near-linear) components to the initial tilt. By varying the differential acceleration of the particles at different distances z along the pulse, the control over the initial $v_z(z)$ is responsible for bunch compression with pulse shaping. The effect of such control is

best seen by Figs. 2-40 and 2-41. These figures show two final pulse shapes attained at 5-fold power amplification, given the accompanying initial velocity tilt of the beam from head to tail shown in Figs. 2-40(a) and 2-41(a).

Although our summary does not elucidate all the physics or all possible shapes of final pulses, many pulse shapes are technically possible. More complicated shapes require more linear components to $v_z(z)$. Some engineering limits may, however, be imposed on this process.

When a desired pulse shape and compression factor is reached, our present scenario requires some action of time-dependent final focus or induction accelerator pulses to remove residual velocity tilts that are still highly ordered but somewhat compressed in length. [Because similar pulses generated the initial tilt $v_z(z)$, we might expect that they could also be used to remove at least some of the residual tilts.] Further studies of the electrical engineering of this process are needed. We are also looking at alternative physical processes for pulse shaping.

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Plasma Physics

Introduction

The goal of the ICF Plasma Group is to develop our understanding of plasma processes needed to optimally use present and future experimental facilities for ICF and other applications. A substantial part of our effort is devoted to collaboration with other parts of the ICF program on the design and interpretation of experiments on the coupling physics. These results are reported in other sections of the *Laser Annual*. The remainder of our effort is approximately equally divided between plasma theory and simulation, including the development of improved computational models. Our progress in these areas is reported in the following articles.

We have developed a new model that significantly improves our description of electron energy transport in plasmas. This

fluid-like description self-consistently accounts for departures of the isotropic part of the electron-velocity distribution from a Maxwell-Boltzmann distribution. The resulting heat flow is both nonlocal and naturally flux-limited. Earlier results of time-consuming Fokker-Planck calculations are recovered in detail. The new transport model, suitable for implementation in 1-D target-design codes, is discussed in the next article, "Nonlocal Heat Transport by Near-Maxwell-Boltzmann Electron Distributions."

Our understanding of laser-driven instabilities in long-scalelength plasmas continued to advance. As detailed in Ref. 64, many important observations of Raman scattering of laser light are in accord with our calculations. The features include an increase in the scattering with density scalelength, hot-electron generation⁶⁵ correlated with this scattering, and collisional suppression⁶⁶ of the Raman instability in high-Z plasmas irradiated with short-wavelength light. However, several puzzles remain unsolved, including the detailed frequency and angular distribution of the Raman scattered light and the quantitative noise levels and threshold intensities.

We continue to study the $2\omega_{pe}$ instability, a potential source of superhot electrons that appears to be operative at a lower level than is understood on the basis of current experiments. Recent work on the $2\omega_{pe}$ instability and on Raman scattering in our long-scalelength experiments (thin, foil targets) is reported in "ZOHAR Simulations of Two-Plasmon Decay," "Parametric Instabilities in Inhomogeneous Media," and "Simulations of Raman Backscattering in a Parabolic Density Profile." We also seek an improved understanding of Brillouin scattering because future target implosions will be quite sensitive to the symmetry of the drive and, hence, the location of the energy deposition. Recent work in this area is reported in "Nonlinear Features of Stimulated Brillouin Scattering."

The calculations reported in "Monte Carlo Atomic Physics Calculations" are improving our understanding of x-ray conversion in high-Z plasmas. Our work this year also included the simulations presented in "Electromagnetic Effects in

Relativistic Electron-Beam-Plasma Interactions." One area of application is to laser particle accelerators.

Our 2.5-dimensional, electromagnetic, direct-implicit particle code is now operational and running a variety of problems. Present operation shows that we can significantly expand the parameter regime of problems to which our particle codes can be applied. We now routinely run problems with rapidly varying, strongly inhomogeneous profiles. The code has run stably and productively with a time step $25 \times$ larger than the corresponding explicit code. Our work in this and other areas of advanced code development is reported in "Implicit Particle Code Development," "Darwin Field Model for Charged-Particle Source Studies," and "Electron-Electron Collision Operator in Low-Z Transport." Part of our advanced code development was supported this year by the internal research program of the Laboratory.

Eight invited papers were presented this year by members of the Plasma Group. In addition, we collaborated on the interpretation of laser-plasma experiments with other laboratories, including Los Alamos, KMS Fusion, Inc., Ecole Polytechnique, NRC Canada, the University of Rochester, and UCLA.

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Nonlocal Heat Transport by Near-Maxwell-Boltzmann Electron Distributions

The need for an efficient and self-consistent treatment of high-flux electron transport in laser irradiated plasmas has been acute for some time. Intense research has identified and resolved important issues in the Coulomb collisional theory.⁶⁷⁻⁷⁰

We now understand that the failure of classical transport theory occurs when the mean free path of the high-energy, heat-carrying electrons is not much smaller than the (temperature) scalelength of interest. This relation implies that the transport fluxes are not locally determined. Indeed, idealized systems of interest have

been successfully described by a fully multigroup (in energy) diffusion (in space) treatment.⁷⁰ This reduced Fokker-Planck treatment accounts for the nonthermal distribution induced at high energy by the transport itself, which, in turn, acts to alter the transport from classical. Heat fluxes associated with steep temperature gradients are reduced from classical, while there is preheating ahead of a temperature front that exceeds classical values. This understanding has been gained by analysis of state-of-the-art numerical calculations that are still too inefficient for application to many problems.

In practice, ad hoc fluid schemes have been employed to treat high-flux electron transport. The most successful scheme limits the heat flux to a fraction of its so-called free-streaming value $Q_{\parallel} = F_{\parallel} n (T/m)^{1/2} T$. Recently a fluid scheme possessing nonlocal phenomenology has been employed and normalized against certain first-principles calculations.⁷¹

In this article, we report numerical calculations with a new, nonlocal fluid treatment of Coulomb collisional transport that self-consistently accounts for the non-Maxwell-Boltzmann, high-energy electrons occurring in laser plasmas. The source of the nonthermal electrons is the spatial transport of thermal electrons with a range that is not short compared to the temperature scalelength. The associated Fokker-Planck calculation has been carried out analytically and permits the first-principles formulation of a nonlocal description (in space) of the evolution of the plasma temperature.

Our theory requires that the number and energy of the nonthermal electrons be small. Their number and energy fluxes are not restricted and may dominate the total fluxes. Typical results are heat fluxes significantly different from those of either classical or flux-limited fluid transport. Indeed, we have recovered in detail the results of previous calculations of fully multigroup diffusion laser absorption and heat transport. Our transport scheme is suitable for implementation in 1-D laser-plasma simulation codes and is expected to be efficient enough for general use.

We give an approximate solution of the reduced Fokker-Planck equation that governs the electron distribution in position and energy⁷⁰ as follows:

$$\frac{\partial f_0}{\partial t} - \frac{\partial e\phi}{\partial t} \frac{\partial f_0}{\partial \epsilon} - \frac{1}{v} \frac{\partial}{\partial x} \frac{v^2 \lambda_{90}}{6} \frac{\partial f_0}{\partial x} = \frac{df_0}{dt_{ee}} \quad (33)$$

Here $mv^2/2 = \epsilon + e\phi(x, t)$ defines the usual total energy ϵ , and the electrostatic field is $E = -\partial\phi/\partial x$. The scattering mean free path is $\lambda_{90} = (mv_e)^2/2\pi e^4 \times (Z^2 N \ln \Lambda_{ei} + n \ln \Lambda_{ee})$. Equation (33) neglects the hydrodynamics of background ions.

The electron number and energy

$$\left(n, \frac{3}{2} nT\right) = 4\pi \int_{v\phi}^{\infty} \frac{d\epsilon v}{m} \left(1, \frac{mv^2}{2}\right) f_0 \quad (34)$$

obey the associated energy moments of Eq. (33):

$$\frac{\partial n}{\partial t} + \frac{\partial \Gamma}{\partial x} = 0 \quad (35)$$

and

$$\frac{\partial}{\partial t} \frac{3}{2} nT + \frac{\partial Q}{\partial x} + e\Gamma E = 0. \quad (36)$$

Here the number and energy fluxes Γ and Q , are

$$\begin{aligned} \Gamma, Q = & -4\pi \int_{v\phi}^{\infty} \frac{d\epsilon}{m} \left(1, \frac{mv^2}{2}\right) \\ & \times \frac{v^2 \lambda_{90}}{6} \frac{\partial f_0}{\partial x}. \end{aligned} \quad (37)$$

Electron-electron (self) collisions d/dt_{ee} conserve both number and energy of the species. The electric current is $J = -e\Gamma$.

The electrostatic potential is determined by the requirement of quasineutrality, $n = ZN$. In the present application, this becomes the vanishing of the number flux $\Gamma(v\phi) = 0$.

If the self-collision term df_0/dt_{ee} on the right-hand side of Eq. (33) dominates, then the distribution takes its local-

thermodynamic-equilibrium form, the Maxwell-Boltzmann (MB), as follows:

$$f_0 = f_{MB} = \frac{n}{(2\pi T/m)^{3/2}} e^{-\frac{\epsilon + e\phi}{T}}. \quad (38)$$

Employing f_{MB} in Eqs. (35) through (37) yields the classical transport treatment wherein $Q = -\partial T/\partial x$; this closes the fluid moment equations in n , T , and $e\phi$.

The transport itself upsets the MB distribution whenever the scalelength of interest (L) is not much longer than the stopping length, $\lambda_s = (2\lambda_{90} \lambda_e/3)^{1/2}$. This may be seen by estimating $df_0/dt_{ee} \sim (v/4\lambda_e) f_0$ and comparing it with the transport term $\sim (v\lambda_{90}/6L^2) f_0$. Here $\lambda_e = (mv^2)^2/2\pi e^4 n \ln \Lambda_{ee}$ is the energy-loss mean free path due to electron-electron collisions. Because of the strong energy dependence of λ_e , the higher-energy electrons are the first to be altered by the transport. Inspection of Eq. (37) reveals that these higher-energy electrons are also the ones that carry the fluxes. Thus a non-thermal distribution at high energy must, in turn, be accounted for in calculating the transport itself.

The preceding discussion motivates our solution of Eq. (33). We write $f_0 \approx f_{MB} + \delta f$ and solve for δf at high energy, where Eq. (33) becomes

$$\begin{aligned} & -\frac{1}{v} \frac{\partial}{\partial x} \frac{v^2 \lambda_{90}}{6} \frac{\partial f_{MB}}{\partial x} - \frac{1}{v} \frac{\partial}{\partial x} \frac{v^2 \lambda_{90}}{6} \frac{\partial \delta f}{\partial x} \\ & = \frac{2mv^2}{\lambda_e} \frac{\partial \delta f}{\partial \epsilon}. \end{aligned} \quad (39)$$

Here f_{MB} is annihilated by d/dt_{ee} . On the right-hand side, we neglect the thermalization of high-energy electrons on the bulk as small (by $T\partial \ln \delta f/\partial \epsilon$) compared with the energy loss. The temporal variation of the distribution and potential is assumed to be slow compared with the transport and energy loss retained for analysis.

Equation (39) makes explicit that the source (possibly negative) of nonthermal electrons δf is the spatial transport of thermal electrons f_{MB} . Because the temperature is characteristic of the thermal electrons, we see that the nonthermals emerge when the temperature scalelength is not much longer than the stopping length.

We solve Eq. (39) in the limit $\epsilon > -e\phi$, so that $m v^2/2 \sim \epsilon$ and find

$$f_0 = \int d\xi' \int_{\epsilon}^{\infty} d\epsilon' \frac{e^{-\frac{(\xi' - \xi)^2}{\epsilon'^4 - \epsilon^4}}}{[\pi(\epsilon'^4 - \epsilon^4)]^{1/2}} \times \frac{f_{MB}(\xi', \epsilon')}{T(\xi')}. \quad (40)$$

In solving the diffusion Eq. (39), ϵ^4 is the time-like variable, and $d\xi = dx/\tilde{\lambda}_s$ defines the space-like variable. We use $\tilde{\lambda} = \lambda/(mv^2)^{1/2}$. We have not taken account of spatial boundary conditions here but leave them for introduction later. (We must give the distribution or its spatial gradient on boundaries.) We have also passed a factor of $Z^{1/2}$ through $\partial/\partial x$ in defining ξ , and thus require that the ionization state vary more slowly than the temperature.

Equation (40) makes clear that nonthermal electrons at ξ and ϵ have come from all ξ' and all $\epsilon' > \epsilon$ by stopping transport. We observe that f_0 of Eq. (40) is determined by only the density and temperature of the plasma via f_{MB} . The total number and energy, δn and $\delta 3nT/2$, respectively, of nonthermals δf are ignored as small in the moment Eqs. (35) and (36) so that the density and temperature of the plasma become those of the bulk MB distribution. Thus, substitution of f_0 from Eq. (40) into Eq. (37) closes the moment equations and naturally yields a nonlocal fluid treatment of the transport.

In Eq. (37) we have energy integrals of $\partial f_0/\partial x = (1/\tilde{\lambda}_s) \partial f_0/\partial \xi$. This differentiation acts on the Gaussian kernel in Eq. (40) for f_0 . We replace $\partial/\partial \xi$ by $-\partial/\partial \xi'$ and integrate by parts in ξ' to cast the spatial derivative onto f_{MB}/T . Once again, we do not yet take up the spatial boundary conditions. The electric field enters explicitly upon differentiation of the exponential factor of f_{MB} :

$$\frac{\partial}{\partial \xi'} \frac{f_{MB}}{T} = \left(\frac{1}{n} \frac{\partial n}{\partial \xi'} - \frac{5}{2} \frac{1}{T} \frac{\partial T}{\partial \xi'} \right) \frac{f_{MB}}{T} + \frac{\epsilon + e\phi}{T} \frac{1}{T} \frac{\partial T}{\partial \xi'} - \frac{1}{T} \frac{\partial e\phi}{\partial \xi'} \frac{f_{MB}}{T}. \quad (41)$$

It is convenient to split the potential into local (l) and nonlocal (nl) parts, $e\phi = e\phi_l + e\phi_{nl}$, where

$$\frac{\partial e\phi_l}{\partial \xi'} = \frac{T}{n} \frac{\partial n}{\partial \xi'} - \frac{5}{2} \frac{\partial T}{\partial \xi'}. \quad (42)$$

Finally $\epsilon > -e\phi$ is exploited by letting $-e\phi \rightarrow 0$ everywhere except under differentiation by ξ' . Then Eq. (37) becomes:

$$\Gamma, Q = -\frac{8\pi}{m^2} \frac{\tilde{\lambda}_s}{\tilde{\lambda}_t} \int d\xi' \int_0^{\infty} d\epsilon \epsilon^3 (1, \epsilon) \times \int_{\epsilon}^{\infty} d\epsilon' \frac{e^{-\frac{(\xi' - \xi)^2}{\epsilon'^4 - \epsilon^4}}}{[\pi(\epsilon'^4 - \epsilon^4)]^{1/2}} \times \left(\frac{\epsilon'}{T} \frac{1}{T} \frac{\partial T}{\partial \xi'} - \frac{1}{T} \frac{\partial e\phi_{nl}}{\partial \xi'} \right) \frac{f_M}{T}. \quad (43)$$

Here $f_M = f_{MB}(e\phi \rightarrow 0)$.

The double energy integrals have been computed to obtain

$$\Gamma, Q = \frac{(\tilde{\lambda}_{90}/\tilde{\lambda}_t)^{1/2}}{4\pi(3m)^{1/2}} \int dx' n T^{-1/2} (1, T) \times \left\{ \frac{\partial T}{\partial x'} [K(\theta), K(\theta)] - \frac{\partial e\phi_{nl}}{\partial x'} [J(\theta), L(\theta)] \right\}. \quad (44)$$

Here we have made the transformation $d\xi' \partial/\partial \xi' = dx' \partial/\partial x'$ and introduced

$$\theta = \left| \frac{\int_{x'}^x \frac{dx''}{\tilde{\lambda}_s}}{T^2(x')} \right|, \quad (45)$$

which is the number of stopping lengths from the source point x' to the field point x at an energy equal to the source temperature. The nonlocal transport propagators I, J, K , and L are functions of θ only.

We have computed and tabulated the nonlocal transport propagators for rapid use in numerical calculations. Figure 2-42 is a plot of the four functions vs θ .

The propagators $P = (I, J, K, L)$ have the following form:

$$P(\theta) = \theta^{2+2\beta} \int_0^\infty dy y^\beta e^{-\theta^{1/2} y^{1/4}} \times \int_0^1 dy' \frac{y'^\alpha}{(1-y')^{1/2}} e^{\frac{y'}{y(1-y')}} \quad (46)$$

For small θ the propagators are linear in θ , taking the form $P(\theta) \rightarrow P(0) + P'(0)\theta$ as $\theta \rightarrow 0$; for large θ they decay exponentially such that

$$P(\theta) \rightarrow 32 \left(\frac{2\pi}{5} \right)^{1/2} A \theta^\gamma e^{-\left(\frac{5}{4^{1/5}} \right) \theta^{2/5}} \quad (47)$$

Fig. 2-42. Transport propagators for non-local calculation of fluxes decrease with number of stopping lengths between the source and field points.

The constants $\alpha, \beta, P(0), P'(0), A, \gamma$, and the propagator normalization integrals are shown in Table 2-5. The range of the propagators is $\theta \sim \text{few}$; the propagators decay exponentially as $\theta^{2/5}$ because of the

energy-dependent stopping and moments taken in the fluid treatment.

Two limiting temperature profiles are of particular interest and illuminate the physics that is in effect. We first consider the limit of gradients much longer than the stopping length. In this limit, the non-local heat flux reduces to the classical value because only x' near x contributes in the integrals of Eq. (44). We may let $dx' \rightarrow T^2 \lambda_s d\theta$ so that

$$\Gamma, Q \rightarrow \frac{1}{4\pi^{1/2}} \frac{2}{3} n \left(\frac{T}{m} \right)^{1/2} \hat{\lambda}_{n0} T^2 \left(\frac{1}{T}, 1 \right) \times \left[\frac{\partial T}{\partial x} \int_{-\infty}^{\infty} d\theta (I, K) - \frac{\partial e\phi}{\partial x} n l \int_{-\infty}^{\infty} d\theta (J, L) \right] \quad (48)$$

Taking $e\phi_{nl}$ from $\Gamma = 0$ yields the classical results:

$$\begin{aligned} \frac{\partial e\phi}{\partial x} &= \frac{\partial e\phi_2}{\partial x} + \frac{\partial e\phi_{nl}}{\partial x} \\ &= \frac{T}{n} \frac{\partial n}{\partial x} + \frac{5}{2} \frac{\partial T}{\partial x} \end{aligned} \quad (49)$$

and

$$Q = - \frac{64}{(2\pi)^{1/2}} n \left(\frac{T}{m} \right)^{1/2} \lambda_{mlp} \frac{\partial T}{\partial x} \quad (50)$$

In Eq. (50) $\lambda_{mlp} = T^2 \hat{\lambda}_{90}$.

Our nonlocal heat flux is naturally self-limiting. The integral of the delta function that arises from a step in temperature is finite. For a temperature step from hot (H) to cold (C) over a distance much shorter

Table 2-5. Properties of the propagators $P = (I, J, K, L)$.

Propagator	Constants				Normalization Integral	
	α	β	$P(0)$	$P'(0)$	A	γ
I	0	0	48	$-8\pi^{1/2}$	1	1
J	0	$-1/4$	16	$-8\pi^{1/2}$	$4^{-1/5}$	$3/5$
K	$1/4$	$1/4$	$24\pi^{1/2}\Gamma(5/4)/\Gamma(7/4)$	$-16\pi^{1/2}$	$4^{2/5}\Gamma(5/4)$	$13/10$
L	$1/4$	0	$96\pi^{1/2}\Gamma(5/4)/\Gamma(7/4)$	$-8\pi^{1/2}$	$4^{1/5}\Gamma(5/4)$	$9/10$

than a stopping length, we find

$$\Delta e\phi_{ni} \rightarrow 3\Delta T \quad (51)$$

and

$$Q \rightarrow 1.285 \left(\frac{\ln \Lambda_{ee}}{Z \ln \Lambda_{ei} + \ln \Lambda_{ee}} \right)^{1/2} \times \frac{n}{m^{1/2}} (T_i^{3/2} - T_e^{3/2}). \quad (52)$$

Note that the actual temperature within the transition interval does not enter. The limit of the nonlocal treatment is, itself, nonlocal. The flux-limit prescription does not exhibit this nonlocal phenomenology;

it is nonphysical in that it is exactly pointwise local; recall $Q_{ii} = F_{ii}(T/m)^{1/2}T$. (When the interval becomes much shorter than the scattering mean free path, the collisionless limit obtains.)

We have solved the heat transport Eq. (36) numerically and compared our results with those of previous direct solutions⁷⁰ of Eq. (33). Figure 2-43 shows some results of illuminating a stationary plasma with a laser beam of constant intensity. To this end, we have introduced an energy source κ on the right-hand side of Eq. (36). (The laser intensity obeys $dl/dx = -\kappa l$; the laser opacity is $\kappa = K(l, n, T)\kappa_c$. Here K is the reduction factor given by Langdon,⁷² and κ_c is the classical opacity.) The plasma is initially

Fig. 2-43. Nonlocal transport treatment gives significantly different results from classical in a calculation of a laser-heated plasma. Classical absorption $K = 1.0$.

Fig. 2-44. Previous results for fully multigroup diffusion transport⁷⁰ are recovered by the nonlocal treatment when nonlinear reduction of laser opacity is introduced. $K \sim 0.7$.

isothermal at $T = 100 \text{ eV}$. The calculation was terminated when the heat flux into the overdense plasma nearly equaled the absorbed laser flux so that the coronal plasma had reached a quasisteady condition.

A comparison of Figs. 2-43(c) and 2-43(d) shows that the signatures of the nonlocal transport are:

1. A hot absorption layer.
2. A steep temperature gradient separating the laser-heated plasma and the conduction-heated plasma.
3. A nonisothermal corona.
4. A preheat foot on the temperature profile penetrating into the colder, denser plasma.

The reduced absorption in Fig. 2-43(a) is associated with signature 1 because $\kappa_c \propto T^{-3/2}$. This feature and signature 2 are typical of strongly flux-limited simulations required to reproduce results of laboratory experiments. The nonisothermal corona results from nonlocal transport

from there toward the colder, denser plasma. This physics is also responsible for the preheat foot. The bottling up of absorbed laser energy is due to the stopping of electrons as they transport toward the colder, denser plasma.

Figure 2-44 shows that including nonlinear reduction of the laser opacity⁷² brings the absorbed energy, heat flux, and temperature into agreement with the previous calculations. Indeed, in Fig. 2-45 we show that at high energy the underlying nonlocally determined distribution function, f_0 of Eq. (44), is in agreement as well. In our transport treatment, we have neglected the effect of the non-Maxwell-Boltzmann distribution induced at low energy by laser heating.⁷² The heating is communicated to the high-energy transporting electrons via the bulk temperature T ; self collisions mediate between laser heating at low energy and transport at high energy.

Fig. 2-45. Nonlocally calculated electron-energy distribution of Eq. (40) compares favorably with previous results for fully multigroup diffusion.⁷⁰

The effect of the nonlocally determined electric field is shown in Fig. 2-43(d). The field acts to reduce the heat flow from that due to the temperature gradient. The reduction is substantial and is not constant: it can be either larger or smaller than the classical limit. In particular, in the cold, preheat region only a relatively small electric field is required to establish the return current of dense, cold electrons required to obtain zero total current against the small number of high-energy, nonlocal, heat-transporting electrons. Thus only modest reduction of the nonlocal heat flux due to the temperature gradient occurs there.

Figure 2-46 shows that neglect of thermalization of high-energy electrons on the bulk electrons is justified. Only modest errors of our nonlocal treatment are indicated.

Energy conservation in the laser-heated plasma calculations requires that the heat

flux vanish at the vacuum boundary. This implies the condition $\partial f_0 / \partial x = 0$ at the boundary; this is straightforward but tedious to implement as an (infinite) series of images. Provided the entire system is many stopping lengths in extent, it is sufficient to include a single-image plasma on the low-density side of the fluid calculation, as shown in Fig. 2-47. Extending the Γ and Q integrations over the image plasma ensures that $\Gamma, Q = 0$ is obtained correctly at the low-density boundary.

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ZOHAR Simulations of Two-Plasmon Decay

We have extended our two-dimensional kinetic simulations of two-plasmon ($2\omega_{pe}$) decay to longer durations and to include

Fig. 2-16. Thermalization of high-energy electrons on the bulk is small by $T \partial \ln \delta f / \partial \epsilon$ compared with slowing down. Only modest quantitative errors are indicated by neglect of thermalization in the nonlocal transport treatment.

Fig. 2-47. Single image on the low-side of the laser-heated-plasma calculation is included in the non-local flux calculations to ensure that no heat flows into the vacuum.

other parameters using new diagnostics. We have checked nonlinear saturation mechanisms leading to weak heating in the underdense plasma commonly observed in experiments. These extensions using the ZOHAR code help to clarify the role of ion response in the saturation.

Our choice of parameters was motivated by recent experiments with CH foils on Novette⁷³ and Nova. The typical electron temperature was 2 keV, and illumination intensities corresponded to $I\lambda^2 \approx 10^{15}$ W/cm², where λ is the wavelength in μm . The spatial scalelengths ($\sim 200c/\omega_0$) were chosen to place the instability only slightly above its density gradient threshold, to bring the instability absorption into closer agreement with experiment, and as a shortcut to the steeper profile that evolves in the simulation.

The laser is polarized in the plane of the simulation, and the density range is $0.19n_c$ to $0.29n_c$. The plasma is initialized as a planar self-similar flow, shifted in velocity such that the $0.25n_c$ surface remains at a fixed position. Plasma flows into the simulation region at the high-density end with speed $< c_s$, and out at the low-density end with speed $> c_s$. Due to this velocity

gradient and the small values of $v_0^2/v_i^2 \propto I\lambda^2/T_e$, Brillouin backscatter is below threshold except as noted.

The role of the ion response in saturation and long-term evolution of $2\omega_{pe}$ is emphasized by the results of recent simulations. Saturation occurs first by damping due to ion fluctuations driven by $2\omega_{pe}$ waves, in the manner we discussed previously,^{74,75} and by electron trapping and Landau damping on the heated electrons.⁷⁶ If the latter mechanisms were dominant, then simulations with ion motion frozen would be similar to simulations with normal ion motion. We chose a mass ratio of $m_i/Zm_e = 3600$ for realistic comparison. Other than the mass ratio, there are no differences in ZOHAR code setup between the two runs.

The runs for fixed and mobile ions are most similar in the initial burst of $2\omega_{pe}$. In both runs, the absorption is largest and the electron spectrum is most energetic in this first burst. Even here, however, the results differ considerably. Furthermore, the initial burst is not representative of long-term behavior. Figure 2-48 compares the $2\omega_{pe}$ wave energies. It is clear that ion response to $2\omega_{pe}$ wave pressure is dominant even in the initial saturation. Ion-density fluctuations of $\sim 20\%$ are observed, whereas the density profile changes negligibly. Thus, it is the fluctuations that arrest the $2\omega_{pe}$ instability.

The later evolution is more important from an experimental point of view. With $m_i/Zm_e = 3600$, the $2\omega_{pe}$ wave energy remained near the level shown in Fig. 2-48 for the duration of the run (to $t = 12\,000 \omega_0^{-1}$). Absorption became fairly constant at 8%, with $T_{hot} \approx 40$ keV. These results are consistent with those for CH-foil experiments on Novette suggesting that, while

the quarter-critical surface is present, the $2\omega_{pe}$ instability contributes to the observed hot electron production. (In Ref. 73, the dominant instability absorption is attributed to Raman at $\sim 0.1n_c$.) In addition, our ZOHAR results for instability absorption and T_{hot} are consistent with the highest intensity ($\sim 10^{16}$ W/cm²) point in the Novette gold-disk irradiations.⁷⁷

The results for these and other simulations shown in Table 2-6 emphasize the role of ion response in the long-term behavior of the instability. (In short-wavelength experiments,^{78,79} the observed correlation between $3/2 \omega_0$ emission and Brillouin scatter also indicates such a connection.) In long-wavelength experiments,⁸⁰ the same details of ion response and instability saturation are directly observed.

Fig. 2-48. In simulations differing only in ion mass, plasmon energy with fixed ions is more than twice that with $m_i/Zm_e = 3600$.

Run	$I\lambda^2$ (W/cm ²)	m_i/Zm_e	T_i/ZT_e	Absorption (%)	Brillouin backscatter (%)	T_{hot} (keV)
1	1×10^{15}	3600	4.3	8.0	~ 0	40
2	1×10^{15}	∞	—	27.0	—	$> 100^b$
3	2×10^{15}	100	0 ^c	0.2	< 14	—
4	2×10^{15}	3600	0 ^c	4.3	< 10	—

^a Raman backscatter is negligible.

^b NonMaxwellian: peaks at 100 keV and extends past 500 keV.

^c Initially $T_i = 0$ but increases to $< ZT_e/100$ due to collisions.

Table 2-6. Summary of parameters for computer data.^a

With a shallower initial density gradient, the absorption would be much higher, at least until the self-steepening evolved. The shallower the initial gradient, the longer the time before steepening can reduce instability absorption. The improved agreement of absorption and T_{hot} with the steepened profile suggests that steepening may play a role in the experiment.

An intriguing feature is evident near the time that the mobile-ion run was suspended. As the density profile evolves under the action of $2\omega_{\text{pe}}$ wave pressure and the flow, a density depression forms that is localized both in the flow direction and transverse to it. The $2\omega_{\text{pe}}$ waves are somewhat concentrated at this depression. Some observers have speculated that this run may be showing the early stages of Langmuir wave collapse.

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Parametric Instabilities in Inhomogeneous Media

Parametric processes, such as the two-plasmon decay instability, stimulated Brillouin scattering, and stimulated Raman scattering have long been recognized as significant mechanisms in the interaction of high-intensity lasers with matter. The simplest model for these processes that includes the two primary stabilization mechanisms, namely damping on the product waves and background gradients in the plasma, was introduced by Rosenbluth.⁸¹ The model consists of the following pair of coupled equations for the wave action amplitudes a_i ($i = 1, 2$):

$$\left(\frac{\partial}{\partial t} + V_1 \frac{\partial}{\partial x} + \nu_1\right) a_1 = \gamma_0 a_2 e^{i\phi(x)} \quad (53)$$

and

$$\left(\frac{\partial}{\partial t} + V_2 \frac{\partial}{\partial x} + \nu_2\right) a_2 = \gamma_0 a_1 e^{-i\phi(x)}, \quad (54)$$

where the group velocities V_i and damping rates of the modes ν_i are treated as constants, γ_0 is the homogeneous growth

rate, and $\phi(x)$ is the phase mismatch arising from the inhomogeneity.

Spatial variation of the k vectors of the coupled modes is defined by

$$\phi(x) = \int^x [k_0(x') - k_1(x') - k_2(x')] dx'. \quad (55)$$

In this article we extend previous analyses of this model and consider particular cases in which the matching point, where ϕ' vanishes, lies at or near an extremum of the plasma profile. Our extension is of particular relevance in the interpretation of exploding-foil experiments.

The generic case of a single matching point x_0 , where $\phi'(x_0) = 0$, but $\phi''(x_0) \equiv \kappa'$ does not vanish, was analyzed in Rosenbluth's original paper⁸¹ and is typically applicable when the plasma profiles are monotone and both the coupled waves are far from their turning points. Rosenbluth found that for both forward scattering ($V_1 V_2 > 0$) and backscattering ($V_1 V_2 < 0$), only convective instability was supported and that growing eigenmode solutions with outgoing boundary conditions were not permitted. By outgoing wave conditions, we mean that the amplitude of the incoming waves asymptotically vanishes; that is, $a_i \rightarrow 0$ as $x \rightarrow -\infty \text{sign}(V_i)$.

Rosenbluth derived the corresponding convective gain coefficient C_G , the definition of which we now make precise. Without loss of generality, we consider the gain coefficient for wave 1 and take the group velocity to be positive. If wave 1 were undamped, we could define the intensity gain as the ratio $C_G = \lim_{x \rightarrow \infty} |a_1(x)|^2 / |a_1(-x)|^2$, with $|a_2(in)|^2 = 0$. If wave 1 is weakly damped, this definition must be modified to factor out the damping. A suitable definition is then

$$C_G = \lim_{L \rightarrow \infty} \exp\left(\frac{\nu_1 L}{V_1}\right) \frac{|a_1(L/2)|^2}{|a_1(-L/2)|^2}. \quad (56)$$

For strong damping, such that $\nu_1 > \gamma_0(V_1/V_2)^{1/2}$, no net spatial growth can occur; therefore, this possibility can be

eliminated as uninteresting. In the generic case, Rosenbluth found the convective gain to be given by

$$C_G = \exp \frac{2\pi\gamma_0^2}{\kappa' V_1 V_2} \quad (57)$$

and to be independent of the damping ν_2 on wave 2.

For sufficiently strong damping on wave 2 such that $\nu_2 > \gamma_0(V_2/V_1)^{1/2}$, a general expression can be readily obtained for C_G (Ref. 82). From Eqs. (53) and (54) in steady state, we obtain

$$\left(V_1 \frac{d}{dx} + 2\nu_1 \right) |a_1|^2 = \gamma_0 a_2 a_1 e^{i\phi} + \gamma_0 a_2 a_1 e^{-i\phi} \quad (58)$$

and

$$\left(V_2 \frac{d}{dx} - i\phi' V_2 + \nu_2 \right) a_2 e^{i\phi} = \gamma_0 a_1 \quad (59)$$

On neglecting the derivative term in Eq. (57), we obtain a single equation for the action density $|a_1|^2$

$$\begin{aligned} \left(V_2 \frac{d}{dx} + 2\nu_1 \right) |a_1|^2 \\ = \frac{2\gamma_0^2 \nu_2}{\nu_2^2 + V_2^2 \phi'^2} |a_1|^2 \end{aligned} \quad (60)$$

and thus

$$C_G = \frac{2\gamma_0^2}{V_1} \int_{-\infty}^{\infty} dx / [\nu_2^2 + V_2^2 \phi'(x)^2], \quad (61)$$

which yields Eq. (57) for the generic case $\phi'(x) = \kappa' x$.

However, if the damping is sufficiently strong, the resonant coupling is broadened, and the spatial variation of V_2 (and ν_2) can no longer be neglected.⁸² In this case, the gain must be computed directly from Eq. (61) and not from Eq. (57).

Johnston⁸³ has considered the gain formula of Eq. (61) in the context of an exploding-foil plasma, where the matching point x_0 may be at or near the density maximum. In this situation, $\phi'(x)$ takes

the form

$$\phi'(x) = \frac{1}{2} \kappa'' (x^2 - x_0^2). \quad (62)$$

We then obtain

$$\begin{aligned} \ln C_G = \left(\nu_0 / \kappa''^{1/3} \sqrt{V_1 V_2} \right)^{3/2} \left(\frac{\nu_0 \sqrt{V_2/V_1}}{\nu_2} \right)^{1/2} \\ \times 2\pi f \left(\frac{\kappa'' x_0^2 V_2}{2\nu_2} \right), \end{aligned} \quad (63)$$

where

$$f(p) = \left[\frac{(1 + p^2)^{1/2} + |p|}{1 + p^2} \right]^{1/2}. \quad (64)$$

The matching point x_0 is related to the difference $\Delta\omega$ from the frequency of the waves that would match at the density maximum by

$$\frac{1}{2} \kappa'' x_0^2 = \Delta\omega \left(\frac{1}{V_2} - \frac{1}{V_1} \right) \quad (65)$$

and thus p , the argument of f in Eq. (63) is given by

$$p = \frac{\Delta\omega}{\nu_2} \left(1 - \frac{V_2}{V_1} \right). \quad (66)$$

When $p \gg 1$, the two resonant regions at $x = \pm x_0$ are uncoupled and have effective linear slopes given by $\kappa' = \kappa'' x_0$. In this limit, the gain becomes independent of ν_2 and is given by

$$\ln C_G = 4\pi \gamma_0^2 / \kappa' x_0 V_1 V_2, \quad (67)$$

where $p \rightarrow \infty$, and x_0 is determined from Eq. (65).

For $p < 1$ the two resonant regions merge, and the modes amplify with a spatial gain coefficient of $2\gamma_0^2 / V_1 \nu_2$ over a region of effective size $\simeq (\nu_2 / V_2 \kappa'')^{1/2}$ giving

$$\ln C_G = 2\pi \gamma_0^2 / V_1 (\kappa'' V_2 \nu_2)^{1/2}. \quad (68)$$

In the case of a linear mismatch, the gain coefficient in Eq. (61) is derived by assuming strong damping $\nu_2 \gg \gamma_0(V_2/V_1)^{1/2}$. Then C_G is independent of ν_2 and is in

fact valid even when $v_2 \rightarrow 0$. Such is not the case for the quadratic mismatch; in addition, the possibility of absolute instability arises for the case of backscattering where $V_1 V_2 < 0$.

We are thus led to consider weak damping, where $v_2 < \gamma_0 (V_2/V_1)^{1/2}$. If $p \gg 1$ so that the resonances are distinct, Eq. (63) still applies, and we are left to consider the case where $x_0 = \Delta\omega \approx 0$. Assuming an $e^{i\omega t}$ time dependence, combining Eqs. (53) and (54), and eliminating the first derivative, we obtain

$$\left\{ \frac{d^2}{dx^2} + \frac{1}{4} \left[\phi' - \omega \left(\frac{1}{V_1} - \frac{1}{V_2} \right) - i \left(\frac{V_1}{V_1} - \frac{V_2}{V_2} \right)^2 \right] - \frac{\gamma_0^2}{V_1 V_2} + \frac{1}{2} i \phi'' \right\} \alpha = 0. \quad (69)$$

A Wentzel-Kramers-Brillouin (WKB) analysis of this equation gives the convective gain for forward scattering as

$$\begin{aligned} \ln C_G &= 2 \int_{x_1}^{x_2} dx \\ &\times \left(\gamma_0^2 / V_1 V_2 - \frac{1}{16} \kappa''^2 x^4 \right)^{1/2} \\ &= A \left[\gamma_0 / \kappa''^{1/3} (V_1 V_2)^{1/2} \right]^{3/2}, \end{aligned} \quad (70)$$

where $A = \sqrt{\pi} (-3/4) / (3/4)! \approx 6.992$, and $x_T = 2(\gamma_0 / \kappa'' \sqrt{V_1 V_2})^{1/2}$ is the turning point. We can thus summarize the results for the forward scattering gain by

$$\begin{aligned} \ln C_G &= (\gamma_0 / \kappa''^{1/3} \sqrt{V_1 V_2})^{3/2} \\ &\times \min \left[A, 2\pi \left(\gamma_0 \frac{\sqrt{V_2/V_1}}{V_2} \right)^{1/2} \right] \\ &\times 2\sqrt{2}\pi \left(\frac{\gamma_0 \sqrt{V_2/V_1}}{\Delta\omega \sqrt{V_1 V_2}} \right)^{1/2} \end{aligned} \quad (71)$$

For backscattering (where $V_1 V_2 < 0$) the two resonant regions are coupled by the counterpropagating waves and give rise to the possibility of absolute instability. A series of such modes with discrete frequency shifts occurs. The shifts correspond to the

total phase shift around the feedback loop, which is an integral multiple of 2π . Following this line of reasoning in the manner of Koch and Williams,⁸⁴ we obtain the following thresholds in the absence of damping:

$$\frac{\gamma_0^2}{\kappa''^{2/3} V_1 V_2} \approx \frac{\ln 2}{2\pi} \left(\frac{3}{2} \phi_n \right)^{1/3}, \quad (72)$$

where the resonance points and frequency shifts are given by

$$\frac{2}{3} \kappa'' x_0^3 = \phi_n \approx 2n\pi \quad (73)$$

and

$$\begin{aligned} \Delta\omega &= \frac{1}{2} \frac{|V_1 V_2| \kappa''^{1/3}}{|V_1| + |V_2|} \left(\frac{3}{2} \phi_n \right)^{2/3} \\ &- \frac{|V_1 V_2| \kappa''^{1/3}}{|V_1| + |V_2|} \text{Re}(F_n), \end{aligned} \quad (74)$$

Table 2-7 shows the thresholds and eigenfrequencies obtained from a numerical determination of the eigenvalues of the differential Eq. (69). Equation (72) is an excellent approximation for all mode numbers n , however, Eq. (73) naturally fails for $n = 0$.

Rosenbluth derived an expression for the growth rate of this instability, which we here generalize to $n = 0$ and to include the effects of damping. Dimensional analysis readily shows that the growth rate and frequency shift must be given by an expression of the following form, with the complex function F_n to be determined:

$$\begin{aligned} \omega &= -i \frac{v_1 |V_2| + v_2 |V_1|}{|V_1| + |V_2|} \\ &+ \frac{\kappa''^{1/3} |V_1 V_2|}{|V_1| + |V_2|} F_n(\Gamma), \end{aligned} \quad (75)$$

where the pump parameter Γ is

$$\Gamma = \frac{\gamma_0}{\kappa''^{1/3} |V_1 V_2|^{1/2}}. \quad (76)$$

Figure 2-49 shows the real (Re) and imaginary (Im) parts of the functions $F_n(\Gamma)$ for $n = 0, 1, 2$, and 3. This figure

Mode number, n	Numerical threshold		Analytic threshold from Eqs. (72) to (74)		Analytic threshold from Eq. (77)	
	Γ	$\text{Re}(F_n)$	Γ	$\text{Re}(F_n)$	Γ	$\text{Re}(F_n)$
0	0.386	0.691	0.000	0.000	0.397	0.794
1	0.495	2.380	0.483	2.231	0.825	1.651
2	0.548	3.633	0.542	3.541	1.160	2.320
3	0.583	4.706	0.580	4.641	1.452	2.900

Table 2-7. Instability thresholds.

 Fig. 2-49. (a) Scaled growth rate $\text{Im}(F_n)$ and (b) frequency $\text{Re}(F_n)$ vs pump parameter Γ .

compares the Rosenbluth analytic results with our numerically obtained results, where

$$F_n(\Gamma) \approx \left(n + \frac{1}{2}\right) \Gamma^{1/2} + i \left[2\Gamma - \left(n + \frac{1}{2}\right) \Gamma^{-1/2} \right]. \quad (77)$$

The threshold and frequency shifts at threshold obtained from Eqs. (72) through (74) appear in Fig. 2-49 as asterisks.

Despite the fact that the most unstable mode $n = 0$ is derived by assuming $\Gamma \gg 1$, accurate results for the growth rate are obtained from Eq. (77) down to the threshold at $\Gamma \approx 0.4$. However, the approximation in Eq. (77) fails approaching threshold, as expected for larger n ; but it is in this region that the complementary results in Eq. (72) apply. In practice, we are likely to be most concerned with the fastest-growing mode ($n = 0$).

If damping is small compared to $\gamma_0(V_2/V_1)^{1/2}$, then there is no significant regime of convective gain before absolute instability occurs at the threshold $\Gamma = 0.386$ determined by inhomogeneity. If the opposite is the case and damping is large, then at pump levels below the

threshold given by Eq. (77) essentially $\gamma_0 < v_2(V_1/V_2)^{1/2} + v_1(V_2/V_1)^{1/2}$ significant convective gain is possible. The convective gain is then given by Eq. (63).

Conclusions. We have completed our analysis of the Rosenbluth coupled-mode model for parametric instabilities in an inhomogeneous plasma for which the matching point lies close to an extremum of the plasma profile. For forward scattering, only convective growth is possible and is given by Eq. (71). For backscattering, absolute instability occurs above a threshold $\text{Im}(\omega) > 0$ determined from Eq. (75), where the function $F_n(\Gamma)$ has been determined numerically and approximated analytically in Eq. (77).

Our investigations of higher-order modes ($n > 0$) are of interest in the linear theory, but these modes may or may not play a role in the nonlinear evolution of the instability. Below the absolute threshold, significant convective growth for backscattering is possible only if the damping is sufficiently strong that the damping length for the idler (second) mode is short compared to the homogeneous growth length. In this case and as expected physically, the distinction between forward scattering and backscattering is lost, and the convective

gain takes the same form as that for forward scattering.

The primary limitation of our mode-coupling model is that it does not account for possible reflections of the waves from their respective turning points. Such reflections have been found to be an important feedback mechanism that can give rise to instability at lowered thresholds, provided the turning point(s) are close to the matching point.⁸⁴ The coupled-mode theory therefore does not apply to Raman or Brillouin sidescattering, to the Raman instability near quarter critical density, to stimulated Brillouin scattering near critical density, or to the two-plasmon decay instability. Each of these exceptional cases must be treated individually.^{84,85}

An issue of broad practical interest is the extent to which the interaction physics of large, reactor-scale plasmas can be investigated using available lasers and exploding-foil targets. The optimist's view is that it should be sufficient to match the local plasma conditions, such as density, temperature, laser intensity, and scale-length. Our analysis shows that for 3-wave parametric processes, we can achieve effective scalelengths in foil experiments many times the curvature length. For instance, to obtain one e -folding for forward Raman scattering in a 2-keV, 0- to $1-n_c$ plasma irradiated with 10^{15} W/cm² of 0.53- μ m light, a (disk) scalelength of 2700 μ m would be required. However, we can achieve the same gain in a foil plasma with a curvature length of only 60 μ m; such a plasma requires considerably less energy to create.

The scattering is localized in the foil within a small volume around the density maximum and to a corresponding small frequency bandwidth. In the equivalent disk profile, the scattering would extend over a much broader band of densities and frequencies. We must therefore be cautious in extrapolating the nonlinear consequences of such instabilities (e.g., superhot electron generation in the case of Raman forward scattering) from foils to equivalent disks because such nonlinear effects may not be purely locally determined. An analogy is the sensitivities of particle-code simulations to their boundary conditions.

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Simulations of Raman Backscattering and Forward Scattering in a Parabolic Density Profile

In burn-through foil experiments, Raman scattering has been explored in underdense plasmas with density scalelengths comparable to those expected in future high-gain targets. An important aspect of these experiments is the reduction of the intensity threshold of the Raman instability near the peak density of the expanding plasma.⁸⁶⁻⁹⁰ To assist in the interpretation of such experiments, we have carried out simulations of Raman backscattering and forward scattering in plasmas with a parabolic density profile.

In our simulations, we use a 1.5-D, electromagnetic relativistic particle code to propagate laser light from vacuum through a slab of underdense plasma. Because the Raman instability involves only the electrons, ion motion is frozen. The density profile was symmetric, and the plasma size was fixed at 128 vacuum wavelengths (λ_0) of the light. Peak density n and electron temperature T_e were selected, and the Raman scattering was measured in a series of simulations in which intensity was varied.

The onset of Raman backscattering and forward scattering as the intensity increases is evident in Fig. 2-50. In the simulations, the density had a parabolic profile ranging from 0.075 to 0.125 n_{cr} and back to 0.075 n_{cr} in 128 λ_0 and $T_e = 1$ keV. Here, n_{cr} is the critical density for the laser light. At low intensity, the fraction of light scattered with frequency downshifted by a plasma frequency is small and does not vary significantly with intensity. We interpret this as Thomson scattering⁹¹ from the thermal level of electron plasma waves in the simulation in which 80 000 electrons were used. As expected, the level of this scatter is considerably lower for forward scattering because the plasma wave then has a significantly longer wavelength. Both backscattering and forward scattering grow suddenly as the intensity increases.

It is difficult to accurately pinpoint the threshold intensity from these data, but our estimate is based on the intensity at which the signal strongly increases. For Raman backscattering, the threshold intensity is estimated to be about

$10^{14} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$; for forward scattering, we estimate a threshold intensity of about $3 \times 10^{14} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$. The threshold values are within a factor of 2 of the theoretical predictions discussed in the previous article.

In other simulations, we varied the peak density between 0.075 and $0.2 n_{cr}$ and the electron temperature up to 2 keV . The estimated threshold values are again found to be within a factor of 2 of predictions (typically higher). We plan more detailed comparisons, with particular attention to delineating regimes of convective and absolute instability for backscattering.

As shown in Fig. 2-51 by the phase-space plot from our simulation, the Raman instability occurred near the top of the density profile, where the gain is a maximum.⁸⁸⁻⁹⁰ The amplitude of the plasma waves near the top of the parabola is generally not symmetric. For a gain large enough that pump depletion occurs, the electron plasma waves peak toward the left, where the laser enters. The electron plasma waves extend to lower density on the right side compared to the left. Plasma waves generated on the top of the density profile readily convect to the right and to lower density, where they become progressively more damped.

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Nonlinear Features of Stimulated Brillouin Scattering

Stimulated Brillouin scattering (SBS) of laser light is a potentially important process in long-scalelength plasmas. A better understanding of the process is needed to calculate where laser energy is deposited. The deposition can, in turn, strongly affect the symmetry of explosions.

In this article, we present simulations of Brillouin scattering that illustrate several features observed in recent experiments.^{92,93} First, Brillouin scattering can reduce Raman scattering. Second, the Brillouin-generated ion wave can produce a significant second-harmonic component, in agreement with simple theory. To address a large region of plasma, we consider 1-D simulations that suffice to explore backscatter of laser light.

Fig. 2-50. Fraction of light backscattered and forward scattered vs intensity in a series of simulations.

Fig. 2-51. Electron phase space from a simulation of Raman scattering in a parabolic density profile. Density ranged from 0.1 to $0.2 n_{cr}$ and back to $0.1 n_{cr}$ in $128 \lambda_0$.

In the simulations, we use a 1.5-D (2 velocities) electromagnetic, relativistic particle code to propagate light from vacuum through a slab of underdense plasma with initially uniform density and temperature. Particles that reach the boundaries are re-emitted with their initial temperature, and reflected and transmitted light waves are allowed to freely exit the system. The scattered and transmitted light waves and plasma behavior are then monitored.

Figure 2-52 shows the evolution of the wave number k spectrum of the fluctuations in electron density n_e/n_0 . In this simulation, laser light with $I\lambda_0^2 = 2.5 \times 10^{15} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$ was propagated through a $128\lambda_0$ slab of plasma, with $T_e = 1 \text{ keV}$, $T_e/T_i = 30$, and $n/n_{cr} = 0.05$. Here I is laser intensity, λ_0 is laser wavelength, T_e is electron temperature, T_i is ion temperature, n is electron density, and n_{cr} is critical density. The brief line at $kc/\omega_0 \approx 1.7$ represents Raman backscattering, which occurs strongly only before the Brillouin

Fig. 2-52. Wave number spectrum of electron density n_e/n_0 . Time increases from bottom to top of the figure.

scattering becomes large. This dramatic effect has been experimentally observed.⁹²

The fraction of the Raman backscattered light was measured to be about 0.04 between 114 and 370 laser cycles. At later times (from 489 to 1001 cycles), the Raman reflection was sufficiently weak (about 0.004) that the associated density fluctuation becomes obscured in the noise. In an identical simulation, but with ions fixed, the wave number spectrum exhibits pulsations in the Raman scattering as noted before.⁹⁴ The peak fraction of ~ 0.04 is absorbed, and a fraction of about 0.016 is Raman reflected within times of 481 to 993 laser cycles.

The density fluctuation shown in Fig. 2-52 at $kc/\omega_0 \approx 1.95$ is due to Brillouin backscattering. The fluctuation is largest between 300 and 600 laser cycles, during which time the second harmonic at $kc/\omega_0 = 3.9$ can be clearly seen. The evolution of the amplitudes of the Brillouin-generated ion wave and its second harmonic is shown in Fig. 2-53. The fundamental peaks at an amplitude of $\delta n(k)/n_0 \approx 0.8\%$, while the amplitude of the second harmonic reaches $\delta n(2k)/n_0 \approx 0.2\%$. The percentage of light that is Brillouin backscattered varies with time but averages about 23% in this simulation.

A simple calculation accounts for the observed harmonic generation. We consider an ion sound wave with amplitude $\delta n(k)$ and wave number k . Using the two-fluid equations and iterating to second order in the amplitude $\delta n(k)$, we readily obtain the

Fig. 2-53. Spectral amplitude of the Brillouin (a) fundamental wave and (b) second harmonic wave vs time.

amplitude $\delta n(2k)$ of the second harmonic driven by the ion wave:

$$\frac{\delta n(2k)}{n_p} = \frac{1/2 [\delta n(k)/n_p]^2}{1/(1 + k^2 \lambda_D^2) - 1/(1 + 4k^2 \lambda_D^2)}, \quad (78)$$

and

$$\frac{\delta n(2k)}{n_p} \approx \frac{1/6 [\delta n(k)/n_p]^2}{k^2 \lambda_D^2}, \quad (79)$$

where $k^2 \lambda_D^2 \ll 1$.

If we substitute the peak $\delta n(k)$ from Fig. 2-53 and $k \lambda_D \approx 0.39$ into Eq. (78), then $\delta n(2k)/n_{cr} \approx 0.0015$, in reasonable agreement with the peak value observed in the simulation. For small $k \lambda_D$, the above iteration scheme breaks down because $\omega_i(2k) \approx 2\omega_i(k)$, and harmonics are too resonantly driven. For $\delta n(k)/n_p \gg 6k^2 \lambda_D^2$ estimated from Eq. (79), a rich spectrum of harmonics is expected and calculated.⁹⁵⁻⁹⁷ Harmonics of Brillouin-generated ion fluctuations have been directly observed⁹³ by Walsh and Baldis at NRC, Canada and by Pawley and Luhmann at UCLA.

The reduction of Raman scattering and/or the heated electron temperature when the Brillouin instability is operative at a significant level is illustrated in Table 2-8 for several simulations with different plasma densities. Note that the reduction is greater as the initial ratio of electron temperature to ion temperature is increased, i.e., as the level of Brillouin scattering is enhanced. An even stronger reduction (especially for densities near $0.25 n_{cr}$) might occur in 2-D simulations because side-scattered ion waves can then efficiently dampen both the plasma wave and the scattered light wave. Two effects may explain the decrease in Raman scattering: pump depletion by Brillouin scattering and competition from the shorter-wavelength Brillouin ion waves.

In the case for $n/n_{cr} = 0.05$, the time-averaged Brillouin scattering was about 23%. The Raman backscattering is proportional to exponent of the Raman gain,

n/n_{cr}	T_e	T_e/T_i	Raman	Raman forward	Brillouin	
			backscattering	scattering	T_{hot}	backscattering
0.05	1	30	0.04, 0.004 ^a	0.00009	4.5	0.23
0.05	1	fixed	0.04, 0.016 ^a	0.00007	6	—
0.1	1	3	0.18	0.005	12	0.25
0.1	1	30	0.12	0.001	10	0.31
0.1	1	fixed	0.26	0.13	17	—
0.2	3	5	0.33	0.07	42	0.02
0.2	3	fixed	0.33	0.18	60	—
0.2	3	30	0.23	0.044	40	0.15

^aLong-term (time > 500 cycles) Raman reflection.

where the gain is

$$\frac{k_{epw}^2 l}{8k_s} \frac{l \lambda_D^2}{1.4 \times 10^{18}} \frac{\omega_p^2}{\omega_i \omega_{epw}}, \quad (80)$$

Here k_{epw} and k_s are wave numbers of the electron-plasma wave; l is system length; and ω_p , ω_{epw} , and ω_i are plasma frequency, frequency of the electron plasma wave, and electron Landau damping, respectively. A reduction in pump intensity could make a difference, by one order of magnitude, in the scattering even if damping remained constant. Of course, damping does not remain constant; thus, the second effect obtains—that is, damping^{89,90} due to the shorter-wavelength ion waves. The reduction in the heated electron temperature is also attributed to generation of slower plasma waves via interaction with the ion waves.

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Monte Carlo Calculation of Non-LTE Line Radiation

Absorbed laser energy is converted to soft x radiation in high-Z plasmas in a thin density/temperature layer. The process is dominated by line emission from collisionally excited ions. Local reabsorption is weak, so that a source of radiation results. First-principles calculations of the radiation spectra in these non-local-thermodynamic-equilibrium (non-LTE) plasmas are not yet common. Here we describe a technique that appears to be efficient and complete enough to form the basis for such calculations.

Three difficulties must be overcome with respect to present schemes. One is that of computing atomic structure; another is

Table 2-8. Eight kinetic simulations showing reduction of Raman T_{hot} and fraction of light scattered with moving and fixed ions.

that of computing the collisional and radiative transition rates of the bound electrons. The XSN routine,^{100,101} for example, uses a screened hydrogenic atomic structure model that is far too crude for comparison with even modest-quality spectrographic data. Indeed, a fully self-consistent relativistic structure model appears to be both necessary and sufficient to successfully model the complex spectra of interest. An improved structure model naturally requires the associated transition rates. These issues, and techniques for overcoming the indicated difficulties, are discussed in the 1984 Laser Annual.¹⁰²

Because the number of excited ionic species making up the ensemble in the plasma is large at any (time, density, temperature) point, and because a large number of points must be computed, any scheme must be extremely efficient if it is to be useful. This is the third difficulty, and the subject of this article. The virtues and defects of the XSN scheme are that the average occupations of the levels of an average atom are evolved in time. Thus, at the expense of completeness and accuracy, the XSN scheme is extremely efficient.

Here we describe a Monte Carlo scheme that simulates in time the transition events of a single ion in the ambient plasma. The time-averaged values of such quantities as ionization state and emitted radiation spectrum describe the steady state of the system. This is usually the information of interest because the ambient plasma typically evolves more slowly than the atomic processes. (To simulate nonsteady systems,

we must evolve enough atoms to describe the ensemble dynamics.) The Monte Carlo scheme is readily implemented for any atomic structure and transition rate model. Furthermore, the important case of LTE is just a special case of the general non-LTE routine. In fact, the LTE case admits certain simplifications, which are described at the end of this article.

The outstanding virtue of our application of the Monte Carlo notion is its well-known efficiency in finding the bulk of the probability distribution of interest. This efficiency is especially significant in solving non-LTE problems. Thus we discover the radiated line spectrum according to relative strength: in studying how the energy is converted in laser plasmas, this is exactly the way in which we would like to proceed. Weak lines of diagnostic interest are probably better investigated by other means. The difficulty of our approach is the required knowledge of large amounts of detailed microphysical information, such as atomic energy levels and transition rates. This difficulty is inherent in any scheme that can produce results of the desired detail; our scheme reduces this burden to a near-minimum, and the burden may be further mitigated by storing such information for reuse.

Before taking up our Monte Carlo treatment of screened hydrogenic atomic physics, we elaborate somewhat on the screened hydrogenic atomic structure model. An ion is modeled as levels that are numbered $n = 1, 2, \dots, N$ whose instantaneous occupation is $N_n \in [0, \omega_n]$.

Fig. 2-54. (a) An ionic configuration vector is composed of the occupations of the bound levels. (b) The associated transition rate array includes every event that the ion can experience. The total rate is used in casting the lifetime of the configuration.

Here ω_n is the hydrogenic degeneracy, $2n^2$. The total energy of the ion is

$$E_{\text{ion}} = \sum_{n=1}^i N_n \varepsilon_n, \quad (81)$$

and $\varepsilon_n = -e^2 Z_n^2 / 2a_0 n^2$ is the screened hydrogenic level energy, where a_0 is the Bohr radius and

$$Z_n = Z - \sum_{n'=1}^i \sigma_{nn'} N_{n'} \quad (82)$$

is the screened charge experienced by an electron in level n . The $\sigma_{nn'}$ are strictly constant screening factors. A transition energy is simply the change in the total ion energy when an electron is moved from one level to another, including the continuum or free-electron levels not explicitly indicated.

Figure 2-54 shows the information necessary to advance a test atom in time. The ionic configuration (the list of the number of electrons occupying each atomic level) is shown in Fig. 2-54(a); the associated transition rates are shown in Fig. 2-54(b). For storage, we assign each configuration and its rates a number s . This permits fast reuse and avoids recomputation as the configuration recurs in the calculation. Recall that knowledge of the configuration and ion energy function implies knowledge of all transition energies. Of course, the transition energies and transition rates vary with the configuration itself, and each of these in turn requires evaluation as indicated. The XSN rates suggest that the transitions be displayed as a matrix as shown in the figure. Radiative and collisional rates for upwards (excitation and ionization) and downwards (deexcitation and recombination) transitions between levels are included. For example, R_{54}^s is the radiative decay rate for an electron from level 5 to level 4, and C_{4c}^s is the collisional ionization rate from level 4. The structure of any rate is $T_{mn}^s = N_m^s(\omega_n - N_n^s) t_{mn}^s$, where t_{mn}^s is the underlying single-electron rate. Many zeroes appear in the matrix, corresponding to full or empty levels.

To advance the test atom to a new configuration and to fully diagnose its history,

it is necessary to determine the incremental time dt_s spent in the present configuration s before a transition occurs and to identify the transition that occurs. The time dt_s is found from

$$dt_s = - \frac{\ln(\text{Ranf})}{\sum_{m,n} T_{mn}^s}, \quad (83)$$

Equation (83) realizes the exponential decay of the probability that the atom has made no transition. Here Ranf is a random number uniform in $[0,1]$. The associated rate is the total rate for all transitions.

The transition that occurs is found via the "pointer" rate

$$\hat{T} = \text{Ranf} \sum_{m,n} T_{mn}^s \quad (84)$$

Figure 2-55 illustrates the use of $\hat{T}$ in "looking up" the transition that occurs. The magnitudes of the individual rates are displayed, so that Eq. (84) realizes the relative probability of events as proportional to the relative magnitudes of the transition rates. The interval bounding the pointer is the associated event that is implemented. All information concerning the event is recorded and a new configuration s' is produced by moving an electron from level m to level n ; $N_m^{s'} = N_m^s - 1$, $N_n^{s'} = N_n^s + 1$. This procedure is repeated until satisfactory data of whatever kind are accumulated; typically, slowly varying time averages are accepted as a criterion for "steady state." In particular, the probability of the configuration N^s is given by

$$P(N^s) = \frac{\sum_{\text{events}} dt_s}{\sum_s \sum_{\text{events}} dt_s} \quad (85)$$

Fig. 2-55. The transition that occurs to the configuration is realized according to the relative magnitudes of the rates.

It should be clear that the notions developed here are readily implemented for any atomic structure and transition rate model. The bookkeeping required varies with the details of the model used. For example, the transition-rate matrix becomes sparse if forbidden transitions are ignored in a model that resolves subshell splitting. In general, straightforward procedures are readily found that manage the information well.

Figure 2-56 shows a result of our Monte Carlo XSN work. The emitted line spectrum for $5 \rightarrow 4$ transitions, which dominate the conversion process, is shown. No attempt to model subshell splitting has been made; only the transition energies of the screened hydrogenic atom appear. The data are from a non-LTE simulation in which the upwards radiative processes of excitation and ionization are inactive (the corresponding rates are set to zero). Table 2-9 gives average-atom data from this simulation. Table 2-10 shows the most

probable ions in the plasma and their $5 \rightarrow 4$ transition energies.

Table 2-11 gives data from a simulation chosen to exhibit the difference between emission from our non-LTE plasmas and emission from a plasma with LTE level populations. The average LTE ion is more ionized and more highly excited than the average non-LTE ion. This is reflected in the higher average $5 \rightarrow 4$ transition energy and in distinct spectral contrasts with the non-LTE data (see Fig. 2-57). There is also strong emission in the $4 \rightarrow 3$, $5 \rightarrow 3$, and $6 \rightarrow 5$ transitions.

The LTE data shown in Table 2-11 may be obtained by means other than direct simulation, and it is instructive to understand this application of the Monte Carlo scheme. A calculation of line emission may be viewed as comprising two parts: one, the determination of the probabilities of the configurations; and two, the application of the radiative line rates to each configuration. Of course, all the transition

Fig. 2-56. Time-averaged $5 \rightarrow 4$ line spectrum for the non-LTE Monte Carlo XSN run summarized in Table 2-9.

		$\bar{N}_3$ 18.00	$\bar{N}_4$ 10.13	$\bar{N}_5$ 0.38	$\bar{N}_6$ 0.16	$\bar{N}_7$ 0.12	$\bar{Z}^*$ 40.21
ϵ_{mn} (keV)	m	$n = 3$	4	5	6	7	c
	4	2.44	—	—	—	—	—
	5	3.24	0.78	—	—	—	—
	6	3.50	1.09	0.30	—	—	—
	7	3.75	1.26	0.48	0.18	—	—
	c	4.24	1.71	0.93	0.63	0.46	—
$\bar{T}_{mn}$ (keV/s-ion)							
	4	1.1×10^{11}	—	—	—	—	—
	5	1.6×10^9	4.5×10^{12}	—	—	—	—
	6	2.4×10^8	6.4×10^{11}	4.2×10^{11}	—	—	—
	7	1.3×10^8	1.9×10^{11}	1.1×10^{11}	1.2×10^{11}	—	—
	c	1.5×10^8	7.4×10^{10}	2.9×10^{10}	1.5×10^{10}	9.0×10^9	—
$\bar{R}_{mn}$							
	4	4.51×10^{10}	—	—	—	—	—
	5	3.13×10^8	5.71×10^{12}	—	—	—	—
	6	5.11×10^7	5.80×10^{11}	1.41×10^{12}	—	—	—
	7	1.46×10^7	1.50×10^{11}	2.26×10^{11}	6.78×10^{11}	—	—
	c	6.83×10^6	4.44×10^{10}	3.03×10^{10}	2.42×10^{10}	2.03×10^{10}	—
$\bar{C}_{mn}$							
	3	—	4.33×10^{10}	1.69×10^9	4.11×10^8	1.58×10^8	5.92×10^7
	4	1.64×10^8	—	6.46×10^{12}	5.57×10^{11}	1.37×10^{11}	4.02×10^{10}
	5	5.77×10^5	6.28×10^{11}	—	3.65×10^{12}	2.62×10^{11}	2.51×10^{10}
	6	8.87×10^4	3.05×10^{10}	2.07×10^{12}	—	4.72×10^{12}	3.77×10^{10}
	7	2.70×10^4	6.26×10^9	1.23×10^{11}	3.92×10^{12}	—	6.84×10^{10}
	c	4.82×10^3	8.89×10^8	5.70×10^9	1.50×10^{10}	2.26×10^{10}	—

Table 2-9. Average-atom values from a non-LTE Monte Carlo XSN run. The run produced 10 000 ionization/recombination events in 922 262 total events that subtended 2.9×10^{-8} s. (Gold, $\rho = 0.1$ g/cm³; $T_e = 450$ eV.) $\bar{N}_1 = 2.00$; $\bar{N}_2 = 8.00$.

N_4	N_5	N_6	N_7	ϵ_{54} (keV)	$P(N)$	N_4	N_5	N_6	N_7	ϵ_{54} (keV)	$P(N)$
10	—	—	—	—	0.124	11	1	1	—	0.76887	0.00537
11	—	—	—	—	0.119	8	1	1	—	0.61196	0.00474
9	—	—	—	—	0.0869	13	—	1	—	—	0.00447
12	—	—	—	—	0.0707	11	1	—	1	0.77113	0.00432
10	1	—	—	0.78751	0.0511	6	—	—	—	—	0.00396
8	—	—	—	—	0.0436	13	—	—	1	—	0.00366
9	1	—	—	0.80197	0.0424	8	1	—	1	0.81439	0.00353
11	1	—	—	0.77277	0.0374	7	—	1	—	—	0.00321
13	—	—	—	—	0.0341	12	2	—	—	0.74999	0.00300
8	1	—	—	0.81615	0.0245	12	1	1	—	0.75395	0.00297
12	1	—	—	0.75776	0.0218	7	2	—	—	0.82135	0.00271
10	—	1	—	—	0.0208	6	1	—	—	0.84369	0.00249
11	—	1	—	—	0.0168	12	1	—	1	0.75616	0.00246
9	—	1	—	—	0.0168	15	—	—	—	—	0.00235
10	—	—	1	—	0.0158	14	1	—	—	0.72689	0.00231
7	—	—	—	—	0.0139	10	—	1	1	—	0.00221
11	—	—	1	—	0.0131	7	—	—	1	—	0.00221
9	—	—	1	—	0.0123	9	—	1	1	—	0.00211
14	—	—	—	—	0.0107	11	—	1	1	—	0.00205
12	—	1	—	—	0.0104	7	1	1	—	0.82577	0.00195
7	1	—	—	0.83005	0.00965	10	—	2	—	—	0.00143
8	—	1	—	—	0.00896	9	—	2	—	—	0.00136
9	2	—	—	0.79364	0.00894	8	—	1	1	—	0.00130
13	1	—	—	0.74246	0.00871	7	1	—	1	0.82325	0.00129
12	—	—	1	—	0.00833	10	2	1	—	0.77545	0.00129
10	2	—	—	0.77937	0.00803	9	2	1	—	0.78963	0.00128
9	1	1	—	0.79788	0.00737	11	—	2	—	—	0.00124
10	1	1	—	0.78351	0.00725	13	—	1	1	—	0.00112
8	2	—	—	0.80764	0.00621	14	—	1	—	—	0.00112
8	—	—	1	—	0.00619	8	2	1	—	0.80353	0.00107
9	1	—	1	0.80025	0.00559	14	—	—	1	—	0.00103
10	1	—	1	0.78583	0.00557						

Table 2-10. Of the 750 ionic configurations produced by the calculations summarized in Table 2-9, here are the 64 configurations with probability $> 0.1\%$ and their 5 -- 4 transition energies. In all cases $N_1 = 2$, $N_2 = 8$, and $N_3 = 18$.

Plasma Physics

rates themselves determine the configuration probabilities in the non-LTE case, but this is not so in the LTE case. The principle of microscopic reversibility or detailed balance may be applied in the LTE case to recognize that the underlying microphysical transition rates are irrelevant to the configuration probabilities. This is the

well-known physics that leads to the Saha distribution of the configurations. Thus we can obtain the LTE configuration probabilities without knowledge of the true bound-free rates or of any bound-bound rates at all. In practice, we set the collisional rates and bound-bound radiative rates to zero and set all radiative recombination rates to

Table 2-11. Average-atom values from an LTE Monte Carlo XSN run. Same plasma conditions as the non-LTE run of Table 2-9. $\bar{N}_1 = 2.90$; $\bar{N}_2 = 8.00$.

		$\bar{N}_3$ 16.99	$\bar{N}_4$ 1.37	$\bar{N}_5$ 0.28	$\bar{N}_6$ 0.16	$\bar{N}_7$ 0.12	Z^* 50.07
ϵ_{mn} (keV)	n	$n = 3$	4	5	6	7	c
	4	2.61	—	—	—	—	—
	5	3.56	0.93	—	—	—	—
	6	3.99	1.36	0.43	—	—	—
	7	4.25	1.62	0.69	0.26	—	—
	c	4.97	2.32	1.38	0.96	0.70	—
T_{mn} (keV/s ion)	4	3.2×10^{13}	—	—	—	—	—
	5	1.9×10^{12}	7.2×10^{12}	—	—	—	—
	6	4.2×10^{11}	1.7×10^{12}	9.8×10^{11}	—	—	—
	7	1.4×10^{11}	5.8×10^{11}	2.7×10^{11}	2.8×10^{11}	—	—
	c	6.9×10^{10}	2.1×10^{11}	6.8×10^{10}	4.4×10^{10}	2.7×10^{10}	—

Fig. 2-57. Time-averaged $3 \rightarrow 4$ line spectrum for the LTE Monte Carlo XSN run summarized in Table 2-11.

unity. The principle of detailed balance gives us the photoionization rates from the radiative recombination rates. The Monte Carlo calculation then determines the configuration probabilities. Application of the more easily obtained line emission rates gives the emission spectrum.

We illustrate the preceding discussion by examining a model atom with a single (bound) level. Extension of this analysis to multilevel atoms is straightforward. The XSN nomenclature previously introduced is used. Only ionization and recombination transitions are possible, and the LTE probability that N electrons occupy the level obeys

$$\begin{aligned} \frac{dP(N)}{dt} = 0 = & -P(N)Nt_{\text{bf}}(N) \\ & + P(N-1)(\omega - (N-1))\xi t_{\text{fb}}(N-1) \\ & - P(N)(\omega - N)\xi t_{\text{fb}}(N) \\ & + P(N+1)(N+1)t_{\text{bf}}(N+1). \end{aligned} \quad (86)$$

Note that the probability is both increased and decreased by both ionization and recombination; b and f denote bound and free, respectively. Here ξ is the degeneracy factor of the free electrons; in nondegenerate plasma, $\xi = n_e h^3 / (2\pi m T)^{3/2}$.

The principle of detailed balance may be applied to solve Eq. (86) by pairwise cancellation:

$$\begin{aligned} P(N) = P(N-1) \\ \times \frac{\omega - (N-1)}{N} \xi \frac{t_{\text{fb}}(N-1)}{t_{\text{bf}}(N)}. \end{aligned} \quad (87)$$

In LTE, the ionization and recombination rates are related by

$$t_{\text{bf}}(N) = t_{\text{fb}}(N-1) \exp\left(-\frac{I_{N,N-1}}{T}\right) \quad (88)$$

The ionization/recombination energy is $I_{N,N-1} = E_{\text{ion}}(N-1) - E_{\text{ion}}(N)$.

Combining Eqs. (87) and (88), we obtain the Saha distribution:

$$\begin{aligned} P(N) = C \frac{\omega!}{N!(\omega - N)!} \\ \times \exp\left[-\frac{E_{\text{ion}}(N) + (Z - N)\mu}{T}\right] \end{aligned} \quad (89)$$

Here C is the normalization constant, determined by $\Sigma P(N) = 1$. The chemical potential of the free electrons is $\mu = T \ln \xi$. The free energy of the free electrons is $F = (Z - N)\mu$.

From our derivation of Eq. (89), it is evident that the actual ionization and recombination rates do not determine $P(N)$. The Monte Carlo treatment of Eq. (86) requires the relative rates for the terms representing transitions from the configuration N , namely $t_{\text{bf}}(N)/t_{\text{fb}}(N)$ from the terms with the negative signs. Equation (88) must be obeyed, but we are free to adopt

$$t_{\text{fb}}(N') = 1 \quad (90)$$

for all N' without compromising the essential physics or the essential result, Eq. (89). Now Eq. (88) with Eq. (90) give the information required for the Monte Carlo simulation, which effectively evaluates $P(N)$ of Eq. (89) according to relative magnitude. By this we mean that the larger values of P are accurately determined first during the simulation run.

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Electromagnetic Effects in Relativistic Electron-Beam-Plasma Interactions

Beam-plasma instabilities are of interest in many applications, including plasma heating, space and astrophysical plasmas, the generation of radiation, and perhaps in some laser particle accelerators. The large number of previous theories and computer simulations of beam-plasma instabilities were primarily motivated by applications to plasma heating.¹⁰³⁻¹⁰⁵

Our work is motivated by recent experiments that measured the electromagnetic radiation due to a relativistic electron beam interacting with a plasma.¹⁰⁶ When the beam density was comparable to the plasma density, it was found that a significant fraction (on the order of 0.2%) of the beam energy was converted to microwaves. These waves contain a broad range of frequencies, extending to many times the electron-plasma frequency ω_{pe} , rather than simply containing frequencies of ω_{pe} and its harmonics.

To investigate the beam-plasma interaction, we have used a 2-D, electromagnetic, relativistic particle code.¹⁰⁷ The beam travels in the x direction, and the evolution of $\mathbf{E}_x$, $\mathbf{E}_y$, and $\mathbf{B}_z$ is followed in the x - y plane. We started with the simplest problem and followed the temporal evolution of an electron beam in a uniform plasma with no guiding magnetic field. The boundary conditions were doubly periodic; the simulation box was rectangular, with $L_x = 10.9 c/\omega_{pe}$ and $L_y = 50.2 c/\omega_{pe}$, where c is the velocity of light and ω_{pe} is the electron-plasma frequency.

The results changed by less than about 20% when both L_x and L_y were decreased by one-half. The ions were usually treated as a fixed, neutralizing background, and 800 000 electrons were used. The beam initially had a kinetic energy of approximately 0.5 MeV (i.e., $\gamma = 2$) with a narrow thermal spread of 50 eV. The background electrons had a temperature of 2.5 keV and were initialized with a drift to balance the initial beam current. The beam density (n_b) ranged from 0.05 to $0.5 \times$ the total electron density (n_i).

The beam-plasma interaction in our simulations is dominated by two instabilities. The first is the well-known electron-beam instability in which electrostatic waves are driven unstable.

For a beam with a narrow spread in energy and angle, and for $n_b/n_p\gamma_b < 1/20$, the results for the frequency and growth rate of the most unstable mode are

$$\frac{\omega}{\omega_{pe}} = 1 - \frac{1}{2} \left(\frac{n_b}{2n_p\gamma_b} \right)^{1/3} f(\theta) \quad (91)$$

Fig. 2-58. Evolution of the electrostatic field energy in a 2-D simulation ($\gamma_b = 2$, $n_b = n_p$).

and

$$\frac{\gamma}{\omega_{pe}} = \sqrt[3]{3/2} \left(\frac{n_b}{2n_p\gamma_b} \right)^{1/3} f(\theta), \quad (92)$$

where

$$f(\theta) = \left(\sin^2 \theta + \frac{\cos^2 \theta}{\gamma_b^2} \right)^{1/3} \quad (93)$$

Here, $n_b(\gamma_b)$ is the density of the beam, n_p is the plasma density, v_b is the beam velocity, θ is the angle between the wave vector k and v_b , and $k \cdot v_b \approx \omega_{pe}$.

The second instability corresponds to a purely growing, spatially varying magnetic field transverse to the beam. The maximum growth rate occurs for $\theta = 90^\circ$ and is

$$\frac{\gamma}{\omega_{pe}} = \left(\frac{n_b}{n_p\gamma_b} \right)^{1/2} \frac{k v_b}{(k^2 c^2 + \omega_{pe}^2)^{1/2}} \quad (94)$$

This Weibel instability occurs because parallel currents attract. Hence, beam electrons tend to bunch together or coalesce into filaments as do the background electrons carrying the return current. Of course, the Weibel instability also produces an electrostatic field that expels background electrons from a filament to preserve quasi-neutrality.

Simulation Results. To illustrate the simulation results, consider a robust example in which $n_b = n_p$. The electrostatic field energy EL vs time is shown in Fig. 2-58. At early time there is rapid growth ($\gamma \approx 1/2\omega_{pe}$), followed by a period of slower growth as the beam nonlinearly develops a significant spread in energy and angles. Finally, the wave energy saturates at about 12% of the initial kinetic energy of the beam as the velocity distribution becomes monotonic. The wave energy thereafter slowly decays.

The magnetic field energy $B_z^2/8\pi$ vs time is shown in Fig. 2-59. The magnetic field due to the Weibel instability initially grows rapidly ($\gamma \approx 1/2\omega_{pe}$) and then grows more slowly as the beam spreads in energy and angles. Finally, the magnetic field saturates at a large value that corresponds to $\omega_{ce}/\omega_{pe} \approx 0.5$, where $\omega_{ce} = eB/mc$ is the electron-cyclotron frequency for a nonrelativistic particle.

As the unstable waves grow and saturate, electromagnetic waves are generated,

as shown in Fig. 2-60. In this example, the transverse electric field energy EM reaches about 2% of the kinetic energy of the beam (KE_b). These electromagnetic waves have a significant range of frequencies rather than only a frequency of ω_{pe} and its harmonics. Figure 2-61 shows the Fourier transform of B_z at a given location in the plasma.

Note the large peak at $\omega = 0$ due to the Weibel instability, followed by a signal extending from ω_{pe} to 3 or $4 \times \omega_{pe}$ due to light waves. A broad range of frequencies is generated, a feature noted in the experiments. However, the amplitude of the light waves falls off more rapidly with frequency than was experimentally observed.

We have also simulated weaker beam densities. For a beam of $\gamma_b = 2$ and with $n_b/n_t = 0.3$, the peak magnetic field is $\omega_{ce}/\omega_{pe} \approx 0.17$; about 0.3% of the kinetic energy of the beam is converted to transverse electric fields. For $n_b/n_t = 0.05$, the peak magnetic field is $\omega_{ce}/\omega_{pe} \approx 0.017$, and the fraction converted into transverse electric fields is only about 0.06%. In all these cases, the beam loses about 40% of its energy to the plasma. The beam finally contains a spread in velocity comparable to its initial velocity as well as a broad range of angles.

Although we obtain many of the qualitative features observed in experiments, our simulations are highly idealized. The transient temporal problem does not do justice to effects such as those associated with ion motion. For example, note in Fig. 2-58 that the electrostatic field rises and decays in a time of $\approx 100 \omega_{pe}^{-1}$, which is too brief for the ions to respond. A spatial problem in which the beam is continually injected would be more realistic and would also admit boundary conditions closer to those in experiments. Allowance for spatial variations in the background plasma density and beam quality may also be important for details of the radiation generation and the power spectrum.

As a result of such considerations, we recently carried out simulations of an improved 2-D model. The plasma and beam were again periodic in the direction of the beam velocity, but were given a finite spatial extent in the direction transverse to the beam. The finite spatial extent allowed us to examine radiation that escapes the

Fig. 2-59. Evolution of the magnetic field energy in a 2-D simulation.

Fig. 2-60. Evolution of the transverse electric field energy EM in a 2-D simulation.

Fig. 2-61. Fourier transform of the magnetic field at a specific location in the plasma ($\gamma_b = 2$, $n_b = n_p$).

plasma into a surrounding vacuum. The temporal problem was again considered; i.e., the beam was initialized at $t = 0$, and the evolution followed. The gross features were the same as before, but we obtained a few notable exceptions. Radiation was

observed with frequencies less than that of the electron-plasma frequency of the central region of the plasma. In other words, the lower-density plasma along the edges also radiated. In addition, the spectrum of the radiation above ω_{pe} was somewhat flatter than that shown in Fig. 2-61.

Our simulations indicate that beam-plasma interactions (electrostatic and electromagnetic) might play a role in schemes such as the beat-wave accelerator.^{108,109} The effect of these instabilities will depend on factors such as the beam density and energy, and the spatial width and structure. In experiments currently planned at UCLA,¹¹⁰ it is estimated that $n_b/n_0 \approx 10^{-2}$ and $\gamma_b \approx 20$. The growth rate of the electron-electron beam instability is then $\approx 0.05 \omega_{pe}$, and that of the Weibel instability is $\approx 0.02 \omega_{pe}$. For future applications, the beam density (and γ_b) will be significantly larger. However, it is not clear that these instabilities need have a deleterious effect in ultimate applications.

The finite radial extent of the beam can be a stabilizing influence, limiting the wave vectors of the unstable waves. In addition, the finite extent of the beam bunch may allow the beam to "outrun" unstable fluctuations,¹¹¹ particularly in view of the fact that the fluctuations have wave vectors at significant angles to the beam. Improved simulations can help clarify potential limitations and guide future experiments.

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Implicit Particle Code Development

Electromagnetic Direct Implicit Plasma Simulation. During 1985 we made significant advances in the development of a Cartesian (x,y) electromagnetic (EM), direct implicit, particle-in-cell (PIC) code. The essence of an implicit PIC algorithm is to use part of the yet-to-be-determined fields to advance the particle/fields to the next level so that stability with a large time step can be maintained.¹¹²⁻¹¹⁴ The major impediment to the development of implicit methods to PIC techniques has been the difficulty of solving the field equations rapidly and robustly in the parameter regimes of general interest.

The basic idea of the direct implicit approach is to use an intermediate time level for both particle position and velocity, which is obtained by using all of the particle dynamical equations except that due to the fields at the advanced time level. By exploiting the definitions of plasma charge and current densities as sums over the particle shape factor, the advanced plasma source terms can be expressed in two parts. One part is obtained by summing over the intermediate level; the second part is expressed as a linear operator on the advanced field. These expressions can then be used directly in the field solution. Once the advanced fields are obtained, the last piece of the particle advances from the intermediate level to the new positions, and velocities can be completed.

As with the moment method,¹¹⁵ an inconsistency can exist (depending on the care taken in the preparation of coefficients multiplying the advanced fields) between the new fields and the source terms that could be obtained from the new positions and velocities. We have found it unnecessary to employ an iteration over this final correction of particle coordinates.

Implementation details now used in the AVANTI code include the representation of quantities on an interleaved mesh, the simultaneous splitting technique, and the methods used to provide relatively general boundary conditions. Code applications are still largely confined to those problems for which the physics is well understood so that operating characteristics of the model can be established. Our studies include a more careful analysis of the code behavior on the well-known problem of electron-beam filamentation. Preliminary results are given in Ref. 116.

Overview of the Direct implicit PIC Algorithm. The particle integration scheme in the direct implicit algorithm is implemented in two conceptual steps. First, we advance the particles to the n -level using everything but the advanced E field. This step is called the "pre-push." For each particle we use

$$\begin{aligned} \mathbf{v}_{n+1/2} &= \mathbf{v}_{n-1/2} + \frac{1}{2} \Delta t \\ &\times \left[\bar{\mathbf{a}}_{n-1} + \frac{1}{2} (\mathbf{v}_{n+1/2} - \mathbf{v}_{n-1/2}) \frac{q \mathbf{B}_n(\mathbf{x}_n)}{mc} \right] \end{aligned} \quad (95)$$

and

$$\tilde{\mathbf{x}}_{n+1} = \mathbf{x}_n + \Delta t \tilde{\mathbf{v}}_{n+1/2} \quad (96)$$

followed by the accumulation of some particle-derived source terms at the n -level. The n -level source terms are constructed such that the advanced source terms can be obtained from the source terms when the new $\mathbf{E}$ field is known. The expressions we use are

$$\rho_{n+1}(\mathbf{X}_j) = \rho_{n+1}(\mathbf{X}_j) - \nabla \cdot \chi \cdot \mathbf{E}_{n+1} \quad (97)$$

and

$$\begin{aligned} \Delta t \mathbf{J}_{n+1/2}(\mathbf{X}_j) &= \Delta t \mathbf{J}_{n+1/2}(\mathbf{X}_j) + \chi \cdot \mathbf{E}_{n+1} \\ &- c \Delta t \nabla \times \zeta \cdot \mathbf{E}_{n+1}, \end{aligned} \quad (98)$$

where ρ , $\mathbf{J}$, χ , and ζ result from summations over n -level particle quantities. The precise form of the tensors χ and ζ is discussed elsewhere.^{116,117} To solve for the advanced fields, these expressions are substituted into the time differencing D_1 scheme, for which the Maxwell field equations are:

$$\begin{aligned} \mathbf{E}_{n+1} - \mathbf{E}_n &= c \Delta t \nabla \times \mathbf{B}_{n+1/2} \\ &+ 4\pi \Delta t \mathbf{J}_{n+1/2}, \end{aligned} \quad (99)$$

$$\mathbf{B}_{n+1/2} - \mathbf{B}_{n-1/2} = c \Delta t \nabla \times \bar{\mathbf{E}}_n \quad (100)$$

$$\mathbf{B}_{n+1} - \mathbf{B}_{n+1/2} = \frac{-1}{2} c \Delta t \nabla \times \mathbf{E}_{n+1}, \quad (101)$$

and

$$\bar{\mathbf{E}}_n = \frac{1}{2} (\mathbf{E}_{n+1} + \mathbf{E}_{n-1}), \quad (102)$$

Given these particle-derived source terms, the solution of the field equations becomes a numerical task independent of the particle equations. Our new field solution technique is described below.

Assuming that the advanced fields have been calculated, the particle advance can now be completed with these new fields. The final advance for each particle (called

the "final push") is given by

$$\begin{aligned} \delta \mathbf{v} &= \frac{\Delta t}{2} \left[\frac{q}{m} \mathbf{E}_{n+1}(\tilde{\mathbf{x}}_{n+1}) \right. \\ &\left. + \frac{\delta \mathbf{v}}{2} \times \frac{q \mathbf{B}_n(\mathbf{x}_n)}{mc} \right], \end{aligned} \quad (103)$$

$$\mathbf{x}_{n+1} = \tilde{\mathbf{x}}_{n+1} + \Delta t \delta \mathbf{v}, \quad (104)$$

$$\mathbf{v}_{n+1/2} = \tilde{\mathbf{v}}_{n+1/2} + \delta \mathbf{v}, \quad (105)$$

and

$$\bar{\mathbf{a}}_n = \frac{1}{2} \left[\bar{\mathbf{a}}_{n-1} + \frac{q}{m} \mathbf{E}_{n+1}(\tilde{\mathbf{x}}_{n+1}) \right]. \quad (106)$$

Further discussion of this method can be found in Refs 114, 116, and 117.

Solution of the Implicit Field Equation.

The field equations are obtained by substituting the expressions for ρ_{n+1} and $\mathbf{J}_{n+1/2}$ into Eqs. (99) through (102). By rearranging, we can replace Eq. (99) with

$$\begin{aligned} \mathbf{E} - \frac{1}{2} c^2 \Delta t^2 [\nabla^2 \mathbf{E} - \nabla (\nabla \cdot \mathbf{E})] \\ + 4\pi (\chi \cdot \mathbf{E} - c \Delta t \nabla \times \zeta \cdot \mathbf{E}) = \mathbf{Q} \end{aligned} \quad (107)$$

and

$$\begin{aligned} \mathbf{Q}' &= \mathbf{E}_n - 4\pi \Delta t \mathbf{J} + c \Delta t \nabla \times \mathbf{B}_{n-1/2} \\ &- \frac{1}{2} c^2 \Delta t^2 \nabla \times \nabla \times \bar{\mathbf{E}}_{n-1}, \end{aligned} \quad (108)$$

where the $\mathbf{E}$ without subscript is $\mathbf{E}_{n+1}$ and, for the moment, $\mathbf{Q} = \mathbf{Q}'$. This equation for $\mathbf{E}$ governs the time evolution of both the electrostatic and inductive parts of the electric field.

The solution of Eq. (107) with $\mathbf{Q} = \mathbf{Q}'$ does not satisfy

$$\nabla \cdot (\mathbf{I} + 4\pi \chi) \cdot \mathbf{E}_{n+1} = 4\pi \rho_{n+1}, \quad (109)$$

which is our best statement of Gauss's law $\nabla \cdot \mathbf{E}_{n+1} = 4\pi \rho_{n+1}$. By taking the divergence of Eq. (107), we see that $\mathbf{E}_{n+1}$ satisfies instead $\nabla \cdot (\mathbf{I} + 4\pi \chi) \cdot \mathbf{E} = \nabla \cdot \mathbf{Q}'$, in which $\nabla \cdot \mathbf{Q}' \neq 4\pi \rho_{n+1}$ due to the method of forming $\bar{\mathbf{J}}$ and to incomplete convergence of the final push at earlier times. This nonconservative aspect of the area-weighted $\bar{\mathbf{J}}$ was noted previously.¹¹⁸

Ignoring this issue permits $\nabla \cdot \mathbf{E}$ to drift far away from ρ because ρ does not appear in Eq. (107).

We first consider an adjustment to $\mathbf{Q}$ that is analogous to one used successfully in several explicit EM codes.¹¹⁹ Because ρ_{n+1} provides information only on the longitudinal part of $\mathbf{Q}$, we adjust only that part of $\mathbf{Q}$, as follows:

$$\mathbf{Q} = \mathbf{Q}' - \nabla\psi. \quad (110)$$

The equation for ψ , obtained by substitution into

$$\nabla \cdot \mathbf{Q} = 4\pi\rho, \quad (111)$$

is

$$\nabla^2\psi = \nabla \cdot \mathbf{Q} - 4\pi\rho. \quad (112)$$

The effect of this adjustment to $\mathbf{Q}$ is quite localized in space and is not propagated by the EM waves.

In the context of implicit codes, this correction has been observed to cause two problems. The first problem arises in simulations with large values of $\Omega\Delta t$. This field $\mathbf{E}_{n+1}$ corresponds to spuriously large currents across the magnetic field because the correction $\nabla\psi$ was computed without information concerning the anisotropy of the mobility caused by $\mathbf{B}$. The second problem arises in simulations with large and abrupt changes in density, such as those at a plasma-vacuum interface. We have found spuriously large fields $\mathbf{E}_{n+1}$ in the low-density region. Large values of the right-hand side of Eq. (112) in the high-density region lead to finite "corrections" in the near vicinity, including regions where χ is much smaller, and therefore $\mathbf{E}_{n+1}$ can be large.

We were motivated by such observations to include χ in the correction process. Our present solution to both problems is to take a new form for the correction, as follows:

$$\mathbf{Q} = \mathbf{Q}' - (1 + 4\pi\chi) \cdot \nabla\psi. \quad (113)$$

Substitution into Eq. (111) yields an elliptic equation for ψ , such that

$$\nabla \cdot (1 + 4\pi\chi) \cdot \nabla\psi = \nabla \cdot \mathbf{Q} - 4\pi\rho. \quad (114)$$

Corrections similar to those in Eq. (113) were first suggested by Barnes¹¹⁴ to achieve the qualitative feature that the correction current $\propto (\mathbf{Q} - \mathbf{Q}')$ is biased toward flowing along $\mathbf{B}$. This correction affects both the solenoidal and irrotational parts of $\mathbf{Q}$.

The improved robustness of the code, when large and abrupt density changes arise, can be illustrated as follows: Consider the interface between a high-density ($\chi \gg 1$) plasma and a vacuum ($\chi \ll 1$). Although the two forms of the correction are similar in magnitude in the plasma, the solution ψ of Eq. (114) is smaller than that of Eq. (112) by about the factor χ . In the vacuum, ψ is the solution of a Poisson equation connecting continuously to ψ in the plasma; ψ in the vacuum is thus not larger than ψ at the plasma edge. The vacuum correction $-\nabla\psi$ is therefore much smaller with the second form of correction.

Numerical Implementation of the Field Solution. The solution of Eq. (107) is the most interesting aspect of the field solution. The field equations, Eqs. (100) through (103), can be simply evaluated in a point-by-point manner once $\mathbf{E}_{n+1}$ is known. Equation (107) involves a coupling of $\mathbf{E}$ components and plasma source terms that has proven to be difficult to solve if not properly arranged. Tests reported last year¹¹⁶ revealed the advantage of solving all components simultaneously with a linear system solver. Because a direct 2-D solution would require too much storage, an iterative solution has been implemented. We have achieved good results in limited-parameter regimes with an alternating direction implicit (ADI) method, which solves all three component equations simultaneously in each direction.

The parameter regimes for which the code runs well have been greatly expanded by our efforts this year. The difficulty previously encountered arose when the coefficient $1/2c^2\Delta t^2$ divided by $\Delta_i^2 = \min(\Delta x^2, \Delta y^2)$ is large. In this case, the term $\nabla \times \nabla \times$ dominates the equation. The convergence rate slows—presumably because the operator $\nabla \times \nabla \times$ does not fit symmetrically into the splitting scheme. Because we expect to simulate strongly inhomogeneous plasmas or vacuum regions, the algorithm must effectively deal with regions in which $1/2c^2\Delta t^2/\Delta_i^2 \gg \chi$ and for which

the collisionless skin depth is well resolved.

To render the term $\nabla \times \nabla \times$ more tractable in these regions, we employ the traditional technique of splitting $\mathbf{E}$ into irrotational and solenoidal parts, denoted by $\mathbf{E}_i$ and $\mathbf{E}_s$, respectively. Because $\mathbf{E}_i = -\nabla\Phi$, where Φ is the electrostatic potential, we must self-consistently compute the $n + 1$ time level of four 2-D quantities. Equation (107) still determines the total advanced electric field, although it only provides three scalar equations. An additional equation is needed to serve as a constraint condition on Φ .

We deal with this problem by solving Eq. (107) and using its divergence as a constraint equation

$$\begin{aligned} \mathbf{E}_i - \nabla\Phi - \frac{1}{2}c^2\Delta t^2\nabla^2\mathbf{E}_i \\ + 4\pi\chi\cdot(\mathbf{E}_i - \nabla\Phi) \\ - c\Delta t\nabla \times \zeta\cdot(\mathbf{E}_i - \nabla\Phi) = \mathbf{Q} \end{aligned} \quad (115)$$

and

$$\nabla\cdot(\mathbf{I} + 4\pi\chi\cdot)(\mathbf{E}_i - \nabla\Phi) = \nabla\cdot\mathbf{Q}. \quad (116)$$

These equations provide a set of four coupled, elliptic equations that offer symmetry in the terms associated with each coordinate direction. The essential feature of this decomposition is that we can obtain the correct cancellation of the terms $\nabla^2\mathbf{E}_i - \nabla(\nabla\cdot\mathbf{E}_i)$ analytically if we use an interleaved mesh.¹¹⁶

To summarize, the strong coupling of these equations makes simultaneous

solution necessary; an ADI procedure has been devised that seems to work well in parameter regimes of interest for this code. We generally require one part in 10^3 for all components, and the necessary precision is usually obtained with only 10 to 15 iterations. The details of this splitting procedure can be found in Ref. 117.

Test Cases. We have used electron-beam filamentation as one test problem for our code. Preliminary results were reported last year in the limit of many mesh points in a collisionless skin depth. In the opposite limit, the algorithm reported last year failed to deliver useful results. The present algorithm works well in both regimes; we have recovered filamentation growth rates that are in good agreement with simple theory.¹¹⁷ Our code provides useful results with time steps larger than 100 times the explicit stability limit. We find no stability limit on the time step.

To obtain a more standard comparison of the computational effort expended on achieving implicitness, we extrapolate our results for the plasma-vacuum to the more familiar, homogeneous plasma. If we assume 15 particles of each of 2 species in every cell on a 40×32 mesh, our present code uses 60% of the effort in the field solution and 22% for the particle push. The remaining 18% is given to overhead associated with plotting. Our implicit particle pusher (unoptimized Fortran) takes about 50% longer than the standard (unoptimized Fortran) explicit pusher.

The field solution in an explicit code is generally negligible compared to the particle advance. Thus, the overhead of our implicit field solution increases our

Fig. 2-62. 1-D profiles of initial configuration for magnetic reconnection studies. (a) Ion and electron densities, (b) ion and electron temperatures, (c) confining fields, and (d) electron drift velocity.

computational effort by a factor of 4 per time step over the explicit method. The gain, of course, is that the typical time step for an explicit code ($\omega_{pe} \Delta t = 0.2$) can often be increased by one to two orders of magnitude, depending on the problem.

In a further application, we have begun to apply the code to basic physics questions concerning magnetic reconnection (supported by IRD) as a means of testing its behavior in the presences of large, zero-order magnetic fields. Starting with a

Fig. 2-63. Progression of magnetic reconnection simulated by AVANTI. The left column contains vector plots of the $\mathbf{B}$ field in the plane of simulation showing the emergence of reconnection; the right column shows contours of the normal component of the magnetic vector potential or magnetic flux.

uniform ion charge density $\rho_{\text{ion}} = 1$, and with the 1-D initial temperature profiles

$$T = (T_{\text{max}} - T_{\text{min}}) \times [1 - \tanh^2(\alpha x^2)] + T_{\text{min}} \quad (117)$$

for both electrons and ions, we construct the electric field E_x necessary to confine the ions

$$E_x = \partial_x (\rho T_{\text{ion}}) / \rho_{\text{ion}}, \quad (118)$$

assuming that the ions carry no current. Gauss's law provides the electron density required to produce this electric field

$$\rho_{\text{elect}} = -\rho_{\text{ion}} + \partial_x E_x, \quad (119)$$

and pressure balance

$$\partial_x \left(\frac{B_y^2 - E_x^2}{8\pi} \right) + \frac{\rho_{\text{elect}} T_{\text{elect}}}{q_{\text{elect}}} + \frac{\rho_{\text{ion}} T_{\text{ion}}}{q_{\text{ion}}} = 0 \quad (120)$$

gives the confining magnetic field. A curl of this magnetic field determines the current the electrons must carry. These initial profiles are shown in Fig. 2-62.

As time progresses in the simulation, the opposed magnetic field lines reconnect across the field null, as shown in the vector plots of $\mathbf{B}$ and the associated contours of constant magnetic flux in Fig. 2-63. Recalling that magnetic field vectors are everywhere tangent to these contours gives a clean view of the reconnection process. These simulations were run with $\Delta t = 5\omega_{pe}^{-1}$ (a value approximately 25 times larger than could be used in an explicit code), with a mass ratio of 200, and with 40×32 cells. We are now investigating several new phenomena that have emerged from our more complete model.

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Darwin Field Model for Charged-Particle Source Studies

The quality of a charged-particle beam is a strong function of the performance of

the particle injector, the device that generates the input beam to the accelerator. In particular, the brightness (ratio of current to phase space volume) is strongly dependent on the anode-cathode geometry. For long-pulse operation, steady-state codes have been used to model the complex interrelationships among applied voltage, cathode and anode position, output current, and brightness. Recent interest in high-frequency pulsed operation requires proper modeling of the transient behavior as well.

A standard approach that encompasses the transients is to make use of fully electromagnetic (EM) models that include all time-dependent, EM physics. Another model, the Darwin model, may also include the relevant physics at significantly less expense when purely EM modes do not play a significant role. In the Darwin model^[20,121] plasma coupling to EM modes is eliminated by a modification to Maxwell's equations. Neglecting the solenoidal part of the displacement current (the part that has no divergence, which we denote by the subscript l) allows the particle Lagrangian to be accurate through terms of the order v^2/c^2 , while neglecting plasma source-term coupling to purely EM modes. The resulting EM field equations, are

$$\nabla \times \mathbf{B} = \frac{4\pi}{c} \mathbf{J}_l, \quad (121)$$

$$\nabla \cdot \mathbf{B} = 0, \quad (122)$$

$$\nabla \times \nabla \times \mathbf{E} = -\frac{4\pi}{c^2} \mathbf{J}_l, \quad (123)$$

and

$$\nabla \cdot \mathbf{E} = 4\pi\rho, \quad (124)$$

where we have used the fact that the charge continuity equation provides cancellation of the irrotational parts of the current and the displacement current. (We denote the irrotational part, which has no curl, by the subscript l .)

This modification also changes Maxwell's equations from hyperbolic to elliptic in character, reflecting the instantaneous propagation speeds characteristic of the Darwin model. The instantaneous propagation of field information necessitates a

fully implicit solution of the EM fields consistent with the plasma source terms at any instant of time. This requires a more complex field solution algorithm and the accumulation of some additional particle moments.

Instantaneous propagation does, however, provide several benefits. The most obvious benefit is that the time step Δt is no longer constrained by the CFL condition:

$$c \Delta t < \min(\Delta x, \Delta y), \quad (125)$$

which expresses the idea that light cannot propagate over more than one computation cell in one time step. Our model eliminates the CFL constraint in Eq. (125) on isotropic light propagation by using the Darwin approximation; we are left with a similar, but less restrictive, condition on the particle velocity. (Injector ions are expected to have velocities significantly less than c .) Added savings can now be realized because these ion velocities are not isotropic; we may resolve ion trajectories using elongated mesh cells in the direction of ion propagation. Further benefits may also be achieved due to the elimination of the spurious energy in purely electromagnetic modes resulting from the inherently large fluctuations in PIC representations.

We have recently completed work on a Darwin model PIC code, referred to as DPC, for use in modeling electron injectors.^{122,123} The DPC code also has the capability of modeling ion injectors. Figure 2-64 shows one of the cathode geometries that is possible with this code.

Early versions of the code contained troublesome requirements on boundary conditions; the most difficult problem arises in the solution of Eq. (123). The equation for the inductive $\mathbf{E}_i$ requires developing an implicit representation for $\mathbf{J}_i$. Following Ref. 120, $\mathbf{J}_i$ can be obtained by accumulating additional particle quantities associated with the kinetic energy tensors of each species, so that Eq. (123) can be written

$$\begin{aligned} \nabla^2 \mathbf{E}_i &= \frac{4\pi}{c^2} \mathbf{J}_i \\ &= \frac{4\pi}{c^2} [\mathbf{K} + \mu(\mathbf{E}_i + \mathbf{E}_e) + \zeta \times \mathbf{B}], \end{aligned} \quad (126)$$

where $\mathbf{K}$ represents divergence of the kinetic energy tensor summed over species

$$\mathbf{K} = - \sum_s q_s \sum_p \mathbf{v}_p \mathbf{v}_p \cdot \nabla \mathbf{S}(\mathbf{x}_s - \mathbf{x}_p), \quad (127)$$

and s is the particle shape factor. The terms μ and ζ are given by

$$\mu = \sum_s \frac{q_s \rho_s}{m_s} \quad (128)$$

and

$$\zeta = \sum_s \frac{q_s \mathbf{J}_s}{m_s c}. \quad (129)$$

Because the right-hand side of Eq. (126) explicitly contains $\mathbf{E}_i$ in an expression that must be vector decomposed, iteration is required to obtain an implicit solution. Introducing the concept of iteration on the superscript n , the traditional algorithm is:

1. Form $\mathbf{J}$ using the latest $\mathbf{E}_i^n$.
2. Vector decompose (Ref. 125) $\mathbf{J}$ to find $\mathbf{J}_i$ using the procedure

$$\nabla^2 \Psi = - \nabla \cdot \mathbf{J}, \quad (130)$$

followed by

$$\mathbf{J}_i = \mathbf{J} + \nabla \Psi. \quad (131)$$

(Note that the boundary condition required for the Poisson equation involves specifying the normal component of $\mathbf{J}_i$ on the boundaries.)

3. Solve

$$\nabla^2 \mathbf{E}_i^{n+1} = \frac{4\pi}{c^2} \mathbf{J}_i^n. \quad (132)$$

4. Vector decompose $\mathbf{E}_i^{n+1}$ using a procedure similar to Step 2 to eliminate any spurious irrotational part.

5. Check convergence, and, if inadequately converged, return to Step 1.

Choosing the appropriate boundary conditions for the normal component of $\mathbf{J}_i$ on all boundaries has proven to be quite difficult because the physical interpretation of the boundary condition is not intuitive. Large, time-dependent currents do flow into the cathode, and simplified assumptions produce large errors. Even rough estimates are difficult to specify with realistic

cathode arrangements, such as the one shown in Fig. 2-64.

We have made recent progress with another approach to this equation. Instead of solving Eq. (123), we consider the related equation

$$\nabla^2 \xi + \frac{4\pi}{c^2} \mathbf{j} \cdot \nabla \xi = 0 \quad (133)$$

We solve Eq. (133) using boundary conditions that require $\xi = E_r$ on the boundary and then decompose ξ to obtain

$$\mathbf{E}_r = \frac{E_r}{r} \mathbf{r} \quad (134)$$

Iteration is still required because E_r is needed for the right-hand side of Eq. (133). Since decomposition is not required on the right-hand side of Eq. (133), we may speed the convergence by moving some of the explicit dependence on E_r to the left-hand side. Displaying explicitly the functional dependencies, we have

$$\nabla^2 \xi^{(n+1)} = \frac{4\pi}{c^2} \mu \xi^{(n)} + \frac{4\pi}{c^2} (\mathbf{K} + \mu \mathbf{E}_r + \zeta \times \mathbf{B} - \mu \xi^{(n)}) \quad (135)$$

At convergence, the solution is obtained by a divergence-cleaning of ξ to obtain E_r . The new iteration for E_r consists of the following steps:

1. Form the right-hand-side of Eq. (135).
2. Solve Eq. (135) for $\xi^{(n+1)}$ using E_r Dirichlet conditions on ξ .
3. Decompose ξ by specifying ξ_r on the boundaries. Since the properties of ξ_r are irrelevant except for the boundary conditions, we use the zero condition to preserve the desired E_r boundary condition.
4. Check convergence of ξ . If inadequately converged, return to Step 1.

No decomposition of $\mathbf{j}$ is required, thus obviating the need for the awkward boundary condition on $\mathbf{j}_r$. We have in its place a requirement to decompose ξ with boundary conditions no more arduous than specifying E_r on the boundaries—a considerable improvement. We are presently incorporating this new system of equations into the DPC code.

Authors: D. W. Hewett and J. J. Boyd

Fig. 2-64. Typical axisymmetric r - z region for a DPC simulation of an electron gun showing the cathode and anode placement.

Differencing the Electron-Electron Collision Operator in Low-Z Transport

It is frequently desirable that macroscopic thermodynamic properties of the electron transport equation be retained even when zoning is not sufficiently fine to produce high accuracy in all respects.¹²⁵ In Ref. 126, the first two angular moments in velocity of the electron Fokker-Planck collision integral in Landau form are expressed in difference equations that retain almost all of the following properties of the Landau term:

- Conservation of density, momentum, and energy.
- Galilean invariance to first order.
- The Maxwellian is a stable equilibrium.
- Preservation of Onsager symmetries near equilibrium.

These properties are algebraically exact for general zoning in energy but may be limited in practice by the accuracy with which the equations can be solved. The matrix equations to be solved are simpler than would be obtained from a finite-element-method formulation.

Authors: A. Decoster (CEA) and A. B. Langdon (LLNL)

LASNEX and Atomic Physics

Introduction

During 1985 we addressed several problem areas in the LASNEX model for which we felt that significant improvements could be made. The diffusion algorithm used in LASNEX to model electron and ion heat conduction, radiation flow, and neutron and charged-particle transport is successful in meshes that are not highly skewed.

However, in meshes with short-wavelength distortion, such as those in the advanced stages of hydrodynamic instability growth, observed diffusion fluxes are incorrect by factors of 2 to 5.

To obtain better results, we have developed a finite-element-based diffusion model that is now running in LASNEX. Test problems have shown that the model corrects the problem of diffusion through distorted meshes and gives more accurate answers (for a given mesh size) on problems with meshes that are less skewed.

We use the XSN atomic model in LASNEX to generate opacities for radiation transport. This average-atom model makes several crude simplifying assumptions about atomic structure to quickly obtain qualitatively correct values. For certain applications, such as those involving x-ray lasers and diagnostics that measure atomic line radiation, we must simulate line transport. Thus, it is important to use a more realistic and accurate atomic model even at the expense of an increase in run time by a factor of 10. We have placed the newly developed detailed-configuration atomic model into LASNEX to fill this need.

The discrete-ray method and spherical-harmonic-radiation transport models currently in LASNEX simulate only nonrelativistic-limit Compton scattering (Thomson scattering) in which the photon changes its direction without a change in photon energy. To model relativistic Compton scattering, we require the Compton scattering kernel (the average of the Klein-Nishina cross section over a relativistic Maxwellian distribution). We decided not to use the Fokker-Planck approximation (which assumes that the photon frequency shift is very small compared to photon frequency) because the approximation has been shown to be inadequate for many design problems of interest. A survey of the literature revealed that the available methods for calculating the Compton scattering kernel were slow and inaccurate. We have developed a new method that is much faster (12.8 μ s per cross section on a Cray XMP) and has guaranteed accuracy throughout the parameter space.

LASNEX has both a 1-D and a 2-D upwind, implicit, discrete-ray scheme for solving the radiation-transport equation.

A persistent problem with these schemes has been their inaccuracy in the thermal wave limit when the thickness of zones is many mean free paths. We have been using the Wilson-Lund-Castor factor, which was proposed as a cure for this problem; however, the factor has proved to be inadequate in real design problems. Later in this section, we present a new upwind, implicit, discrete-ray scheme, which behaves correctly in the thermal wave limit.

To maintain and improve LASNEX, we must keep track of 80 000 lines of code that are modified by a dozen individuals located in different buildings. We load subroutines from a dozen different libraries and use two different compilers. We maintain equivalent versions of the code on five Cray computers. All versions of the code are archived, and records are maintained of all changes. The following article, "LASNEX Code Development System," describes how we have automated the procedures necessary to make changes in the code.

Author: D. S. Kershaw

LASNEX Code Development System

The LASNEX Code Development System is a set of utilities that aids computational physicists and computer scientists in making changes to LASNEX. The utilities automate, when possible, the creation, coordinating, updating, checking, and archiving of new versions of LASNEX. The utilities are:

- **QUE**, which creates new versions of LASNEX.
- **GETBOOK**, which reserves public libraries for updating.
- **UPDATE**, which updates public libraries.
- **HERMES**, which checks that all public libraries are consistent.
- **STASHEM**, which archives compiler listings and libraries.

Our code development is a complex task due to the number of people working on LASNEX and the size of the code. Five physicists working in two buildings and five computer scientists are currently committed to LASNEX full time. More than

70 000 lines of LRLTRAN, which is an extended Fortran; more than 500 subroutines, functions, and cliches; and more than 10 subroutine libraries are used in LASNEX.

Typical Change to LASNEX. To illustrate how the utilities are used to introduce a change to LASNEX, we offer a typical example starting from problem detection and ending with user availability of a corrected LASNEX code. Suppose that on line 130 in subroutine XXXX, the variable YYYY should be changed to ZZZZ. To change the variable, one of the LASNEX computational physicists or computer scientists would add the following lines to the alter files:

```
\ALTER. XXXX PUB.  
RP130!YYYY!ZZZZ
```

To create a version of LASNEX with this change, the user would type:

```
LASLIB QUE PUB. SUBS. XXXX / t v
```

QUE produces a new LASNEX version, which is tested to determine whether or not a problem has been corrected.

After determining that the change has been correctly made, the new version of LASNEX becomes available to all users by typing:

```
LASLIB GETBOOK PUB. / t v
```

GETBOOK reserves the public libraries that are needed to obtain an update. A reservation is required to ensure that two users do not simultaneously update the public libraries.

Once the libraries are reserved, the user would type:

```
LASLIB UPDATE / t v
```

The UPDATE utility updates the LASNEX group's public libraries so that the changes are available to the next user of LASNEX.

At 30-min intervals, the HERMES utility runs on all Crays to check that the LASNEX public files are correct. If a public file is incorrect, HERMES requests a correct file from another machine and uses the correct file as a replacement.

At midnight, the STASHEM utility is run on all Crays to archive the LASNEX public libraries, compiler-listing files, alter files, source-code files, and binary files. The files are sent to a special LASNEX account by UPDATE, when an update is performed, so that STASHEM can archive them later.

QUE Utility. The QUE utility is the compilation controller. LASNEX code developers use QUE to compile and load a new version of LASNEX with their changes added. QUE knows the names and location of the LASNEX cliche files, binary files, alter files, source files, loader files, and subroutine libraries that are needed to compile and load a new version of LASNEX. The new version is tested until the code developers are satisfied that they have corrected a problem or added an enhancement. If their testing shows that further refinement is needed, they change the alter file and rerun QUE. The process is repeated until changes are complete.

QUE has the following options that the user can specify:

- Which subroutine, function, or cliche to alter.
- Which code to load (public, experimental, or both).
- Which compiler to use
- What type of compiler listings should be produced.
- To which output device the compiler listings should be diverted.

The user can choose to load LASNEX without altering a module or alter a module without loading.

QUE maintains a log, referred to as a time stamp, for every file that is used to create the new version of LASNEX. A time stamp is the date and time at which the file was last modified. The time stamp log is passed to UPDATE to verify that all previous changes by other code developers are accounted for and contained in the current version.

QUE produces a library file, named lasfiles, which contains all files needed for a public update. The lasfiles library contains the:

- New version of LASNEX.
- Loader map.
- Altered binaries.
- Time stamp file.
- Alter files.
- Altered source files.

GETBOOK Utility. In the past, a book was required for a user to update a LASNEX public library. The GETBOOK utility was created because LASNEX group members are currently located in two buildings. GETBOOK replaces the book and sets or releases a reservation on the LASNEX public libraries. The reservation must be set to update any LASNEX public library and ensures that two people do not simultaneously update that public library.

GETBOOK allows the user to:

- Reserve a single public library.
- Reserve all libraries to update a public or experimental code.
- Learn who has reserved a specific library.
- Release the user's own reservations. (Other users' reservations cannot be released.)
- Obtain a help package (for novice users).

GETBOOK tracks reservations by using an XPORT directory. If a library is not reserved, the directory is

.C:BOOK:library

If a library is reserved, the directory is

.C:BOOK:library:RESERVED:WHOmmuuuuu

where "library" is the reserved library, "m" is the machine on which the library was reserved, and "uuuuu" is the user number of the individual who reserved that library.

UPDATE Utility. The UPDATE utility integrates the code developer's changes into the LASNEX public libraries by performing the following steps:

1. Determining whether a user is a LASNEX code developer. If not, UPDATE quits immediately.
2. Checking for reservations on the libraries that are needed to update. If the libraries are not reserved, UPDATE quits immediately.
3. Verifying that current changes do not invalidate changes made by other group members by checking information in the time stamp file against the times and dates in the public libraries. If the current update is in conflict, UPDATE quits immediately.
4. Requesting a history file to determine what was changed and the reason for the change.

5. Merging the binary, alter, executable, and loader files.

6. Copying the new libraries into public libraries.

7. Giving a copy of the libraries, compiler listings, and altered sources to LASNEX user number 489063. STASHEM then archives these files at midnight.

HERMES Utility. The HERMES utility runs every 30 min on each of the Cray computers and checks all LASNEX public libraries for the proper checksums. If one of the libraries contains an erroneous checksum, HERMES requests a correct library from HERMES on another machine and copies the correct library into public.

HERMES also creates a file of times that the libraries were last modified. This file is sent to another HERMES, and the files on both machines are compared. HERMES determines which public library is the most recent, then transports and copies the new library into public.

STASHEM Utility. The STASHEM utility runs on each of the Cray computers and archives the files that were sent to the special LASNEX user number 489063. (Recall that the files were sent to user 489063 when UPDATE was run.) At present, the following files are sent:

- Renamed copies of the LASNEX public libraries.
- Compiler listings.
- Altered sources.

STASHEM archives the public libraries in a directory structure. The naming convention of this directory structure informs the user of the library, the machine (Cray 1/S or Cray XMP), and the date. The compiler listings are archived in a directory structure that informs the user of which subroutine, function, or cliché was altered; the public or experimental code; and the date it was altered. The altered sources are updated at midnight into a library of current LASNEX sources.

Summary. Automation in the processes of creating, coordinating, updating, checking, and archiving new versions of LASNEX has substantially reduced the time needed to incorporate changes. Built-in checks of the automated systems allow many users to work in parallel and to increase overall productivity.

Author: L. M. Stringer

Simple Evaluation of the Compton Scattering Kernel for Radiative Transfer

The radiative transfer equation is, in general, an integral equation with a kernel given by the photon-matter scattering cross section.¹²⁷ For radiation interacting with a relativistic, nondegenerate gas of free electrons, this kernel is the Compton scattering kernel (CSK). The CSK is defined as the Klein-Nishina cross section averaged over a relativistic Maxwellian distribution. The CSK has been studied by various authors, originating with Stone and Nelson,¹²⁸ in the form of a complicated multiple integral. The integral is then evaluated numerically by means of Gaussian quadrature.^{129,130} The complicated nature of the CSK has led some to a perturbative solution of the radiative transfer equation, known as the Fokker-Planck treatment of scattering.¹³⁰⁻¹³² The treatment is valid for photon energies and electron temperature $\ll$ electron rest energy.

We have been motivated by the need to implement Compton scattering in the radiation transport packages of LASNEX; we have found that the CSK can be analytically reduced to a single integral that can then be rapidly evaluated in a variety of ways. This article presents what is, to our knowledge, the first correct computation of the CSK. Our calculations have uncovered significant errors in the Stone-Nelson numerical evaluations.

The quantum mechanical cross section for the scattering of a photon by a free electron moving with velocity $\mathbf{v}$ is given by the Klein-Nishina formula:

$$\sigma(\gamma \rightarrow \gamma', \mathbf{\Omega} \rightarrow \mathbf{\Omega}', \mathbf{v}) = \frac{r_0^2}{2\gamma v \lambda} \times \left[1 + \left(1 - \frac{1 - \xi}{\lambda^2 D D'} \right)^2 + \frac{(1 - \xi)^2 \gamma \gamma'}{\lambda^2 D D'} \right] \times \delta \left(\xi - 1 + \lambda \frac{D}{\gamma} - \lambda \frac{D'}{\gamma} \right), \quad (136)$$

where

$$D \equiv 1 - \mathbf{\Omega} \cdot \frac{\mathbf{v}}{c}, \quad (137)$$

$$D' \equiv 1 - \mathbf{\Omega}' \cdot \frac{\mathbf{v}}{c}, \quad (138)$$

$$\lambda \equiv \left(1 - \frac{v^2}{c^2} \right)^{1/2}, \quad (139)$$

$$\gamma \equiv \frac{h\nu}{m_0 c^2}, \quad (140)$$

$$\gamma' \equiv \frac{h\nu'}{m_0 c^2}, \quad (141)$$

$$\xi \equiv \mathbf{\Omega} \cdot \mathbf{\Omega}', \quad (142)$$

$\gamma, \mathbf{\Omega}(\gamma', \mathbf{\Omega}')$ is the incoming (outgoing) photon energy (in units of electron rest energy $m_0 c^2$) and direction, respectively, and r_0 is the classical electron radius. We follow the notation of Ref. 127 in Eqs. (136) through (145).

The Dirac delta function in Eq. (136) enforces energy-momentum conservation for the scattering process. The Maxwellian distribution for a gas of free, relativistic, nondegenerate electrons at temperature T_e is given by

$$f(v) = \frac{\lambda^3 e^{-\lambda/\tau}}{4\pi\tau c^3 K_2(\tau^{-1})}, \quad (143)$$

where

$$\tau \equiv \frac{kT_e}{m_0 c^2}, \quad (144)$$

K_2 is the modified Bessel function of the second kind, and k is the Boltzmann constant. The CSK is defined as the average of the Klein-Nishina cross section in Eq. (136) over the Maxwellian distribution in Eq. (143):

$$\sigma_s(\gamma \rightarrow \gamma', \xi, \tau) = \frac{N_e r_0^2}{2\gamma v} \int f(v) \frac{1}{\lambda} \times \left[1 + \left(1 - \frac{1 - \xi}{\lambda^2 D D'} \right)^2 + \frac{(1 - \xi)^2 \gamma \gamma'}{\lambda^2 D D'} \right] \times \delta \left(\xi - 1 + \lambda \frac{D}{\gamma} - \lambda \frac{D'}{\gamma} \right) d\mathbf{v}, \quad (145)$$

where N_e is electron density.

To evaluate Eq. (145), previous authors have taken the z axis of integration along the incoming photon direction $\mathbf{\Omega}$, leading to a rather intractable expression. Instead, we take the z axis of integration along the

photon momentum transfer, such that

$$\int (d\mathbf{v}) = \int_0^{\tau} v^2 dv \int_1^{\tau} d(\mathbf{n} \cdot \mathbf{\Omega}_e) \int_0^{2\pi} d(\phi_e), \quad (146)$$

where

$$\mathbf{v} \equiv v \mathbf{\Omega}_v, \quad (147)$$

$$\mathbf{n} \equiv \frac{(\gamma' \mathbf{\Omega}' - \gamma \mathbf{\Omega})}{q}, \quad (148)$$

and

$$q \equiv \sqrt{\gamma^2 + \gamma'^2 - 2\gamma\gamma'\xi}. \quad (149)$$

The Dirac delta function in Eq. (145) then fixes $\mathbf{n} \cdot \mathbf{\Omega}_v$; the requirement $|\mathbf{n} \cdot \mathbf{\Omega}_v| < 1$ along with $v > 0$ implies that

$$\begin{aligned} \lambda \geq \lambda_+ &\equiv \frac{\gamma' - \gamma}{2} \\ &+ \left\{ \left[1 + \gamma\gamma' \frac{1 - \xi}{2} \right] \right. \\ &\left. + \left[1 + \frac{(\gamma - \gamma')^2}{2\gamma\gamma'(1 - \xi)} \right] \right\}^{1/2}. \end{aligned} \quad (150)$$

The remaining calculations are straightforward. The result is

$$\begin{aligned} \frac{\sigma_s(\gamma - \gamma', \xi, \tau)}{\Sigma_0} &= \Psi + \int_0^{\infty} e^{-\rho/\tau} \\ &\times \left\{ \left[\Lambda_+ - \frac{\rho + \rho_-}{\tau(1 - \xi)} \right] R_+^{-1/2} \right. \\ &\left. + \left[-\Lambda_- + \frac{\rho + \rho_-}{\tau(1 - \xi)} \right] R_-^{-1/2} \right\} d\rho, \end{aligned} \quad (151)$$

where

$$\begin{aligned} \rho_{\pm} &\equiv \pm \frac{\gamma + \gamma'}{2} + \left\{ \left[1 + \frac{\gamma\gamma'(1 - \xi)}{2} \right] \right. \\ &\times \left. \left[1 + \frac{(\gamma - \gamma')^2}{2\gamma\gamma'(1 - \xi)} \right] \right\}^{1/2}, \quad (152) \\ \Lambda_{\pm} &\equiv -\gamma\gamma' + \frac{2}{1 - \xi} + \frac{2}{\gamma\gamma'(1 - \xi)^2} \\ &- \left(\frac{\pm 1}{\tau} \right) \left(\frac{1}{\gamma} + \frac{1}{\gamma'} \right) \frac{1}{(1 - \xi)^2}, \end{aligned} \quad (153)$$

$$\Sigma_0 \equiv \frac{N_e r_0^2}{4\gamma v \tau K_2(\tau^{-1})} e^{-\lambda_+/\tau}, \quad (154)$$

$$R_{\pm} \equiv (\rho + \rho_{\pm})^2 + \omega^2, \quad (155)$$

$$\omega^2 \equiv \frac{1 + \xi}{1 - \xi}, \quad (156)$$

$$\alpha_{\pm} \equiv (\rho_{\pm}^2 + \omega^2)^{-1/2}, \quad (157)$$

and

$$\begin{aligned} \Psi &\equiv \frac{2\tau\gamma\gamma'}{q} + \frac{1}{(1 - \xi)^2} \left(\frac{1}{\gamma} + \frac{1}{\gamma'} \right) \\ &\times (\alpha_+ + \alpha_-) + \frac{(\rho, \alpha_+ - \rho, \alpha_-)}{1 - \xi}. \end{aligned} \quad (158)$$

The integral in Eq. (151) cannot be reduced further.

We can now make three observations. First, the right-hand side of Eq. (151) is manifestly symmetric in γ and γ' so that

$$\begin{aligned} \gamma^2 \sigma_s(\gamma \rightarrow \gamma', \xi, \tau) e^{-\gamma/\tau} \\ = \gamma'^2 \sigma_s(\gamma' \rightarrow \gamma, \xi, \tau) e^{-\gamma'/\tau}, \end{aligned} \quad (159)$$

which is simply a statement of detailed balance.

Second, there are delicate cancellations of terms in Eq. (151) for the following regions:

- $\xi \rightarrow 1$.
- $\gamma, \gamma' \rightarrow 0$ with γ/γ' fixed.
- γ fixed and $\gamma' \rightarrow 0$, and vice versa.
- $\tau \rightarrow 0$, which is important when evaluating σ_s numerically.

Third, the essential features of σ_s are contained in $\Sigma_0(\tau\gamma\gamma'/q)$, particularly in the factor $e^{-\lambda_+/\tau}$. Everything else is smoothly varying compared to this exponential. In fact, because $(1 - \xi)(\lambda^2 DD')^{-1}$ is 1 - the cosine of the photon scattering angle in the electron rest frame, the expression within the brackets of Eq. (145) is bounded between 1 and $(2 + 4\gamma\gamma')$, which implies

$$1 \leq \frac{\sigma_s(\gamma - \gamma', \xi, \tau)}{\left[\Sigma_0 \frac{\tau\gamma\gamma'}{q} \right]} \leq (2 + 4\gamma\gamma'). \quad (160)$$

The factor $e^{-\lambda_+/\tau}$ has a simple physical explanation. The term $\lambda_+ m_0 c^2$ is the

smallest electron energy for which a photon scattering with given γ, γ' and ξ is possible. For this scattering to occur, an electron of probability $e^{-\lambda/\tau}$ or lower must be sampled. It follows that it is highly inaccurate to expand the ξ dependence of σ_s in Legendre polynomials, as has been done by a several authors,¹²⁷ because one can choose physically relevant values of γ, γ' and τ for which the factor $e^{-\lambda/\tau}$ changes by many orders of magnitude with arbitrarily small variations in ξ . In particular, for $\gamma, \gamma' \gg \tau$, the factor $\Sigma_s \tau \gamma \gamma' / q$ approximates the Dirac delta function in Eq. (136) with $\mathbf{v} \rightarrow 0$.

Some of the many ways of numerically computing Eq. (151) can be found in Ref. 133. Our fastest method of computing σ_s for a given γ, γ', ξ , and τ requires 12 μs on a Cray XMP for an accuracy of 1 part in 10^3 .

Authors: D. S. Kershaw and M. K. Prasad

Radiation Diffusion with the Finite-Element Method

We have developed a numerical scheme based on the finite-element (FE) method to solve the multigroup radiation-diffusion equations in LASNEX. The new method gives significantly better results on problems with irregular meshes. Particular applications are grids occurring in the later stages of a Rayleigh-Taylor instability.

In this article, we provide a brief description of the application of FE and its compliance with existing LASNEX requirements. We present two examples to demonstrate the performance of FE. The application of the method in LASNEX will be described more fully in an LLNL report to be published soon.

The scheme that is presently implemented time advances the photon energy density u_ν and the electron temperature T_e by the partial-temperature method. In this formulation, T_e is advanced in N_ν steps, where N_ν is the number of frequency groups. During the i th step, T_e and $u_{\nu i}$ are advanced using fully implicit methods. The relevant equations are

$$\partial_t u = \nabla \cdot (D \nabla u) + \gamma [B(T) - u] + P_1 \quad (161)$$

and

$$c_v \partial_t T = -\alpha \Delta v \gamma [B(T) - u] + P_2, \quad (162)$$

where $D = cl_v/3$ and $\gamma = c/l_v$ are functions of the frequency-dependent mean free path l_v , $B(T_e)$ is the Planck function, $\alpha = 0.00211$ is a normalizing constant, and c is the speed of light.

In Eqs. (161) and (162) and in the following discussion, we omit the subscripts ν and e for brevity. Physical processes such as radiation PdV work and Compton scattering (denoted by P_1) are solved by operator splitting. Processes such as electron-ion coupling P_2 are solved as before. Their solution is not described here, but the effects are incorporated in the FE scheme.

The LASNEX mesh coordinates are labeled by logical K, L variables: $Z(K, L)$, $R(K, L)$. In (K, L) space, the grid is rectangular; in (Z, R) space, the mesh consists of quadrilaterals. Azimuthal symmetry is assumed. The finite-difference (FD), diffusion package computes average zonal variables in each quadrilateral.¹³⁴ In contrast, the FE package computes nodal values of u and T . Variables such as the specific heat c_v and l retain their zonal definitions. Each quadrilateral is bisected into triangles; u and T are piecewise linear functions over the entire domain and continuous across the triangle sides.

The FE scheme is more accurate, and greater flexibility is allowed for the boundary conditions. At the boundaries, the user defines one of the following:

- The radiation temperature, hence u .
- The radiation flux $-D \partial u / \partial n$.
- A Milne condition $u + (2/3) \partial u / \partial n$,

where n is the unit outward normal. Reflection or symmetry conditions are specified using null values for the flux.

As applied, the FE equations conserve the same quantities as do the differential equations. Precise analogues are computed for the change in radiation energy, the energy transported, the energy coupled to the electrons, etc. A limitation on the energy check is imposed by the error tolerance prescribed to the linear system solver ICCG (Ref. 135). Experience has shown that an error criterion too crude leads to poor energy conservation.

Because the FE solution of Eqs. (161) and (162) is distinct from that of the FD

method, we have made considerable effort to incorporate FE within the existing LASNEX Lagrangian architecture. For example, zonal u and T values are computed a posteriori by integrating the piecewise linear functions over a quadrilateral. Changes in the zonal T values (e.g., from electron-ion coupling) are mapped to points. The radiation energy density is advected by the fluid.

The FE method is sensitive to the grid geometry: in particular, triangles with large obtuse angles may generate negative energy densities. Therefore the mesh

movement is monitored, and, if necessary, quadrilaterals are retriangulated by an algorithm using Lagrange multipliers that minimizes numerical diffusion and conserves energy. However, cases in which quadrilaterals "bow tie" give unphysical results. Details are given in a forthcoming LLNL report.

We provide two examples that demonstrate the FE and compare it with the analytic solution. In both cases, the mesh is fixed but "random." All distances are given in centimeters; time is measured in units of $t_0 = 10^{-8}$ s; hence, $c = 300$. Temperatures and frequencies are in units of keV. The purpose is to determine what effect the mesh has on expected symmetry of the solution.

In our first example, $\gamma = 0$ and $l = 0.005$, i.e., only linear radiation diffusion is modeled. At $t = 0$, $u = 0$. Boundary conditions are $\partial u / \partial n = 0$ along $K = 1$ and $K_{\max}$, and $u + (2/3)\partial u / \partial n = 0$ along $L = L_{\max}$. Along $L = 1$, the radiation temperature is set to unity. For this problem, $K_{\max} \sim L_{\max} = 33$. Only one frequency group is used. The steady-state numerical solutions and the outgoing flux S (at $L = L_{\max}$) are compared with the analytic solution. The problem is spherically symmetric. If $X^2 = R^2 + Z^2$, the domain is $X(L = 1) = X_0 = 1$, $X(L = L_{\max}) = X_1 = 2$, and $0 \leq \tan(R/Z) = \pi/2$. The steady-state solution is $u = -(a_2/D) \times X^{-1} + b_2$, where a_2 and b_2 are constants. The flux $S \approx X^{-2}$.

Figure 2-65 shows the domain and grid, while Fig. 2-66 shows the fluxes along X_1 for both the FD and FE methods; they should be compared with the exact value $S = 0.0034193$. Note the greater variation for the FD method. This problem demonstrates the better accuracy obtained by FE when it is used to approximate the diffusion term $\nabla \cdot (D \nabla u)$ on irregular grids. Indeed, the power of FE was first demonstrated by its ability to solve elliptic equations on arbitrary domains using triangular grids.¹³⁶

An analogous slab-symmetric problem was run on a similar random mesh. The steady-state solution is linear in Z ; hence, FE computes the exact solution and the exact outgoing flux. For this case, FE is a dramatic improvement over the FD method because grid effects are quite

Fig. 2-65. Domain and grid for spherically symmetric problem.

Fig. 2-66. Spherically symmetric problem. Outgoing flux vs K along $X = 2$.

noticeable in the latter solution and are absent in the former.

Our second problem is a typical thermal wave. It includes electron coupling; the Planck function introduces nonlinearity. The analytic solution is slab-symmetric. Consider one frequency group, constant l and c_v , and neglect P_r in Eqs. (161) and (162). The analytic equations are

$$\partial_t u = D \partial_{zz} u + \gamma [B(T) - u] \quad (163)$$

and

$$c_v \partial_t T = -\alpha v_x \gamma [B(T) - u], \quad (164)$$

where

$$v_x B(T) = \int_0^{T^*} v^3 dv / (e^{v^2} - 1). \quad (165)$$

We solve Eq. (164) for u , differentiate, and substitute into Eq. (163).

We assume:

- Strong coupling, $l = \epsilon \ll$ problem dimension.
 - The characteristic time $\tau \gg t_0$.
 - Small (but not too small) specific heat ratio.
- Since

$$v_x B(T) = T^4 \int_0^{T^*} x^3 dx / (e^{x^2} - 1) \quad (166)$$

for large v_x (≈ 20) and $T \leq 1$,

$$\int_0^{T^*} x^3 dx / (e^{x^2} - 1) \approx \int_0^x \pi^4/15. \quad (167)$$

The integral has units of σ/c , where σ is the Stefan-Boltzmann constant. Thus, $v_x B(T) \approx T^3$.

With these assumptions, Eqs. (163) and (164) reduce to

$$\begin{aligned} \partial_t T &= \frac{\alpha v_x}{c_v} D \partial_z [B'(T) \partial_z T] \\ &= a \partial_z (T^3 \partial_z T), \end{aligned} \quad (168)$$

where $a = 4(\alpha/c_v)(cl/3)\pi^4/15$ is a constant.

Fig. 2-67. Domain and grid for thermal wave problem. The grid is assembled from two random meshes.

The program is tested by solving Eqs. (161) and (162) on the domain: $1 \leq R \leq 2$, $0 = Z_0 \leq Z \leq Z_2 = 2.09375$, $K_{\max} = 33$, and $L_{\max} = 68$. The grid displayed in Fig. 2-67 is assembled from two random meshes: one for $0.0625 \leq Z \leq Z_1 = 1.0625$, and another for $1.09375 \leq Z \leq Z_2$. The numerical flux is computed along the straight grid line $Z = Z_1$.

The solution for a plane source $T(t=0) = Q\delta(z)$ is self-similar, and is given analytically.¹³⁷ It is characterized by a front position $z_f(t) \approx t^{1/5}$ and the value at the symmetry plane $T_c \approx z_1^{-1}$. We choose $l = 0.001$ and $c_v = 50$; hence, $D = 0.1$, $\gamma = 3 \times 10^5$, and $a = 1.096 \times 10^{-4}$. Symmetry conditions ($\partial u / \partial n = 0$) are imposed along $R = 1$ and 2 and along Z_0 . Milne conditions are imposed along $Z = Z_2$.

The numerical solution loses physical significance when the front arrives at $Z = Z_2$. At $t = 0$, T and u are nonzero only along $Z = Z_0$. Initially, the functions are in equilibrium, are independent of R , and $B(T^0) = u^0$. Setting $T = 1$ along $Z = Z_0$ gives $Q = 0.0312586$.

Figures 2-68 and 2-69 compare the FE with the analytic solution. Figure 2-68 shows temperature vs Z along $R = 1$ at $t = t_{\max} \approx 2.019 \times 10^9$ s. The profile agrees with the analytic function. Figure 2-69 plots the time history of the flux across $Z = Z_1$. The numerical results of Fig. 2-69 display the radiation flux vs time, where

$$f = -2\pi\alpha \int_1^2 R dR \int_0^{T^*} D \nabla u dv. \quad (169)$$

Fig. 2-68. Thermal wave problem. Temperature profile vs Z along $R = 1$ at time $t = 2.019 \times 10^9$.

Fig. 2-69. Thermal wave problem. Radiation flux vs time at $Z = 1.0625$.

The assumptions following Eq. (164) imply $u \approx B(T)$. Using Eq. (166) and the definitions of D ,

$$f = -4\pi\alpha(\pi^4/15)ctT^3 \partial_z T. \quad (170)$$

The flux corresponding to Eq. (168) is $S = -aT^3 \partial_z T$. Thus, $f = 3\pi c_v S$. Along $Z = Z_1$, the analytic flux peaks at $t = 7.522 \times 10^8$ s, and $S_{\max} = 2.1751 \times 10^{-12}$. The FE solution $f_{\max}$ is within 3%. For this problem, FD also compares well with the analytic solution. We hypothesize that the strong coupling in this problem, $\gamma \gg D$, reduces the importance of accurately computing the transport term.

In conclusion, we expect the FE scheme to provide better results for problems with mesh distortions due to Lagrangian hydrodynamics. However a caution is in order: the FE scheme is based on triangles, hence FE will not maintain the exact symmetry

that the FD scheme preserved when that scheme was used on orthogonal meshes.

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A Correct 1-D, Upwind, Discrete-Ray-Method, Radiation Transport Algorithm in the Strong Coupling Limit

Introduction. Upwind, implicit, discrete-ray-method schemes for solving the radiation transport equation have many advantages. The advantages include positivity, stability, algebraic energy conservation, and upper-triangular linear systems that are trivial to solve. A persistent problem with such schemes has been their inaccuracy in the thermal wave limit when zones are many mean free paths thick. The Wilson-Lund-Castor factor,¹³⁸ proposed as a cure for this problem, has proved to be inadequate in real design problems.

In this article, we present a new upwind, implicit, discrete-ray-method scheme that has all the above-mentioned virtues and behaves correctly in the thermal wave limit as well. We show how the thermal wave limit works in the analytic case and then describe the new scheme and study its thermal wave limit.

Thermal Wave Limit: An Analytic Case. For simplicity of exposition, we assume grey opacity and use a single frequency group. Frequency dependence is a nonessential complication when studying the thermal wave limit. In the thermal wave limit, the radiation field is in steady state, and the coupled electron temperature and radiation equations are

$$c\mu \frac{\partial I}{\partial z} = \frac{\eta}{\epsilon} (S - I) \quad (171)$$

and

$$C \frac{\partial T}{\partial t} = -\frac{1}{2} \int_{-1}^1 d\mu \frac{\eta}{\epsilon} (S - I), \quad (172)$$

where c is the speed of light, μ is the cosine of the angle between the photon direction and the z axis, I is the radiation

intensity, $\eta = c\rho\kappa$ is the coupling coefficient, $S = aT^4$ is the source function, C is the electron specified heat, T is the electron temperature, and t is time. The term ϵ is a formal parameter set to 1 at the end of the calculation.

In the thermal wave limit, η is large, thus we assume that I may be expanded in powers of ϵ such that

$$I = \sum_{n=0}^{\infty} \epsilon^n I_n. \quad (173)$$

By equating terms of the same order in ϵ in Eq. (171), we find

$$I_0 = S, \quad (174)$$

$$I_n = -\frac{c\mu}{\eta} \frac{\partial I_{n-1}}{\partial z}, \quad (175)$$

$$I_1 = -\frac{c\mu}{\eta} \frac{\partial S}{\partial z}, \quad (176)$$

$$I_2 = (c\mu)^2 \frac{1}{\eta} \frac{\partial}{\partial z} \frac{1}{\eta} \frac{\partial S}{\partial z}, \quad (177)$$

and so forth. We see from examination of the individual I_n terms that our expansion is really an expansion in powers of the radiation mean free path divided by the gradient length of the source.

Substituting this into Eq. (172), we find to lowest order in ϵ (the second order)

$$C \frac{\partial T}{\partial t} = \epsilon^2 \frac{\partial}{\partial z} \frac{c^2}{3\eta} \frac{\partial S}{\partial z}, \quad (178)$$

which is the correct thermal wave equation (setting $\epsilon = 1$).

Computational Algorithm. If we assume some specific form for the source function throughout a zone, we can calculate the intensity at each node exactly for each ray direction. We assume that η is constant within each zone. We label the zone boundaries $z_k (z_{k+1} > z_k)$, and the zone center points $z_{k+1/2}$. We choose a set of discrete ray directions $-1 < \mu_j < +1$, with $j = 1, 2, \dots, 2n$, with $\mu_j = -\mu_{2n-j+1}$, and ray weights $\delta\mu_j$, satisfying $\sum \delta\mu_j = 2$,

$\sum \mu_j \delta\mu_j = 0$, and $\sum \mu_j^2 \delta\mu_j = 2/3$. Then for a given ray direction μ_j and for $z_{k+1} > z > z_k$, we find by integrating Eq. (171) that

$$I_j(z) = e^{\left(\frac{z - z_k}{\tau_{j,k+1/2}}\right)} I_{j,k} + \frac{1}{\tau_{j,k+1/2}} \times \int_{z_k}^z dz' e^{\left(\frac{z - z'}{\tau_{j,k+1/2}}\right)} S(z') \quad (179)$$

if $\mu_j > 0$, and

$$I_j(z) = e^{\left(\frac{z_{k+1} - z}{\tau_{j,k+1/2}}\right)} I_{j,k+1} + \frac{1}{\tau_{j,k+1/2}} \times \int_z^{z_{k+1}} dz' e^{\left(\frac{z - z'}{\tau_{j,k+1/2}}\right)} S(z') \quad (180)$$

if $\mu_j < 0$. Here

$$I_j(z) = I(\mu_j, z), \quad (181)$$

$$I_{j,k}(z) = I(\mu_j, z_k), \quad (182)$$

and

$$\tau_{j,k+1/2} = \frac{c\mu_j}{\eta_{k+1/2}}. \quad (183)$$

For a given source function, our solution for the intensity is exact (we have done the integration analytically). Therefore, if we make the representation of the source function within a zone accurate enough, we will certainly obtain the correct thermal wave limit. We find that a zero-order representation (S flat within a zone) is not good enough, but a linear representation works well. Therefore, let S_k be the value of the source at the k th zone boundary. Interior to a given zone, $z_{k+1} > z > z_k$, we represent S by

$$S(z) = S_k + (S_{k+1} - S_k) \left(\frac{z - z_k}{z_{k+1} - z_k} \right). \quad (184)$$

Substituting this linear fit into Eq. (179), we find that if $\mu_j > 0$,

$$I_j(z) = e^{-\xi} I_{j,k} + (1 - e^{-\xi}) S_k + \frac{(S_{k+1} - S_k)}{\alpha} (\xi - 1 + e^{-\xi}), \quad (185)$$

where

$$\xi = \frac{(z - z_k)}{\tau_{j,k+1/2}} \quad (186)$$

and

$$\alpha = \frac{(z_{k+1} - z_k)}{\tau_{j,k+1/2}}, \quad (187)$$

If $\mu_j < 0$, we find

$$I_j(z) = e^{-\xi} I_{j,k+1} + (1 - e^{-\xi}) \times S_{k+1} + \frac{(S_k - S_{k+1})}{\alpha} \times (\xi - 1 + e^{-\xi}), \quad (188)$$

where now

$$\xi = \frac{(z - z_{k+1})}{\tau_{j,k+1/2}}, \quad (189)$$

and

$$\alpha = -\frac{(z_{k+1} - z_k)}{\tau_{j,k+1/2}}, \quad (190)$$

In particular, we set $z = z_{k+1}$ in Eq. (185) and $z = z_k$ in Eq. (188) and obtain the following simple recursion formulae for calculating the zone boundary intensities:

$$I_{j,k+1} = e^{-\alpha} I_{j,k} + S_{k+1} - e^{-\alpha} S_k + \frac{S_k - S_{k+1}}{\alpha} (1 - e^{-\alpha}) \quad (191)$$

if $\mu_j > 0$, and

$$I_{j,k} = e^{-\alpha} I_{j,k+1} + S_k - e^{-\alpha} S_{k+1} + \frac{S_{k+1} - S_k}{\alpha} (1 - e^{-\alpha}) \quad (192)$$

if $\mu_j < 0$. The term α is defined as in Eqs. (185) and (188).

We have now discretized Eq. (171). To complete our scheme, we need a discrete version of Eq. (172). Because we know the behavior of $I_j(z)$ and $S(z)$ in the interior of a zone, we also know the emission minus the absorption, which is

$$\frac{1}{2} \sum_j \delta\mu_j \eta(z) [S(z) - I_j(z)] \quad (193)$$

everywhere.

To discretize Eq. (171), we simply attribute all the electron-radiation coupling between $z_{k-1/2}$ and $z_{k+1/2}$ to the electron temperature change at zone boundary k . Thus we have for the discrete version of Eq. (172),

$$(z_{k+1/2} - z_{k-1/2}) C_k \frac{\partial T_k}{\partial t} = -\frac{1}{2} \sum_j \delta\mu_j \int_{z_{k-1/2}}^{z_{k+1/2}} \eta(z) \times [S(z) - I_j(z)], \quad (194)$$

where $\eta(z) = \eta_{k+1/2}$ if $z_k < z < z_{k+1}$. Because Eqs. (185) and (188) are the exact integrals of Eq. (171), we can perform the z integral in Eq. (194) by integrating both sides of Eq. (171) to obtain

$$\int_{z_{k-1/2}}^{z_{k+1/2}} \eta(z) [S(z) - I_j(z)] = c\mu_j [I_j(z_{k+1/2}) - I_j(z_{k-1/2})], \quad (195)$$

Using Eqs. (195), (185), and (188) in Eq. (194) we find

$$(z_{k+1/2} - z_{k-1/2}) C_k \frac{\partial T_k}{\partial t} = -\frac{1}{2} \sum_{\mu_j > 0} \delta\mu_j (c\mu_j) \times \left\{ e^{\left(\frac{\alpha_{j,k+1/2}}{2}\right)} (I_{j,k} - I_{2n-j+1,k+1}) + \left[\frac{2}{\alpha_{j,k+1/2}} - \left(\frac{2}{\alpha_{j,k+1/2}} + 1 \right) e^{\left(\frac{\alpha_{j,k+1/2}}{2}\right)} \right] \right\}$$

Detailed Configuration Accounting Kinetic Code for Non-LTE Plasmas

$$\times (S_{k-1} - S_k) - e^{\frac{\alpha_{j,k-1/2}}{2}} (I_{j,k-1} - I_{2n-j-1,k}) \\ - \left[\frac{2}{\alpha_{j,k-1/2}} - \left(\frac{2}{\alpha_{j,k-1/2}} + 1 \right) e^{\left(\frac{\alpha_{j,k-1/2}}{2} \right)} \right] \\ \times (S_{k-1} - S_k) \Bigg\}, \quad (196)$$

where the sum is over only positive values of μ .

In the strong coupling limit, when zones are many mean free paths thick, α becomes large so that the terms multiplying $e^{-\alpha/2}$ are negligible compared to the others. Neglecting these terms, Eq. (196) becomes

$$C_k \frac{\partial T_k}{\partial t} = - \frac{1}{2(\alpha_{k-1/2} - \alpha_{k+1/2})} \\ \times \sum_{\mu_j=0} \delta\mu_j(c\mu_j) \left[\frac{2}{\alpha_{j,k-1/2}} (S_k - S_{k+1}) \right. \\ \left. - \frac{2}{\alpha_{j,k+1/2}} (S_{k-1} - S_k) \right]. \quad (197)$$

Using

$$\alpha_{j,k-1/2} = \frac{(\alpha_{k-1} - \alpha_k) \eta_{k-1/2}}{(c\mu_j)} \quad (198)$$

and performing the sum on μ , we find

$$C_k \frac{\partial T_k}{\partial t} = \frac{c^2}{3(\alpha_{k-1/2} - \alpha_{k+1/2})} \\ \times \left[\frac{(S_{k+1} - S_k)}{\eta_{k+1/2}(\alpha_{k+1} - \alpha_k)} \right. \\ \left. - \frac{(S_k - S_{k-1})}{\eta_{k-1/2}(\alpha_k - \alpha_{k-1})} \right]. \quad (199)$$

This is the preferred difference scheme for Eq. (178). Our algorithm indeed recovers the thermal wave equation correctly in the strong coupling limit.

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Introduction. In the analysis of emission spectra from laser-produced plasmas and simulation of x-ray laser experiments, we need to know the charge-state distribution and the ionic level populations in a plasma. Such detailed information cannot be obtained from an average-atom model, such as XSNQ.^{1,39} Until now, XSNQ has been the only model available in LASNEX for non-local thermodynamic equilibrium (non-LTE) plasma simulation. Recently, we developed and installed in LASNEX a detailed configuration accounting kinetic code, called TDIB. This code calculates the time-dependent ionization balance and ionic level populations in non-LTE plasmas.

The TDIB package has two options for calculating these level populations. The default option treats all the ionization stages in a plasma using hydrogenic data. The other option treats certain ionization stages using detailed nonhydrogenic energy levels and rate coefficients.

In the default option, we assume that each energy level couples with all the levels of the same ionization stage, but only to the ground state of the next ionization stage through ionization and recombination processes. This approximation allowed us to develop a particularly efficient method of solving for the population.

We generate the ground-state and excited-state energy levels for each ion stage by using screening constants. There is one energy level per principal quantum number. The following atomic processes are included:

- Electron collisional excitation and deexcitation.
- Radiative emission and absorption.
- Electron collisional ionization and photoionization.
- Radiative recombination.
- Three-body recombination.
- Dielectronic recombination.

All the rate coefficients are scaled from the results for hydrogenic ions. Using the default option, TDIB is useful for calculating charge-state distribution and radiative power of arbitrary non-LTE plasmas.

In the second option, the code obtains the energy levels and both collisional and

radiative data for certain ionization stages from a data file. These data files, which are generated by the code ADAM,¹⁴⁰ are read by the x-ray laser design code XRASER.¹⁴¹ This option makes it possible to apply TDIB to analyze emission spectra of non-LTE plasmas and to simulate laboratory x-ray laser experiments.

TDIB is self-consistently coupled to LASNEX. We use the electron temperature, electron density, and photon density from LASNEX to compute the ionic level populations. The average ionization state, emissivity, absorption coefficient, and binding energy are calculated using the new ionic level populations. LASNEX uses these quantities for radiation transport, hydrodynamics, and electron energy balance to update the electron temperature, electron density, and photon density for the next time step.

Hydrogenic Approximation. In the hydrogenic approximation, the populations for the ground state N_i^1 and excited state N_i^l ($l > 1$) of the ion at its i th ionization stage are described by the following equations:

$$\begin{aligned} \frac{dN_i^l}{dt} = & -R_i(l)N_i^l \\ & + \sum_k \tilde{t}_i^{lk} N_i^k + S_{i-1}(l)N_{i-1}^1 \\ & + \sum_k [R_{i-1}(k)N_{i-1}^k - S_i(k)N_i^1] \delta_{l1}, \quad (200) \end{aligned}$$

where

$$\tilde{t}_i^{lk} = n_e C_{kl}^e + B_{kl} U_{hv}, \quad \text{when } l > k, \quad (201)$$

$$\tilde{t}_i^{lk} = n_e C_{kl}^d + A_{kl} + B_{kl} U_{hv}$$

$$\text{when } l < k, \quad (202)$$

$$\begin{aligned} \tilde{t}_i^{ll} = & - \sum_{k=1} (n_e C_{lk}^d + A_{lk} + B_{lk} U_{hv}) \\ & - \sum_{k=1} (n_e C_{lk}^e + B_{lk} U_{hv}), \quad (203) \end{aligned}$$

n_e is electron density, A_{kl} is the Einstein coefficient for spontaneous transition

between the two states k and l , and B_{kl} is the Einstein stimulated emission coefficient for a transition induced by radiation of energy density $U_{hv} d(h\nu)$ in the energy range $h\nu$ to $h\nu + d(h\nu)$.

The matrix $\tilde{t}_i$ couples all the energy levels within the i th ionization stage by collisional or radiative transition. The terms C_{kl}^e and C_{kl}^d are the collisional excitation and deexcitation rate coefficients, respectively, for transition between level l and level k . The term $R_i(l)$ represents the total ionization rate from level l , including electron collisional ionization and photoionization. The total recombination rate $S_i(l)$ includes radiative recombination, three-body recombination, and dielectronic recombination. The term $\sum_k R_i(k)N_i^k$ represents the increase of the ground-state population due to ionization. The term $S_i(l)N_i^l$ represents the decrease of the ground-state population due to recombination. These two terms exactly cancel each other when the plasma reaches steady-state equilibrium. At steady-state the total recombination rate of a given ion stage is equal to the total ionization rate of the lower ion stage. Although in this model each excited level is coupled to the ground state of the next ionization stage only, the coupling to the excited states of the next ionization stage is, in general, negligible for electron densities less than 10^{22} cm^{-3} .

The equation for excited-state level populations can be rewritten as

$$\begin{aligned} \frac{dN_i^l}{dt} = & -R_i(l) + \sum_k \tilde{t}_i^{lk} N_i^k \\ & + S_{i-1}(l)N_{i-1}^1 \\ & + n_e (C_{il}^e + B_{il} U_{hv}) N_i^l > 1, \quad (204) \end{aligned}$$

where the matrix $\tilde{t}_i$ couples the excited states. The coefficients C_{il}^e and B_{il} give the electron collisional excitation rate and the radiative absorption rate, respectively, for the transition from ground state to level l .

By solving Eq. (204), we obtain the following difference relation for the

excited-state level population:

$$(N_i^l)^{t+\Delta t} = (c_i^l)^t$$

$$+ \Delta t [a_i^l (N_i^l)^{t+\Delta t} + b_i^l (N_{i-1}^l)^{t+\Delta t}], \quad (205)$$

where Δt is the time step for the kinetic calculation. The term c_i^l is given in terms of the level populations. Both the coefficients, a_i^l and b_i^l are expressed in terms of the matrix $\tilde{r}_i$ and the ionization and recombination rate coefficients.

Substituting Eq. (205) into Eq. (200) gives

$$(N_i^l)^{t+\Delta t} = \frac{\Delta t [F_i^0 + F_i^1 (N_{i-1}^l)^{t+\Delta t} + F_i^2 (N_{i-1}^l)^{t+\Delta t}]}{1 - \Delta t F_i^2}, \quad (206)$$

where the coefficients F_i^0 , F_i^1 , F_i^2 , and F_i^3 depend on a_i^l , b_i^l , c_i^l , and the ionization and recombination coefficients. Equation (206) for the ground-state populations can be easily solved using any method for a tridiagonal matrix.

The energy levels are calculated using the screened hydrogenic model.¹⁴² The ionization potential for level l of the i th ionization stage is taken as

$$I_i^l = \epsilon_n^0 - \frac{Q_n^2 e^2}{2a_0 n^2}, \quad (207)$$

where

$$r_n = \frac{a_0 n^2}{Q_n}, \quad (208)$$

$$\epsilon_n^0 = \frac{e^2}{2r_n} \sigma(n, n) P_n, \quad (209)$$

Fig. 2-70. Contour plot of Z_{DCA}/Z_{AA} for an LTE cesium plasma in (T_e, ρ) space.

and

$$Q_n = Z_n - \sum_{m \geq n} \sigma(n, m) P_m - \frac{1}{2} \sigma(n, n) (P_n - 1). \quad (210)$$

The symbol a_0 is the Bohr radius; n is the principal quantum number for the valence electron of level l ; $P_n = 0, 1, 2, \dots, 2n^2$ is the number of electrons in the n th shell; and $\sigma(n, m)$ represents the hydrogenic screening constants. There is one energy level for each principal quantum number; the energy levels depend only on the principal quantum number n . Each level of an ion is described by a set of populations P_m . For example, the ground level of a neon-like ion is represented by $P_1 = 2$, $P_2 = 8$, $P_m = 0$, for $m \geq 3$. The first excited level is represented by $P_1 = 2$, $P_2 = 7$, $P_3 = 1$, $P_m = 0$, for $m \geq 4$. This model neglects all $\Delta n = 0$ transitions and assumes that the sublevels are in statistical equilibrium.

The coefficients for electron collisional excitation and ionization rate are scaled from the results of Golden and Sampson for hydrogenic ions.¹⁴³ The inverse rates

are calculated using the detailed balance relation. We use the Burgess and Merts formula to calculate dielectronic recombination rates for both the $\Delta n = 1$ and $\Delta n = 0$ transitions.^{144,145} In calculating the $\Delta n = 0$ transitions, we use the Post et al results for both the average excitation energy and oscillator strength.¹⁴⁶ The radiative recombination rate is computed with the classical formula of Kramer¹⁴⁷ together with a bound-free Gaunt factor. The radiative spontaneous emission rates are calculated using hydrogenic oscillator strengths. All rate coefficients are represented by simple analytic formulas.

Because TDIB is intended to coexist with an average-atom (AA) model in LASNEX, it is of interest to compare the results of both. In Fig. 2-70, we show a contour plot for the ratio $Z_{\text{DCA}}/Z_{\text{AA}}$ in the (T_e, ρ) plane for a cesium plasma in LTE. For $T_e > 20$ eV, the agreement is $\pm 20\%$. The largest discrepancies occur at high densities and low temperatures, where we expect pressure ionization and multiple excited states to be important.

Figure 2-71 shows the results of our calculation for a steady-state argon plasma at an electron density of 10^{19} cm^{-3} . The fractions of argon charge-states Ar X, Ar XI,

Fig. 2-71. Argon charge-state distribution vs electron temperature at a fixed electron density of 10^{19} cm^{-3} . Solid curves represent the results of our calculation; dashed curves are results of a coronal equilibrium calculation.

and Ar XII are plotted as a function of electron temperature. We also show the results of a coronal equilibrium calculation. The comparison shows that our results give the corresponding charge-state distributions at lower electron temperatures than those for the coronal equilibrium calculation. The enhanced ionization in our calculation is due to the contribution from the excited levels, which are populated by collisional excitation from the ground state. At these temperatures and densities, both the collisional excitation rate and excited-state ionization rate are very large compared to the ionization rate of the ground state.

Figure 2-72 compares the charge-state distribution of a steady-state krypton plasma to the experimental measurement. Both the electron temperature and density were estimated using the K-shell spectroscopy of phosphorus impurities seeded in the plasmas. The charge-state distribution was inferred from the L- and M-shell emission spectra.¹⁴⁸ We also show the results of a calculation using the code XRASER. Our results agree reasonably well with experimental measurement and the XRASER theoretical calculation.

Because one of the motivations in developing TDIB was to design x-ray laser experiments, we show the results of a prototype Nova foil target.¹⁴⁹ The design is a 500-Å selenium foil illuminated by a 0.5-μm laser of 4-ns FWHM; duration at an intensity of $5 \times 10^{13} \text{ W/cm}^2$. In Fig. 2-73(a) we show the resulting LASNEX (T_e, n_e)

time histories at the peak of the laser pulse, where the electron temperature is approximately 900 eV and the electron density is approximately $5 \times 10^{20} \text{ cm}^{-3}$. In Fig. 2-73(b), we show the resulting charge-state distributions, which have reached steady state by the peak of the laser. Finally, we compare in Fig. 2-74 the output x-ray energy of the average-atom result to that of the DCA (hydrogenic) model. Although the integrated outputs agree to within 10%, there is much more detail in the DCA results; these details are important in interpreting experimental results.

Nonhydrogenic Approximation. TDIB obtains the ionic level populations for the ionization stages, which are described using nonhydrogenic atomic data from the solution of the following rate equation:

$$\frac{dN_i}{dt} + \sum_k (W_{ki} N_i - W_{ik} N_k) = S_i, \quad (211)$$

Fig. 2-72. Charge-state distribution for krypton plasmas at an electron temperature of 800 eV and an electron density of $3 \times 10^{20} \text{ cm}^{-3}$.

Fig. 2-73. (a) Time history of (T_e, n_e) for two zones of an x-ray laser-foil design. (b) Time dependence of charge states Se XXIII to Se XXVII for the x-ray laser-foil design.

Fig. 2-74. Comparison of DCA and AA output spectra for an x-ray laser-fall design.

Fig. 2-75. Charge-state distributions of a selenium plasma calculated using different atomic models at a fixed density of $5 \times 10^{20} \text{ cm}^{-3}$.

Fig. 2-76. Average ionization states vs electron temperature of selenium plasmas calculated with different atomic models at a fixed electron density of $5 \times 10^{20} \text{ cm}^{-3}$.

where W_{lk} is the collisional or radiative rate for the transition between levels l and k . The code reads a data file to obtain the energy levels, oscillator strengths, electron collisional excitation and ionization-rate coefficients, and photoionization cross sections. The level populations for the remaining ionization stages are calculated using Eqs. (205) and (206) in the hydrogenic approximation. The term S_l gives the contribution to the population of level l due to ionization or recombination from the remaining ionization stages.

Figure 2-75 shows the charge-state distribution of a selenium plasma at an electron density of $5 \times 10^{20} \text{ cm}^{-3}$ and an electron temperature of 1000 eV. The curve labeled hydrogenic model represents the results of a calculation using the default option of TDIB. The curve labeled neon-like ion represents the results of a calculation in which we use nonhydrogenic data for a neon-like ion. The remaining curve represents the results of a calculation in which we use nonhydrogenic data for neon- and fluorine-like ions. The atomic model has 58 energy levels for the neon-like ion and 113 energy levels for the fluorine-like ion.

The energy levels for the $n = 2$ and $n = 3$ states of neon- and fluorine-like ions are computed using the multiconfiguration Dirac-Fock code YODA. The oscillator strengths for the transitions between these states are also computed using YODA.

Figure 2-76 shows the average ionization state of a selenium plasma at an electron density of $5 \times 10^{20} \text{ cm}^{-3}$ and as a function of electron temperature. The calculation using a hydrogenic model produces results that agree reasonably well with the calculation using a more detailed atomic model.

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The wetted-foam shell is a new concept in target design. The shell (shown at right) acts as a wick for the liquid DT, holding it and preventing gravitational slumping. For initial experiments, silica aerogel foam having a density of 0.11 gm/cm^3 was machined into hemishells that were bonded together (shown below, left). For future experiments, polymer foam with a lower atomic number is needed. Early foams contained spherulite structures because of the slow cooling rate (right, middle). Rapid cooling produced uniform $5\text{-}\mu\text{m}$ cell structures (right, bottom).

Experiments with wetted-foam targets require complex cryogenic systems. The vessel shown below (both photos) is part of a system that is used to fill cryogenic foam targets with liquid deuterium-tritium mixtures. Experimental capsules are mounted in this vessel and surrounded with helium gas to remove heat produced by the radioactive decay of tritium. Electrical and gas feed-throughs are used to monitor pressures and temperatures during these experiments.

3

Target Development and Fabrication

Target Development and Fabrication

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Target Development and Fabrication

Introduction

J. H. Campbell

The Target Development Group has responsibility for turning theoretical target designs into working target systems. The unclassified targets under development this past year generally fell into two broad categories: (1) ICF-related capsules and physics targets and (2) x-ray laser targets. We are also developing an advanced target assembly and prealignment system. This system will aid in the final assembly of the target components and expedite positioning and alignment of the target within the Nova target chamber. It is important to note that the Target Development Group also supports the classified ICF efforts, and thus the work reported here represents only one portion of our activities.

During 1985 considerable research and development has gone into a new direct-drive, ICF capsule design. This particular concept, which closely approximates an ideal ICF capsule, makes use of a hollow, low-Z, foam shell to hold a cryogenic liquid deuterium-tritium (DT) fuel layer in place. The requirements on the foam material are quite demanding. A thorough survey of commercial foam materials showed that none were available that met all the requirements. As a consequence, a sizable development effort has been initiated to produce the desired material.

We are also continuing development of a method for forming pure (i.e., not contained in a foam) liquid fuel layers in Nova-sized capsules. Currently, the most promising method makes use of a thermal gradient across the capsule; the thermal gradient induces a corresponding gradient in liquid surface tension forces leading to the formation of a uniform liquid layer within the capsule.

Last year we also began an effort to produce spin-polarized DT mixtures for possible fusion applications. In 1982, a paper by Kulsrud et al.¹ showed that the cross section for nuclear fusion of deuterium and tritium is increased by a factor of 3/2 if all the spins are aligned along an external magnetic field (i.e., polarized). Thus by using spin-polarized DT fuel it would be possible, in principle, to reduce the driver energy by one-third and still achieve the same fusion energy output. It is clear that the reduced driver size would have a significant impact on fusion reactor costs.

X-ray laser targets were developed and fabricated for two major campaigns this year, one on the Phoenix laser and the other on the Nova two-beam system. The objective of these campaigns was to develop a better understanding of the plasma physics and scale effects on the exploding-foil x-ray laser. The targets needed for the Phoenix campaign required coatings of a number of pure elements, mixtures of elements, and compounds. These coatings were deposited as 75-nm-thick microdots on a polymer film that was then assembled into the final target. The Nova x-ray laser experiments called for targets varying in length from 0.5 to 5.0 cm with lasant foils of selenium and yttrium. Also developed were coatings of oxides of higher-Z elements, particularly europium, gadolinium, and samarium.

A New Direct-Drive ICF Capsule

High-gain, direct-drive, single-shell capsules for reactor application employ levitated layers of condensed-DT fuel distributed very uniformly on the interior surface of a spherical capsule. This fuel layer may be either solid or liquid depending on the density of DT vapor required inside the capsule. Production of uniform cryogenic fuel layers inside such capsules represents a severe target fabrication and fielding challenge. The use of cryogenic capsules for power generation in reactors requires production techniques compatible with high-repetition-rate operation (injection of 1 to 10 capsules per second), and the costs must not be prohibitive.

A primary fabrication problem in forming uniform cryogenic layers is slumping of the fuel layer under the influence of gravity. This problem exists for both solid and liquid layers; during cooling from ambient temperature to temperatures below the DT triple point (19.7 K), the fuel will be present in the liquid phase long enough to undergo appreciable slumping. Several techniques have been proposed that either eliminate gravitational sag (e.g., fuel layer formation in free fall) or counteract it (e.g., thermal-gradient-induced flows). For application in reactors, the layering technique should be passive and require no active control or characterization of the fuel layer.

A new target that uses a passive mechanism to produce uniform layers has been proposed. It employs a rigid sponge matrix to support liquid DT in a spherical shell configuration. Capillary forces, generated by wetting the foam matrix with DT, stabilize the liquid against gravitational flow. Liquid hydrogen has been observed to wet most materials² with a zero contact angle. For liquid in a capillary, a pressure difference exists across the liquid/vapor meniscus. The capillary pressure is

$$P = \frac{2\gamma \cos(\theta)}{a}, \quad (1)$$

where γ is the surface tension of the liquid, θ is the contact angle ($\theta = 0$ for hydrogen) and a is the radius of the

capillary. The column height h of liquid that can be supported against gravity by this capillary pressure is

$$h = \frac{P}{g\rho}, \quad (2)$$

where g is the gravitational acceleration and ρ is the density of the liquid. For a random porous medium, the capillary pressure, expressed in terms of the specific surface area (surface area per unit volume) S and porosity ϕ of the foam material is

$$P = \frac{\gamma S}{\phi} \quad (3)$$

This allows us to define an effective hydraulic radius $r = \phi/S$ for the foam, which is loosely correlated with the pore or cell size of the foam. For a wetted-foam capsule, the pores must be small enough to support a liquid column about the height of the capsule diameter (≤ 0.5 cm) under the maximum acceleration the capsule will experience. With a capsule accelerator length of 50 m, pellet injection into a reactor at 1000 m/s requires an acceleration of ~ 1000 g. To remain stable at this acceleration with liquid DT wetting the foam, a capsule pore diameter of ~ 1 μ m is required. A little caution is in order at this point. The surface tension of DT is inversely proportional to temperature; it has a maximum value of 4.36 dyn/cm at the triple point (19.7 K) and vanishes at the critical point (39.4 K).³ Capillary pressure calculations must use the correct surface tension value, which is determined by the liquid temperature at the liquid/vapor interface. In foams with a small pore size, the stresses generated by capillary pressure can be high enough to cause material failure in the foam. Consequently, there exists a lower as well as an upper bound on allowable foam pore size.

A possible reactor capsule is illustrated in Fig. 3-1. The outer surface of the foam shell is covered with a thin permeation/pressure barrier to prevent DT boiloff. This capsule combines the traditionally distinct regions of ablator and fuel into one low-density, low- Z foam shell. Depending on its position within the shell, the DT-wetted foam serves either as the ablator or as the fuel. By varying the temperature,

the density of DT vapor in the capsule interior can be adjusted over a wide range of values. The presence of DT vapor in the capsule interior permits vapor hot-spot ignition, and consequently places less stringent requirements on the driver pulse shaping. Filling the capsule relies on permeation at ambient temperature through the thin polymer overcoat. For a capsule with a foam shell thickness equal to half the outer radius, we need to fill by diffusion to a gas density approximately 80% that of liquid DT. This corresponds to a pressure of ~ 1000 atm at 300 K. To avoid buckling or compression failure of the shell, the pressure must be increased in a controlled manner. Slow filling to extremely high final pressures can take several weeks or more. Radiation damage to the foam and overcoat is a major materials issue that requires minimizing exposure times and using radiation-resistant materials. The potentially long fill times are a concern because large tritium inventories are associated with batch-filling 10^4 to 10^6 capsules. Therefore, it is desirable to make capsules from mechanically strong foam and overcoating materials to increase the buckling threshold. The overcoat should be very permeable to hydrogen at room temperature (perhaps 100 to 1000 times the permeability of polystyrene, for example), but yet have reduced permeability at cryogenic temperatures and high resistance to tritium attack. These properties reduce the radiation dose by shortening the fill time. The use of damage-resistant materials and smaller batch fills with reduced high-pressure tritium inventories would also reduce radiation damage.

An alternative fill technique is possible. It eliminates the overcoat and thus makes diffusion filling unnecessary. This method places somewhat more stringent requirements on the foam properties and capsule fielding. In this approach, the foam is filled with liquid DT, which, in the absence of an overcoat layer, is free to communicate with the surrounding environment. However, evaporative transport of DT from the capsule in the reduced partial-pressure environments of the capsule injector and target chamber would cause problems that would require further investigation.

Delivery of the capsule in the proper cryogenic configuration to the center of

Fig. 3-1. Reactor capsule design combines ablator and fuel regions.

the target chamber is a concern relevant to the design of all cryogenic reactor capsules. For wetted-foam capsules, deviations from optimum temperature when the capsule is imploded would change the relative amounts of liquid and vapor DT in the capsule; this produces underfilling or overfilling, both of which are undesirable. Because times for capsule warming are short (0.1 to 1 s), dynamical thermal/mass transport issues may be relevant. For the non-overcoated capsule, one could conceivably utilize evaporative removal from the outer surface of the capsule to reduce or prevent capsule warm-up.⁴

The production of foam capsules for power reactors requires high-rate or batch processes with throughput of 1 to 10 capsules per second. Potential production pathways are droplet generation for high-rate shell formation or molding on leachable spherical mandrels for batch production. Approaches to generating rigid, open-cell, small-pore foams are described in the following article, "Polymer Foam Developments."

Wetted-foam targets potentially affect plant design and economics. It is necessary to evaluate the possible impact of the larger tritium inventory required by the use of wetted-foam capsules. Since DT is utilized as both fuel and ablator in this capsule, the effective burnup fraction will be relatively lower than in a capsule design where the DT is present solely as fuel. Aside from such issues as tritium inventory, we view the wetted-foam capsule as a very promising technical route to producing cryogenic targets economically for use in commercial ICF power production.

Author: D. H. Darling

Polymer Foam Developments

We are conducting research on foam production and developing overcoating techniques for the new direct-drive ICF target described in the preceding article. This target design calls for a spherical shell of foam that can be wetted with liquid DT. In this target, the liquid DT is both fuel and ablator; the foam serves to hold the DT in a spherical-shell configuration which is easily compressed to high density by the laser energy.

The foam must have a density of 0.025 g/cm^3 with less than 1% of the material composed of elements of higher atomic number than silicon and with less than 10% composed of elements of intermediate atomic number between nitrogen and silicon. The foam must tolerate exposure to the radiation environment associated with tritium without changing dimensions during the time needed for the target to fill and stabilize before the shot. The foam must be machinable or moldable to the required spherical shell configuration and must accept an overcoat of approximately $5 \mu\text{m}$ of a hydrocarbon (CH) or other low-Z material. The foam specifications and our foam development efforts are discussed in detail in the following articles

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Polyethylene Foams

Introduction. We have been working on the development of foams made from ultrahigh-molecular-weight polyethylene (UHMW PE) gels since April 1984. The goal is to produce a CH polymer foam for use in a new direct-drive foam shell containing liquid-DT fuel. For initial experiments, the foam must have a density of 0.05 g/cm^3 , a cell size of $1 \mu\text{m}$, and be uniform and strong enough mechanically that it can be machined and wetted with liquid DT without changing its dimensions.

Last year, we learned to handle UHMW gels and to characterize them for structure by optical microscopy. This year, we

improved and simplified our methods and can make PE gels in a matter of a few hours without oxidative degradation. Several new methods have improved foam: uniformity, and we routinely make easily handled, machinable, cohesive foams in densities ranging from 0.04 to 0.5 g/cm^3 with cell sizes of 2 to $5 \mu\text{m}$ at the lowest density.

We make foams by first forming an entangled solution, generally in tetralin, of UHMW PE, 2 to $6 \times 10^6 \text{ g/mol}$. The solution is gelled by cooling it, crystallizing the polymer. Isopropyl alcohol replaces tetralin in the first solvent exchange. Finally, the gel is flushed and permeated with liquid CO_2 and then supercritically dried.

Thermal characterization with differential scanning calorimetry (DSC), thermal mechanical analysis (TMA), and measurements of heat transfer on bulk pieces of gel were used to study the crystallization process. We found that we could improve the uniformity of the gels by reducing thermal gradients. We characterized cell structure by scanning electron microscopy (SEM), optical microscopy, and BET (Brunauer-Emmett-Teller) surface-area measurements. We used the solvent-exchange step to characterize the gels for cell structure by measuring the diffusion rate. Foam uniformity was measured with x-ray radiography. Our recent work has concentrated on improving the uniformity, mechanical properties, and characterization of foams.

Polymer Dissolution. The UHMW PE we use is a fine powder with a particle size of $\sim 30 \mu\text{m}$. Initially, we dissolved the polymer by heating it with the solvent in a sealed ampule for several weeks. The polymer was distributed throughout the solution by diffusion with some occasional shaking. Shaking or stirring the solution introduces bubbles that take days to rise through the viscous solution. Another disadvantage is that stirring also causes chain alignment. Even after several weeks, the gels may not be homogeneous.

In the quiescent sealed-ampule dissolution technique, the polymer settles to the bottom of the ampule as the solvent is heated. To dissolve the polymer more readily and to make the gel uniform, the solvent and powdered polymer are first mixed to a uniform suspension using

Fig. 3-2. Scanning electron micrographs of PE foams with densities of: (a) 0.035, (b) 0.10, and (c) 0.35 g/cm³.

ultrasonic agitation. The test tube, which contains polymer dispersed throughout the solvent, is plunged into a hot-oil bath. Within seconds, the particles, which were slowly rising in the dense, cold tetralin, start falling as the tetralin density drops with increasing temperature. Within one minute, all motion stops as sufficient polymer dissolves to thicken the solution. Polymer gels are made in a matter of two hours by this technique; the diffusion process requires several weeks. Polymer foams with densities ranging from 0.035 to 0.35 g/cm³ made with the quiescent sealed-ampule technique are shown in Fig. 3-2. Cell size of the foams decreases as the density increases. These foams are more uniform than those made without bubbles using other techniques.

Homogeneity During Dissolution. With radiography, we observed that the foam density is higher in the center axis of the sample. This density inhomogeneity probably results from the drop in the density of tetralin with increased temperature, as discussed in the preceding paragraph.

We tried to match the density of the solvent to PE at 150°C by making solutions of PE in mixtures of trichlorobenzene (TCB) and tetralin that ranged from 5 to 50 wt% TCB. The mixture of 5 wt% TCB in tetralin appeared to match the density of our PE. Some of these gels have been extracted to foams, but the SEM analyses have not been completed. We calculated the fraction of TCB required for density matching as 20.9 and 22.8 wt% from the respective experimental and theoretical values of the pure solvent densities. The estimated ratio of TCB to tetralin that is required to match PE density is much higher than the experimental value. However, we did not account for the dynamics of the melting and

heating processes, which include spatial and time variations in both PE and mixed-solvent density.

Density at each temperature was determined by analysis of the weight in solvent of a plumb bob of known volume. At 150°C, the densities of TCB and tetralin are about 1.3 and 0.87 g/cm³, respectively.

We predicted the pure solvent densities by first calculating critical temperatures from group contribution theory.⁵ Then the densities were estimated at 150°C by a method given in Perry⁶ as 1.29 ± 0.03 for TCB and 0.86 ± 0.02 for tetralin.

Chain Extension and Entanglements in Solution. Once the particles are melted, the polymer chains and the solvent interact. The structure and degree of chain extension of the polymer in the gel state are important factors in the structure of the foam. Viscometry measurements of the extension of the polymer chain molecules were used to calculate the concentration at which the molecules become entangled. Dilute-solution capillary-flow viscometry measurements of the limiting viscosity number $[\eta]$ were used to calculate the mol wt. The polymer solution viscosity is also a function of the solvent. In a good solvent, the chain is swollen with solvent and extended, which increases the viscosity; whereas, in a poor solvent, the chains become more compact.

During our viscometric measurements of PE solutions, we noticed a reduction in viscosity with time. We suspect that this reflects the reduction in mol wt caused by oxidation of the polymer molecules in the hot (135°C) solutions. An antioxidant, 4,4'-methylenebis(2,6-di-*tert*-butylphenol), was added to control oxidation. Figure 3-3 shows the relative viscosity as a function of time. The upper curve represents a

Fig. 3-3. Relative viscosity as a function of time for (a) a PE solution with no antioxidant and (b) a piece of foam dissolving with no antioxidant added.

Fig. 3-4. (a) Without antioxidant, the foam morphology is lamellar. (b) With antioxidant, the mol wt is kept high, which produces a fibrous morphology and a stronger foam.

dissolving piece of foam with antioxidant, and the lower curve shows a dissolved polymer without antioxidant. Without antioxidant, the viscosity drops rapidly. The mol wt falls from 2.5×10^6 to 4.0×10^5 g/mol in 24 h. This drop in mol wt reduces the number of chain entanglements; the resulting foam has a larger cell size and poorer mechanical strength. It takes approximately 2 h to dissolve the polymer with antioxidant. The solution can then be held at high temperature for more than 20 h without a drop in mol wt. The morphology of foams made with and without antioxidant is compared in Fig. 3-4. The foam without antioxidant is predominantly lamellar. Foam with antioxidant is much more fibrous, which may result from the greater number of chain entanglements caused by the higher mol wt.

The viscosity of polymer solutions can be used to evaluate the chain volume in solution. Einstein's viscosity equation⁷

$$\eta = \eta_0(1 + 2.5\phi) \quad (4)$$

relates the solution viscosity η to the solvent viscosity η_0 and the volume fraction of the particles ϕ . This equation can be rewritten in terms of limiting viscosity number, partial molar volume of the polymer V_2 , and the mol wt M as

$$[\eta] = 2.5 \frac{V_2}{M} \quad (5)$$

For polyethylene with a mol wt of 3×10^6 g/mol in tetralin solution, the experimental value for $[\eta]$ was 20 dL/g. The partial molar volume found for these conditions is

Fig. 3-5. The viscosity of tetralin solutions of PE rises rapidly at 80°C. Dodecane solutions of PE also show gelation, but at higher concentrations than tetralin solutions.

$3.98 \times 10^{-15} \text{ cm}^3/\text{molecule}$. If we assume the particle is spherical, the radius of the particle is $9.83 \times 10^{-6} \text{ cm}$.

The average radius of a polymer molecule can be calculated, assuming the polymer takes on a random flight configuration in solution. Then the mean square radius is⁸

$$r^2 = nl^2, \quad (6)$$

where n is the number of repeating units and l is the length of each bond. For $3 \times 10^6 \text{ g/mol PE}$, n is 2.1×10^5 , l is $1.54 \times 10^{-8} \text{ cm}$, and the mean radius is $7.06 \times 10^{-6} \text{ cm}$, which is slightly lower than the experimental value of $9.83 \times 10^{-6} \text{ cm}$.

Actual molecules do not assume truly random flight configurations in solution because bond angles and steric effects restrict flexibility and tend to increase the size of the molecule in solution. The mean square radius⁸ of a chain with a bond angle θ is

$$r^2 = nl^2 \frac{(1 - \cos \theta)}{(1 + \cos \theta)}. \quad (7)$$

For PE in which the polymer is linked by carbon-carbon single bonds, the angle

is the tetrahedral angle 109.5° . The mean radius calculated with the restricted bond angle is $9.98 \times 10^{-6} \text{ cm}$, which is almost exactly the same as the experimentally determined value of $9.83 \times 10^{-6} \text{ cm}$.

We measured the extent of entanglement of the UHMW polymer chains at the concentration used to make foams with vibrational viscometry in the gel state. The probe, a 1/2-in.-diam cylinder, oscillated at approximately 600 Hz with an amplitude of $25 \mu\text{m}$. Chains in UHMW PE solutions align along the stream lines made by stirring. In this case, the extremely low amplitude of oscillation avoided any stirring motion that might have induced polymer-chain alignment. Solutions of PE in both dodecane and tetralin over a range of concentrations were analyzed for viscosity-temperature behavior. The results are graphed in Fig. 3-5. As the temperature is reduced from 150°C , viscosity rises. Viscosity increases sharply for solution concentrations high enough to produce chain entanglement, and is determined by the amount of chain entanglement.

The concentration at which the viscosity shows a large rise is defined as the gel concentration C_R , which is given for both

dodecane and tetralin solutions in Table 3-1. If we assume that gelation is the condition at which the polymer chains just touch, then the polymer chain volume is given by

$$V_p = 0.74 V_t \quad (8)$$

where V_t is the total volume of the solution. The total polymer-particle volume equals the product of the volume of an individual particle times the number of particles

$$V_p = \frac{4\pi r^3}{3} \left(\frac{C_R V_t N_0}{m} \right) \quad (9)$$

where N_0 is Avogadro's number. Equations (8) and (9) can be combined to give

$$r = \left(\frac{0.177 m}{C_R N_0} \right)^{1/3} \quad (10)$$

Table 3-1. The radius r of the polyethylene molecule determined from experimental and theoretical techniques.

Solvent	Dilute solution		Concentrated gels	
	$[\eta]$ (dL/g)	r (cm)	C_R (mg/dL)	r (cm)
Dodecane	5	6.19×10^{-6}	30	3.08×10^{-6}
Tetralin	20	9.83×10^{-6}	7.5	4.88×10^{-6}

Random-flight radius = 7.06×10^{-6} cm.

Random flight with bond-angle restriction = 9.98×10^{-6} cm.

For tetralin and dodecane, the two solvents used, the gel concentrations are 7.5 and 30 mg/cm³, respectively. This gives a polymer-particle radius of 4.88×10^{-6} cm in tetralin and 3.08×10^{-6} cm in dodecane. These radii are approximately half those calculated by dilute solution viscometry. The higher solution concentrations may change the polymer dimensions, or the polymer chains may require overlap before gelation occurs. This approach gives us the ability to calculate a concentration for gelation based on dilute-solution viscosity. Alternatively, gel concentrations can be calculated for any listed solvent from the Mark-Houwink coefficients.⁹

Crystallization Kinetics and Crystal Morphology. The strength and structure of a foam depend directly on the crystalline structure induced during cooling to gelation. We studied the kinetics of crystallization and crystal size to aid our interpretation of gel structure and the heat transfer required during gelation.

Free-energy equations predict that the size of the most probable primary crystal will decrease as the temperature of crystallization is decreased. The final size of the

Fig. 3-6. Faster cooling lowers the crystallization temperature of PE/tetralin gels.

crystals also depends on secondary nucleation processes, but the initial size is determined by the primary nucleus. Thus, crystallization at lower temperatures should give smaller foam structures. We estimated the size of primary nuclei for PE in a solution of decalin to be 37 Å at 30°C and 89 Å at 90°C (Ref. 10).

We modeled the growth rate of the crystals that form the structure of high-mol-wt PE gels using Fick's law with Mandelkern's equation¹¹ modified for dilute solution. Two rate processes, transport and heterogeneous nucleation, are included in the model. We assumed 1000 nuclei per cm^3 . In the region of $T/T_{\text{melt}} = 0.8$, which corresponds to $T \sim 30^\circ\text{C}$ for our PE solutions, the overall rate of crystallization reaches a maximum because there is a maximum in the rate of homogeneous nucleation at that point. Above $T/T_{\text{melt}} = 0.9$ at $T \sim 70^\circ\text{C}$, homogeneous nucleation does not occur.

From DSC scans, we learned more precisely the effect of cooling rates on the temperature of crystallization of 5% PE in tetralin. When the cooling rate was increased from 1 to 30°C/min, the onset of crystallization dropped from 87 to 76°C, as

shown in Fig. 3-6. For these conditions, heterogeneous nucleation is occurring during gelation. However, in a simulation of the cooling of a thin section (0.305 cm) of gel, the temperature decreased from 120 to 20°C within 25 s, a rate of 240°C/min. This implies that we can increase the number of nucleation sites in thin sections with homogeneous nucleation and achieve a smaller cell structure.

Because the cooling rate determines the temperature at which crystallization begins, it should also determine the range of crystal sizes and perhaps the types of crystals produced. We used DSC to investigate this possibility at cooling rates of 0.1, 1, 10, 30, 60, and 320°C/min. In all cases (Fig. 3-7), a higher melting point and a better defined melting peak were observed when the sample was cooled more slowly. Very fast cooling (320°C/min) resulted in double peaks; cooling at 10°C/min resulted in a left shoulder. These results are consistent with others obtained from samples subjected to different rates of cooling under less controlled conditions than those provided by DSC. More rapid cooling should result in smaller morphology. It

Fig. 3-7. The rate of cooling of an ultra high-mol-wt PE in tetralin gel affects the morphology as seen by DSC.

Fig. 3-8. The crystal morphology of PE gels in tetralin becomes finer as the cooling rate is increased.

Fig. 3-9. Effect of the temperature of crystallization on the crystal structure of PE/tetralin gels.

can also cause dimorphism—the formation of two distinct crystals or the formation of crystals of different size.

To examine the effects of the temperature of crystallization, DSC was performed on three gel samples of 5% PE in tetralin; the samples were quench-crystallized at 99, 45 and -5°C . Optical micrographs of these three samples in Fig. 3-8 show that the gel morphology obtained at 45°C is

smaller than that at 99°C , and the morphology at -5°C is smaller than that at 45°C .

The gel quenched to 99°C shows less evidence of crystal dimorphism (Fig. 3-9) than expected from the cooling rate studies. However, the sample quenched to -5°C has less than half the crystallinity of samples quenched to 45 and 99°C . During handling, we noted that this gel was obviously

Fig. 3-10. Thermal mechanical analysis shows that PE foams melt at 130°C and that cycling below the melt leads to larger crystal structures.

weaker than gels quenched to higher temperatures. It might be possible to quick-quench the samples to -5°C and then anneal them at 80°C to increase crystallinity.

A TMA probe applying a force of 100 mN in a square waveform with a frequency of 0.05 Hz to the sample of PE foam obtained the results shown in Fig. 3-10. In the scan of sample 1, the temperature was increased from 30 to 160°C at a rate of $5^{\circ}\text{C}/\text{min}$. The deformation rate and modulus changed at $\sim 110^{\circ}\text{C}$; the melting point was at 130°C . With a fresh sample of PE foam under the same conditions, the temperature was increased to 110°C at $5^{\circ}\text{C}/\text{min}$; sample 2 was cooled rapidly and then heated again. After the first cycle, sample 2 showed a height change and modulus change, as before, but in the second cycle, it retained the higher modulus and reduced height reached at the end of the first cycle. We infer that the inflection at 110°C results from the melting of small, imperfect crystals and their recrystallization into higher density, thicker crystals. This is not easily seen on the DSC.

Crystal Structure. The thickness of the folded polymer chain lamella that make up the walls of the cellular foam can be measured by x-ray diffraction techniques.

For a foam whose walls are relatively perfect crystalline lamella, the crystallite thickness can be calculated from the Scherrer equation (Ref. 12),

$$L_{hkl} = \frac{K\lambda}{\beta_0 \cos \theta} \quad (11)$$

In the Scherrer equation, L_{hkl} is the crystallite dimension perpendicular to the hkl plane, β is the full width at half-maximum intensity of the reflection occurring at a diffractometer angle of θ corresponding to the hkl reflection, and λ is the radiation wavelength. Accurate analysis of crystal dimensions by the Scherrer equation is complicated by the broadening caused by lattice distortions. However, comparisons between similar materials provide useful qualitative information on crystal thickness. The thickness of polyethylene lamella is known to be a function of the magnitude of the undercooling during crystallization.⁷ We compared two samples of 3.5×10^6 g/mol polyethylene: one was melted and slowly cooled as a film; the other was a foam prepared as a gel and then dried. Analysis of the diffractometer scans from the two samples gave a crystal thickness of 17 nm for the melt and 8 nm for the foam. Because of

lattice distortion, both values may be low by a factor of 2. Since a cell-wall thickness of ~ 35 nm was calculated from the foam structural model, the cell walls must therefore be composed of two to four layers of polyethylene lamella. The relatively thin lamella present in the foam are consistent with the dimensional changes we see from recrystallization and lamellar thickening when foams are heated.

Fig. 3-11. Model foam structures. A typical ball chain is pictured only at the top front edge of the last model; these chains occupy all edges of the unit cube.

Table 3-2. Summary of model foam properties.

Model	Reduced foam ^a density (ρ/ρ_s)	Specific surface area ξ (m ² /g)	Foam density
Spherical shell	$4.44 T/L$	$2/\rho_s T$	$8.88/\rho_s L$
Sphere	—	$6/\rho_s T$	—
Cubic (closed)	$3 T/L$	$2/\rho_s T$	$0.00/\rho_s L$
Cubic (open)	$2.36(T/L)^2$	$4/\rho_s T$	$13.44 T/\rho_s L^2$
Tetrahedral (closed)	$7.35 T/L$	$2/\rho_s T$	$14.7/\rho_s L$
Tetrahedral (open A)	$16.37(T/L)^2$	$4/\rho_s T$	$65.48 T/\rho_s L^2$
Tetrahedral (open B)	$6.66(T/L)^2$	$4/\rho_s T$	$26.64 T/\rho_s L^2$
Cubic (ball chain)	$1.57(T/L)^2$	$6/\rho_s T$	$9.42 T/\rho_s L^2$

^a Density of foam divided by the density of the foam's structural material in g/cm³ expressed in terms of T/L , the reduced cell dimension.

Gel Pore Structure. We measured the time it takes to replace tetralin in the gels with isopropyl alcohol, a step that extends the time required to produce a foam. We also used this data to calculate diffusion coefficients, which are related to foam morphology.

A cylindrical PE gel, permeated with tetralin, was placed into a beaker of isopropyl alcohol. The refractive index of the stirred solution was measured over time and used to calculate the concentration of tetralin in the solution. In all cases, complete exchange had occurred within 48 h. The process was modeled using Fick's law for unsteady-state diffusion from an infinite cylinder.

Diffusion coefficients D_{ab} were a strong function of time, possibly because we assumed an infinite cylinder in the model. At times close to zero, all D_{ab} are about 6×10^{-4} cm²/s; after 3 h, the D_{ab} range from 1.7×10^{-5} for a sample cooled over a period of days to 2.8×10^{-5} cm²/s for samples made without antioxidant. Liquid-liquid diffusion coefficients are generally on the order of 10^{-5} cm²/s, while D_{ab} for liquid-through-solid-polymer diffusion range anywhere from 10^{-7} to 10^{-12} , depending on whether the polymer is in the rubbery or the glassy state. The pores in our PE foams are large enough (greater than ~ 5 nm)¹³ that the resistance to transport is due to the tortuosity of the pore structure and liquid-liquid transport. For long diffusion times, tetralin is probably diffusing from the smallest pores and out of the polymer itself, so the D_{ab} may be decreasing with time.

Foam Structure. The calculation of foam densities from several foam-model structures guided foam development and characterization and determined the relation of cell size and cell-wall thickness to ideal model structures. The volume of material contained in a unit cell is obtained with simple geometry. From the volume of material in a particular foam structure, both the mass of material in the unit cell and the foam density can be calculated. Modeling results were extended to foam-specific surface areas, and proved very useful in interpretation of the characterizations performed by BET surface-area analysis and mercury-intrusion porosimetry.

Fig. 3-12. Foam densities calculated for several cell shapes and cell dimensions. Typical PE foam has a density of 0.6 g/cm^3 , a plate thickness of $0.1 \text{ }\mu\text{m}$, and a cell thickness of $9.5 \text{ }\mu\text{m}$.

Table 3-3. Foam properties with values for wall thickness T and cell size L calculated from model structures.

The foam-model structures are shown in Fig. 3-11. The first structure is simply a close-packed array of spherical shells. Cubic (and tetrahedral) models have two distinct variations, closed cell or open cell. In closed-cell structures, the faces of a unit cell of length L are membranes of thickness T . The membranes are the solid portion of the structure. The void within the cell is not connected to adjacent voids in neighboring unit cells. In the open-cell model, the edges of a unit cell are composed of fibers with diameter T , and all void area is continuous. The cubic ball-chain model simulates an aerogel.

The equations for reduced foam density ρ_f/ρ_s and specific surface area ξ are summarized in Table 3-2. Using the equation for foam density in column 1, the specific surface-area equation in column 2 is rewritten in terms of foam density in column 3. Figure 3-12 is a plot of ρ_f/ρ_s as a function of reduced cell dimension (T/L). For most polymer foams, the density of the structural material is $\sim 1.0 \text{ g/cm}^3$. The magnitude of ρ_f/ρ_s is, therefore, close to the actual foam density. At constant density, foams with open structures can be

Material	Density (g/cm^3)	Surface area (BET) (m^2/g)	Model	T (μm)	L (μm)
PE 147-1	0.024	45.5	cubic (closed)	0.047	5.49
PE 147-2	0.035	50.5	cubic (closed)	0.043	3.39
PE 147-3	0.050	52.9	cubic (closed)	0.041	2.27
PE 142-2	0.053	54.9	cubic (closed)	0.039	2.06
PE 142-4	0.078	61.6	cubic (closed)	0.035	1.25
PS UL-2	0.10	16.2	spherical shell	0.11	5.5
$\text{SiO}_2\text{-Lu}$	0.10	513	cubic (ball chain)	0.0053	0.031

composed of thicker structural elements, which increases mechanical strength.

The foams we are making from UHMW PE can be modeled well as closed-cell cubes, although they are not actually closed-cell foams. Polyethylene foams are made up of polymer lamella, which are plates about 2 to $5 \text{ }\mu\text{m}$ across and $0.1 \text{ }\mu\text{m}$ thick. A data point marked on Fig. 3-12 represents a typical PE foam with density of 0.06 g/cm^3 ($\rho_f/\rho_s = 0.063$), cell size of $5 \text{ }\mu\text{m}$, and plate thickness of $0.1 \text{ }\mu\text{m}$ ($T/L = 0.02$).

Specific surface areas were measured with BET for several polymer foam samples (Table 3-3). Using the density and specific surface area from the model, we calculated the structural properties, wall thickness T and cell size L , for a variety of

foams. These calculations match SEM measurements. The foam models and the derived equations, which relate density to cell size and structural-element thickness, and those which relate specific surface area to density and structural-element thickness, have been very useful in developing foam materials. With the models, we can assess the feasibility of preparing foams by various techniques. We can also calculate parameters that are difficult to measure, such as plate thickness, from a knowledge of cell size and density.

We considered the use of x-radiography to investigate the uniformity and void structure of PE foams. X-radiography has been used in the past to measure the concentricity of coatings on microshells, but those coatings are full density. We were unsure of the applicability of our equipment for foams with a density less than one-tenth that of the full density materials studied previously. Test pieces of PE foam were irradiated with 7 keV x rays for 12 min at an electron current of 10 mA. The radiograph in Fig. 3-13(a) is of a cylinder with a higher density at the center. The radiograph in Fig. 3-13(b) is the profile of a foam made in a test tube. In addition to the radial density variation, a higher density region also shows up at the bottom of the tube. We believe both the radial and longitudinal density variations are settling problems caused by the drop in solvent

density during heating. We are applying several techniques described above in the paragraphs on density matching to study these variations in uniformity. X-radiography has given us extremely valuable information about the uniformity of these foams, even in the relatively early stages of its use. Using step wedges to calibrate each exposure, the radiographic information can be analyzed quantitatively for density and small-scale density fluctuations.

Mechanical Properties. Our early attempts to machine PE foam deformed it plastically instead of cutting it. Foam pieces that were diamond turned exhibited a very smooth surface finish. However, when the piece was fractured and examined under the SEM, we found that a fibrous compacted mat had formed on the surface, which is evidence of polymer-chain alignment. When the surface was bombarded by electrons in the microscope, the surface layer contracted and cracked. Electron-induced heating annealed the aligned chains and allowed them to recrystallize. The piece of PE foam in Fig. 3-14, which was drilled in the ion-milling machine, has a crack on the left side of the cut that formed during observation.

We improved the machinability of the PE foams by (1) using the rate-of-cooling effects, (2) adding submicron carbon fibers, and (3) irradiating the foams to

Fig. 3-13. X-radiography of PE foams shows a higher density in the center axis and at the bottom of the cylindrical samples.

Fig. 3-14. Ion milling produces well-defined holes in diamond-turned PE foams.

Fig. 3-15. Stress-strain plots for polystyrene inverse-emulsion foam, PE foam, and carbon-fiber PE foam.

produce crosslinking. Slowly cooling the gels increases the mechanical strength of the foam because it increases the thickness of the lamellar plates composing the cell walls. The effect of cooling on other foam properties, such as cell size and uniformity, has not been studied. Figure 3-15 shows stress-strain plots for polystyrene inverse-emulsion foam, PE foam, and carbon-fiber PE foam. The slowly cooled PE foam has a modulus E of 2.21 MPa with an elastic limit at 10% deformation. Carbon fibers were added to PE foams in various relative amounts, hoping that the composite foam would have improved mechanical properties. However, materials with a high carbon-fiber content are very fragile and crumble easily. The modulus of carbon-fiber-loaded foams is 0.11 MPa. The material with the highest modulus (8.46 MPa) is the polystyrene inverse-emulsion foam, discussed in "Inverse-Emulsion Foams" later in this section.

The first irradiation experiments oxidized the PE. Three samples of foam and a PE film were irradiated to total doses of 10, 100, and 500 Mrad of ^{60}Co gamma radiation. Infrared spectroscopic examination of the PE films showed a large absorption at 1750 cm^{-1} , which corresponds to the carbonyl ($\text{C}=\text{O}$) group absorption. The foams became progressively more crumbly as the radiation dose increased. To prevent the undesirable oxidation that accompanies irradiation in air, later irradiations were done in hydrogen or in vacuum. One outstanding feature in the infrared spectra of the irradiated PE film is an absorption at 970 cm^{-1} , which is the vinyl-hydrogen rocking mode. This

indicates that we are cleaving chains and leaving a high population of terminal vinyl groups.

An attempt was made to use mercury-intrusion porosimetry to characterize the PE foams for cell size and specific surface area. Foam compression during the intrusion of mercury is evident from the smaller size of the foam piece after the experiment. The foam appeared to fill with mercury at ~ 30 psi, which is the pressure necessary to fill a pore with a diameter of $5.5\text{ }\mu\text{m}$. The mechanical properties of PE foam under compression include elastic

behavior up to a strain of 10% and a stress of 0.2 MPa, which is equal to a pressure of 29 psi. The intrusion data must not be taken as a true measure of the pore structure of the foam but as confirmation of data on its mechanical strength.

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Rigid-Rod Foams

We have made a number of rigid-rod foams by crosslinking copolymers of polyisocyanates or polyisocyanitriles. Although our research may not be immediately applicable, we have established a method that allows us to specify the chemical composition and structure of a foam. Polymers of this type can be made with one enantiomeric sense of the twist to circularly polarize light; the process may also increase the likelihood of the formation of liquid crystals.

Crosslinking is a common procedure to increase the strength and cohesiveness of polymers. Most polymer chains take on vastly different conformations depending on their thermodynamic environment. In a solvent that matches the polymer chemical potential, the polymer chain is extended. Thus, a crosslinked polymer system in this solvent swells and the density of the polymer in the gel is uniform and low. But on removal of the solvent, by any means, the network collapses on itself. Therefore, rigid chains are used to make a low-density foam.

In this work, we make foams by crosslinking copolymers of polyisocyanates or polyisocyanitriles, polymers that have a rigid-rod, helical configuration and are soluble because of a pendant group whose functionality can be changed. Two comonomers are used. One is inert to the crosslinking procedure. The concentration of the other, the reactive comonomer, determines the spacing between crosslink junctions, and thus, the void diameter. The chemical structure of each monomer influences the rigidity of the polymer chain, e.g., a hydroaromatic functionality in the

crosslinking monomer lends more rigidity than a linear aliphatic functionality. The advantages of this scheme are that an inherent chain stiffness, a specific chemical composition, and a specific void diameter can, in principle, be built into the structure.

These two families of helical, rigid-rod polymers are of interest for other reasons. They can be made to have one enantiomeric sense of the twist, which will cause the rotation of circularly polarized light and should make the formation of liquid crystals more likely. Green and his students at The Polytechnic Institute of New York were the first to develop an optically active polyisocyanitrile with a solketal functionality (a functionality that is easily derivatized). They work in conjunction with LLNL to develop the copolymer systems most suitable for synthesis of foams.

Thus far, we have shown that foams of copolymerized, crosslinked polyisocyanitriles or polyisocyanates can be made. We have made 30 to 50 gels using 4 different comonomer systems over a range of molar ratios: hexyl isocyanate/hexane diisocyanate (HIC/HDIC), butyl isocyanate/hexane diisocyanate (BIC/HDIC), α -methyl benzylisocyanitrile/4,4'-diisocyanitrile dicyclohexyl methane (MBIN/DINDCHM), and *n*-butylisocyanitrile/1,6-diisocyanitrile hexane (BIN/DINH).

We removed solvent from the gel systems by supercritical extraction. During extraction, the foam densities increased from the original 0.2 g/cm³ to final densities in the range 0.3 to 0.5 g/cm³. Scanning electron micrographs of two of the foams are shown at two magnifications in Fig. 3-16. An HIC/HDIC foam with a density of 0.27 g/cm³ [Fig. 3-16(a) and (b)] has fibers 900 Å in diameter that form a woven structure with an average cell size of ~ 1 μ m. The foam itself is white and tough. The structure shown in Fig. 3-16(c) and (d) is a MBIN/DINDCHM foam of 0.4 g/cm³. It is globular with cell sizes on the order of 0.1 μ m. This foam is reddish brown, brittle, hard, and somewhat friable.

With the computer, we generated atomic scale pictures of the poly(hexylisocyanate) and also modeled the reaction kinetics of the isocyanate systems. A short segment of the chain backbone is pictured as

Fig. 3-16. (a) and (b) Polyisocyanate foam of density 0.27 g/cm^3 from HIC/HDIC has an open-weave, fibrous structure; (c) and (d) polyisocyanate foam of density 0.4 g/cm^3 from 4,4'-MBIN/DINDCHM has a small, globular structure.

Fig. 3-17. The helical, rigid-rod structure of a segment of poly-(hexyl)isocyanate is clear in this computer-generated molecular model.

spheres in Fig. 3-17; an n-hexyl group is on each nitrogen and an oxygen on each carbon. There are eight repeating units of RNCO per turn of the helix; the pitch of the helix is $5.17 \text{ \AA}$, and the angle between repeating units is 135° .

We modeled the crosslinking, using the material density and number of rods of a given length, by creating a box with periodic boundary conditions to contain the system. Then for N rods, N initiation sites are chosen within the box in a random fashion. Random directions are picked for each site for the rod growth. One monomer is then added to each site in the direction of rod growth. A specified fraction of these monomers are diisocyanate. The

distance between the rod position of each crosslinker and the heads of each growing rod is determined. If the distance is less than the length of a crosslinker, we then consider the possibility of a crosslinking reaction. When the distance criterion is met, but the crosslinker is headed in the wrong direction, a geometric factor of $1/6$ is applied. That is, a random number is generated and if it is greater than $1/6$, we say that the crosslinker was in the wrong direction and no reaction can occur. More monomer is then added to the rod. The process continues until the rods have reached the desired length.

The model oversimplifies several features. First, there is no excluded volume,

meaning that rods can pass through each other. In terms of a gross measure of structure and number of crosslinks, this is probably not a serious problem since the calculation is purely statistical; bent rods would give the same results. A more serious flaw is the assumption that if a crosslinking is geometrically possible, it will occur. It is likely in many cases that the competition with additional monomer will lead to simple rod extension rather than to crosslinking. In addition, we have not allowed for rod mobility, probably fairly important for short rods, but less so for longer ones.

We find that for systems of 100 rods of 100 monomers each at a density of 0.05 g/cm^3 with 20% crosslinkers and a crosslinking length appropriate to 6 carbons, very little crosslinking occurs. We get fewer than one crosslink per rod, meaning that 95% of the crosslinkers go unreacted. Later reaction is unlikely because of the immobility of the rods. This result is not unreasonable, since crosslinking requires the intersection of a dangling crosslinker by a growing rod. At densities of 0.05 g/cm^3 this is not a common occurrence. However, it may not take very many crosslinks to significantly increase the strength of what is probably a "log-jammed" structure.

Our experiments show a possible connection, because we get gels when no crosslinking monomer is used. For homopolymers synthesized to have chain lengths of 50, 125, and 400 repeating units, and at a density of 0.2 g/cm^3 , gels formed during synthesis. Gel stiffness increased with chain length, indicating further the likelihood of a "log-jammed" structure.

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Inverse-Emulsion Foams

Introduction. Emulsions are two-phase mixtures in which one phase is dispersed as small droplets in the other, immiscible phase. Surfactants (surface-active agents)

reduce the interfacial tension between the two phases, and may be used to stabilize the dispersed droplets in an emulsion. The magnitude and direction of the interfacial curvature characterize an emulsion: the curvature of the interface determines the droplet size; the direction of the curvature determines whether an emulsion is regular or inverse. In a regular emulsion, the phase with the smallest volume is the dispersed phase. An inverse emulsion is one in which the major phase is encapsulated within the minor phase.

We are investigating the use of inverse emulsions. These emulsions have been reported in the literature many times.¹⁴⁻¹⁶ The novelty of our work is that we are polymerizing the emulsion and then removing the aqueous phase to produce a foam. We are following a threefold approach to the development of inverse-emulsion foams for use in direct-drive ICF targets: academic study of the process at Case Western Reserve University, industrial development by Unilever Research Laboratory, and work in our own laboratory. Workers at Unilever developed an emulsion process for making foams.¹⁷ Their patented process is similar to that discovered independently at Case by Litt and Krieger,^{18,19} who have synthesized special surfactants to produce inverse emulsions. Inverse-emulsion experiments have also been conducted at LLNL.

The foams made during the past year by several variations of the inverse-emulsion process have good strength and machining characteristics. However, these foams exhibit density variations and voids of entrained air or coalesced water droplets, which could cause uneven fuel distribution. Because of its closed-cell structure, the Case foam is not as easily wetted as the Unilever or LLNL foams. However, the Case foam is well behaved in that exactly one monolayer of surfactant is adsorbed at the interface; this allows us to predict and control the cell size. Some combination or hybrid of these processes may enable us to produce a foam that has both the open cell structure and the predictable cell size required for use in direct-drive ICF targets.

Inverse Emulsion Formation. Emulsions are very familiar items; many food products, such as margarine, whipped cream, ice cream, milk, and mayonnaise

are emulsions. Historically, emulsions have been described as oil-in-water (O/W) or water-in-oil (W/O). If the "oil" phase (neither phase need be oil) occupies 10% of an O/W emulsion, the solution is composed of many small droplets of oil. Droplet size is governed by the interfacial tension and the amount of mechanical energy (e.g., stirring) used to emulsify the mixture. A surfactant molecule with both a hydrophobic end and a hydrophilic end reduces the interfacial tension, and a smaller amount of mechanical energy is required to stabilize the emulsion. The interface at which the surfactant molecules have adsorbed can have a different interfacial tension at the water and oil sides. The emulsion will curve towards the side with the greater interfacial energy, thus reducing its area. For example, in a 10% oil emulsion, if the water end of the surfactant has greater interfacial energy than the oil end, then the interface curves toward the water phase. This would result in an inverse, W/O emulsion.¹³

The work at Case postulates that the size and shape of the surfactant, as well as the interfacial tension, determine the direction of interfacial curvature. To investigate this premise, Litt and Krieger synthesized a range of surfactants that were block co-oligomers of selected polarities and geometries. Early results showed that it is possible in this way to produce dispersions of minute droplets of water in polystyrene where the water occupies as much as 95% of the volume. Initially, they used diblock co-oligomers prepared by cationic ring-opening polymerization from monomers that were substituted oxazolines. A hydrophobic block of bulky geometry and low surface energy was coupled to a more hydrophilic block. They varied the nature of the monomers and the lengths of the blocks to determine the

effect on the structure of the emulsion and the droplet size of the dispersed phase.

Of the 22 block copolymers synthesized and tested so far, 5 show promising results as inverse emulsifiers that produce solid foams from water-in-styrene (W/S) emulsions. For any given W/S ratio, there is a maximum amount of surfactant that produces stable emulsions; larger amounts break the emulsion, allowing the droplets of the dispersed phase to settle out under the combined effects of sedimentation and coalescence. The maximum effective weight of surfactant was found to increase as the styrene-to-water ratio increases. This trend is depicted in Fig. 3-18, where the volume of styrene used is plotted against the amount of surfactant. A boundary for the formation of solid foams can be clearly delineated. Stable foams are made only with formulations plotted above the line.

The emulsions were polymerized with a potassium persulfate initiator at 60°C and then dried in a vacuum oven. Scanning electron micrographs of three dried foams are shown in Fig. 3-19(a), (b), and (c). All foams were made from emulsions with

Fig. 3-18. Emulsion stability and structure are determined by the volume of styrene and the concentration of surfactant (in this case, PAE-04).

Fig. 3-19. Microstructure of inverse emulsions influenced by the surfactant: (a) 0.1 wt% surfactant, (b) the same surfactant at 0.05 wt% concentration has a larger cell structure, and (c) a different surfactant at 0.5 wt% has a thin cell wall.

a W/S ratio of 90/10. The foams shown in (a) and (b) were made under the same conditions except that (a) has twice the surfactant concentration of (b), and the total cell area in (a) is about twice that in (b). Calculations of the surface area of the foam indicate that the surfactant spreads to exactly one monolayer thickness at the water/styrene interface. In contrast to the Case method, the Unilever and LLNL processes use a large excess of surfactant. The Case emulsion, however, becomes unstable when excess surfactant is added. Figure 3-19(c) shows polymerized emulsions prepared under the same conditions as (a) and (b) except for the substitution of a different surfactant. The cell size is similar to those in Fig. 3-19(a) and (b); however, the cell wall structure is very different. The cell walls in Fig. 3-19(c) appear to be thin membranes, whereas the cell walls in (a) and (b) are thicker. We found that the first two foams dried much more slowly than the third. The thickness of the cell wall in the first two samples decreased the rate of water permeation. In general, these foams can be characterized as cellular materials with few holes in the cell walls.

Two foam samples made by inverse-emulsion polymerization at Unilever Research Laboratory were machined; diamond turning worked very well. Figure 3-20(a) is a SEM of a foam with a density of 0.10 g/cm^3 and a cell size of $\sim 3 \mu\text{m}$. The cell structure is remarkably different from the Case foam in that the cell walls contain numerous holes interconnecting the cells. This foam has many bubbles, ranging from a few microns to a millimeter across [Fig. 3-20(b)], which

would interfere with wicking of the liquid hydrogen fuel. By carefully degassing the components and avoiding turbulence during stirring, we may be able to remove the bubbles. We are working on an agreement with Unilever to collaboratively develop the inverse-emulsion process for the production of foams for laser targets.

The Unilever foams were characterized for surface area by the BET surface-area adsorption method and for pore size by mercury-intrusion porosimetry. The specific surface area of an emulsion foam with a density of 0.1 g/cm^3 was measured as $16.2 \text{ m}^2/\text{g}$. Using the spherical shell model discussed in a preceding article, "Polyethylene Foams," we calculated the wall thickness as $0.11 \mu\text{m}$ and the cell diameter as $5.5 \mu\text{m}$. Both these calculations match SEMs of the cell structure. Mercury-intrusion porosimetry showed that the pores filled when the pressure was sufficient to intrude a $1.5\text{-}\mu\text{m}$ -diam pore. This diameter corresponds to the size of the holes in the walls of the emulsion cells.

We have developed a process in our own laboratory for making inverse-emulsion foams. Rosen's work¹⁵ indicated that a surfactant with an HLB (hydrophilic-lipophylic balance) rating from 3 to 6 is required to form a W/S inverse emulsion. We chose a surfactant, sorbitan mono-oleate, with an HLB of 4.3 (Ref. 20). Styrene and surfactant were mixed first; then water (with potassium persulfate) was slowly added to the styrene while stirring continuously. After several minutes of stirring, the mixture thickened to a paste-like consistency. Next, the emulsion was poured into a tube or mold and polymerized

Fig. 3-20. Inverse emulsion foams made by Unilever (a) with holes in the cell walls and (b) with numerous voids from $20 \mu\text{m}$ to 2 mm in diameter.

Fig. 3-21. (a) An inverse emulsion foam made at LLNL with the same general structure as the Unilever foam in Fig. 3-20. (b) X-ray radiograph of a 2.5-cm-diam by 1-cm-thick section of this foam shows several internal voids.

by heating for several hours at 60°C. Finally, the emulsion was dried in a vacuum oven for several days. The structure of this foam, shown in Fig. 3-21, is similar to that of the Unilever foam. Figure 3-21(b) shows a radiograph of a cylindrical piece of the same foam. Several millimeter-size bubbles are visible. We had no difficulty machining the inverse-emulsion polystyrene made at LLNL on a vibratome.

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Extractable Pore-Former Foam

Recent results obtained at Porelon Corporation suggested that their extractable pore-former process might be used to produce CH foams to meet ICF target requirements. Because the Porelon foam did not have the strength at low densities needed for ICF target shells, we briefly examined the process in our laboratory to see if its limitations could be overcome. Although we were able to modify this process for ICF applications, the problem of fragility at low density remains.

Workers at Porelon Corporation mixed polyethylene and NaNO_3 powders in a grinder, compressed them at elevated temperature to form a continuous polymer network, and then extracted the NaNO_3 with water. The foams with a density of 100 mg/cm³ or less tended to be very

fragile because of a lack of structural connections, and although most cells were less than 10 μm in diameter, cells with diameters as large as 50 μm formed frequently because large pieces of NaNO_3 were in the powder.

Taking advantage of the fact that the chemistry and physics that determine cell size can be largely separated from the manufacture of the foam material itself, we first demonstrated that an air-jet grinder could produce a finer powder. We produced a NaNO_3 powder with an average particle size of a few microns by recirculating the powder in two colliding gas jets until it was fine enough to escape a cyclone separator.

We next explored alternative methods of dispersing the foam material: in a polymer slurry and in a dry mix of polymer with the pore-former material. First, we formed a slurry in polymer solution, evaporated it to near dryness, and then compressed it into a pellet. This worked to some extent, but it was difficult to dry the slurry without squeezing out most of the polymer solution; the solution had to remain tacky enough for the polymer to form a continuous network. In a subsequent variation of this method, we evaporated the slurry to dryness before pelleting and then assured a continuous network by heating the pellet to the softening point of the polymer. The pressed pellets had a porosity of about 30% for rapid leaching. We evacuated the pellets before adding them to the water to eliminate air bubbles that could cause vapor lock.

Using the dispersion method above, we made foams from polystyrene and from a 1:2 mixture of a polystyrene-butadiene copolymer and fine carbon fiber. We attained cell sizes of a few microns and densities of about 100 mg/cm^3 . The polystyrene foam was sufficiently plasticized by the water that it had to be freeze-dried to prevent shrinkage. The carbon-fiber-based foam did not shrink perceptibly when dried just in a vacuum oven; it did tend to shed dust upon handling. In situ polymerization with a cross-linking agent would eliminate plasticization of the polystyrene foam and a higher fraction of polymer would probably make the carbon foam less powdery, but neither of these options was pursued.

Extension of the technique to smaller cell sizes with either a cross-linking agent or a higher fraction of polymer is feasible and requires only a finer particulate. However, the variable thickness and some lack of connections means that foams produced by this technique are not as strong at a particular density as those with a more perfect structure, such as the foams produced by the inverse-emulsion process described in the preceding article "Inverse-Emulsion Foams." In fact, we were not successful at producing a lower-density material with adequate strength.

In a second approach, we mixed a solid polymer or precursor with the NaNO_3 in the grinding process. The material used in this case was a silicone resin because it was available in very fine particle form, but any brittle, solid resin would work.

Fig. 3-22. Foam with a density of 140 mg/cm^3 made from the methyl silicone resin GR-650 by the extractable pore-former process.

The resulting pellet was then heated to 120°C to cure the resin before extraction of the NaNO_3 . A foam with a density of 140 mg/cm^3 and a structure similar to that of foams made with the slurry method is shown in Fig. 3-22. This foam was more flexible than the others, but it shrank about 5% in length during drying, which would make precision forming difficult. With this dry-mixing technique, it should also be possible to make a foam with a controlled density gradient by varying the ratio of resin and pore-former with time in an aerosol deposited on a substrate. The deposit could then be compressed into a sheet and the pore-former extracted.

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Condensed Fuel Layers

The development of techniques for forming condensed fuel layers on the interiors of hollow microspheres is a continuation of our work from previous years. This year, we are concentrating on thermal-gradient studies, which seem more promising than solid layers formed by cooling the microspheres while in free-fall. The initial experiments with the noncontact thermal-gradient technique for forming and maintaining liquid layers and the work with solid layers were reported in the 1984 Laser Annual.²¹

This year, Professor Kim and his co-workers at the University of Illinois made more measurements of the thermal gradient needed to support liquid hydrogen-deuterium (HD) layers. Uniform layers of liquid HD were obtained from an improved apparatus, but the results still did not agree with the two-component model developed at the University of Illinois.

In a new experiment, thermal gradients were developed in the microspheres by cooling the spheres with cold-gas jets and differentially heating them at the same time with a graded-intensity laser beam. This system produced uniform layers in both D_2 - and DT -filled microspheres. We found that these layers could be frozen while maintaining the thermal gradient; we also made preliminary experiments with a

thermal-gradient apparatus constructed at LLNL.

Temperature-Plate-Induced Thermal Gradient

Parallel plates at different temperatures apply a temperature gradient to a hydrogen-filled microsphere in the University of Illinois apparatus described in the 1984 Laser Annual.²¹ The microsphere is placed between the plates and immersed in low-density helium gas to conduct the thermal gradient to the sphere. Last year, 600- μ m-diam microspheres filled with H_2 - D_2 were inserted between 4- μ m-diam plates placed 2.5 or 10 mm apart. Uniform liquid layers could be obtained with either plate separation, but the apparent temperature gradient required was different. This led us to believe that the gradient was not simply linear, as required, but that a higher-order thermal field was present. In 1985, we increased the diameter of the plates to 11 mm and worked at plate separations of 2.5 and 5 mm to minimize edge effects.

Measurements in a 600- μ m-diam sphere are shown in Fig. 3-23. Temperature gradients across the microsphere were

calculated from the temperature differences between the plates, taking into account the thermal transport properties of the helium coupling gas, the sphere materials, and condensed hydrogen. We now see no systematic difference in the data whether the plates are 2.5 or 5 mm apart. However, the data are scattered and disagree quantitatively with a two-component model published previously. We do not understand the scattering, especially since data from D_2 -filled spheres do not exhibit it.²¹ Results may disagree with the model because there are three components in the condensed hydrogen: D_2 , H_2 , and HD. A three-component model is presently being developed.

Gas-Jet-and-Laser-Induced Thermal Gradient

Kim et al. demonstrated in 1978 that condensed layers could be maintained in a microsphere that was held in a vacuum but otherwise exposed to ambient radiation. To do this, they directed cold helium gas jets onto the sphere.²² Kim and Rieger also showed that they could manipulate the distribution of liquefied hydrogen

Fig. 3-23. Thermal gradients required to form a uniform liquid layer inside an HD-filled microsphere.

Condensed Fuel Layers

inside a microsphere by directing a laser beam onto portions of it.²³ We have now combined these two techniques in a single apparatus. Cold helium gas jets cool a microsphere in a room-temperature vacuum chamber while a graded-intensity laser beam changes the temperature and temperature distribution of the microsphere,

thereby manipulating the condensed hydrogen into a uniform layer.

The apparatus is pictured in Fig. 3-24. A single copper cryostat cools ultrahigh-purity helium gas and delivers it to two adjustable gas-jet assemblies mounted in the cryostat head. A separated vacuum chamber houses the cryostat and gas jets,

Fig. 3-24. Diagram of the cold gas-jet apparatus.

Fig. 3-25. Schematic of diagnostic and laser heating optics.

with the cryostat head forming one of the dividing walls. The target under test is mounted on a glass stalk attached to a nylon screw. The screw is attached to a micrometer-controlled xyz-translation assembly. This configuration allows us to easily load and manipulate the target independently of the cryogenics and gas jets.

The optical arrangement for heating and analyzing the microsphere is shown in Fig. 3-25. We use an argon-ion laser as a heating source. A smaller He-Ne laser is part of an interferometer that is the primary diagnostic of layer uniformity. The two laser beams travel the same path through the microsphere and the interferometer. At the output of the interferometer, a beam splitter forms two images to be viewed by two cameras. The first image is filtered so that only the He-Ne image is observed. A second image, produced at the second camera by the much brighter argon-ion laser, is attenuated and reduced and the sphere appears small against the profile of the laser light. This argon-ion image provides information about the

graded light influencing the target-liquid layers and is also useful in the alignment of the target within the two opposing gas jets.

Uniform layers in a 200- μm -diam sphere filled with D_2 were formed using this apparatus (Fig. 3-26). The centered interference fringes in the sphere indicate a uniform layer. In Fig. 3-26(a) the cryostat temperature is 14.3 K. The sphere temperature is unknown, but it results from a balance between the heating power of the laser and the cooling power of the gas jets. In Fig. 3-26(b) the cryostat has risen to 15.9 K, thinning the condensed layer and thus producing fewer interference rings. No changes in the heating laser profile or power were made, and the layer remained uniform. In Fig. 3-26(c) the layer thickness was increased again by turning down the heating laser power from 15 to 10 mW, which produced an increase in the number of fringes. The intensity distribution was kept constant while the laser power was decreased, thus maintaining layering uniformity. Finally, by decreasing the laser

Fig. 3-26. Interference micrographs of uniform condensed layers in a 200- μm -diam D_2 sphere with a 3.17- μm wall thickness and 133-atm fill pressure: (a) at a cryostat temperature of 14.3 K, (b) at 15.9 K, the layer is thinner, i.e., fewer fringes, (c) at 15.9 K, a thicker layer from reduced heating laser power, and (d) a frozen layer.

(a) $T = 14.3$ K (b) $T = 15.9$ K
Heating laser = 15 mW (c) $T = 15.9$ K
Heating laser = 10 mW (d) Frozen layer

(a) $T = 14.8$ K (b) $T = 12.2$ K

Fig. 3-27. Interference micrographs of uniform condensed layers in a 465- μm -diam DT sphere with a 6.3- μm wall thickness and fill pressure of 64 atm with (a) cryostat temperature 14.8 K and (b) 12.2 K.

power to zero, we were able to freeze the layer shown in Fig. 3-26(d). The frozen layer is uniform and apparently fairly smooth.

We also experimented with a 465- μm -diam sphere filled with a $\text{D}_2\text{-T}_2\text{-DT}$ mixture. The thermal gradient is induced in the opposite direction from that for D_2 . Results with the cryostat at 14.8 K and 12.2 K are shown in Figs. 3-27(a) and (b).

Fig. 3-28. Photograph of the cryogenic chamber for the LLNL thermal-gradient apparatus.

LLNL Thermal Gradient Experiment

We have constructed a parallel-plate version of a thermal gradient apparatus at LLNL. The plate diameters are 30 mm, and the separation is 5 mm. We expect to obtain a good linear thermal gradient with no edge effects. A picture of the cryogenic chamber, which is mounted on the tail of a variable-temperature cryostat, is shown in Fig. 3-28. The temperatures of the plates are controlled by heaters and diode thermometers in a closed-loop system. Windows in the chamber provide dual perpendicular-path optical access for two-view holographic interferometry, the technique we use to characterize the fuel layer.²⁴ The cryostat is securely mounted to an optical table where the holographic interferometer has been assembled according to the schematic in Fig. 3-29. We use an image analysis system to enhance the TV images of the interferograms and to aid in quantitative evaluation of the data. In our preliminary experiments, we placed a DT-filled microsphere of bare glass between the plates. The 200- μm sphere was suspended between ultrathin Mylar films held by a thin Teflon ring. We were able to manipulate the liquid layer from a puddle in the bottom of the sphere [Fig. 3-30(a)] to a reasonably uniform layer [Fig. 3-30(b)], and then to an "inverted"

Fig. 3-29. Diagram of the two-view holographic interferometer. The He-Ne laser forms the hologram. The argon-ion laser vaporizes the fuel without heating the chamber.

Fig. 3-30. Effect of plate-temperature changes on DT-liquid layers formed in the LLNL apparatus: (a) initial puddled liquid, (b) uniform layer, and (c) inverted layer.

layer collected at the top of the sphere [Fig. 3-30(c)]. In this run, we were using only one view of the sphere. The measurement of the triple point for the fuel was about a degree less than it actually is, so either there was a systematic error in the thermometry or the sphere was being subjected to radiant heat from the windows. We are working to correct this situation.

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Spin-Polarized Deuterium-Tritium

Introduction

The use of spin-polarized DT fuel could lead to a large saving in laser energy in fusion reactors. It decreases driver energy requirements by a factor of 3 (Ref. 25), and calculations show that the polarization could survive long enough to increase the performance of both magnetic fusion and ICF reactors.^{25,26}

In spin-polarized DT, which is defined in detail in the *Laser Program Annual Report 1984*,²⁷ the nuclear spins have a net alignment in a direction defined by an externally applied magnetic field. A solid sample of electronic and nuclear spins in a high magnetic field has an electronic polarization $P_e \approx 1$, indicating that all the electron spins point in one direction. The nuclei, however, point up or down at random, and have nearly zero polarization.

A microwave field can be applied at low temperatures to provide photons (of a frequency determined by the difference in Zeeman energies before and after the flip) that force the simultaneous flip of a pair of electron-nuclear spins. In this process, called dynamic nuclear polarization, DNP, one electron can flip many nuclear spins, provided that the electronic relaxation time of the electron T_e is much shorter than the spin-lattice relaxation time of the nucleus T_n . If the properties of the sample are favorable, this leads to a nuclear polarization equal to the electron polarization.

Normal DT is unsuitable for spin polarization²⁷ and pure molecular DT must be used. A chemical synthesizer was built in 1984; collision-induced infrared spectroscopy shows that the product is 96 mol% DT, 0.6 mol% T_2 , 1.8 mol% D_2 , and 1.6 mol% HD. Because making pure DT is an expensive process, work this year was done with normal DT whenever possible.

For reasons outlined in the 1984 Laser Annual,²⁷ our efforts to produce spin-polarized DT have focused on measuring T_e , T_n , and the number of free electrons N_e , which are the end result of an energetic electron produced during tritium decay. The concentration of free atoms that will be reached in a sample of DT is unknown and must be measured. To maximize nuclear polarization, we need many electrons, a short electronic relaxation time, and a long nuclear relaxation time.

Our program consists of three steps:

1. Measure T_n with nuclear magnetic resonance (NMR).

2. Measure T_1 and N_1 with electron paramagnetic resonance (EPR).

3. Test DNP at 3.4 T.

We selected a magnetic field of 3.4 T for our first measurements of polarization because the cost of scaling NMR experiments to this field is reasonable, and it appears possible to build a flexible microwave pumping system for this magnetic

Fig. 3-31. Longitudinal nuclear relaxation time T_1 for protons and deuterons in solid HD with $J = 1$ H₂.

Fig. 3-32. First measured triton T_1 experiment at 27 MHz on 25% T₂, 50% DT, and 25% D₂.

field. Our first effort was to measure T_1 using NMR.

Nuclear Magnetic Resonance Measurements

During 1985, we completed a system for measurements of T_1 using pulsed NMR. The system consists of a flow cryostat operating between 4.7 and 20 K, an iron-core magnet running at 0.66 T, a DT-gas handling system, and a commercial NMR spectrometer. In mixtures of hydrogen isotopes that do not contain tritium, the spin-lattice relaxation time is determined by the concentration $c(1)$ of rotating ($J = 1$) molecules. Figure 3-31 shows T_1 for HD as a function of mixed $c(1)$ H₂. Tritium in DT molecules should behave in a similar way as a function of $J = 1$ T₂ concentration. The minimum around $c(1) = 0.01$ occurs because the electric quadrupole moment of the $J = 1$ molecules produces a fluctuating magnetic field that has a maximum spectral density at the NMR frequency at this concentration.

Clearly, we would like to get $c(1)$ as low as possible. One way is to synthesize molecular DT and start with a low concentration of T₂ molecules. Another way is to catalyze the $J = 1$ to $J = 0$ conversion process. This is normally done with a paramagnetic catalyst, but the decay of tritium produces paramagnetic atoms in the solid that also serve this purpose. Figure 3-32 displays a measurement of T_1 as a function of time in normal DT at 5 K in a magnetic field of 0.6 T. Catalysis, which reduces $c(1)$ from about 0.18, explains the rapid initial drop of T_1 . As the concentration falls, the sample goes into the hole around $c(1) = 0.01$ and appears to start up the other side before reaching an equilibrium value. The concentration will not drop to zero because some of the atoms produced by radiation will recombine into $J = 1$ T₂.

The equilibrium value of $T_1 = 0.15$ s may be dictated by the equilibrium value of $c(1)$; however, the paramagnetic impurities are also a relaxation mechanism and they may dominate below some value of $c(1)$. Determining the dominant mechanism is important because polarization is favored if the paramagnetic impurities are the strongest nuclear relaxation mechanism.

As the temperature of the DT is reduced, the number of paramagnetic impurities in the sample at equilibrium increases rapidly,²⁸ and we expect $c(1)$ to drop with decreasing temperature. Measurements of T_1 , coupled with measurements of N_e as a function of sample temperature, should tell us under what conditions the paramagnetic impurities become the dominant relaxation mechanism.

Electron Paramagnetic Resonance System

A major part of our effort during 1985 was devoted to the construction of an EPR system that will allow us to measure T_1 and N_e . The system is an iron-core magnet running at 0.36 T, a Dewar, and a 1.5 to 20 K variable-temperature insert containing a microwave cavity, a data acquisition and control system, a vacuum-support system, a DT-gas handling system, and a 10-GHz microwave spectrometer.

The variable-temperature Dewar insert probe is basically a vacuum can immersed in liquid helium with support plumbing attached to a room-temperature flange. The flange mates to a Janis Dewar mounted on our magnet. Inside the vacuum can, we have a 10-GHz TE₀₁₁-mode microwave cavity. A sapphire cell is mounted inside the EPR cavity to hold the solid DT. Its design is very similar to the cell we are using for NMR work in our flow cryostat. When microwave power is applied to the cavity, a magnetic field is generated with the correct orientation with respect to the external magnetic field to drive electron spin flips in the DT. When the magnetic field is adjusted by sweeping the magnetic power supply so that the spin-flip energy corresponds to the resonant frequency of the cavity, energy is absorbed and this is detected by the EPR spectrometer. An identical coaxial line attached to an NMR coil will allow us to extend our NMR measurement into the temperature range between 1.4 and 5 K. Figure 3-33 is an assembly drawing of the EPR cavity.

In addition to providing a resonant circuit for the 10-GHz EPR measurements, the cavity holds liquid helium that can be pumped down to 1.4 K and temperature controlled between 1.4 and 20 K. This requires a liquid-helium input line and a

Fig. 3-33. Assembly drawing of the EPR cavity.

pumping line. A shutoff valve between the main helium bath and the cold plate has been included to isolate the DT from the main bath in the event of a leak. The main limitation of this refrigerator is that total heat leak must be below ~ 10 mW. The 25 mg of DT that we use supplies ~ 5 mW, so the total heat leak from the supports and electrical leads has to be below ~ 5 mW.

Another line into the cavity for the DT has a vacuum jacket that is an extension of the vacuum can to the room-temperature flange. Thus raising the temperature of the cavity to 20 K will warm the DT line and prevent frozen DT from blocking it. There is also a heat switch between this line and the 4-K bath to prevent excess heat leak into the cavity. When the switch is closed, the line is cooled to 4 K. To open the switch, a heater is turned on to overcome the thermal link to the bath. This allows DT to be moved in and out of the system with liquid helium in the cryostat. The vacuum can must be evacuated to prevent thermal contact between the variable-temperature cavity and the 4.2-K bath. The can is designed so it can be filled with helium and nitrogen gas during the cooldown from 300 K.

Fig. 3-34. Schematic of the EPR microwave spectrometer.

Figure 3-34 is a schematic of the microwave spectrometer. The klystron is tuned to the resonant frequency of the cavity and an automatic frequency control loop operating at 61 kHz locks it to the adsorption maximum of the cavity. Frequency modulation of the klystron produces sidebands 61 kHz above and below the main klystron frequency. These sidebands are reflected from the cavity and detected. Their relative amplitude is a measure of the difference between the klystron frequency and the cavity frequency. A feedback signal is applied to the klystron reflector voltage to keep the difference at zero.

Klystron power is directed around the circulator and into the probe. With proper coupling, no power is reflected back into the arm connected to the magic tee until there is EPR adsorption, so the circulator

functions as a microwave bridge. A fraction of the microwave power is directed through a phase shifter and attenuator in the opposite arm of the magic tee to control the bias on the detector diodes on the remaining arms. The two detector outputs are differentially amplified and sent to a lock-in amplifier. The lock-in reference drives the modulation coil in the cavity. This modulation scheme detects the derivative of the adsorption as a function of magnetic field as the field is swept through the EPR transitions of the sample in the cell.

During the coming year, this EPR system and the measurements of T_1 will let us estimate the ratio of electronic to nuclear polarization^{27,29} that can be achieved. The EPR system will also give us our first look at the NMR behavior of DT below 4 K.

With this information, we will begin to know how difficult the dynamic nuclear polarization of DT will be.

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Targets for the X-Ray Laser Campaign

Introduction

In 1985, the x-ray laser project was divided into two major campaigns, microdot experiments and exploding-foil x-ray laser experiments. Using the systems and processes described in this article, we delivered thin-film microdot coatings of Se, Fe, Ta, Ge, Mo, Si, Eu, Al, CaSi_2 , CaSiO_4 , NaCl, Se + CaSi_2 , and Mo + CaSi_2 on Formvar films for the x-ray laser microdot experiments on the one-beam Phoenix laser facility. In doing so, we acquired a quantity of rare-earth sputtering targets for use in future x-ray laser target development and fabrication. A total of 124 microdot targets on washers and 24 with exposed edge-on views were coated and constructed. We also made a significant effort to develop coatings of Eu, Gd, and Sm to extend the exploding-foil x-ray laser scheme to shorter wavelengths. Related theory and experiments are described in "X-Ray Laser Experiments" in Sec. 4.

During the first half of 1985, we prepared for and performed experiments with the Phoenix laser. These experiments were done to gain a better understanding of the important physics issues that apply to the exploding-foil x-ray laser that was successful in 1984.³⁰ Because the Phoenix experiments were a study of the more fundamental processes of the x-ray laser, the targets needed to be simple from a physics standpoint.

Early in the second half of 1985, we conducted exploding-foil x-ray laser experiments that repeated and extended last year's study.³⁰ The target chamber used in the 1984 Novette x-ray laser experiments was permanently installed in the Nova facility and two of the Nova beams were diverted to it, which allowed us to continue

the x-ray laser experiments under the same conditions that existed in Novette. The primary differences this year are in the lasant materials used and in the length of the targets. Previously, the length of the foil support structure was fixed and the foils were cut to the specific length used for a particular shot. We now can vary the length of the structure to match foil lengths ranging from 0.5 to 5.0 cm. We repeated the shots for the neon-like schemes using selenium and yttrium as the lasant foil in the same way as last year. We are continuing to study the collisional-excitation x-ray laser and are extending it to a nickel-like ionization scheme that requires lasant materials of higher atomic weights. The higher-Z elements used were europium, gadolinium, and samarium, deposited in metallic or oxide form.

Microdot Targets

The experiments on the Phoenix laser system were designed to study the plasma conditions of the exploding foils employed in earlier x-ray laser experiments and those of foils to be used in the future. The targets for these experiments were microdots of candidate x-ray laser material, 20 to 50 μm in diameter and 75 nm thick, which were coated onto free-standing Formvar films. Because the size of the dot makes it difficult to see in the Phoenix alignment system, cross hairs were added to help in coarse alignment.

Two basic configurations were used for the microdot target support. The simpler of the two was the dot and cross hairs coated on Formvar and mounted on a metal washer (Fig. 3-35). Another design, which was to be used in interferometer studies of the exploding microdot, required the microdot to be visible edge-on without interference from the support washer. For this target, we used a U-shaped support base that only contacted the substrate above and below the microdot, leaving both the front and back of the dot visible when viewed from the edge.

To mount the targets on a washer, we coated a glass slide with NaCl, which acts as a release agent. Then we coated the slide with Formvar. The microdots and

Fig. 3-35. Photograph of 20-μm-diam microdot target.

Table 3-4. Materials for microdot targets.

Single elements		Compounds	Ternary mixtures
Se	Mo	CaSi ₂	Se + CaSi ₂
Fe	Si	CaSiO ₄	Mo + CaSi ₂
Ge	Eu	NaCl	
Ta	Al		

cross hairs were deposited through a mask that was in direct contact with the Formvar. Then the Formvar, with the dots and cross hairs, was released from the slide and mounted across various washers with the microdot suspended in the center of the opening. For targets with exposed edges, we had to first mount the Formvar on the support and then coat through the mask onto the free-standing substrate. The first technique is preferable because it allows intimate contact between the mask and Formvar substrate, resulting in a flat thickness profile and good edge definition of the microdot. However, when it was necessary, the second method was used to make exposed-edge targets. The coating process is described in greater detail below.

The geometrical requirements dictated that the microdot targets be free-standing films 75 nm thick and 20 and 50 μm in diameter. The most nearly free-standing film that is practical to produce is one deposited on a thin polymer substrate. We used Formvar films 150 to 350 nm thick as

substrates; the thicker films were needed for coatings that applied higher stresses on the substrate.

We coated some of the required metal films on free-standing Formvar films for last year's x-ray laser experiments,³¹ but not in the extremely small microdot geometry. The two possible methods for achieving a 20- or 50-μm microdot are (1) photoresist chemical etching of the dot material on the Formvar film and (2) coating the dot material onto the Formvar through an appropriate mask. Because we were uncertain about how the chemicals used in photolithography would interact with our films and substrate, and because we expected some handling problems with photoresist etching, we decided to use the masking approach. The materials deposited for the x-ray microdot experiments (see list, Table 3-4) are unusual and difficult to work with.

The masks were designed by the FTF Microfabrication Group in several versions to meet changing microdot specifications, and were prepared by the Electron and Solid State Devices Group of the Electronics Engineering Department. The masks were silicon wafers about 50 mm long, 40 mm wide, and 100 μm thick. Rows of 20- and 50-μm square holes and appropriate cross hair slots were chemically etched into the masks. This method is the same one that is employed in the fabrication of components for the resonant photoexcitation x-ray laser targets described in the 1984 Laser Annual.³² Although a circular microdot was preferable, the type of anisotropic etching used can yield only square patterns. One advantage of this type of etching is that the etched plane forms an angle of 36° with the mask surface, minimizing shadowing during coating (see Fig. 3-36).

We used four different coating processes to deposit the microdots: co-evaporation, sputtering, co-sputtering, and evaporation. The first three systems were modified specifically to meet the needs of the microdot experiments.

Co-evaporation. A co-evaporation system, originally dedicated solely for selenium work, was used for Se, Si, and Ca. It employs two resistance-heating evaporation sources. The pure selenium (Se) and the selenium and calcium silicide

Fig. 3-36. Chemical etching of silicon mask reduces shadowing effect.

(CaSi_2) mixtures were evaporated from a temperature-controlled, covered evaporation boat (Fig. 3-37). The boat was heated to $\sim 200^\circ\text{C}$ to obtain an evaporation rate of 15 to 20 Å/s as determined by calibration. In production runs, a shutter located above the boat was used to time the coating cycles. A thickness monitor indicated the coating rate and film thickness of the selenium being deposited. Small temperature adjustments were usually needed to maintain a constant evaporation rate. The shutter was closed when the thickness reached 750 Å.

One target required the formation of a thin-film microdot that was a uniform mixture of Se, Ca, and Si. After deciding to use CaSi_2 to provide two of the three required elements, we experimented with the evaporation of CaSi_2 in the selenium coating system. Although the reported melting point for CaSi_2 is 1020°C , we never observed melting. Little coating occurred until the sublimation temperature was reached; then the material disappeared abruptly with little control on the rate of coating. Analysis of the film by energy-dispersive spectroscopy confirmed the presence of calcium and silicon, with a trace of impurities. The CaSi_2 film turned transparent after exposure to air, which we believe to be the result of oxidation. We had problems controlling the composition ratio of the elements during co-evaporation. Even with a heat shield to isolate the sources from one another, operation of the CaSi_2 boat changed the evaporation rate of the selenium boat from its calibrated value. Since the single-thickness monitor responded to both sources simultaneously, the evaporation rates could not be adjusted individually. The film thickness was maintained at

Fig. 3-37. Schematic of selenium evaporation system.

750 ± 50 Å. We delivered the Se-Ca-Si microdot films without measuring the composition ratio. Because selenium was added to the CaSi_2 , the ternary compound did not turn transparent after exposure to air.

Sputtering. The second method used to prepare microdots was sputtering. Because it is not a thermal evaporative process, sputtering is ideal for depositing thin films of high-melting-point elements such as tantalum and molybdenum.³³ We also used sputtering to produce the Al, Eu, Fe, and Ge microdots. We used two general-purpose sputtering systems equipped with Sloan sputtering guns that use cylindrical sputter targets. The majority of our existing sputtering targets were cylindrical. The sputtering systems were originally used in establishing coating techniques and parameters for the various microdot materials.

The production of Al, Ge, Mo, Fe, and Ta films is usually routine. However, the thickness of the microdot configuration could not be determined with a witness slide; the slide's film thickness does not correlate well with microdot thickness because there is shadowing from the mask in the microdot coating. Since the argon

sputtering gas scatters the coating molecules, they reach the substrate at various angles rather than at the desired 90° . The rate at which coating molecules pass through the mask depends on the size and cone angle of the mask opening and on the mask thickness. Since the mask geometry and coating conditions are uniform over the mask area, we scanned a sample microdot on the substrate with the profilimeter to determine the thickness of the set of microdots on a given slide.

The preparation of europium microdot thin films was difficult, chiefly because the metal is so readily oxidized in air. Europium belongs to the lanthanide group, commonly known as the rare-earth metals, and it has a strong affinity for oxygen. Normally, europium oxides are in the Eu_2O_3 form. When the target was exposed to air, oxidation began immediately. After a few minutes in air, a deep green-yellow powdery substance, presumably $\text{Eu}(\text{OH})_3$ with a trace of incomplete oxidation products, came off the surface. Because time did not permit the use of special handling procedures and vacuum or inert-gas storage devices, preservation of europium microdots in the metallic state was impossible. The microdots were probably delivered as $\text{Eu}(\text{OH})_3$ (see "X-Ray Laser Targets" below).

Co-sputtering. A co-sputtering system with both dc and rf capabilities was used to coat compound materials. It contains a sputter gun that uses a simple disk target configuration. Most of the required microdot materials can be molded into this shape. Co-sputtering is an alternative way of producing the Mo, Ca, and Si ternary compounds in the form of $\text{Mo} + \text{CaSi}_2$. The molybdenum was sputtered from the cylindrical gun, using deposition parameters from our previous calibration. Calcium silicide can be sputter-coated with either dc or rf. We used the rf disk gun with auto-tuning for more stable operation. After calibration, the ternary compound was easily deposited.

Evaporation. A small, dedicated thermal evaporation unit was used for all salt (NaCl) release layers and for the NaCl microdots. The thickness of the deposit was controlled by weighing the proper amount of salt (determined from calibration runs) and the salt was completely

evaporated for each run. Because the solubility of salt in water prevented us from floating the Formvar with the microdots from a glass substrate, we had to coat these microdots on the Formvar after it had been stretched across the target holder. The free-standing Formvar did not necessarily have intimate contact with the coating mask; however, evaporation is a line-of-sight deposition process and imposes a less stringent requirement on mask-to-substrate distance than does sputtering. Because the salt microdot was deposited directly onto a free-standing Formvar film, the thickness of a microdot was derived from calibration and not from direct measurement.

X-Ray Laser Targets

The mechanical design of x-ray laser targets for 1985 was virtually the same as for 1984. The primary changes are in the design and fabrication of the support structure for the foils, foil-cutting method, target lengths, and target-base design.

The support frame has been changed so that the length of each target is a function of the length of the laser foil. The angle of the upper and lower support bodies was reduced to allow visualization of the laser-gain channel by more diagnostics. With the new computerized milling machine, we can machine all x-ray laser components in-house; this reduces design and fabrication time and gives greater flexibility in overall target design.

Last year, the coated x-ray laser foils were cut to the desired length with a $5\text{-}\mu\text{m}$ -diam glass fiber. This fiber was adequate for controlling target length, but the quality of the cut ends was improved this year by using the focal spot of an argon-ion laser beam instead of the fiber. A microscope objective, used in common with the laser focusing optics and the optics for a video camera, allows us to view the focal spot (where cutting occurs) on a TV monitor. The position of the focal spot is fixed in space and the material to be cut is moved relative to the spot with a 2-axis micromanipulator. This process gives sharp edge definition and controls foil length to an accuracy of $\sim 25\text{ }\mu\text{m}$. Targets were cut to lengths ranging from 0.5 to 5.0 cm.

The mounting base of the target was modified to improve accuracy and simplify attachment of the target to the target positioner. This kinematic arrangement allows repeatable target mounting in three linear and three rotational axes. Figure 3-38 shows 1- and 5-cm-long targets on the new mounting bases. In 1985, we constructed a total of 74 exploding-foil x-ray laser targets; 37 of these were eventually shot on Nova.

The requirements for coating the x-ray laser materials were generally the same as in 1984; films of lasant materials from 50 to 150 nm thick were coated onto thin Formvar films. However, considerable effort was required in order to develop alternative methods for coating x-ray laser targets with some lasant materials, because of their chemical and material properties. Films of lasant materials 25 to 150 nm thick with a tolerance of ± 5 nm were deposited by a physical vapor-deposition process on a thin Formvar film. The films were supported either by wire hoops or by the aluminum x-ray laser support frames. A critical requirement was that the surface finish should be flat and free of wrinkles. During the x-ray laser process, the laser-gain medium is within the cylindrically expanding plasma of lasant material. The gain region is confined to the diameter of plasma where the density gradients are low ($\sim 50 \mu\text{m}$ across during lasing). Wrinkles of approximately that dimension or greater would cause refraction and seriously degrade the x-ray laser operation. This stringent condition posed a serious problem for some of the lasant elements chosen.

The seven elements selected for lasant coating thus far were Se, Ge, Mo, Eu, Gd, Y, and Sm. Three of the elements, Mo, Gd, and Y, have an unfavorable combination of high melting point and low vapor pressure for thermal evaporation (Table 3-5), so the sputtering process was used instead. As mentioned earlier in the discussion of microdot targets, the high susceptibility to oxidation of this group of elements, particularly europium, makes coating onto thin-film substrates very difficult.

The coating processes for the seven individual lasant elements are discussed below.

Selenium. Selenium was the material selected for the majority of the targets. It has

Fig. 3-38. X-ray laser targets, 1 and 5 cm long, on the new mounting bases.

Element	Melting point (K)	Temperature (K) (10^{-3} Torr)
Selenium	490	450 ^a
Europium	1090	805 ^a
Germanium	1213	1505 ^a
Gadolinium	1585	1700 ^b
Yttrium	1799	1748 ^a
Molybdenum	2890	2575 ^a
Samarium	1345	920 ^a

^a Calculated values from Ref. 40.

^b Interpolated value for Gd from Ref. 41.

Table 3-5. Temperatures used in the evaluation of adaptability of seven elements to the evaporation process.

the lowest melting point and a relatively high vapor pressure (Table 3-5); therefore, it has the best properties for evaporative deposition of the lasants chosen. Indeed, we were able to deposit a 50-nm thickness of Se on 25-nm Formvar films while maintaining a flat surface. The evaporation conditions for Se have been described in the previous section on microdot targets. The Se coatings were quite delicate only when deposited on the thinnest (25-nm) Formvar films supported by a wire hoop. During handling and in storage, these particular targets had to be protected from air currents and vibrations.

In the hydrodynamics of an exploding foil, an important property of the x-ray laser film is its ρt , where ρ is the density and t is the thickness. We examined two areas of uncertainty in this property for selenium. First, the red color of the film indicated that it is amorphous selenium with a density of 4.26 g/cm^3 , in agreement with published results³⁴ for our deposition conditions. The deposited Se crystallizes at

a rate that depends on the temperature and substrate. At room temperature, the crystallization occurs within a few weeks on a glass substrate, but takes several months on a Formvar substrate. The glass appears to have more nucleation sites for crystallization. Second, we were concerned about the rate at which the Se re-evaporates in the target chamber. If the vapor pressure is known, the re-evaporation rate R can be calculated from the Langmuir-Knudsen equation

$$R = 2.1 \times 10^8 \frac{P}{\rho} \left(\frac{M}{T} \right)^{1/2}, \quad (12)$$

where R is in nm/h, P is the vapor pressure (Torr), ρ is density of the material (g/cm^3), M is the molecular weight of the evaporating species, and T is the temperature of the evaporant (K). Preliminary calculations and an extrapolation of literature results³⁵ indicated that re-evaporation may be large enough to be of concern, but the uncertainty was high. We determined the evaporation rate by measuring the thickness of a selenium film with a profilometer before and after exposure to vacuum for 64 h. Re-evaporation of 0.25 ± 0.06 nm/h was obtained at room temperature, indicating that film thickness loss would be significant if the sample were evacuated for several hours. Moreover, the re-evaporation rate increases by 2 to 3 times for every 10°C increase in temperature,^{34,36} so the sample is fairly sensitive to illumination by intense light while in the target-chamber vacuum.

Germanium. Germanium, the next material used for target construction, has a higher melting point and lower vapor pressure than Se, Eu, and Sm, but evaporation deposition can still be used provided that the thickness of the Formvar substrate is >100 nm. The Ge deposits were normally flat and not susceptible to oxidation at room temperature.

Molybdenum. Molybdenum, which was used for shorter-wavelength x-ray laser targets, is a high-melting-point refractory metal, which makes sputtering more convenient for deposition. We first tried dc sputtering at 200 W of beam power with an argon pressure of 10 mTorr. Although the sputtering process proceeded well, the coatings on 150-nm Formvar films

mounted on target support frames were very wrinkled. To see whether depositing the metal in the oxide or nitride form would solve the problem, we first tried reactive sputtering with oxygen, but there was no apparent improvement in the surface quality of the coating. Further reactive sputtering experiments with nitrogen significantly reduced wrinkling of the film. The Mo coatings prepared as nitrides stayed wrinkle-free in air at room temperature indefinitely, indicating that the compound of molybdenum and nitrogen included in the coating must be nonhygroscopic and chemically stable. According to Samsonov and Umanskiy,³⁷ Hagg established the presence of three intermediate phases in the Mo-N system: the β phase, Mo_3N (5.4 wt.% N), the γ phase, Mo_2N (6.75 wt.% N), and the δ phase, MoN (12.73 wt.% N). The authors also commented that the solubility of nitrogen in molybdenum is very low and that when heated in air, molybdenum nitrides are oxidized to MoO_3 . The second comment may imply that the molybdenum nitrides are otherwise stable in air at room temperature, in agreement with our observation.

Europium. Europium, according to the data in Table 3-5, should be the second easiest element of the seven to deposit by evaporation. Our experiments in a simple evaporation system equipped with resistance-heating confirmed this. Before the system was opened, the Eu coating had a typically metallic appearance with few or no wrinkles. But as soon as the system was opened, the Eu films began to oxidize. Eu films several hundred angstroms thick were completely oxidized in seconds or minutes; those a few microns thick oxidized in a few days. During the oxidation processes, the color first turned bright yellow followed by a pale yellow or yellow-green color. Eventually the film turned white, with some pink coloration in spots. Our x-ray diffraction measurements detected the initial formation of EuO , which quickly converted to $\text{Eu}(\text{OH})_3$ with a trace of intermediate EuOOH . Infrared spectroscopy indicated that absorbed water was present in the initial stages, and hydroxyl and carbonated products were observed at all times. The relation between the colors and the intermediate products is not clearly understood, but one source³⁸

attributes the yellow color to $\text{Eu}(\text{OH})_2 \cdot \text{H}_2\text{O}$.

Oxidation of the Eu films caused more severe wrinkling, and also led to large increases in film thickness. In one case, the film had a nominal thickness of 100 nm, but thickened to 600 nm after oxidation was complete, indicating that the final $\text{Eu}(\text{OH})_3$ was quite porous. Such severe distortion made the films unusable as x-ray laser foils.

In view of the Eu oxidation problem and our success with the reactive sputtering of molybdenum, we tried rf sputtering of Eu_2O_3 . Although the oxide deposits were stable at room temperature, the coatings were wrinkled, largely because of heating induced by the sputtered atoms. Unsintered rare-earth oxides form carbonated hydroxides when they are exposed to the atmosphere.³⁹ These processes occur faster with europium than with the other laser candidate materials. Hence, even if we can successfully deposit the less dense oxide by sputtering, there will still be a volume expansion from the absorption of water and CO_2 once the deposit is put into a room environment.

Gadolinium. Gadolinium was another candidate laser material. From the data listed in Table 3-5, we determined that our simple evaporation system would be incapable of handling this element, and we resorted to the sputtering process. In the first coating run with a cylindrical sputtergun, we found that gadolinium could easily be sputtered in either the rf or the dc mode. Unfortunately the Gd-coated Formvar film was severely wrinkled. We tried reactive sputtering with both oxygen and nitrogen with mixed results. For example, dc sputtering under argon pressure of 9.3 mTorr and oxygen pressure of 0.1 mTorr produced a Gd film that showed no visible wrinkles, but it began to sag and wrinkle after 4 hours exposure to air. Sputtering with nitrogen pressure at 0.5 mTorr and argon pressure at 9.8 mTorr reduced the severity, but wrinkling of the Gd film still persisted. The Gd coating deposit is unstable, whether the metal is in the metallic state or is doped with oxygen or nitrogen because of the absorption of water and CO_2 from the room environment.

Yttrium. Yttrium was one of two elements shown to produce soft x-ray lasing

in the 1984 x-ray laser experimental series. The yttrium films were wrinkled, however, causing some uncertainty about the exact location of the plasma plane when the yttrium-coated Formvar was illuminated by the Novette laser beam. The yttrium films were smooth and mirror-like before the sputtering chamber was vented to air, indicating that wrinkling was caused by oxidation. To meet the present requirement of a wrinkle-free surface, we attempted reactive sputtering with oxygen and nitrogen under a variety of conditions of power and pressures. In all cases, films of either yttrium oxide or nitride were produced. However, the Formvar films could not survive a Y_2O_3 coating ~30 nm thick without rupturing. Examination of the unbroken area of the Y_2O_3 -coated Formvar film showed that wrinkling was reduced, but still persisted. Similar results were obtained from reactive sputtering with nitrogen. We have not exhausted all the combinations of the operational parameters of the sputtering system, and are continuing developmental work.

Samarium. Samarium has a moderate melting point and a sufficiently high vapor pressure to be evaporated in our simple system. Our first efforts with this lanthanide have been mixed. The metal-coated Formvar films were smooth before exposure to air, and the wrinkling was somewhat less severe than that of other rare-earth coatings. The Sm coatings retained their metallic luster for several days to a few weeks, depending on the thickness of the coating.

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Physics Targets

Direct-Drive Shock Plates

Introduction. We have developed two approaches to fabricate direct-drive shock plates: two-step deposition with an in situ movable mask and micro-machining an aluminum layer deposited on a precision-machined mandrel. Shock plates produced

by both methods are suitable for measuring laser power in current experiments. Potentially, the precision-machined mandrel approach can yield higher quality shock plates at lower cost. We will decide which approach to adopt when more data are available to compare the properties of the shock plates produced by each method.

When a laser beam impinges on the surface of a material, very high pressures can be generated by the accompanying shock wave. The rate at which the wave travels in the material is an indication of the pressure generated by the laser pulse. We use a streak optical pyrometer (see "Laser Plasma as an X-Ray Backlighting Source" in Sec. 4) to record the onset of light emission created by the shock breakout in two known thicknesses of aluminum—the witness or shock plate. Given the time differential between the onset of light emission in the two layers, we can calculate the pressure and indirectly the laser power.

To provide samples for these shock-plate measurements, we developed coating and machining processes to manufacture well-characterized aluminum disks. The aluminum disks are made of two layers of different thickness with a sharp transition step between the layers. The design of the samples, which considers the laser energy, instrument resolution, and known material properties, optimizes the time differential of light emission in the layers. Since a $1\text{-}\mu\text{m}$ discrepancy in the thickness causes a significant error in the measurement, the surface finish must be less than $1\text{ }\mu\text{m}$ thick. The material also needs to be full density to apply the known equation-of-state data.

We previously used either a two-step physical vapor deposition process or chemical etching to produce shock plates. We had difficulties with both methods.

During the two-step deposition, we had to break the vacuum to move the mask inside the chamber that created the step. The two aluminum layers frequently separated during subsequent processing, resulting in a low yield for the two-step deposition process. We attributed the layer separation to the oxide formed on the first aluminum layer during the vacuum break. We also experienced problems with the chemical etching process. It produced a poor surface finish and made the aluminum foil too soft to use in later assembly processes.

Two-Step, Movable Mask Deposition Process. In the first of two new approaches, the oxide formation at the interface is eliminated with the use of an in situ movable mask; no vacuum break is necessary during deposition. A mechanized mask over the substrate is driven by a vacuum-sealed electric motor inside the vacuum chamber. The mask movement is designed to maintain a specified clearance and thus avoid damage to the surface of the metal coating already deposited. Multistep shock plates can also be made with this approach. This movable mask with an electric motor drive was constructed and proved successful; the design is shown in Fig. 3-39. We used both copper and copper-coated glass substrates. The mask is positioned at the inside edge of the $1 \times 3\text{-in.}$ substrate for the first deposition; the clearance from the substrate is set approximately $50\text{ }\mu\text{m}$ higher than the first deposition thickness. The mask is then driven to cover part of the first layer before the second aluminum layer is deposited, forming the required step. A well-defined step (Fig. 3-40) with less than $50\text{-}\mu\text{m}$ spread can be obtained.

We used an e-beam evaporation system for the aluminum depositions. The deposition system was calibrated to achieve a uniformity of 5% thickness within a 3-in.-square area of the aluminum layer from which the shock plates are to be cut. The thickness variations within a millimeter-size shock plate should be 0.1%, a differential too small for measurement. We measured the density of the deposited aluminum foil by fluid immersion methods and confirmed that it is indeed full density.

We still have some problems. In spite of a nearly 90° incident-coating angle and high vacuum (5×10^{-6} to 5×10^{-7} Torr

Fig. 3-39. In situ movable mask process for making shock plates by the two-step deposition of aluminum coating on a substrate.

Fig. 3-40. (a) Top and side-section views of shock-plate geometry. (b) The SEM shows a well-defined thickness step in the aluminum deposition layers.

pressure), the deposited aluminum still has a columnar structure, giving rise to a milky appearance on the surface. The resultant surface roughness is $1\text{-}\mu\text{m rms}$ for the thicker deposits, which barely meets the allowable shock plate tolerance.

We have produced 30/50-, 50/80-, and 75/150- μm aluminum layers to be machined into final shock plate configurations. The electric discharge machine (EDM) was used to cut out various shock plates while the aluminum deposit was still attached to the substrate. The copper interface was chemically etched away later to obtain the shock plate. We found the copper substrates better than glass substrates for cutting by the EDM process. Occasionally, a grain was dislodged from the aluminum surface during the etching process, leaving a $1\text{-}\mu\text{m}$ hole on the surface.

Deposition on Machined-Mandrel Surface.

Our second approach was to provide a machined-mandrel surface on which a thick aluminum layer is deposited. By following a prearranged procedure, the mandrel can be re-indexed to obtain the proper thickness when the steps are back-machined from the deposited aluminum layer. The copper mandrels for our initial trial were 2-in. disks with a flat machined surface for the aluminum deposition. The disk is $1/2$ in. thick and the indexing is accomplished by gauging at the circumference.

The main uncertainty of this mandrel approach was the machinability of the deposited aluminum. Pure aluminum is soft and a difficult material for micro-machining. Pure aluminum must be used,

however, because only the pure aluminum material constants are well characterized for hydrodynamic calculations. The use of pure aluminum also prevents the introduction of O_2 or N_2 gases that can retard columnar growth during the e-beam deposition of aluminum. We used the e-beam evaporator to deposit an $80\text{-}\mu\text{m}$ -thick aluminum layer on a flat surface mandrel for re-indexing and machinability testing. The result was encouraging; the mandrel could indeed be re-indexed to cut thicknesses accurately and the machined surface was better than the deposited aluminum surface. We were still concerned about the columnar structure, however. Because of this uncertainty, we decided to establish an alternative to e-beam evaporation for depositing the thick aluminum layer.

A high-rate magnetron sputtering system became operational in November 1985. The system has a 5-in. sputtering target that has an aluminum deposition rate of $1\text{ }\mu\text{m/min}$ with better than 85% uniformity on a 3-in.-square substrate 2 in. from the source. Our experience with a 1-1/2-in. sputtering target of the same design showed that a thick ($>100\text{-}\mu\text{m}$) aluminum layer can be deposited with no columnar growth. Preliminary experiments with the 5-in. gun showed similar results. We have deposited a $150\text{-}\mu\text{m}$ aluminum layer onto a flat-surface copper mandrel for back-machining.

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Target Prealignment

It became apparent, with the activation of the Nova laser, that some form of prealignment was necessary to reduce the alignment time in the target chamber. The goals of prealignment are to reduce the amount of time required to align a target in the target chamber and to improve the accuracy of alignment. The "prealignment" task has evolved to include consideration of every step in the development of a target up to the time it is shot: design, where provision for multiple alignment features must be made; micro-machining of flags, fiducials, and shields; assembly of target components; inspection; emulation of Nova target alignment viewers (TAVs); and finally, TAV image matching. TAV emulation was chosen over other alignment concepts because it was a logical evolutionary step that imposed few special requirements on existing equipment.

Target Alignment Difficulties

Currently, there are three TAVs on each of the two Nova target chambers. The TAVs define a functional target chamber center, defined by periodically aligning the TAVs on a small object placed near the geometrical center of the chamber. Once the TAVs are aligned, or aimed at a common point, any object can be brought to that point and an alignment technician can verify that it is at the functional target chamber center. This approach to aligning targets is somewhat subjective if the actual target is hidden by flags and shields from the view of the TAVs. Careful planning and skillful fabrication are required to prevent obscuration of some target features by others. Frequently, there is no single point on a target that can be viewed by all the TAVs. The only distinctive target features visible to the TAVs can be as much as 1000 times larger than the desired positioning accuracy. The depth of field of the TAVs can also cause some uncertainty, since it is at most a few millimeters. TAV images of a properly aligned target often change from shot to shot because target

configurations are changed frequently. The subjectivity and uncertainty that stems from these difficulties can result in time-consuming alignment procedures and unknown alignment accuracy, thus some improvements were essential for quality experiments.

TAV Emulation

We elected to build a system of TAV emulators after review of the few proposals for improving the target alignment process. The emulators created few special requirements on target construction and minimized the changes required on the target chambers. The three TAV emulators, built on an optical table (Fig. 3-41), provide views of the target that duplicate the views from the TAVs on the target chamber. The emulator optics match the optics of the target chamber TAVs as closely as possible. The target is mounted on a precisely calibrated 4-axis stage that allows a technician to locate the target at a center point defined by the emulators. A fifth axis allows the target to be moved so that TAV emulators can simulate the views from most diagnostics.

Target images from each of the three TAV emulators will be digitally recorded and stored at various magnifications on a dedicated minicomputer. The recorded images are delivered with the target to the Nova facility on a removable tape or hard disk, and displayed on three corresponding TAV monitors with the use of another minicomputer. Software and specialized digital image-processing hardware (Imaging Technology FG-100) overlay the images of the TAV emulator reticle and the TAV reticle on the monitor. Once the reticles coincide, the Nova target positioner is used to move the target until the TAV image of the target matches the emulator image. Image subtraction is used to show the differences between stored and real-time images.

Target Assembly Process

Tolerances on each target component (flag, backlighter, shield) become critical because several components must be

Fig 3-41. Three TAV emulators, on an optical table with a rigid frame, duplicate the TAV views of a properly oriented target in the Nova target chamber.

assembled in close proximity without interfering with TAV or emulator views. Future targets will have more components than past targets and will require even smaller linear and angular tolerances. To improve the micro-machining facilities and procedures that produce these components, a 5-axis calibrated manipulator was designed for a micro-machining drill press that gives a 10-fold improvement in the accuracy of hole placement.

Precision target assembly has relied heavily on the exceptional skill of a very few technicians. These people have learned to assemble targets with greater accuracy than the nominal precision of the equipment they work with. However, to consistently achieve the tolerances of current and future targets, we need to make improvements in the assembly equipment.

One of the most difficult aspects of what we perceive as the next step in mechanizing the assembly process is to orient each component with respect to a set of reference coordinates. Microscope

reticles, which have provided a convenient optical reference in the past, will remain an essential part of the assembly equipment. However, new methods of certifying the orientation of each component are being reviewed to improve the accuracy of orientation with respect to three axes of rotation and the corresponding linear axes. Methods that excessively inhibit the assembly process must be avoided. Precision fixtures that will reduce the required number of component manipulation axes and speed the assembly process are being studied.

Measurements made during the assembly process can be used to assure that the target is properly located at a point defined by the TAV emulators. Figure 3-41 shows the 4-axis positioner that places the target in the correct orientation for viewing by the TAV emulators. These 4 axes correspond to the 4 axes of the target positioner on the Nova 10-beam chamber. Note that the target must mount to the positioner in a way that guarantees proper rotation about the N/S and E/W axes of the

Target Prealignment

10-beam chamber because there is no active control over that mode of rotation. A large goniometer, beneath the 4-axis emulator positioner, allows the target to be rotated so that one or more of the TAV emulators can simulate a diagnostic or beam view of the target. Design, fabrication, and assembly of the TAV emulators is

complete and certification of emulator alignment accuracy is now under way.

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The photo, right, shows a time-resolved measurement of a shock velocity using a stepped aluminum witness plate. The other side of the plate is irradiated by two $0.35\text{-}\mu\text{m}$ laser beams. The shock breakout is seen first on the right side from the $12\text{-}\mu\text{m}$ -thick aluminum and then after 0.67 ns , it breaks out on the left side from the $111\text{-}\mu\text{m}$ -thick side. This shock velocity measurement gives a laser ablation pressure of 66 Mbar . A timing fiducial on the right-hand side also gives the velocity of the shock and shows it is constant.

(Opposite page, below) High-resolution transmission electron microscope (TEM) image of alternating $\sim 50\text{-\AA}$ -thick layers of molybdenum (dark bands) and silicon (light bands). This multilayer structure provides efficient reflectivity for soft x rays at normal incidence.

(Below) High-resolution TEM image of the interface between a layer of Si_3N_4 deposited on a single-crystal silicon substrate. The atomic lattice planes of the crystalline silicon are evident.

In subsequent processing, the silicon substrate is etched away leaving the thin Si_3N_4 membrane supporting a multilayer structure as discussed above. This device serves as a soft-x-ray beamsplitter.

Experiments and Diagnostics

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Experiments and Diagnostics

Introduction

E. M. Campbell

This year the primary activity of the Laser Experiments and Advanced Diagnostics Program (LEAD) has been to prepare the Nova target areas for the sophisticated inertial fusion, weapons physics, and extreme ultraviolet (XUV) laser experiments that will be fully under way in early 1986. Our main tasks include the design, fabrication, installation, and activation of numerous diagnostics on both the ten-beam and two-beam target chambers. When complete, Nova will be the world's most sophisticated and versatile experimental facility for the study of high-energy, high-density plasma physics.

We have made sufficient progress in these tasks to initiate a series of experiments that serves to activate the laser facility and diagnostic systems and to address some relevant physics issues. These Nova experiments are described in "Laser-Plasma Experiments." The first three articles discuss the physics of high-Z plasmas irradiated with both 0.53- and 0.35- μm laser radiation. The multiterawatt capability of Nova at both wavelengths allows us to address such issues as absorption, x-ray conversion efficiency, suprathermal electron production, and parametric instabilities in large plasmas (with scalelengths of 10^2 to 10^3 laser wavelengths) at intensities of interest to laser fusion. These experiments also serve to activate several x-ray and optical diagnostics and allow us to compare the results with those obtained at the Novette facility.

We are continuing with the successful experiments that use exploding foils to study parametric instabilities, such as stimulated Raman scattering in long-scalelength, fully underdense plasmas. In addition to activating many optical diagnostics, our goals for such experiments, described in "Exploding-Foil Experiments," are to study collisional damping of stimulated scattering at 0.35- μm irradiation using both plastic and gold foils and to clearly identify the occurrence of forward Raman scattering.

The early experiments on the Nova two-beam target chamber have been primarily focused on XUV laser studies. The experiments to date, described in "X-Ray Laser Experiments," serve to activate numerous sophisticated diagnostics and are a continuation of the successful experiments conducted on Novette in 1984. We have extended the length of selenium exploding-foil targets to 2.5 cm, and we have continued to observe exponentiation of the intensities of the 206- and 209- $\text{\AA}$ lines from neon-like selenium that were observed previously. Due to the longer amplification length of the plasma and improved detector sensitivity, we have also observed gain at three additional lines from the same 3p manifold at 180, 220, and 260 $\text{\AA}$. The line at 180 $\text{\AA}$ is of particular interest because its upper state has $J = 0$ and was predicted before the experiments to have the largest gain of any of the 3p to 3s transitions in neon-like selenium. The 180- $\text{\AA}$ line was not unambiguously observed in earlier experiments.

One of the goals of the laboratory XUV program is to extend lasing to shorter, more penetrating wavelengths. One scheme that we are examining is collisional excitation of excited states in nickel-like ions. Sophisticated simulations indicate that there are numerous transitions in nickel-like gadolinium that have significant gain when irradiated with moderate-intensity

($I \sim 0.5$ to 1×10^{14} W/cm²) 0.53- μ m light. The lines resulting from transitions between the 4d-4p states have wavelengths in the range of 80 to 160 Å. We have initiated a series of experiments to determine the spectrum of exploding-foil gadolinium targets in the vacuum ultraviolet (VUV) region of the spectrum and to confirm the presence of nickel-like ions by observing the ~ 1.5 -keV line emission from $n = 4$ to 3 transitions.

We have continued our successful collaboration with KMS Fusion, Inc. by conducting joint experiments to measure the electron-density profile using four-frame holographic interferometry from a variety of exploding-foil target designs and irradiation conditions. This work, described in "Plasma Profile Experiments," serves as important checks on the hydrodynamic modeling of the XUV laser design and permits quantitative analysis of the effects of refraction on radiation propagation through the plasma.

Experiments to characterize laser-driven, x-ray backlighting sources at <100 ps have also been performed as a joint effort at KMS Fusion, Inc. A variety of future hydrodynamics and spectroscopic experiments will require intense x-ray sources at <100 ps and at a number of photon energies. The short duration of the source is required to minimize motional blurring of experimental data. These experiments demonstrated bright x-ray bursts with time durations <50 ps at photon energies as high as 4.7 keV and are presented in "Laser Plasma as an X-Ray Backlighting Source."

A knowledge of the plasma conditions of a laser-heated target is essential for experiments in laser fusion, weapons physics, and x-ray lasers. To determine such quantities as electron density, temperature, and ion charge-state distribution, we have begun a program to develop absorption spectroscopy using laser plasmas as the x-ray source. The early, encouraging results are discussed in "X-Ray Absorption Spectroscopy." Finally, in the article, "Laser-Driven Impulse Experiments," we present results from a set of experiments conducted on the two-beam Nova target area. Here, we studied the impulse imparted to aluminum targets irradiated with intensities in the range of 10^{10} to 10^{12} W/cm² in a nanosecond pulse of 0.53- μ m light.

Several new technologies are presented in "Diagnostics." These articles describe our efforts to place diagnostic systems on the Nova target chambers and to develop advanced x-ray and neutron instrumentation and technology. Highlights include the fabrication and testing of multilayer mirrors for x-ray laser cavities, the development of small, bright low-power x-ray sources, and the operation of a direct-recording, subkilovolt, x-ray charge-coupled device (CCD) camera. The latter two technologies will be extremely useful in characterizing new x-ray optical components.

We are also continuing to devote sizable resources to developing the capabilities discussed in "Time Resolved Measurements." Of particular note are the advances made in x-ray framing cameras and in an optical fiducial system for x-ray and optical streak cameras.

High-performance capsules on Nova will require sophisticated neutron diagnostics to characterize the imploded fuel state. Our intensive programs to develop diagnostics that can measure the spectra, time history, and spatial extent of penetrating fusion products are presented in "Neutron Diagnostics."

The experimental programs on Nova demand a system for the rapid processing and archiving of data; the Data Management and Analysis Group (DMAG) in LEAD provides this service. DMAG activities for the year

are summarized in the final group of articles in this section. Highlights include the activation and on-line operation of a computer-based electronic shotbook that enables scientists in the program to have rapid access to experimental data from Nova in a standardized presentation format.

Laser-Plasma Experiments

Nova Experiments on Laser-Matter Interaction

Introduction. In 1985, experiments on Nova investigated laser-matter interactions while initializing newly installed diagnostics and subsystems. These include beam diagnostics, a target-positioning system, and equipment for laser focusing and pointing. Our experiments used frequency-tripled 0.35- μm laser light to study the wavelength scaling of light coupling in high- Z plasmas. Use of this frequency extends earlier studies of absorption and underdense plasma started on Novette at 0.53 and 0.26 μm . Preliminary results agree with the scaling derived from Argus and Novette experiments.

Energy-balance experiments, in which 0.35- μm light interacted with thick gold disks, indicate high absorption and low suprathermal-electron production arising from the high intensities and larger spot sizes available on Nova. Absorption remained higher than 90% at 10^{15} W/cm^2 , indicating that stimulated Brillouin scattering is not effectively competing with inverse bremsstrahlung absorption. Also, only a small fraction of the laser energy is absorbed by parametric instabilities that produce suprathermal electrons. These results are described in the article "Suprathermal Electrons."

Preliminary measurements show that stimulated Raman scattered light yields correlate with fast-electron yields. These measurements agree with Novette results, indicating that stimulated Raman scattering (SRS) is the principal source of fast electrons. Experiments to date have shown SRS efficiencies to be in the range 10^{-3} to 10^{-2} . These experiments are described in the article, "Stimulated Raman Scattering from Gold-Disk Targets."

Low-energy x-ray conversion efficiencies with 0.35- μm laser light are slightly higher than yields at 0.53 μm measured on Novette. These results, predicted by wavelength scaling, are described in the article "X-Ray Conversion Efficiency."

Preliminary exploding-foil experiments at 0.35 and 0.53 μm investigated thresholds for underdense plasma instabilities, particularly SRS. These are described in the article "Exploding-Foil Experiments." Of particular note is our observation of forward SRS during high-intensity, 0.53- μm experiments. This phenomenon has not been unambiguously observed in earlier experiments. The observation of both up-shifted (anti-Stokes) and down-shifted SRS components unequivocally established the presence of forward SRS. Supporting evidence came from suprathermal electrons ($>130 \text{ keV}$) produced by high-phase-velocity plasma waves, which the forward SRS would be expected to generate.

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Suprathermal Electrons. The production of suprathermal electrons was studied in earlier work¹ and found to be related to SRS, which is dependent on laser intensity and wavelength. No Nova experiment to date has focused primarily on suprathermal-electron production; however, some preliminary data were obtained during a recent series intended to measure x-ray conversion efficiency. (See the article "X-ray Conversion Efficiency," later in this section.)

In this series, gold-disk targets, typically 20 μm thick, were irradiated with 0.53- μm light at intensities of about $4 \times 10^{14} \text{ W/cm}^2$ and with 0.35- μm light at intensities from 2×10^{14} to $4 \times 10^{15} \text{ W/cm}^2$. For the 0.53- μm measurements and one 0.35- μm measurement, 1-ns Gaussian pulses were used. The remaining 0.35- μm measurements employed 1-ns square pulses. All

shots were at diverging focus ($f/4$), giving more uniform target irradiation. Typical energies were 1 to 2 kJ. Spot sizes used in

Fig. 4-1. FFLEX high energy x-ray spectrum; f_{hot} is the fraction of suprathermal electrons, and T_{hot} is the suprathermal-electron temperature.

Fig. 4-2. Inferred fraction of laser energy in suprathermal electrons as a function of average laser intensity.

the intensity calculation came from x-ray microscope pictures ($h\nu \sim 1$ to 216 V).

The filter-fluorescer diagnostic (FFLEX)² measured the x-ray spectrum from 9 to several hundred keV. The resulting data were analyzed as in previous Novette experiments.³ Analysis involved fitting a two-temperature Maxwellian spectrum that had thermal and suprathermal components, and extracting the suprathermal component (effective temperature of 20 to 35 keV). Figure 4-1 shows a typical spectrum exhibiting this component. (In addition to thermal and suprathermal components, the data analysis showed evidence of a third, "superhot" component, which must be further analyzed.) Suprathermal-electron production was calculated using a thick-target bremsstrahlung formula.

Results of the Nova measurements were plotted with previously obtained Novette data³ as shown in Fig. 4-2. The preliminary 0.35- μm results are widely scattered. The 0.53- μm results are consistent, but indicate somewhat more suprathermal-electron production than observed on Novette. The experiment with the 0.35- μm Gaussian pulse produced a low proportion of suprathermal electrons, but the result is still within the scatter of the other 0.35- μm data.

The wide variation in the suprathermal-electron fraction is possibly related to the fact that these experiments are close to the SRS threshold. (Hot-electron fractions are well correlated with SRS, as discussed in the article "Raman Scattering from Gold-Disk Experiments," later in this section.)

Differences between Nova and Novette data might be related to differences in beam uniformity or to the use of converging beams on Novette and diverging beams on Nova.

It is important to recognize that the Nova results are preliminary—obtained during experiments that were optimized for a different measurement. More specific experiments are required for a better understanding of suprathermal-electron production.

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Raman Scattering from Gold-Disk Targets. We have irradiated gold-disk targets using 1-ns Gaussian or square pulses of 0.35- or 0.53- μm light. We previously reported measurements of Raman scattering from gold-disk targets irradiated at 0.53 μm by Novette⁴; this article reports new measurements from Nova using 0.35- μm irradiation. The experiments discussed here are also described in detail in the immediately preceding and following articles.

We measured the amount of Raman scattering with approximately 40 filtered, calibrated photodiodes. The diodes were filtered to detect broadband Raman scattering at wavelengths from 550 to 700 nm. Although it is possible to obtain Raman scattering from wavelengths below 550 nm, the time-resolved Raman spectra obtained on a few of these gold-disk shots did not show significant scattering at such short wavelengths. The diodes measure the J/sr of Raman light at several locations. The angular distribution of the scattering will be discussed in future reports. We assigned a solid angle to each diode (based on its distance from adjacent diodes) to obtain the total energy of the Raman light. We divided this energy by the incident laser energy to obtain the fraction of Raman scattered light.

Figure 4-3 shows the fraction of hot electrons as a function of the fraction of Raman scattered light. The hot electron fraction is defined and discussed in the preceding article. Most of the data were obtained using 0.53- μm irradiations with Novette; these data are reported in Ref. 4. We have added data from gold-disk targets irradiated at 0.35 μm by Nova during 1985. All but one of the Nova experiments used a square, 1-ns pulse. The Nova data lie along or slightly above the curve (solid line in Fig. 4-3) for the Novette data.

The essential result of our early Nova experiments using 0.35- μm light is that Raman scattering and hot electron production are both still present and well correlated. We believe Raman scattering to be the dominant source of hot electrons in all these experiments. Based on our 0.35- μm data obtained from Nova to date, we require about twice the laser intensity that

was required at 0.53 μm to produce a given fraction of Raman scattering.

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Fig. 4-3. Hot-electron fraction from x-ray measurements vs Raman scattered light fraction. Arrows indicate that some Raman detectors were saturated, and the data point shows a lower limit.

X-Ray Conversion Efficiency. The efficiency of converting incident laser energy to low-energy x rays from gold-disk targets is a measure of the coupling efficiencies of the laser energy with high-Z plasmas. We measured conversion efficiencies at laser wavelengths of 0.53 and 0.35 μm to check out the Dante spectrometer on Nova and to obtain preliminary results for planning future experiments.

The Nova Dante system measures x-ray fluence below 1.5 keV, the region in which most of the x-ray energy is emitted. We compared results at 0.53 μm with measurements made at Novette under similar irradiation conditions to verify the instrument's operation.

Preliminary measurements with Nova at 0.35 μm continue an investigation of the effects of wavelength on conversion efficiency. Earlier work took place at Argus and Novette.^{5,6} The Nova results indicate that x-ray conversion efficiency at 0.35 μm agrees with the earlier wavelength-scaling studies.

Experiment Parameters. All the conversion-efficiency experiments employed single beams of the 10-beam Nova laser, ranging in energy from 0.5 to 2.0 kJ. In most experiments, we used 1-ns square pulses. These provided nearly constant beam power on target, simplifying later analysis and optimizing the conversion of the 1.05- μm beam to the 0.35- μm wavelength. In other experiments, we used Gaussian pulses at 0.53 and 0.35 μm for comparison with previous data.

We varied the beam spot diameter on target from 200 to 1500 μm by defocusing

the input f/4 lens from 0.8 to 4 mm. This range of spot sizes yielded average intensities from 6×10^{13} to 2×10^{15} W/cm². To determine spot sizes, we compared x-ray images from the 8 $\times$ microscope (correcting for microscope viewing angles) with sizes predicted by geometrical optics. Average intensity was derived by dividing the incident laser energy by the total irradiation area and pulse width.

For all experiments, we used a diverging beam because intensity on the diverging side of best focus is more uniform than on the converging side. Variations in incident intensity due to the spatial structure of the beam were not assessed.

The targets were gold disks, 2 mm in diameter and 25 μm thick. The beam was incident at angles varying from 0 to 40° with respect to the disk normal. Nova Dante spectrometers viewed the disk normal at angles ranging up to 66°, depending on the disk orientation.

Dante Spectrometer. The Dante spectrometer at Nova is a broadband transmission filter and x-ray diode system that was used on Argus, Shiva, and Novette to measure the low-energy x-ray spectrum from laser-produced plasmas.^{7,8} The filters and detectors comprising the system's ten low-energy x-ray channels are listed in Table 4-1, and channel response as a function of energy is plotted in Fig. 4-4(a). The channels are designed so that most of their response occurs in a narrow energy band just below the absorption-edge energies of the filters (also listed in Table 4-1).

The four lowest-energy channels are equipped with x-ray mirrors to suppress signals from energies above the absorption edge of the filter. However, x rays with energy above the absorption edges can contribute significantly to the signal detected by higher-energy channels 5 through 8. Gold M-band emission occurs in the region from 2 to 4 keV, at which most of the filters become transparent.⁹ Figure 4-4(b) shows the responses of channels 5 and 8. It is not practical to use x-ray mirrors on these channels because the mirrors would have to be very long and set at a highly grazing angle of reflection. These requirements would make fabrication and alignment difficult.

To correct for the high-energy response, we designed "background" channels

Table 4-1. Nova Dante configuration.

Channel No.	XRD cathode	Filter edge (keV)	Filter material	Filter thickness ($\mu\text{g}/\text{cm}^2$)	Mirror material	Mirror angle
1	Aluminum	0.19	Parylene Boron	17	Carbon	5.0°
2	Aluminum	0.28	Parylene	880	Beryllium	3.5°
3	Aluminum	0.52	Parylene Vanadium	25	Carbon	2.4°
4	Aluminum	0.58	Parylene Chromium	17	Carbon	1.9°
5	Aluminum	0.71	Parylene Iron	17		
6	Aluminum	3.79	Parylene Cobalt	17		
7	Chromium	0.94	Parylene Copper	17 1010		
8	Nickel	1.02	Parylene Zinc	17 1625		
9	Nickel	1.30	Parylene Magnesium	23 5860		
10	Nickel	1.56	Parylene Aluminum	23 5020		
11	Aluminum	—	Parylene Iron Chromium	535 360 365		
12	Aluminum	—	Parylene Cobalt	1035 1165		
13	Nickel	—	Parylene Zinc Cobalt Vanadium	515 535 910 490

Fig. 4-4. X-ray energy response of the Nova Dante channels: (a) in the low-energy region and (b) in the high-energy tail.

(11 through 13 in Table 4-1). These channels do not respond in the region of interest, but do have a response in the 2- to 4-keV range similar to that of the other channels. Two of the background channels are shown in Fig. 4-4(b) compared with their corresponding low-energy channels. In data reduction, corrections obtained from the background channels were applied to signals from channels 5 through 8. In some experiments, corrections greater than 50% were required for channels 5 and 6. Corrections of 20% or less were required for channels 7 and 8.

Data Analysis. We analyzed the data by unfolding a spectrum from the corrected channel-response signals.¹⁰ The integrated fluence in the low-energy portion of the spectrum was derived by integrating the spectrum from 0.1 to 1.5 keV. The angular distribution of the x rays was not measured. To infer the total energy in low-energy x rays emitted from the disk, we assumed that the angular distribution was emitted from a planar region or acted as a Lambertian radiator. If $F(\beta)$ is the integrated fluence measured at angle β to the disk normal, the total x-ray fluence F is given by

$$F = \frac{\pi F(\beta)}{\cos \beta} \quad (1)$$

where $F(\beta)$ is in joules per steradian, and F is in joules. The assumption of a Lambertian radiator should be valid because, for high-power lasers like Nova, the beam spot size is much larger than the plasma-expansion distance. To minimize uncertainties, we kept the viewing angle near 60° with respect to the disk normal

when possible. At that angle, the integration is less sensitive to assumptions about angular distribution.

Results. Figure 4-5 shows x-ray conversion efficiencies demonstrated at Nova, Novette, and Argus for three wavelengths. The efficiency values are calculated with respect to incident-light intensity rather than absorption because absorption was not measured on Nova for $0.53\text{-}\mu\text{m}$ light. At the $0.53\text{-}\mu\text{m}$ wavelength, Nova and Novette results agree within the experimental error. Differences between these two sets of results, all within the error range, appear to be systematic: Nova data are systematically higher. These differences may be due to uncertainty in x-ray diode detector calibrations on Novette, which changed during the course of the Novette experiments.

The effect of gold M-band emission was not considered in the analysis of the Novette experiments. In an earlier study, the efficiency of gold M-band production was shown to increase at shorter wavelengths.⁹ At $1.06\text{ }\mu\text{m}$, M-band production was measured to be less than 0.2% and would not, at this level, change the interpretation of the x-ray conversion efficiencies measured at Shiva. At shorter wavelengths (0.35 and $0.53\text{ }\mu\text{m}$), M-band production was measured at more than 3% of the incident laser energy, according to the same study. The Nova results now lead us to conclude that the Novette x-ray conversion-efficiency data for $0.53\text{-}\mu\text{m}$ laser light at intensities around $2 \times 10^{14}\text{ W/cm}^2$ are about 10% too high because of M-band effects. At higher intensities, the effect of M-band emission on the spectrum may be even greater.

Fig. 4-5. Preliminary x-ray conversion efficiency results from Nova, plotted with results from Novette and Argus for comparison.

The x-ray conversion efficiency measured at Novette for 0.26- μm light is also shown in Fig. 4-5. Preliminary estimates indicate that these measurements are also too high because of gold M-band effects. The Novette data are being reanalyzed to assess the correction required for gold M-band emission.

Preliminary Nova results for x-ray conversion efficiency at a wavelength of 0.35 μm are also shown in Fig. 4-5. The figure includes data taken at incident angles less than 30° with respect to the disk normal. Larger angles of incidence were also measured. Some variation in efficiency with angle of incidence was observed, but the data trends have not been analyzed. The conversion efficiency in general is slightly higher at 0.35 μm than at 0.53 μm at most intensities, although the differences are within experimental error. If normalized to absorbed laser energy, the data at the two wavelengths would overlap because absorption is slightly higher at the shorter wavelength.¹¹ At intensities above 10^{15} W/cm^2 , the 0.53- μm conversion efficiencies appear higher than the 0.35- μm data, although the 0.53- μm Novette data were not corrected for gold M-band effects.

In Fig. 4-5, the 0.35- μm Nova results can also be compared with Argus results,¹ which came from experiments that used much less energy, a smaller spot size, and a Gaussian pulse of about 0.7 ns FWHM. Above 10^{14} W/cm^2 , the Nova data agree with those from Argus. Below 10^{14}

W/cm^2 , the Nova data continue on an upward trend, while the Argus data remain constant. This comparison is similar to that observed for Novette at the 0.53- μm wavelength: Novette and Argus data agree above 10^{14} W/cm^2 but diverge below 10^{14} W/cm^2 .

Summary: The Nova Dante systems were activated and reproduced previous Novette results. During the activation, gold M-band emission was found to produce unwanted signal in some Dante channels. This emission affects the interpretation of some Novette results. Past analysis of the Novette data is being reviewed. The Nova Dante systems have been improved to account for gold M-band effects. Preliminary measurements with 0.35- μm light indicate x-ray conversion efficiency to be slightly higher than that at 0.53 μm , but much of the difference could be due to increased absorption. No significant angular dependence was observed for angles up to 40° incident to the disk normal. In the next year, the effects of intensity and the angular dependence will be investigated more thoroughly.

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Exploding-Foil Experiments. In our 1985 experiments, we used 0.35- μm light from Nova to irradiate foil targets. This was the first use of "blue" light in exploding-foil studies. We irradiated targets that were thin enough to become completely underdense during the laser pulse. The resulting plasmas had large-scale-length density gradients, which made them ideal for producing and studying nonlinear laser-plasma interactions.¹² These studies were an extension of 1984 experiments on Novette.¹³ We also conducted a brief series of experiments with 0.53- μm laser light that confirmed the presence of forward Raman scattering.

Several improvements in the diagnostics were made since the Novette experiments. These include the addition of a second time-resolved optical spectrometer

and the use of square 1-ns laser pulses.¹⁴ Square pulses make interpretation of the experiments easier because the laser intensity is approximately constant throughout the pulse.

The principal diagnostics for the new experiments were two time-resolved optical spectrometers. In addition, the following diagnostics were used in all experiments:

- Numerous energy-balance modules (EBMs), which are spectrally filtered, time-integrated photodiodes designed to measure scattered light energies and angular distributions.

- An optical/x-ray (OX) streak camera with discrete channels for Raman, Brillouin or unabsorbed laser light, and $3\omega/2$ light.

- FFLEX, the filter-fluorescer diagnostic, which measures hard x-ray bremsstrahlung produced by suprathermal electrons.

An optical pyrometer was also used on a few shots to measure the timing of target burn-through.

Our first series of experiments was designed to survey a variety of conditions: laser intensity, angle of incidence, and wavelength. We used only one Nova beam, either at $0.35\text{ }\mu\text{m}$ and 1.5 kJ or at $0.53\text{ }\mu\text{m}$ and 2 kJ. During these experiments, we recorded several measurements that warrant further investigation:

- The $\omega/2$ scattered light during irradiation at a 50° angle of incidence. This geometry implies that the incident laser light was reflected from the plasma at a point near quarter-critical density.

- Suprathermal electron transport.

- Collisional damping with laser light at $0.35\text{ }\mu\text{m}$. In the only collisional-damping experiment conducted so far on Nova, a thin ($0.15\text{-}\mu\text{m}$) gold foil yielded strong SRS when irradiated at $4 \times 10^{15}\text{ W/cm}^2$. We plan additional experiments similar to those done previously at $0.26\text{ }\mu\text{m}$.¹²

We also observed strong SRS when $4\text{-}\mu\text{m}$ -thick plastic (CH) targets were irradiated with $0.35\text{-}\mu\text{m}$ light at an intensity of $2 \times 10^{15}\text{ W/cm}^2$. The SRS was measured to be on the order of 2% of the laser energy. Figure 4-6 shows the time-resolved SRS spectrum from one experiment. The maximum SRS wavelength clearly decreases with time, an observation

Fig. 4-6. Time-resolved SRS data shows the scattered-light wavelength decreasing with time as density decreases.

consistent with the decrease in maximum density as the foil burns through. This spectrum provides high-quality, absolutely time-resolved data on the maximum density present. (Time resolution is achieved through the use of an auxiliary fiducial timing pulse, not shown in Fig. 4-6.) In future experiments with silicon- or sulfur-doped foils, we will use this information to cross-calibrate density data obtained by time-resolved x-ray line spectroscopy. Such experiments should provide insight into some of the atomic physics problems encountered in reducing x-ray spectroscopy data.¹⁵

We also conducted a brief series of experiments with $0.53\text{-}\mu\text{m}$ light that were designed to look for forward Raman scattering. To avoid any possible ambiguity due to refractive effects, we looked concurrently in the forward direction for upshifted light (the laser frequency plus the plasma frequency) and the usual downshifted light (the laser frequency minus the plasma frequency). It has been predicted¹⁶ that an upshifted (anti-Stokes) component should be observed in forward Raman scattering because the plasma wave, which has a relatively small wave number, is near resonance for this process. Forward SRS was observed, both upshifted and downshifted. It was the first time both observations have been reported.

The amount of energy in the forward SRS was relatively low—two to three orders of magnitude lower than the amount of backscattered SRS. However, enough very hard x rays ($>100\text{ keV}$) were generated to be detected by the FFLEX diagnostic. The x-ray spectrum is shown in Fig. 4-7.

Fig. 4-7. Bremsstrahlung spectrum from forward SRS experiment shows a superhot (~ 130 keV) distribution of electrons.

These x rays indicate the presence of highly energetic electrons (characterized by a "temperature" of 130 keV). Such energetic electrons are expected to be produced by the high-phase-velocity plasma waves associated with forward SRS. Superhot electrons like these, in sufficient numbers, could pose a significant preheat problem for ICF targets. It is important to learn to control and limit their production.

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X-Ray Laser Experiments

Introduction. The primary tasks for x-ray laser experiments during 1985 were construction of the Nova two-beam laser facility and a complete upgrade of the x-ray laser diagnostics used previously at the Novette laser facility (see "Diagnostics Systems on the Two-Beam Chamber"). Nevertheless, we completed nearly 50 target shots by the end of 1985. Our primary goals in these experiments were to

- Fully understand the physics of the selenium exploding-foil amplifier and the laser kinetics.
- Achieve a saturated and well-characterized selenium x-ray amplifier.
- Demonstrate amplification at a wavelength as short as was feasible.

We focused our early efforts on activating diagnostics and following up on the previous Novette experiments.^{17,18} We primarily studied the selenium x-ray laser lines at 182.4, 206.3, 209.6, 220.2, and 262.3 Å. In "Plasma-Length Studies with Neon-Like Systems," we report on curves

for intensity vs length for these lines and the gain per cm we determined for each one. With present cylindrical lenses, we have been able to achieve gain $\times$ length products as great as ~ 10 to 12, although the values remain below saturation. We are presently installing new lenses that will allow us to reach improved gain $\times$ length products. We then should be able to determine both the onset of saturation (enabling us to quantify losses) and whether or not there is any difficulty with propagating an x-ray laser beam through an amplifier of this length.

We initiated experiments on other schemes to achieve soft x-ray amplification at shorter wavelengths. In addition to planned studies of higher-Z neon-like systems (yttrium and molybdenum), we have initiated experiments on a nickel-like gadolinium scheme that could produce amplification with lines extending from 100 to 60 Å. We describe our initial data obtained on this system in "Experiments on Nickel-Like X-Ray Lasing Systems."

We also performed experiments to unravel the complicated physics involved in the selenium x-ray amplifier and made progress in both theory and modeling of amplifier performance. We worked with KMS Fusion, Inc. to perform a set of measurements to determine the electron density of the exploding-foil amplifier under a variety of laser irradiance conditions and foil thicknesses. This work was done primarily to validate the accuracy of the LASNEX/XRASER modeling codes and to understand some peculiarities in performance noted in measurements completed last year. A summary of our results is presented in "Plasma Profile Experiments" and "Progress in the Analysis of Selenium X-Ray Laser Targets" (see Section 2).

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Plasma-Length Studies with Neon-Like Systems. Activation of the Nova two-beam target chamber has allowed us to continue with experiments on neon-like systems that were begun on Novette.¹⁹ Our studies of neon-like selenium ($Z = 34$) in

Fig. 4-8. Axial spectrum from a 2-cm-long selenium amplifier.

exploding-foil plasmas have been extended to plasmas with larger length due to the greater flexibility of the Nova two-beam optics. Novette had cylindrical correcting lenses producing a 12-mm-long line focus. The Nova two-beam optics has pairs of cylindrical lenses on each beam that can be rotated to give any length up to a 52-mm-long line focus. The line foci can be narrowed to a 100- μm width and the opposed beams superposed to better than 50 μm on the target.

Nova has been able to provide up to 2500 J per beam of 0.53- μm light in a 500-ps pulse, allowing irradiation of selenium-coated foils with lengths up to 5 cm. The foils, consisting of 750-Å selenium on 1500-Å Formvar ($\text{C}_{10}\text{H}_{10}\text{O}_3$), were irradiated at $I_L \approx 4 \times 10^{13} \text{ W cm}^{-2}$ in a 300- μm -wide line focus of variable length. Most data was taken at lengths less than 3 cm.

A typical spectrum for a 2-cm-long plasma taken from the microchannel plate incident grating spectrometer (McPIGS) spectrograph is shown in Fig. 4-8. The two $2p^33p$ ($J = 2$) $\rightarrow$ $2p^33s$ ($J = 1$) lines at 206.3 and 209.6 Å are overexposed so that

the weaker lines are visible. Of particular interest is the third $J = 2 \rightarrow 1$ transition at 262.9 Å, the $J = 1 \rightarrow 1$ transition at 220.3 Å, and the $J = 0 \rightarrow 1$ transition at 182.4 Å. The latter line was not conclusively identified on Novette due to underlying carbon emission and its observed small amplification. However, the small amount of gain on the line (on the order of 1 cm^{-1}) has made its identification easier for longer targets than was possible on Novette. As has been observed previously,¹⁹ the $J = 0 \rightarrow 1$ line does not exhibit the large gain (on the order of 10 cm^{-1}) predicted by earlier theory.²⁰

The lines at 206.3 and 209.6 Å, observed on Novette to exponentiate with lengths up to 1 cm, continue such behavior up to 3 cm. Figure 4-9 shows spectrally integrated intensity vs amplifier length for the 206.3-Å line. The 209.6-Å line shows essentially identical behavior and seems to possess the same intensity. The data shown in Fig. 4-9 are from the McPIGS spectrograph; more data are still being analyzed. Preliminary scrutiny of the data at 4 and 5 cm indicates that the amplifier may have begun

to saturate at about 1 cm. These data will be published after more detailed analyses and diagnostic calibration. Experiments are planned that will determine whether saturation is due to gain saturation or to an apparent saturation due to refraction in the plasma preventing x rays from propagating the full length of the target.

Figure 4-10 shows the time history of the 206.3-Å line taken from a transmission grating streaked spectrograph (TGSS) streaked spectrum. The time duration is on the order 200 ps for the 2-cm target. The 206.3-Å line reaches maximum

Fig. 4-9. Spectrally integrated intensity of the selenium $J = 2 \rightarrow 1$ transition at 206.3 Å vs amplifier length.

Fig. 4-10. Time history of the selenium 206.3-Å line.

intensity ~ 300 ps after the onset of continuum emission (usually assumed to coincide with the presence of the pump laser light), which may be consistent with emission during the interval of heating by the pump laser. As yet, we have no absolute timing fiducial on our streaked data, but an optical fiducial system will be incorporated in our experiments soon (see "Optical Fiducial for X-Ray Streak Cameras," later in this section).

During experiments on Novette, we observed that the equivalent neon-like system in yttrium ($Z = 39$) could be inverted using irradiances on the order of 10^{14} W/cm^2 and targets of 800-Å yttrium on 1000-Å Formvar. On Nova, we irradiated such foils of up to 3 cm in length using foci from 100 to 200 μm in width. We used the same 500-ps pulse length that was used for selenium. A typical axial spectrum from a 2-cm yttrium target is shown in Fig. 4-11. The main feature is the line at 154.8 Å that has been overexposed to emphasize the other lines. The predicted wavelength of 155 Å is one of the $J = 2 \rightarrow 1$ transitions of neon-like yttrium ($2p^5 3p \rightarrow 2p^5 3s$) and is equivalent to the 206.3-Å line in selenium.²¹

The wavelength of ~ 154.8 Å is also the predicted position of the $J = 0 \rightarrow 1$ transition, which is equivalent to the 182.4-Å

line in selenium. The 154.8-Å line is much more intense than either of the other two $J = 2 \rightarrow 1$ lines at 157 Å (equivalent to 209.6 Å in selenium) and 218 Å (equivalent to 262.9 Å in selenium). This outcome contrasts with the behavior in selenium for which the 206.3- and 209.6-Å lines are of equal intensity and gain. It is possible, therefore, that the line visible at 154.8 Å is the $J = 0 \rightarrow 1$ line or a combination of the $J = 0 \rightarrow 1$ and one of the $J = 2 \rightarrow 1$ lines.

Spectra taken using the full resolving power ($\lambda/\Delta\lambda \approx 1800$) of the McPIGS has failed to show the 155-Å line as a doublet. Thus, if a resonance occurs between the $J = 0 \rightarrow 1$ and $J = 2 \rightarrow 1$, then it is better than 0.1 Å. A limited set of data has been taken with yttrium; estimates of the gain on the 154.8-Å line are on the order of 4 cm^{-1} and are much less for the other lines. The time history measured on TGSS indicates a duration of about 200 ps for the 154.8-Å

Fig. 4-11. Axial spectrum from a 2-cm-long yttrium amplifier.

Fig. 4-12. Axial spectrum from a 2-cm-long molybdenum amplifier.

line. The duration and relative timing of the 157-Å line has not yet been measured.

Foils of 2000-Å molybdenum ($Z = 42$) coated on 1500-Å Formvar have been irradiated with $2 \times 10^{14} \text{ W} \cdot \text{cm}^{-2}$ at lengths up to 2 cm. Such foils produce gains in several neon-like transitions. Figure 4-12 shows a molybdenum spectrum taken from a 2-cm target. The lines at 131.0 and 132.7 Å are the $J = 2 \rightarrow 1$ transitions, equivalent to 206.3 and 209.6 Å in selenium. Like selenium, the two lines show the same intensity and gain.

The $J = 0 \rightarrow 1$ transitions at 106.4 and 141.6 Å have been detected, but have low gain. This outcome is similar to that for selenium, where the other $J = 0 \rightarrow 1$ transition (at 169 Å) has only now been detected (see Fig. 4-8), but with low gain. The $J = 1 \rightarrow 1$ transition, equivalent to 220.3 Å in selenium, may have been detected at 139.4 Å. The third $J = 2 \rightarrow 1$ transition, equivalent to 262.9 Å in selenium, is predicted to be at 196 Å and outside the spectral range of McPIGS as it was set up for these shots.

Despite the observed increase in required laser intensity with target Z , it may be possible to extend the study of neon-like schemes on the Nova 2-beam chamber in the near future to silver ($Z = 47$), which is calculated to have a $J = 2 \rightarrow 1$ transition at less than 100 Å (Ref. 21).

Further work must be done before accurate gain numbers can be obtained for the three elements we have studied. However, exponentiation has been observed in all three elements. The difference in relative intensities of the two main $J = 2 \rightarrow 1$ lines in yttrium is not yet explained in view of the equal intensities of the two lines in molybdenum and selenium.

Author: B. J. MacGowan

Experiments on Nickel-Like X-Ray Lasing Systems. We have performed initial measurements on a gadolinium x-ray

laser target to produce an x-ray laser with wavelengths shorter than 100 Å. Lasing is expected from the nickel-like gadolinium ions in the closed electronic shell, analogous to the lasing action observed for the closed-shell neon-like ions of selenium.²² For nickel-like gadolinium, lasing is predicted from 74.7 to 76.99 Å and from 94.66 to 101.38 Å (Refs. 23 and 24). The strongest predicted $4f \rightarrow 4d$ and $4d \rightarrow 4p$ transitions are shown in detail in Table 4-2.

Targets were made in the exploding-foil geometry, with 700 Å of gadolinium coated on one side of 1000 Å of Formvar. Because the method of fabricating gadolinium targets has not yet been perfected, the coating process produced wrinkling of the Formvar. The two targets used in the experiment were both 1 cm in length.

We studied the targets on two consecutive shots. The peak intensity for the first shot was 10^{14} W/cm^2 of 0.53-μm light in a Gaussian shaped pulse with a FWHM of 400 ps. The intensity was doubled for the second shot.

Figure 4-13 shows a spectrum obtained by the on-axis McPIGS. Lines can be identified as emission from copper-, nickel-, and zinc-like ions.^{25,26} All of the expected lasing lines lie within the region of spectral coverage. However, comparison with the off-axis McPIGS gives no indication of gain.

Figure 4-14 shows the two spectra observed by the streaked crystal spectrometer (SCS) at the peak of the incident laser pulse and at two different irradiation intensities. [Note that the y axes—the scales for relative intensities of observed spectra—are quite different for Figs. 4-13, 4-14(a), and 4-14(b).] Line emissions from iron-, cobalt-, nickel-, and copper-like ions are identified in Fig. 4-14. The structures on the higher-wavelength side of the identified lines for the ions in the nonclosed electronic shell (i.e., ions other than nickel-like ions) are probably due to satellite transitions from lower ionization stages.²⁷ The strength of these lines can be taken as an indication of the abundance of each ion averaged over the plasma. For the higher-intensity shot shown in Fig. 4-14, the relative intensities of nickel-like lines decrease, while the relative intensities for lines from more highly charged ions increase. This outcome suggests that

Table 4-2. Predominant calculated $4d \rightarrow 4p$ lasing transitions in nickel-like gadolinium.

Wavelength (Å)	Upper configuration	Lower configuration	$J_u \rightarrow J_l$
97.33	$3d_{5/2} - 4d_{5/2}$	$3d_{5/2} - 4p_{3/2}$	$2 \rightarrow 1$
101.38	$3d_{5/2} - 4d_{5/2}$	$3d_{5/2} - 4p_{3/2}$	$1 \rightarrow 1$
76.99	$3d_{3/2} - 4d_{3/2}$	$3d_{3/2} - 4p_{1/2}$	$1 \rightarrow 1$
74.71	$3d_{3/2} - 4d_{3/2}$	$3d_{3/2} - 4p_{1/2}$	$2 \rightarrow 1$
94.66	$3p_{3/2} - 4f_{7/2}$	$3p_{3/2} - 4d_{5/2}$	$2 \rightarrow 1$

Fig. 4-13. The 4 - 4 spectrum for on-axis x-ray emission in the region from 95 to 105 Å.

Fig. 4-14. The 4 - 3 spectra for iron-, cobalt-, nickel-, and copper-like gadolinium ions. Relative nickel-like emission is increased for the lower-energy shot shown in (b).

the doubled intensity of this shot, compared to the intensity of the first shot, was overly effective in producing significant ionization beyond the nickel-like stage.

In conclusion, preliminary data have been obtained for the nickel-like lasing system in gadolinium; no lasing was observed. We offer three possible reasons for the absence of lasing:

- Wrinkling of the target foil, which would produce a nonuniform plasma.
 - The fact that the target was only 1 cm long.
 - The possibility that, at even the lowest intensity, the incident flux was sufficiently high to overionize the plasma.
- The SCS data show that target length could be doubled to at least 2 cm while still leaving sufficient intensity to produce an abundance of nickel-like gadolinium ions.

**Authors: J. G. Woodworth
and M. S. Maxon**

Plasma Profile Experiments. Following the successful 1984 experiments on Novette and in an effort to optimize and understand x-ray laser output, we have examined an enlarged parameter range of target and irradiation conditions. Our collaborative experimental program was undertaken at KMS Fusion, Inc., Ann Arbor, Michigan.

The program addressed the following four scientific issues:

1. Experiments on one-sided vs two-sided irradiation. During the Novette experiments, we observed an unpredicted decrease in amplifier performance when irradiating from only one side with the pump laser. KMS Fusion researchers

Fig. 4-15. Shot matrix of parameters.

performed holographic interferometric measurements²⁸ to determine the temporal evolution of plasma electron density profiles for two-sided vs one-sided irradiation of selenium-coated Formvar targets. We also studied variations of the profile with laser pulse length and intensity and with thickness of the selenium layer.

2. Double-foil x-ray laser targets. To design even more suitable amplifier characteristics, we studied a double-foil x-ray laser geometry. The design of this target was undertaken to:

- Ensure the production of extremely flat electron-density profiles.
- Improve target-absorption characteristics.
- Increase the time duration available for amplification.

3. Measurements on laser light absorption by the exploding-foil target. To further our understanding of the amplifier heating process, we obtained a substantial number of measurements of laser light absorbed by the exploding-foil target as a function of three parameters:

- Target thickness.
- Laser pulse length.
- Laser intensity.

4. X-ray spectroscopic measurements. We obtained spatially resolved spectral

Fig. 4-16. Relationship of four holographic probe pulses relative to the main laser pulse.

(a) 300 ps

(b) 425 ps

(c) 550 ps

(d) 675 ps

Fig. 4-17. Typical interferograms from Type B target with two-sided irradiation at 1.6×10^{14} W/cm² and medium pulse length of 400 ps.

measurements of the selenium x-ray emission in the vicinity of the target. These measurements serve as a guide in evaluating target and plasma conditions and in determining the presence of certain spectral lines that were indicative of pumping by dielectronic recombination.²⁹

The targets used for our experiments can be categorized by thickness, as follows:

- Type A: 500 Å selenium on 500 Å Formvar.
- Type B: 750 Å selenium on 500 Å Formvar.
- Type C: 1500 Å selenium on 500 Å Formvar.
- Type D: 750 Å selenium on 1500 Å Formvar.
- Type E: 1500 Å selenium on 1500 Å Formvar.
- Double-foil type-D targets spaced 200 to 410 μm apart with selenium sides facing each other.

A summary of all the shots attempted is shown in Fig. 4-15; probe beam timing for interferometry is shown in Fig. 4-16. The method of performing interferometry and spectroscopy is discussed in Ref. 30.

Results. Almost all of the irradiated two-sided targets provided well-balanced and symmetric on-axis target density profiles. Figure 4-17 shows the results of four timeframe interferograms of a typical two-sided irradiation. Axial density profiles for single-sided vs double-sided irradiation have been obtained. In almost all cases, one-sided irradiation results in an asymmetric density profile; the effect is more pronounced as target thickness increases. In addition, the profile is steeper on the side opposite that irradiated by the pump laser. The peak density is shifted away from the region of the original plane of the target, an occurrence not predicted by theory. This observation could explain why we did not observe as large a gain coefficient when using "effectively" single-sided irradiation (horizontally displaced beams) to obtain 2.2-cm gain regions.²⁹ For similar irradiance conditions, the peak density decays more rapidly for a two-sided shot than for a one-sided shot, as shown in Fig. 4-18.

Two major benefits are predicted for a double-foil target. First, the longer time duration for the plasma at a given density is

Fig. 4-18. Peak density vs time for two-sided and one-sided irradiation.

Fig. 4-19. On-axis density profile for a double-foil target.

Fig. 4-20. Typical spatially resolved x-ray spectrum from the double-foil target.

an advantage for multipass cavity development. Second, the large density plateau minimizes possible refraction effects.

Interferograms obtained at 425 and 550 ps show a flat density plateau with a width of approximately 400 μm . Figure 4-19 shows the on-axis density profile at 375 ps for a double-foil target with 410- μm separation. This profile was obtained from an Abel inversion of the data about the axial midplane. The plasma is larger by a factor of 6 than that obtained with a single foil. In future work on double foils, we will vary laser irradiance conditions and foil separation. Such work should enable us to optimize the hydrodynamic performance of this target before attempting tests for amplification.

The total absorption of laser light by the targets appears to be fairly high. Irradiations of short pulse length generally yield lower absorption than do those of

longer pulses, as was predicted by theory. Values range from approximately 40 to 80% for the irradiation conditions that we studied.

The x-ray spectroscopic measurements attempted during these experiments were designed to study the variations in the relative proportions of resonant line radiation from neon-like, fluorine-like, and sodium-like selenium transitions as a function of laser and target parameters. Imaging spectrographs were used to provide 1-D pictures of spectrally resolved x-ray plumes normal to the target surface. Figure 4-20 shows spectral results from 7.8 to 8.5 $\text{\AA}$ for a double-foil target with 410- μm separation and irradiated by $1.1 \times 10^{14} \text{ W/cm}^2$. Analysis of these data should help determine ionization balance and electron temperatures in the spatial region of interest.

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Laser Plasma as an X-Ray Backlighting Source

Short, bright pulses of x-ray line emission are needed for flash radiography of laser-fusion targets. Using the Chroma laser at KMS Fusion, Inc., Ann Arbor, Michigan, we conducted x-ray radiography experiments with 0.53- μm light and nominal pulse lengths of 60 and 120 ps.³¹ We obtained M $\rightarrow$ L spectra for yttrium (2.00 to 2.35 keV), palladium (2.95 to 3.46 keV), and cesium (4.43 to 5.15 keV); and L $\rightarrow$ K spectra for chlorine (2.790 keV, He α $^1\text{P}_1$), titanium (4.749 keV), manganese (6.180 keV), and nickel (7.804 keV).

We used a time-resolved spectrograph to measure the duration of the x-ray line emission and an absolutely calibrated, time-integrated spectrograph to measure the absolute brightness. With the 60-ps laser pulse, the x-ray flashes from yttrium, palladium, and titanium were bright and shorter than, or comparable in duration to, the laser pulse length. For palladium, we measured the x-ray conversion efficiency into 4π sr to be $\xi_x = 2 \times 10^{13}$ photons/(J $\cdot$ sphere) over laser intensities I ranging from 6×10^{14} to $2 \times 10^{16} \text{ W/cm}^2$. For titanium, we measured $\xi_x = 3$ to 6×10^{11} photons/(J $\cdot$ sphere) at $I \approx 10^{16} \text{ W/cm}^2$.

Scaling arguments based upon an exact two-parameter transformation group of the hydrodynamic equations suggest that it might be possible, in light of these experimental results, to generate 10- to 20-ps x-ray flashes at energies as high as 6 keV with intense 0.26- μm laser light. The hydrodynamic equations upon which the transformation group is based assume flux-limited classical electron heat transport and an ideal-gas equation of state.

Experiments at the University of Rochester¹² have shown that with a 0.35- μm irradiation wavelength, some x-ray lines can be as much as an order of magnitude brighter than with 1.06- μm irradiation. Matthews et al.³³ also noted substantially higher conversion efficiency at shorter wavelengths for several lines.

Experimental Results. Table 4-3 summarizes our experiments. The targets were thick disks irradiated by a converging

Table 4-3. Summary of the KMS x-ray flash backlighter characterization experiments.

Materials and x-ray lines	For the experiment with maximum ξ_x			Range of the x-ray FWHM for all experiments	
	ξ_x [photons/(joule·sphere)]	Laser intensity (W/cm ²)	Laser FWHM (ps)	X-ray FWHM (ps)	(ps)
At 120-ps FWHM pulse length:					
Cl K (2.79 keV ¹ P ₁ Heα)	105×10^{11}	8.0×10^{14}	124	126	82-139
Pd L (2.95-3.46 keV NeI)	380×10^{11}	$0.6-1.6 \times 10^{15}$	108	157	64-157
Ti K (4.75 keV ¹ P ₁ Heα)	13×10^{11}	1.6×10^{15}	102	115	80-117
Cs L (4.43-5.15 keV NeI)	16×10^{11}	$2.3-4.6 \times 10^{15}$	115	110	83-110
Mn K (6.18 keV ¹ P ₁ Heα)	3.4×10^{11}	6.0×10^{14}	130	121	98-127
Ni K (7.80 keV ¹ P ₁ Heα)	0.85×10^{11}	$1.3-2.5 \times 10^{16}$	104	103	87-100
At 60-ps FWHM pulse length:					
Yt L (2.00-2.36 keV NeI)	240×10^{11}	1.6×10^{15}	~60	67	57-67
Pd L (2.95-3.46 keV NeI)	190×10^{11}	$0.9-1.6 \times 10^{15}$	80	60	51-66
Ti K (4.75 keV ¹ P ₁ Heα)	6.1×10^{11}	1.1×10^{16}	57	49	48-50
Mn K (6.18 keV ¹ P ₁ Heα)	1.6×10^{11}	$0.9-1.8 \times 10^{16}$	70	57	56-57

Table 4-4. Summary of the M → L emission-band experiments at 0.53 μm .

Element	Shot No.	FWHM laser pulse width (ps)	Laser energy (J)	FWHM spot diam from the x-ray pinhole camera (μm)	Focal offset (μm)	Laser intensity (W/cm ²)	Conversion efficiency, ξ_x [photons/(joule·sphere)]	X-ray FWHM (ps)
Yttrium	7291	~60	17.9	133 (sat.)	-200	$0.5-1.1 \times 10^{16}$	1.9×10^{13}	57
	7292	~60	12.2	126	-360	1.6×10^{15}	2.4×10^{13}	67
Palladium	7280	80	14.4	84 (sat.)	0	$0.9-1.6 \times 10^{16}$	191×10^{11}	66
	7271	79	13.0	64 (sat.)	0	$0.8-1.5 \times 10^{16}$	169×10^{11}	65
	7270	64	3.6	40	0	4.5×10^{15}	142×10^{11}	—
	7273	83	16.2	130 (sat.)	-200	$2.7-9.2 \times 10^{15}$	143×10^{11}	67
	7279	~70	1.6	35	0	2.4×10^{15}	70×10^{11}	—
	7281	~70	11.9	165	-360	8.0×10^{14}	93×10^{11}	59
	7282	63	21.1	284	-580	5.3×10^{14}	135×10^{11}	—
	7283	65	19.0	281	-580	4.7×10^{14}	230×10^{11a}	59
	7276	68	7.6	173	-360	4.8×10^{14}	134×10^{11}	51
	7278	65	17.4	328	-770	3.2×10^{14}	68×10^{11}	55
	7241	104	42.8	118 (sat.)	0	$2.0-3.8 \times 10^{16}$	156×10^{11}	118
	7232	115	37.5	105 (sat.)	0	$1.6-3.0 \times 10^{16}$	204×10^{11}	131
	7233	118	33.1	140 (sat.)	-200	$4.2-9.9 \times 10^{15}$	204×10^{11}	118
	7234	117	45.7	200 (sat.)	-360	$1.8-5.4 \times 10^{15}$	265×10^{11}	151
	7242	~120	36.2	168 (sat.)	-360	$1.4-4.2 \times 10^{15}$	232×10^{11}	140
	7236	108	39.6	247 (sat.)	-580	$0.6-1.6 \times 10^{15}$	378×10^{11}	157
	7237	114	34.6	287	-920	4.7×10^{14}	204×10^{11}	131
	7238	116	3.0	99	-360	3.4×10^{14}	110×10^{11}	68
	7239	116	8.5	266	-810	1.3×10^{14}	40×10^{11}	64
Cesium	7213	~120	26.1	97 (sat.)	0	$1.1-2.0 \times 10^{16}$	11×10^{11}	—
	7212	~120	24.1	124 (sat.)	-200	$3.0-7.1 \times 10^{15}$	—	90
	7218	115	37.9	190 (sat.)	-360	$2.3-4.6 \times 10^{15}$	16×10^{11}	110
	7223	91	32.9	269 (sat.)	-580	$0.6-1.6 \times 10^{15}$	12×10^{11}	83

^a Obtained from ratio of time-resolved data. Time-integrated data were not acquired.

Element	Shot No.	FWHM laser pulse width (ps)	Laser energy (J)	FWHM spot diam from the x-ray pinhole camera (μm)	Focal offset (μm)	Laser intensity (W/cm^2)	Conversion efficiency, ξ_x [photons/(joule $\cdot$ sphere)]	X-ray FWHM (ps)
Chlorine	7230	127	31.3	56	0	1.0×10^{16}	19×10^{11}	111
	7227	109	33.1	99	-200	3.9×10^{15}	34×10^{11}	139
	7224	125	35 ^a	130	-360	2×10^{15}	55×10^{11}	122
	7225	124	25.9	182	-580	8.0×10^{14}	105×10^{11}	126
	7229	113	4.6	102	-360	5×10^{14}	50×10^{11}	77
	7226	115	24.3	287	-920	3.3×10^{14}	77×10^{11}	98
	7228	134	25.2	500	-1450	1.0×10^{14}	—	110
	7231	125	24.9	500	-1450	1.0×10^{14}	27×10^{11}	82
Titanium	7286	62	19.0	58 (sat.)	0	$1.5\text{--}2.9 \times 10^{16}$	5.2×10^{11}	48
	7284	60	9.9	39	0	1.4×10^{16}	3.5×10^{11}	—
	7288	57	19.1	61	-360	1.1×10^{16}	6.1×10^{11}	49
	7287	59	19.0	76	-200	7.3×10^{15}	4.1×10^{11}	50
	7289	54	17.4	139	-612	2.1×10^{15}	small	—
	7217	117	34.8	64 (sat.)	-200	$> 0.3 \times 10^{15}$	4.2×10^{11}	105
	7210	~ 120	34.2	86 (sat.)	-200	$0.4\text{--}1.0 \times 10^{16}$	—	107
	7208	~ 120	12.1	48	0	5.6×10^{15}	8.5×10^{11}	—
	7211	144	29.3	89 (sat.)	-200	$3.1\text{--}7.2 \times 10^{15}$	—	117
	7214	113	24.1	80	-200	4.2×10^{15}	3.5×10^{11}	106
	7220	112	38.2	162	-580	1.7×10^{15}	12.2×10^{11}	112
	7219	118	37.3	159	-360	1.6×10^{15}	11.8×10^{11}	120
	7216	102	32.9	162	-360	1.6×10^{15}	13.0×10^{11}	115
	7222	117	24.2	267	-810	3.7×10^{14}	5.7×10^{11}	80
Manganese	7266	~ 70	12.3	37	0	1.6×10^{16}	1.0×10^{11}	57
	7269	74	13.9	55 (sat.)	0	$0.9\text{--}1.8 \times 10^{16}$	1.6×10^{11}	56
	7268	~ 70	10.3	37	0	1.4×10^{16}	1.2×10^{11}	—
	7267	~ 70	11.8	96	-200	2.3×10^{15}	0.8×10^{11}	—
	7243	122	34.5	63 (sat.)	0	$1.4\text{--}2.6 \times 10^{16}$	1.5×10^{11}	125
	7244	110	37.1	64 (sat.)	-200	$> 1.1 \times 10^{16}$	1.4×10^{11}	116
	7245	127	31.2	156 (sat.)	-360	$1.1\text{--}3.4 \times 10^{15}$	3.1×10^{11}	127
	7246	130	38.6	252	-580	6.0×10^{14}	3.4×10^{11}	121
Nickel	7247	119	31.2	281	-810	4.2×10^{14}	1.2×10^{11}	98
	7252	~ 550	35.9	123 (sat.)	— ^b	$0.6\text{--}6.0 \times 10^{15}$	2.5×10^{11}	350
	7248	102	24.7	60 (sat.)	0	$1.2\text{--}2.3 \times 10^{16}$	—	97
	7249	127	40.8	115 (sat.)	-200	$0.5\text{--}1.1 \times 10^{16}$	0.62×10^{11}	100
	7250	107	28.9	80 (sat.)	0	$1.3\text{--}2.5 \times 10^{16}$	0.84×10^{11}	87
	7253	104	28.2	82 (sat.)	0	$1.3\text{--}2.5 \times 10^{16}$	0.85×10^{11}	103
	7254	103	40.6	95 (sat.)	-200	$0.6\text{--}1.4 \times 10^{16}$	0.80×10^{11}	97

^a Laser energy was not measured because of computer failure; the value of 35 J is assumed because amplifiers were set to this value.

^b Data omitted because of conflicting records.

Table 4-5. Summary of the 1 — K emission-band experiments at $0.53 \mu\text{m}$.

Fig. 4-21. Conversion efficiency vs laser intensity for He-like and Li-like chlorine $2p \rightarrow 1s$ transitions. Neither Ko nor L α lines were summed over. Energy limits on the integration were typically 2.725 to 2.805 keV.

beam focused by an $f/1.8$ paraboloidal mirror. The target normal was at $\phi = 0^\circ$, $\theta = 10^\circ$. We used a time-integrated spectrograph at $\phi = 28.8^\circ$, $\theta = 62.0^\circ$, where θ is the polar angle from the paraboloidal mirror and ϕ is an azimuthal angle. We also used a time-integrated spectrograph with its line of sight between 62° and 72° from the target normal.

Tables 4-4 and 4-5 contain data obtained in individual experiments employing a variety of target materials. In these tables, the laser pulse widths under 100 ps are accurate to about ± 10 ps, and those over 100 ps are accurate to about $\pm 10\%$.

The focal-offset values in Tables 4-4 and 4-5 represent the distance the target was moved away from the best-focus position of the beam. The negative signs on the offsets indicate that the targets were in the converging part of the $f/1.8$ laser beam. For each focal offset, we computed the average spot size and standard deviation for the measurements yielding good x-ray pinhole images. These were $44 \pm 7 \mu\text{m}$ FWHM for a zero offset, $76 \pm 16 \mu\text{m}$ FWHM for a $-200\text{-}\mu\text{m}$ offset, $130 \pm 34 \mu\text{m}$ FWHM for a $-360\text{-}\mu\text{m}$ offset, $223 \pm 51 \mu\text{m}$ FWHM for -560- to $-612\text{-}\mu\text{m}$ offsets, and $292 \pm 26 \mu\text{m}$ FWHM for -770- to $-810\text{-}\mu\text{m}$ offsets. These values were used to calculate the laser-intensity ranges in Tables 4-4 and 4-5 whenever all the x-ray pinhole images for an experiment were exposed to a density exceeding 5.115 (the maximum our microdensitometer can handle).

None of the x-ray pulse widths in the tables have been corrected for the 45-ps resolution of the x-ray streak camera, which was determined by the sweep rate and slit width.

X-ray conversion efficiencies for the target elements chlorine, palladium, titanium, and manganese are plotted vs laser intensity in Figs. 4-21 to 4-24, respectively. All four plots at 120 ps show the conversion efficiency peaking at an intermediate intensity. This result may be due to higher absorption fractions at lower intensities. (The absorption fractions were not measured.) The longer scalelengths over which heat has to be transported at the higher intensities may also be part of the explanation.

We studied the M — L emission bands for elements with atomic number Z ranging from 39 (yttrium, 2.00 to 2.35 keV) to 68 (erbium, 7.1 to 8.25 keV). The palladium results are given in Fig. 4-22 and Table 4-4. Yttrium and cesium results are also given in Table 4-4. The total integrated brightness of the low-energy L lines is considerably greater than that of the K lines at nearly the same energies. At higher x-ray energies, only the K lines can be generated. At 120 ps, the cesium L lines were the highest-energy L lines that we observed. The neon-like L lines, which were strong for yttrium and palladium, were absent for cesium even at high intensity. We

Fig. 4-22. Conversion efficiency vs laser intensity for the palladium M — L lines. Energy limits on the integration were typically 2.87 to 3.48 keV.

Fig. 4-23. Conversion efficiency vs laser intensity for the titanium 2p — 1s lines for helium- and lithium-like ions. The energy limits on the integration were typically 4.66 to 4.78 keV.

looked for the erbium L lines, which we expected to be in the energy range 7.1 to 8.25 keV, but saw only the continuum emission.

At high x-ray energies, the lines from an L — K line emitter are brighter than those from an M — L line emitter. We can show why this is so if we neglect screening and

Fig. 4-24. Conversion efficiency vs laser intensity for the manganese 2p → 1s lines for helium- and lithium-like ions. Energy limits on the integral were typically 6.02 to 6.25 keV.

Fig. 4-25. Spectrum and time history of the 2p → 1s line emission for chlorine (shot 7225), titanium (shot 7248), and manganese (shot 7245).

use the simple nonrelativistic hydrogenic energy levels: $E_n = -Z^2 e^2 / (2a_0 n^2)$, where e is the proton charge, a_0 is the Bohr radius, and n is the number of the Rydberg shell. It then takes 4/5 of the M → L transition energy to remove the last M-shell electron, while it takes only 1/3 of the L → K transition energy to remove the last L-shell electron.

All the 2p → 1s lines we observed were primarily from helium- and lithium-like ions. The Kα line was always comparatively weak. Representative spectra from chlorine, titanium, and manganese are shown in Fig. 4-25, and those for M → L transitions from yttrium, palladium, and cesium are shown in Fig. 4-26. Neon I M → L line energies were identified using Mosely fits to experimental results from the Naval Research Laboratory.³⁴ The

2p — 1s transition energies were obtained from Scofield's calculations.³⁵ The K α line emission comes from cold material collisionally excited by suprathermal electrons. This material is either laterally outside the heated spot or below it. To produce the component of K α emission from below the heated spot, the suprathermal electrons do not have so far to diffuse as the pulse length is decreased; thus this component should increase in brightness at shorter pulse lengths. The K α line is indeed more intense at 60 ps for both titanium and manganese (see Table 4-6).

Diagnostics. The x-ray diffraction crystals used in the time-integrated spectrograph were calibrated at KMS. Table 4-7

gives their angle-integrated reflectivities R_c and also gives the calibrations for Kodak Direct Exposure Film (DEF) at the energies of the x-ray lines that we observed. We used a microdensitometer with a 10 $\times$ objective and matching input and output numerical apertures (N.A.) of 0.25 to obtain the film calibrations. Calibrations for the titanium, manganese, and nickel He α lines were obtained directly from the measurements of Rockett et al.,³⁶ and are in terms of the diffuse density. The diffuse density D_{diffuse} was first converted to specular density for 0.1 N.A. optics and then to specular density for 0.25 N.A. optics. We converted to $D_{0.1 \text{ N.A.}}$ by using a fit to Fig. 4-26 in Ref. 36. Conversion to $D_{0.25 \text{ N.A.}}$

Table 4-6. Chlorine, titanium and manganese line ratios.

Element	Shot No.	Intensity (W/cm ²)	Laser FWHM pulse width (ps)	K α $\frac{I_{K\alpha}}{I_{P_1 \text{ He}\alpha}}$	Li-like ^a $\frac{I_{K\alpha}}{I_{P_1 \text{ He}\alpha}}$	$\frac{I_{P_1 \text{ He}\alpha}}{I_{P_1 \text{ He}\alpha}}$	He doublet $\frac{I_{He}}{I_{P_1 \text{ He}\alpha}}$
Chlorine	7230	1.0×10^{16}	127	—	0.34	0.30	0.46
	7227	3.9×10^{15}	109	—	0.29	0.20	0.20
	7224	$\sim 1.7 \times 10^{15}$	125	0.005	0.34	0.19	0.10
	7225	8.0×10^{14}	124	0.008	0.39	0.42	0.07
	7229	5×10^{14}	113	—	0.34	0.30	0.00
	7226	3.3×10^{14}	115	0.014	0.57	0.42	0.00
	7231	1.0×10^{14}	125	—	0.58	—	0.00
Titanium	7217	$\sim 10^{16}$	117	0.02	0.61	0.55	0.015
	7208	5.6×10^{15}	120	0.02	0.60	0.54	0.00
	7214	4.2×10^{15}	113	0.05	0.82	0.66	0.00
	7222	3.7×10^{14}	117	0.04	1.37	0.69	0.00
	7286	$1.5-2.9 \times 10^{16}$	62	0.23	0.73	0.52	0.027
	7284	1.4×10^{16}	60	0.13	0.72	0.60	0.025
	7288	1.1×10^{16}	57	0.21	0.78	0.55	0.025
	7287	7.3×10^{15}	59	0.17	0.65	0.59	<0.02
Manganese	7243	$1.4-2.6 \times 10^{16}$	122	0.056	0.78	<0.63	—
	7244	$>1.1 \times 10^{16}$	110	0.00	0.60	<0.61	—
	7245	$1.1-3.4 \times 10^{15}$	127	0.00	1.42	—	—
	7246	6.0×10^{14}	130	0.00	2.01	—	—
	7247	4.2×10^{14}	119	0.00	3.40	—	—
	7266	1.6×10^{16}	~ 70	0.41	0.89	<0.70	—
	7269	$0.9-1.8 \times 10^{16}$	~ 74	0.41	0.83	<0.55	—
	7268	1.4×10^{16}	~ 70	0.68	1.10	<0.63	—
	7267	2.3×10^{15}	70	1.03	1.27	<0.65	—

^a Notation from A. H. Gabriel, *Mon. Not. R. Astron. Soc.* **160**, 99 (1972), Table V.

Spectrograph center energy (keV)	Crystal	R_c (radlans)	Assumed DEF film calibration (photons/cm ²)	
2.1	(002) FET	1.12×10^{-4}	4.9×10^7	$D_{0.1 \text{ N.A.}} \times (1 + 0.0131 D_{0.1 \text{ N.A.}}^3)$
2.8	(111) Si	5.4×10^{-5}	4.4×10^7	$D_{0.1 \text{ N.A.}} \times (1 + 0.0169 D_{0.1 \text{ N.A.}}^3)$
3.2	(111) Si	4.65×10^{-5}	4.67×10^7	$D_{0.1 \text{ N.A.}} \times (1 - 0.106 D_{0.1 \text{ N.A.}} + 0.031 D_{0.1 \text{ N.A.}}^2)$
4.75	(200) LiF	1.5×10^{-4}	3.37×10^7	$D_{\text{diffuse}}^{0.87} \times (1 + 0.11 D_{\text{diffuse}})$
6.2	(200) LiF	2.1×10^{-4}	4.03×10^7	$D_{\text{diffuse}}^{0.87} \times (1 + 0.084 D_{\text{diffuse}})$
7.8	(200) LiF	2.34×10^{-4}	4.84×10^7	$D_{\text{diffuse}}^{0.87} \times (1 + 0.07 D_{\text{diffuse}})$

Table 4-7. Crystal and film calibrations for the time-integrated spectrograph.

Fig. 4-26. Spectrum and time history of the M → L line emission for yttrium (shot 7292), palladium (shot 7236), and cesium (shot 7218).

was done on the basis of measurements made within the LEAD Program. Film calibrations for the yttrium M → L lines, the chlorine He α lines, and the palladium M → L lines were provided by the Naval Research Laboratory and assume 0.1 N.A. specular density. Again, we had to convert to 0.25 N.A. specular density values.

The conversions used were:

$$D_{\text{diffuse}} = (0.478 + 0.0236 D_{0.1 \text{ N.A.}}) \times D_{0.1 \text{ N.A.}} \quad (2)$$

and

$$D_{0.08 \text{ N.A.}} = (1.28 - 0.077 D_{0.25 \text{ N.A.}}) \times D_{0.025 \text{ N.A.}} \quad (3)$$

If we assume that $D_{0.08 \text{ N.A.}} \approx D_{0.1 \text{ N.A.}}$, then

$$D_{\text{diffuse}} = 0.61 D_{0.25 \text{ N.A.}} \quad (4)$$

All densities were measured from the fog level.

The (200) LiF and (111) Si x-ray diffraction crystals were calibrated at the Ti K α / β and C1 K α lines, respectively. The results were scaled to the appropriate energies by using the $R_c(\theta)$ curves of Gilfrich et al.,³⁷ for (200) LiF and the $R_c(\phi)$ curve of Burek³⁸ for (111) Si. Burek's double-crystal data were converted to single-crystal reflectivity after scaling by assuming a perfect crystal with zero absorption. The (002) PET crystal was not calibrated. The R_c value of 1.12×10^{-4} radians was taken from Gilfrich et al.³⁷

We used a Ken Tech streak camera, with a cesium iodide photocathode and a 0.75-mm-wide slit. The streak speed was measured to be 27 ps/mm by irradiating a brass wire with a train of five 100-ps (FWHM) laser pulses separated by 218 ± 4 ps. We also checked the timing simply by changing the trigger cable delay by

420 ps, and obtained a measurement of 28 ps/mm. The camera was operated at the No. 4 sweep setting for all these experiments. The magnification was measured to be 2.27 by taping a washer over the slit. Within several percent, this is the magnification needed to match the calculated spectral line positions at the photocathode. The slit width of 0.75 mm represents a time resolution of 45 ps. None of the FWHM x-ray pulse widths in Tables 4-3 to 4-6 have been corrected for this resolution. Many of the x-ray pulse widths for the 60-ps experiments are near the 45-ps time resolution and must actually be considerably shorter.

We also used an x-ray pinhole camera at $\phi = 0^\circ$, $\theta = 52^\circ$. It had four pin-holes $17 \mu\text{m}$ in diameter, which were filtered with 2.5 mils of beryllium. The target-to-pinhole distance was 5.5 cm. This camera took four images with $7.1\times$ magnification.

Physics of the X-Ray Line Generation. X-ray flashes that are both short and bright can be achieved only over a finite wavelength or energy range. At low x-ray-photon energy, one possible problem in producing a short x-ray flash is that highly ionized species might take too long to recombine. For instance, hydrogenic ions could recombine at a later time and produce $L \rightarrow K$ lines. At high energy, the problem is that there might not be sufficient time for the atoms to strip down to the required electron configuration.

Let us estimate the collisional ionization time τ_{ik} at the critical electron density n_c . We will sidestep the determination of the electron temperature. We assume that by choosing the right intensity, any electron temperature can be obtained. However, the maximum electron density at which the x-ray line generation can occur seems to be at most a few times the critical density for the laser frequency. Experimentally, it seems to be difficult to heat up much higher densities than this. With a fixed electron density n_c , there is an electron temperature T_e that minimizes τ_{ik} . Using Seaton's result⁴³ for the near-threshold collisional ionization rate, we find

$$\frac{1}{\tau_{ik}} = \left(\frac{256\sqrt{\pi}}{9} \right) \bar{g} a_0^2 \alpha n_c \times \left(\frac{E_H}{T} \right)^{3/2} \frac{1}{n} \left[\sqrt{\frac{T}{kT}} E_2 \left(\frac{I}{kT} \right) \right], \quad (5)$$

where g is the Gaunt factor for photoionization, which we take to be ~ 0.42 ; α is the dimensionless electromagnetic fine-structure constant of $1/137.036$; e is the speed of light; E_H is the hydrogen ionization energy of 13.61 eV; I is the ionization energy; T is the temperature; k is Boltzmann's constant, and the exponential integral $E_2(x)$ is given by

$$E_2(x) = \int_1^\infty \frac{e^{-xt}}{t^2} dt. \quad (6)$$

The factor in square brackets in Eq. (5) has the limits:

$$\left[\sqrt{\frac{T}{kT}} E_2 \left(\frac{I}{kT} \right) \right] kT \ll I$$

$$\sqrt{\frac{kT}{I}} \exp \left(-\frac{I}{kT} \right), \quad (7)$$

$$\left[\sqrt{\frac{T}{kT}} E_2 \left(\frac{I}{kT} \right) \right] kT \gg I \sqrt{\frac{T}{kT}}. \quad (8)$$

The factor in square brackets in Eq. (5) carries all the temperature dependence of τ_{ik} . It has a maximum value of 0.2541 at about $kT = 3.75I$.

For example, consider the 7.804-keV nickel ^{19}Ni He α line. The energy required to strip nickel from $2s\ 1s^2$ to $1s^2$ is 2398 eV. Thus,

$$(\tau_{ik})_{\min} = 141 \text{ ps} \left(\frac{10^{21} \text{ cm}^{-3}}{n_c} \right). \quad (9)$$

The n_c for 1.064- μm light is 0.985×10^{21} electrons/cm³. To completely strip the L-shell requires a time of about $\ln_{0.5} \times \tau_{ik}$, because the population of eight electrons in the L-shell decreases exponentially with τ_{ik} . At least 400 ps are actually required to strip the nickel to give $n_e = 10^{21} \text{ cm}^{-3}$. LASNEX simulations⁴⁰ indicate that at $n_e = n_c$ most of the electrons are in the cold population, which has a temperature of only about 2 keV, even at an irradiance of 10^{10} W/cm^2 . For $kT/I = 0.8$, the factor in square brackets in Eq. (5) is 0.116, 46% of its peak value. It is true that the low-density corona is very hot, but n_e is then much lower. Flash radiography at <100 ps and 1.064 μm , using the nickel He α lines, is therefore not expected to be possible.

It is easy to show that three-body recombination prevails over photo-recombination. The speed at which cooling and recombination take place determines how quickly x-ray emission will cease after laser irradiation stops. Our calculations indicate that cooling is effected primarily through the hydrodynamic expansion of the plasma rather than through heat transport into the dense plasma or by radiative losses.

The three-body recombination rate can be derived from detailed balance with collisional ionization. It can also be derived directly by considering classical electron-electron collisions in the Coulomb field of an ion. No quantum mechanics beyond the Bohr model is needed. The derivation is sketched in the next part of this article. It shows that the recombination rate into the n th Rydberg shell is proportional to n^4 . If n is not too high, then radiative decay to the ground state is rapid. For high n , the problem is much more complex. Low-momentum-transfer collisions will transfer electrons in high- n states preferentially to near-neighbor- n states. Mathematically, the transfer resembles diffusion, as discussed by Zel'dovich and Raizer.⁴¹

For simplicity, we will consider only those n states for which radiative decay is rapid. This constraint will give a lower limit to the three-body recombination rate. We have

$$\tau_{\text{RR}} = \frac{83 \text{ ps} \left(\frac{T}{1 \text{ keV}} \right)^2 \left(\frac{10^{22} \text{ cm}^{-3}}{n_e} \right)^2 Z^2}{\sum_{n=0}^{n_{\text{max}}} n^4} \quad (10)$$

We need to know what to use as n_{max} . Bethe and Salpeter⁴² give the summed oscillator strengths

$$F_{n'n} = \sum_{l,l',m,m'} f_{nlm}^{n'l'm'} \quad (11)$$

for the spontaneous transition rate of hydrogen from the Rydberg shell n to the lower n' . The spontaneous Einstein A_E

coefficient is then

$$(A_E)_{n'n} = \left(\frac{F_{n'n}}{n^2} \right) k^2 r_0 c \quad (12)$$

or

$$(A_E)_{n'n} = \frac{1}{n^2} \left[\frac{2^6}{3 \times 3\pi} \left(\frac{1}{n'^2} - \frac{1}{n^2} \right)^3 \right. \\ \left. \times \left(\frac{1}{n^3} \right) \left(\frac{1}{n'^3} \right) \right] k^2 r_0 c. \quad (13)$$

The rate $(A_E)_{n'n}$ is averaged over all possible l, m values for shell n and summed over all possible l', m' values for shell n' . It is an approximate result for nonrelativistic, hydrogenic ions. Here k is the wave number of the emitted photon, and $r_0 = 2.818 \times 10^{-13} \text{ cm}$ is the classical electron radius. We have divided the $F_{n'n}$ by n^2 to get an averaged rather than a summed oscillator strength for the upper shell n . The wave number k equals $[\bar{Z}^2 E_{II}/h c][(1/n')^2 - (1/n)^2]$, where h is Planck's constant divided by 2π , and $\bar{Z}$ is the effective charge of the ion, taking into account screening by the other electrons.

A measure of the rate of the radiative cascade to lower levels is given by:

$$\frac{d}{dt} \left\langle \frac{1}{n^2} \right\rangle = \sum_{n'=1}^{n-1} \left(\frac{1}{n'^2} - \frac{1}{n^2} \right) \\ \times (A_E)_{n'n} \approx \frac{1}{53 \text{ ps} \left(\frac{n^5}{\bar{Z}^4} \right)} \quad (14)$$

Let n_{max} be that value of n for which the righthand side of Eq. (14) equals τ_{RR}^{-1} . Substituting that value for n_{max} into Eq. (10), we obtain the three-body recombination rate as limited by spontaneous radiative decay:

$$\tau_{\text{RR}} \approx 150 \text{ ps} \left(\frac{T}{1 \text{ keV}} \right) \left(\frac{10^{22} \text{ cm}^{-3}}{n_e} \right) \frac{1}{\bar{Z}^2}. \quad (15)$$

The $2p \rightarrow 1s$ line emission can persist after the laser pulse ends only if the electron temperature stays high for a time or if hydrogenic ions that are created take a long time to recombine. The hydrodynamic

cooling time should be comparable to the laser pulse width for one-dimensional expansion or possibly even much shorter for three-dimensional blowoff.

The $2p \rightarrow 1s$ helium-like emission was observed to persist only for chlorine, which also had strong H α emission. If no hydrogenic ions are created, it does not matter how slowly recombination occurs, provided the electron temperature drops rapidly enough. Hydrogenic ions in the very underdense corona recombine on a long time scale, but their number is small.

Classical Calculation of the Three-Body Recombination Rate Into a Rydberg Shell. Our treatment of three-body recombination considers the collision of two free electrons in the Coulomb field of an ion. The electron-electron collision itself is treated completely classically. Our criterion for an electron to be captured into the n th Rydberg shell is that it have a resultant binding energy between $-\bar{Z}^2 E_H/2(n+1)^2$ and $-\bar{Z}^2 E_H/2n^2$.

Let us explain why this criterion is reasonable and what approximations we must make. We assume that the kinetic energies of the colliding electrons are much greater than the magnitude of the binding energy of the n th Rydberg shell and that n is much greater than 1. This assumption allows us to consider that the electron-electron collision occurs at a point, since it actually occurs mostly in a region much smaller than that of the Bohr orbit.

We compute the phase-space volume as a check on our criterion. The n th Rydberg shell accommodates $2n^2$ electrons and so should correspond to a phase-space volume of $2n^2 h^3$, where h is Planck's constant. Classically, if the angular momentum is zero, the maximum distance the electron can be from the nucleus is $r_{\max} = 2a_0 n^2 / \bar{Z}^2$. We approximate the phase-space volume as the product $V_x V_p$, where V_x is the volume of a sphere of radius $r_{\max}$ and V_p is the volume of a shell in momentum space corresponding to a maximum kinetic energy of $\bar{Z}^2 E_H/2n^2$ and a minimum kinetic energy of $\bar{Z}^2 E_H/2(n+1)^2$. The kinetic energy limits are obtained from the virial theorem, which states that the average kinetic energy equals the magnitude of the binding energy. $V_x V_p$ is then

calculated to be $(16.3\pi)n^2 h^3 = 1.7n^2 h^3$, which is close to $2n^2 h^3$.

Let $\mathbf{v}_1$ and $\mathbf{v}_2$ be the individual velocities of the colliding electrons. Then $\mathbf{V} = 1/2(\mathbf{v}_1 + \mathbf{v}_2)$ is the velocity of the center of mass, and $\mathbf{g} = \mathbf{v}_2 - \mathbf{v}_1$ is the velocity difference. Only the direction of $\mathbf{g}$ changes in the collision. After algebraic manipulation, the differential cross section is found to be:

$$\frac{d\sigma}{d\Delta E}(\mathbf{v}_1, \mathbf{v}_2) = \frac{2\pi e^4}{|\Delta E|^3} \times \left[\frac{1}{g^2} - \left(1 + \frac{2\Delta E}{m\mathbf{V} \cdot \mathbf{g}} \right) \frac{(v_2^2 - v_1^2)^2}{4g^4} \right], \quad (16)$$

where m is the mass of an electron, $\Delta E = 1/2m(v_2^2 - v_1^2)$, and $\mathbf{v}'_1$ and $\mathbf{v}'_2$ are the individual velocities after the collision.

The differential transition rate is

$$\frac{dw}{dE'_1} = \int \frac{d\Omega_1}{4\pi} \int \frac{d\Omega_2}{4\pi} \times \left[n_{e1} n_{e2} \frac{d\sigma}{d\Delta E}(\mathbf{v}_1, \mathbf{v}_2 | \mathbf{g}) \right], \quad (17)$$

Averaging is done over the directions of $\mathbf{v}_1$ and $\mathbf{v}_2$. We find

$$\frac{dw}{dE'_1} = 2\pi e^4 n_{e1} n_{e2} \times \left(\frac{1}{mv_2} \right) \left[\frac{E'_1 + \frac{E_1}{3}}{|E_1 - E'_1|^3} \right] \text{ for } v_1 < v_2 \\ \times \left(\frac{1}{mv_1} \right) \left[\frac{\frac{4}{3}E_2 + E_1 - E'_1}{|E_1 - E'_1|^3} \right] \text{ for } v_1 > v_2, \quad (18)$$

where E_1 is the energy of an electron with speed v_1 and E'_1 is the energy of an electron with speed v'_1 .

We now assume a distribution of electron energies $r_e(E)$. If the collision occurs

at distance r from an ion, this distribution is modified to become

$$n'_v(E) = \sqrt{\frac{E}{V(r)}} \times [n_v(E - V(r))]. \quad (19)$$

Let $E_b(n)$ be the magnitude of the binding energy of the n th Rydberg shell. We shall use $E_b(n)$ as a variable of integration with the limits $E_b(n)_{\max} = Z^2 E_H / n^2$ and $E_b(n)_{\min} = Z^2 E_H / (n+1)^2$.

The three-body recombination rate into the n th Rydberg shell is

$$\begin{aligned} \frac{1}{\tau_{3R}} = & \int_{E_b, \min}^{E_b, \max} dE_b \int_0^{\infty} 4\pi r^2 dr \\ & \times 2\pi e^4 \left\{ \int_{E_2=0}^{\infty} dE_2 \int_{E_1=0}^{E_2} dE_1 \right. \\ & \times \left[\left(n_v(E_1) n_v(E_2) \frac{1}{\sqrt{2mE_2}} \right) \right. \\ & \times \left(\frac{\frac{4}{3}V(r) - E_b + \frac{E_1}{3}}{[E_1 + E_b]^3} \right) \\ & \times \left(\sqrt{\frac{E_1 + V(r)}{E_2}} \right) \left(\sqrt{\frac{E_2 + V(r)}{E_2}} \right) \\ & + \left. \int_{E_2=0}^{\infty} dE_2 \int_{E_1=E_2}^{\infty} dE_1 \left[\left(n_v(E_1) n_v(E_2) \frac{1}{\sqrt{2mE_1}} \right) \right. \right. \\ & \times \left(\frac{\frac{4}{3}V(r) + E_b + \frac{4}{3}E_2 + E_1}{(E_1 + E_b)^3} \right) \\ & \times \left. \left. \left(\sqrt{\frac{E_1 + V(r)}{E_1}} \right) \left(\sqrt{\frac{E_2 + V(r)}{E_2}} \right) \right] \right\}, \end{aligned} \quad (20)$$

where $\varepsilon = E_1 - V(r)$ and $\varepsilon_2 = E_2 - V(r)$. Assuming that $n_v(\varepsilon)$ is a Maxwellian distribution with $kT \gg E_b(n)$, we find

$$\begin{aligned} \frac{1}{\tau_{3R}} \approx & 5.22 n_e^2 a_0^5 c \left(\frac{n^4}{Z^2} \right) \times \left(\frac{E_H}{kT} \right)^2 \\ & \times \left\{ \frac{1}{2} \int_0^{\infty} \frac{\exp \left[\left(\frac{-2E_b}{kT} \right) y \right]}{(y+1)^{3/2}} dy \right\}, \end{aligned} \quad (21)$$

where 5.22 is the result of evaluating the expression $2048\pi\alpha/9$, in which α is the fine-structure constant. This result is valid only if $E_b(n) \ll kT$.

Let us compare this result with that calculated by detailed balance with collisional ionization using the Thomson formula⁴³ for the collisional ionization cross section:

$$\begin{aligned} \frac{1}{\tau_{3R}} \approx & 192\pi^2 a_0^2 m_e^2 c \left(\frac{n^4}{Z^2} \right) \left(\frac{E_H}{kT} \right)^2 \\ & \times \left[\int_0^{\infty} \exp \left[\left(\frac{-E_b}{kT} \right) t \right] \frac{1}{(t+1)^2} dt \right]. \end{aligned} \quad (22)$$

If we use this method of calculation, f would equal $1/3$. The result of evaluating $192\pi^2\alpha/3$ would then be 4.61, which is close to the corresponding coefficient of 5.22 in Eq. (21).

Two-Parameter Transformation Group for the Hydrodynamic Equations. We can show, from very general arguments, why shorter-wavelength irradiance can produce shorter x-ray flashes at higher energies. We will derive scaling relations from the hydrodynamic equations, which include either flux-limited or classical heat conduction and the deposition of the laser's beam energy into the plasma. The conservation equations for mass, momentum, and energy in one dimension are:

$$\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x}(\rho u) = 0, \quad (23)$$

$$\frac{\partial}{\partial t}(\rho u) + \frac{\partial}{\partial x}(\rho u^2 + p) = 0, \quad (24)$$

$$\begin{aligned} & \frac{\partial}{\partial t} \left[\rho \left(\varepsilon + \frac{u^2}{2} \right) \right] \\ & + \frac{\partial}{\partial x} \left[\rho u \left(\varepsilon + \frac{u^2}{2} \right) + pu + J \right] \\ & = P(x, t), \end{aligned} \quad (25)$$

Here t and x are time and distance, ρ is density, u is the velocity of matter, p is pressure, ε is the internal energy per gram, J is the heat flux in ergs/(cm²·s), and $P(x, t)$ is the laser beam energy deposited per unit volume and per unit time.

We assume $P(x, t)$ has the form

$$P(x, t) = I(t) f(\rho)(x, t) \frac{\partial \rho(x, t)}{\partial x}. \quad (26)$$

The function f is normalized so that $\int f(\rho) d\rho = 1$. Thus $I(t)$ is the incident laser intensity, and $f(\rho)d\rho$ is the fraction of the laser energy that is deposited in the density interval $\rho - 1/2 d\rho$ to $\rho + 1/2 d\rho$. The heat flux J has one of the following forms:

$$J = \begin{cases} \pm f \sqrt{\frac{k T_e}{m}} k T_e n & \text{flux-limited} \\ \frac{2}{7} \lambda_e' \frac{\partial}{\partial x} T_e^{7/2} & \text{classical.} \end{cases} \quad (27)$$

Here, f is the flux-limiter, a dimensionless number usually taken to be between 0.02 and 0.1, and λ_e' is a constant that is independent of both the density and the temperature.

The internal energy per gram ε is assumed proportional to temperature. The pressure is assumed proportional to the product of the density and the temperature. The effects of radiation cooling and radiative heat transport are neglected.

We find three-parameter scaling groups for purely classical heat conduction and for purely flux-limited heat conduction. However, flux-limited heat conduction is highly unphysical. The heat flow cannot be flux-limited everywhere and particularly cannot be flux-limited in only one direction. The case of purely classical heat conduction is of interest only for short laser wavelengths and low intensities. Usually, the heat flow is flux-limited, but only in certain regions and at certain times. Fortunately, there is a two-parameter scaling group that applies whether the heat flow is classical, flux-limited, or partially both. Following are the scaling groups for each of these cases.

For flux-limited heat conduction, if $\rho(x, t)$, $u(x, t)$, $T(x, t)$, $I(t)$, and $f(\rho)$ constitute a solution set, then so do

$$\begin{aligned} \rho_A(x, t) &= A \rho(x, t) \\ u_A(x, t) &= u(x, t) \\ T_A(x, t) &= T(x, t) \end{aligned} \quad (28)$$

$$I_A(t) = A I(t)$$

$$f_A(\rho_A) = A^{-1} f(A^{-1} \rho_A)$$

and

$$\begin{aligned} \rho_B(x, t) &= \rho(x, \sqrt{B} t) \\ u_B(x, t) &= \sqrt{B} u(x, \sqrt{B} t) \\ T_B(x, t) &= B T(x, \sqrt{B} t) \end{aligned} \quad (29)$$

$$I_B(t) = B^{3/2} I(\sqrt{B} t)$$

$$f_B(\rho_B) = f(\rho_B)$$

and

$$\begin{aligned} \rho_C(x, t) &= \rho(Cx, Ct) \\ u_C(x, t) &= u(Cx, Ct) \\ T_C(x, t) &= T(Cx, Ct) \end{aligned} \quad (30)$$

$$I_C(t) = I(Ct)$$

$$f_C(\rho_C) = f(\rho_C).$$

For classical heat conduction, the alternate solutions are

$$\begin{aligned} \rho_A(x, t) &= A \rho(Ax, At) \\ u_A(x, t) &= u(Ax, At) \\ T_A(x, t) &= T(Ax, At) \end{aligned} \quad (31)$$

$$I_A(t) = A I(At)$$

$$f_A(\rho_A) = A^{-1} f(A^{-1} \rho_A)$$

and

$$\begin{aligned} \rho_B(x, t) &= B \rho(x, B^{1/3} t) \\ u_B(x, t) &= B^{1/3} u(x, B^{1/3} t) \\ T_B(x, t) &= \sqrt{B} T(x, B^{1/3} t) \end{aligned} \quad (32)$$

$$I_B(t) = B^{7/4} I(B^{1/3} t)$$

$$f_B(\rho_B) = B^{-1} f(B^{-1} \rho_B)$$

and

$$\begin{aligned}\rho_C(x, t) &= C \rho(x/C, t/\sqrt{C}) \\ u_C(x, t) &= \sqrt{C} u(x/C, t/\sqrt{C}) \\ T_C(x, t) &= C T(x/C, t/\sqrt{C}) \\ I_C(t) &= C^{5/2} I(t/\sqrt{C}) \\ f_C(\rho_C) &= C^{-1} f(C^{-1} \rho_C).\end{aligned}\quad (33)$$

For flux-limited heat transport, solution A [Eq. (28)] shows what happens if the laser wavelength is changed so that absorption takes place at A times the original density. Laser intensity must be multiplied by the same factor A to maintain the same electron-temperature profile. But with the change in wavelength, the density for each temperature is A times as great. We are heating A times as much matter at A times the density. Solution C [Eq. (30)] shows that if the pulse is made shorter, the length and time scales for the hydrodynamic variables shorten correspondingly. Except for this change of scale, the hydrodynamic variables have the same shapes and the same peak values.

For classical heat conduction, it is again solution A [Eq. (31)] that tells what happens if the laser wavelength is changed so that the absorption is at A times the original density. Both the length and time scales shrink for $A < 1$. With a laser pulse at $1/\sqrt{A}$ times the original wavelength, A times the original intensity, and $1/A$ times the original pulse width, we heat the same amount of matter as before, but at A times the density.

The two-parameter scaling group is given by

$$\begin{aligned}\rho_{(A, B)}(x, t) &= A \rho\left(\frac{x}{AB^2}, \frac{t}{\sqrt{AB^{3/2}}}\right) \\ u_{(A, B)}(x, t) &= \sqrt{AB} u\left(\frac{x}{AB^2}, \frac{t}{\sqrt{AB^{3/2}}}\right) \\ T_{(A, B)}(x, t) &= AB T\left(\frac{x}{AB^2}, \frac{t}{\sqrt{AB^{3/2}}}\right) \\ f_{(A, B)}(\rho_{(A, B)}) &= A^{-1} f(A^{-1} \rho_{(A, B)}) \\ I_{(A, B)}(t) &= A^{5/2} B^{3/2} I\left(\frac{t}{\sqrt{AB^{3/2}}}\right).\end{aligned}\quad (34)$$

We do not have the freedom we had in the previous cases to independently change the density at which absorption occurs, the laser intensity, and the laser pulse width. Only two of these three quantities may now be specified. We then obtain a new solution for the hydrodynamic variables.

This scaling predicts that if we observe a line conversion efficiency ξ_x at intensity I_1 , laser pulse length τ_1 , and laser wavelength λ_1 , then at wavelength $\lambda_2 < \lambda_1$, intensity $I_2 = I_1(\lambda_1^2/\lambda_2^2)$, and pulse length $\tau_2 = \tau_1(\lambda_2^2/\lambda_1^2)$, the x-ray conversion efficiency ξ_x should be unchanged. There are a number of important invariants in this scaling. The scaling is obtained by letting $A = \lambda_1^2/\lambda_2^2$ and $B = A^{-1}$. The distance and time scales are shortened by the factor λ_2^2/λ_1^2 , while the density scale is increased by the factor λ_1^2/λ_2^2 . The Lagrangian mass coordinate $m = \int \rho dx$ is uncharged. Fluid elements at the same Lagrangian mass coordinate are unchanged in temperature and velocity. Because $(\lambda)^2$ is not changed, resonance absorption produces the same distribution of suprathermal electrons. The suprathermal electrons deposit the same amount of energy per unit areal mass in the same fluid element of a given Lagrangian mass coordinate. Although the inverse bremsstrahlung absorption coefficient increases by a factor of λ_1^2/λ_2^2 because of the higher density, the amount and energy deposition of inverse bremsstrahlung are preserved since the distance scale is correspondingly shortened. Even though less time is available for stripping the ions, the increased density exactly compensates if we assume that electron collisions and not photoionizations are predominant. Other invariants are $\rho\tau$ and $\rho L/u$. (L is the length scale for density, and u is the velocity of matter.) L/u is a measure of the time the ion spends in the hot, dense region where line generation occurs before it is convected into the corona.

Several physical processes are missing from this model. First, L/λ is not invariant, but rather proportional to λ . At the shorter wavelength, fewer wavelengths of plasma are present for Brillouin scattering. Also, three-body recombination occurs more efficiently at the shorter wavelength because $\rho^2\tau$ is proportional to λ^{-2} . If, at the longer wavelength, the x-ray line emission persists because of incomplete recombination before the ions are convected into the

corona, the relative intensity of this tail will be lower at the shorter wavelength and the tail will dissipate more quickly. If heat flow is limited by ion turbulence, it is likely to be less so at the shorter wavelength and higher density, since v_{ei} is proportional to the electron plasma frequency ω_{pe} or to $\sqrt{\rho}$, whereas v_{ei} is proportional to the ion density or to ρ . Thus, at the shorter wavelength, more heat is conducted to densities above the critical density.

Therefore, under conditions yielding bright 50-ps x-ray flashes from 2ω irradiation, we would expect to achieve, using 4ω irradiation, x-ray flashes four times shorter and four times more intense but with the same time-integrated x-ray emission per unit area on the source. Of course, the irradiance at 4ω would have to be four times as great and last only one-fourth as long.

Conclusions. Bright x-ray flashes 50 ps or less in duration were demonstrated for the yttrium L lines (2.00 to 2.35 keV), palladium L lines (2.95 to 3.46 keV), and titanium K lines (4.75 keV 1P_1 He α). Based upon the two-parameter transformation group described in this article, we have reason to believe that bright 10- to 20-ps flashes at these same energies should be

possible with 4ω irradiation. X-ray flash radiography with point sources is simple and promising for moderate resolutions of about 10 to 30 μm . Higher resolutions of 1 to 2 μm probably will require a large-area backlighter and an x-ray microscope.

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X-Ray Absorption Spectroscopy

The ionization balance in laser-produced plasmas is normally determined by x-ray emission spectroscopy techniques that use Bragg crystals configured to give space or time resolution. Fractional ion abundances, densities, and temperatures are determined from best-fit treatments including line profiles⁴⁴ and line intensity ratios. Inherent in these methods is a prior knowledge of the relative, or preferably absolute, response of the spectrometer as

Fig. 4-27. Diagram of the arrangement of absorber and backlighter plasmas.

a function of photon energy. But this approach invariably leads to large uncertainties in the conclusions.

Self-absorption techniques can also provide information on plasma properties.⁴⁵ However, higher-quality, time-resolved data can be obtained by using an absorption technique allowing the backlighting source and the absorbed spectrum (recorded with the same instrumentation) to be compared at the same photon energy. The viability of such a technique has been demonstrated on the LLNL Janus laser system.

The geometry for this experiment is illustrated in Fig. 4-27. An aluminum slab target (250 μm thick and 2 mm wide) was irradiated by a 12-J, 1-ns pulse at 1.06 μm in a defocused spot 1.5 mm in diameter. The resulting plasma was probed by x rays emitted from a quasi-point backlighting source. The x rays were generated by irradiating a 20- μm -diam aluminum wire with a 15-J, 0.8-ns pulse at 0.53 μm in a spot 100 μm in diameter at the position of the point backlighter. The peak of the backlighting x-ray line emission (hydrogen- and helium-like lines) was delayed 0.5 ns with respect to the peak of the main plasma line emission and verified by simultaneously recorded x-ray streaks from each plasma.

Two x-ray pinhole cameras gave data on the separation of the plasmas (~ 4 mm), the spatial extent of the quasi-point source in the flow direction of the main target

plasma (~ 10 μm), and the uniformity of the irradiation distribution on the main slab target. Two miniature PET crystal spectrometers⁴⁶ were used to view the backlight emission with crystal-limited spectral resolution. One was placed behind the main target and the other in front, with a line of sight through the expanding coronal plasma. The absorbing plasma was projected by means of a Bragg crystal onto direct-exposure x-ray film at a magnification of 20.

Figure 4-28 shows the higher-order helium-like lines of aluminum and the hydrogen-like Ly- β (1s $\rightarrow$ 3p) lines for Shot PO829536. Close to the target surface, line reversal is evident on the 1S $\rightarrow$ 3¹P, 1S $\rightarrow$ 4¹P, and 1S $\rightarrow$ 5¹P lines, but not on the Lyman lines. The 1S $\rightarrow$ 2¹P line was not in range. On this shot, there was no self-emission visible from the main aluminum plasma because the irradiance was too low (3×10^{11} W/cm²). The line reversal occurred because the low-density, cold, absorbing plasma has a spectral line profile that is relatively narrow compared to the backlighting profile. (In absorption, the large source size of the main plasma does not cause broadening.) Further away from the plasma, no reversal is seen, presumably because the density of absorbing atoms drops too low.

Self-emission was evident on shots at slightly higher irradiance. Figure 4-29 shows a shot similar to that in Fig. 4-28 except that the irradiance was 3×10^{12} W/cm².

Fig. 4-28. The emission from the 12- μm , aluminum point backlight viewed through the main plasma at 10^{12} W/cm² at 1.06 μm . There is no emission visible from the main plasma. Positions along the spectral lines correspond to viewing distances from the target surface.

Fig. 4-29. Self-emission from the main plasma at a higher irradiance.

The lineout shows both self-emission, which is wide in spectral direction because of source broadening, and the much narrower point backlighting line superimposed on it.

To analyze these data, the film density was converted to exposure using a calibration by Rockett.⁴⁶ Spectral lineouts well away from the plasma were assumed to be unaffected by the plasma and were used to define the backlighting spectrum $I_0(\lambda)$. These profiles were then used in a simple radiation-transport equation to predict the line shape $I_l(\lambda)$ that would be measured after transmission through a homogeneous slab plasma:

$$I_l(\lambda) = \int_{-\infty}^{\infty} I_0(\lambda') F(\lambda' - \lambda) \times \exp - \tau(\lambda') d\lambda'. \quad (35)$$

Here

$$\tau(\lambda) = n_1 l \left(\frac{\pi e^2}{2\epsilon_0 mc} \right) y(\lambda) + n_e l \sigma_c \quad (36)$$

is the opacity, n_1 is the ground-state ion density, l is the absorption path length, $y(\lambda)$ is the line-shape function, σ_c is the slowly varying continuum absorption, and f is the normalized instrument function. Stark-broadened line profiles for the aluminum atoms were calculated and then

Doppler-broadened to give $y(\lambda)$.⁴⁷ The instrument broadening was due to the source size of the backlighting pin and the width of the Bragg diffraction peak. The instrument profile, taken from the apparent spectral width of the dielectronic satellites to the Al XIII Ly- α line, was 3.5 mÅ (FWHM). The other lines are spectrally broadened by the high electron density of the backlighting plasma.

From Eq. (35), $I_l(\lambda)$ depends on $n_1 l$ and implicitly on the electron density n_e and the temperature as they affect the line profile. The electron temperature of the main plasma can be obtained from the slope of the recombination continuum. On shot P422517 (Fig. 4-29), it is 120 eV. On shot PO829536, the plasma temperature was lower because there was no detectable self-emission. A value of 100 eV was used for the electron temperature, giving an insignificant Doppler broadening of 0.6 mÅ. Consequently, the absorption depends on two parameters, $n_1 l$ and n_e .

Lines of best fit in $n_1 l$ and n_e parameter space are shown in Fig. 4-30. An additional constraint for aluminum is that

$$n_e > (Z - 2) \frac{N_h}{f_i}, \quad (37)$$

where Z is the atomic number, f_i is the fraction of atoms in the helium-like state

Fig. 4-30. Lines of best fit for Al XII: 1-3 (beta), 1-4 (gamma), and 1-5 (delta).

of ionization, and N_a is the atom density. This function is also plotted in Fig. 4-30 for $f_i = 1$ ($l = 1$ mm). As shown in the figure, the fits indicate an electron density of 10^{21} cm^{-3} . The value of $n_i l$ for which the lines are consistent implies $f_i \sim 0.3$ from Eq. (37).

To obtain a good fit for the Al VIII $1S \rightarrow 5^1P$ transition, it was necessary to use the continuum absorption cross section σ_c appropriate to absorption above the edge. This cross section was needed because the spectral position of the continuum-lowered K edge for the lithium-like state of ionization is below the Al XIII $1S \rightarrow 5^1P$ line, but above the Al XIII $1S \rightarrow 4^1P$ line.

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Laser-Driven Impulse Experiments

Under the Strategic Defense Initiative (SDI), the effectiveness of high-energy, pulsed laser beams as defensive weapons is being evaluated. The mechanical impulse that can be delivered to a target determines the lethality of such beams against threat targets. Experiments were performed in Nova's two-beam target chamber to measure mechanical-

impulse coupling coefficients for precisely defined laser-target interactions.

In addition, two instrumentation packages were evaluated in the two-beam chamber to determine their feasibility for use in underground tests at the Nevada Test Site (NTS). The packages consisted of:

1. Piezoresistive pressure sensors to measure pressure induced in targets during laser-target interactions.
2. Linear-velocity transducers (LVTs) to measure mechanical impulses delivered to targets during shots.

This article describes the SDI laser-target interaction experiments and the NTS instrumentation feasibility studies discussed above. The experiments were funded by the Weapons Physics Laser Project based in LLNL's Defense Programs Military Applications Division.

SDI Experiments. The SDI experiments measured mechanical-impulse coupling coefficients between aluminum targets and 2ω , 1.0-ns Gaussian beams at fluences up to 18.6 kJ/cm^2 . The targets were $0.020 \times 2.0 \times 2.0$ -in. foils of Al 6061 attached to a compound pendulum. Fluences between 89.0 J/cm^2 and 18.6 kJ/cm^2 were delivered to the targets via Nova's Beam 7 by varying the spot size of the beam and the amplification available from the front end of the laser.

With the target chamber open to the atmosphere, a 2ω alignment beam was used to position the cylindrical focusing lenses for Beam 7 until the required spot size was obtained on the target. The target was positioned on the diverging side of best focus. The chamber was then pumped down to a vacuum of approximately 10^{-5} Torr.

One target could be shot at a time, and a total of 11 shots were performed. The interaction of the laser with the target imparted momentum to the target and, therefore, to the supporting pendulum, causing the pendulum to swing. The deflection of the target was recorded on videotape by a closed-circuit television camera mounted on the outside of the target chamber. The mechanical impulse was calculated from the maximum deflection of the pendulum. The impulse is generated by a combination of light reflection, vapor blowoff, and plasma pressure.

Table 4-8. Nova SDI test series. Mechanical-impulse coupling coefficients (C) for 1.0-ns Gaussian beam interaction with Al 6061 at 2ω in vacuum.

Shot No.	Laser energy (J)	Spot size (cm ²)	Fluence (J/cm ²)	E (dyn·s)	Momentum (dyn·cm ²)	Impulse, I (dyn·cm ²)	Pendulum sensitivity (deg/dyn·s)	Change in target mass (mg)	C (I/E)	I/E^3
2B2 ω -1	150 $\pm$ 50 ^b	1.68	89.3	335	199	2.911×10^{-3}	—	—	2.23	4.48
2B2 ω -2	238	1.68	151.7	404	241	2.911×10^{-3}	0.30	1.70	3.84	—
2B2 ω -3	393 ^b	1.68	234.0	531	316	2.911×10^{-3}	0.37	1.35	3.25	—
2B2 ω -4	2060	1.68	1226.0	2246	1337	2.911×10^{-3}	1.20	1.09	1.87	—
2B2 ω -5	2378	1.68	1416.0	2495	1485	2.911×10^{-3}	1.55	1.05	1.78	—
2B2 ω -6	2237	0.30	7457.0	2031	6770	2.16×10^{-3}	11.40	0.91	1.02	—
2B2 ω -7	1015	0.30	3383.0	920	3067	2.16×10^{-3}	3.48	0.91	1.34	—
2B2 ω -8	385	0.30	1283.0	415	1383	2.91×10^{-3}	—	1.08	1.84	—
2B2 ω -9	1875 ^b	0.16	1.17×10^4	1810	1.13×10^4	2.91×10^{-3}	19.80 ^b	0.97	0.94	—
2B2 ω -10	2047	0.16	1.28×10^4	1975	1.23×10^4	1.99×10^{-3}	16.50	0.97	0.92	—
2B2 ω -11	2972	0.16	1.86×10^4	2758	1.72×10^4	1.99×10^{-3}	18.90	0.93	0.82	—

^aCoupling estimates used by the lethality task group of the Defense Nuclear Agency.

^bCalorimeter failure—energy inferred from 1ω injection.

The mechanical impulse coupling coefficient C was calculated from

$$C = I/E, \quad (38)$$

where I is the impulse delivered to the target and E is the energy delivered per pulse.

Mechanical-impulse coupling coefficients from 0.91 to 2.23 dyn·s/J were measured. This test series verified predictions of impulse coupling to aluminum for unambiguously one-dimensional plasma expansion at fluences or fluxes relevant to SDI. A summary of these results is provided in Table 4-8.

Simulated NTS Experiments. The two-beam target chamber was also used to simulate NTS underground test conditions. Targets for these experiments were aluminum and contained two piezoresistive pressure gauges made of Manganin alloy. The front gauge was embedded 0.5 mm behind the front surface of the target, and the rear gauge was embedded 1.0 mm behind the front gauge. The target was positioned on the diverging side of best focus in the chamber. Fluences from 2.7 to 10.1 kJ/cm² at 2ω were delivered to the target by the same techniques of adjusting energy and spot size as in the SDI experiments.

At a fluence of 2.7 kJ/cm², the front and rear gauges measured 39.0 and 18.0 kbar, respectively. At a fluence of 10.1 kJ/cm², the front pressure gauge was destroyed and the rear gauge measured 68.0 kbar.

LVTs were used to measure the mechanical impulses resulting from selected laser-target interactions. An LVT consists of an aluminum rod that initially rests against the front surface of a gauge. When contacted by the laser beam, the rod is driven by mechanical impulse a known distance to a back stop. As the rod moves, it moves a connecting rod in a magnetic field. The rod motion was calibrated so that the induced voltage corresponded to mechanical impulse. Fluences from 2.6 to 61.0 kJ/cm² produced impulses from 4.1×10^3 to 1.3×10^5 dyn·s/cm².

In summary, the two-beam chamber test series achieved the following:

1. Established the feasibility of using Nova for laser-effects testing with beam parameters relevant to SDI.

2. Demonstrated Nova's utility for measuring key target responses to simulated underground-test sources.

Similar tests will continue in 1986. Future tests will utilize more sophisticated instrumentation so that exact physics conditions can be more precisely determined.

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Diagnostics

Introduction

In 1985, LEAD Program resources were largely dedicated to installing and activating diagnostic systems on the ten-beam

and two-beam target chambers of Nova. These diagnostics enable us to measure photon energies from 1.2 eV to 500 keV with respect to time, frequency, and space. They allow us to conduct a wide range of experiments dealing with laser-plasma interaction, hydrodynamics, implosion, radiation physics, and XUV lasers. The next two articles in this section describe our present diagnostic systems.

We are committed to developing advanced diagnostics and improving our present systems while supporting the technology base required to fulfill this commitment. Our efforts in the past year focused on three major areas:

- X-ray diagnostics.
- Time-resolved measurements.
- Neutron diagnostics.

X-Ray Diagnostics. The articles "McPIGS Spectrograph Improvements" and "The Transmission-Grating Streak Spectrometer" describe improvements in the diagnostics fielded in our XUV laser experiments. For the McPIGS, a microchannel-plate grazing-incidence spectrometer, these improvements include increased sensitivity and imaging capabilities. For the TGSS, a transmission-grating streak spectrometer, reliability has been improved, and we are developing an absolute energy-calibration technique that employs well characterized, continuous x-ray sources.

We are also expanding our capability to apply x-ray optics to XUV laser cavities by using synthetic multilayer mirrors. We fabricated and calibrated mirrors that can potentially serve for cavity feedback and as output couplers in the 100- to 200-Å spectral range. These efforts are described in "X-Ray Laser Cavities."

Our experiments require a fast, compact, flat-response calorimeter/detector for x rays in the 0.1- to 2.0-keV range. The article "Solid-State X-Ray Detectors" describes our efforts to develop such an instrument. We are exploring the potential of several solid-state devices such as photoconductive and pyroelectric detectors for this application.

To develop advanced x-ray diagnostics, we need a technology base that includes radiation sources, detectors, and fabrication capabilities. Our progress in building this technology includes the development

of two new instruments: a simple, bright x-ray diode and a sensitive, direct x-ray recording camera for $h\nu > 500$ eV that employs a charge-coupled device (CCD). These instruments are described in the articles "The X-Ray Diode Source" and "Development of X-Ray Charge-Coupled Device Camera." We also established a state-of-the-art microfabrication laboratory that now uses optical lithographic techniques and will use x-ray lithographic techniques to produce a variety of diagnostic-related structures. Fabrication of these structures typically demands accuracy to micron, and often submicron, tolerances. The article "Microfabrication Laboratory" outlines the capabilities of this facility.

Time-Resolved Measurements. Present and future Nova experiments will depend on our ability to field quantitative, time-resolved diagnostics. Many experiments require two-dimensional recording instruments such as framing cameras with high spatial resolution, high sensitivity, sufficient contrast, and framing times from 50 to several hundred picoseconds. A critical aspect in the development of such instrumentation is the ability to switch high-voltage pulses rapidly. To achieve switching with the requisite speed, we are investigating two kinds of technology. In the article "The X-Ray Framing Camera," we describe progress in developing laser-driven Auston switches and characterizing voltage-pulse propagation in transmission lines. The framing camera design based on this technology requires rapid switching of 5 kV. We are also evaluating the use of overdriven silicon diodes to gate microchannel plates. Lower voltages (≤ 1 kV) are required for this design, which is described in "The Gated Microchannel-Plate X-Ray Pinhole-Camera Detector."

As fast, compact x-ray detectors are developed, it becomes increasingly important to evaluate appropriate transmission and recording media. Therefore, we are developing directional couplers in which laser light is modulated between two closely spaced waveguides by an applied electric field (such as that produced by a fast detector). The laser light is transmitted by single-mode fiber optics to an optical streak camera. The combined response of

the transmission and recording medium can exceed 20 GHz. The article "Directional Couplers" discusses our progress in this area.

For many experiments, well characterized x-ray streak cameras are imperative. For example, quantitative hydrodynamic experiments require a detailed knowledge of the contrast transfer function of the camera. As explained in the article "Spatial and Temporal Resolution of LLNL X-Ray Streak Cameras," we thoroughly characterized our cameras and identified their performance limitations.

Because increasing numbers of x-ray and optical streak cameras are employed in experiments, a means of cross-timing these diagnostics is essential. To achieve cross-timing, we developed a fiber-optic-based distribution system that employs a laser synchronous with Nova's main oscillators. This laser delivers harmonically converted $0.53\text{-}\mu\text{m}$ light to all our streak cameras. The light directly illuminates the photocathode of the optical cameras. For the x-ray cameras, a $140\text{-}\text{\AA}$ aluminum coating on the camera end of the fiber bundle serves as the photocathode. Photoelectrons are produced by multiphoton processes in the aluminum. This fiber-optic system, which is described in "Optical Fiducial for X-Ray and Optical Streak Cameras," will allow for absolute correlation of streak cameras to better than 100 ps.

We also improved our recording of streak-camera output by employing CCD readout systems. A prototype unit, consisting of a commercial CCD camera lens coupled to the output phosphor of an optical streak camera, has demonstrated excellent linearity, sensitivity, and spatial resolution. This system, described in "CCD Readout for Optical Streak Cameras," is being installed on all of the optical streak cameras that measure Nova performance.

Neutron Diagnostics. We investigated the development of several advanced neutron diagnostic instruments and techniques for measuring fuel areal density, fuel spatial distribution, and fusion burn (confinement) time. When implemented, these diagnostics will prove invaluable in compression experiments.

The article "Neutron Spectroscopy" describes a neutron time-of-flight

spectrometer consisting of many channels that detect single events. Variations of this instrument will measure the fuel ion temperature and the areal density.

We thoroughly analyzed a thick Fresnel zone-plate camera proposed for sensitive neutron imaging (see "Neutron Imaging with Zone Plates.") Our analysis shows that this coded aperture can serve as a camera with enough sensitivity to detect $<10^{11}$ neutrons, high resolution ($10\text{ }\mu\text{m}$), and a sufficiently large field of view.

For the measurement of fuel areal density, we also investigated neutron-activation techniques that employ a krypton tracer in the fuel. Experiments conducted at the University of Rochester indicate that the neutron-activated krypton can be collected with materials that have been used in earlier implosion experiments. These experiments are described in the article "Krypton Activation Method for Measuring Fuel $\langle pr \rangle$." If successful, this extremely sensitive technique will allow us to measure fuel-yield-area-density products as low as 10^9 neutrons $\cdot\text{g}/\text{cm}^2$.

Author: E. M. Campbell

Activation of the Nova Ten-Beam Target Chamber for Target Experiments

Introduction. During 1985 we prepared Nova for ICF and weapons physics experiments that will use the ten-beam target chamber. In addition to preparing this chamber, we undertook a variety of experiments to study the physics of large lasers and the performance of Nova. We also began the construction project for the two-beam target chamber and subsequent experiments.

To pursue all the projects, we needed many shots from the Nova laser. Figure 4-31 shows the cumulative system shots delivered by Nova beginning at the end of the construction project. The performance of Shiva and Novette is also shown. A system shot is any shot in which all amplifiers are fired on at least one laser chain. Nova has consistently delivered 40 to 50 shots per month, which is more than either Shiva or Novette were able to

Fig. 4-31. Cumulative system shots vs months after the end of construction for Shiva, Novette, and Nova.

produce so soon after construction. The success rate of these shots has been 85 to 90%, which is comparable to the rates obtained by Shiva and Novette.

Chronological Summary. To prepare a laser facility for complex target experiments, we must fire many simple target shots to test, calibrate, and debug target diagnostics. There is no other source of signals that is bright enough and no other way to test system integration. In addition, we must fire shots to test target alignment and other operating procedures. Figure 4-32 shows the target shots fired by Shiva, Novette, and Nova as they came into operation after construction. Because of our multiple priorities, we did not fire as many target shots on Nova during this interval as we did on Novette during the first few months after construction. However, by late 1985 we produced more total target shots per month on Nova than Novette during a comparable interval. The Nova ten-beam chamber (one of three major experimental areas) yielded target shots at 80% of the rate obtained by the only experimental area in Novette. Both Nova and Novette were much more productive than Shiva.

The activation of the Nova ten-beam chamber consisted of three phases with distinctly different target-shot rates, as can

be seen in Fig. 4-32. The first phase, from January through April 1985, included few target shots. During this interval, we prepared the ten-beam target chamber for diagnostics activation. During 25% of the shifts worked, our top-priority effort was the completion and activation of the chamber vacuum and the vacuum controls system. By May 1985, we were able to routinely operate the chamber and diagnostic vacuum systems for target experiments. Throughout the remainder of the year, many details of these systems were completed, and many additional diagnostics were added. During this first phase, target positioning, viewing, and alignment were the top priority one-third of the time. One of the difficulties we had to overcome was the fact that the target chamber flexes upon pumpdown, changing the alignment of some instruments by up to 1 mm. All target shots were taken with 0.53- μm light. In April, we installed the target-plane imager, allowing us to begin to shoot targets with 0.35- μm light in May.

The second phase extended from May through August. A major emphasis was on the laser-alignment systems. We greatly increased our ability to accurately position the focus lenses for target shots at 0.35 and 0.53 μm . By the end of the activation effort, we were able to routinely align the

Fig. 4-32. Cumulative target shots vs months after the end of construction for Shiva, Novette, and Nova. The two curves for Nova show the total target shots into both target chambers and the shots into the ten-beam target chamber.

laser beams and the target to a relative accuracy of $\pm 100 \mu\text{m}$. This accuracy is better by a factor of 2 than that achieved on Shiva and Novette and is less by a factor of 2 than that ultimately required for high-density implosions.

Another major emphasis during our second phase was the laser diagnostic systems. We activated several electronic and optical subsystems, bringing us closer to measurement of the incident $0.35\text{-}\mu\text{m}$ energies, incident pulse shapes, reflected energies, and reflected pulse shapes. We also completed the first set of measurements on the instrument that is described in the next section (see "Opposed Port Diagnostics"). This instrument measures the spatial distribution of energy in the high-power beam near best focus. In addition to these efforts, we fired target shots to activate several target diagnostics, and we installed a number of new target diagnostics that are described later in this article.

The final phase began in September 1985. We obtained an average of 18 target shots per month and used them to activate the target diagnostics needed for hohlraum characterization and other experiments addressing target physics and interactions between light and matter. Our strategy was to execute a sequence of brief, simple experiments. Each experiment

required a different partial set of target diagnostics and provided a strong motivation to obtain valid data from our diagnostics. The experiments included:

- X-ray conversion efficiency from gold-disk targets.
- Absorption, drive, and preheat of hohlraum targets.
- Raman scattering from exploding-foil targets.
- Ablation pressure produced on directly-driven aluminum targets.
- Neutron yield produced by directly-driven implosions. Preliminary results of these experiments can be found in Sections 2 and 3.

Target Diagnostics. We required 5 to 10 shots per diagnostic instrument to prepare for complex target experiments on Novette. Our experience during the Nova ten-beam activation was similar. We fired 160 shots and activated 18 diagnostic instruments. Figure 4-33 shows a summary of our progress in diagnostic activation. The two curves in the figure show the cumulative number of diagnostics that were mechanically installed and the number fully activated. A diagnostic is mechanically installed when the major mechanical hardware is in place on the Nova target chamber. After mechanical installation, we may require vacuum, air, and nitrogen plumbing

Diagnostics

as well as checkouts of electrical wiring and controls. Our attempts to obtain data often reveal the need for minor modifications. Eventually we obtain reliable data from an appropriate test shot and evidence that background signals from unintended sources are not a problem. The diagnostic is then ready to produce data for physics experiments and is fully activated.

Fig. 4-33. Cumulative number of target diagnostics installed and activated on the Nova ten-beam chamber vs time, beginning at the end of the Nova construction project.

By the end of activation, we required an average of 2.5 months after mechanical installation to activate a diagnostic.

Table 4-9 is a brief description of each of the target diagnostics we installed during Nova activation. The table lists the diagnostic name, the type of detector, the number of channels, and the type of measurement. Nova contains several multi-channel diagnostics and many diagnostics using streak cameras. Despite the need to build entirely new instrumentation and controls for most diagnostics to fit the Nova target chamber and the Nova controls environment, Nova now has a capability comparable to that of Novette to diagnose the results of target experiments. (The Novette target chamber had the most complete set of target diagnostics in the history of the ICF Program.)

Using these diagnostics, we can measure the emissions from a target using imaging, spectroscopy, time-resolution, and absolute measurements, as shown in Table 4-9. We use the x-ray diagnostics to see where the laser beams strike the target, to measure the spectrum and distribution of soft

Table 4-9. Summary of target diagnostics on the Nova ten-beam chamber.

Diagnostics	Detector	Number of channels	Type of measurement
X ray			
8 x east microscope	Film	4	Image 1-5 keV
8 x west microscope	Film	4	Image 1-5 keV
X-ray pinhole camera #1	Film	8	Image 2-10 keV
Reentrant minispectrometer	Film	1	Spectrally resolved image 1-4 keV
Filter-fluorescer hard x-ray spectrometer	Photomultiplier tube	20	Absolute spectrum 3-250 keV
Dante U spectrometer	X-ray diode	10	Absolute spectrum 0.1-1.5 keV
Dante D spectrometer	X-ray diode	10	Absolute spectrum 0.1-1.5 keV
Henway #1 crystal spectrometer	Film	4	Absolute spectrum 2-15 keV
Spatially discriminating streaked spectrometer #1	Streak camera	1-4	Time-resolved 1-D image or spectrum 0.1-1.2 keV
Streaked crystal spectrograph	Streak camera	1	Time-resolved high-resolution spectrum 0.5-10 keV
OX x-ray	Streak camera	1-4	Time-resolved low-resolution spectrum 5-50 keV
Optical			
Energy balance	Diode	180	Absolute J/sr <350-1100 nm
Streaked optical spectrometer #1	Streak camera	1	Spectrum 300-1100 nm
Streaked optical spectrometer #2	Streak camera	1	Spectrum 200-800 nm
Streaked optical pyrometer	Streak camera	1	Absolute, time-resolved 1-D image 200-800 nm
OX (optical)	Streak camera	4	Relative timing of various wavelengths 200-800 nm
Neutron			
Long copper activation	Photomultiplier tube	1	Neutron yield >10
Short copper activation	Photomultiplier tube	1	Neutron yield >4 x 10 ⁸
Neutron time of flight	Photomultiplier tube	1	Ion temperature >0.2 keV

x rays produced by the target for x-ray drive, to measure the spectrum of hard x rays produced by suprathermal electrons in the target, to measure the relative timing of various x-ray emissions and other events, and to measure the x-ray transmission of the target in backlighting experiments. We use the optical diagnostics to measure the absorption of laser light by the target, to detect the emergence of shock waves from materials, and to detect evidence of the parametric processes in the plasma. We use the neutron diagnostics to measure the neutron yield from and ion temperature of an imploded capsule.

In addition to other target diagnostics, we installed an optical fiducial system described in "Optical Fiducial for X-Ray and Optical Streak Cameras." This system provides an optical signal to all the x-ray and optical streak cameras in Nova. It allows timing of all target-diagnostic data relative to the laser pulse with an accuracy of ± 50 ps. Such timing is critical for accurate analysis and interpretation of experimental results.

Author: R. P. Drake

Diagnostic Systems on the Two-Beam Chamber

We have configured the target diagnostics for the first experimental series on the Nova two-beam chamber for x-ray laser experiments. Most diagnostics for this series are upgraded versions of the instruments used on the final Novette x-ray laser experiments. The objectives of the upgrades are to improve instrument alignment capabilities and diagnostic sensitivity and to minimize operational maintenance.

Among our instruments are two x-ray pinhole cameras that record the quality of beam overlap and line foci. Two Henway spectrometers record the time-integrated soft x-ray emission between 4 and 14 Å. A chamber isolation valve was added to both the pinhole cameras and Henway spectrometers to facilitate film changes between shots. A target changer is being added so that the chamber need not be vented to change targets.

The LLNL streak camera used on the Novette TGSS was characterized temporally and spatially at the monojoule laser facility. The streak camera cathode and ellipsoidal mirror were also characterized before the TGSS was installed at one of the two on-axis locations. A detailed description of the TGSS calibrations is provided in the article, "Calibration of TGSS," later in this section.

We have installed an upgraded diagnostic for the McPIGS opposite the TGSS to record the on-axis lasing lines. A second McPIGS (unmodified from Novette) was mounted off-axis to give information on the nonlasing x-ray emission lines. The sensitivity of both McPIGS was enhanced by adding new fiber optics, a film cassette, and a repeller grid.

The current TGSS configuration gives continuous temporal information over a 41-Å window within a range from 120 to 260 Å ($E/\Delta E \approx 200$). The McPIGS complements the TGSS by providing a wider spectral window (typically 149 to 270 Å) and improved spectral resolution ($E/\Delta E \approx 1500$) but has limited and non-continuous temporal information using the Auston switch gating system. (See "Auston Switch System," in the next section for details of that system.) By changing the grating, the McPIGS spectral range can be extended down to 70 Å.

We plan to use a five-channel diode array to calibrate both McPIGS and TGSS, using continuum emission from a gold target as a reference spectrum. The diode array will be built from a 10-channel Dante spectrometer from Novette. The array will be modified to provide us with the capability of mounting the instrument in several port locations.

In addition to the upgraded Novette diagnostics, we have fielded several new instruments. An x-ray crystal streaked spectrograph measures off-axis emission in the range of 4 to 9 Å. The spectral window can be adjusted from <1 to 2.5 Å, with the corresponding resolution ranging from $E/\Delta E \approx 4000$ to 500. Two spectral regions or two different dispersions can be covered simultaneously.

Finally, an optical imaging camera was built to give us information on the

line-focus quality and beam position on the target. Several new diagnostics, including a gated crystal spectrometer and an x-ray framing camera, are being built and will be fielded in early 1986.

Authors: D. G. Nilson and D. L. Matthews

X-Ray Diagnostics

McPIGS Spectrograph Improvements.

The McPIGS were successfully used in x-ray laser experiments on Novette.³⁸ During 1985, we modified McPIGS spectrographs to obtain added versatility.

Our first modification was the use of a cylindrical mirror, which was placed in front of the spectrometer to correct for the astigmatism of its Rowland circle geometry. Figure 4-34 shows the setup in a trial experiment on the Janus laser. By bringing light from the source to a stigmatic focus on the microchannel plate, the cylindrical mirror produced spatial resolution in the sagittal direction (along the spectral lines) with a large magnification ($\times 5.5$ in the example shown) and normal spectral resolution (typically 1500).

The mirror also served as a collecting optic, allowing the spectrometer to subtend a greater solid angle at the source and, hence, to increase the sensitivity of the instrument. For the setup shown in

Fig. 4-34, the enhancement due to the mirror was a factor of 100 for a $100\text{-}\mu\text{m}$ source. The x rays were incident on the 25.84-mm -radius cylindrical mirror at a grazing angle of 4.3° .

The mirror is a cylindrical negative lens with a $100\text{-}\text{\AA}$ gold coating. This mirror was estimated to have a reflectivity of 80% in the wavelength region from 140 to $260\text{ \AA}$ covered by the spectrometer. The mirror was placed 203 mm from the center of the Janus target and 1118 mm from the center of the microchannel plate detector. The mirror was shielded to give an 8-mm -wide aperture and thus had a depth of focus yielding more than $40\text{-}\mu\text{m}$ spatial resolution over most of the spectral range. Spatial resolution was verified by spectra taken on Janus from irradiation of aluminum-slab targets.

We aligned the sagittal mirror in situ using a fiber-optic emitting He-Ne light. We placed the light at chamber center, looking optically for a focus in the zero-order reflection from the spectrometer grating at the same distance from the grating as the center of the microchannel plate. First-order diffracted x rays come to a stigmatic focus on the microchannel plate. The microchannel plate had three 2.1-mm -wide strips of MgF_2 -coated photocathodes that could be pulsed at time scales of 200 ps . When operated in this spatially resolving mode, only one strip is useful,

Fig. 4-34. Schematic of McPIGS spectrograph showing the 25.84-mm cylindrical mirror with sagittal focusing.

thereby making the spectrometer a single-frame camera.

We plan to use a sagittal mirror during x-ray laser experiments on the Nova two-beam chamber to image the x-ray laser emission from exploding-foil targets. The mirror will also be used as a collimator to measure the divergence of the x-ray laser beam.

Our second modification of McPIGS was a biased electric field above the surface of the microchannel plate to deflect electrons emitted from the surface back onto the plate. The deflected electrons can be amplified and can contribute to the observed signal.⁴⁹ When a microchannel plate is operated at grazing incidence, most electrons are emitted from the webbing surrounding the channels. Only electrons emitted in the holes, or close to the rim, are attracted down the channel and contribute to an electron avalanche. By providing an electric field to bias the electrons back to the plate, we found that the sensitivity of the microchannel plate could be increased by a factor of 6. The enhancement is due to the deflected electrons preferentially entering a channel on their return to the plate due to the potential gradient inside the channel. A curved metal plate was fixed 15 mm away from the microchannel plate and held at a potential of -10 kV. The microchannel plate cathode was operated at a maximum potential of -1 kV. The large bias field was required due to edge effects near the edge of the 2.1-mm-wide cathode strip. The microchannel plate beside the cathode strip remains at ground potential while the cathode is

pulsed, which modifies the effect of the bias field.

Our two modifications have given the McPIGS spectrographs greater sensitivity, when required. The imaging capability of the mirror will provide data on the location of the gain region in x-ray laser plasmas.

Author: B. J. MacGowan

Transmission-Grating Streak Spectrometer. The TGSS was installed and is operational on Nova's two-beam target chamber for x-ray laser experiments. Its primary functions are to measure the absolute time history and the absolute intensity of the x-ray laser output. The TGSS has been upgraded since its previous use in x-ray laser experiments, and a new, absolute, energy-calibration scheme has been initiated.

The TGSS⁵⁰ consists of an elliptical grazing-incidence mirror, a free-standing gold transmission grating with a 2000-Å period, and an LLNL soft x-ray streak camera. As shown in Fig. 4-35, light from the x-ray laser is refocused by an elliptical mirror onto the photocathode of the streak camera. This light is spectrally dispersed across the cathode by the transmission grating, thus providing high time resolution (20 to 30 ps) and moderate wavelength resolution ($E/dE \sim 200$).

By adding and reworking certain components, we improved the operation and durability of the TGSS. We altered the stray-light shield so that only light focused by the mirror through the grating reaches the streak camera. We also placed a

Fig. 4-35. TGSS schematic.

holder for interchangeable debris filters in front of the mirror to protect it and the grating from target debris. Such protection was shown to be necessary by our past experience. In the x-ray laser experiments of 1984, the grating was destroyed during a target shot, probably because of target debris. When thin beryllium and aluminum filters were used in the holder, they lasted for only two to five shots. Now we are using a 1000-line-per-in. nickel mesh to stop debris—with good results. This mesh appears to be much more robust than the thin metal filters.

We improved the alignment accuracy of the TGSS and the x-ray laser target by increasing the magnification of the alignment telescope by a factor of 10. Increased magnification also allows us to inspect the

grating, filters, and mesh for damage sites as small as $20\text{ }\mu\text{m}$.

Obtaining an exact time history of the laser transitions is crucial in understanding the physics of the laser-pumping scheme. Therefore, we improved the time resolution of the TGSS and installed an absolute timing fiducial. The time resolution was increased by using a higher sweep rate on the streak camera and by adjusting the spatial resolution of the streak camera to better match the resolution achieved by the grating, mirror, and source size. Operating the camera with a $200\text{-}\mu\text{m}$ spatial resolution results in a 20-ps time resolution. We can achieve even better time resolution by decreasing the spatial resolution.

An absolute timing fiducial will be operating on Nova in March 1986. It consists of a $0.53\text{-}\mu\text{m}$ light pulse with a width less than 100 ps. The pulse is delivered to the streak-camera cathode on a fiber optic. Its timing is known relative to the main Nova laser pulse. Preliminary testing of the TGSS streak camera has shown that its spatial resolution and general operation was unaffected by the presence of the fiber optic in its cathode. More information about the timing fiducial can be found in the article "Optical Fiducial for X-Ray and Optical Streak Cameras," later in this section.

We are also developing a new calibration technique for the TGSS. Previous energy calibrations used a laser plasma for

Fig. 4-36. Dynamic-vs-static sensitivity plot for the TGSS streak camera.

Fig. 4-37. Time history of the selenium 206-Å x-ray laser line.

Fig. 4-38. Time history of the selenium 209-Å x-ray laser line.

the soft x-ray source.⁵¹ While sufficiently bright for a dynamic calibration, the transient plasma is difficult to characterize accurately. An alternative is to use well characterized, continuous x-ray sources such as those developed by Henke⁵² or Manson.⁵³ This technique requires measurement of the relative dynamic-vs-static sensitivity of the streak camera. Then the absolute calibration performed with a continuous exposure can be corrected for any dynamic effects in the camera that could affect the sensitivity.

As required for this type of calibration, we measured the relative dynamic-vs-static sensitivity of the streak camera. In this measurement, the streak camera viewed a laser-produced tantalum plasma of ~ 200 -ps duration. We used an aluminum photocathode with a beryllium filter to allow detection of the continuum from above 50 eV. A second streak camera with a similar photocathode and filter was used as a reference detector. With both cameras operating in streak mode, we determined their relative sensitivity to a wide range of x-ray intensities. Then, we operated the TGSS streak camera with no voltage on the sweep plates to achieve cw operation. By comparing the cw exposures obtained on the TGSS camera with the dynamic exposures obtained on the reference camera,

we determined the relative static-dynamic sensitivity. Figure 4-36 shows the results of these measurements. The relative sensitivity is linear over the range of exposures anticipated for our x-ray laser experiments, and the static sensitivity is basically the same as the dynamic sensitivity.

We also made preliminary energy calibrations of the streak camera by using a Manson source. The absolute sensitivity at 277 and 109 eV agrees with previous calibrations.⁵¹ Calibrations at lower energies, with the mirror and grating in place, are under way.

The TGSS has already provided time-resolved data for selenium and yttrium x-ray lasers. Figures 4-37 and 4-38 show the time history of the 206- and 209-Å selenium laser lines. Figure 4-39 shows the 155-Å yttrium laser line. All three have durations of 180 ps. Both lasers appear to turn on ~ 200 ps after the continuum. The exact timing relative to the main Nova laser pulse will soon be determined by using the fiducial.

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Major Contributor: T. W. Phillips

X-Ray Laser Cavities. Recent demonstrations of amplified spontaneous emission (ASE) at soft x-ray wavelengths have

Fig. 4-39. Time history of the yttrium 155-Å x-ray laser line.

spurred interest in normal-incidence, multilayer structures for use as x-ray laser (XRL) cavity mirrors. Calculations indicate that cavity mirrors can significantly enhance XRL output and drive the oscillations to the saturation limit. Novel ideas for cavity-output coupling might expand the utility of XRL devices and encourage the marriage of lithographic and multilayer technologies for the general benefit of x-ray optics. Preliminary tests of a double-pass cavity have already been conducted.

Recent experiments at LLNL demonstrated the generation of significant levels of ASE at soft x-ray wavelengths.⁵⁴ In these experiments, the ASE emerging from the ends of a gain-medium channel produced multikilowatt bursts of line radiation at 206.3 and 209.6 Å. These soft x-ray laser lines are from 3p → 3s transitions in neon-like selenium ions. In these experiments, the plasma gain medium was produced by an intense pulse of green laser light ($\sim 5 \times 10^{13}$ W/cm²) from a frequency-doubled Nd:glass laser. The duration of the driving laser pulse was ~ 450 ps, and the duration of the soft x-ray ASE was less than 200 ps. The total length of the gain channel was nominally 1 cm. The total small-signal gain length αL was approximately 6.

Mirrors for XRL Cavities. The reliable production of a macroscopic amplification medium for soft x rays strongly motivates the use of normal-incidence multilayer mirrors in an XRL cavity. Such a cavity would significantly increase the XRL output. The unsaturated output of the

transient emission could be expected to increase by the factor

$$F = T \sum_{k=0}^N R_o^i R_b^j e^{k\alpha L}, \quad (39)$$

where

- T = output coupling efficiency,
- R_o = reflection coefficient of the output mirror,
- R_b = reflection coefficient of the back mirror,
- α = net gain of the amplification medium,
- L = length of the gain medium,
- N = total number of single bounces the amplified beam can be expected to make during the finite gain lifetime,
- $k = i + j$, where $i = j$ for k even, and $i = j + 1$ for k odd.

We conducted a preliminary investigation of the viability of XRL cavity concepts during the LLNL experiments. A single, flat, molybdenum/silicon multilayer mirror was used at normal incidence to produce double-pass amplification of the ASE from the gain medium. (Table 4-10 lists the mirror's characteristics.) The parameters for the double-pass XRL experiment were as follows: $T = 1$, $N = 1$, $i = 1$, $j = 0$, $R_b = 0.25$, $\alpha L = 6$, yielding $F = 100$ according to Eq. (39).

Figure 4-40 compares the x-ray spectra obtained in this experiment with and without the multilayer mirror. Signal enhancement of 2 to 3 × was observed—an

increase much less than expected. Possible explanations for this result include mirror damage, gain-medium dynamics, or mirror misalignment.

We are investigating the important issues of mirror alignment, calibration, and survival. Alignment requirements are not severe. Milliradian precision is required in aligning the optically opaque cavity mirrors normal to the axis of the plasma-gain channel. As for calibration, multilayer x-ray mirrors must be characterized at precisely the wavelength of the x-ray laser because of the uncertainty in published optical constants in that wavelength range. Such measurements are best made at or near normal incidence. Survival issues are important because cavity mirrors must perform in a very hostile x-ray environment. The lifetime of a multilayer x-ray mirror in such an environment requires experimental investigation.⁵⁵

If the issues of mirror alignment, calibration, and survival can be resolved, calculations indicate that significant increases in single-pass XRL output can be achieved by using a multilayer mirror cavity. Figure 4-41 shows calculations of reflectivity and transmission vs x-ray energy for the two mirrors of an XRL cavity under consideration. Figure 4-41(a) illustrates the properties of the partially transmitting output mirror. Figure 4-41(b) shows the properties for the totally reflecting back mirror of the cavity. The mirrors represented in this figure are Mo/Si multilayers with a thin overcoat of carbon. The output mirror is assumed to be a freestanding membrane of the multilayer materials. Detailed parameters for the mirror design are listed in

Table 4-11. In calculating the curves in Fig. 4-41, we used the following optical constants from Ref. 56: $f_1^{\text{Mo}} = 17.0$, $f_2^{\text{Mo}} = 2.5$, $f_1^{\text{Si}} = 3.8$, and $f_2^{\text{Si}} = 0.5$. These constants provide best fits to recent experimental data.⁵⁷

This cavity design is expected to produce the following XRL performance characteristics: $T = 0.09$, $R_o = 0.4$, $R_b = 0.6$, $\alpha L = 6$, $L = 2$ cm, and $t \geq 200$ ps ($N = 3$). Performance with these characteristics would yield an enhancement factor F well in excess of 10^5 . Such amplification could be expected to exceed small-signal gain conditions and drive the

High-Z material	Mo
Low-Z material	Si
Fraction of high-Z material (f)	0.38
Number of layers (n)	20
Thickness of single Mo/Si layer pair (d)	115 Å
Overcoat material	Carbon
Overcoat thickness	10 Å
Substrate material	Polished Si wafer

Table 4-10. Multilayer-mirror characteristics for double-pass XRL cavity.

Fig. 4-40. X-ray spectra obtained with and without a multilayer XRL-cavity mirror.

Fig. 4-41. Calculated reflectivity and transmission for the two mirrors of the XRL cavity described by Table 4-11.

	Output mirror	Back mirror
High-Z material	Mo	Mo
Low-Z material	Si	Si
Fraction of high-Z material (f)	0.1	0.4
Number of layers (n)	20	20
Thickness of single Mo/Si layer pair (d)	107 Å	112 Å
Overcoat material	Carbon	Carbon
Overcoat thickness	50 Å	5 Å
Substrate	None	Polished quartz mirror ($f = L$)

Table 4-11. Mirror ▲ parameters for XRL cavity design.

Fig. 4-42. Concepts for output coupling from an XRL cavity.

Pinhole diameter	75 μm
Laser "footprint" diameter	250 μm
Output mirror reflectivity (outside pinhole region)	0.6
Net reflectivity of output mirror (R_0)	0.55
Transmission of output mirror (T)	0.09

Table 4-12. Design parameters for pinhole output-coupling mirror.

cavity oscillations into saturation. It would then be possible to analyze and characterize the saturation conditions of the gain medium. [Equation (39) assumes unsaturated gain conditions, so the above calculation of F is intended only to demonstrate that significant amplification can be achieved, even with relatively inefficient mirrors and only a few round trips of the beam.]

Output Coupling from XRL Cavities. In addition to an increased output signal and the opportunity to observe saturated behavior in media of limited gain-length, XRL cavities provide opportunities for greater flexibility in output coupling. This potential may encourage the development of applications and diagnostic efforts to characterize the laser itself.

We are investigating a number of concepts for output coupling from an XRL cavity. Some of these ideas are illustrated in Fig. 4-42. Figure 4-42(a) shows the simplest output-coupling mechanism, the pinhole. In this case, the fraction transmitted is merely the ratio of the area of the pinhole to the multimode XRL "footprint" on the mirror. Examples of the parameter values under consideration are listed in Table 4-12. The simple pinhole is a fairly efficient, low-technology approach to multimode output coupling.

Figure 4-42(b) shows a "beam splitter"—the partially reflecting output mirror discussed in connection with Fig. 4-41. While slightly less efficient and more difficult to fabricate than the pinhole, the beam splitter has the advantage of spatially uniform feedback across the radial profile of the gain medium. This property could be important under transient gain conditions, which might preclude radial-mode mixing and make pinhole output coupling impractical.

The most challenging and potentially useful output coupling scheme is illustrated in Fig. 4-42(c). In this case, a diffraction pattern is lithographically produced on or in the multilayer structure itself. The pattern may provide amplitude modulation or phase modulation (for greater efficiency) of the mirror reflection, and it may be periodic or aperiodic (for focusing). The reflected x rays are diffracted

into multiple orders R_i , providing both cavity feedback R_0 and diffracted output R_i , where $i \neq 0$.

Fabrication and Testing of XRL Cavity Components. Recent progress has been made in fabricating and characterizing x-ray optical components for use in XRL cavities. Some of these achievements are listed below:

1. Fabrication and testing of x-ray beam splitters at normal incidence over the wavelength range 100 to 300 Å.
 2. Testing of multilayer x-ray mirrors at near-normal incidence (within 0.5°) over a wavelength range of 100 to 300 Å. Peak reflectivity of 25% was achieved at 206 to 209 Å.
 3. Fabrication and testing of high-dispersion (2000-Å period) gratings on multilayer mirrors.
 4. Fabrication of zone-plate patterns on multilayer mirrors.
- In addition, experiments scheduled for Spring 1986 will test cavity-mirror survival in the hostile x-ray laser environment.

We are committed to the use of multilayer technology for XRL cavity development. The prospects for increasing the output power of soft x-ray ASE devices look promising. The coupling of submicron lithographic and multilayer technologies will spawn a new class of x-ray optical components that promise greater flexibility and wider application for XRL devices.⁵⁸

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Solid-State X-Ray Detectors. The ICF and Nuclear Test Programs use x-ray diodes (XRDs) extensively for x-ray diagnostics. XRDs are photodiodes based on photoelectron generation and collection in vacuum. They have a typical time response of 0.2 to 1 ns. Because these detectors are relatively large and slow, we are investigating solid-state devices for use as x-ray detectors. In addition to being smaller than XRDs, solid-state x-ray detectors will have better spatial and temporal resolution. They may eventually supersede

vacuum XRDs in the same way semiconductor transistors made thermionic-emissive vacuum electron tubes obsolete.

We are investigating two types of solid-state x-ray detectors with flat spectral response, fast temporal response, and good spatial resolution: photoconductors and pyroelectric detectors. Their properties will allow us to use these devices not only as fast x-ray or γ -ray detectors, but also as fast x-ray calorimeters. Because of their small size, one-dimensional arrays of such detectors should be useful for time-resolved spectrometers, and two-dimensional arrays could serve in time-resolved imaging applications. A fast x-ray calorimeter would be invaluable because it could measure the total x-ray emission from the ICF capsule.

When used in conjunction with the directional coupler described later in this section, solid-state detectors could miniaturize many standard diagnostic systems. For instance, the "Dart"—a filtered XRD or a filter-fluorescer XRD broadband x-ray spectrometer routinely used in the test program—could be transformed to a "needle," relatively speaking.

Photoconductors. Upon absorbing radiation, semiconductors and insulators may become temporarily conductive electrically. According to Shockley,⁵⁹ the average energy required to produce an electron-hole pair in a solid is independent of the energy and the mass of the incident particle. The number of electron-hole pairs produced is proportional to the absorbed energy of the incident particle. Thus, the photoconductor theoretically has a flat spectral response to x rays. Unlike semiconductor junction detectors (such as pin diodes, surface-barrier diodes, or lithium-drifted detectors), photoconductors can be made free of dead layers; thus they can be used as low-energy x-ray bolometers. The spectral response of an InP:Fe photoconductor appears to be flat for 0.7- to 3-keV x rays within experimental uncertainties.⁶⁰

Different photoconductors vary widely in sensitivity and in temporal response. A response to optical excitation in as short a time as 4 ps FWHM has been reported for amorphous and ion-implanted silicon.⁶¹

Fig. 4-43. Typical x-ray pulse from Nova target observed with a photoconductor detector.

Fig. 4-44. Typical x-ray pulse from Nova target observed with a pyroelectric detector.

Some photoconductors can be made far more sensitive to x rays than the XRD and far more sensitive to γ rays than the vacuum Compton diode. Our goal is to study systematically the sensitivity, linearity, spectral response, and temporal response of various photoconductors to x rays and γ rays. We will also investigate the effects of impurity doping and radiation damage. So far, we have obtained or fabricated detectors made of InP:Fe (neutron-damaged and undamaged), GaAs (neutron-damaged and undamaged), GaAs:Cr, and polycrystalline silicon.

Pyroelectric Detectors. When a pyroelectric crystal is heated, its spontaneous dipole moment changes. The heating of a pyroelectric crystal due to the absorption of incident particles thus induces a displacement current that can be detected by means of an appropriate electric circuit.

The spectral response of a pyroelectric detector to low-energy x rays is inherently flat because the x rays are totally absorbed by the crystal and converted into heat. The temporal response of the pyroelectric detector can be fast. It is ultimately limited by the period of phonon vibrations of the crystal lattice, which is ~ 1 ps.⁶² In practice, the detector speed is limited by the RC constant of the circuit.

A major advantage of the pyroelectric detector is the fact that it is a passive device; no bias voltage is required. Its sensitivity is rather low ($\sim 10^{-7}$ A/W), a disadvantage for some applications. However, for the high-flux measurements we make, a relatively low sensitivity can be advantageous. We have obtained or fabricated LiTaO_3 , $\text{Sr}_{0.5}\text{Ba}_{0.5}\text{Nb}_2\text{O}_6$, and LiNbO_3 detectors.

Detector Testing. We are testing the two types of solid-state detectors with x rays and γ rays produced at various facilities:

- 1-ns low-energy x rays from Nova.
 - 3- and 30-ns, 100-keV x rays from the LLNL Febetron electron-beam generators.
 - 50-ps, 100-MeV γ rays from electron linear accelerators at LLNL and Santa Barbara.
 - 100-ns low-energy x rays from Zapp, the LLNL Z-pinch plasma source.
- Tests at the LLNL Monojoule, Janus, and 8-ps short-pulse laser facilities are also planned.

Figure 4-43 shows a typical x-ray pulse observed with a $1\text{-cm} \times 1\text{-cm} \times 1\text{-mm}$ InP:Fe photoconductor biased at 100 V and situated 4.4 m from a gold-disk target. A 1.06- μm , 1.5-kJ, 1-ns Nova beam was focused on the target at 3×10^{14} W/cm², and the resulting x rays were filtered with a 25- μm aluminum filter. Figure 4-44 shows a typical x-ray pulse observed with a $1\text{-cm} \times 1\text{-cm} \times 50\text{-}\mu\text{m}$ LiTaO₃ pyroelectric detector situated 4.4 m from a gold target. A 0.35- μm , 1.67-kJ, 1-ns Nova beam was focused onto the target at 2×10^{15} W/cm², and the x rays were filtered with a 0.39- μm aluminum filter. Clearly, pyroelectric detectors are sensitive enough for this type of application.

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The X-Ray Diode Source. We developed and characterized a new soft x-ray source called the x-ray diode source (XRDS). When illuminated with visible or near-infrared light, a multialkali photocathode in the XRDS produces electrons, which are

accelerated onto a thin metal film anode to generate characteristic x rays and bremsstrahlung. The novel feature of this design is the proximity of the photocathode and the anode. They are close enough to yield x-ray emission that is a spatial and temporal replica of the input light. In this article, we describe our measurements of the x-ray brightness and energy spectrum and the minimum source size. We also discuss unique applications of the XRDS.

Design. The XRDS, shown schematically in Fig. 4-45, is a compact, disk-shaped device approximately 2 in. in diameter and 1 in. wide. The simple design incorporates a multialkali (Sb, K, Na, Cs) transmission photocathode mounted 3 mm from a transmission anode. The photocathode is deposited on a glass input window by a process that reduces the sheet resistance to less than 1 k Ω /square. To avoid contamination of the photocathode, the XRDS is permanently sealed under high vacuum. The anode consists of a thin film of metal deposited on a 5-mil-thick beryllium output window. The window diameter of 1 cm constrains the source area to a maximum of 0.8 cm². The XRDS is specified to operate at a maximum dc voltage of 12 kV.

We fabricated two XRDSs, one with a titanium anode 4000 Å thick and the other with a gold anode 500 Å thick. The

thickness of the metal films was chosen to be greater than the range of a 10-keV electron in these materials. The gold film was kept thin to minimize self-absorption of the x rays.

A unique feature of the XRDS is that the x-ray emission is intimately related to the input light signal. The spatial configuration of the x-ray source is a replica of the illumination pattern on the photocathode. Similarly, any temporal modulation of the input light signal is directly reproduced in the x-ray emission.

Characterization. We investigated the performance of the XRDS to determine the emission brightness, energy spectrum, and minimum source size (point spread function). Figure 4-46 shows the configuration

Fig. 4-45. Schematic depiction of the XRDS.

Fig. 4-46. Configuration of the XRDS test station.

of the test station used to evaluate the XRDS. The XRDS was mounted on a flange at one end of a vacuum flight tube. A glass window in the flange permitted the XRDS to be illuminated with $0.63\text{-}\mu\text{m}$ light from a helium-neon laser outside the vacuum.

The x rays emitted by the XRDS were detected with an x-ray CCD camera mounted at the opposite end of the flight tube. The camera, which was custom-built at MIT,⁶³ contains a front-illuminated, virtual-phase CCD made by Texas Instruments. The CCD contains an array of 584×390 pixels, each $22\text{ }\mu\text{m}$ square. During operation, the CCD is cooled to about -60°C and is scanned slowly to reduce the total noise to about 20 electrons per pixel—well below the signal produced by a single x ray. (The number of electrons N produced by an x ray with an energy of $E\text{ eV}$ is approximately given by $N = E/3.6$.) Furthermore, a 5-mil-thick beryllium window installed in front of the camera prevents scattered light from reaching the CCD. Figure 4-47 shows the CCD's calculated x-ray detection efficiency. The combination of low noise and high efficiency makes the CCD camera ideally suited for

experiments that require both imaging and photometric analysis.

We measured the XRDS emission spectrum while the photocathode was uniformly illuminated at an intensity of $\sim 0.1\text{ mW/cm}^2$. Monitoring of the current generated between the photocathode and the anode revealed that the photocathode efficiencies at $0.63\text{ }\mu\text{m}$ were $25\text{ }\mu\text{A/mW}$ for the XRDS with the titanium anode and $32\text{ }\mu\text{A/mW}$ for the one with the gold anode. X rays were collected by the CCD camera for an exposure time of several milliseconds. Exposure was controlled by an electronic shutter located in the laser beam path. The x rays were detected as single events in a 200×200 -pixel subarray on the CCD. To minimize charge spreading during readout, we selected the location of the subarray to be as close as possible to the output amplifier on the CCD. We obtained x-ray spectra by constructing a histogram of the single x-ray events as a function of the charge deposited in the CCD pixel. When the charge was distributed between two or more pixels, an event was discarded. We subsequently adjusted the spectra to take into account the detection efficiency of the CCD and transmission through the beryllium window.

The x-ray emission spectra of the two XRDSs operating at 10 kV are shown in Fig. 4-48. The energy resolution at 4.5 keV is approximately 200 eV. The emission spectrum from the titanium anode in Fig. 4-48(a) is dominated by the $K\alpha$ and $K\beta$ lines at 4.5 keV, which are not separately resolved. In contrast, the spectrum from the gold anode exhibits broadband bremsstrahlung emission across the entire 2- to 10-keV range, with an additional

Fig. 4-47. X-ray detection efficiency of the CCD camera.

Fig. 4-48. (a) X-ray emission spectrum of the titanium-anode XRDS operated at 10 kV; (b) x-ray emission spectrum of the gold-anode XRDS operated at 10 kV.

Fig. 4-49. (a) Pinhole image of the XRDS source spot generated by projecting a 25- μm light spot onto the photocathode; (b) lineout across a diameter of the pinhole image. The dashed line is a theoretical fit with $\sigma = 20 \mu\text{m}$.

contribution from gold M-band transitions at 2.0 to 2.4 keV.

The brightness of the x-ray emission is indicated on the vertical scale of the spectra in Fig. 4-48. The ordinate values were normalized to an anode current of 1 mA and an emission energy bandwidth of 1 keV. The brightness was determined by integrating the total number of x rays detected in the subarray during an exposure. In the case of the XRDS with the titanium anode, approximately 50% of the total x-ray power is emitted into the K lines. The ratio of the power associated with the photoelectrons (as determined from the anode current and operating voltage) to the x-ray power emitted into the titanium K lines implies a conversion efficiency of 3×10^{-4} .

We varied the input light intensity to determine the maximum illumination at which the photocathode could be operated safely. In dc operation, the anode current density was stable up to $\sim 5 \text{ mA/cm}^2$. With more intense illumination, the anode current decreased irreversibly over time, indicating permanent degradation of the photocathode. These results imply that the XRDS can be operated safely at an emission brightness up to $\sim 10^{12} \text{ photons/s} \cdot \text{cm}^2 \cdot \text{sr}$.

We projected an image of the x-ray source spot of the titanium-anode XRDS onto the CCD camera through a pinhole 50 μm in diameter at a magnification of 10. The purpose of this procedure was to measure the spatial dispersion of the photoelectron beam inside the XRDS and thus

determine the minimum source size.

Figure 4-49(a) shows the image of the source spot generated by projecting a 25- μm spot of light onto the photocathode and operating the XRDS at 10 kV. The spreading that occurs in the transfer of the input spot to the output spot is associated with the transverse energy E_t of the photoelectrons as they are emitted from the photocathode. It has been shown⁶⁴ that the transverse energy distribution $P(E_t)$ of the electrons emitted from a multialkali photocathode can be described by

$$P(E_t) = P(0) \exp\left(\frac{-E_t}{\bar{E}_t}\right), \quad (40)$$

where $\bar{E}_t = 0.076 \text{ eV}$ for excitation with 0.60- μm light. A point illumination of the photocathode produces an x-ray source spot $S(r)$ having a Gaussian profile,

$$S(r) = S(0) \exp\left(-\frac{r^2}{\sigma^2}\right), \quad (41)$$

where

$$\sigma^2 = \frac{4d^2\bar{E}_t}{eV}, \quad (42)$$

Here d is the gap width, e is the electron charge, and V is the operating voltage. For XRDS operation at 10 kV, Eq. (41) predicts a point response of $\sigma = 16 \mu\text{m}$. Figure 4-49(b) shows a lineout across a diameter of the image in Fig. 4-49(a). The dashed line represents the best fit to the lineout

Diagnostics

Table 4-13. Operating parameters of x-ray sources.

	Maximum brightness (photons/s · cm ² · sr)	Minimum source size (μm)	Typical power (W)	Water-cooling
XRDS	10 ¹²	<40	0.1-1.0	No
Conventional source	10 ¹⁴	~3000	1000-5000	Yes

Fig. 4-50. Schematic of possible applications of the XRDS: (a) coded imaging by generating a complex x-ray source pattern; (b) phase-sensitive x-ray measurements by temporal modulation of the input light signal; (c) calibration of gated x-ray diagnostics (e.g., the x-ray framing camera) by means of picosecond x-ray pulses.

generated by convoluting Eq. (41) with a $25\text{-}\mu\text{m}$ input spot and a $50\text{-}\mu\text{m}$ pinhole. The best fit was obtained by using $\sigma = 20\text{ }\mu\text{m}$. These results imply that the minimum source size of the XRDS, achieved by illuminating the photocathode at a single point, is a Gaussian spot with an FWHM of $33\text{ }\mu\text{m}$. The average E_f of the photoelectrons generated by $0.63\text{-}\mu\text{m}$ light is 1.1 eV .

Applications. It is instructive to compare the XRDS with a conventional laboratory x-ray source. For example, Ref. 65 describes a typical bright laboratory source that produces x rays by bombarding a water-cooled anode with a focused electron beam generated from a hot filament. Table 4-13 compares some of the operating parameters of this conventional source with the XRDS. The conventional source can be up to 100 times brighter than the XRDS, but its spot size can be no smaller than 1 to 10 mm (Ref. 64). Thus it must be operated at high power. In

contrast, the XRDS can have a very small spot size, enabling it to achieve high brightness at a low power. The absence of water cooling and high-power voltage supplies makes the XRDS much simpler to operate and to incorporate into laboratory experiments. For applications in which source brightness and not just total x-ray flux is important, the XRDS is a desirable alternative to the conventional source.

The utility of the XRDS, however, extends beyond conventional laboratory applications. The ability to directly control the x-ray emission through the input light signal could inspire some unique applications. For example, when an image is projected onto the photocathode, a replica of the image is produced as the x-ray source. Thus it is possible to generate x-ray emission in complex geometric patterns. Three-dimensional coded imaging could be produced by a coded source such as a Fresnel zone-plate.^{66,67} The object to be studied would be placed between the XRDS and the detector as depicted in Fig. 4-50, and a coded image of the object would be produced on film. The three-dimensional real-space image could be reconstructed digitally or by using standard holographic techniques.

Temporal modulation of the x-ray emission by means of the input light signal leads to novel applications such as phase-sensitive x-ray detection, depicted in Fig. 4-50(b). The intensity of the input light signal is modulated sinusoidally at frequencies up to $\sim 1\text{ MHz}$ by an acousto-optic modulator. Weak x-ray signals, such as the higher-order reflections from multilayer mirrors, can be analyzed in phase with the reference frequency to improve the signal-to-noise ratio by many orders of magnitude.

As an extreme example of temporal modulation, it should be possible to generate ultra-short x-ray pulses by applying picosecond light pulses to the XRDS input. Figure 4-50(c) illustrates how the XRDS can be used to measure the temporal resolution of a gated x-ray detector such as an x-ray framing camera. The picosecond light

Fig. 4-51. Prototype model of thermoelectrically cooled CCD camera head.

Table 4-14. Characteristics of the thermoelectrically cooled x-ray CCD camera.

Property	Demonstrated characteristic
Energy resolution (FWHM)	118 eV at $E_x = 0.68$ keV
Operating temperature	$-10^\circ\text{C} < T < -30^\circ\text{C}$
Low-energy limit	$E_x = 0.28$ keV (Carbon K_α)
Gate oxide/depletion layer transition-zone thickness	< 800 Å
Readout noise	11 electrons rms
Charge-transfer efficiency	
parallel	> 0.99999
serial	> 0.99998
Dark current	10 pA/cm ² at $T = 25^\circ\text{C}$
Partial-charge collection events	$< 5\%$ for $E_x = 2.6$ keV

pulses can optically trigger the camera as well as drive the XRDS. The frame duration of ~ 50 ps is readily measured by monitoring the integrated response of the camera while varying the relative delay between the x-ray pulse and the trigger pulse.

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Development of X-Ray Charged-Coupled Device Camera. Last year we reported on a direct x-ray imaging charge-coupled device (CCD) camera that covers a spectral range from 0.5 to 15 keV. The CCD camera is cooled by liquid nitrogen to operate in a temperature range from -60°C to -100°C . It currently operates with an equivalent noise figure of ~ 30 electrons rms and exhibits a dark current on the order of 50 pA/cm².

Our work during 1985 focused on two goals. The first goal was to increase the operating temperature so that thermoelectric cooling could be used to replace the liquid nitrogen system. This improvement would result in a smaller, lighter, and easier-to-operate camera unit. The second goal was to decrease the lower energy limit of the CCD camera sensitivity so that it could be used in subkilovolt x-ray applications. We have achieved significant progress in both areas.

Figure 4-51 shows a prototype model of the thermoelectrically cooled CCD camera head. The design is considerably more compact (a volume of ~ 1.5 vs ~ 30 L) and considerably lighter (a weight of a few pounds vs ~ 70 pounds) than the previous

camera head that was cooled by liquid nitrogen. In addition, the CCD cooldown time has been decreased by more than one order of magnitude.

Table 4-14 lists the characteristics of the new thermoelectrically cooled CCD camera. In addition to the increase in operating temperature, thereby allowing use of the thermoelectric cooler, we have made advances in reducing readout noise to 11 electrons rms and reducing dark current to 10 pA/cm². The success in noise reduction and the reduced oxide layer on the CCD surface (to less than 800 Å) has improved the low-energy sensitivity and energy resolution of the camera.

Figure 4-52 shows a nondispersive energy spectrum of characteristic x rays from a Mylar target measured with the CCD camera. The carbon K_α x rays at 277 eV can be clearly resolved. In the future, we will focus on further reduction of CCD readout noise and on increased low-energy sensitivity of the CCD camera.

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Microfabrication Laboratory. Microstructure fabrication has grown increasingly

Fig. 4-52. The x-ray spectrum from a Mylar target measured with a CCD camera.

important in developing and evaluating diagnostics for the LEAD Program. During the past year, we used optical lithographic techniques to produce a variety of structures. Their fabrication had one requirement in common: accurate control over some critical dimension to micron, and often submicron, tolerances.

As our microfabrication laboratory becomes more heavily involved in x-ray optics, dimensional control on the nanometer scale, and often on the angstrom scale, is required. For example, a Fresnel zone-plate lens with a minimum feature size of 1000 Å needs accurate line placement to some fraction of the minimum line width, or about 100 Å, and for spectroscopic resolving powers of 100 or greater, an x-ray diffraction grating or an x-ray multilayer mirror typically needs accurate period control to a few angstroms. To meet such demanding specifications, several new pieces of equipment have been brought on-line:

- X-ray lithography system for submicron replications. Built in our laboratory, this system employs a novel, interchangeable x-ray anode shaped like an inverted cone for efficient cooling. Electrostatic

force maintains contact between the lithographic mask and wafer, and a capacitance meter is used to monitor this contact by measuring the capacitance between the mask and wafer. This process has been used to replicate submicron structures.

- Argon ion beam. We purchased a commercial argon ion beam system from Ion Tech, Inc., Fort Collins, Colo., and modified it to save space in the laboratory. Now in operation, the system is capable of bombarding samples mounted on a water-cooled, rotating platform with a uniformity of $\pm 5\%$ over a 4-in. diameter. We can obtain ion energies from 50 to 2000 eV and ion current densities up to 5 mA/cm². The purpose of the ion beam is to etch materials—usually to transfer patterns from one material to another. For example, we used it to transfer a 2000-Å-period diffraction grating in resist to a layer of gold <1000 Å deep. The ion beam system is used in many of the laboratory's daily procedures.

- Cesium iodide (CsI) thermal evaporator. Some of the photocathodes we routinely produce consist of CsI deposited onto a substrate, such as 1000 Å of Formvar. For CsI deposition, we now have

a dedicated thermal evaporator equipped with a crystal thickness monitor.

- Reactive-ion etching (RIE) system. We built an RIE system that can handle up to five etching gases simultaneously. It is designed for anisotropic etching of materials in an rf plasma to produce patterns with straight sidewalls and high aspect ratios. Etching of submicron patterns in similar systems has been reported in the literature.⁶⁸ Materials etched so far in our system include Si, SiO₂, Si₃N₄, Mo, and polyimide. We produced anisotropic etches in polyimide with straight sidewalls and a minimum line width of 5 μ m (Fig. 4-53).

With the above facilities and those reported earlier,⁶⁹ the microfabrication laboratory supplies several different types of structures to the LEAD program. Our most time-consuming responsibility is to fabricate appropriate photocathodes for all experiments and to store replacement photocathodes. Typically these are thin layers of aluminum, gold, or CsI on Formvar in various shapes.

Besides photocathode work, the microfabrication laboratory does the following:

- Provides thin (typically submicron) membranes of polyimide, Formvar, Si, and Si₃N₄ for filters and vacuum barriers, and as support structures.
- Deposits thin metal layers for applications such as photocathodes, neutron-activation equipment, x-ray filters, and fiber optics for x-ray streak-camera fiducials.
- Fabricates resolution patterns for x-ray streak cameras (typically with a minimum size of 50 μ m in either 1- to 5- μ m-thick gold or 5- to 20- μ m-thick polyimide).

- Produces gold grating and zone-plate structures for a variety of x-ray diagnostic techniques such as spectroscopy and zone-plate coded imaging.

- Deposits grating and zone-plate patterns on x-ray multilayer mirrors (see Fig. 4-54). These can be used as output couplers for an x-ray laser.

- Produces x-ray beam splitters on thin layers (300 to 1500 Å) of Si₃N₄ to serve as output couplers for an x-ray laser.

Next year's program calls for fabrication of x-ray multilayer mirrors and their coupling with submicron structures to produce 3-D structures. To this end, we have designed and started to assemble a multihead magnetron sputtering system that will enable us to produce these mirrors on flat and curved surfaces. The sputtering system will employ a closed-loop monitor to guarantee the accuracy of the

Fig. 4-53. Polyimide, 2.5 μ m thick, after reactive-ion etching in an oxygen plasma. 500-Å chromium was used as a mask. The minimum line width is 5 μ m.

Fig. 4-54. (a) A 6- μ m-period, 1000-Å-thick, gold diffraction grating on an x-ray multilayer mirror. (b) Off-axis, gold, Fresnel zone plate, 1000 Å thick, on an x-ray multilayer mirror. The minimum line width is 2.5 μ m.

layer thickness. Our x-ray lithography system—the other major component for fabrication of 3-D submicron structures—is now in operation. We also plan further characterization of our RIE system, which will allow increased flexibility in fabricating these 3-D structures. In 1986, the microfabrication laboratory will play an expanding role in filling the LEAD Program's x-ray and optical diagnostic needs.

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Time-Resolved Measurements

The X-Ray Framing Camera. This article reports our progress in developing an ultrafast x-ray framing camera to be used in laser fusion experiments on Nova. The camera will provide high-fidelity images with frame durations less than 100 ps. The first prototype camera will generate a single image approximately 1 cm^2 , with a spatial resolution of $30 \mu\text{m}$ at the image plane. Coupling the framing camera to our $22 \times$ x-ray microscope will provide resolution of $\sim 2 \mu\text{m}$ at the source, a value limited by the resolution of the microscope. The camera will be optically triggered by a laser pulse that is synchronous with the Nova driver beams.

The schematic design of the framing camera is shown in Fig. 4-55. The principles of its operation have been described previously.⁷⁰ The camera is a highly specialized suspended-strip transmission line.

Its width of 25 mm and gap width of 1 mm yield a characteristic impedance of 15Ω . A photoconductive switch is directly incorporated into the transmission line. The switch electrically isolates a pulse-shaping charge line section that is biased to about -10 kV .

A -5 kV voltage pulse of less than 100-ps duration is generated by illuminating the photoconductive switch with a short pulse (10 ps) of intense light. The high-voltage pulse, propagating along the transmission line, traverses an x-ray photocathode and gates the photocathode "on" for the pulse duration. During this time, a photoelectron replica of the x-ray image is accelerated across the transmission-line gap.

A passive microchannel plate, mounted in the ground plane directly opposite the photocathode, acts as a direction-sensitive filter. Because its channels are oriented at an angle of approximately 4° from the normal, and the x rays impinge upon the semitransparent photocathode at near normal incidence, the plate is able to block the large number of x rays that pass through the photocathode. The photoelectrons, however, are deflected the requisite 4° by the magnetic field associated with the electromagnetic wave in the transmission line, and they pass through the microchannel plate. The gated photoelectron image is directly recorded by a thinned charge-coupled device (CCD) situated 0.5 mm from the back surface of the microchannel plate.

Although the framing camera is conceptually quite simple, the unique constraints of its design pose technical challenges related particularly to its state-of-the-art

Fig. 4-55. Design of the x-ray framing camera.

components: photoconductive switches, microchannel plates, and CCDs. Consequently, in developing the framing camera, we have gained expertise in a variety of important technologies.

Fabrication of Framing-Camera

Components. One of the project's major challenges was to fabricate the camera's transmission-line assembly. The difficulty lies in building the photoconductive switch, the photocathode, and the microchannel plate directly into the transmission line without degrading its electrical performance.

We constructed a transmission line composed of a conducting strip suspended over a gold-plated, copper ground plane. The strip is made of a machinable ceramic (MACOR) overlaid with gold and resistor films. It is fabricated by thick-film techniques. The ceramic is painted with Dupont No. 5715 gold ink and then fired at high temperature. An advantage of this technique is that the gold layer can be deposited continuously around sharp corners. A band of Dupont No. 1613 resistive ink is deposited at one end of the strip to serve as a termination resistor. After firing, the resistor is ground to a thickness that yields a resistance of exactly $15\ \Omega$. The final thickness of the gold and resistor films is approximately $25\ \mu\text{m}$. These films adhere excellently to the MACOR substrate.

The suspended strip is composed of two sections that are precisely aligned. The break between the sections occurs at the site of the photocathode. The photocathode is constructed from a beryllium ring 18 mm in diameter with a 2-mm wall thickness, capped on one end by $50\text{-}\mu\text{m}$ -thick beryllium foil. The foil is coated

by thermal evaporation with $1000\ \text{\AA}$ of cesium iodide. The photocathode is mounted into the transmission line by a compression fit between the two ceramic sections. This modular design makes it possible to replace the photocathode quickly and to vary the gate pulse width by interchanging transmission line sections having charge lines of different lengths.

The microchannel plate is mounted in the copper ground plane directly opposite the photocathode (see Fig. 4-55). The custom-made, rimless plate has a diameter of 25 mm and is 0.24 mm thick—about one-half as thick as standard microchannel plates. Its channels are $12\ \mu\text{m}$ in diameter, and their centers are separated by $15\ \mu\text{m}$ in a close-packed hexagonal array. The plate is cut so that the channels are biased at approximately 4° from the normal. Both surfaces of the microchannel plate are coated with a standard nichrome layer. In addition, the front surface is coated with a $1\text{-}\mu\text{m}$ -thick layer of gold followed by a $1\text{-}\mu\text{m}$ -thick layer of beryllium. The gold and beryllium layers are designed to provide the excellent conductivity required of the ground plane and at the same time behave as inefficient photoelectron emitters.

The photoconductive switch presented another challenge in the development of the framing camera. Building this switch into the transmission line is complicated by two important requirements. The switch must be mounted so that the optical trigger pulse can be efficiently coupled to its active area. The switch must also not perturb the electrical properties of the transmission line. To satisfy these requirements, we fabricated photoconductive switches as sketched in Fig. 4-56(a). The

Fig. 4-56. (a) Design of the silicon-on-sapphire photoconductive switch. (b) Incorporation of the photoconductive switch into the transmission line.

switch is composed of a layer of semi-insulating crystalline silicon approximately $1\text{ }\mu\text{m}$ thick, grown on a transparent sapphire substrate. Silicon films are commercially available on sapphire wafers 3 in. in diameter and 0.046 mm thick. A wafer is cut to yield approximately ten rectangular switches 7 mm wide and 25 mm long. The long edges are ground to a 45° bevel, giving the switch a trapezoidal cross section. Gold contacts are then deposited along the long edges, and a strip, either 1.5 or 2.5 mm wide, down the middle of the switch is left to serve as the gap.

The photoconductive switch is built into the suspended transmission line as shown in Fig. 4-56(b). The switch is mounted in a rectangular groove so that the silicon film is flush with the conducting surface of the transmission line. The grooves created by the beveled edges of the switch are filled with conductive epoxy, assuring good electrical contact with the transmission line.

Table 4-15. Operating parameters of the RCA SID501 CCD.

Pixel array size	512 $\times$ 320
Pixel size	30 $\times$ 30 μm
Total image area	15.36 $\times$ 9.6 mm
Output noise	<150 e^-
Dark current	$\sim 1.5\text{ nA}$
Full well capacity	600 000 electrons
Dead layer	<1 keV (estimated)

Fig. 4-57. CCD header assembly.

A fiber-optic light guide is inserted into a slot directly behind the photoconductive switch. The spot-to-line light guide is composed of step index fibers that are randomized so that the output illumination pattern is spatially uniform. Light is coupled onto the active area of the photoconductive switch through the transparent sapphire substrate.

As the detector for the framing camera, we selected the RCA SID501 CCD. Several of its most important operating parameters are listed in Table 4-15. The CCD is thinned to $8\text{ }\mu\text{m}$, making it extremely sensitive to electron bombardment at its back surface (i.e., the surface opposite the gate structure). The combination of a minimal dead layer and a low dark current at room temperature should provide sensitivity adequate to detect individual 5-keV electrons.

The thinned CCD was mounted by RCA on a custom-made header, as shown in Fig. 4-57. The CCD is epoxied to a glass frame, which in turn is epoxied to a thin Kovar plate. A rectangular aperture in the Kovar plate provides direct access to the back side of the CCD. The Kovar plate is attached to a ceramic support plate, and electrical connections to the CCD are made via wires bonded to aluminum strips on the ceramic plate. A push-pin connector links the controller cable to the ceramic plate. The essential feature of the header is that it facilitates access to the back surface of the CCD. The header is inserted into a rectangular pocket machined in the ground plane behind the microchannel plate. When the Kovar plate is in contact with the back surface of the microchannel plate, the CCD and the microchannel plate are separated by only 0.5 mm. Such proximity is essential for achieving high spatial resolution in the photoelectron image.

The CCD is controlled by a Model V camera built by Princeton Scientific Instruments. The camera is customized to operate the CCD at a rate of approximately 500 000 pixels/s while maintaining a digital pixel-by-pixel readout. The fast readout limits the accumulation of dark current and allows the CCD to be operated at room temperature. In addition, the software and control system (including an IBM AT microcomputer) that support the

Model V camera provide many useful image-processing capabilities.

Characterization of Framing-Camera

Components. Characterization of the framing camera involves two distinct phases. First the camera's components must be characterized individually so that the physics governing their performance is well understood. This phase is particularly important because many of the physical mechanisms on which the operation of the instrument relies have not been previously investigated. The second stage of characterization involves the assembly and testing of the camera as a whole. We are now in the middle of the first stage—measuring the behavior and performance of the camera's components.

A crucial aspect of the framing camera's operation is the ability to generate and propagate high-voltage pulses of short duration and minimal spatial variation along the transmission line. The duration of the pulses determines the temporal resolution of the camera. The variation of the voltage along the transverse dimension of the transmission line affects the fidelity of the image. To understand the performance of the transmission line, it is necessary to map out the electromagnetic field inside the line both spatially and temporally on

a picosecond time scale. Mapping is achieved by using electro-optical sampling.

We built an electro-optical sampling system based on a technique originally developed by Valdmánis et al.⁷¹ A synchronously pumped, mode-locked dye laser is amplified to provide ~ 8 -ps pulses with energy of ~ 1 mJ/pulse at a wavelength of $0.61 \mu\text{m}$ and a repetition rate of 10 Hz. Most of the laser pulse is coupled into the fiber-optic light guide and used to activate the photoconductive switch. The light intensity inside the light guide can be as high as 10^9 W/cm^2 . At this intensity, we measured a throughput efficiency of 50% and no appreciable temporal broadening of the pulse.

The charge line is pulse-biased to 2 kV—the maximum voltage it can sustain in air without electrical breakdown. When the photoconductive switch is illuminated by the laser, a ~ 1 -kV pulse of short duration is generated and propagates along the transmission line. A small fraction of the pulse is split off and used to sample the electric field inside the transmission line. This splitting is achieved by linearly polarizing the light and focusing it through a small lithium tantalate crystal suspended in the transmission line gap. The polarization rotates by an amount proportional to

Fig. 4-58. Two-dimensional spatial scan of the high-voltage gate pulse inside the transmission line. The scan was produced by the electro-optical sampling system.

the instantaneous electric field strength inside the birefringent crystal. The rotation is proportional to the sine of the electric field strength. (The half-wave voltage of the lithium tantalate crystal is 2200 V.) The rotation is analyzed by measuring the individual components of the polarization with two diodes. The output signals from the diodes are then subtracted and sent to a gated integrator, which averages the data over approximately 100 laser pulses to optimize the signal-to-noise ratio. The averaged results are collected and reduced in real time on an IBM XT microcomputer. Temporal scans are made by varying the relative delay between the firing of the photoconductive switch and the probing of the lithium tantalate crystal.

The entire transmission line is mounted on moving x and y stages. Spatial scans are made by translating the transmission line about the fixed position of the crystal. The temporal resolution of the sampling system, determined by the convolution of the laser pulse width and the transit time through the crystal, was 11 ps. The spatial resolution, 1.0 mm, is determined by the size of the crystal.

We used this electro-optical sampling system to study the behavior of the framing camera's transmission line, complete with the in-line photocathode and microchannel plate. Figure 4-58 shows a two-dimensional spatial scan of the voltage pulse at the instant the pulse is centered at the photocathode. The electric field is strongly confined to the interior region of the transmission line. It is reasonably uniform across the transverse dimension of the waveguide, but exhibits variations of 10 to 20%. Temporal scans of the voltage pulse are shown in Fig. 4-59. These scans were measured at the center of the transmission line at various positions downstream from the photoconductive switch. In this figure, it is apparent that the pulse propagates with minimal attenuation and dispersion. The scans demonstrate that the photocathode and microchannel plate do not adversely affect the electrical properties of the transmission line.

An unexpected feature observed in our study of the transmission line was significant ringing in the voltage pulse at a characteristic frequency of approximately 40 GHz. We attribute the ringing to LC

Fig. 4-59. Temporal scans of the high-voltage gate pulse measured at several positions along the transmission line. The positions refer to the distance from the photoconductive switch. The time axes are relative.

oscillations that arise from the parallel inductance L and capacitance C of the photoconductive switch. This behavior can be explained in terms of the simple lumped-circuit diagram shown in Fig. 4-60(a). The capacitance, inductance, and resistance $R(t)$ of the photoconductive switch are connected in series to the transmission line impedance Z . Initially the resistance of the switch is infinite, and a bias voltage V_b is applied to the charge line. If the resistance of the switch vanishes instantaneously at time $t = 0$, then the output voltage V_o is given by

$$V_o = \frac{V_b}{2} \left[1 - e^{-\gamma t} \left(\cos \delta t + \frac{\gamma}{\delta} \sin \delta t \right) \right], \quad (43)$$

where

$$\gamma = \frac{1}{4ZC}$$

and

$$\delta = \left[\frac{1}{LC} - \gamma^2 \right]^{1/2}.$$

Figure 4-60(b) shows the oscillations in the voltage pulse predicted by Eq. (43) using the values of $C = 0.6$ pF and $L = 35$ pH, which best reproduce the observed behavior. Physically, the ringing arises because the capacitance of the photoconductive switch prevents the initial distribution of charge on the charge line from being uniform. The effect of the oscillations on the performance of the framing camera, however, is not serious. The trajectory of any particular photoelectron is determined by the electromagnetic force on the electron, integrated over its ~ 40 -ps transit time across the transmission line gap. Computer calculations show that the integrating effect of the transit time significantly suppresses the effect of the voltage pulse oscillations on the photoelectron trajectories.

In another characterization experiment, we determined the x-ray throughput of the microchannel plate. Successful operation of the framing camera requires that the microchannel plate perform as an extremely efficient direction-sensitive filter. Photoelectrons with trajectories at 4° from the normal must pass through the

microchannel plate, while the ungated x-ray signal at normal incidence is strongly attenuated. Our x-ray throughput experiment is shown schematically in Fig. 4-61.

The microchannel plate was mounted inside a vacuum chamber, in a manipulator that allowed rotation about an axis ϕ normal to the surface of the plate and about an axis θ in the plane of the plate. Collimated x rays struck the plate's front surface, and a proportional gas counter directly behind the plate measured the throughput. The entire manipulator could be moved out of the x-ray beam so that the incident x-ray flux could be measured and absolute throughput values obtained.

The throughput at x-ray energies of 2.0 and 4.5 keV is shown in Fig. 4-62 as a function of the angles ϕ and θ . It is apparent in Fig. 4-62(a) that the throughput has a relatively broad maximum as a function of ϕ . It is also evident that the throughput away from the maximum is greater at lower energies, probably because of critical-angle reflectance from the walls of the microchannels. The x-ray throughput as a function of θ is shown on a

Fig. 4-60. (a) A lumped-circuit representation of the photoconductive switch within the transmission line. (b) Oscillatory behavior of the voltage as predicted by the model.

Fig. 4-61. Experimental configuration for the microchannel plate throughput measurements.

Fig. 4-62. X-ray transmission through the microchannel plate as a function of the plate's orientation.

logarithmic scale in Fig. 4-62(b). Transmission reaches a maximum at 3.8° from the normal, corresponding to the bias angle of the channels. The maximum observed throughput of 50% is close to the theoretical value of 55%, calculated from the ratio of open surface area to total surface area. X-ray throughput decreases strongly at angles away from the maximum such that attenuation is greater than four orders of magnitude for x rays at normal incidence. These results indicate that a microchannel plate can indeed perform as an effective x-ray attenuator in the framing camera.

Other experiments in progress will complete the characterization of the framing-camera components. We are measuring photoelectron throughput of the microchannel plate and photoelectric

emission due to absorption of x rays by the microchannel plate. Photoelectrons emitted from the microchannel plate are a potential source of noise and must be suppressed. We are also setting up a facility to thoroughly characterize the response of the RCA SID501 CCD to direct electron bombardment. A Nanolab 2100 scanning electron microscope is used as the source of electrons ranging in energy from 500 to 30 000 eV. The electron beam is gated with a high-speed beam blarker, so that a single line can be written onto the CCD at fluxes as low as one electron per pixel. We intend to measure the line response of the CCD as a function of electron flux and energy. It is also important to characterize the sensitivity, linearity, and dynamic range of the CCD under electron bombardment.

Characterizing the components of the x-ray framing camera has yielded much important information. When this task is completed in the next year, we plan to assemble and test the prototype camera as a unit.

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The Gated Microchannel-Plate X-Ray Pinhole-Camera Detector. The success of the gated microchannel-plate grazing-incidence spectrometer, combined with advances in electronic pulse generation, has allowed us to build a reentrant gated x-ray detector. Although not as fast as the framing camera described in the preceding article, this detector is electronically triggered and so does not need a laser. Its compact design allows it to approach to within 12 in. of the target in Nova. Consequently, simple instruments such as an x-ray pinhole camera or a crystal spectrometer can be used in front of the detector.

The detector itself is a microchannel plate with an active area of 40×15 mm, sandwiched between two thin gold coatings. It serves as the dielectric for a 6.25- Ω transmission line. A matched feed is provided by connecting the ends of the line to four parallel, 25- Ω , coaxial transmission lines.

X-ray photons cause photoemission from the inside of the detector's channels. The photoelectrons are amplified and then converted to light as in a conventional image intensifier. Time gating is achieved by applying a voltage pulse to the 6.25- Ω transmission line. The electron transit time in the channel is about 150 ps.⁷² The open time of the detector is the gate time minus the transit time.

To produce a satisfactory gating pulse close to the center of the Nova target chamber, we use a silicon diode as the switch of a charged transmission line. It has been known for some time^{73,74} that when silicon diodes are rapidly reverse-charged, they can avalanche with a fall time of a few hundred picoseconds. By charging the diode very quickly, voltage

transitions with a fall time of less than 100 ps have been achieved.⁷⁵

We are using the circuit shown in Fig. 4-63 to produce flat-topped voltage pulses. The 12.5- Ω transmission line shown in this figure is charged in a few nanoseconds to approximately 2 keV. The relatively slow time scale of this charging allows the microwave triode pulser to be outside the Nova target chamber, thus simplifying operation. The diode breaks down in an avalanche mode during the charging, to produce a 1.4-keV output voltage pulse for approximately twice the transit time of the charged line. The shunt inductor is tuned to give a square-topped voltage pulse. The 25- Ω output is then transformed to 6.25 Ω by a transmission line transformer and fed to the microchannel plate by the four 25- Ω coaxial cables. The output of this impedance transformer is shown in Fig. 4-64 on a sampling oscilloscope. The time jitter of the pulse is less than 50 ps, and the top of the pulse is flat to about 4% of the amplitude.

The diode assembly and the detector package are held by a cart that follows a rail system inside a tube. The rails allow the detector cart to be inserted into the target chamber through a load lock and to roll forward and stop close to the laser target. In its forward position, the cart

Fig. 4-63. Schematic of pulse former.

Fig. 4-64. Output from pulse former. Output impedance is 6 Ω , 100 V/div., 200 ps/div.

engages the dc power lines and the charging line for the diode.

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Directional Couplers. In 1985 we built and tested a high-speed directional-coupler traveling-wave modulator (DCTWM) that operates at a laser wavelength of $0.81\ \mu\text{m}$. The DCTWM converts high-speed electrical pulses to high-speed optical pulses, which can be recorded by high-bandwidth optical streak cameras. When developed to its potential bandwidth capability of $>10\ \text{GHz}$, this device will be versatile enough to make time-resolved measurements of all types of radiation that can be detected with small, fast, radiation-to-electrical signal detectors.

The impetus for this work lies in the fact that the time resolution of an optical streak camera recorder ($<10\ \text{ps}$) is far superior to that of the other components typically in a system for detecting, transmitting, and recording radiation. A direct process for converting radiation to light has been studied for many years by

others. This approach has not produced the speed and light intensity required for applications that demand superior time resolution. Our approach was to divide the objective into two parts: a radiation-to-electrical signal detector and an electrical signal-to-laser light converter. Both parts can be made to be extremely fast. The development of fast radiation-to-electrical signal detectors centers on photoconductive devices. Work on these is described earlier in this section in the article "Solid State X-Ray Detectors." The electrical-to-laser light step is described in this article. Research on both steps is a joint effort among several divisions at LLNL, aided by scientists from Bell Laboratories, Holmdel, New Jersey.

Electro-optic modulators have found considerable application in fiber-optical communication systems for high-speed processing.^{76,77} Their use has focused on the wavelength region around $1.30\ \mu\text{m}$ to minimize the attenuation of optical light in the fibers. Unfortunately, streak-camera photocathodes are extremely insensitive at that wavelength, and there is little hope for improvement over the next few years. Thus, we are forced to operate at $0.81\ \mu\text{m}$, the other wavelength region in which streak cameras have reasonable sensitivity and for which inexpensive laser diodes are readily available.

Figure 4-65(a) shows a schematic of a directional coupler. These devices are fabricated by diffusing $500\text{-}\text{\AA}$ -thick strips of titanium into a z-cut, y -propagating lithium niobate (LiNbO_3) crystal at 1025°C for 8 h in an atmosphere of flowing oxygen bubbled through water. After diffusion, a glass buffer layer, $2000\text{ }\text{\AA}$ thick, is sputtered to reduce electrode-loading loss for the TM polarization. The device is then annealed in flowing oxygen at 350°C .

A plan view of the titanium waveguides is shown in Fig. 4-65(b). Titanium increases the index of refraction of the guide in the LiNbO_3 material and forces the laser light to follow that path. If two waveguides are separated by a small distance d so that their evanescent fields overlap, then laser light entering one of the guides will completely transfer its energy to the adjacent guide along the coupling length L . This transfer can be deterred, however, by exploiting the electro-optic effect: making a

Fig. 4-65. Schematic of the DCTWM: 3-D and plan views.

slight change in the index of refraction in one guide with respect to the other by means of an applied electric field. The index change causes a phase mismatch between the two guides. Thus a fast electrical signal applied to the electrodes can control the amount of light transferred and can cause a modulated laser output that replicates the electrical pulse in time. This process has been shown to be very fast (>13 GHz) for modulators designed to operate at $1.30\text{ }\mu\text{m}$.

Without electrodes and for phase-matched conditions, the transfer efficiency from one guide to the other is⁷⁸

$$\eta = \sin^2(KL + 2\phi), \quad (44)$$

where K is the coupling coefficient, and L is the interactive length. The term 2ϕ represents the finite coupling in the bend region. Figure 4-66 shows the transfer efficiency as a function of L for couplers that we constructed in which the waveguide width W is $3.1\text{ }\mu\text{m}$ and d is $4.8\text{ }\mu\text{m}$.

The coupling length l is related to the coupling coefficient K by⁷⁹

$$l = \frac{2\pi}{K} \quad (45)$$

and

$$K = K_0 \exp\left(-\frac{d}{\gamma}\right), \quad (46)$$

where K_0 and γ are coefficients that depend on various waveguide parameters.⁸⁰ The coupling length vs d for several waveguide widths is shown in Fig. 4-67. As expected, the data are linear on a semi-logarithmic plot. The data from the $3.1\text{-}\mu\text{m}$ width are substantially lower than we anticipated based on the theoretical guidance from Ref. 78. We have yet to fully understand this result. It may have to do with the stoichiometry of the LiNbO_3 , which varies from batch to batch.

We did preliminary work in two other areas of directional coupler development. We measured the susceptibility of LiNbO_3 to photorefractive damage by $0.81\text{-}\mu\text{m}$ light and concluded that we can operate safely above $300\text{ }\mu\text{W}$ of continuous laser power. For actual use, we want to operate at about 1 mW in the pulsed mode and

Fig. 4-66. Transfer efficiency as a function of L for $W = 3.1\text{ }\mu\text{m}$, $d = 4.8\text{ }\mu\text{m}$.

Fig. 4-67. Coupling length vs waveguide separation, d , for $W = 3.1, 3.5$, and $4.0\text{ }\mu\text{m}$.

now feel confident about being able to achieve that.

We used a network analyzer and an electrical pulser to measure the frequency response of the electrode structure. For a -3-dB bandwidth, the network analyzer measured 7 GHz . In pulsed mode, the response was 16 GHz . Measurements of actual laser-beam modulation using a fast electrical pulse are planned for near-term experiments.

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Spatial and Temporal Resolution of LLNL X-Ray Streak Cameras. In 1985 we characterized the spatial and temporal resolution of the LLNL model 2 x-ray streak cameras. The information we obtained is useful in determining the feasibility of any experiment that relies on a high degree of resolution in either space or time, or in both simultaneously.

An appropriate figure of merit for the spatial resolution of a camera is the contrast transfer function (CTF). A resolution grid with ideal square-wave transmission

Diagnostics

Table 4-16. Resolution characteristics of streak cameras.

Camera No.	Best horizontal focus			Best vertical focus		Magnification
	Focus (V)	Width (μm)	CTF, 50% (line pairs/mm)	Focus (V)	Width (μm)	
20	915	1400	2.4	555	231	1.19
21	820	733	4.2	575	124	1.20
22	849	514	4.3	582	139	1.22
23	760	663	1.8	620	364	1.20
24	670	436	4.9	605	340	1.24
25	855	560	3.7	560	163	1.24
26	905	620	3.0	615	300	1.16

Table 4-17. Timing characteristics of streak cameras.

Camera No.	Sweep speed (ps/mm)	Internal delay (ns)
20	103	18
21	118	22
22	127	22
23	123	15
24	123	15
25	124	19
26	122	16

contains regions of 100% transmission alternating with regions of zero transmission. For such a grid, the CTF is defined as

$$\text{CTF} = \frac{I_{\max} - I_{\min}}{I_{\max} + I_{\min}}, \quad (47)$$

where $I_{\max}$ and $I_{\min}$ are the maximum and minimum measured intensities.

For temporal resolution, a corresponding function is not so readily measured. Instead we measured the line spread function by obtaining the width of the image of a known photocathode slit without applying the sweep voltage.

To determine the CTF for a streak camera, we measured the image created by an x-ray source incident on a resolution grid placed before the x-ray photocathode of the camera. To distinguish the electron optics contribution to the CTF from the contribution due to the camera's image intensifier, we separately characterized the intensifier (ITT model 4113) by forming an image with a light source incident on a similar resolution grid placed on the input face of the intensifier. Besides determining CTFs, we measured the internal delay, vertical resolution, horizontal magnification, and sweep speed for each camera. Tables 4-16 and 4-17 show these results.

All measurements were made on the LLNL Monojoule laser facility using a tantalum target. Sweep speeds were measured by using an etalon with a pulse

separation of 800 ps. Vertical resolution was measured by grounding the deflection plates of the cameras.

The resolution grid used for the CTF measurements was a 3- μm -thick, self-supporting, gold, linear zone plate with zone widths varying from 1 mm to 50 μm (0.5 to 10 line pairs/mm). The photocathodes we used consisted of $200 \pm 30 \text{ \AA}$ of aluminum evaporated on 1000 $\text{\AA}$ of Formvar. The photocathode slit in each case measured $100 \pm 10 \mu\text{m}$. Kodak Tri-X film was used as the recording medium.

To measure the CTF of the part of the streak camera system consisting of the image intensifier and the film, we removed the intensifier from the camera and placed it below a dim ($3 \times 10^{-9} \text{ W/cm}^2$) microscope illuminator. We placed a resolution grid similar to that used for the camera-CTF measurements on the input fiber-optic plate of the intensifier. The grid consisted of a linear zone-plate pattern of chromium deposited on glass. The zone plate had an initial clear zone 1 mm wide. Neutral-density filters placed between the grid and the light source produced a range of exposures on each photograph. The intensifier was gated on for about 100 μs . We used both Tri-X and RXP film in a wide range of exposures.

The resulting images for the intensifier and streak cameras were digitized on the photodigitizer-microdensitometer in our Laser Experiment Analysis Facility (LEAF). Figures 4-68 and 4-69 show the CTFs for the cameras and for the intensifier, respectively, as functions of spatial frequency. The electron optics of the cameras introduces an astigmatism into the focusing; thus the best horizontal (spatial) and vertical (temporal) focusing occur with different voltages on the focus cone, as shown in Table 4-16. The camera CTFs indicate

Fig. 4-68. CTFs for representative LLNL streak cameras.

that three of the streak cameras approach what might be an optimum spatial resolution for their design. Improving the electron optics would result in the camera's being limited by the resolution of the intensifier. Experiments requiring good spatial and temporal resolution simultaneously must use a compromise focus setting that is not optimum for either.

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Optical Fiducial for X-Ray and Optical Streak Cameras. To characterize the optical fiducial system for the LEAD Program's

x-ray and optical streak cameras, we tested the system's components at the Monojoule laser facility. Our tests determined pulse dispersion, fiber-damage thresholds, and the energy required to produce photoelectrons in x-ray streak cameras.

Our dispersion measurements showed that a 30-ps, 3ω input pulse through 20 m of 100- μ m-diam fused silica fiber broadens to ~ 130 ps. The FWHM of this pulse proved to be repeatable and could be used as a fiducial with an accuracy of ± 25 ps.

In a test to determine the energy threshold for producing photoelectrons, we passed a 30-ps pulse of 3ω laser light

Fig. 4-69. CTF for an ITT Image Intensifier. The two magnifications correspond to resolution at the input to the intensifier (1.0) and to resolution at the input to the streak camera (1.2).

Fig. 4-70. Relative fiber-optic cathode response with 3ω light.

through 4 m of fiber. A plot of the relative fiducial exposure vs the fiber output energy reveals a threshold energy of $\sim 0.05 \mu\text{J}$ (see Fig. 4-70). The nonlinearity of the plot indicates a multiphoton electron-production process. On Nova, a 100-ps pulse would pass through 20 m of fiber, in contrast to the 30-ps pulse and 4-m fiber used on the test facility. Because photoelectron production appears to be intensity-driven, the longer pulse and fiber length were expected to increase the energy threshold to $0.27 \mu\text{J}$. When this threshold requirement and a 41.5-dB loss from the fiber-optic distribution system⁸¹

are taken into account, the input energy required at the fiber-optic array face amounts to 4.0 mJ (0.37 J/cm^2)—well within the capability of the Nova fiducial laser beam.

Unfortunately, the high intensity required with 3ω light damaged the face of the fiber array when the fiducial system was used on Nova. The damage caused the output energy at the streak-camera (cathode) end of the fiber bundle to dip well below the required threshold for photoelectron production.

Instead of delaying the program to rebuild the fiber array, we decided to reconfigure the system to use 2ω light. To produce photoelectrons at this wavelength, we had to coat the fiber bundle with a thin layer of aluminum at the cathode end. Despite this extra step, 2ω light offered the following advantages over 3ω :

- Its transmission in the fibers is superior. We could reduce intensity at the array face, yet avoid the large transmission losses experienced with 3ω light, which amounted to 80% through 20 m of fiber.
- Pulse dispersion in the fiber was reduced.
- It is visible; thus the fiducial laser was easier to align with the fiber array face.

We recognized that electron production with 2ω light would depend on a multiphoton process. Previously acquired evidence of such a process (Fig. 4-70) gave us confidence in our approach. Despite the fact that Nova's fiducial laser beam was not optimized for 2ω production, it produced enough energy to yield a sufficient number of photoelectrons.

With $1 \mu\text{J}$ of 2ω light in fibers coated with 100 to $145 \text{ \AA}$ of aluminum at the cathode end, we produced an abundance of photoelectrons on Nova. The signal was observed to increase nonlinearly as the 2ω energy in the fiber was increased. The threshold for electron production was about $1 \mu\text{J}$ with either $100\text{-\AA}$ or $145\text{-\AA}$ aluminum coating. These energy measurements, made with a diode on a reference fiber in the bundle, are accurate to about $\pm 50\%$.

Figure 4-71 shows fiducial data recorded at Nova on an x-ray streak camera equipped with a $3 \times$ card that has a sweep speed of $\sim 125 \text{ ps/mm}$. Nova's fiducial

pulse starts out at 1ω with an FWHM of 85 ps. It is compressed as it passes through frequency conversion crystals to 2ω and then disperses as it is transported through the fibers. After the pulse travels through 20 m of fiber to reach the x-ray streak camera, we observed that its width expands to 139 ps FWHM (Fig. 4-72). Figure 4-73 shows the 58-ps FWHM photoelectron pulse produced by the aluminum-coated fibers in the x-ray streak camera. The pulse compression from 139 ps of 2ω light to about 60 ps of photoelectrons is consistent with nonlinear photoelectron production. The short duration of the photoelectron signal makes it an ideal fiducial for x-ray streak cameras on the two Nova target chambers.

The $\sim 1 \mu\text{J}$ of 2ω light required in a 100- μm -diam fiber leads to energy density of $\sim 0.013 \text{ J/cm}^2$ on both ends of the fiber bundle. This energy density is well inside the damage threshold for both the uncoated end and the aluminum-coated end—estimated to be ~ 0.1 to 0.2 J/cm^2 (Ref. 82). During the course of about 40 shots at comparable or slightly higher energy densities, no damage could be detected. A few shots at $10 \mu\text{J}$ did damage the coating on one fiber.

This energy-density requirement implies that for a bundle of fibers 1 mm in diameter,

0.1 mJ of 2ω light is needed. The Nova fiducial pulse amplifiers have routinely produced 20-mJ energies of 1ω light, and conversion efficiencies up to 30% are attainable. (See the article "Optical Fiducial System" in Sec. 5 for more information on the fiducial system's capabilities.) The use

Fig. 4-71. X-ray streak camera fiducial record.

Fig. 4-72. 2ω optical streak camera record after the pulse has traveled through 20 m of fiber.

Fig. 4-73. X-ray streak camera record after the pulse has traveled through 20 m of fiber.

of 2ω light for streak camera fiducials should enable the fiducial beam to be run well within its design limits.

We believe that photoelectrons are produced in the aluminum (which has a work function of 4.28 eV) because of the metal's interaction with two or more 2ω photons ($h\nu = 2.35$ eV). The nonlinear behavior of the photoelectron signal, also observed with 3ω light, is consistent with such a multiphoton process. This process would include the effects of an aluminum oxide layer about $50 \text{ \AA}$ thick, which is expected to be on the aluminum surface because no precautions were taken to prevent oxidation by exposure to the atmosphere.

Obviously, more work needs to be done to understand the photoelectron production process. But the fiducial system works with 2ω light. In the near future, all Nova x-ray and optical streak cameras will be equipped with a fiducial so that data from each camera can be compared temporally within ± 50 ps.

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CCD Readout for Optical Streak

Cameras. In the ICF Program, streak cameras are used as the principal diagnostic instrument to monitor fast optical and x-ray phenomena associated with laser-plasma experiments.^{83,84} The visible output from the streak camera is recorded either on photographic film or by a CCD camera. Film is used when accuracy is more important than fast turnaround time, as in the case of laser-target diagnostics. For these purposes, film provides good

resolution, high sensitivity, and a large dynamic range. However, for applications requiring fast information turnaround, such as the adjustment of a laser pulse shape, the time required to process and digitize the film is prohibitive.

For several years, an LLNL-designed streak-camera readout system^{85,86} was used when fast information turnaround was important, but this system sacrificed image resolution and dynamic range when compared to photographic film. With the development of our new streak-camera readout system, resolution and dynamic range approach that of film. As a result, electronic readout systems can be used in more applications in the future.

System Overview. The block diagram of Fig. 4-74 shows the overall streak-camera recording system. The system includes an LLNL-designed streak camera and a readout system consisting of a CCD camera and a computer-based camera controller. In a typical streak camera diagnostic application, a transient optical or x-ray signal is focused onto the input slit of the streak camera. The camera converts this signal to a two-dimensional output image on which time is displayed vertically and spatial or spectral information is displayed horizontally. This visible image is then relayed through a lens to the CCD camera where it is captured by the CCD sensor in a matrix array of photoelements (pixels). This camera digitizes the image and transfers it to the camera controller for local display, image processing, analysis, and storage.

CCD Camera. Our new streak-camera readout system^{87,88} is based on a CCD camera obtained from Photometrics, Ltd., Tucson, Arizona. This camera incorporates state-of-the-art technology to provide analog-to-digital conversion (ADC) and cooling for the CCD. The CCD is an almost ideal sensing device for streak-camera readout. Its matrix array of pixels provides a readout sensor that is small, geometrically stable, and inherently linear to light input. The CCD we have chosen (Thomson-CSF, TH-7861) is a silicon array composed of 576 rows by 384 columns of charge-coupled pixels. Each pixel is $23 \mu\text{m}$ square, resulting in an image area of $13.2 \times 8.8 \text{ mm}$.

The CCD works by absorbing visible-light photons that create hole-electron

Streak camera readout system

Fig. 4-74. Streak-camera recording system consisting of an LLNL-designed streak camera and a readout system.

Fig. 4-75. The LLNL streak camera with CCD camera attached.

pairs in the silicon. The electrons are trapped in potential wells (pixels) by gate electrodes on the surface of the CCD. The light from the phosphor of the streak camera lasts less than 100 μ s. During this time, a charge image corresponding to the streak camera output is built up in the CCD, thereby capturing the light image. Then, the charge in each pixel is shifted to a single output by clocking electrode voltages. The capacity of a pixel is 250 000 electrons. A three-stage thermoelectric cooler is used to cool the CCD to -30°C , reducing the CCD's noise level to about 50 electrons (rms). The resultant theoretical dynamic range (pixel capacity divided by rms noise) is 5000. The CCD camera is compatible with many different CCD types. Its flexibility will accommodate evolving CCD technology.

Streak-Camera-CCD Coupling. The streak camera's output image can be coupled to the CCD sensor with fiber optics or relayed by a lens. Because the CCD in our older streak-camera readout system was not cooled, maximum coupling efficiency was required. For this reason, we used fiber-optic coupling. By cooling the CCD in the new system, we greatly reduce its noise and increase its sensitivity. Improved sensitivity makes it possible to couple the streak camera's optical output to the CCD readout system by means of a lens. This approach offers ease of maintenance, mechanical reliability, and the possibility of readily coupling the CCD system to other streak cameras.

System Configuration. The mechanical design of this new streak-camera readout system permits adjustment of image size

and focusing in a very stable light-tight mounting configuration. The streak camera's output is relayed to the CCD array through a 56-mm, $f/1.2$, finite-conjugate lens. This lens reduces the image from the output intensifier of the streak camera by about 3:1 to match the size of the CCD. Figure 4-75 shows the overall construction of the combined streak camera and CCD camera.

The camera readout system can be configured either to stand alone or to link with the Nova laser system by means of a front-end processor. Figure 4-76 shows a block diagram of the system as used in the stand-alone configuration. Components such as the direct memory access (DMA) interface, the parallel interface, and the central processing unit (CPU) are interconnected by means of a Digital Equipment Corporation (DEC) Q-Buss. Commands

Fig. 4-76. Block diagram of the camera control system.

and status signals are transferred to and from the camera controller through the parallel interface.

Image Processing. To capture a single image, the system receives a command to enter integration mode, in which the CCD is continually being cleared while waiting for a trigger signal. When the trigger arrives, the CCD camera integrates the image from the streak camera's image-intensifier output. For the streak camera's output image to be captured, the streak-camera trigger must arrive while the CCD is in integration mode. The image is then digitized in the camera controller with a precision of 14 bits (16 384 ADC units full scale) by using a technique known as "correlated double sampling" to reduce noise. The data are then transferred by DMA into video memory. This configuration is capable of capturing an image and transferring it at a rate of 50 000 pixels per second. The captured image can be viewed immediately on a color monitor through the analog processor module. The video memory and analog processor (made by Imaging Technology, Woburn, Massachusetts) provide color enhancement, zoom, pan, and scroll features. The captured image can be manipulated while resident in video memory and can be transferred to disk for archiving.

The custom interface, required for proper interconnection and signal levels between the camera controller and the computer system, is the only component in this imaging system that is not commercially available. The interface is a convenient place for test points and monitor lights to indicate system status for troubleshooting.

System Configuration on Nova. On Nova, the readout system is operated in a front-end processor configuration. In the diagram of this configuration (Fig. 4-76), the Novanet interface replaces the disk and console. The software to operate the readout system is downloaded via Novanet⁸⁹ from the Nova VAX-11/780 mini-computer. This configuration reduces the cost of the system by eliminating the console and the attached mass-storage device. During a laser experiment, the streak camera's output image is captured in the video memory and then transferred, via Novanet, to the Nova VAX for data analysis.

Related Software. Software routines operate the readout system in two modes: a single-shot mode for streak-camera readout and a repetitive mode for immediate feedback during relay-lens focusing and adjustment of image size. Besides the software devoted to system operation, a wide range of software was written to perform

the image processing and analysis necessary to recover the original streak camera intensity-vs-time data. For image recovery, two CCD images must be captured and processed. The processing involves subtracting a background (dark-field) image from the signal image. The background image is very stable because of CCD cooling and need not be acquired frequently; once a day may be often enough. It must be acquired while the CCD is in total darkness and must have the same integration (exposure) time as the signal image. This background subtraction process removes the CCD-pixel dark current and yields an image containing only the data of interest.

Performance Characteristics. The critical performance characteristics of a streak-camera readout system are dynamic range, sensitivity, and resolution.⁹⁰ Dynamic range is the ratio of the maximum signal that can be recorded to the minimum signal that can be detected (also known as sensitivity). The important feature of our new readout system is that its dynamic range exceeds that of the streak camera. The dynamic range plot in Fig. 4-77 shows data for the new CCD camera-streak-camera combination.

The signal at the input to the streak camera is frequently very weak, but the readout system is sensitive enough to detect individual photoelectrons generated at the streak camera's photocathode. Therefore, the sensitivity of the system is limited only by that of the photocathode.

The temporal and spatial resolution of the streak camera will be degraded if the readout system has inadequate spatial resolution. An ideal readout would have a resolution much higher than that of the streak camera (including the intensifier). It is in this area that electronic readouts are likely to be the limiting factor because the number of pixels is limited by current technology to about 500×500 . (Higher-resolution CCD's are available but very costly.) The resolution of an optical system is defined by the modulation transfer function (MTF). The MTF of the CCD camera alone and that of the streak-camera-CCD-camera combination is shown in Fig. 4-78. The CCD camera has a resolution that is better than that of the streak camera alone. But it is still not as good as film

Fig. 4-77. Dynamic-range plot.

Fig. 4-78. Spatial-resolution plot.

and produces a noticeable degradation in the overall system resolution. This degradation is largely due to lens demagnification and could be reduced or eliminated with a more direct coupling scheme.

Application. Our new streak-camera readout system is being used to monitor the temporal profiles of shaped pulses from the master oscillator on Nova. Laser pulses with a square shape are often desired for experiments. To produce the desired shape, master-oscillator pulses injected into Nova's amplifier chain must be shaped in such a way as to compensate for gain saturation. Pulse shaping simplifies the experimental analysis of laser-target diagnostics. The need for a streak-camera readout system that provides a large dynamic range and good linearity is important in this application. A typical streak-camera image of the shaped master-oscillator input to the amplifier chain (as recorded by the new CCD system) is shown in Fig. 4-79(a). The high-energy output of a

Fig. 4-79. (a) Single lineout of the shaped Nova master oscillator as captured by the streak-camera readout system; (b) the high-energy output of a Nova laser beam.

Nova laser beam is shown in Fig. 4-79(b). Note the low noise prior to the rise of the signal, the wide dynamic range, and the obvious effects of gain saturation.

Summary. A continuing effort to develop instrumentation to monitor laser performance produced our second-generation streak-camera readout system. We are now using the first of these systems in the LEAD Program, and plan to install several more in the coming year. We intend to continue system characterization and eventually to replace all the older CCD readout systems with the new ones. We also plan to explore the possibility of replacing photographic film in applications employing x-ray streak cameras.

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Neutron Diagnostics

Neutron Spectroscopy. Neutron time-of-flight techniques can be used to measure fuel ion temperatures.^{91,92} This article describes how such measurements can be made on Nova by using new instrumentation techniques. These techniques also have the potential for yielding other types of diagnostic information, in particular, fuel areal density ($\langle \rho r \rangle$) and fusion burn width.

Technique. Thermal broadening of the neutron energy spectrum emitted during fusion burn causes a dependence of the observed neutron line width on ion temperature. If we assume a Maxwellian distribution, this relation is $\text{FWHM} \approx 176 \sqrt{\theta_i}$ keV for a 14-MeV neutron from a DT plasma, where θ_i is the ion temperature in keV. This width has been measured for a 45-m flight path with a relatively slow (4-ns FWHM) time-of-flight detector operating in "current mode."⁹¹ A detector that operates in current mode (for example, one consisting of a single large scintillator and a photomultiplier tube) creates a signal that depends on the neutron flux reaching the detector. Because of its non-uniform response to the neutrons, an unfolding technique may be needed to interpret the detector's signal.

For Nova, we propose that an array of detectors operating in "pulse mode" (single-particle counting) be used to make fuel ion temperature measurements. In such "single-hit" arrays, each detector element records the time of arrival of the first neutron and produces a signal. Pulse mode is the standard mode of operation for accelerator-based experiments, so associated timing techniques are well developed. System timing resolution on the order of 100 ps is routine, and even higher resolution has been achieved for some applications. The recent development of detectors and electronics for high-energy, nuclear, and medical physics makes possible the construction of large arrays with pulse-mode timing.

Increased resolution allows the length of the neutron flight path to be reduced. On Nova, for example, neutrons measured by a detector located at the chamber wall would have a 2.3-m flight path. If the

detector's time resolution were 100 ps (FWHM), it would have a satisfactory energy resolution of 64 keV (FWHM) for DT neutrons. Thus a single-hit array offers the advantage that construction of a long flight path and associated collimation are not required.

Another major reason for developing the single-hit array technique is its potential adaptation to other diagnostic measurements based on neutron spectroscopy, for example, fuel $\langle \rho r \rangle$ through measurement of fuel-scattered neutrons.

Detector. For measuring ion temperatures on Nova, silicon charged-particle detectors made neutron-sensitive by use of a polyethylene neutron-to-proton (n-p) converter appear to be the best choice. In this type of device, neutrons elastically scatter protons from the polyethylene into the detector. These detectors have been selected for the following reasons:

1. Ease of fabricating an array. It is possible to divide a single silicon wafer into many individual detector elements (called channels, hereafter) through use of standard integrated circuit technology. An example is the "microstrip" detector currently coming into common use in high-energy physics. Figure 4-80 shows a design for use on Nova with 192 channels or a single detector.

2. Excellent fast-timing characteristics. Silicon charged-particle detectors exhibit low noise, resulting in low electronic jitter.

3. Relative insensitivity to x-ray backgrounds. X-ray backgrounds can be a problem for any type of neutron detector, but silicon detectors are relatively insensitive to even suprathermal hard x rays that penetrate several centimeters of lead shielding. Calculations indicate that 5 cm of lead will adequately shield these detectors against even the hardest x-ray spectra observed on Nova. Experimental measurements are under way to characterize this potential problem.

Electronics. A standard single channel of electronics consisting of an amplifier, a discriminator, and a time-to-digital converter (TDC) is shown in Fig. 4-81. The amplifiers and discriminators are commercially available as integrated circuits (typically four channels to an IC). Plans

Fig. 4-80. Silicon microstrip charged-particle detector: 192 channels are produced on a single wafer, which is mounted in a printed circuit board.

call for mounting them on small printed circuit boards called chamber cards. The chamber cards, in turn, will be mounted near microstrip detectors in packages called front-end modules. Each front-end module will consist of a single microstrip detector and enough chamber cards to support the 192 channels. Timing signals are sent to remote TDCs through ribbon cables. TDCs with 100-ps timing resolution are commercially available and in routine use.

To achieve 100-ps timing, however, we must correct for the dependence of timing on pulse amplitude (walk). There are several standard means to do this. The one we are developing is a chamber card with a constant fraction discriminator (CFD). A CFD timing output is electronically corrected for walk.

System Design. Given a number of channels that each measures the arrival time of a neutron and is then "dead," we must determine the best channel configuration for measuring ion temperature. Also we must determine how the observed data should be interpreted. These calculations require definitions for the following quantities to be determined for each detector:

N_c is the number of channels in detector D_i ($N_c = 192$ for proposed detectors).

$\epsilon_i = \epsilon_i F_i$, where ϵ_i is the intrinsic detector efficiency and F_i is the single-channel solid angle.

N_{hi} is the number of hits (signal-producing neutrons) in a detector.

Fig. 4-81. Four front-end modules sharing a common bank of TDCs. Inset shows a single channel of electronics.

N_i is the number of channels fired in a detector (not to be confused with hits).

$$S \text{ (sensitivity)} = N_h/N_c = \epsilon_1 Y_n$$

Y_n is neutron yield.

$$P_i \text{ (the fraction of channels fired)} = N_i/N_c$$

P_p is the fraction of channels fired that received multiple hits (pile-up).

The distribution of neutrons is described by the binomial distribution:

$$P(H_1, N_h, N_c) = \frac{N_h!}{H_1!(N_h - H_1)!} \times \left(\frac{N_c - 1}{N_c} \right)^{N_h - H_1} \left(\frac{1}{N_c} \right)^{H_1}, \quad (48)$$

where H_1 is the number of hits in a channel and P is the fraction of channels with H_1 hits for a given N_h, N_c .

Equation (48) can be used, for example, to derive an expression for P_p , the fraction of channels fired ($H_1 \neq 0$):

$$P_i = 1 - P(H_1 = 0, N_h, N_c) = 1 - \left(1 - \frac{1}{N_c} \right)^{N_h} \quad (49)$$

or

$$P_i \approx 1 - e^{-S}. \quad (50)$$

Equation (50) can be shown to be good for values of N_h and N_c of interest and gives the intuitively expected result that the fraction of channels fired should depend only on N_h/N_c .

A similar derivation gives an expression for P_p , the fraction of channels fired that had multiple hits ($H_1 > 1$).

$$P_p = 1 - \left[\frac{P(H_1 = 1, N_h, N_c)}{P_i} \right] \quad (51)$$

or

$$P_p \approx 1 - \frac{Se^{-S}}{1 - e^{-S}}. \quad (52)$$

Further useful relations are:

$$S = -\ln(1 - P_i) \quad (53)$$

and

$$N_h = -N_c \ln(1 - P_i). \quad (54)$$

For a given detector, D_i , there is a range of values of S that should give useful data. (Note that S is determined by Y_n and by ϵ_1 , which can be varied by choice of the thickness of the polyethylene converter and solid angle of the detector.) For S too small, the number of channels fired (N_i) becomes small, and the statistical error in

determining the FWHM of the neutron line becomes large. For example, for $S \approx 0.1$ and $N_c = 192$ ($N_i \approx 20$), the statistical error on θ_i is about 25%. If S is too large, even in an ideal detector, all the channels would go dead after being hit by early neutrons, and none would be left to record the arrival of later neutrons. In reality, a channel of the microstrip detector is not dead upon the arrival of an additional neutron. Instead later neutrons distort the time observed for the initial neutron. This effect, called pile-up, sets a practical limit on how large S can be, as shown by Eq. (52).

For a CFD experiencing pile-up, the time of arrival observed becomes a time intermediate to the actual neutron arrival times. Large amounts of pile-up cause a narrowing of the observed neutron line width (early and late neutrons "average out" to the mean). Computer simulations show that this effect is completely negligible for $P_p \leq 20\%$ and not a major concern even up to $P_p \approx 50\%$. From Eq. (52) it can be seen that $P_p \approx 23\%$ for $S = 0.5$.

The requirement to work in a given S range (e.g., $0.1 < S < 0.5$) determines the dynamic range of neutron yields Y_n that can be covered for a given value of ϵ_i . Since the opposing requirements for a large statistical sample and a small degree of pile-up make it desirable to work in a small S range, the obvious way to increase dynamic range with respect to Y_n is to field several detectors with different values of ϵ_i . For economy, these detectors can be made to share the same TDCs (see Fig. 4-81). Appropriately delaying the more sensitive detectors gives the least sensitive detector "first chance" at TDC channels. Since this detector will probably have a relatively small number of hits, many channels would still be available for the next detector. This process would continue through the series of detectors, each more sensitive than the last.

This concept is clarified by Table 4-18, which gives the data expected for a range of Y_n values from a system of four front-end modules sharing a bank of common TDCs. In this example, the microstrip detectors, each with 192 channels, are located 2.3 m from the target, and the detector efficiencies are chosen such that D_1 , the most sensitive detector, has the maximum detection efficiency possible for

this technique. (The n-p converter thickness cannot be increased without limit, because the protons ultimately stop in the polyethylene. Consequently, the detectors have a maximum intrinsic n-p conversion efficiency of about 4×10^{-3} .) The remaining detectors have efficiencies such that $\epsilon_1(D_2) = \epsilon_1(D_1)/5$, $\epsilon_1(D_3) = \epsilon_1(D_2)/5$, and $\epsilon_1(D_4) = \epsilon_1(D_3)/5$. In Table 4-18, the value associated with D is the number of TDC channels remaining for detector i —effectively the number of channels in that detector. The corresponding values of N_i give the number of remaining channels that fire.

The configuration illustrated in Table 4-18 effectively spans several orders of magnitude in Y_n . If a Y_n value becomes typical for a given experimental series, sensitivities can be adjusted such that larger samples for that Y_n are obtained at the expense of dynamic range. A specific example would be the case of low neutron yields ($Y_n < 10^8$). All detector efficiencies could be increased to the maximum, improving N_i by nearly a factor of four over the N_{ii} values shown in Table 4-18.

Other Applications. The neutron time-of-flight measurement technique using arrays of single-hit channels might be a suitable means to obtain other diagnostic information. Of particular interest is the measurement of fuel $\langle \rho r \rangle$ by counting

Table 4-18. Example of expected data as a function of neutron yield.

Y_n	Quantity	i			
		4	3	2	1
10^8	D	192	192	191	188
	S	6.4×10^{-4}	3.2×10^{-3}	0.016	0.08
	P_i	6.4×10^{-4}	3.2×10^{-3}	0.016	0.08
	N_i	0	0	3	14
	P_p	—	—	8.0×10^{-3}	0.04
10^9	D	192	191	185	158
	S	6.4×10^{-3}	3.2×10^{-2}	0.16	0.8
	P_i	6.4×10^{-3}	0.03	0.15	0.55
	N_i	1	6	27	87
	P_p	—	0.02	0.08	0.35
10^{10}	D	192	180	131	26
	S	6.4×10^{-2}	0.32	1.60	8.0
	P_i	0.06	0.27	0.80	1.00
	N_i	12	49	105	26
	P_p	0.03	0.15	0.60	1.00
10^{11}	D	192	101	4	0
	S	6.4×10^{-1}	3.20	16.0	80.0
	P_i	0.47	0.96	1.00	—
	N_i	91	97	4	—
	P_p	0.28	0.87	1.00	—

fuel-scattered neutrons in DT plasmas.

There are two potential methods:

1. Spectral measurements below 14.1 MeV (Ref. 93). Full-energy neutrons can elastically scatter from deuterium or tritium nuclei and form a low-energy component in the neutron spectrum. The size of this component is dependent on fuel $\langle \rho r \rangle$.

2. Spectral measurements above 14.1 MeV (Ref. 94). Full-energy neutrons elastically scattering in the fuel produce energetic (several MeV) deuterium and tritium nuclei that can, in turn, undergo an energetic fusion reaction ("beam" interaction) and produce neutrons up to 30 MeV. This component of the spectrum is also dependent on $\langle \rho r \rangle$.

While both measurements are straightforward applications of time-of-flight techniques to different parts of the energy spectrum, they are made difficult by the small amount of fuel scattering expected with Nova targets ($\langle \rho r \rangle$ about 50 to 500 mg/cm²). For a target that results in $\langle \rho r \rangle = 50$ mg/cm², the downshifted component amounts to only about 1% of the neutron energy spectrum. The 14-MeV neutrons fire most of the detector channels, resulting in the need for a large number of channels if some are to still be live when these late neutrons arrive. (Gating schemes have been investigated, but most suffer from either long time constants for the gating process or a lack of sensitivity in the detection method.) Still, there is every expectation that this technique will work if a sufficiently large number of channels is fielded.

Measuring the upshifted component of the neutron spectrum does not require such a large number of channels because the neutrons of interest arrive first. But the measurement is based on a second-order process (scattered neutrons produce an energetic ion that must undergo another interaction); therefore, the upshifted component will be extremely small ($\leq 10^{-4}$ for 50 mg/cm²). If a sufficiently sensitive detector is developed or neutron yields are high, this method might prove to be feasible for measuring the fuel $\langle \rho r \rangle$.

Another possible application of neutron time-of-flight analysis using arrays of single-hit channels is measurement of the fusion burn width. Timing resolution as

low as 15 ps (FWHM) has been achieved with silicon charged-particle detectors.⁹⁵ This resolution is good enough that the entire burn history, estimated to be ~ 100 ps wide, could be followed. Because of thermal broadening of the neutron line, measurements must be taken very close to the target—5 to 10 cm away. Hard x rays could make this application difficult.

Summary. We are developing an instrumentation technique that will allow measurement of fuel ion temperatures through neutron time of flight. This technique is potentially useful for other neutron diagnostic applications, which are also being explored.

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Neutron Imaging with Zone Plates. We are investigating techniques for imaging the thermonuclear burn region of fusion targets irradiated by the Nova laser. Because the compressed core diameters of fusion targets are expected to measure tens of microns, our goal is to obtain images with better than 10- μ m resolution.

Alpha-particle images provided the first direct information about thermonuclear burn geometry in thin-walled, exploding-pusher targets.⁹⁶ In future high-density target experiments, only highly penetrating radiation like 14-MeV neutrons will escape the target core to provide information about the burn region.⁹⁷ One possible means of measuring this radiation is with a neutron pinhole camera; however, this would require a 10- μ m pinhole through about 10 cm of material and 10^{14} to 10^{15} source neutrons.⁹⁸

Zone-plate coded imaging (ZPCI) techniques are more promising for imaging small objects like the compressed core of a laser fusion target. We have been using ZPCI techniques to image nonpenetrating radiation like x rays and alpha particles for about 10 years.⁹⁹ The techniques are well developed.

Imaging penetrating radiation like 14-MeV neutrons with ZPCI techniques has two major advantages. The large solid angle subtended by the zone plate might substantially reduce the neutron yield

needed to produce a useful image. Also a zone-plate system with 10- μm resolution might be easier to fabricate and characterize than a pinhole system. Over the past year, we determined the conditions under which ZPCI techniques could be used with penetrating radiation.

Zone-Plate Coded Imaging. ZPCI is a two-step process. First a coded image is formed by letting a source of incoherent radiation cast a geometrical shadow of a Fresnel zone plate onto an image plane. Then the coded image (or shadowgraph) is processed to create a transparency, which produces an optical image of the original source when illuminated with coherent light.

A zone plate, illuminated with coherent light, contributes to the encoding/decoding process by producing a focused spot, the axial position of which depends on the size of the coded image. The coherent-light diffraction process performs a Fresnel transform, which serves to unfold the source-aperture convolution achieved in creating the coded image.

Figure 4-82 illustrates a neutron zone-plate system. The zone plate is a section of a set of concentric right-circular cones with a common apex at the center of the source distribution. The system is characterized by a source-to-zone plate distance S_1 , the zone plate-to-detector distance S_2 , the detector efficiency η_{det} , and the zone-plate characteristics: thickness (or length)

D , number of zones N , outer zone width ΔR_N , and macroscopic neutron cross sections Σ_i and Σ_o for the "transparent" and "opaque" zones.

The neutron zone plate differs from an ideal zone plate in three significant ways: it is thick, tapered, and has poor signal modulation. To cast a shadow with penetrating radiation, its opaque zones must attenuate the radiation enough to create modulation in the detector plane. For 14-MeV neutrons this requirement forces the zone plate to be several centimeters thick. For a zone plate of this thickness, the open (or transparent) zones might need to be filled with a solid material for support and ease of construction. Vacuum evaporation techniques could be used to construct such a zone plate and at the same time allow in situ characterization. With this method, alternating layers of opaque and transparent materials would be evaporated onto a rotating cylinder.¹⁰⁰ The materials used to construct the zone plate should be

Fig. 4-82. Neutron zone-plate system. A neutron zone plate is thick and tapered.

Fig. 4-83. Zone-plate transfer function. (a) Ideal zone plate produces sharply imaged zone boundaries. (b) Neutron zone plate gives sharp boundaries for point sources near the focus and fuzzy boundaries for sources away from the focus.

selected to minimize the neutron cross section in the transparent zones and maximize it in the opaque zones.

The boundaries between the opaque and transparent zones of a thick, tapered neutron zone plate cast sharp shadows when a point source is near the focus (apex of cones). The coded image achieves sharpness similar to that produced by a thin zone plate. Standard reconstruction techniques can be applied to such images. However, when a point source is away from the focus, the imaged boundaries appear fuzzy (see Fig. 4-83).

Signal modulation for a neutron zone plate is much poorer than for an ideal zone plate because the zones in the neutron zone plate are neither totally transparent nor totally opaque. The transparent zones cause some source attenuation and the opaque zones allow some transmission. Consequently the system's signal-to-noise ratio is decreased.

Modulation. The signal-to-noise ratio (SNR) depends on the modulation of the zone-plate transfer function and is calculated for a thin zone plate by dividing the expression for the reconstructed signal amplitude by the expression for the standard deviation of the signal amplitude. This SNR is also valid for sources near the focus of a thick zone plate. Rearranging the SNR equation gives an equation for the source strength required to produce an image. The required source strength depends on the desired SNR, the system dimensions, the detector efficiency η_{det} , the zone-plate materials, and the spatial distribution of the source. The number of events N_0^i needed in the i th source resolution element may be written as

$$N_0^i = \left(\frac{\pi}{2\sqrt{2}} \right)^2 (\text{SNR})^2 \left(\frac{4S_1}{d_{zp}} \right)^2 \left(\frac{1}{\eta_{\text{det}}} \right) \times \left(\frac{2}{P_1} \right) \left[\frac{(1+f)}{(1-f)^2} \right] \left(\frac{N_0}{N_0^i} \right). \quad (55)$$

Here d_{zp} is the diameter of the zone plate; P_1 is the transmission probability of the transparent zones; f is the ratio of opaque-zone transmission to transparent-zone transmission, $[\exp(-\Sigma_o D)/\exp(-\Sigma_t D)]$; and N_0 is the total source strength.

A similar relation can be written for an ideal pinhole system:

$$N_0^i = (\text{SNR})^2 \left(\frac{4S_1}{d_{\text{PH}}} \right)^2 \left(\frac{1}{\eta_{\text{det}}} \right), \quad (56)$$

where d_{PH} is the diameter of the pinhole.

The SNR, solid angle, and detector efficiency affect source-strength relations for pinhole and zone-plate systems in the same way. The advantage of a zone plate over a pinhole with similar resolution is its larger solid angle, which depends only on the number of zones. The ratio of a zone-plate solid angle to that of a pinhole system is

$$\left(\frac{d_{zp}}{d_{\text{PH}}} \right)^2 = \left(1 - \sqrt{\frac{N-1}{N}} \right)^2, \quad (57)$$

where N is the number of zones.

The constant in Eq. (55), which equals 1.23, is due to reconstruction coherence effects. Three additional factors associated with Eq. (55) decrease the advantage gained by the larger solid angle:

1. $P_1/2$ is the fraction of source events incident on the zone plate that pass through the transparent zones. It represents the effective open area of the zone plate.
2. The factor $(1+f)/(1-f)^2$ accounts for the zone-plate modulation caused by transmission through the opaque zones. In general, the better the signal modulation the lower the required source strength.
3. The term N_0/N_0^i shows that the required strength of the i th source resolution element depends on the ratio of the total number of source events to the number of events in the source resolution element of interest. For a uniformly emitting source, this expression is just the number of equal-strength source elements.

Using a thick neutron zone plate is a compromise between minimizing the required source strength and providing the required field of view, as defined by Fig. 4-84. For example, for alternating aluminum and gold zones ($\Sigma_1 = 0.1 \text{ cm}^{-1}$ and $\Sigma_o = 0.3 \text{ cm}^{-1}$), the minimum source strength occurs for a thickness of 9.3 cm but varies by only 10% between thicknesses of 6.7 and 12.7 cm.

The field of view decreases with increasing thickness D . Zones that are thin

behave like collimators. Thus to achieve the widest field of view, we would minimize the zone-plate thickness. For an aluminum-and-gold zone plate 5 cm thick, the field of view is 86% greater than for a zone plate 9.3 cm thick, and the required number of source events is increased by less than 40%.

The neutron zone-plate system offers better image detection than a pinhole system for imaging small sources because of its larger solid angle. The solid angle of a zone plate consisting of 100 zones is 40 000 times that of a pinhole with equivalent resolution. Figure 4-85 shows the ratio of the source strength required for a pinhole system to that required for a zone-plate system as a function of the diameter of a uniformly emitting disk source. The curve shows the advantage of a zone plate (aluminum and gold, 100-zones, 5 cm thick, 5- μm outer zone) over a pinhole 11 μm in diameter. For small sources, the zone plate system requires a factor of 1000 fewer source neutrons than a pinhole to achieve the same SNR. For source diameters larger than 170 μm , the source strength advantage is less than a factor of 10, and for sources larger than 540 μm , the pinhole system is actually superior. In a zone-plate system, unlike a pinhole, detected events from all source points contribute to the noise associated with the source point of interest. A zone plate offers substantial advantages for small, localized sources, but can actually be a poorer choice than a pinhole for large, distributed sources.

Fuzzy Zone Boundaries. The zone boundaries of a thick zone plate look fuzzy to penetrating radiation originating at locations other than the zone-plate focus. We have been using Monte Carlo techniques to study the effect of fuzzy boundaries on system characteristics like the field of view, planar resolution, tomographic resolution, quantum noise, and distributed sources. First we select a source distribution. We use point sources to evaluate the effect of various system parameters on the point spread function (PSF) and flat, uniformly emitting disks to investigate extended sources. The position and direction for each source event are randomly selected. Then each neutron is transported

Fig. 4-84. Definition of field of view α .

Fig. 4-85. Ratio of source strengths required to obtain equivalent images from pinhole and zone-plate systems vs source diameter for a uniformly emitting disk.

through the zone plate to the detector to form the coded image. If the neutron has a collision in the zone plate, it is removed from the calculation. Scattering events degrade the system modulation by contributing to the uniform background of the coded image. For the systems considered in this article, the change in modulation by scattering was negligible.

We developed a numerical reconstruction technique for analyzing these computer-generated coded images. The technique is similar to the laboratory method of illuminating the coded image with a coherent laser beam. Illumination is simulated by launching a spherical wave from each image point and calculating the resulting field amplitude over the source region of interest. The wavelength used is the same as that used for an ideal zone-plate system. Once the reconstructed source distribution is formed, lineouts are taken to determine quantities such as FWHM and peak amplitude.

We determined the PSF for a point source at the focus of the zone plate and for off-axis point sources located in the plane perpendicular to the zone-plate axis

Diagnostics

Fig. 4-86. (a) PSF amplitude and (b) FWHM show little deterioration for distances less than α .

Fig. 4-87. Lineouts from simulated images of source-point pairs located at the focus and at the edge of the field of view.

and containing the focus. Figure 4-86(a) shows the peak amplitude of the PSF vs the distance off axis for a 5-cm-thick, aluminum-and-gold zone plate with an outer zone width of $5 \mu\text{m}$. The zone plate is 50 cm from the source and 250 cm from the image plane. The distance scale is plotted in units of a field-of-view quality parameter α —the radial distance from the focus at which source radiation no longer has a direct path to the image plane through the outermost transparent zone of the zone plate. Figure 4-86(b) shows the PSF FWHM vs the source distance off axis. Based on these curves, we quantified the field of view by defining its edge to be at a point $[(S_1 + D/2)/D]\Delta R_N$ off axis. Our simulations show less than a 10% decrease in amplitude and less than a 15% broadening in the FWHM for distances less than α when α is smaller than the central zone radius. For distances greater than α , an image may be observed, but with poorer SNR and spatial resolution. For a 5-cm-thick zone plate, the field of view is $110 \mu\text{m}$.

To study the planar (or two-point) resolution of the neutron image, we used a resolution standard similar to the Rayleigh criterion to allow easy comparison with other imaging techniques. We selected a point-source pair and calculated the reconstructed intensity distribution. The point pair is considered marginally resolved when the reconstructed valley-to-peak ratio is 0.735, similar to the Rayleigh criterion for incoherent sources. Figure 4-87 shows lineouts taken through images of point-source pairs at the zone plate focus and at the edge of the field of view. For points 10, 11, and $12 \mu\text{m}$ from a point at the zone-plate focus, the zone-plate system classically resolves two points spaced $11 \mu\text{m}$ apart. For points 10, 11, and $12 \mu\text{m}$ from a point $53 \mu\text{m}$ off axis (at the edge of the field of view), the system resolves the points with $12\text{-}\mu\text{m}$ separation. These two-point resolutions agree well with the $10\text{-}\mu\text{m}$ resolution expected for an ideal zone plate.

To study the depth of field, we determined the PSF for points along the system axis. We found that source points within $S_1(S_1 + D/2)/2ND$ of the zone plate focus produce a good quality PSF when α is less than the central zone radius. This distance is the z-axis position at which source radiation no longer has a direct path through the outermost transparent zone to the image plane. We use this distance to signify the edge of the depth of field. For systems with S_1 between 2.5 and 200 cm and a magnification of 6, the PSF amplitude for points at the edge of the field has more than 90% of the amplitude of a point at the system focus, and the PSF FWHM increases by less than 20%. For a 5-cm-thick zone plate with 100 zones situated 50 cm from the target, the depth of focus

extends 2.5 cm on either side of the zone-plate focus.

We examined the effect of thick zone plates on tomographic resolution by looking at the PSF of a point source reconstructed in various axial planes. A measure of the tomographic capability of a zone-plate system is the distance in source space to the first null in the intensity distribution.³⁹ For a thin zone plate, the first null occurs at

$$\Delta = \left(\frac{S_1}{N} \right) \left[\frac{(S_1 + S_2)}{S_2} \right]. \quad (58)$$

This equation is useful as a tomographic figure of merit. The peak-to-null distances measured in Monte Carlo simulations agree well with Eq. (58) for source points at the focus of various systems and agree within a factor of 2 for points at the edge of the depth of field when α is less than the central zone radius.

Detector Element Resolution. It is possible to use different types of media for recording coded neutron images: scintillator arrays, film detectors, or track detectors. Each type of detector has a spatial resolution associated with the detection of an event. We examined the effect of detector resolution on zone-plate imaging by calculating the PSF for systems with detector arrays containing different size elements. When the elements are smaller than the outer zone width multiplied by the system magnification, the system performs nearly as well as if the exact position of each event were known. The signal amplitude decreased to 90% and the FWHM showed no change. For element dimensions twice the outer zone width multiplied by the magnification, the signal amplitude and the FWHM begin to deteriorate rapidly. A zone plate for Nova might have a 5- μm outer zone and a magnification of 6. Thus a 400- $\times$ 400-element array of 30- μm -square detector elements would be adequate for a 100-zone system. It should be noted that the number of detector elements required to achieve a specified system resolution depends only on the number of zones in the zone plate.

Conclusion. We examined the use of a zone plate to image penetrating radiation such as 14-MeV neutrons. Our calculations

indicate that it might be possible to produce high-quality images of neutron source distributions using a thick zone plate made of alternating layers of transparent and opaque materials such as aluminum and gold. The zone plate must be thick to optimize modulation, and it must be tapered to approximate an ideal zone plate for sources within its field of view and depth of field.

The primary effect of the zone-plate thickness is to cause the zone boundaries to appear fuzzy in the coded image of events coming from points other than the zone-plate focus. We developed concepts of field of view and depth of field for thick zone plates. For sources within the field of view and depth of field, image characteristics such as planar and tomographic resolution are nearly identical to the values calculated for a thin zone plate. For a field of view greater than the central zone, the PSF distortion depends on the number of zone-plate zones.

We demonstrated a numerical reconstruction procedure. In the past, all our zone-plate work involved laboratory reconstruction whereby the coded image was illuminated with laser light to form an image on film. In this work we numerically reconstructed the source distribution from computer-simulated coded images.

The zone plate has the potential for superior collection efficiency compared with a pinhole system of similar dimensions and resolution. With 100 zones, the zone plate has a solid angle 40 000 times larger than that of an equivalent pinhole system. Consider a 100-zone, 5-cm-thick, aluminum-and-gold zone plate with a 5- μm outer zone width, situated 50 cm from a target and 250 cm from a detector with 1% efficiency. This system requires a yield of 2.5×10^{12} neutrons to image a 30- μm -diameter burn region with 10- μm resolution and an SNR of 10. The zone plate has a significant advantage over a pinhole for small localized sources but provides little advantage over (and indeed may not be as good as) a pinhole system for imaging distributed sources that are larger than the central zone of the zone plate.

Much work remains before we can field a thick zone plate system for 14-MeV neutron imaging. A tapered zone plate must be

built, and an alignment procedure devised. An efficient, easy-to-read detector with sufficient resolution must also be developed.

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Krypton Activation Method for Measuring Fuel (ρr). Direct measurement of the mean areal fuel density (ρr) averaged over the duration of the fuel burn in ICF implosions has been accomplished by detection of high-energy deuterons and tritons (knockons) that arise from scattering of 14-MeV neutrons in the DT fuel.¹⁰¹ Use of this method is limited to experiments in which the mean areal densities of the fuel and ablator/pusher are small compared to the range of the knockons. The higher densities expected from implosions produced by the Nova laser require a method based on detection of highly penetrating radiation (e.g., neutrons or γ rays) or on detection of an activation product.

We have been studying the parameters required to develop an activation method for determining fuel (ρr) based on the reaction $^{80}\text{Kr}(n,2n)^{79\text{m}}\text{Kr}$. If ^{80}Kr is homogeneously distributed in the DT fuel with the atom fraction f , then, for a simple spherical implosion,

$$(\rho r) = \frac{N(^{79\text{m}}\text{Kr})}{Y\sigma(A_0/A)[f/(1-f)]}, \quad (59)$$

where Y is the DT fusion yield, σ is the cross section for the $(n,2n)$ reaction at 14 MeV, A_0 is Avogadro's number, A is the mean atomic weight of the DT fuel, and $N(^{79\text{m}}\text{Kr})$ is the number of reaction products formed in the implosion. The cross section has been found to be 0.415 b (Ref. 102). For fuel $\rho r \geq 50 \text{ mg/cm}^2$, as little as 0.1 at.% krypton in the fuel could be determined with a fusion yield as low as 10^9 if at least 1% of the $^{79\text{m}}\text{Kr}$ atoms formed in the implosion are collected and counted.

The choice of krypton in this method was dictated by the requirement that the tracer be chemically inert with respect to possible shell materials, such as glasses or plastics, and by the need for a diagnostic applicable to cryogenic targets. Because the tracer must be soluble in liquid DT for cryogenic applications, only the rare gases can be considered. Although the solubility of krypton in liquid hydrogen has not been measured, thermodynamic calculations suggest solubilities in the range of 0.7 to 3×10^{-3} mole fraction at 24.6 K (Ref. 103). The stable isotopes of the lighter rare gases do not lead to activation products with the necessary decay characteristics. LASNEX calculations of implosions on generic cryogenic targets suggest that fusion yields are not seriously affected by the presence of krypton at concentrations ≤ 0.1 at.% in the fuel.

The principal problem we face in developing this diagnostic is the collection of krypton after the implosion. Because of the absence of chemical binding, the only readily applicable collection methods are cryogenic absorption or ion implantation into a substrate, which exploits the kinetic energy of krypton ions in the expanding plasma or the kinetic energy resulting from acceleration in an external field.

We studied the characteristics of krypton collection on various substrates following implosions of glass-filled microballoons at the Omega laser in the University of Rochester's Laboratory for Laser Energetics (URLLE). The microballoons were filled with natural krypton and irradiated with thermal neutrons to produce 4.48-h $^{85\text{m}}\text{Kr}$, which served as a radioactive tracer for estimating collection efficiencies. The collectors included metallic foils of aluminum, titanium, beryllium, and zirconium and a number of organic films deposited onto metallic substrates. Typical characteristics of the microballoons and laser irradiations are listed in Table 4-19.

We found that the most reliable collection substrates were titanium and aluminum in the form of rolled metallic foils or as sputtered layers with very rough surface texture. Assuming that emission of all debris and quanta from the target takes place isotropically, we found that the intrinsic collection efficiency ϵ , defined as the probability of collection per incident ion, varied

Table 4-19. Target and laser irradiation conditions at URLLE.

Targets		Laser irradiations	
Shell material	Glass	Number of beams	24
Shell diameter	300–400 μm	Wavelength	0.35 μm
Shell thickness	1–2 μm	Pulse width	600 ps
		Pulse shape	Gaussian
Gas fill	Natural krypton	Total energy on target	1–2 kJ
Gas pressure	3–7 atm	Absorbed fraction	~ 0.50

substantially in implosions with the same nominal irradiation and collector characteristics. Nevertheless, a reasonable correlation of ϵ with total energy fluence or irradiance on the collector surface was observed with aluminum and titanium substrates. These data are shown in Fig. 4-88. The total energy fluence includes all sources of energy radiating from the target (e.g., particle kinetic energies, scattered laser light, and x rays).

Most of the collector foils were located 7 to 20 cm from the targets and at angles of 60 to 80° between the target surface normal and the radius vector from the target to the midpoint of the collector. We have not been able to observe correlation of ϵ with variation in these parameters. The absence of such relations suggests that variations in the dynamics of implosions are the dominant sources of scatter in the data. The overall results suggest that, for the experimental conditions shown in Table 4-19, sufficiently high intrinsic collection efficiencies can be achieved to meet the requirements for the diagnostic if total energy fluence can be maintained below about 0.1 J/cm².

The sensitivity of ϵ to the details of the implosion and the total energy fluence is consistent with an ion implantation mechanism that has the following characteristics:

1. The ion energy-distribution function varies.
2. Mechanisms that cause erosion of the collector surface, such as ablation and evaporation, limit the fraction of incident ions that implant beneath a stable surface layer.

Figure 4-89 shows mean ranges for krypton ions in aluminum and titanium calculated in the low-temperature limit according to the formulation of electron and nuclear scattering summarized in Ref. 104. For the mean krypton ion energies expected from these implosions ($T_{\text{ion}} \leq 1$ keV), the maximum depth of penetration at normal incidence is on the order of 100 Å, and it is clear that the collection efficiency must be highly sensitive to surface erosion.

We examined the erosion effect by determining the thickness of titanium eroded from a mirror surface. These collectors were prepared by vacuum evaporation of titanium onto silicon wafers. They

Fig. 4-88. Measured intrinsic collection efficiency for krypton vs total energy fluence at the collector surface.

Fig. 4-89. Ranges of ⁸⁴Kr in aluminum and titanium in the low-temperature limit.

were situated in the Omega target chamber at a mean distance of 11 cm from the target and with their surface normal at an angle of 82° with respect to the target-to-collector radius vector. The mirrors were partially covered with a thick titanium foil to prevent erosion of a portion of the surface. Following the implosion of a krypton-filled microballoon, the cover foil was removed and the step height between the original and eroded surface was measured with a profilometer. Figure 4-90 shows the results from two experiments. These data suggest that surface erosion may be a dominant factor in limiting the fraction of krypton that can be collected. At the higher energy fluences used in these experiments ($I \sim 0.5$ J/cm²), krypton ion energies on the order of 100 keV would be required for implantation into the stable surface if:

1. The eroded material were present at the time of arrival of the debris from the implosion.
2. No other mechanism were available for collection and trapping of the krypton.

Fig. 4-90. Thickness of titanium removed as a function of the total irradiance on a mirror surface.

(The assumption that energy-loss rates of krypton ions can be taken in the low-energy limit is probably conservative. The maximum possible temperature of the eroded titanium is in the range of 10 eV, and the actual krypton-ion energy loss rates¹⁰⁴ can be somewhat higher than assumed here.)

Some evidence that such high ion energies are indeed present was obtained in another series of experiments. Titanium foils covered with thin layers of Formvar or carbon were placed adjacent to uncovered foils at target-to-collector distances large enough that erosion of the titanium surface would be negligible ($I \leq 0.2$ J/cm²). The data from these measurements are consistent with an ion-energy distribution in which 1 to 3% of the ions possess energies as high as 0.2 MeV. While not conclusive, these results agree with the collection efficiency data shown in Fig. 4-88 if we assume surface erosion is not sensitive to the surface finish of the collector.

We conclude that an ion implantation scheme for collecting activated krypton ions from implosions of glass-filled microballoons can be made viable if surface erosion can be minimized or if external fields can be used to imbed ions in the stable surface of a collector. Additional measurements are planned to define the mechanisms controlling collection and to obtain more conclusive data on the ion-energy distribution function. Experiments to determine the parameters required for efficient collection under Nova conditions are now being designed.

Authors: S. G. Prusslin and S. M. Lane

Data Management and Analysis

Introduction

Hardware and software enhancements in 1985 made LEAD data processing faster and more efficient than ever before. From the first Nova target shot, all hardware and software components of the Laser Experiment Analysis Facility (LEAF) were ready to acquire, analyze, and archive data. Besides managing target shot data, the LEAF supported Nova operations by characterizing laser performance.

Most laser-characterization and target-diagnostic data are recorded on film. To convert these data to computer-readable form, the LEAF added an on-site photodigitizer-microdensitometer system (PDS), which became operational early in the year. For the first time, data could be digitized and analyzed within a few hours after an experiment. So far, the PDS has been used to digitize over a thousand pieces of film in support of experiments and diagnostic development. Linked for data transmission to the LEAF VAX computer, the PDS has proven indispensable for fast, automated data processing. The article "PDS Facility Update" describes this system.

Another major development this year was the completion of the "electronic shotbook," a software system designed to store and quickly retrieve shot data. This system, described in the article "Electronic Shotbook," brought the LEAD Program into the electronic information age. Instead of searching through hard-copy shotbooks, scientists can now select and view results on the screen of a computer terminal. The electronic shotbook displays the results from 15 different diagnostics and provides shot summary data.

Software of increasing complexity and a growing community of users imposed heavy demands on the LEAF VAX hardware. To maintain fast response times, more computing power became necessary. The CPU, formerly a VAX-11/780, was upgraded to an 11/785, providing a 50% increase in performance. Physical memory was also added to decrease processing time. The VAX is now

linked to three other computer systems used in data processing and analysis:

1. A PDP-11/34 controlling the PDS is linked via DECNET, the Digital Equipment Corporation (DEC) inter-computer network.

2. The Nova target diagnostics computer system is linked with a one-way, fiber-optic link.

3. The Octopus system, consisting of Cray and CDC-7600 computers, is linked via the Octoport network.

This third and most recent link allows transfer of data from the VAX to Octopus for analysis by complex codes. The article "System Update" contains more information on hardware enhancements.

Development of application software in 1985 focused on two areas:

1. Integration of film data processing codes with the ORACLE relational data base management system.

2. New and updated codes for the analysis of film data.

Data base development methodology called for simple user interfaces and a two-tiered system of usage: scientists responsible for a set of data store it in the ORACLE data base, while others, not familiar with the underlying physics, are able to perform data processing tasks.

Much of the new application software is used to analyze data for laser pulse characterization. These data include far-field beam images from the Nova opposed-port diagnostic and pulse histories from incident- and reflected-beam streak cameras. Other new software was written to support diagnostic development, especially in calculating the response of x-ray and optical diagnostics. The article "Application Software Update" describes these programming achievements.

In addition to data processing software, a new beam alignment code, GEOBEAM, was developed and installed at LEAF this year. By graphically depicting beam patterns, GEOBEAM enables scientists to position the 10 beams of Nova to achieve desired irradiation patterns. The article "GEOBEAM for Nova" describes this code.

**Authors: J. M. Auerbach
and E. M. Campbell**

System Update

The LEAF—where Nova experimental data are stored and analyzed—was the site of major upgrading in 1985. This year's modifications brought more powerful hardware, an improved version of the DEC VAX operating system, and new communications links.

The principal hardware changes were made on the facility's VAX. We replaced the computer's 11/780 CPU with an 11/785 and enlarged the memory from 5.5 to 20 megabytes. The CPU upgrade increased the computer's processing power by nearly 50%, and the larger memory allows more users to run larger codes at higher speeds.

Other hardware additions this year are:

- A second Random Element Processor¹⁰⁵ for more time-efficient plotting of graphs.

- Sixteen terminal lines.

- An electrostatic printer for the PDP-11/34.

Figure 4-91 is a diagram of the present facility.

We updated the VAX operating system by installing two new versions, VMS 4.0 and 4.1. Version 4.0 represents a major upgrade by DEC and required extensive changes in application software. Advantages of the revised operating system include improved security and file management.

In the area of communications, we added two new links to the network that carries data within the LEAF and connects it with other systems:

- A DECNET connection between the PDP-11/34 and the VAX.

- A connection between the VAX and the Octopus system.

The DECNET connection allows digitized data to be transferred from the PDS (Ref. 105) to the VAX for storage and analysis. This link saves the time and effort formerly required to write the data from one computer onto magnetic tape, and then read the tape into the other computer.

The connection to Octopus is a secure link from the LEAF to the Octoport gateway. It enables computer files to be transferred directly between the LEAF and any of the Octopus system's Cray mainframes. It also allows indirect transfers via the Crays to

Fig. 4-91. LEAF computer system components and logical connections.

other Octopus computers such as the CDC 7600s. An important advantage of the Octopus connection is terminal emulation by the LEAF VAX. With the VAX acting as an Octopus terminal, users who also have Octopus accounts can log on to any of the Cray computers from a LEAF terminal. This aspect of the network saves time and effort for the many users of both systems. A software interface was written to give users access to all network capabilities and allow them to diagnose the state of the network at any time.

Author: J. E. Randolph

Application Software Update

In 1985, application software development focused on:

- Processing Nova laser-characterization data.
- Processing Nova data from target diagnostics.
- Supporting diagnostic development.

Much of the development effort centered on an extensive software package for analyzing laser performance. This software is used to analyze laser intensity, power histories, and energy absorption by the

target so that laser-target interactions can be evaluated.

All codes developed to process Nova data, whether from target diagnostics or laser-characterization shots, now transfer results to the ORACLE data base management system or to standard-format data files. This advance has made code output available to users through the electronic shotbook.

Processing Laser-Characterization

Data. Software used to analyze laser focal-spot images was made more interactive and now provides files containing intensity-distribution data for the electronic shotbook. Earlier versions of this software subtracted background exposure and calculated the center of the focal spot automatically. If intensity varied over a few orders of magnitude, this approach gave misleading results; therefore, defining the background and spot center was made interactive. Users now determine these parameters by moving a Ramtek trackball cursor on a computer display. Also included in the upgrade are an improved line contour plot and optional scaling on 2-D profile plots. Figure 4-92 shows a color display of the intensity distribution of a Nova focal spot. Figure 4-93 shows the corresponding calculated power distribution.

New software for calculating laser energy balance was completed this year. It uses the output of a photodiode processing program¹⁰⁶ to calculate the total energy scattered and absorbed by the target. The new software produces optional plots of fluence viewed in three different coordinate systems: chamber spherical coordinates, spherical coordinates relative to the beam normal and polarization vectors, and spherical coordinates rotated relative to the target surface normal. Figure 4-94 shows fluence plotted as a 3-D "sky map." This plot shows light scattered in the incident and specular directions as well as sidescattering. The software also provides summary tables containing fluence-integration data and transfers fluence and angularly integrated data to the ORACLE data base for access by the electronic shotbook.

Also installed this year was a new data base for storing and retrieving

Fig. 4-92. Color-enhanced intensity distribution at a Nova beam focal spot from an opposed-port diagnostic film record.

calibration and configuration data for the energy-balance modules. As a result, the code processing photodiode data, which provides input for the data base, was upgraded too. This program reads a file containing counts from an analog-to-digital converter and extracts calibration parameters from the ORACLE data base to the radiant intensity (in joules/steradian) measured by each detector. Enhancements include a report on the configuration of all photodiode detectors and warning messages indicating saturated and faulty detectors with suggested adjustments for improving the detector signal.

For every Nova shot, streak cameras are capable of recording 20 time-resolved pulse shapes—two produced by the incident and reflected light from each of the laser's 10 beams. Given this large volume of data, it is important to calculate histories of incident and reflected laser power in a timely, automated fashion. A program was written to analyze each streak record and produce profiles of incident and reflected laser power vs time. The software reads a digitized film-density image array, converts density to exposure over the entire image, and takes lineouts through each streak record. All parameters needed to convert the data to a power-vs-time profile (e.g., laser energy, camera sweep speed, streak identifier, location, width, and delay) are stored in and retrieved

Data Management and Analysis

from ORACLE data base tables. The program uses the processed film data, shot parameters from ORACLE, and an absolute time reference from a timing fiducial to create power-history profiles on demand. Because ORACLE relates all this information,

Fig. 4-93. Radial power distribution calculated from the intensity distribution in Fig. 4-92.

Fig. 4-94. Sky map of fluence data from a Nova scattered-light photodiode.

users enter only the shot number to obtain a power history.

Processing Data from Target Diagnostics. Important new routines were added to the software used to process data on target diagnostics connected to a CAMAC-based data-acquisition system. Unlike previous control systems, the Nova target-diagnostic control software enables technicians to electronically map the links between diagnostic sensors and the recording module connections by using the control system's interactive menus. After the shot, the mapping is written into the last section of the Nova shot data file. In this form, the mapping consists of a few thousand lines of ASCII characters—a size that is impractical for fast data processing. To reduce this volume, a new program named MAP_PROCESS was written. It parses the mapping data and writes it into ORACLE data base tables reserved for this purpose.¹⁰⁶ The mapping data is reduced from a few thousand records to a few hundred rows in the data base. In this form, it is usable by all existing processing codes for data from CAMAC-based instrumentation. Codes for CAMAC data written this year include:

1. A program to process and archive data from a system that monitors streak-camera trigger timing. The software allows scientists to track the trigger-timing stability for all streak cameras on a shot-by-shot basis.

2. A program to process data from bolometer and pyrometer x-ray detectors. Output from these broad-band x-ray detectors is recorded on CAMAC-compatible waveform digitizers. The integral under the waveform corresponds to total x-ray emission from the target. Figure 4-95 shows a sample waveform plot from the new program.

An important software-development task this year integrated the ORACLE relational data base management system with most of the film-data processing codes. These codes are used to process digitized film data from streak cameras, spectrometers, and imaging diagnostics such as the x-ray

microscope and the pinhole camera. Data base tables were created to hold instrument set-up data for each shot and lists of instrument components, such as filters, mirrors, and gratings, which give the instrument its spectral or spatial resolution. The data base tables are updated by the scientist responsible for the diagnostic. Afterwards, the processing code can be run by others who are not familiar with the details of the instrument, because the software and the data base contain all information required. Figure 4-96 shows the flow of information to a typical film-processing code. Note that all film-processing codes produce files containing

SHOT: 15112207

BOLOMETER CHANNEL WAVEFORM

SYSTEM:
BOLO

DET NO:
4

CHANNEL ID:
PU000104

r7912 ID
28

ATTN:
2

DELAY (NS):
0.00

SCOPE SERIAL NO:
882

Fig. 4-95. Waveform plot produced by the bolometer diagnostic processing software.

Fig. 4-96. Function block diagram of the LEAF film-data processing codes.

calculated intensity distributions, which serve as data sources for the electronic shotbook and for software that analyzes film data.

In addition to film-processing codes for specific diagnostics, LEAD programmers are developing a general, interactive code capable of analyzing film-data files from different sources. This program allows a user to view a color-enhanced display of film density (or intensity) on a Ramtek

Fig. 4-97. Color-enhanced display produced by an image-analysis program, showing a selected lineout.

Fig. 4-98. Detailed plot of the lineout in Fig. 4-97.

monitor. Users can analyze the corresponding data by selecting options from menus on the screen or by specifying positions with a trackball cursor.

A menu of analysis options allows the user to:

- Obtain single or cross-average lineouts in horizontal, vertical, or arbitrary directions.
- Obtain radial (azimuthally averaged) profiles.
- Obtain histograms and histogram running integrals of a selected region of the intensity/density data.
- Do uniform and nonuniform background subtractions from the intensity distribution.
- Obtain integrals of a specified portion of the distribution.

All profiles can be plotted in hard copy. Axis types (log or linear), titles, and scale limits are either automatically selected or user-specified. Locations for profiles can be specified by a keyboard entry or with the trackball cursor. Figure 4-97 shows a display of x-ray microscope data produced by irradiating a gold disk at Nova. A user-selected, tilted lineout is also displayed. Figure 4-98 shows the corresponding graph of the lineout data.

An essential phase of converting film density to quantities representing conditions at the source is to unfold the film response. This unfolding involves use of the D -log E relation, which provides the film density corresponding to a given exposure. This year new software was installed to calculate the D -log E relation for film from:

- Rattle pairs: spot arrays in which the intensity ratio from one spot to another is known.

- Two-pattern sets: two arrays of film density data created by identical spatial sources differing in intensity by a known ratio.

- Arbitrary patterns in which the ratio of density to incident light is not constant.

Work began in 1985 on unfolding from image data not only film response and the spectral response from instruments, but also spatial response. For the first time, a Floating Point Systems array processor connected to the LEAF VAX is being used for applications such as fast Fourier transform (FFT) filtering and deconvolution of instrument response. The array processor can do the calculation-intensive transforms many times faster than the VAX. Deconvoluting the spatial response of streak cameras is necessary for data reduction in forthcoming hydrodynamic-instability experiments.

Supporting Diagnostic Development.

Software for calculating the response of x-ray diagnostics has also improved this year. The Henke library of atomic f_1 and f_2 scattering factors used in filter-transmission and mirror-reflectivity calculations has been extended to an energy range of 30 eV to 10 keV. This range is used for designing diagnostics for x-ray laser experiments. Software used to calculate the response of the streak crystal spectrometers has been expanded to include convex and flat crystals as well as elliptical crystals.

**Authors: J. M. Auerbach
and D. S. Montgomery**

GEOBEAM for Nova

A new computer program modeled after the Shiva code GEOBEAM¹⁰⁷ was written to aid in the setup of Nova experiments. The new GEOBEAM helps researchers position each of the laser's 10 beams to

achieve the desired irradiation patterns on the target. It also aids in target placement and orientation. The program allows the user to specify:

- Beam profile parameters (diameter of the circular beam cone as a function of distance from the best focus spot).

- The offset of the beam's best focus spot from the center of the target chamber.

- The target size and type (disk or sphere).

- The orientation of the target normal for planar targets.

- The angle from which the target is viewed.

The user can also select any of the following graphics devices for output: a Ramtek color display, a Retro-Graphics high-resolution graphics terminal, or a Versatec plotter. Nongraphic information—beam parameters and plot data—can be directed to a file and stored for later use.

Graphic output shows the target surface marked with the intercepts of each beam cone. The intercepts are color-coded on the Ramtek display and are drawn with different symbols for each beam on black-and-white Versatec plots and on the monochrome Retro-Graphics display. The target is drawn from a user-specified viewpoint. For spherical targets, views of the near and the far hemispheres are shown.

The program works within the geometry of the Nova target-chamber coordinate system (Fig. 4-99). This is a spherical coordinate system designed for the positioning of focusing lenses and diagnostics mounted on the chamber. The radius of this system is defined as the distance from the target chamber center. The polar angle θ is defined as the angular distance from the east axis. It varies from 0° to 180° (the west axis). The azimuthal angle ϕ is the angular distance from the up axis, varying from -180° to $+180^\circ$ (with positive defined as being toward the south axis).

Although the program works within the spherical system, it can interpret input parameters expressed as spherical, cylindrical, or Cartesian coordinates. The displacement of a focusing lens, which is specified as the displacement of the best focus spot of the beam, can be expressed in terms of any of these coordinates. Target-displacement parameters are most

Fig. 4-99. The Nova target-chamber coordinate system showing Nova beam coordinates.

conveniently entered as Cartesian coordinates. Regardless of which coordinate system is used for data entry, the program always plots in a cylindrical system.

GEOBEAM operates in three phases:

1. It prompts for setup information, which the user enters via a series of menus. It also reads setup data from disk files.

2. It calculates a grid of intersection points where each beam meets the target surface.

3. It creates the plots and sends them to the selected graphic device.

Figure 4-100 shows the menus that prompt user entries. Users can cycle through these menus as often as desired before starting the calculation.

The main calculation routines of the program establish an array of beam-and-target intersection points. The program calculates the intersection of the beam axis and the target for each beam. A disk target lies in a given plane offset from the origin by $\text{offset}(x)$, $\text{offset}(y)$, $\text{offset}(z)$ with

its normal toward (θ_T, φ_T) . A line along the beam axis runs through the points $(\Delta_{\text{up}}, \Delta_{\text{south}}, \Delta_{\text{east}})$ and $[(\sin \theta_B \cos \varphi_B) - \Delta_{\text{up}}, (\sin \theta_B \sin \varphi_B) - \Delta_{\text{south}}, \cos \theta_B - \Delta_{\text{east}}]$. The intersection of the line and the target plane is given by equating the line

$$\frac{x - \Delta_{\text{up}}}{\sin \theta_B \cos \varphi_B} = \frac{y - \Delta_{\text{south}}}{\sin \theta_B \sin \varphi_B} = \frac{z - \Delta_{\text{east}}}{\cos \theta_B} \quad (60)$$

and the plane

$$z' = \Delta z', \quad (61)$$

where

$$z' = xR_T(3, 1) + yR_T(3, 2) + zR_T(3, 3), \quad (62)$$

BEAM PARAMETERS MENU

Beam No.	Used	Yc	Scout	Beam
1	Yes	59.	81.	0.
2	Yes	69.	-81.	0.
3	Yes	95.	-31.	0.
4	Yes	95.	31.	0.
5	Yes	0.	100.	0.
6	Yes	0.	-100.	0.
7	Yes	95.	-31.	0.
8	Yes	95.	31.	0.
9	Yes	59.	81.	0.
10	Yes	69.	-81.	0.

To move around the display use the arrow keys.
PF1 => Calculate, PF3 => Previous menu, PF4 => Next menu
CTRL Z => Exit Novabeam
All distances are in microns and all angles in degrees.
Displacements are from chamber center.
C => Cylindrical, X => Cartesian, D => Defocus

TARGET PARAMETERS MENU

Disk	Radius	Yc	Scout	Beam
500.	0.	0.	0.	0.

To move around the display use the arrow keys.
PF1 => Calculate, PF3 => Previous menu, PF4 => Next menu
CTRL Z => Exit Novabeam
All distances are in microns and all angles in degrees.
Displacements are from chamber center.
D => Disk, S => Sphere

VIEWPOINT MENU

Viewpoint	Beam
130.	54.

To move around the display use the arrow keys.
PF1 => Calculate, PF3 => Previous menu, PF4 => Next menu
CTRL Z => Exit Novabeam
All distances are in microns and all angles in degrees.
Displacements are from chamber center.

FILE RESULTS OUTPUT OPTIONS

No	No	Yes	Yes	No	1

To move around the display use the arrow keys.
PF1 => Calculate, PF3 => Previous menu, PF4 => Next menu
CTRL Z => Exit Novabeam
All distances are in microns and all angles in degrees.
Displacements are from chamber center.

Fig. 4-100. GEOBEAM data-entry menus.

and

A spherical target is given by

$$\begin{aligned} \Delta z' &= [\text{offset}(x)]R_T(3, 1) \\ &+ [\text{offset}(y)]R_T(3, 2) \\ &+ [\text{offset}(z)]R_T(3, 3). \end{aligned} \quad (63)$$

R_T is the rotation matrix from the chamber coordinate system to the target coordinate system. For the line to intersect the target (and not just the target plane), the following condition must also be met:

$$x'^2 + y'^2 \leq R_{\text{disk}}^2, \quad (64)$$

where x' and y' are rotated to the target coordinate system, and R_{disk} is the radius of the target. This expression generates the first point in the intersection array. Subsequent points are found by incrementing x' and y' until the edge of the beam is found.

$$\begin{aligned} R_{\text{sphere}}^2 &= [x - \text{offset}(x)]^2 \\ &+ [y - \text{offset}(y)]^2 + [z - \text{offset}(z)]^2. \end{aligned} \quad (65)$$

The beam axis may intersect the sphere in two places. The proper intersection is determined by the program, depending on the incident beam. Because the sphere is most easily described in spherical coordinates, intersections are converted to r , θ , and ϕ for the starting point in calculating an intersection array, as

$$\begin{aligned} r &= \sqrt{[x - \text{offset}(x)]^2 + [y - \text{offset}(y)]^2} \\ &+ [z - \text{offset}(z)]^2}^{1/2}, \end{aligned} \quad (66)$$

$$\theta = \arccos \left[\frac{z - \text{offset}(z)}{r} \right], \quad (67)$$

$$\varphi = \arctan \left[\frac{y - \text{offset}(y)}{x - \text{offset}(x)} \right]. \quad (68)$$

Subsequent points are found by incrementing θ and φ until the edge of the beam cone is found.

The points in this array that lie near the surface of a given beam cone mark the outline of the intersection. For graphic output, this outline is projected into the plane, perpendicular to the user-specified line of sight.

Each beam is represented by an array of beam diameters as a function of displacement from the position of best focus. Currently all 10 beams are described by the same function, which represents an $f/4.22$ focusing lens,¹⁰⁸ converging to a beam with a minimum diameter of 100 μm . Points that lie within the cone of rotation

formed by this array are considered to lie within the beam. As the characteristics of each beam are measured, more realistic beam profiles may be substituted.

To represent a disk target, GEOBEAM's plotting routines draw an ellipse that is offset from the center of the plotting area by a user-specified distance. The view of the target represented by the ellipse is determined by an angle value entered via the "Viewpoint Menu." To represent a spherical target, the plotting routines draw a view of a sphere projected into a plane perpendicular to the viewing angle.

After plotting the target, GEOBEAM plots the beam outlines, distinguishing them by color coding or by drawing with unique symbols. Adjacent to the plot, the program displays the following parameters for each beam and for the target:

- The angle between the beam and target normal (labeled "Beta").
- The inverse of the area of the beam spot in inverse square centimeters (labeled "1/Area").
- The target type, size, and normal orientation.
- Beam offsets in terms of the cylindrical coordinate system.
- Target offsets in microns.
- The angle of view.

A legend showing the color or symbol for each beam with the appropriate beam number is also displayed. Figure 4-101 shows a typical GEOBEAM plot of a disk target. Figures 4-102(a) and 4-102(b) show the two hemispheres of a spherical target.

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Fig. 4-101. Plot showing beam intersections on a disk target.

Fig. 4-102. Plots showing beam intersections on a spherical target: (a) near hemisphere and (b) far hemisphere.

Electronic Shotbook

With the start of the Nova-activation shot series in February 1985, the electronic shotbook¹⁰⁹ went into operation on the LEAF VAX. During the remainder of the year, shotbook software development kept pace with the addition of new diagnostic devices. Diagnostic-related displays, menus, and ORACLE¹¹⁰ update forms were

ready for users as shot data became available. The shotbook now displays data from 15 different types of diagnostics.

Shotbook software uses the DEC Forms Management System (FMS)¹¹¹ to produce its menus. The software is designed around the ORACLE relational data base management system. ORACLE is used to define the unique set of Sequel Query Language (SQL) statements required to

Fig. 4-103. Menus showing the update and selection functions in the shotbook.

(a) ** LEAF Electronic SHOTBOOK **

UPDATE INFORMATION UPDATE OPTIONS

Return to PREVIOUS Menu

- . DISK
- . FOIL
- . FUEL BALL
- . HOLEPUNCH
- . X-RAY LASER

(b) ** LEAF Electronic SHOTBOOK **

DIAGNOSTIC SECTION UPDATE OPTIONS

Return to PREVIOUS Menu

X-RAY DIAGNOSTICS

- . DANTE-Low Energy Spectrum
- . FFLEX-High Energy Spectrum
- . HELMUT-Crystal Spectrograph
- . KEVIN-EB Microscope
- . LEO-Streamed Microscope
- . MOPICS-MCP Cray. Inc. Spectr.
- . OI-Optical/X-ray Streak Camera
- . PINHOLE-Pinhole Imaging Camera
- . SCS-Stra. Crystal Spectrograph.
- . SSSS-Spat. Disor. Streamed Spectr.
- . SLS-Stra. Grating Inc. Spectr.
- . TSSS-Trans. Grating Strk. Spectr.

SCATTERED LIGHT DIAGNOSTICS

- . OPD-Dioped Port Diagnostic
- . SPP-Streamed Optical Pyrometer
- . SSS-Streamed Optical Spectrograph

(c) LEAF Electronic SHOTBOOK

SHOTBOOK SELECTION PARAMETERS

Shot Series :

Shot Subseries :

Shot Number :

Target Chamber :

Diagnostic System :

(d) ***** SHOTBOOK DIAGNOSTIC USER/GROUP ASSIGNMENTS *****

Page 1 of 2

Block 1

USER NAME	DIAGNOSTIC TYPE	PULSE SHAPE
ED	SUN	SOS
ED	EM	TSSS
HEMWAY	MOPICS	SCS
FFLEX	OPD	SSPS
KB XON	OX	SPP

* Enter a Y to assign a DIAGNOSTIC TYPE authority to USER.

***** AVAILABLE DIAGNOSTIC TABLE PRIVILEGES *****

Block 2

TABLE NAME	S

Char Model: Replace Page 1 Count: 40

(e) LEAF Electronic SHOTBOOK

CREATE OR MODIFY A SHOT SUMMARY

- You are about to create or modify a shot summary using the EDIT editor. To save the edited summary, exit EDIT with the EXIT command. The QUIT command ends the editing session without saving your changes.
- Please enter SHOT number: 507000 and CHAMBER ID: 1 and press RETURN (or press PPF twice to return to the previous menu)
- Enter a REF shot number:
- Press CSETUP to enter EDIT editor (or press PPF to return to the previous menu)

gather the data for each of the shotbook displays. The commercial package DISSPLA¹¹² provides graphics at the terminal and in hard-copy form.

Hardware and software upgrades in 1985 led to substantial improvements in shotbook performance. The speed of the shotbook software increased notably after the LEAF VAX physical memory was enlarged, and the time required to produce hard copy was cut significantly by the addition of the Versatec Random Element Processor. The article "System Update," earlier in this section, describes these upgrades. The shotbook is now widely used by the LEAD scientific and engineering staff. It has achieved its design goal of providing timely, consistent Nova experimental results to ICF scientists.

Shotbook User Interface. The utility of the shotbook was improved by adding a new target section, which stores and retrieves detailed data about each target used in a laser experiment. Users can retrieve these data by target type and shot number. This section was required to maintain a record of the targets used for each shot at the two Nova chambers. Figure 4-103(a) shows the target-section menu. Similarly, the diagnostics-update section of the shotbook was greatly expanded. Its menu is shown in Fig. 4-103(b).

Nova's dual target chambers can produce two distinct sets of results for each series of laser shots. So it is vital to be able to distinguish data from each of the two target chambers. The 1985 programming

effort guarantees separate storage of and access to these data. Users can now specify the desired target chamber on the shotbook menu shown in Fig. 4-103(c).

Shotbook has been made more user-friendly by the addition of stringent internal checking. System checks ensure that, at each of the many levels of data request, the information sought is both existent and consistent. Another user-friendly feature is the ease of entering commands. All shotbook commands consist of a single keystroke except "EX" (for exit from the program).

To provide data base integrity and security, a new access-control system was completed this year. Through specially designed tables in the ORACLE data base, the system allows each experimentalist to update, insert, or delete the data for which he or she is responsible. At the same time, it maintains a "select only" environment for all other users. Figure 4-103(d) shows the entry screen for assigning user-access privileges.

Detailed shot summaries are written by using EDT, the VAX on-line editor. The summary can fill a single page or, more typically, be a lengthy, structured document. Writing summaries has been simplified by a new option on the menu in Fig. 4-103(e), which allows the experimentalist to begin with an existing summary and modify it. This option saves time because summaries are often similar from one shot to the next.

Shotbook Displays. Displays are produced on VT-100 terminals equipped with a Retro-Graphics enhancement.¹¹³ The resulting screen resolution of 1024 × 780 raster points offers high-quality graphics. Because most displays are multi-paged, a new option was added to allow direct access to specific pages.

Two-dimensional data displays are flexible. Users can make changes, such as adjusting the scaling of the x or y axis, without affecting the data in the data base. Results of such changes can be displayed and printed.

The shotbook provides a custom display for each diagnostic on Nova. These displays were designed by using Yourdon structure charts.^{114,115} Shotbook data are linked (in a virtual sense) with the processing codes by a common file format to

Table 4-20. Electronic shotbook displays.

Name	Diagnostic or type of display
TSMURY	Shot summary
EBSUM	Energy-balance-module summary
LBEAMS	Laser beam pulse shapes
OPD	Opposed-port diagnostic
DANTE	Low-energy x-ray spectrum
FFLEX	High-energy x-ray spectrum
KB	3 × , 8 × K-B x-ray microscopes
22X	Streak x-ray microscope
HNWY	X-ray crystal spectrograph
MCPIGS	Multichannel-plate grazing-incidence spectrometer
PINHOL	X-ray pinhole camera
SCS	Streak crystal spectrograph
SOP	Streak optical pyrometer
SOS	Streak optical spectrometer
TGSS	Transmission-grating streak spectrometer

ensure consistency between the related processing-code output and the data input. The routines in the shotbook for reading these files are RD_CONTUR_FILE (for contour data) and RD_PROFILE_FILE (for multiple-profile data). To aid future

modification, all display software is written with complete internal documentation. The name and purpose of each variable are defined, and a narrative header page precedes every subroutine. Beyond giving the requested technical information, the

Fig. 4-104. Sample displays of data from diagnostics processing codes.

displays themselves include titles and labels that thoroughly explain the source of the information. For example, each display shows a shot number, series number, subseries, chamber identifier, and diagnostic identifier. Additional space on plot displays is reserved for three 80-column lines of technical comments. The shotbook offers 15 such displays, which are listed in Table 4-20. Figure 4-104 shows representative examples.

Work in Progress. Work now under way will add the following features to the shotbook:

- Display sections for new diagnostics that come on line.
 - Selection, by single keystroke, of the previous shot, next shot, table of contents, or current display.
 - Separate hard-copy plots of lineouts in addition to the usual plots of a lineout superimposed on a contour.
 - Target-fabrication data, which will eventually include target-characterization data such as x-ray radiographs.
 - A target-usage display to provide ready access to information about the types of laser targets used "to date."
 - An electronic preshot module based on the electronic shotbook architecture. This menu-driven, ORACLE-based environment will permit easy, consistent definition of parameters and diagnostics for forthcoming experiments and provide hard-copy reports containing this information for the staff. The definition will be identified by a preshot number.
 - The capability to store the Preshot parameters in the ORACLE data base for later review by general users or for reuse by a shot physicist as the basis for defining a new experiment.
 - The capability to connect Preshot parameters with post-shot data in the shotbook. This connection will be made at the first level of post-shot data processing when the Preshot number is related to the actual shot number. Then a subset of the data in the ORACLE Preshot tables will be inserted in the ORACLE tables used by the shotbook.
- Conclusion.** The electronic shotbook gives timely, reliable access to shot data and offers an easy way to review all experimental results. Any data displayed at the terminal can be printed in hard

copy. Through its ORACLE-based access-control system, the shotbook guarantees the integrity of the experimental data base and the consistency of shot data shown in its displays.

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PDS Facility Update

A Perkin-Elmer PDS 1010GM microdensitometer system, purchased in late 1984 (Ref. 116), was brought on-line to digitize film data from Nova. The system consists of a PDS 1010GM granite-based scanning unit, a Class-100, laminar-flow clean hood, a DEC PDP-11/34 computer for control and data acquisition, control consoles, and hardware connections. The PDS is housed in the LEAF computer area. Perkin-Elmer installed the system and performed an initial checkout. It was then turned over to the LEAD Program's data management and analysis staff for testing and evaluation.

Testing and evaluation concentrated on two areas:

1. The linear-positioning accuracy of the PDS stage motors and linear optical encoders.
2. The densitometric accuracy of the data-conversion system: the photomultiplier tube (PMT), analog-to-digital converter (ADC), and log amplifier.

Linear-Positioning Accuracy. The PDS is a flatbed, scanning microdensitometer with stepping motors that move the sample stage in the *x* and *y* directions. To locate and control the actual stage position accurately, linear optical encoders are installed on both axes. A glass reticle was used to measure the positioning accuracy of the encoders and the stage-position logic in SCANSALOT, the software that controls the densitometer. The locations and lengths of the scribe lines on the reticle were measured with a He-Ne laser interferometer, and their positions certified to within 0.1 μm . The reticle was fastened to the PDS stage and moved to selected *x, y* coordinates with the SCANSALOT MOVE command. The reticle was placed in the center and in all four corners of the

250 × 250-mm stage. The maximum error demonstrated by any of these tests was less than 2 μm over a stage-traveling distance of 120 000 μm . This stage test is repeated by the PDS system manager every two months.

Densitometric Accuracy. When film is scanned, the PMT generates an electrical signal that is proportional to the optical transmission of light passing through the film. This signal is digitized by an ADC and translated to optical-density data by a logarithmic amplifier. To test the densitometric accuracy of the PMT/ADC/log-amp system, a gray scale with known, discrete steps (Fig. 4-105) was produced in the LEAD Program's microfabrication facility. The scale is made of chromium deposited on glass in patches, each with a different, uniform thickness or optical density (OD). These patches are deposited in order of increasing density. There are two areas of patches, one ranging from 0.03 to 2.50 OD and the other, from 2.00 to 4.50 OD. Between the patches, the first area is clear (OD = 0.03), and the second area has a minimum density of 2.00. This series of increasingly larger density transitions permits us to measure the system response under conditions similar to those anticipated in actual operation.

The gray scale was scanned while conditions were varied over a range of aperture sizes, incremental step sizes in the x and y directions, and speeds from 200 mm/s to 1.25 mm/s. Two different scan modes were also tested, the edge mode and the raster mode. In edge scanning, the data are taken in one direction only, and the stage returns to its initial scan position before taking another line of data. In raster scanning, the stage scans in both the positive and negative directions on the x axis. After the end of a scan line in either mode, the stage moves in the y direction to the next scan line, and the process is repeated.

Film scanned with various apertures over a range of density transitions yielded little variation in optical density compared to that scanned at various speeds over the same range of density transitions. At the same location on the gray scale, different scanning speeds gave different density values. These discrepancies were due primarily to an inadequate response time by the PMT/ADC/log-amp circuitry.

In testing at the different speeds, optical density was compared to that at the lowest speed, 1.25 mm/s. The criterion for passing the test was a difference <0.1 OD at a position within two pixels of the edge of a patch on the gray scale. The scans over

Fig. 4-105. Gray scale developed to test the densitometric accuracy of the PDS.

Fig. 4-106. Number of images digitized per month by the LEAF PDS microdensitometer system for the first nine months of operation.

small density transitions, from 0.0 to 1.5 OD, were accurate at speeds up to 100 mm/s. With density transitions of 0.0 to 2.5 OD, the speed had to be reduced to 50 mm/s to pass the test, and for transitions from 0.0 to 3.5 OD, the speed could be no higher than 25 mm/s.

Another discrepancy was noticed when the raster scans and the edge scans were compared. In a raster scan, adjacent lines showed different densities for adjacent y positions at the same x position. The measured discrepancy was similar to that found in the speed tests, but more pronounced. Although plots of density vs (x, y) produced by a raster scan differed from those produced by an edge scan for adjacent scan lines, they were nearly identical for alternate scan lines. This difference in measured density at the same location was first thought to be a positioning error, but positioning tests proved otherwise. The difference was finally attributed to the slow response time of the ADC/log-amp circuitry. Consultation with users of other PDS 1010-GM systems with scanning capacity in the 0.0- to 5.0-OD range supported this conclusion.

The results of testing the densitometric accuracy of the PDS led to restrictions on the scan mode and on the maximum scan speed. Only the edge (or unidirectional) mode is now used. The scan speed limit was set at 25 mm/s for most data in the 0.0- to 3.5-OD range. For higher optical densities or for scan sizes less than

1000 μm long, the speed has been reduced to 4 mm/s or less. The reduction in speed for short scans is due to the overhead time taken by the system to reach the scanning speed and to return to the scan starting point between lines.

A possible solution to the response-time problem is now under development: Perkin-Elmer is testing a faster ADC/log-amp circuit. If successful, this new circuit can speed up the system by a factor of 2 to 3. Initial tests of a prototype model showed improvements over the present 0.0- to 5.0-OD circuit. A final upgrade from Perkin-Elmer is expected in early 1986.

PDS Performance. The final testing of the system and training of the staff was completed in April 1985, and the system has been in operation ever since. Figure 4-106 shows the number of images per month scanned and transferred to the LEAF VAX during the first nine months of operation. The typical image (512×512 pixels) takes an average of 10 to 12 min to set up and scan. After a scan, the resulting data are transferred from the PDP-11/34 to the VAX via the inter-computer network DECNET. It takes another 4 min to transfer a 512×512 -pixel image file and convert it to image-display format on the VAX.

When results are needed immediately to modify an experimental setup, scanning and transfer for one sample are done in sequence, taking approximately 15 min. However, the normal procedure is to digitize images for an entire day and transfer the resulting data to the VAX the following morning while the PDS electronics warm up (typically for 30 to 45 min).

To improve the performance of the PDS system, a Class-100, laminar-flow clean hood was added to cover the scanning assembly. It encloses the scanner in an area continuously flushed with air that has been passed through two stages of filters. These filters reduce the quantity and size of dust particles that fall onto the digitizer during a scan. Fewer than 100 particles per cubic foot greater than $0.5 \mu\text{m}$ in diameter can penetrate. Because dust particles can range from 50 to $500 \mu\text{m}$ in diameter, and the typical scanning aperture is $60 \mu\text{m}$ square at the sample plane, the clean hood makes an important contribution to the quality of the digitized data.

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Nova's beamlines enter through the switchyard wall (right). Turning mirrors direct each beam to the target chamber (below) where each beam is converted from its first-harmonic frequency to its second or third harmonic.

Nova's computer control network consists of fifty LSI-11/23 microcomputers linked with four VAX-11/780 computer systems. The control system was designed to satisfy control and data-acquisition functions in power conditioning, target and laser alignment, and target and laser diagnostics. One of the four operator consoles is shown below. Common hardware and software provide the flexibility to transfer operations and functions from one computer to another in the network.

In the small photo at lower right, a reticle is positioned in the target plane as part of target alignment procedures in the target chamber.

5

Laser Systems

Laser Systems

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Laser Systems

Introduction

R. J. Foley and J. T. Hunt

The major accomplishments of the Laser Systems Group in 1985 were completing activation of the Nova two-beam target chamber and developing the capability to do sophisticated, well-characterized target irradiation experiments in both the ten- and two-beam target chambers. The Nova laser system, depicted in Fig. 5-1, consists of the ten laser beams, the optical switchyard used for directing the beams and diagnosing beam quality, and the two target chambers.

During 1985 our main effort was on preparing the system for initial target campaigns and understanding the performance limitations of the laser and the frequency-conversion system. We demonstrated that a single amplifier chain could produce 15 kJ of 1.05- μm light. To quantify the factors that limit Nova's routine operation at high energy and high power, we set up experimental stations at several positions along the laser chain to conduct full- and partial-beam experiments. One of the stations was used to test newly developed platinum-free, neodymium-doped phosphate laser glass, which is required to achieve full potential from Nova on a routine, damage-free basis.

We conducted additional experiments to determine the cause of the lower-than-expected 2ω and 3ω frequency-conversion efficiency and to evaluate stimulated Raman scattering effects created by high laser beam fluences in long air paths. The results obtained from the experiments are presented in detail in this section.

After activating the integrated laser and target systems, we began target experiments in the Nova ten-beam chamber to extend our knowledge of inertial confinement fusion (ICF) physics and weapons-physics concepts.

On July 31, 1985, we delivered approximately 4 kJ of 2ω light in a 400-ps pulse to the two-beam chamber center in a line focus. X-ray laser experiments

Fig. 5-1.

will be conducted in this target chamber to determine the scaling parameters as drive laser power and line focus length are increased. The development of x-ray lasers with wavelengths below 100 Å will also be investigated.

Nova has routinely conducted three highly diagnosed target shots a day at energies of 20 kJ on target with 0.35- μm light. These experiments have enabled us to identify system modifications that will increase Nova's output energy to 50 kJ at 0.35 μm .

Nova Laser

Nova Laser Activation and Performance

On December 19, 1984, Nova delivered ten beams of frequency-converted 3ω light to the target chamber center, marking the successful completion of its construction. As expected, 1985 was a year of intense activity as we readied Nova for its first target campaigns and learned the performance capabilities of the laser and the frequency conversion system.

During the year, our primary emphasis was on developing the capability to do sophisticated, well-characterized target-irradiation experiments. Engineering construction continued in the target bay with the completion of the target vacuum system and controls and with the installation of target-diagnostic and target-alignment instrumentation. We performed target-irradiation campaigns that were carefully planned to contribute to the primary goals of activating the target diagnostic instrumentation and of understanding the operation of the laser system. We studied laser performance to gain a better understanding of the system's present energy and power limitations—an understanding that will enable us to increase output in the future and to routinely irradiate targets at Nova's full inherent power and energy capabilities.

Specific accomplishments during the year, with references to the articles where they are described in more detail, are given below.

1. We developed the laser target focusing capabilities on the ten-beam target chamber and determined typical beam characteristics at representative target planes near focus ("Opposed-Port Diagnostics" and "Target Plane Imager," Sec. 5).

2. We developed the target diagnostic capabilities on the ten-beam target chamber ("Activation of the Nova Ten-Beam Target Chamber for Target Experiments," Sec. 4).

3. We finished construction and activation of the two-beam target chamber and began experiments to continue our studies of x-ray lasers ("X-Ray Laser Experiments" and "Diagnostic Systems on the Two-Beam Chamber," Sec. 4 and "Two-Beam Target Chamber," Sec. 5).

4. We began simple target-irradiation experiments in the ten-beam target chamber to aid in activating the target diagnostic instruments and to produce publishable data to demonstrate Nova's capabilities ("Nova Experiments on Laser-Matter Interaction," Sec. 4).

5. We demonstrated that a Nova amplifier chain has the intrinsic capability of providing more than 15 kJ of 1.05- μm light to its frequency-conversion array.

6. We determined the factors that now limit unique Nova laser energy and power output and harmonic conversion efficiency ("Nova Laser Glass," and "Nonlinear Propagation and Frequency Conversion," Sec. 6).

7. We explored options for correcting Nova laser performance limitations and have begun to implement some of these corrections ("Plans to Replace Nova Laser Glass," "SRS Effects on Beam Propagation," and "Beam Depolarization Effects on Frequency Conversion," Sec. 6).

As the above list shows, we have progressed in all aspects of making Nova a unique experimental facility. We expect to be able to start the first series of sophisticated target experiments early in 1986.

Early in 1985, the Laser Systems Group explored the output capabilities of the Nova laser in both single-beam and full-system tests. These tests demonstrated that Nova has an intrinsic ability to exceed

our output goals at $1.05\text{ }\mu\text{m}$, and after we correct a small amount of depolarization in the output beam, should reach the upper limit of our energy goals at 0.53 and $0.35\text{ }\mu\text{m}$. In full-system tests, the maximum harmonic energy delivered to the target chamber center occurred in a shot in which 105 kJ in a 2.5-ns square pulse were incident on the KDP (potassium dihydrogen phosphate) arrays. This energy was converted to 2ω in the arrays and produced 44 kJ of $0.53\text{-}\mu\text{m}$ energy as measured by the calorimeters in the target chamber. Maximum power on target at 3ω was produced in the last quarter of the year in an experiment in which 60 kJ of $1.05\text{-}\mu\text{m}$ radiation was incident on the KDP arrays in a 1.0-ns square pulse; this provided 21 kJ of 3ω radiation in the target chamber.

We conducted tests with a single beamline to determine the upper limit of Nova's capabilities: A beamline delivered more than 15 kJ of $1.05\text{-}\mu\text{m}$ radiation in a 2.5-ns square pulse. In this series of experiments, we converted the laser output to its second and third harmonics in a KDP array mounted near the laser output in the switchyard and measured the converted energy directly in a large-aperture calorimeter. Figure 5-2 shows the results of this series of measurements. Up to 7.2 kJ of 2ω radiation and 7.8 kJ of 3ω radiation were delivered to the calorimeter. Higher power shots with shorter, 1-ns square

pulses produced up to 9.0 kJ output at $1.05\text{ }\mu\text{m}$. In these tests the conversion efficiency for harmonic generation in the KDP arrays was lower than expected, with the discrepancy increasing with increasing $1.05\text{-}\mu\text{m}$ power input to the KDP array. We conclude that the Nova amplifier system is intrinsically capable of providing more than 150 kJ of $1.05\text{-}\mu\text{m}$ radiation to its frequency-conversion arrays in a 2.5-ns or longer pulse, or more than 80 kJ in a 1-ns pulse. With present beam characteristics and at the longer pulse length, the frequency-conversion arrays should provide more than 50 kJ on target at either the second or third harmonic.

The damage that platinum inclusions in the glass cause to the laser disks presently prevents the laser from routinely operating at its full output capability. As discussed in "Nova Laser Glass," in Sec. 6, the threshold for damage initiated by platinum inclusions is equal to 3 J/cm^2 at 1 ns. To stay below the damage threshold, and thereby keep damage-induced, beam spatial noise at an acceptable level, requires restricting the output of an amplifier chain to less than 6 kJ in a 1-ns pulse. We intend to limit the output of Nova to this energy until we can replace the laser amplifier disks with disks made of platinum-free glass. With the current amplifier disks, the harmonic energy on target will be limited to less than 25 kJ. This is considerably below Nova's potential but does provide

Fig. 5-2. Harmonic output energy vs incident $1.05\text{-}\mu\text{m}$ drive energy. The KDP array was mounted at the output of beam 10 in the Nova switchyard.

adequate energy for meaningful initial target-irradiation experiments.

Early frequency-conversion experiments at 2ω and 3ω showed a lower-than-expected conversion efficiency [Fig. 5-3(a)] and for conversion to the third harmonic a pronounced decrease in conversion efficiency with increasing $1.05\text{-}\mu\text{m}$ drive power above 1.7 GW/cm^2 [Fig. 5-3(b)]. While searching for the cause of this low

Fig. 5-3. Conversion efficiencies from Nova beams at (a) 2ω and (b) 3ω vs the incident $1.05\text{-}\mu\text{m}$ intensity. The expected performances for Gaussian and flat temporal pulses are also shown.

efficiency, we discovered that the Nova output beam was partially depolarized. (The details of the depolarization are discussed in "Polarization Changes at High Intensity," Sec. 6.) The results of the depolarization alter the intensity mix ratio in the KDP crystals from the optimum. This causes a premature back-conversion of the harmonic wave, resulting in a lower high-power conversion efficiency to 2ω and a limit to the useful power that can be used to drive the KDP crystals for 3ω conversion. At energies above 6 kJ in a 1-ns pulse, further increases in the drive energy do not result in more 3ω output. Figure 5-4 illustrates the case for 3ω , showing the 3ω energy delivered to the target chamber from one of the Nova beams vs the 1ω drive to the conversion arrays.

We corrected the depolarization by repolarizing the beam at the output of the laser in both partial-aperture and full-aperture tests. This resulted in a greatly improved 3ω conversion efficiency. Figure 5-5 shows the results of subaperture shots done with and without beam repolarization. During these tests the frequency-conversion crystals were mounted in the Nova switchyard at the output of the laser chain. Similar results were obtained with the full-aperture beam and full Nova KDP arrays mounted in the switchyard and on the ten-beam target chamber. From this data we conclude that correcting the depolarization of the output beam will result in a substantial increase in the frequency-conversion efficiency of Nova beamlines. Both changing the laser glass to increase its damage threshold and correcting the depolarization are necessary to allow us to increase 3ω power on target.

A third effect that can limit the energy Nova delivers to a target is rotational stimulated Raman scattering (SRS) in N_2 in the beam air path to the target chamber. (A detailed summary of our data on SRS is given in "SRS Effects on Beam Propagation," Sec. 6.) SRS has its most serious adverse effect on high-power operation at pulse lengths longer than 0.5 ns and for the longest beam paths. Its effect can be eliminated by replacing the air in the beam pipes with a gas with a lower cross section for SRS, such as one of the rare gases. We have made a provision for doing this in the long propagation paths to the

two-beam target chamber, where SRS can have a serious effect at current operating levels. We are not yet sure whether replacing the air in the longest paths to the ten-beam chamber will be necessary for future higher-energy operation; however, such a change should be relatively simple to implement.

During 1985, our effort on Nova concentrated on installing and activating the system's target diagnostics and on exploring the present limits of the Nova laser amplifier and harmonic conversion performance. At year's end, the diagnostics were ready that will allow sophisticated target irradiation experiments to begin in early 1986. Several effects that limit laser and harmonic conversion performance—including glass damage from platinum inclusions, depolarization of the output beam, and rotational SRS—have been identified, and methods to correct these deficiencies will soon be implemented.

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Nova Optics

Introduction. At the start of 1985, the basic optics required for Nova's operation were complete. As a result, we concentrated our efforts this year on the following areas:

- Analysis of a crazing problem on the coatings of large polarizers.
- Initiation of a study to develop a sol-gel polarizer coating.
- Development of full-beam (74-cm) polarizers to evaluate beam depolarization problems, which result in a lower conversion efficiency for the KDP crystals.
- Design of a cylinder lens for the ten-beam chamber. Three different focal lengths were necessary to meet the requirement for a 2.5-cm line focus.
- Development of the Nova Optics and Components Management Net, a network designed to provide record-keeping functions for optics and assembly. The system is designed to handle inventory control,

Fig. 5-4. From Nova beam 9, the 3ω energy on target vs $1.05\text{-}\mu\text{m}$ drive energy for flat temporal pulses of 1-ns duration. At drive energies above 6 kJ, the 3ω energy does not increase.

Fig. 5-5. Tests done with 5- and 25-cm-diam segments of Nova beam 10 confirm that 3ω conversion will improve after repolarizing the output beam.

procurement tracking, and historical data analysis.

- Of particular significance was the fabrication, delivery, and installation of optical components for the two-beam target chamber. These components included focus lenses, mirrors, debris shields, cylinder lenses, and KDP crystals.

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Polarizer Coatings.

Crazing. In the first few months of 1985, we observed a modulation pattern in photographs of the output near field of several Nova beams. The modulation appeared as a full-aperture radial pattern (see Fig. 5-6), which we traced to the 20.8-cm polarizers. Close inspection of the coatings on the 20.8-cm polarizers revealed fine radial

cracks or craze lines. The crazing, not characteristic of normal laser-induced damage, was found on all 20.8-cm polarizers made by a particular manufacturer after May 1982. The parts had been damage

Fig. 5-6. Typical near-field radial pattern observed on Nova. The pattern was caused by crazing on 20.8-cm polarizers.

Fig. 5-7. Fused-silica polarizer (90 cm) used in conversion efficiency tests. The coating is on the beam output surface.

tested by the manufacturer. We then selectively verified that parts had thresholds less than 8.3 J/cm^2 . In addition, only parts exposed to laser radiation exhibited crazing; other polarizers from the same lot were not crazed after two years of storage. Polarizers coated by this manufacturer prior to 1982 were undamaged after tests were conducted under the same fluence conditions using an even greater number of shots than had caused the polarizers to craze.

A review of the coating process revealed that the vendor had changed the vapor deposition angle in early 1982 to improve the yield. The vendor suggested that the crazing may be stress-relief induced, so we examined the polarizers under $50\times$ magnification to examine the failure mode in more detail. The examination revealed right-angle crack intersections characteristic of stress-relief-induced failures. Scanning electron microscope (SEM) studies of polarizer coatings deposited on witness plates both before and after May 1982 are not conclusive but suggest a different titania (TiO_2) morphology for the two deposition configurations.

The coating manufacturer measured test witness pieces produced in 1985 coating runs, using both the pre- and post-1982 deposition configurations. Using interferometry, they conducted the tests on 1-in.-diam witnesses to determine the bending direction of the coated samples. The old process coating, which was stable, showed a compressive stress, while the new coating, which had crazed, showed a tensile stress. The vendor is now coating all polarizers using the old deposition-angle configuration. Any residual tensile stress remaining in polarizers can be reduced by an air bake at 200°C .

The 20.8-cm aperture parts were not the only polarizers coated after 1982. Our 31.5-cm aperture parts were also affected. Crazing was not evident until we removed the parts to inspect for a highly localized damage site. We saw "bullet hole" damage on the 31.5-cm polarizers from beams 6 and 8, as well as on their respective rotators. We carefully inspected those polarizers from beam 8, S/N 28 (input) and S/N 31 (output), which were both coated in June 1983. The region of elliptical damage was

more severe on the input part of the polarizer and is offset 45° on the second part, which corresponds to the polarizer offset.

The output polarizer on beams 6 and 8 showed crazing like that seen on the 20.8-cm polarizers. While searching for a reason for the crazing, we noted that the output side of the rotator assembly was in a dry nitrogen environment while the input side was in air. However, this same configuration existed for the 20.8-cm rotator assembly in which both polarizers had crazed. To study the effect of low humidity on this coating, we are humidity testing a witness piece from a 20.8-cm polarizer run containing parts that crazed. Results are forthcoming.

Full-Aperture Polarizers. Two major issues confronted us this year. The first dealt with the use of full-aperture polarizers to increase frequency conversion efficiency. The conversion efficiency of the large-aperture Nova crystal arrays was below the value previously obtained on smaller-aperture lasers. One possible cause, deconversion, could result from residual s-polarized light.

A second issue then confronted us: Should we continue to produce polarizers on BK-7 optical glass or could we produce them on fused silica? The damage threshold of BK-7 is lower than that of fused silica, but the low thermal expansion coefficient of fused silica may cause the polarizer coatings to craze. Fused silica has one-twentieth the expansion of BK-7 and could create a tensile stress in the polarizer coating. BK-7, on the other hand, exhibits damage thresholds that are at best marginal for this application.

To resolve these questions we coated two fused-silica, full-aperture polarizers (90 cm each) and installed them in beam 10, immediately in front of a Nova crystal array. Performance improved to approximately the level of the small-aperture conversion, suggesting that deconversion was responsible for causing the lower, less-than-anticipated harmonic efficiencies. However, the coatings were damaged, probably late in the tests during the 3-J/cm² shots. The damage existed over the beam aperture (see Fig. 5-7); it seemed independent of the orientation of the coated surface. One polarizer, S/N 103, was

Coating run	Substrate	Damage threshold (J/cm ²)	Design and process parameters
Test run #1	Fused-silica "ear"	6.3 ± 1.0 ^a	No barrier
		6.0	Normal temp
Test run #2	BK-7 (10 × 10 cm)	5.4	Barrier
	Fused silica	5.4	Normal temp
	Fused silica (5 cm)	6.6	
Test run #3	BK-7 (10 × 10 cm)	6.0	Barrier
	Fused silica (10 × 10 cm)	6.5	Reduced temp ^b

^a This measurement made on LLNL's absolute damage test laser.

^b Temperature about 50°C less than normal.

Table 5-1. Laser-damage thresholds of LLNL-supplied substrates measured at LLNL Comparative Damage Test Facility.

Fig. 5-8. Microscopic view of damage on 90-cm fused-silica polarizer. (A human hair, ~100 μm, is used for scale.)

positioned with the coating on its beam input surface; the second, S/N 102, was oriented with the coating on the beam output surface.

We are reviewing three properties of polarizer coatings on BK-7 and fused-silica substrates rejection ratio, damage threshold, and microscopic damage morphology. Our progress to date is outlined below.

- Values for T_p and T_s on the damaged full-aperture polarizer (transmission coefficients for p-polarized and s-polarized radiation) were measured to be 96.3% ± 1% and 0.24% ± 0.1%, respectively, yielding a rejection ratio of 400:1. The initial T_p and T_s values were 99.23% and 0.29% as measured by the vendor. The damage introduced a scattering loss of about 3%.

- Damage threshold measurements are shown in Table 5-1 for this coating and other polarizer coatings using barrier layers and reduced temperatures.

- We observed microscopic damage as 20- to 25-μm pits, uniform in size but nonuniformly distributed (see Fig. 5-8). This suggests that damage originated at the surface.

Polarizer Design and Tolerancing for Sol-Gel Coatings. We calculated a polarizer design and tolerancing information for developing future sol-gel coatings.

We recalculated the standard design to include index data more representative of the sol-gel materials, specifically $\text{SiO}_2 = 1.22$ and $\text{TiO}_2 = 1.85$. Typical refractive indices for these materials would be 1.45 and 2.2, respectively.

Figure 5-9 shows both the standard polarizer design and the modified standard design. Solid curves represent the standard design, which uses indices for evaporated TiO_2 and SiO_2 . Dashed curves represent the modified standard design, which uses sol-gel indices of 1.85 and 1.22 for TiO_2 and SiO_2 , respectively. The sol-gel design has a larger rejection width region, making it easier to build, and will therefore be studied for future applications.

Author: G. R. Wirtenson

Fig. 5-9. Comparison of standard polarizer coating design using typical vs sol-gel refractive indices.

Table 5-2. Original cylinder lens design for Nova ten-beam chamber.

Nova beam number	Focal length (m) ($\lambda = 0.63 \mu\text{m}$)	Number of lenses required	Line length (cm)
3, 4, 7, 8	-201	4	1.0
1, 2, 9, 10	-252	4	1.0
5, 6	-311	2	1.0

Table 5-3. Revised cylinder lens design for Nova ten-beam chamber.

Nova beam number	Focal length (m) ($\lambda = 0.63 \mu\text{m}$)	Number of lenses required	Line length (cm)
3, 4, 7, 8	-158	8	2.48
1, 2, 9, 10	-201	8	2.50
5, 6	-252	4	2.46

Cylinder Lens Design for the Ten-Beam Chamber.

Modifications to Baseline Design. In 1985 we modified the initial Nova ten-beam baseline design, increasing the overall maximum line length for each lens to 2.5 cm and incorporating a zoom capability similar to that of the two-beam chamber. The initial baseline design proposed a single cylinder lens, located inside the target chamber in the debris shield holder, to provide a fixed line length of 1.0 cm. Because of the 50° open-cone geometry, the design required cylindrical lenses of three different focal lengths. Table 5-2 lists the beam locations, the focal lengths for the required cylinder lenses, and the resulting line-focus lengths of the original ten-beam system. The cylinder lenses in each group of beam locations are tipped at specific angles, with respect to the optical axis of the Nova beam, so that the line foci lie flat on an x-ray target.

To provide the zoom line-focus capability, we have proposed two equal focal length cylinder lenses in a modified debris shield holder in each Nova beam. The cylinder lenses are counter-rotated to provide a variable line focus from 0 to 2.5 cm in each beam. The lens holder also provides a variable tilt angle to align the line focus onto the x-ray target.

Use of Two-Beam Cylinder Lenses in Ten-Beam Chamber. The Nova two-beam chamber currently uses six cylinder lenses (four plus two spares) with focal lengths of -158 m to generate a variable line focus nominally from 0 to 2.5 cm. We ordered a new set of six cylinder lenses to extend this capability from 0 to 5.0 cm, thereby making the original two-beam cylinder lenses available for use on the ten-beam chamber. Some of the cylinder lenses with focal lengths of -201 and -252 m have already been fabricated; therefore, the revised ten-beam design uses these new lenses plus the original two-beam lenses. Table 5-3 shows the revised design, which provides the zoom capability on the ten-beam chamber.

Note that the two-beam lenses are used in the first set of locations and that those lenses originally designed for the first and second locations are now used in the second and third locations. This scheme reduces the number of new cylinder lenses

that we must purchase by six, reducing costs significantly.

Author: L. G. Seppala

Nova Optics and Components Management Net. The Nova Optics and Components Management Net is a local area network designed to replace or supplement many record-keeping functions that are now done manually. The system will provide three services: procurement tracking, inventory control, and historical data record keeping and analysis. While such records are now kept in different locations, the Management Net will coalesce the data into a single set that is distributed among the terminals on the Net and is made available to other users on demand.

Each terminal, connected to others in the Net by coaxial cable, has a 500-kbyte memory and a 30-Mbyte disk. A system of five terminals is planned for the optics, clean room, and electrical areas. With a transfer rate along the Net of 2.5×10^6 bits/s, data is quickly available to other terminals on demand.

The Net's procurement tracking service will start when a requisition is issued. The purchase order may be interrogated at any time and updated as items are received, rejected, or cancelled. When an optic is delivered to the Nova organization from a vendor or from somewhere else on site, it moves into the inventory system and receives a serial number. Thereafter, every move or change of condition is recorded. The serial-numbered optic is tracked as it moves from the Optics Depot, to a position in a serial-numbered frame assembly in the clean room, and eventually to an arm on the laser. Thus, a continuous bill of materials is maintained for the Nova laser, which may be called up on the system at any point in time.

The first service to be offered in the area of history data and analysis will be that of assisting in analyzing damage to an optic. Relevant data from the shot record will be maintained on the network. If an optic is damaged, inventory control data and the history data will be combined to access the data for every shot imposed on the glass. A three-dimensional computer-aided design facility will aid in analyzing the damage. Each terminal has a floating point processor to support any numerical analysis required.

The first phase of the project, the inventory control of serial-numbered optics, is complete and operational. Lists of parts, vendors, the locations of parts in the system, and the conditions of the parts have been established. All of the data from the Optics Depot files have been entered into the system. Inventories were conducted at all locations given by the files and were printed. A physical inventory was then conducted to confirm the accuracy of the computerized inventory. We are now preparing the clean room for its network software.

Author: R. J. Foley

Beam Alignment and Diagnostics

Introduction. Nova was in operation for most of 1985. Early in the year emphasis was on construction and activation, but as the year progressed we began to test target diagnostics and to irradiate targets. We continued to conduct a number of laser alignment and diagnostic activities to more fully develop the potential of the alignment and diagnostic systems.

This year we extended alignment control capabilities to include algorithms for centering and pointing the 1ω alignment beam with respect to the KDP array. In addition, we incorporated KDP tuning in the automatic alignment program, and are using digital filters in the image-analysis software to effectively eliminate fixed-pattern noise from the output charge-injection device (CID) cameras.

All of the output-sensor reflected diagnostic components have been positioned and adjusted, and measurements of light reflected from the target within an angle of 1.5 mrad from the beam axis are now performed regularly. Calibration of the reflected-light measurements is accomplished by collecting the light reflected from an uncoated optic across the chamber from each beam.

Our 5-cm absorbing-glass calorimeters are now traceable to the National Bureau of Standards (NBS). The electrical calibration values we obtained agreed to within 0.5% with NBS optical calibration values for pulse durations up to 6 s. We are doing our own optical calibrations of large-diameter

calorimeters, both at atmospheric pressure and under vacuum, using an industrial YAG laser as the source and a 5-cm-diam calorimeter as the reference.

The Nova system operation generates a large volume of data. By the end of 1985, we had stored approximately 40 000 measurements. The ORACLE data base package provides a mechanism for relating and comparing data from different shots. For example, system operators regularly construct plots of performance parameters vs time.

Fiber optics provide a very versatile way to multiplex signals from several sources onto a single streak camera. However, standard fibers exhibit enough dispersion at short wavelengths to cause unacceptable pulse distortion. Measurements on a special fiber have shown much less dispersion but higher transmission loss.

The remainder of this subsection describes the activities of the year in greater detail. We anticipate that as Nova matures and as the purposes for which it is utilized evolve, additions and improvements will continue in the alignment and diagnostic area.

Author: E. S. Bliss

Closed Loop Alignment. Work on the closed loop alignment system during 1985 involved extending automatic alignment to the ten-beam target chamber. We designed and implemented algorithms for pointing and centering the 1ω beam to the target chamber. We also added the capability to scan KDP crystals to obtain optimum 2ω conversion efficiency. The increasing amount of command and parametric information needed to characterize approximately 300 loops required that we modify the methods used to access this data. (A detailed description of the digital image processing used in the Nova alignment system is presented in a separate report.¹⁾

New Automatic Functions. Final alignment of the main beam to the target chamber is measured by viewing the KDP crystal at the target chamber with the output sensor. The KDP frequency converter is an array of nine crystal slabs. Because the cluster image is complex, the simple centroid algorithm used to make most

pointing measurements is unsuitable for determining the location of the cluster. Instead, pointing measurements are made by viewing, in the far field, the cluster of spots reflected from the nine crystal surfaces. Histogram equalization is applied to the image to increase the contrast of the dimmer spots, thereby increasing and making more consistent the apparent structure of the cluster. A reference image, or template, is then correlated to the sensor image to locate the position of the cluster. The reference to which the KDP cluster is aligned is found by locating the reflection of the corner cube instead of the camera reticle.

Centering of the main beam to the target chamber is measured by comparing a preceding cross hair to the shadow cast by the KDP array mounting structure [Fig. 5-10(a)]. To observe the shadow, the centering mirror is inserted behind the array to direct the light back to the output sensor. The image is then inverted to make the shadow lines bright. Because the dark background on the image exterior inverts to white and the large white area will compete for best correlation, the exterior of the image is zeroed. The alignment reference is the center of the parallelogram formed by the diagonal KDP mounting bars. This structure is also located using the template-matching technique. The diagonal lines have a width approximately equal to the finest resolution of one pixel. In this case, the image cannot be low-pass-filtered to attenuate noise. The lines are best enhanced by calculating the localized mean value of the non-zeroed portion and subtracting it from the image. Normally, the mean-value calculation is performed globally on the entire image.

The artificial template used in the template-matching process is made by generating four intersecting lines that are equivalent to those formed by the mounting bars [Fig. 5-10(b)]. Extending the lines to the edge of the image field improves the sharpness of the match. For this type of image, employing thin lines to form the template works better than using a template prepared from the original image.

This year, we integrated KDP crystal array alignment for 2ω operation, which is described in the 1984 Laser Annual,² into

Fig. 5-10. (a) Sensor TV image of the KDP mounting structure. The diagonal mounting bars form the alignment reference, which is the center of the parallelogram. (b) The artificial template formed by generating four intersecting lines that are equivalent to those formed by the mounting bars.

the automatic alignment program to provide control sequencing and configuration definitions for component selection, device reservations, and task parallelism. Previous attempts to fit the scan data to the expected sinc-squared function failed because of large errors introduced by random noise. To process the noisy data, we used a new curve-fitting routine based on a modification of the Levelberg-Marquardt algorithm.³

Improvements in Processing Dim Images. Some images viewed by the output sensor had marginal signal-to-noise levels when the continuous wave beam was fully distributed ten ways. Images of pinholes and images reflected from the KDP array and centering mirror were particularly affected. Because of this, we studied and improved image acquisition from the CID cameras to permit reliable automatic alignment of all beams in parallel. We created a menu-selected toolkit of image-processing functions, which has proved beneficial for investigating these improvements and for testing our latest algorithm designs.

Several of the image-processing algorithms, such as the centroid method for locating focused spots, require as a processing step the application of a fixed threshold function to the image. Images are now normalized to the full 8-bit gray scale. The threshold function can then be applied relative to the white level instead of to the black level. This technique yields the consistent after-threshold image shapes necessary for uniform processing results.

The CID camera, while eliminating the severe problem of blooming that is experienced

with all charge-coupled device (CCD) cameras, posed a different challenge in the processing of dim images. The CID generates a fixed-pattern background noise in the video signal, which introduces a significant spatial aliasing problem when digitized by the video-processing system. Appearing in the image as fine vertical lines with a peak intensity of about 10% of full-scale brightness, this noise most seriously affects algorithms that are sensitive to simple intensity or to vertically oriented picture information.

The alias signal is not a problem in many analyses (such as most cross-hair processing) because the signal is not strongly correlated to the image information of interest. Images such as those obtained for pinhole alignment have been improved by median low-pass filtering. The median filter, in contrast to low-pass filters of finite-impulse response, not only attenuates the higher spatial frequencies but preserves image edges quite well. Analysis of the alias signal shows it is composed of a single spatial frequency that occupies a part of the spectrum not important to the accurate representation of typical alignment-image information. Because of this, a band-stop filter can be used to attenuate the alias signal without significant loss of image resolution. Furthermore, the signal appears in only the horizontal image dimension, so that a more simple one-dimensional digital filter can be employed.

Future Efforts. In 1986, we will concentrate on completing the capabilities of the closed loop system as operational

experience dictates. Candidates for this extension include adding harmonic alignment of beams on target, KDP scans for 3ω operation, and closed loop operation on the two-beam chamber.

Authors: P. J. Van Arsdall and E. S. Bliss

Major Contributor: R. J. Reeves

Reflected-Beam Diagnostics. This year we activated the reflected-beam diagnostics on the even-numbered (east cluster) Nova beams. Light reflected from the target chamber is collected by the focusing lens and travels back along the incident beam path to the output sensor (Fig. 5-11). Diagnostic measurements of reflected energy and the temporal pulse shape of the reflected beam at 2ω and 3ω are performed in the splitter module of the output sensor, which is shown in Fig. 5-12. The splitter-module configuration allows for reflected-beam diagnostics with an approximate field of view of ± 1.5 mrad. This corresponds to a 9-mm-diam circle at target chamber center. (Operation of the 1ω module, which includes 1ω reflected diagnostics, is described in detail in the 1980 and 1982 Laser Annuals.^{4,5})

Turning mirrors for 2ω and 3ω operations are coated such that by the time the 3ω light reaches the diagnostic turning

mirror its intensity is low enough that the BK-7 substrate of the mirror will not be solarized. The reflected light passes through the diagnostic turning mirror, reflects off the auxiliary mirror, and enters the input telescope of the output sensor. For 2ω shots, a full-aperture beam dump is inserted between the auxiliary mirror and diagnostic turning mirror, reducing the beam energy to prevent damage to optics in the output sensor.

Upon entering the splitter module, shown in Fig. 5-12, the 2ω or 3ω light is reflected from the reflected-beam splitter into the reflected diagnostic path of the module. (The 1ω reflected light would travel directly through this splitter and into the 1ω module.) Because some optics in the splitter module are different for 2ω and 3ω , they must be manually installed for the desired wavelength. For 3ω operation, two lenses are required: the first lens recollimates the 3ω light from the input telescope, and the second focuses the beam at the correct location for the streak-camera fiber. For 2ω , only one lens is required to focus the beam in the correct location, but because of the intensity of the beam, an additional attenuating filter must be installed in the beam path to lower the beam energy to an acceptable level.

For both 2ω and 3ω operation, the

Fig. 5-11. Nova output diagnostic system. Light reflected from the target chamber travels back to the output sensor for measurement. A fused-silica reflector is used to provide a reference signal for calibration.

reflected-energy measurement is made by an absorbing glass calorimeter located behind the calorimeter beam splitter, and the signal used for the pulse length measurement is directed to a fiber optic for transmission to a centrally located streak camera. A filter wheel, located before the fiber optic, allows computer-controlled adjustment of the streak-camera signal level.

On previous lasers with opposing-beam geometries, reflected-beam diagnostics on one arm have been calibrated using the incident diagnostics of the opposing arm. By firing the opposing beam at reduced energy and setting the two focus lenses to be confocal, the incident energy of one beam will travel down its opposing beam as a known amount of "reflected energy." This method of calibration depends on the accuracy of the incident beam diagnostics. Nova, with its non-opposing beam geometry, does not allow the calibration of reflected-beam diagnostics in this manner.

To supply an equivalent diagnostic calibration source, an uncoated fused-silica reflector has been designed to be mounted

on the 44-cm calorimeters in the target chamber. Both surfaces of this fused-silica reflector share the same center point of curvature. If this point is placed at the focus of the focus lens, the Fresnel reflections from the front and rear surfaces will return along the incident beam path and can be used as a calibration source for the reflected-beam diagnostic package. The light that is transmitted through the reflector (92.7% at 3ω) to the 44-cm absorbing glass calorimeter is used as the reference for the incident beam energy. This method of calibration gives an accurate measure of reflected energy and safely calibrates the reflected diagnostics without the risk of propagating a high-energy laser pulse back down the beamline.

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Major Contributors: C. E. Annese,
L. G. Seppala, and B. D. Cline

Calorimeter Calibration. We have established an optical calibration facility

Fig. 5-12. Reflected 2ω and 3ω diagnostic measurements are made in the splitter module of the output sensor.

Fig. 5-13. Our new facility for optically calibrating large-aperture calorimeters uses a pulsed Nd:YAG industrial welding laser with an output of 250 W.

(Fig. 5-13) in which we can calibrate our large-aperture (34-, 44-, and 74-cm diam) calorimeters against a 5-cm-diam reference calorimeter.

Our reference calorimeter uses absorbing glass with a 5-cm clear aperture. It is designed to be calibrated by passing a known amount of electrical energy into a built-in heating element.⁶ The calorimeter sensitivities we obtained by electrical calibration agreed to within 0.5% with these values obtained in recent NBS calibrations.⁷

While electrical calibration is accurate for small-aperture calorimeters, it has proven unreliable for our large-aperture calorimeters. The large-aperture calorimeters are designed to measure the 10-kJ output energy of a Nova amplifier chain. Since the Nova facility is not a feasible calibration source and no other accessible 1- μ m, 1-ns laser is powerful enough to be of use, we have chosen to base our calibration facility on a pulsed Nd:YAG industrial welding laser with an output of 250 W at 1.064 μ m. It takes this laser 4 s (132 pulses) to generate the 1 kJ of output energy needed to calibrate our calorimeters.

To accurately measure the energy being deposited in a large-aperture calorimeter, a 5-cm calorimeter measures approximately 1% of the beam that passes through a partially transmitting mirror. Measurements before and after each calibration sequence enable us to determine the entire laser output energy. The 4-s optical heating time is acceptable since large-aperture calorimeters respond slowly, requiring 20 to 30 s to achieve maximum output. The 5-cm calorimeter samples the same 4-s gate of pulses from the industrial laser. Because of its smaller size, it responds more quickly than the large calorimeter. NBS measurements showed that the 5-cm calorimeters remained accurate to within 0.5% for optical irradiation times of up to 6 s.

We recently added a vacuum vessel to the calibration facility to test the 44-cm calorimeters that are used inside the Nova target chamber. This enables us to observe and document, off-line, the effects of vacuum on these calorimeters. All calibrations can be made under similar conditions that exist when the calorimeter is used on Nova. Overall, the calibration facility has proven

to be a valuable means of quickly and accurately verifying calorimeter sensitivities.

Authors: J. B. Richards and D. W. Myers

Major Contributors: J. A. Prior and B. D. Cline

Laser Diagnostic Data Base. For each Nova system shot, the laser diagnostic system records energy data from several hundred sensors distributed along the ten laser chains. The reduced data is archived in large data files on VAX computers. The archival file format is difficult to access by operators and experimenters and provides no general mechanism for relating information gathered from several shots. Therefore, the ORACLE relational data base package, already in use at the Laser Experiments Analysis Facility (LEAF) to store physics data, is being used to store and access this laser diagnostic shot data.

Figure 5-14 shows a block diagram of the laser diagnostic data base. Energies are measured by both sensors and sensor calibrating instruments. Shot data is stored in two ORACLE tables of identical format: the shot data table and the calibration data table. Shot number, sensor location,

hardware information, sensor parameters, measured voltages, gains, and calculated energies are stored as a separate table entry for each sensor. The largest amount of data is actual sensor data; it is stored in the shot data table. The rest of the data is maintained in the calibration table. A third table, shot description, defines each shot by recording the laser wavelength, wave plate settings, and shot date and time. An additional entry names the shot definition data base record, which is described in the article "Shot Definition File," later in this section. The shot definition details the laser and diagnostics configuration for the shot on a beam-by-beam basis.

The laser diagnostic data base is updated semi-automatically the morning after a shot. The data utilities on the laser diagnostics VAX read each stored shot file into the global section memory and translate the data format into a compact data file that is compatible with ORACLE. The positions of the wave plates, recorded by the alignment system for each shot, are also transferred. DECNET is then used to transmit the compact file to the development VAX where ORACLE is installed. At the same time, the compact data file is sent over the one-way link to the LEAF

Fig. 5-14. Laser diagnostic data base.

Fig. 5-15. Layout of the optics used to measure temporal dispersion of optical fibers allows selection of several path lengths; the optical path shown is 41 m long.

VAX where it is later read into the LEAF target diagnostics ORACLE system. The compact file is loaded into a temporary ORACLE table and then sorted into the correct data base tables, as described above. A history file of shot data transactions records all semi-automatic activity with the data base.

The data base operator selects utility functions from a menu displayed on the computer terminal (shown as "user functions" in Fig. 5-14). Data may be queried, entered, or modified using convenient fill-in-the-blank forms. In addition, summary reports can be prepared using the relational capabilities of ORACLE. The summary report can be in the form of a single shot report, a shot history report for one sensor, location, or a calculation of sensor calibrations. Data necessary for calculating sensor gains and energy balance can be written to intermediate files.

Two utilities use the data stored in ORACLE to calculate system parameters needed for new shot setups. Sensor gains,

which are loaded into the diagnostic energy processor before each shot, must be optimized so that the best signal-to-noise performance is realized and out-of-range conditions are avoided. The ORACLE data is selected from a previous, similar shot. This is a convergent process, in which new gain calculations improve our initial best-guess estimates. A second utility accesses shot energies at critical points in the laser chains, along with associated wave plate positions, to prepare input data for the energy balance software described in the article "Energy-Balance Software," later in this section.

All shot data from 1985 as well as important calibration data from 1984 have been inserted into the data base. This data, now representing approximately 40 000 sensor records, is accessed routinely by several members of the operations staff who are proficient with ORACLE programming. For instance, laser amplifier performance over time is periodically checked. The sorted data output by the

data base is transferred to the Tell-a-graph graphics package for plotting. Following these procedures is well within the capability of many staff members who are not regular computer users.

Author: P. J. Van Arsdaal

Wide Bandwidth Ultraviolet Fiber

Optics. Our goal is to measure the temporal pulse shape of each individual Nova beam at all wavelengths generated by the laser: 1.06, 0.53, and 0.351 μm .⁸ Making this goal realizable in a cost-effective manner, however, relies on multiplexing samples from several beams onto one centrally located streak camera to minimize the number of these very expensive instruments in the system.⁹ From a system's viewpoint, multiplexing must be convenient, reliable, and inexpensive, both in implementation and operation. We also require that the transmitted signals be faithful reproductions of the input for pulse widths as short as 100 ps.

A commercial fiber-optic cable, Siecor 104, provides these characteristics at 1ω and 2ω . We have implemented the cable effectively, transporting the optical beam samples to our streak cameras in Nova. At 3ω , Siecor 104 unacceptably broadens input pulses with full widths at half maximum (FWHM) of about 100 ps. After failing to find a commercial fiber for this 3ω application, we contacted AT&T Research Laboratories, Murray Hill, New Jersey, for assistance and acquired a sample of an experimental fiber that they thought might possess the required 3ω bandwidth.

We made dispersion measurements on both the Siecor and Bell fibers using the intermediate laser system (ILS), which produces 1ω , 2ω , and 3ω , 100-ps FWHM pulses at energies greater than a few millijoules. Experimental apparatus, illustrated in Fig. 5-15, allowed us to compare the fiber output signal with an appropriately delayed sample of the input, which propagated only over an air path. Four percent of the incident pulse was extracted at clear-fused-quartz (CFQ) splitter SP2 and delayed via an optical path made up of mirrors M3 through M8. By varying the number of reflections made between mirrors M4 and M7, we produced air-path

lengths of approximately 27, 41, and 54 m. A 13-m path was obtained using mirrors M3, M7, and M8. We cut the fibers to precisely match the air-path optical lengths. Both air and fiber signals were recorded simultaneously by a streak camera having a total recording time of 1.5 ns and a temporal resolution of better than 10 ps. Simultaneous recording was necessary because of the significant variability of the laser pulse from shot to shot.

We used bandpass and color filters to select wavelengths and neutral-density filters to control intensity. CFQ ground-surface diffusers at the input to the fiber and at the input to the streak camera for the air-path signal provided a spatial average of the entire beam. We prepared each fiber end by first stripping off the buffer material and then cleaving the fiber to produce a flat surface. The fiber output end was then inserted into the slit of the streak camera.

We measured the dispersion and velocity of propagation of both fibers at all three Nova wavelengths. Table 5-4 summarizes the fiber parameters and the velocity of propagation. Figure 5-16 shows the measured 3ω dispersion. The Siecor data represent the true fiber response, 6.2 ps/m, as calculated by

$$\tau_r = (\tau_o^2 - \tau_a^2 - \tau_{sc}^2)^{1/2}, \quad (1)$$

Table 5-4. Summary of Siecor 104 and Bell Labs optical fiber parameters.

Parameter	Siecor 104	Bell Labs
Core/cladding diameter	50/125 μm	50/125 μm
Dopant/intrinsic material	GeO ₂ /SiO ₂	SiF ₄ /SiO ₂
Numerical aperture	0.2	0.1
Index profile	Graded	Graded
Velocity of propagation		
1ω	0.675 ft/ns	0.683 ft/ns
2ω	0.665 ft/ns	0.672 ft/ns
3ω	0.642 ft/ns	0.651 ft/ns

Fig. 5-16. Optical fiber dispersion vs distance at 3ω .

where τ_f = fiber response, τ_o = FWHM of fiber output, τ_a = FWHM of air path (i.e., fiber input), and τ_{sc} = streak-camera response. The streak-camera response time had previously been measured to be approximately 10 ps (Ref. 10). Since the output pulses of the Bell fiber were generally shorter than the input, the data shown are the pulse-width difference ($\tau_o - \tau_a$). Neither fiber exhibited measurable dispersion at 1ω or 2ω , these results being similar to those of the Bell fiber at 3ω . Our measurements were made with fiber lengths equivalent to the four optical air paths for the Siecor fiber and equivalent to the 27- and 41-m air paths for the Bell fiber.

We believe that the source of the observed temporal dispersion is a modal effect rather than the result of nonlinear processes that could shift the transmitted wavelength and thereby cause pulse dispersion. Modal dispersion occurs because different fiber modes propagate at different velocities. Within the range of intensities available to us using the ILS laser, we also

found that neither fiber exhibited any measurable nonlinear dispersion at 3ω when we increased the input intensity by more than a factor of 10 000. These data indicate that the Bell fiber meets the 3ω dispersion requirements.

When making these measurements, we observed that for the same fiber length much more input intensity was required for the Bell fiber than for the Siecor fiber to produce an equivalent signal at the streak camera. Some of this loss was due to a reduced collection angle (i.e., numerical aperture) and was regained by more careful alignment of the fiber input. Nonetheless, at the 41-m air path length, we found it necessary to focus the laser directly onto the Bell fiber end without using the scatter plate; we were unable to make any measurement at the 54-m length. This suggested a need to make transmission measurements for the different fibers.

We measured the near 3ω ultraviolet (UV) transmission characteristics of both fibers using the synchronous detection

Fig. 5-17. Experimental arrangement for making UV ($\sim 3\omega$) optical fiber attenuation measurements.

scheme shown in Fig. 5-17. The argon-ion laser ran at a multi-line output of 50 mW between 0.3511 and 0.3638 μm . A CFQ scatter plate provided a diffuse source for coupling into the fiber. We carefully chose the distance between the fiber input and diffuser so that the UV spot at the diffuser was smaller than the projected collection spot of either fiber. We aligned the fiber input to the UV spot by projecting an He-Ne signal backwards through the fiber and then centering the two spots. The distance between the fiber output and detector PD2 was chosen so that PD2 would collect all of the signal from either fiber. The higher numerical aperture of the Siecior fiber meant that the fiber was somewhat underfilled. Large area silicon photodiodes were used for all three detectors.

Our technique for measuring transmission involved fixing both the fiber position and laser power level at the input and incrementally cutting back the output end. At each length we measured the transmitted signal, cutting off a few millimeters several times at the major incremental lengths to account for variations caused by imperfect end terminations and positioning. The results were averaged for each major length.

These results (Fig. 5-18) show that the Bell fiber had an attenuation coefficient nearly four times greater than the Siecior fiber, 25.5%/m vs 6.55%/m. This poor transmission clearly limits the length of the Bell fiber that can be conveniently used in the Nova system. The index profile of the Bell fiber is much shallower than that of the Siecior, and the width of its low-index cladding zone is much narrower. These two factors may produce the observed loss for the Bell fiber, rather than optical absorption, which we believe is insignificant. A radial loss out the side of the fiber can easily be seen for the Bell fiber but not for the Siecior when propagating a visible signal through the fiber. The Bell fiber side loss is extremely sensitive to bends with radii less than 8 cm and to the slightest pressure imposed upon the fiber. In both cases, the observed sidescatter is always in the forward direction, indicating that small-angle scattering is predominant.

Fortunately, this loss is not so large as to make the Bell fiber unacceptable for use on Nova. Side losses due to micro-bends

Fig. 5-18. Optical fiber transmission vs distance at 3ω .

will be reduced when the Bell fiber is in cabled form. Additional care in handling and system configuration, however, will still be required to prevent excess loss. Bending does not appear to limit the fiber bandwidth. Even with these limitations, we predict that the Bell fiber will be usable for 3ω measurements on Nova for lengths up to 20 m.

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Evolution and New Capabilities of the Nova Control System

Introduction. We are continually developing and improving the capabilities of the Nova control system to support the number and variety of experiments being conducted on Nova. We extended experiments and operations in two fundamental ways in 1985.

First, a shift in emphasis from laser fusion research to weapons-physics research led to new experiments that support studies on laser-matter interaction, the x-ray laser, and the effect of long-path transmission of high-intensity laser beams through air (which is hindered by stimulated Raman scattering). This led to the addition of a two-beam target chamber and the development of new diagnostic techniques.

The second way in which we extended experiments and operations involves the

extensive use of the laser to obtain data on its own performance. This enables us to better understand the cause of severe optical damage at high energies (arising from platinum inclusions) and low frequency-conversion efficiency (affected by stress-induced birefringence, which causes depolarization). Often entire beamlines were dedicated to laser research, while other beamlines were used for target experiments into the same target chamber on the same shot. Setup and calibration efforts proceeded concurrently in the other chamber.

Continued research in laser technology has extended the performance options of Nova. Pulse shaping and other improvements have been made in the MOR; new target-plane imaging techniques are being employed; and laser diagnostic sensors, as well as a second control room in a secured area for the alignment of classified targets, have been added to the Nova facility.

Each of these improvements depended on new developments within the Nova control system. For example, to operate the second chamber and provide operations from the second control room, we had to replicate existing ten-beam control system hardware and modify control system software. In particular, we extended the capabilities of the alignment subsystem and operator console graphics software.

The variety of target experiments in both chambers and the concurrent experiments to understand laser system performance required frequent changes in the laser configuration. Therefore, overall coordination and automatic reconfiguration of the facility became important. Through these capabilities the control system helps to prevent operational errors that may damage expensive and long-lead-time laser components or waste valuable system operational time.

Several other technological advances depended on the development of new capabilities within the Nova control system. When we began to use different laser chains to support more than one experiment on a single shot, schedule changes became a daily occurrence. Therefore, we developed energy-balance software that would automatically optimize the performance of each laser chain. In addition, new and improved diagnostic instruments depended on higher-resolution charge-

coupled device (CCD) cameras, which in turn required new control and software capabilities. Finally, poor laser performance led to the development of a facility-wide system that monitors air temperature and pressure in the laser bay.

The capabilities of most of our machines and equipment (VAX and LSI-11 computers, networks, disk units) are fully taxed. As a result, we have been adding equipment (more memory and larger disk units) and optimizing software. Praxis language capabilities have also been improved to optimize our machine resources and thereby increase system performance.

Author: F. W. Holloway

Controls Integration. The Nova control system is composed of many subsystems and integrating software.¹¹ This year, for the first time since the control system was activated, a few functions of each subsystem were coordinated and monitored automatically through the integration software. (See the 1984 Laser Annual¹² for a description of the software's general capabilities.) Subsystem software was interfaced to the two major components of the integration system: the shot scheduler and the final countdown software. This interface allowed a number of intersubsystem procedures to be defined and activated.

Shot Scheduler. The shot scheduler provides a mechanism through which any ordered list of subsystem-supported control functions may be defined and performed. Such a list is referred to as an automated checklist. An operator interacts with the shot scheduler through a console display, such as that shown in Fig. 5-19. Some of the subsystem functions that have been synchronized by the shot scheduler are:

- Selecting the alignment device configurations required for different modes of laser alignment (i.e., rod shots, system shots, and automatic alignment).
- Commencing laser-energy-data acquisition by the laser diagnostics subsystem.
- Initiating the selection of laser amplifiers and rotators to be fired by the power-conditioning subsystem.
- Starting and stopping automatic alignment.

Checklists composed of these and other functions are providing operators with reliable and repeatable high-level control of

Fig. 5-19. Console display for controlling and monitoring checklists. Manual controls are provided by the touch-sensitive buttons at the bottom of the display.

frequently performed activities. A prototype "shot checklist," composed of a number of functions that must be performed before every shot, has been developed and successfully tested.

Final Countdown. The final countdown software checks overall subsystem readiness during the few minutes before a laser shot. During the final countdown, subsystems receive data from integration software on the progress of the shot and can submit responses that will suspend it. A log of subsystem readiness is available to operators on a console during the countdown (see Fig. 5-20). Operators can also suspend the countdown, override the suspension, and abort and restart the countdown from another console display.

A few of the subsystem functions synchronized during the final countdown are:

- Initiating and suspending command sequences for charging and firing capacitor banks by the power-conditioning subsystem.
- Starting and monitoring pre-shot configuration of target diagnostic data-acquisition systems.
- Beginning laser streak-camera data acquisition and analysis following a shot.
- Acquiring laser and target-bay humidity and temperature data.

The controls integration software decreases the time required to configure the laser for shots, and its checking procedures help to protect the laser. Several modifications can be made to improve the impact of the integration software on the operation of the laser. For example, an increase in the number of integrated subsystem functions will greatly improve the reliability and efficiency of the control system. Moreover, we can quicken the software's response time to subsystem requests, increase the amount of detailed subsystem status data available to console operators, and simplify the task of implementing or modifying shot scheduler checklists. We can further improve the flexibility of the integration software by inputting shot definition data to the shot scheduler. (See the article "Shot Definition File," later in this section.) The shot scheduler must first be modified to acquire information from the data base and transfer it to subsystems. Because the data base is designed to maintain only minimum specifications, such as laser output energies and pulse parameters, subsystem software must be upgraded to translate shot definition data into tasks that are compatible with existing subsystem function definitions.

To compensate for the variation in amplification of each laser chain, the operator

selects as a reference a previous laser shot whose pulse length and energy are similar to those desired on the next shot. The energy-balance software extracts data from the shot data base on the energy obtained for the reference shot at five points along each of the ten laser chains. The performance desired during the next shot is then used to solve a set of simultaneous equations¹³ to determine the fraction of MOR laser energy that must be distributed into each laser chain to obtain the desired output energy distribution. Then the wave plate cosine-squared functions are solved to obtain the necessary mechanical settings of the wave plates, and a file is prepared that, at the operator's request, is used by the alignment system to automatically move the beam splitter wave plates to the necessary positions.

This software can control the amount of energy injected into each of the laser chains so that the desired output energy balance is obtained. The entire operation, which originally took long periods of time and was error-prone, now takes only a few seconds and is reliable.

Author: F. W. Holloway

Shot Definition File. A control system file has been designed that contains the parametric and descriptive data necessary to characterize complex system shots for concurrent Nova ten- and two-beam activities. The file, which consists of more than 100 pieces of information, is maintained in the ORACLE relational data base. Using the forms capability of the data base, the experimenter enters into the system the data needed by the control system software for a desired shot. The shot definition, which gives pre-shot information, is later combined with the post-shot data to provide the experimenter with a characterization of the shot. (See also "Laser Diagnostic Data Base," earlier in this section.)

The shot definition describes the oscillator pulse, target chamber contents and diagnostics, and the configuration of each beamline. The pulse type, duration, and picket-fence parameters are included for the oscillator subsystem. The function of each beamline, whether for a target shot or some type of calibration, is described.

Included in that description are the target chamber or sensor destination, frequency of laser light, and such target alignment coordinates as focus spot size and location. For the two-beam chamber, the line-focus height and width are specified.

The forms capability of ORACLE provides protection mechanisms for limiting the data values to those that are recognized by the control system. For instance, it verifies that character fields such as the shot type and destination exist as entries in a table of allowable choices. Numeric data are checked for allowable range and precision. Units of measure are specified on the form. These features ensure that the data are consistently entered, even though the definitions may be created by different experimenters. In this manner, the validity of the data is checked before it is needed by the control system.

A menu-operated command utility, similar to the one used to access the laser diagnostics data base, activates the ORACLE forms and prepares the shot definition file. The file may be printed and appended to the shot notebook for use by operating personnel during shot setup. However, other work must be done before we can load the shot file data into the computer systems that are responsible for carrying out the shot event automatically. In addition, each of the control subsystems must provide software to use this data in the individual subsystem shot configuration procedures.

Author: P. J. Van Arsdaal

Major Contributor: F. W. Holloway

Console Alignment. This year we extended the console alignment program (CAP)¹⁴ to include the ten-beam target chamber controls. In addition, we added a major new software product, the alignment dispatcher program, to the Nova alignment system. The dispatcher controls the semi-automatic setup of all alignment devices.

The system configuration (i.e., the commands to stepping motors and binary channels) is defined in a set of easily edited text files. Therefore, the system operators (who are nonprogrammers) can change any of the configuration specifications quickly without having to recompile

the software. In addition, the time required to configure the alignment system has decreased, because the shot scheduler issues commands to the stepping motor controller (SMC) in parallel for all ten chains. The CAP now provides remote manual controls for the target alignment viewers, target positioner, auxiliary aid positioner, and target plane imager.

Figure 5-21 shows how the CAP and dispatcher interrelate with other major software programs in the system. The closed

loop alignment of components such as mirror gimbals and spatial filter pinholes is accomplished by the automatic alignment program (AAP). The dispatcher prepares the configuration of the system for each alignment task. The CAP implements manual controls and serves as the operator interface to the system. One copy of the CAP runs for each active console being used. The graphics system, which is common to the entire Nova control system, interfaces the CAP programs to the operator

Fig. 5-21. Nova alignment subsystem software.

through the console graphics screen. The alignment status program (ASP) shows the alignment system device status on several schematic displays. An additional status program for automatic alignment displays loop activity on a single integrated display. The alignment gate program allows the shot scheduler to control many of the alignment functions from a master integrated control menu.

The dispatcher reads task definition files containing SMC commands and AAP commands. It then issues the specified SMC device commands, checks final device positioning, and reports any errors. Once the SMC commands have been completed and the system is properly set up, additional commands are relayed to the AAP. Automatic alignment is then monitored until all the specified commands are completed. Any automatic alignment error is also reported to the console operator. At any time after the completion of all the specified commands, the dispatcher can verify that final device positions have not changed. It communicates to the other programs through global section memory. Any communication to the SMC is done over the Novanet local area network. Both the console operator and the shot scheduler use the dispatcher to configure the alignment system for shots and other tasks.

The CAP has been modified so that a request to configure the alignment system can now be made through the dispatcher. Prior to this modification, a reduced set of alignment tasks was hard-coded into the CAP. This capability did not include any target chamber equipment for the ten- or two-beam areas. Now, by simply editing the task definition file (maintained in the new file format), the operator can further expand the system.

Alignment tasks are selected by touching programmable function buttons displayed on the configuration menu at the console. To provide additional capacity for the two-beam area, the number of available programmable tasks has been increased to 96. The 96 buttons are arranged in 3 menu pages, with 32 buttons on each page. Another parameter file defines the button name, the name of the task file to be executed, and the area of the laser with which the button is associated. This

mechanism allows similar tasks to be grouped for easier identification by console operators.

When a new task is started, the CAP first stops any currently running task that is in the same area of the laser. It then requests the dispatcher to execute the task file commands for the currently selected laser chains. The button color shows the status of task processing. If the dispatcher detects an error at any time during task execution, the error is reported at the console. The operator can either continue execution, bypass the error, abort execution, or restart the task. If the task is continued after the cause of the error has been identified and corrected, then the command that caused the error will be retried and task processing will resume. If the operator chooses to bypass the error, then the command that caused the error will not be processed further. Once all the errors have been fixed or bypassed by the operator, the task will execute to completion. If the task is aborted, the processing is terminated and all SMC devices and automatic alignment functions for the selected area are de-allocated. At any time during the task execution, processing can be temporarily suspended, aborted, or restarted.

Figure 5-22 shows the manual control page for the target plane imager, target positioner, and target illuminator. The control procedures for these devices are complex, requiring movement of multiple devices.¹⁵ For example, moving the target positioner to the chamber center might first require retracting the target plane imager or the auxiliary aid positioner from the chamber center, opening vacuum gate valves, and moving the target plane imager near-field mirror. The ordered list of motor commands is stored in task files executed by the dispatcher.

The performance of remote manual control of target chamber equipment using the joystick became unsatisfactory because of the extra delay required to process motor commands through the auxiliary high-performance motor drives. Alignment devices in these areas also have a wide range of stepping sensitivities; some devices move very little for a number of steps, while others move a large distance for a single step. The force-operated joysticks used for the last several years,

Fig. 5-22. Alignment manual controls for the Nova ten-beam target chamber.

although well-suited to manual positioning in a tightly coupled system, proved too inflexible to solve the delay and rate problem.

A new joystick was implemented using potentiometers mounted on a two-axis gimbal. The potentiometers are interfaced through analog-to-digital converters to the existing joystick LSI-11 controller. A maximum-rate selector was included to allow operators to adjust the response of the joystick to better match the response of each alignment device. In one mode of operation, the joystick simulates the action of the SMC local control panel slew switches. This mode provides response characteristics similar to those that had become familiar to operators during the early days of system activation when local control was the only option available.

Software throughput was also increased to reduce the delay inherent in communicating the joystick commands through multiple processors to the selected stepping motors. Software in the joystick processor and VAX were optimized for the new joystick. The SMC processors interfaced to Modulynx equipment were replaced with faster LSI-11/73 units. These simple plug-in units increased the speed of

Novanet communications threefold, substantially reducing joystick delay.

The operator console software for the original Nova ten-beam system is complete. Future effort is required to extend this capability to include the additional two-beam system components. Manual controls and status displays need to be added for the two-beam chamber, although the design of the dispatcher already provides full capability for semi-automatic configuration, as described above.

Authors: R. J. Reeves, R. T. Shelton, and P. J. Van Arsdall

Praxis. We continued to extend, improve, and document the Praxis computer language,^{16,17} which is used almost exclusively on the VAX and LSI-11 computers within the Nova control system. Praxis is a modern block-structured language originally designed for Nova by Bolt, Beranek, and Newman, Cambridge, Mass. Praxis has many of the capabilities of the newer Ada language, which is now sponsored by the U.S. Department of Defense with the intent that its use in all embedded control systems will substantially improve performance

at reduced cost. In 1985, Ada became generally available on a number of different computers¹⁸; however, no validated compiler exists that will generate code for the popular PDP-11/LSI-11 series of computers. Praxis remains the best language for this application.

We believe Praxis has other features valuable for embedded control applications, features that either do not exist in Ada or are not easily found in the literature. We have obtained Ada for our VAX computers and are comparing it with Praxis to determine the suitability of each language for future control system applications.

Improvements to Praxis enabled us to design and prepare new software and modify existing software. In 1985, we improved the interface between Praxis and the VAX/VMS operating system so that all of the features of VMS are available to the Praxis programmer. One such feature is particularly valuable for tracing the operation of programs. Called "Symbolic Debugger," it allows the operator to designate as a halt point any statement in any Praxis program. When a program is executed during testing of the Nova control system, it stops at the designated halt point so that the operator can examine and modify symbol variables in the program. A few statements above and below the halt point are displayed on the

terminal screen so that the operator can more easily understand the context of the halt point statement. After the desired examination or modification, program execution is resumed.

Other improvements to Praxis include a direct interface with the VMS run-time library, which has allowed access to thousands of library routines now supported under the newest releases of the operating system; better diagnostic messages; extensive variable range checking; and internal compiler trace messages. In addition, a Praxis user outside LLNL provided us with a complete interface to the VMS file system. Each of these improvements has contributed to a better understanding of both the use of the language on our control systems and the internal operation of the compiler itself.

Praxis was released to the public this year through both the Argonne National Software Center, Argonne, Ill. and the Digital Equipment User's Society (DECUS). Through these releases, we are searching for organizations who will cooperatively share in future developments.

Author: F. W. Holloway

Prototype CCD Camera Test. During 1985 we installed and tested a second generation CCD camera, which is attached to

Fig. 5-23. Hardware configuration of the unit used to evaluate a new CCD camera. The camera is attached to an LLNL-designed streak camera in the MOR.

our standard LLNL-designed streak camera. We evaluated the CCD camera to determine if it would be a suitable replacement for an older CCD camera that can no longer be maintained because of component obsolescence. The new camera has nearly a 14-bit dynamic range and a resolution of 576×384 pixels, which increases performance over the older unit.

A single evaluation unit was built, attached to a streak camera, and installed in the Nova MOR. Figure 5-23 shows a block diagram of the installed hardware. Both a local terminal and local display were placed in the MOR so that from there operators could input commands and view captured images. All software was written in the Praxis language to allow integration into the Nova control system. We also used Novanet, our standard fiber-optic local-area digital communications network, to transfer commands and data to and from the prototype camera. Each image generates over 400 kbyte of data, and Novanet allows fast transfer to the control VAX computer at 10 Mbits/s.

Evaluating the camera in the MOR consisted of several steps. First, a short optical pulse was passed through a 0.5-ns etalon to produce a decaying pulse train with a known time interval. The result, shown in Fig. 5-24, is a three-dimensional plot of the image, the three axes representing time, distance across the streak

Fig. 5-24. Streak calibration plot using a 0.5-ns etalon.

camera slit, and intensity. This procedure served to accurately calibrate the camera in both time and amplitude. Next, the camera was set up to record the output of the Nova MOR. Figure 5-25(a) shows a plot of intensity vs time. The plot was computed for the area bounded by the lines in the color-enhanced photograph [Fig. 5-25(b)]. Note the flat baseline, which has an rms noise level of 8 counts out of a maximum of 16 363 counts.

Because these tests were successful, we have decided to build and retrofit all of our Nova streak cameras with this new CCD camera. During 1985, we have purchased the additional units along with additional disk memory to allow storage and computer analysis of the resultant images.

Author: J. R. Severyn

Major Contributors: B. A. Jones, N. J. Selchow, C. S. Petty, J. W. Taylor, and J. D. Wiedwald

Nova Facility Monitoring System. We developed the Nova facility monitoring system (FACMON) during 1985 to provide a reliable means of relaying continuous information from sensors distributed throughout Nova to the VAX-11/780 computer system and to operator consoles in the central control room. This information is in the form of temperature and humidity data from the laser bay, switchyard, and the two target bays; pressure measurements in the two target chambers; and the moisture content of the dry air used to vent the ten-beam target chamber.

FACMON consists of four Hewlett-Packard 3497A data-acquisition and control units connected to the VAX on RS-232 serial lines. A 3497A is located in each of the four main rooms in the Nova facility (laser bay, switchyard, and the two target bays), and the analog outputs of the various sensors are input to these 3497As. The sensors are sampled every 10 min, and a 7-day log of these readings is maintained in an array, which is written into a series of files once every hour. This information can then be plotted on the Ramtek displays in the control room or on a Versatec printer.

Because this system can be easily expanded, next year we plan to start monitoring the flow rate of nitrogen in the laser

Fig. 5-25. (a) Actual laser-shot intensity profile and (b) corresponding color-enhanced image with time as the horizontal axis.

cooling system, the pressures in the spatial filters, and selected target diagnostic system pressures.

Authors: C. A. Clower and M. W. Kartz

Nova Target Systems

Ten-Beam Target Chamber

Highlights of 1985. Our experimental capabilities have increased immensely in

1985. At the outset of the year, our ability to use the ten-beam Nova target chamber was very limited. For example, in late 1984, we fired all ten beams of Nova through a 4-mm pinhole at the center of the chamber and measured the energy using calorimeters. The target was aligned by a person in the target chamber and the chamber vacuum system was operated manually. No target diagnostics were available to measure target performance.

By the end of 1985, we had made substantial progress in these areas (see

subsequent articles in this section for more detail on these capabilities):

- We could view and align the laser beams and a target to an accuracy of $\pm 100 \mu\text{m}$.
- We could operate the target chamber vacuum system and eight diagnostic vacuum systems with computer control and safety interlocks.
- Fifteen target diagnostic systems provided many channels of data on complex target shots.
- Our laser diagnostics characterized the energy, pulse shape, and spatial structure of the laser beams used to irradiate a target.
- We nearly completed an optical fiducial system to provide timing of all the streak cameras relative to the laser pulse.
- We began to implement a target prealignment system to improve the efficiency and accuracy of target assembly and alignment.

Other Accomplishments.

Training and Operating Procedures.

Many of the target diagnostics and related systems in the Nova target bay are unique instruments with specific operating requirements, and only some of these operations have been automated. As a result, it is not trivial to reliably operate all the equipment in the target area for a two-shift schedule of target shots. We have expended great effort training personnel and developing procedures that will allow us to increase this capability. In addition, we have implemented hundreds of small refinements, additions, and adjustments to make the equipment more usable and reliable.

Target Positioner. We evaluated the target positioner extensively to characterize its motions and determine the accuracy with which set points could be reproduced. We had to make significant changes to the target insertion tower and the positioner's rotary alignment system to make it more usable. The vacuum controls that allow a target to be inserted through an air lock while the chamber is under vacuum have been automated for easier and safer operations.

Alignment Aid Positioner. We fabricated, evaluated, and installed the alignment aid positioner; essentially a duplicate of the target positioner, it is

located 180° away, at the bottom pole of the chamber.

Target Bay Working Platforms. Extensive modifications have been made to the working platform around the target chamber to provide more efficient traffic flow on the main level, space for additional electronic racks, and easier access to various diagnostics.

Interior Chamber Scaffolding. We installed an interior chamber scaffolding that allows easy access to all parts of the inside of the target chamber. It consists of a permanent central frame to which working platforms can be attached as needed.

Dry Air Supply. We installed a dry air supply that allows the chamber and all the vacuum lines to be filled with 15% relative humidity, filtered air when bringing the system to atmospheric pressure. The chamber is also constantly flushed with dry air while it is at atmosphere.

Vacuum System Traps. We added refrigerated traps to the main and auxiliary roughing systems to prevent oil from back streaming into the chamber. Catalytic traps are also being installed on the turbomolecular and cryogenic pumps to provide additional protection against oil contamination.

Authors: R. P. Drake and S. P. Starkey

Vacuum Systems and Controls. In 1983, the Nova ten-beam target chamber and main vacuum system were installed,¹⁹ and in 1984 the vacuum control system was activated.²⁰ This year, we began to activate target diagnostic systems whose vacuum systems are external to the main Nova system. We have developed standardized control configurations for three generic types of vacuum system diagnostics.

The simplest of the three systems, shown in Fig. 5-26, consists of a main chamber gate valve (V1); a vent valve (V2); a roughing valve (V3); a roughing bypass valve (V4) to slow down the evacuation of the diagnostic so that delicate filters and crystals are not damaged; a manifold valve (V5) to isolate the diagnostic from the common auxiliary roughing system, which is shared by all of the Nova diagnostics; and a vacuum gauge (P1) to monitor the pressure inside the diagnostic.

Fig. 5-26. Control configurations for three vacuum system diagnostics. The high-vacuum and streak camera systems build on the configuration(s) that precede them in the schematic.

This year we activated three diagnostics using this generic system: two $8 \times$ microscopes and a filter fluoroscer. This configuration will also be used on the Henway spectrometers, the streak crystal spectrograph, the radiochemistry diagnostics, and on three reentrant manipulators designed to insert and align small diagnostic packages (such as pinhole cameras) into the target chamber close to chamber center.

The second configuration (again, see Fig. 5-26) was developed for diagnostics that have more stringent vacuum requirements. This system contains all of the components of the previous system plus the hardware necessary to maintain the diagnostic pressure below 1×10^{-5} Torr. The additional components include a cryopump, a valve (V6) that connects the pump to the diagnostic, a vent valve (V7), a roughing valve (V8), a pressure gauge for the cryopump (P2), and an additional high-vacuum pressure gauge for the diagnostic (P3). We used this system on the two Dante soft x-ray spectrometers now activated on Nova, and will use it next

year on a reentrant manipulator similar to the one mentioned above.

Figure 5-26 also shows our most complex diagnostic vacuum system configuration. This system contains all of the components of the high-vacuum system just described plus the devices necessary to add a streak camera to the diagnostic. The additional devices include an ion pump, a pressure gauge (P4), a valve (V9) to isolate the streak camera from the diagnostic, a roughing valve (V10), a vent valve (V11) for the streak camera, a valve (V12) to provide additional roughing of the streak camera, and a valve (V13) to isolate the streak camera from the roughing system. We used this configuration to activate the spatially discriminating streak spectrograph (SDSS), and will use it in the future on the optical x-ray streak camera and a second SDSS.

This generic approach to the control of Nova target diagnostic vacuum systems allows us to complete most of the wiring and software modification tasks before the diagnostic is delivered to the target chamber. For this reason, we are able to

activate new diagnostics in a timely and reliable manner.

Author: C. A. Clower

Opposed-Port Diagnostics. We have activated two opposed-port diagnostics (OPD1 and OPD2) on arms 8 and 9 of the Nova ten-beam chamber. The principal function of the OPDs is to provide a magnified image of the Nova beam at four planes in the vicinity of best focus during pulsed-power shots. The OPDs have provided us with the most detailed focal-spot information acquired to date. (See the *Laser Program Annual Report 1984*²¹ for a description of the OPD design.)

In the operating configuration, the end of the OPD telescope is positioned approximately 10 cm from the chamber center. This 10-cm offset is the result of a spacer that was added during installation to prevent mechanical interference with other reentrant diagnostics. The four object planes are located at the end of the telescope. For pictures around best focus, the Nova beam is focused on one of these planes. To determine the best-focus position, we view the alignment beam on a CID camera, which is registered to one of the image planes in the multiple-image camera (MIC). The object planes are nominally separated by 500 μm . Because the image distances in the MIC are unequal and the

effective focal length of the OPD is short (~ 50 mm), the lateral magnification varies from ~ 12 to 14.5 between image planes, causing the actual separation between object planes to deviate slightly from the nominal value. To image a different set of planes in the beam, the Nova focus lens is translated along the focal axis, changing the location of best focus with respect to these planes.

In April, we took preliminary photographs of beam 9 using OPD1. The data sequence comprised 11 shots at 3ω at and near best focus. Upon installation of OPD2, we undertook a target campaign involving simultaneous shots into both diagnostics. This series consisted of 10 shots at 3ω , with energies of about 1500 J in 1100-ps square pulses. Photos were taken at best focus and 3.5, 3.0, 2.5, 2.0, 1.0 and 0.5 mm on each side of focus for both beams. In addition, photos at 7 mm before best focus were taken and full-beam streaks were acquired.

Figure 5-27 is a MIC photograph of beam 8 taken with OPD2 bracketing best focus. The horizontal rows are the four image planes, the bottom and top rows corresponding to the object planes closest to and farthest from the instrument, respectively. The second and fourth rows are staggered to the right to prevent the images from overlapping with ghost reflections caused by the beam splitters for the other two rows. The images on each row are multiple exposures of the first image on the left, each exposure corresponding to a decrease in intensity by a factor of ~ 2 . (The actual ΔE for the first two exposures differs from that of the following exposures due to the difference between the external and internal reflectivities of the MIC rattle mirror.) Unfortunately, due to placement of the MIC rattle mirror in a slowly converging section of the beam, there is a slight focus shift of 80 μm per multiple exposure. The images are effectively moving in the direction of the following row at 80 μm per exposure. Beam intensity from film density ($D - \log E$) information can still be obtained from these images since, as Fig. 5-27 shows, very little change in beam structure occurs over the 80 μm .

The MIC photographs are processed on a densitometer and VAX. Selected images

Fig. 5-27. Opposed-port diagnostic photo of a Nova beam bracketing best focus.

Fig. 5-28. Nova single-arm focal pattern at 1.5 TW, 3ω .

from each photograph are digitized, as described in the *Laser Program Annual Report 1984*.²² Typical pixel size at the film plane is 22 or 60 μm , depending on the size of the image, which translates into 2 to 5 μm at the laser image plane. Care is taken to scan an area that includes the entire image. A D -log E curve for each piece of film is generated using two digitized images from any one row in the photograph.²³ An updated version of the beam analysis code uses this curve to produce a time-averaged spatial intensity profile for each of the digitized images. Several useful plots showing intensity distributions and the beam power as a function of radius are also generated. A sample of the processed data showing the focus profile of beam 8 is presented in Fig. 5-28.

Factors that may limit instrument resolution and therefore determine how the processed data is interpreted include the following: (1) the accuracy with which

beam energy and pulse length are measured, (2) the degree of the exposure and the dynamic range of the recorded images, and (3) mounting stress distortions and spherical aberration introduced by the large telescope mirrors.

Beam energy and pulse length are used during image processing to convert the relative exposures of the beam images to absolute intensities. The relative exposure values are divided by the pulse length to generate time-average relative intensities. By requiring the total volume under the relative intensity profile of the image to equal the total power in the pulse, the absolute intensity profile for the focused spot is generated. A photodiode in the camera module, which is calibrated to the incident beam calorimeter in the diagnostic spool, measures the beam energy to within 5%. Full-beam streak capabilities determine the FWHM pulse length with roughly the same accuracy. These accuracies are consistent

with the current beam diagnostics available on Nova.

The amount of exposure at the film plane is controlled with insertable neutral density filters. Overexposure usually saturates the film, wiping out beam structure and making accurate calculation of D -log E curves impossible. Underexposure limits the dynamic range of the image and results in poorly delineated beam profiles. During processing, the physical dimensions of the beam spot are unknown, so the beam spot is assumed to occupy the entire area of the digitized image, which includes the surrounding film fog. This fog level is not subtracted from the images before processing. When processing underexposed images, a significant amount of the beam energy can be assigned to the surrounding film fog, resulting in a poor measurement of the beam's spot size. Moreover, typical focused beam intensities at the object plane of the OPD have a tremendous dynamic range, varying from 0 to 10^{17} W/cm². If the processed film has a maximum dynamic range of 3 1/2, the background fog level will correspond to an

intensity of $\sim 10^{13}$ W/cm² (see Fig. 5-28), obscuring the fine structure surrounding the focused spot. Controlled exposure and special developing will extend the dynamic range of the film to five or six orders of magnitude, allowing maximum beam information to be extracted from future processed beam images.

Optical resolution of the diagnostic is primarily determined by distortion and spherical aberration introduced by the spherical telescope mirrors. Distortion of 2 to 3 λ s (~ 1500 μ m) was measured in the mounted primary mirror before it was installed. We cancelled this aberration by carefully tilting the two lenses in the camera module. These lenses and the tertiary lens are designed and spaced to partially correct the spherical aberration as well. Measurements of both OPDs at 2ω and 3ω have shown there to be 100 μ m of residual longitudinal spherical aberration. This corresponds to 20- μ m transverse aberration at best focus for an $f/4.3$ beam, which agrees very well with our measurement of a best-focus spot diameter of 20 μ m for a 10- μ m pinhole located at the object plane

Fig. 5-28. Target plane imager.

Fig. 5-30. Target plane imager assembly.

of the instrument. At this level of resolution, the OPD has provided us with valuable data for characterizing the focused Nova beams.

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Target Plane Imager. To correctly align Nova's harmonic beams, we have designed a target plane imager (TPI) that is capable of near-field viewing and far-field viewing at three magnifications and three wavelengths. Mounted on the Nova target chamber, this 4-m-long, 1000-kg microscope relays images from chamber center to a video optics module located on the outside of the chamber (see Fig. 5-29). The entire instrument can be scanned in x , y , z to examine various planes near chamber center.

The system design was driven by the following requirements:

1. Provide remote viewing of focused laser beams at and near chamber center.
2. Provide near-field images of each full-aperture, 74-cm beam.
3. Perform this imaging at each of the three wavelengths.
4. Have sufficient sensitivity to form video images when viewing low-intensity sources.

Figure 5-30 shows the opto-mechanical system that was designed to satisfy these requirements. It consists of two parts: the

optical periscope and the viewing module. The optical periscope forms the images and relays them out of the chamber. It consists of an $f/3$ objective (eight elements), a one-to-one relay optic (four elements), and two field lenses. The viewing module makes it possible to view both the near-field and far-field images at three magnifications. The optical assembly is mounted on an x , y translation stage so that a plane 50×50 mm at chamber center can be examined. In addition, all the optics (and the video camera) can be scanned in the z direction to focus on any plane within ± 25 mm of chamber center. Figure 5-29 depicts one of the laser beams reflected from the viewing mirror into the optical periscope. This is the operating mode used when forming near-field images of the beams. By tipping the viewing mirror and rotating the periscope to the correct angle, each beam can be viewed directly. Optics in the viewing module allow examination of the beams at a plane just outside the chamber. This is where the 74-cm target focus lens and the KDP arrays are located. Therefore, beam alignment to these and other elements of the system can be determined. To directly view a far-field image, the viewing mirror is tipped out of the way and the image of the beam is focused on a frosted reference reticle. This is the most common operating mode.

First-Order Design. The internal diameter (D) of the periscope assembly and its overall length of 3000 mm constrain the first-order design parameters. Given an

object 10 mm in diameter and a front objective magnification M_1 , the maximum diameter of the first internal image is equal to the tube diameter. That is,

$$D = -M_1 \times 10 \text{ mm.} \quad (2a)$$

For a front objective with focal length f_1 , the object distance and image distances (S_1 and S'_1 , respectively) are

$$S_1 = f_1 \left(1 - \frac{1}{M_1}\right) \quad \text{and}$$

$$S'_1 = f_1 (1 - M_1). \quad (2b)$$

The conjugates of the one-to-one relay are equal to S'_1 :

$$S_2 = S'_2 = S'_1 = f_1 (1 - M_1). \quad (2c)$$

The overall length of the periscope is given by

$$\begin{aligned} 3000 \text{ mm} &= S_1 + S'_1 + S_2 + S'_2 \\ &= f_1 \frac{(M_1 - 1)}{M_1} + 3f_1 (1 - M_1) \\ &= \frac{(M_1 - 1)}{M_1} f_1 (1 - 3M_1). \end{aligned} \quad (3)$$

The effective collecting f -number of the front objective is

$$f/\text{No.} = \frac{S_1}{D} = \frac{f_1 \left(1 - \frac{1}{M_1}\right)}{D}. \quad (4)$$

We have generated three equations involving four unspecified parameters: D , M_1 , f_1 , and $f/\text{No.}$. Given a particular $f/\text{No.}$ value, we can solve for the remaining parameters.

Equations (2) and (4) combine to give an expression for f_1 :

$$f_1 = \frac{10 M_1^2 f/\text{No.}}{(M_1 - 1)}. \quad (5)$$

Substituting Eq. (4) into (3) gives

$$3M_1^2 - M_1 = \frac{300}{f/\text{No.}}, \quad \text{or} \quad (6)$$

$$M_1 = \frac{1 - \sqrt{1 + \frac{3600}{f/\text{No.}}}}{6} \approx \frac{-10}{f/\text{No.}} \quad (7)$$

and

$$D = -M_1 \times 10 \text{ mm} \approx \frac{100 \text{ mm}}{\sqrt{f/\text{No.}}}. \quad (8)$$

Therefore, an $f/3$ collecting system can easily be contained within a tube diameter of about 60 mm.

Achromatic vs Uncorrected Chromatic.

Although the periscope must operate at $\lambda = 0.351$, 0.527 , and $1.054 \mu\text{m}$, there is no requirement for simultaneous imaging at all these wavelengths. This is fortunate, since an achromatic $f/3$ front objective would be quite complex. To ensure high optical throughput at 3ω , the front objective was designed using only fused silica. The objective is nearly diffraction-limited at all wavelengths, although the focal length changes slightly. A slight change in front conjugate, obtained by moving the entire periscope assembly, compensates for the shorter focal length at $0.351 \mu\text{m}$.

The one-to-one relay, approximately 1900 mm long, is about 110 mm shorter at $0.35 \mu\text{m}$ than at $1.05 \mu\text{m}$ if it is made entirely of fused silica. Since it is mechanically inconvenient to compensate for this large change, this relay is an achromatic lens with equal length at 0.35 and $1.05 \mu\text{m}$. The small field angle ($+2^\circ$) and high f -number ($f/16$ in and out) permit the design to be implemented as back-to-back, widely air-spaced telescope doublets. The Schott flint elements absorb less than 5% of the energy at $0.35 \mu\text{m}$.

The two field lenses are also achromatic doublets so that the pupil is properly relayed through the system in the wide field-of-view mode.

Performance. To achieve high resolution and high sensitivity, the input to the TPI was made fast ($f/3$). In operation, the

limiting factors on system resolution are the graininess of the diffuse alignment reticles used and the CID resolution. Even so, we have been able to resolve $7\text{-}\mu\text{m}$ features and position focal spots to an accuracy of $\pm 15\text{ }\mu\text{m}$. Although this could be improved by using more sophisticated image analysis techniques, it has proven satisfactory for our application.

System sensitivity is determined by the signal-to-noise ratio in the CID cameras used, the solid angle of the instrument input, and the scattering characteristics of the reticles used. Since the TPI is effectively an $f/3$ system, the input solid angle is about 8.5×10^{-2} sr. Laboratory measurements have determined that the minimum detectable signal (signal-to-noise ratio of 1) for the CID camera used is 2×10^{-7} W/cm² at $1.05\text{ }\mu\text{m}$, 1×10^{-7} W/cm² at $0.53\text{ }\mu\text{m}$, and 1×10^{-5} W/cm² at $0.35\text{ }\mu\text{m}$ (on the chip). Scattering characteristics of the reticle material have been measured also. When a beam is incident on the frosted glass at an angle of 45° , 35 to 40% of the light is scattered forward into angles less than 15° . About 50% of the light is scattered forward into a lobe

pattern that can be approximated by

$$R(\theta) = \frac{7}{2\pi} \frac{P_s}{\cos^6(\theta)} \text{ W/sr}, \quad (9)$$

where P_s is about 50% of the incident power. In addition, 10 to 15% of the light is scattered into a similar pattern centered on the direction expected for specular reflection. The result is that the TPI can view and align focused beams with a power of about 10^{-8} W incident at chamber center for 1.05 and $0.53\text{ }\mu\text{m}$ and about 10^{-6} W at $0.35\text{ }\mu\text{m}$.

System Operation. Aligning the laser beams and the fusion target prior to an experiment is accomplished using several interdependent subsystems at the target chamber. (See Fig. 5-31.) Three target alignment viewers determine an orthogonal reference frame within which target positions are measured. A typical alignment procedure involves positioning a reference reticle near chamber center using the alignment aid positioner. The precise location of the reticle depends on the geometry of the target. Next, the laser beams of the correct wavelength are

Fig. 5-31. Manual target alignment controls.

focused onto the reference reticle as determined by the TPI. Finally, the reference reticle is removed and a fusion target is aligned precisely to chamber center using the target positioner. At this point, all the viewing instruments are commanded to safe configuration before the target is irradiated.

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Target Alignment Viewers and Illuminators. Three of the target alignment viewers (TAVs) used on Novette²⁴ have been reconfigured for use on the Nova ten-beam chamber. The entire objective and reticle module (Fig. 5-32) is housed in a vacuum jacket extending 4.5 ft into the chamber; the first objective lens is held 920 mm from the target chamber center, as it was on Novette. To prevent possible interference with adjacent diagnostics mounted on the outer surface of the target chamber, we eliminated the dogleg in the Novette TAV optical path. (The Novette hardware had been designed to allow mounting in this straight-line configuration.) Thus, the whole optical system for each TAV is mounted on a single 26-in.-diam flange that supports the 4.5-ft reentrant vacuum jacket and an external 6-ft-long support plate for the relay and magnification modules.

Fig. 5-32. Optical layout of ten-beam chamber target alignment viewer.

The equatorial plane of the chamber, containing many of the critical target diagnostics, is oriented vertically within the target bay. A spherical coordinate system has been defined for the chamber such that the ports on this "equator" have a latitude θ of 90° as measured from the east pole ($\theta = 0^\circ$). The laser beam ports are positioned at $\theta = 50^\circ$ and $\theta = 130^\circ$. Longitude ϕ is measured from the target positioner port $\phi = 0^\circ$ at the top of the chamber, down the south side and up the north.

The three TAVs are mounted to give orthogonal views of target chamber center. Two are located at ($\theta = 90^\circ$, $\phi = 45^\circ$) and (90° , 135°). The third and primary viewer, located at the west pole (180° , 0°), is centered within the port circle for the odd-numbered laser beams. Targets are always mounted so that their laser-illuminated surface is visible in the polar TAV, but this surface may not be visible to the equatorial TAVs.

The targets to be viewed by these three TAVs are backlit by three target alignment illuminator (TAI) assemblies, each containing a helium-neon (He-Ne) laser (see Fig. 5-33). In addition, each equatorial TAI provides some front-lighting because its location at either (90° , 225°) or (90° , 315°), directly across the chamber from its TAV, results in forward scattering of light from some portion of most targets into the orthogonal equatorial TAV line of sight.

The optical path for the polar TAV is complicated by the fact that the opposite

Fig. 5-33. Optical layout of ten-beam chamber target alignment illuminator.

pole is occupied by the TPI. The "polar" TAI is actually mounted in port (68°, 0°). Its laser beam is directed to a 13-mm mirror mounted on the TPI and held, for target alignment purposes, about 6 in. from the target. The TPI controls allow this mirror to be adjusted until the illumination strikes the target and centers on the TAV optics.

The TAI and TAV assemblies are operated through the Nova control system. The lasers in the TAs may be turned on and off, the TAs' output intensities may be adjusted by turning a wave plate, and the shutters may be opened and closed from the touch-sensitive panels in the Nova control room. These panels also control the TAV focus, magnification, and field scanning and operate the shutter.

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Target Diagnostics Controls. During the past year, our major effort on the target diagnostics control system has centered on supporting the two-beam chamber. We partitioned off an area near this chamber at the west end of the Nova facility where we installed equipment for diagnostic setup and data acquisition. Among the equipment installed were four data-acquisition front end processors (FEPs) and four vacuum system FEPs.

We installed the data acquisition FEPs, trigger system, and a Ramtek color graphics operator control panel in this newly created area (the vacuum system FEPs were installed in the target chamber room and fan loft), electrically isolating all equipment from the building and then connecting one closely monitored ground cable. All signals into and out of this area are either via fiber optic cable or highly

isolated transformers. This procedure, consistent with that used for past diagnostic systems, helps ensure analog diagnostic signals free from distortion induced by multiple ground connections.

A small video switching system now allows the operator in the main Nova control room to select graphic control menus for either the previously installed Nova ten-beam loft diagnostics or the new two-beam diagnostic equipment. Diagnostic setups can be conducted in both the two- and ten-beam chamber areas, while the actual target shot sequence is performed from the main Nova control room.

During the past year, major software projects that supported data acquisition on the ten-beam system were used to support the two-beam effort. Internal software data tables for the ten-beam system were expanded and data flags added to mark the chamber area location of all diagnostic equipment. The software that controls both diagnostic areas uses the same data tables, which minimizes computer loading; however, to allow operators to set up diagnostic instruments in both the two- and ten-beam chambers at the same time, two sets are needed of that portion of the control software that supports the operator interface.

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Optical Fiducial System. This year we added a small pulsed optical fiducial laser system to the Nova facility. It provides a frequency-doubled or -tripled optical pulse, synchronized within 50 ps of the Nova

master oscillator, as a timing fiducial for all target diagnostic optical and x-ray streak cameras on either target chamber. The second harmonic is being used. At either frequency, the optical pulse width is less than 100 ps, and the streak record is even shorter. This laser also provides a 10-mJ infrared trigger pulse for primary diagnostic instruments used in x-ray laser target experiments on the two-beam chamber.

We designed the optical fiducial laser²⁵ for a 5-mJ output at 3ω . Frequency conversion was accomplished using the same KDP crystal configuration that is used to convert the ten Nova beams at the target chamber.²⁶ In this configuration, either 2ω or 3ω can be generated efficiently with a minor change in the tuning angles of the crystal assembly. Use of KDP as the conversion material was advantageous because we were able to use the sol-gel anti-reflective coatings developed for the large crystal arrays.²⁷

The optical layout of the ten-beam chamber fiducial²⁵ has undergone only minor changes since the system's activation. Because passive transmission (50%) reduced the unamplified MOR pulse to 50 μ J at the input to the crystal assembly and because we had designed the system to drive the crystals with 25 mJ of 1ω energy, we required an overall system gain of 500. This gain has proven attainable on a routine basis. We changed the four-lens relay to a more flexible five-lens combination that places the geometric image of the oscillator aperture to within the Rayleigh range of the crystal, while at the same time projecting the Gaussian beam waist to the crystal. The final pair of lenses along the beamline is a 70% beam expander, which produces a 5-mm ($1/e$) beam diameter at the crystal.

The phase shift resulting from an intensity-dependent change in index of refraction for an amplifier may be written in terms of the output beam intensity:

$$\Delta B = 0.0172 \frac{I(G - 1)L}{G \ln G}, \quad (10)$$

where ΔB = fractional wave-phase shift, I = output intensity in GW/cm², G = gross gain of the amplifier, and L = length of the amplifier. Approximating the output

phase front by a Gaussian of the same width as the intensity profile, the deviation θ of the wave-front normal from the nominal propagation direction is given by

$$\tan \theta = 0.0172 \frac{I_0(G - 1)L}{G \ln G} \exp\left(-\frac{r^2}{a^2}\right), \quad (11)$$

where $a = L/e$ radius of the beam, and I_0 = peak intensity of the Gaussian profile. At the second amplifier, the beam diameter (in the absence of the beam expander) is 3.2 mm.

The maximum deviation of the wave-front normal for this diameter at the required output intensity (3.1 GW/cm²) is 145 μ rad. Expansion of the beam to 5 mm before it reaches the amplifier reduces the maximum beam deviation to 38 μ rad. Because deviations of 100 μ rad reduce conversion efficiency significantly, we placed the beam expander before the second rod amplifier. On the other hand, the edges of this larger beam are clipped by the rod aperture. To prevent the propagation of an interference ring pattern, we have placed a ground-glass apodizer²⁸ between the expander and the amplifier.

The major change in the system design this year involved extending it to the two-beam target chamber. The first use of the fiducial laser was to trigger the x-ray laser primary diagnostics on this auxiliary chamber. To change from one target chamber to the other, we insert a mirror after the diagnostic beam splitter (Fig. 5-34) to divert the 1.05- μ m beam and the frequency-converted beam into an optical beamline that relays and directs them an additional 61 m to the two-beam target chamber bay. Most of the mirrors are dichroic, reflecting both beams. We adjust the beams to be parallel, about 1 cm apart, with a pair of small mirrors located just after the insertable mirror. We adjust pointing and centering of the 1.05- μ m beam in the two-beam target bay by viewing the beam in another alignment sensor. Because of the common optics, the converted beam is then also in alignment. A pair of long focal-length lenses (16 m) relays the image of the crystal assembly to the alignment sensor. The converted beam (now 2ω) is again directed onto the face of a fiber-optic bundle and distributed to the

Fig. 5-34. Two-beam chamber Auston switch trigger and optical fiducial layout.

target-diagnostic streak cameras through individual fused-silica fibers. The $1.05\text{-}\mu\text{m}$ light is split into three beams of $\sim 3\text{ mJ}$, each delayed for correct timing, then directed to three photo-sensitive silicon switch arrays,²⁹ which operate the x-ray laser diagnostics. (See "Auston Switch System," later in this section.)

We changed to the second harmonic optical output because we found that the 3ω fluence (0.63 J/cm^2) on the input end of the fiber bundle was causing sputtering of the fiber support material, distributing it across the optical surfaces. The result was an increase of the optical fiber insertion loss from the design value of 23 dB to an unacceptable 41 dB after just 15 or 20 shots of the fiducial laser. Since the fiducial laser is fired for alignment and timing purposes much more often than the Nova laser, this posed a serious maintenance problem. We investigated 2ω as a possible alternative operating frequency, expecting a reduction in the required input intensity of up to a factor of 10 because of the low-loss transmission of the optical fibers at that frequency. We were uncertain that green light would be as effective as ultraviolet in ejecting electrons from the aluminized output surface of the fiber. However, x-ray streak-camera signals were successfully generated using green light (see "Optical Fiducial for X-Ray and Optical Streak Cameras," Sec. 4) and the conversion

crystal assembly was retuned for maximum green light production.

During activation of the fiducial laser, we observed unexpected shifts of $\sim 1\text{ ns}$ in the relative timing ($\sim 1\text{ ns}$) of the two oscillators. Upon investigation, we found that, although the Q-switch is intended to be a nonresonant structure, it does have a slight resonance close enough to the mode-locker frequency to affect the timing of the oscillator. In the past, this went unnoticed because the oscillator itself had been the source of all critically timed signals. With two such oscillators, it quickly becomes obvious. As long as the arc-lamp and Q-switch controls were left unchanged, the two oscillators were synchronized within the required 50-ps accuracy. However, the master oscillator is set each day for optimum operation by adjusting the arc-lamp current and the Q-switch drive. But when the Q-switch drive is changed, its interaction with the mode-locker also changes: the net effect is a timing shift of the mode-locked pulse train.

We installed a streak camera in the MOR to measure the timing difference before each shot. The switched-out pulses from the two oscillators are actually $\sim 1\text{ }\mu\text{s}$ apart in time. Optical fibers that carry the output pulses from the master oscillator and the fiducial oscillator to the camera have been cut to length so that pulses sent through them arrive at the streak camera

at the same time. The relative timing of the pulses is determined with a resolution of better than 50 ps (Fig. 5-35). If necessary, the length of the master oscillator mode-locker drive cable is adjusted in steps of 150 ps. For fine adjustment, a phase shifter in the rf drive is then used to set the timing to the limit of resolution. Any change in timing detected on the streak record at shot time may be used to interpret the data at the target diagnostic streak cameras. For high-energy shots, we observe timing shifts resulting from the modification of the pulse shape caused by disk-amplifier gain saturation. We compare the MOR streak record with similar records from the laser-diagnostic streak cameras located in the ten Nova output sensors. The laser diagnostic streak cameras also include a fiducial obtained from the first fiducial laser-alignment sensor by transmitting a portion of the infrared fiducial beam through optical fibers to the output streak cameras.

We compared the operational characteristics of the 1-cm rod amplifiers with those of the MOR amplifiers and found them to be very similar. Tests also

revealed a significant amount of thermal birefringence in both types of amplifiers, resulting from the operation of the amplifiers at one pulse every 5 s (Ref. 30). We found that a rod returned to its original condition in 50 s. Data obtained by the tests on Nova arm 10 concerning the effect of birefringence on conversion efficiency (see "Beam Depolarization Effects on Nova Frequency Conversion," in Sec. 6) show that a significant loss of conversion efficiency is to be expected if birefringence causes even 10% of the 1 ω light driving the crystal to be incorrectly polarized. As a result, an input polarizer has been added to the frequency conversion assembly.

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Two-Beam Target Chamber

Introduction. To increase our capacity for target physics experiments, we added an auxiliary two-beam target chamber to Nova (Fig. 5-36). The 6-ft-diam chamber, formerly used on Shiva and then Novette, was reinstalled in the Shiva target facility

Fig. 5-35. MOR streaks of the master and fiducial oscillators.

Fig. 5-36. The two-beam target chamber, 6 ft in diameter, is shown in place in the Shiva laser bay. Sensors, focusing optics, and the frequency conversion crystal arrays are located inside the large cylindrical spools.

according to the budget and schedule described in the 1984 Laser Annual.³¹ While the ten-beam chamber was being configured for more complex experiments, we were able to use the auxiliary chamber for initial x-ray laser and weapons physics experiments. Through the Nova control system, we can redirect 6 and 7 Nova beamlines into the chamber to perform target experiments that are more suited to this smaller, operationally simpler system.

On July 31, approximately 2.8 kJ of 2ω light in a 400-ps pulse were delivered to the chamber center in a line focus. The target system was highly productive soon after activation, with over 100 experiments performed during the last six months of 1985.

Construction and Activation. All construction and activation milestones established at the outset of the project were met (see Fig. 5-37). The criteria for adding the two-beam chamber to Nova, oriented towards x-ray laser experiments requiring temporally short pulses with a line focus spatial energy distribution, included the following:

- Integrating the auxiliary chamber with the main Nova system.
- Delivering 4 to 5 TW of 2ω light on target at 400-ps pulse lengths.
- Achieving a variable-length line focus on target.

Through March we concentrated on making procurements, installing the support frame, and preparing the target area. Preparation included painting, pre-cleaning, and starting to activate the clean-air recirculation system. We also completed installation of support towers for the beam-diverting mirror assemblies in Nova's optical switchyard area. In April, we began installing major mechanical components, beginning with the stainless steel target chamber. We also began full operation of the clean-air filtration system to prepare for the installation of optical components in May.

Early in 1985, experiments with the ten-beam Nova chamber indicated that the propagation of high-intensity, $1.05\text{-}\mu\text{m}$ light through nitrogen in long air paths could lead to beam degradation. In response to this concern, we modified the design of the two-beam chamber addition in two ways. First, we modified the design of the

- Installing the chamber by September 1, 1985.

Start

Milestone shot completion (4 kJ, 2 ω , 400 ps)

beam tubes that enclose the path from the switchyard to the first turning mirror. These tubes can now be evacuated or filled with a gas such as argon, which is less prone to inducing beam degradation than nitrogen. Second, we changed the design of the full-aperture mirror support for beam 7 in the two-beam chamber bay so it can be electromechanically moved out of the beam path. We constructed a diagnostic station behind the normal operating position of this mirror to help characterize the beam quality as the beam is propagated at various intensities through vacuum or gaseous media.

June and July were the peak months of activity. We activated the two-beam chamber for long path experiments in June, marking the first use of light from the Nova beamlines at this station. We also completed installation of all supports and large aperture mirrors and began installation of the target chamber diagnostic instruments. During July, we installed the remaining optics, including the KDP frequency-conversion arrays, fused-silica final focusing lenses, and cylinder lens doublets. In addition, we installed the incident-beam diagnostics, alignment aids, video switching system, target data acquisition system, main vacuum system, and the extensions to the Nova control and data acquisition system. By the end of July, all major components were in place.

During the July 31 shot, beams 6 and 7 delivered 1500 and 1300 J, respectively, of 2 ω light in a 400-ps, temporally Gaussian pulse with a line focus length of 1 cm. X-ray images were captured by two pin-hole cameras mounted on the target chamber. Data were also recorded by two x-ray spectrometers. The target surface was burned, as expected, demonstrating good beam alignment.

In this initial experiment, we acquired incident laser diagnostic data related to beam energy and temporal history. We also demonstrated that the use of incoherent illumination for target viewing during alignment was a significant improvement over previous coherent-illumination strategies.

During August, we concentrated on activating the two-beam chamber for target experiments. We also activated major x-ray laser diagnostics and installed the important Austin switch trigger system for these diagnostics. We completed analysis of the focus spot characteristics, system performance, and diagnostic instruments. We followed this with a series of successful x-ray laser experiments, which began in September and continued through the year.

In the fall of 1985, we began to extend the capabilities of the two-beam chamber. First, the chamber was set up to support, and subsequently perform, a series of laser experiments. Second, we obtained two additional counter-rotating cylinder lens

pairs that would enable us to double the line focus length on targets. We will use this capability in 1986 to support experiments on longer length x-ray laser targets.

Organization. During construction of the two-beam chamber, we were subjected to criteria and constraints beyond those strictly related to performance: The auxiliary two-beam chamber was to be fully integrated into the overall Nova system to minimize the amount of additional manpower required for effective utilization, the level of safety in construction and operation was to be kept high, and the impact on concurrent efforts to activate Nova's ten-beam chamber was to be minimal so that the target experiments could proceed as planned. We responded to these criteria by organizing the two-beam chamber effort in a fashion similar to that of the main Nova project. Performance, budget, schedule, and safety responsibilities were defined using the work breakdown structure shown in Fig. 5-38.

During the design, procurement, and installation phases, we maintained communication with the subsystem teams through design reviews and weekly meetings. As activation commenced, we held more frequent and less formal meetings to coordinate daily installation and activation tasks. Two informal cost and schedule reviews

assisted in keeping us on schedule and keeping costs under budget.

Author: G. J. Suski

Line Focus Capabilities. In the fall of 1985, we began our first major target activity on the two-beam chamber: driving x-ray laser targets. These targets require high-aspect-ratio line foci with lengths varying from 1 to 5 cm and heights from 100 to 300 μm . Also important is that the line be rotated to match precisely the nominally horizontal axis of the target. As with any target shot, the interpretation of an x-ray laser shot depends on accurate measurement of the incident pulse energy.

The line foci are obtained through the use of a matched pair of individually rotatable negative cylinder lenses located ahead of the main focus lens on each beam. The lens bezels are rotated by stepper motors whose controls are cross-coupled so that the two lenses of a pair can be easily rotated equal distances in opposite directions to vary the line focus length or equal distances in the same direction to change the orientation of the line.

If the axes of the two cylinder lenses in a pair are crossed at $\pm 45^\circ$ to horizontal, a round focus occurs beyond the focal plane

Fig. 5-38. Nova two-beam chamber work breakdown structure used to define performance, budget, schedule, and safety responsibilities for the project.

of the main lens. As the $\pm$ angle is reduced toward zero, the net cylinder power increases, causing line foci of increasing length to appear before and after the initial round focus point. The line focus ahead of the round focus point is vertical; the one behind is horizontal. Of course, the distinction between vertical and horizontal is artificial. By simply rotating the two cylinder lenses simultaneously, either line can be brought to the desired orientation. Figure 5-39 summarizes the $\pm$ rotation angle required and the offset from round focus that results as a function of line focus length. Cylinder lenses with two different curvatures have been used on the system, one curvature for lines up to 3 cm long and the other for lines up to 6 cm.

Energy measurements for line focus target shots are obtained by the diagnostic spool, which collects the ghost reflection off the main focus lens.³² With the cylinder lenses installed ahead of the focus lens,

Fig. 5-39. Cylinder lens rotations and focus lens offsets required to produce line foci up to 60 mm long.

Fig. 5-40. Measured conversion efficiency compared to calculated conversion efficiency on (a) arm 6 and (b) arm 7 of the Nova two-beam chamber.

the beam from the ghost reflection becomes elliptical instead of round. Careful checking of alignment in the diagnostic spool is important to ensure that all the light is collected. For long line focus lengths, the spacing between the diagnostic spool and the lens positioner spool must be increased or the ghost reflection will overfill the calorimeter along the major axis of the ellipse. With these adjustments, the calorimetry for line focus shots has proved reliable.

Authors: E. S. Bliss and D. W. Myers

Performance of 2ω KDP Arrays. The KDP arrays for the two-beam chamber employ 1.0-cm input crystals and 1.8-cm output crystals mounted in quadrature,³³ in contrast to the 1.0-cm/1.0-cm scheme used on Nova. This configuration was designed to increase 2ω conversion efficiency over that of the ten-beam Nova arrays, while suffering some loss in performance at 3ω .

Actual 2ω performance of the two-beam chamber arrays has been measured for 1ω drive intensities between 0.5 and 2.5 GW/cm² in 550-ps Gaussian pulses. This data is plotted in Fig. 5-40 along with the calculated conversion efficiency curve for 100- μ rad crystal detuning. For drive intensities up to 1 GW/cm², the data falls just below the calculated curve, but it flattens out above 1 GW/cm² as on Novette and Nova. Conversion efficiency remains relatively flat over the range 1 to 2.5 GW/cm²: 62% for arm 6 and 65% for arm 7.

Fig. 5-41. Schematic of the two-beam target chamber placement within the Nova system.

Experiments to investigate the causes of this flattening are described in Sec. 6, in the subsection entitled "Nonlinear Propagation and Frequency Conversion." These experiments may help improve the maximum conversion efficiency to the 70 to 80% theoretical limit.

Losses at 2ω caused by optics absorption (4%), beam obscurations (arm 6: 4%, arm 7: 9%), and surface reflections (15%) reduce the transmission between the KDP arrays and target chamber center to 77% for arm 6 and 72% for arm 7. This transmission and conversion efficiency plus current limitations on maximum 1ω drive intensity led to a maximum 2ω energy of 2500 J per beam on target in a 500-ps Gaussian pulse.

Authors: D. W. Myers and P. J. Wegner

Major Contributor: M. A. Summers

Mechanical Systems. This was a year of construction and activation on the Nova two-beam chamber. A major portion of the chamber and support equipment was assembled and installed in the first six months of 1985. By August, the two-beam chamber was fully operational (Fig. 5-41). In this article, we describe the construction and activation of its mechanical systems.

To save the considerable time and expense involved with designing and

constructing a new spaceframe, we modified the Novette frame. Judicious welding practice provided a distortion-free frame on which we mounted mirrors and the target system hardware.

Mirrors and Mirror Towers. We installed new mirrors on towers in the switchyard and target room. The switchyard towers were fabricated in weldment segments (Fig. 5-42) so they could be moved through Nova's narrow aisles without disturbing laser operations. Once in the switchyard, the segments were bolted together and mounted to the floor. This method of construction and installation enabled us to continue operating the laser without compromising the standards of the class-10 000 clean room. We were able to construct the target-room tower in a single weldment because of sufficient doorway clearances and crane capacity. Made of hollow structural steel, the mirror towers are designed for vibrational and seismic stability according to the Nova spaceframe stability design philosophy.³⁴

We installed and have successfully operated seven new mirrors: four in the switchyard and three in the new target room. Four of these mirrors are on roller bearing rails.³⁵ Three can be moved out of the beamline to let the beam pass on to the ten-beam chamber and to other diagnostic areas. The largest of the mirrors, 106 cm in diameter, is used to adjust path length

Fig. 5-42. Insertable mirror housing in the Nova switchyard. This was made in weldment segments and then moved into the target bay where it was bolted together.

Fig. 5-43. This 106-cm-diam turning mirror on beamline 6 sits on a carriage to adjust path length.

(Fig. 5-43); it has an adjustment range of approximately 12 in. This mirror is moved infrequently and is hand driven; the other mirrors travel 60 in. to clear the beam path and therefore are motor driven. Of particular importance is our new ability to move these large mirrors and at the same time maintain their optical integrity and achieve alignment routinely.

Beam Pipes. The beam propagation path from the output spatial filter to the target chamber is 298.6 ft on beamline 6 and 346 ft on beamline 7. Each of these paths is longer than any one path on the ten-beam chamber. To minimize the risk of Raman scattering, which can result from such long paths, we constructed 400 ft of 36-in.-diam beam pipe that can be evacuated to less than 1 Torr or filled to atmospheric pressure with helium, xenon, or argon. Figure 5-44, an overall view of the

system, shows the two beamlines. They transport the beam from the switchout mirror in the switchyard to the target chamber. The target chamber and mirror housings can also be seen.

Before conducting laser experiments with the beam tubes filled with gas, full-aperture, thin-quartz windows (coated with an anti-reflective) are attached to each end of the tubes. To fill the beam tubes with gas, it is necessary to evacuate the air with vacuum pumps. Before the air on either side of the window is evacuated, blank-off covers are placed over the thin windows to protect them from failure due to atmospheric pressure. In this way, equal pressure is maintained on both sides of the window during pumpdown and during backfilling with gas. The blank-off covers are then removed and the system is ready for operation.

Installation. The design criteria for the two-beam system required that the target chamber be placed in the target bay, to within 1 mm of the the Nova beams. Since the target chamber turn mirrors and beam

tube mounts can be adjusted easily, it was necessary to position the frames to within only 1 cm of their design points.

To achieve alignment, we established a survey monument on the Nova two-beam chamber floor, directly under what was to become the target chamber vertical centerline. This point was referenced to a master benchmark in the Nova switchyard. The target chamber spaceframe, beam tube supports, and turn mirror boxes were set using this benchmark as a reference. After the spaceframe was set, the target chamber was positioned and a centerline was established between the final turn mirrors through the vertical center of the chamber. After the chamber elevation was set, a helium-neon laser and mirrors, mounted on the chamber beam ports, were used to determine the exact rotation of the chamber within 60 mrad. An optical transit was mounted on the frame at beam centerline to position each spool on the target chamber/laser beam centerline within 1 mm. Several secondary benchmarks were placed strategically in the target

Fig. 5-44. An overhead view of the Nova two-beam target chamber system. Beams 6 and 7 enter the room at the bottom and are directed toward the target chamber by three mirrors. Lenses, frequency conversion crystals, and diagnostic spools are mounted on rails to facilitate maintenance.

Nova Target Systems

Table 5-5. Optical components required for the Nova two-beam target chamber.

Component	Diameter (cm)	Aperture (cm)	Qty
Turning mirror	94	90	6
Turning mirror	109	98	1
Objective lens	80	77	3
Cylinder lens	80	77	12
Beam-tube window	80	77	6
Debris shield	80	77	6
18-mm KDP crystals	27 × 27	—	17
10-mm KDP crystals	27 × 27	—	18

area so that reference surveys can be made in the future.

Authors: C. A. Hurley and H. G. Patton

Optical Components. In 1985 our major successes on the two-beam chamber included fabrication, delivery, and installation of the glass and fused-silica optics and KDP crystals. Table 5-5 lists the optical components required for the two-beam chamber. The target date for delivery of these components was May 31, 1985; however, when components could not be delivered on schedule, substitutions were made from the Nova spares inventory until the new components arrived. Each of these optics was made of fused silica, except for the BK-7 turning mirrors, which were AR-coated using the sol-gel process developed for Nova,³⁶ and the KDP crystals.

Cylinder lenses 80 cm in diameter were manufactured in two focal lengths. These lenses are used as a counter-rotating pair to produce a variable-length line focus. An initial procurement was made for long-focal-length cylinders (72 m) that provide a 0- to 3-cm line focus; these were installed in the laser in July 1985. A second set of 80-cm-diam lenses of shorter focal length (36 m) was purchased and will be available for installation in early 1986. This set will produce a line focus up to 5 cm long. (For a discussion of the cylinder lens design, see "Cylinder Lens Design for the Ten-Beam Chamber," earlier in this section.)

Beam tube windows were manufactured to provide a seal in the beam tubes that run from the switchyard to the first turning mirrors. These windows allow the air to be purged from the long optical path and replaced with an inert gas. The beam tubes windows were debris shield blanks manufactured to a tighter wave-front tolerance.

The two-beam KDP crystal arrays are a modified Nova design with 18-mm-thick input crystals and 10-mm-thick output crystals. Six spare 18-mm crystals were available from Novette, but we had to order and finish an additional 11. To complete the crystal requirements, eighteen 10-mm thick crystals were processed for the output side of the array. These 29 new crystals, and the 6 available from Novette, were diamond-machined in less than eight months and installed in May 1985. Corner crystals were cut from finished squares, which reduced machining costs by \$40 thousand and improved crystal delivery by six weeks. That the optics industry was able to respond to our requirements for state-of-the-art optics within such a demanding schedule was a significant achievement.

Author: F. T. Marchi

Target Alignment Viewers and Illuminators. When we decided to add the two-beam target chamber to Nova, three new TAVs were needed to replace those modified for the ten-beam chamber. The fourth Novette TAV could be used without modification, since the two-beam target chamber is the slightly modified Novette chamber. However, we took this opportunity to replace the He-Ne lasers with incandescent lights.

In the ten-beam chamber, as in the Novette target chamber, fringe patterns are generated by diffraction at the edges of the target and at other apertures in the optics and by interference from internal reflections in the optical trains. These fringe patterns are reduced through the use of a vibrating reflective membrane as the last turning mirror, which directs the laser beam onto the target, effectively reducing the coherence of the light that reaches the target. The vibration is provided by an audio-frequency white-noise generator and a small loudspeaker mounted directly behind the membrane. This technique tends to cause nonuniform illumination because the membranes have resonances within the audio range that set up standing wave patterns in their surfaces. Moreover, the light remains partially coherent and results in a visually noisy presentation on the TV monitor. This visual noise can seriously

Fig. 5-45. Noncoherent illuminator for the Nova two-beam chamber.

reduce a target aligner's ability to detect a feature on the surface of the target that must be placed at a specific point within the field of view.

We decided against using incoherent light in earlier target chambers because of concerns that such a source is several orders of magnitude less bright (depending on the TAV bandwidth) than small He-Ne lasers. Bench tests of a 100-W tungsten-halogen lamp convinced us that by only slightly modifying the Novette-style viewer we could achieve adequate illumination, not only for direct backlighting but for viewing our ground-glass beam-alignment aid³⁷ and front-lit targets at near-specular-reflection angles.

We designed the passband of the new viewer to include the second harmonic of the laser. This allows us to view directly the 2ω Nova alignment laser spot on the alignment aid without having to refocus the TAV. All of the optics in the Novette TAVs were achromatic except for the two 10-cm-diam singlet objective lenses, which focus the image of the target onto the reference reticle. For the Nova two-beam chamber we replaced these lenses with a pair of air-spaced achromatic doublets (Schott BK-7/SF-2) of the same focal length. These give near-diffraction-limited performance at 527 ± 15 nm. The diffraction limit for an $f/10$ optical system at 527 nm is $12.5 \mu\text{m}$. Even when these lenses are used with filters that transmit light between 500 and 580 nm, chromatic aberration results in a negligible blur circle. By comparison, the target plane resolution of the TV camera used to view the image is about $5 \mu\text{m}$ at the highest magnification.

The TAI is new hardware. To speed the design and purchase of the equipment, we used a commercial lamp housing and

condenser lens. The assembly (Fig. 5-45) incorporates three translation stages for manual adjustment of the lamp position. An $f/1$ condenser lens re-images the lamp filament near the chamber center with a cone of light, which fills the $f/10$ aperture of the TAV objective lenses. The image filament is set just far enough from the target plane that the individual wires of the filament coil are blurred to overlapping at the target plane. The 2.3×4.2 -mm lamp filament image is magnified by a factor of 7.1, filling a 16-mm aperture in the target plane that corresponds to the clear aperture of the TAV reticle.

An 80-nm-bandwidth color filter follows the condenser lens, narrowing the spectrum and reducing the effect of residual chromatic aberration on the TV image. This filter, a combination interference and absorption filter, transmits in the infrared portion of the spectrum. A 1-mm-thick BG-18 filter between the lamp and condenser lens eliminates the spectrum for wavelengths longer than 650 nm, reducing the heat load on the lens and the target.

The vacuum window in the illuminator housing is protected from target debris by a control-system-actuated shutter that is pneumatically operated. Lamp intensity is adjusted by a stepper-motor-driven autotransformer, which is also part of the main control system.

The TAVs are mounted in the same ports used on Novette³⁸ and, since the spaceframe that supports the two-beam chamber is the one that was used on Novette, the configuration is the same. One TAI, used as a backlighter for the TAV under the west laser beam, has been displaced by the TPI. It is now mounted on a port ($\theta = 120^\circ$, $\phi = 54^\circ$) that is not directly in line with the TAV. Light from this

TAI is reflected from a mirror at the end of the TPI onto the target and into the TAV aperture, as is the case in the ten-beam chamber. Because of the additional 12-in. distance from condenser lens to target along this optical path, the magnification is greater and the target plane illumination is less intense by a factor of 2.

Author: J. S. Hildum

Auston Switch System. The Auston switch system provides fast (200-ps) gating signals to x-ray spectrometers that use microchannel plates in their detector systems. This switch system was used on Novette, where it played a key role in the success of x-ray laser experiments.³⁹ When the two-beam target chamber was configured to do x-ray laser experiments, we added an Auston switch system to the facility.

One can conceptualize the switch system as containing three parts: the Auston switch assembly, the laser-trigger beam-delivery system, and the bias-pulse generator. We used the the Auston switch assemblies from Novette without modification. The trigger beam is the residual 1ω light from the frequency-converted optical fiducial system, which contains enough energy to deliver about 2 mJ to each Auston switch through a system of mirrors and beam splitters. Each trigger beam has a timing adjustment with a range of ± 4 ns; the adjustment is made by an optical delay line. The trigger-beam timing is critical to the successful operation of the system and must remain stable, to within 100 ps. The Novette bias-pulse generator was redesigned for use on Nova, resulting in high performance and reliability in its first four months of routine operation in the system.

The Auston switch system operates routinely and successfully, gating x-ray spectrometers on the two-beam target chamber. This system should allow us to pioneer the use of the Auston switch technology in new diagnostic systems.

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Target and Alignment Aid Positioners.

With the addition of an auxiliary two-beam target chamber to the Nova system came the need for a target positioner and an alignment aid positioner. These positioners had to be designed, fabricated, tested, and operationally installed over a three-month period. To meet this schedule, the Novette cryogenic target positioner and accompanying cryogenic shroud retractor were retrofitted to support ambient temperature targets and the alignment aid, respectively.

Each of the two Novette cryogenic positioners had four degrees of freedom, one rotational and three linear. The shroud positioner, on the top of the chamber, moved in synchrony with the target positioner, on the bottom of the chamber. To meet the new demands of target alignment, we replaced the cryogenic aspects of both positioners with equipment that provided an additional axis of rotation. The positioners now allow us to move the target or alignment aid on two orthogonal rotational axes and three mutually orthogonal linear axes.

The cryogenic target positioner was disassembled, leaving only the linear and rotary stages, drive mechanisms, and vacuum barrier bellows. By replacing the vertical drive mechanism on the cryogenic positioner with a precision drive screw from a Nova lens positioner we were able to greatly improve the drive characteristics of the cryogenic positioner (see Fig. 5-46). As the screw moves up and down, it is restrained from rotary motion by a spring-loaded, three-wheel trolley, which runs in a track. The drive nut is turned by a small rotary stage. The stage is stepper-motor driven, with rate controlled from 1 to 3000 steps per second. This wide range of speed control allows targets to be moved in and out of position quickly but with submicron resolution.

We added a small goniometer to the end of the target positioner (see insert, Fig. 5-46) to provide rotation about an axis in the horizontal plane, specifically the beam axis. The goniometer is offset from the centerline of the positioner to prevent obscuration of the target alignment viewers. A positioner mount (Fig. 5-46), the complementary part of the target mount, is attached to the goniometer with an equal

Fig. 5-46. The target positioner used in the Nova two-beam target chamber. The insert shows an x-ray laser target mounted on a goniometer, which rotates the target about the beam axis.

but opposite offset so that its centerline is coincident with the vertical axis of the positioner. The new target mount is designed to facilitate construction of targets, simplify placing targets on the positioner, and hasten alignment.

The alignment aid positioner is located on top of the chamber, directly opposite the target positioner. It is designed to position a 1- by 4-cm glass reticle. The plane

of the reticle is vertical, with its long axis running north and south. The reticle is carefully positioned in a bezel that allows it to be rotated about the vertical axis and the beam axis. It is positioned by matching the horizontal centerline of the reticle with a line focus created by passing an alignment laser through the cylindrical lenses on the target chamber. The reticle can now be moved in and out of target chamber

Nova Project Completion

center and is a convenient reference for maintaining beam alignment.

Preinstallation tests with a laser interferometer, 0.1- μm resolution linear probes, and a residual gas analyzer allowed us to certify that the positioners would operate as required. These tests allowed us to discover and correct problems that would have severely limited the operational life of the positioners.

Summary. The Novette cryogenic equipment provided a basis for the development of new target and alignment aid positioners for the Nova two-beam target chamber. New internal mechanisms were designed, fabricated, assembled, and thoroughly tested before the positioners were installed on the chamber. The new internal mechanisms provide a fifth degree of freedom (another axis of rotation). The new target mount, incorporated into the target positioner, has worked well enough that it is being developed for use in the ten-beam chamber. Preinstallation tests resulted in calibration tables and curves

that can be used to judge the translational predictability of the positioners. Modifications that followed the preinstallation tests have resulted in reliable and accurate positioners.

Author: L. V. Griffith

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Nova Project Completion

The U.S. Department of Energy (DOE) authorized two final schedule milestones for the Nova project: completion of construction by December 20, 1984 and completion of activation by February 28, 1985. Completion of construction⁴⁰ was demonstrated at the end of 1984 when ten beams

Fig. 5-47. Expenditures during Nova project completion.

Nova Project Completion

were fired simultaneously into the target chamber and converted to 3ω light. For this demonstration, the energy was delivered into calorimeters mounted within the 15-ft-diam target vacuum vessel.

The DOE defined completion of activation, the final milestone, to mean simultaneous delivery of all ten laser beams into the target chamber, focused on a target at significant energy with diagnostic data collected. We achieved this milestone on February 27, 1985 when Nova's ten beams were focused and fired with 2ω conversion onto a gold disk target; both laser and target diagnostic data were acquired. Data from the target diagnostics was captured by energy-balance modules mounted on the chamber and by two x-ray microscopes. Completion of activation signified the beginning of target experiments and thus Nova's readiness as an operational facility.

The Nova project was budgeted at \$176 million. The DOE authorized the use of project contingency funds to complete construction of an initial set of target diagnostic instruments (\$3 million) and to complete Nova activation (\$3.2 million).⁴¹ All major construction on Nova was completed by the end of 1984, and the remaining activation funds were expended in the first two months of 1985 (Fig. 5-47). Through strict cost and budget control, the

	Baseline	Final costs
Nova laser/target system	\$135 310	\$133 785
Nova laboratory/office building	35 990	35 928
Target diagnostics	—	3 005
Nova activation	—	3 238
Contingency	4 700	—
Total	\$176 000	\$175 956
Unused as of 3/1/85		44

project was completed under budget with approximately \$44 thousand in unused funds. (See Table 5-6.)

The project's major subcontractors were invited to be honored guests at the Nova dedication on April 11, 1985. Approximately 195 major subcontractors had contracts of \$25 thousand or more. Twenty-five of those had contracts of more than \$1 million. In all, the Nova project awarded about 7300 purchase orders whose total value was \$129 million. Small and minority businesses were the recipients of 4880 of those orders, whose value was \$58 million. We attribute the successful completion of the project to the excellent collaboration between these industrial suppliers and the LLNL project team that designed and built the Nova system.

Authors: G. J. Suski and F. J. Holcomb

Table 5-6. Nova project final financial status compared with March 1, 1983 baseline budget (thousands of \$).

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A 46-cm laser disk.

A frequency-conversion array.

Various optical technologies are being pursued to increase the output power and energy of Nova. One limit to operating Nova at high power is optical damage to the laser glass caused by the presence of metallic platinum inclusions. The small photo, lower right, shows a millimeter-sized glass fracture resulting from the evaporation of platinum by a high-power laser pulse. We have developed, in cooperation with our vendors, new production techniques where laser glass will be manufactured without introducing platinum inclusions. The photo on the opposite page, top, shows the profile of a new platinum-free, 46-cm laser disk. The new shape, a change from the half-elliptical shape of previous laser disks, is dictated by the use of an inexpensive polymer-bonded edge cladding.

We are also investigating methods of improving the harmonic conversion efficiency and temporal fidelity of the harmonically-converted pulse. A 74-cm-aperture, Nova frequency-conversion array, is shown on the opposite page, bottom. Another consideration in maintaining optimal beam quality in high-power lasers is the influence of stimulated rotational Raman scattering (SRRS). A profile of an 18-cm section of the Nova beam propagated in air to the two-beam chamber is shown below. The photo shows the disk split and the diffractive features; weak SRRS appears as small, bright spots.

6

Laser and Optical Technologies

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Laser and Optical Technologies

Introduction

J. H. Campbell and R. J. Foley

The Nova laser system has evolved to the point that it is routinely used for conducting target physics experiments in the ten- and two-beam chambers. However, we did encounter several unforeseen problems during the activation period that we are now correcting.

The first major problem we encountered was isolated bulk damage in the large-aperture, Nd-doped phosphate glass disks in the Nova laser amplifiers. Microscopic ($\sim 10\text{-}\mu\text{m}$) metallic platinum (Pt) inclusions, which originate from the platinum crucibles used to melt the laser glass in the manufacturing process, are the source of this damage. Although extremely small, the inclusions vaporize when exposed to the high-power $1.06\text{-}\mu\text{m}$ laser light in Nova and produce fracture sites within the glass. These fracture sites vary in size from a few hundred microns to several millimeters and have been found to grow with repeated irradiation.

In May 1985, a group of ICF scientists and engineers was assembled to address the task of eliminating Pt inclusions in phosphate laser glass. As a result of our work here at LLNL and the efforts at the major glass manufacturers, we have developed a fundamental understanding of the damage mechanisms, developed a procedure that reduces the number of Pt inclusions a thousandfold (to only 1 per disk or less), and developed a reliable inspection tool to verify that Pt inclusions in the disks have been eliminated.

During the course of our work on the Nova laser glass, we also investigated alternate edge-cladding designs for the amplifiers. Disk-edge cladding controls parasitic lasing modes that, if left uncontrolled, can reduce the gain of the disks and, in turn, seriously degrade the system performance. Currently, a glass cladding is poured around the edges of the disks, using a process that requires many complex steps. The current edge cladding gives excellent gain performance, but it is expensive to produce. We have examined less expensive, alternate cladding configurations that employ either a composite polymer-glass design or a monolithic absorbing polymer. The performance of a replacement system must meet or surpass that of the present monolithic glass. We plan to complete this effort by mid-1986, and if successful, to use the new cladding on the replacement Nova disks. The cost savings from either of the new configurations should be substantial.

During the initial phase of operation of the Nova laser, experiments were conducted to evaluate the system performance characteristics. These tests uncovered two nonlinear optical effects that adversely influence system performance: stimulated rotational Raman scattering (SRRS) and an intensity-dependent rotation of the beam's polarization field.

SRRS is excited by the propagation of high-intensity light in the long air-filled beam tubes between the spatial filters and the target chambers. Experiments have been conducted on a Nova beamline to measure the SRRS components as a function of the intensity-length product for two laser pulse durations. The results show that the quality of Nova's beams will degrade at high intensity on the long path length to the two-beam chamber. This effect will be eliminated by evacuating or filling the beam tubes to the

two-beam chamber with an inert gas. SRRS is not a significant issue for normal operation of the ten-beam chamber because of its shorter path length.

During the Nova laser performance evaluations, the efficiency of the polarization-tuned harmonic converter crystals was lower than expected. Experiments showed that the polarization field of the Nova beam undergoes a small intensity-dependent rotation while propagating through the systems. This rotation reduces the efficiency of the conversion process. Experiments are being conducted to evaluate several concepts that can be used to correct the polarization rotation problem.

Nova Laser Glass

Introduction and Overview

The large phosphate-glass laser disks in the Nova amplifiers experienced isolated bulk damage at the sites of metallic platinum (Pt) inclusions. Damage occurred at fluences above 2.5 to 3.0 J/cm² (at normal incidence) for pulse lengths of 1 ns. Although this level is high enough to allow us to operate at about 5 kJ (1- μ m wavelength) per beamline at 1 ns, the glass will need to be replaced to allow continuous damage-free operation at the performance levels specified in Table 6-1.

We have demonstrated full performance for single-shot experiments, and in some operational regimes, exceeded the specified levels. For example, we have measured 15 kJ from one beamline at 2.5 ns, indicating 150-kJ total output at 1 μ m (compared with the 80- to 120-kJ objective). However, in order to prevent further damage to the laser glass, we presently are limiting the output of the laser to ~50% of design capability. We intend to stay at this level until the amplifier disks can be replaced with the newly developed inclusion-free glass. We estimate that the replacement process will take 18 months. Fortunately, we have about 24 months of experiments that require power levels less

than 50% of the design values. Thus, we can fully utilize Nova during the glass-replacement period.

The damage sites caused by the laser irradiation of Pt inclusions have been observed in the Nd-doped phosphate-laser-glass disks of both Nova and its predecessor, Novette. The disks of laser glass contain Pt particles, ranging from <5 μ m (detection limit) up to ~100 μ m in size and from roughly 10/L to as high as 1000/L in concentration. The Pt particles come from the platinum crucibles used to melt the laser glass.

The inclusions in the glass cause fractures when irradiated by laser pulses above a threshold value of ~2.5 to 3.0 J/cm² at 1 ns. Based on experiments and computer modeling work, we have found that the damage threshold corresponds to vaporization of the Pt particles as they are heated to temperatures above the normal boiling point of platinum. Both model and experimental results show that the damage threshold should be independent of particle size (for inclusions >0.5 μ m) and increase approximately as the square root of the laser pulse length.

One of the most disturbing observations associated with Pt inclusion damage is that the damage site continues to grow with repeated irradiations. Thus for all practical purposes, *all* Pt inclusions must be eliminated before the glass can be used in the Nova system. The only exceptions may be a few very small (≤ 1 - μ m) inclusions that seem to asymptotically approach a maximum acceptable damage size ≤ 250 μ m.

A Brief Historical Perspective. Platinum inclusion damage in laser and optical glasses has been a recurrent problem throughout the history of high-power

Table 6-1. Specified performance levels of Nova glass compared with present operating capabilities.

Specified Nova performance			Present on-target capability (as limited by Pt damage)	
Wavelength (μ m)	Energy (kJ)	Pulse length (ns)	Energy (kJ)	Pulse length (ns)
1.06	8-12	0.10	8-12	0.10
1.06	80-120	3	40-60	3
0.53	50-80	3	20-40	3
0.35	40-60	3	20-30	3

lasers. In particular, platinum damage in silicate laser glass has been well known for many years (see, for example, the review in Ref. 1). In some laser systems, Pt inclusions were not a problem. Our previous high-power laser, Shiva, used a Nd:silicate glass and we observed no damage in that system. However, current studies of the Nova glass indicate that Shiva operated just below the Pt damage threshold.

It was anticipated that use of a phosphate-based laser glass would reduce the risk of problems with Pt inclusions, because the solubility of platinum is much greater in phosphate glasses than in silicates. When phosphate glass was selected for use in Nova, five 15- x 15- x 2-cm samples were tested with 1-ns, 1.06- μm pulses at 12 J/cm² using a 6-cm-diam beam produced by an arm of the Argus laser. These experiments indicated that the glass might typically contain one small (<10 μm) inclusion in each 10 cm³. This estimate of inclusion density appears to be correct for all but a few of Nova's disks. However, these qualification studies did not correctly predict the large sizes (up to 100 μm) of some of the inclusions that have been found in glass supplied for Nova or, perhaps more important, the growth of the damaged volume that would result from repeated irradiation of the inclusions.

Extent of Nova Damage. The output stage of a Nova beamline contains groups of Nd:glass amplifiers with apertures of 20.8, 31.5, or 46 cm (Fig. 6-1).^{2,3} Each group of amplifiers is followed by a

vacuum spatial filter. A laser pulse passing along the beamline is amplified at one beam diameter to the desired pulse energy. It is then focused into a spatial filter that removes small-scale spatial ripple and enlarges the beam diameter to properly match that of the next group of amplifiers. A spatial filter at the output of each 46-cm amplifier provides a final cleaning of the beam and increases its diameter to 74 cm. This diameter is maintained over the final path to the target chamber.

As a result of this series of amplifications and beam expansions, the variation of pulse energy or beam fluence with position along the chain is an approximately sawtoothed function (Fig. 6-2) with the largest fluence occurring at the output of the 46-cm amplifier. The number of Nova disks that need to be replaced can be determined by comparing the measured damage thresholds with the average fluences at various positions along the Nova beamline. Such a comparison is given in Fig. 6-2, where the shaded horizontal bands indicate the thresholds at 1 and 5 ns. (The thresholds are displayed by a factor of 1.52 higher to account for the fact that the disks are positioned at Brewster's angle.) It is clear that all the disks in the 46-cm amplifiers and most of those in the 31.5-cm amplifiers will need to be replaced with the inclusion-free glass.

Solution to the Inclusion Problem. In May 1985, a group of LLNL chemists, materials scientists, physicists, and engineers was assembled to address the task of eliminating Pt inclusions in phosphate

Fig. 6-1. Schematic diagram of the final three amplifier stages (20.8-, 31.5-, and 46-cm-diam) and spatial filters that are a part of each of the 10 beamlines on Nova.

Fig. 6-2. Beam fluence at normal incidence as a function of position along the final three amplifier stages of a Nova beamline. Amplifier numbers refer to the numbering in Fig. 6-1.

Nova Laser Glass

laser glass. The five major goals of the group were to:

1. Measure the damage thresholds for selected individual platinum inclusions and the rate of growth of the damaged volume surrounding the inclusions.
2. Develop a quantitative model for the damage caused by laser heating of Pt inclusions.
3. Assess the damage on Nova and determine a safe operating range to use during the period of amplifier repair.
4. Develop a reliable quality-assurance technique for inspecting large volumes of glass for the presence of microscopic Pt inclusions.
5. Work with the glass vendors to quantify the mechanisms for platinum introduction and develop processing conditions to eliminate or minimize it.

The first four tasks are complete,⁴⁻⁷ and the significant progress made on task 5 makes us confident that Pt particles can be eliminated from the laser glass. The following articles summarize the results from each of the five tasks.

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Assessment of Laser Glass Damage on the Nova System

Isolated bulk damage sites caused by laser irradiation of platinum inclusions have been observed in the components of the Nova laser—in particular, in the amplifier disks made of Nd-doped phosphate glass. All 10 beamlines on Nova have been subjected to tens of shots with fluences that exceed the threshold for damage due to platinum inclusions in some sections of the amplifier.

To assess the extent of the damage existing in these Nova amplifier disks, we began a microscopic examination of selected disks from the various beamlines. To minimize interference with the laser operation, we chose a few disks that would be statistically representative of all of the disks. These were removed from the beamline and taken to an inspection station in the Nova class-100 clean room. This test stand permits both back and side illumination of the disk with a focused beam from a high-intensity lamp. We located the various damage sites by observing the scattered light either by eye or with a magnifying viewer. Each site location was marked for later examination with a stereo microscope that has a calibrated reticle in one

eyepiece. Using the microscope, we measured the planar dimensions of the damage or inclusion site. With this process, a trained technician can locate most platinum inclusions or damage sites $>10\text{ }\mu\text{m}$ (maximum dimension). Because there are characteristic ring-fracture patterns about the damage sites, we can easily distinguish between damage sites, bubbles, and undisturbed platinum inclusions. A series of fracture lines in an almost circular pattern perpendicular to the beam direction distinguishes a damage site. There is also some damage in a small lobe in either direction along the beam, but it is not as extensive as the radial pattern. Figure 6-3 is a scanning electron micrograph of the characteristic damage pattern. In our examination of these selected laser disks, we found no undisturbed platinum inclusions. This indicates that all platinum inclusions above some minimal size will become damage sites.

Although the microscopic examination distinguishes damage sites from bubbles and undisturbed inclusions, it cannot differentiate between damage sites caused by platinum and other types of inclusions. By coring a cylindrical sample around the damage site and careful radial fracture through the damage site, we can prepare samples in which the initiation site can be subjected to elemental analysis. We found a high concentration of platinum in these

areas for glass that was processed in platinum crucibles. In some sites, there is also a pure platinum kernel at the center of the damage site.

Figure 6-4 shows results from the microscopic examination of several 46-cm laser disk amplifiers. The data have been normalized as the fraction of the total damage sites that are greater than a given size. This was done because the actual number of sites in a laser disk varied from a few tens to a few thousand (0.01 to 1 site per cm^3). The data show a large spread in the sizes of the damage sites. However, the median damage size for a particular

Fig. 6-3. Scanning electron micrograph of a damage site in a laser disk.

Fig. 6-4. Fraction of the total damage sites greater than a given size vs damage-site size.

amplifier position is fairly consistent. The various amplifier positions in a particular beam see different fluences because they are placed serially in the beam path, and each previous disk increases the fluence on the next disk. Refer to Fig. 6-2 in the preceding article for the spatially averaged fluence at the various amplifier positions for 10-kJ, 1-ns pulses and the damage initiation threshold for a platinum inclusion at these positions for the 1-ns pulse. The damage initiation threshold is also discussed in two following articles in this section, "Modeling Laser Damage Caused by Platinum Inclusions in Laser Glass" and "Development of a Rapid-Scan Laser Glass Inspection Facility." Table 6-2 tabulates the beam energies and number of shots of some of the 46-cm amplifier disks tested. Using Table 6-2 and Fig. 6-5, which is a graphic representation of the median damage size as a function of the fluence at different amplifier positions, we can determine the fluence on each disk for each of these energies. We believe the monotonic increase with fluence shown is due to the limited number of shots on the disks, and expect the median damage size to stabilize for all fluences above threshold as the number of shots increases.

These data represent the cumulative damage to a laser amplifier disk over the course of ≤ 50 shots (estimated lifetime,

15 000 shots). A number of the larger platinum inclusions have grown to sizes much larger than $1000\text{ }\mu\text{m}$. We do not have data to tell us how large the inclusions will grow because a limited number of shots have been made on any one amplifier disk. More information on the growth of individual inclusions is in a later article, "Modeling Laser Damage Caused by Platinum Inclusions in Laser Glass." From the data that we have obtained, it appears that all amplifier disks have damage sites larger than the allowable specification when fluences at the disk position are $> 3.8\text{ J/cm}^2$ at a 1-ns pulse length. These damage sites will continue to grow whenever the fluence exceeds the 3.8 J/cm^2 threshold.

Our evaluation shows that all disks operated at fluences greater than the damage threshold will develop unacceptable damage. Thus, we will have to replace this glass with a platinum-particle-free glass before the laser can be operated for extended periods at full design fluence. However, operation of the disks below this threshold ($\sim 50\%$ of design fluence) will not further endanger the laser for continuing experiments.

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Table 6-2. Position, energy, and number of shots on 46-cm laser disk samples.

Sample	Beamline	Position	Shot energy			
			5 kJ	5-7 kJ	7-9 kJ	9 kJ
HP5-81	2	9	15	5	8	7
HP5-85	5	11	13	7	10	8
HP5-72	2	11	15	5	8	7
HP5-82	4	12	14	5	11	5
HP5-108	2	11	15	5	8	7
HP5-84	9	11	34	9	15	9

Fig. 6-5. Median damage size vs average fluence at four amplifier disk locations.

Large-Aperture Testing of Laser Glass

After the analysis of laser amplifier disks that were removed from the Novette laser and of a few disks removed from the arms of Nova, it became apparent that the growth rate of small platinum inclusions was a substantial problem. Because tests for platinum inclusion damage in a small-beam irradiation facility had not obtained the data to predict the extent of this damage, we needed to irradiate larger sample areas at fluences above damage threshold.

Two stations on the Nova laser were used to provide large-aperture testing capabilities for amplifier glass samples (Fig. 6-6) from small research melts and from large (40-L) melts. For the large samples ($30 \times 30\text{ cm}$), we inserted a Nova 74-cm focus lens ($f/4.3$) in beamline 10 in the

Fig. 6-6. Large-area damage test configuration used to irradiate samples at 7 to 10 J/cm².

Table 6-3. Typical damage data obtained at the Nova beamline 9 and 10 test facilities.

laser switchyard area. After the 74-cm lens, the beam was passed through a negative lens ($f/111$) and collimated to 25 cm (Fig. 6-6). We irradiated a 25-cm-diam spot (about half of a 46-cm disk) at fluences up to 10 J/cm² (5000 J, 1 ns). However, self-focusing and other variations caused small intense spots in the beam. In our initial irradiation at 10 J/cm², we noticed significant bulk damage, especially in the second of two samples placed serially in the beam. This damage obscured the platinum inclusion growth, so we decided to continue testing at ≤ 8 J/cm². Although below design fluence, 8 J/cm² is 2 to 3 times the fluence required for initiation of growth of the platinum inclusion fractures. The smaller test facility was placed in beamline 9 just after the 20-cm amplifier. The resultant 30-cm beam was reduced and collimated to 6.3 or 8.5 cm (Fig. 6-6). This beam was very uniform; several tests were done at 10 J/cm² (500 J, 1 ns) with no apparent bulk damage to the test samples, in contrast to observations at beamline 10.

Data from the samples tested were valuable in guiding glass processing development at Hoya and Schott, the two suppliers. Table 6-3 contains typical data after 10 to 15 irradiations. This test data agreed well with the actual damage found in irradiated laser glass, and the test bed could have been used to qualify replacement glass for the laser. However, the complexity of damage testing on such a large facility and the slow data-return rate (1 laser disk per day) made it undesirable to use our test system for quality control of future laser-glass production.

We also found that preliminary inspection for undamaged platinum inclusions was not a good indicator of laser-glass performance. Table 6-4 lists the number of inclusions that were not detected by our best analysis techniques. Thus, unreliable inspection techniques, a slow data return from the Nova test systems, and the need to examine the total volume of every disk

Site No. ^a	Size (μm) of inclusion (I) or damage site (D) found:		
	Initially	After 8.5-cm-diam test	After 24-cm-diam test
1	10 (I)	10 (I) ^b	10 (I) ^b
2	15 (I)	15 (I) ^b	325 (D)
3	— ^c	25 $\times$ 80 (I) ^d	<2
4	— ^c	— ^b	25 (D)
5	5-10 (I)	5-10 (I) ^b	225 (D)
6	5-10 (I)	5-10 (I) ^b	5-10 (I) ^b
7	5-10 (I)	5-10 (I) ^b	5-10 (I) ^b
8	— ^c	<5 (I) ^b	175 (D)
9	— ^c	<5 (I) ^b	175 (D)
10	— ^c	— ^c	125 (D)
11	— ^c	— ^c	20 (D)
12	— ^c	— ^c	25 (I) ^b

^a Sites 1-4 are from sample NE51652-2, dimensions 28 $\times$ 30 $\times$ $\approx$ 3 cm, volume 2.7 L; sites 5-12 are from sample NE51652-4, dimensions 28 $\times$ 30 $\times$ $\approx$ 3 cm, volume 2.7 L.

^b Outside irradiation zone.

^c Not detected.

^d Non-platinum inclusion.

Sample	Inclusions	Damage sites
NE51652-2	2	3
NE51652-4	3	8
FP20042	0	20
D678-1	0	2

Table 6-4. Number of platinum inclusions found by inspection and the number of damage sites found after laser irradiation.

dictated the construction of the high-repetition-rate test system described in the article, "Development of a Rapid-Scan Laser Glass Inspection Facility," later in this section.

During testing for platinum-inclusion damage, we also attempted to measure inclusion growth rate. However, the limited analytical equipment and the short time interval between shots made on-line measurements difficult. For larger inclusions, where damage became obvious after one or two shots, we obtained some data (Fig. 6-7) showing that inclusion growth does not stop after a few shots. Later samples, made of improved glass, could only be examined for damage sites in the Nova clean room station, and no individual shot data could be obtained. Inclusions in these samples were measured before irradiation, and again after completing a selected number of shots. It appears that the damage from any particle above some

Fig. 6-7. Size of platinum-inclusion damage vs number of shots.

Fig. 6-8. Average number of damage sites per disk vs type of manufacturing process illustrates progress in the elimination of platinum inclusions.

minimal size (1 to 4 μm) will grow larger than the permissible 250 μm .

Figure 6-8 illustrates the remarkable improvement that has been made in reducing the number of observed damage sites in a laser amplifier disk with a 7-L volume. The number of platinum inclusions is several orders of magnitude lower, which comes very close to meeting the required specification for the replacement glass.

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Measurement of Damage Initiation and Growth due to Metallic Platinum Inclusions

We measured the threshold fluences that initiate fracturing in glass that contains platinum inclusions. We repeatedly irradiated inclusions to determine the rate of damage growth to an unacceptable beam-obscurtion size, calculated as 250 μm by W. E. Warren using the beam-propagation code MALAPROP. Irradiation at 2 to 2.5 times the damage threshold caused damage growth to $>250 \mu\text{m}$ around most inclusions tested. The damage threshold was independent of inclusion size at pulse durations of 1.3, 9, and 50 ns.

Selected inclusions in small glass samples were placed in the center of a laser beam and irradiated with 1.064- μm pulses of 1.3-, 9-, or 50-ns duration. For each shot, fluence at the beam center was determined to within $\pm 7\%$. The beam had an almost Gaussian fluence distribution and was about 2 mm in diameter at the e^{-1} level. In such a beam, fluence over an area with diameter of 0.65 mm is $\geq 90\%$ of the peak on-axis fluence, so the beam was large relative to the 250- μm damage size of interest.

Each inclusion received a sequence of irradiations, usually beginning at a fluence near 1 J/cm^2 . After each irradiation, the inclusion was inspected by optical microscopy. If there was no apparent alteration of the inclusion, it was irradiated again at higher fluence. Once damage was initiated, the fluence in the 1.3-ns experiments was increased to between 5 and 9 J/cm^2 , a range typical of the highest fluences anticipated in Nova. The inclusion was repeatedly irradiated until the size of the damaged volume either stabilized or reached a dimension much larger than 250 μm . In experiments with either 9- or 50-ns pulses, the fluences used to observe evolution of damage were selected to be 2 to 2.5 times the threshold for initiation of damage.

The samples containing inclusions were cut from a 20.8-cm phosphate-glass amplifier disk that had been used in Novette for more than 500 shots at fluences of 0.5 to 1.0 J/cm^2 . The major dimension of the inclusions ranged from 4 to 75 μm . None of

these inclusions had caused damage during these 500 shots at low fluence.

Threshold damage was defined as one or more fractures extending a few microns into the glass surrounding the inclusion. These initial fractures had shapes that were peculiar to individual inclusions. During a few irradiations at fluences above threshold, the damage usually evolved into a circular planar fracture oriented perpendicular to the incident beam. The center of the fracture contained structure that obscured the particle itself, and we could not determine by optical microscopy whether the damage expanded along the direction of the beam.

In most experiments, repeated irradiation of the inclusion caused damage growth to sizes larger than $250\text{ }\mu\text{m}$. This is illustrated in Fig. 6-9 by data obtained in the test of the largest inclusion (major dimension, $75\text{ }\mu\text{m}$). The inclusion survived three irradiations at fluences up to 2.4 J/cm^2 , and was damaged during the fourth shot at 3 J/cm^2 . Growth of the damage was rapid during the next two shots at higher fluence, but ceased when the fluence was reduced to $<2.5\text{ J/cm}^2$. Four additional irradiations at higher fluences caused further growth of the damage site to well over $250\text{ }\mu\text{m}$.

For three inclusions, damage first increased and then remained stable under repeated irradiation. Data obtained in the test of a $4\text{-}\mu\text{m}$ particle are given in Fig. 6-10. The particle survived 5 irradiations at fluences as high as 2.5 to 3.7 J/cm^2 . Damage stabilized after the 14th shot at 12 J/cm^2 , and did not change during the last 27 shots.

At all three pulse durations, thresholds were independent of particle size. This was an expected result. The characteristic dimension for heat flow in platinum is $0.15\text{ }\mu\text{m}$ in 1 ns and only $1.1\text{ }\mu\text{m}$ in 50 ns . All inclusions tested exceeded $4\text{ }\mu\text{m}$, so heat could not be distributed during a laser pulse and each point on a surface was relatively unaffected by the presence of the remainder of the particle.

Thresholds measured at the three pulse durations, 1.3 , 9 , and 50 ns , were 2.5 ± 0.3 , 4 ± 0.4 , and $8 \pm 0.8\text{ J/cm}^2$, respectively. As illustrated in Fig. 6-11, these thresholds are well represented by the functional form t^α , where t is pulse duration and α is

Fig. 6-9. Growth in damage around a $75\text{-}\mu\text{m}$ inclusion (initial size) during 12 irradiations by 1-ns , $1.064\text{-}\mu\text{m}$ pulses.

Fig. 6-10. Growth in damage around a $4\text{-}\mu\text{m}$ inclusion (initial size) during 1-ns, $1.064\text{-}\mu\text{m}$ irradiations.

an empirical constant of about 0.3 for our tests. However, all models for damage due to laser irradiation predict that the threshold should vary as $t^{0.5}$, as will be discussed below. The experimental value $\alpha = 0.3$, comes primarily through accommodation of the threshold measured at 1.3 ns . A scaling through the data at 9 and 50 ns yields $\alpha = 0.4$, which is in closer agreement with the theoretical value.

Fig. 6-11. Average threshold fluence of all particles tested at pulse durations of 1.3, 9, and 50 ns; the calculated threshold fluence as a function of pulse duration, $t^{0.3}$ (solid line), illustrates the variation between calculated and observed fluences.

Laser pulse duration (ns)	Test ^a fluence (J/cm^2)	Number ^b of shots	Initial particle size (μm)	Final size of damage	
				Stable (μm)	Growing (μm)
1	6-10	36[9]	4	176	—
1	5-7	20[14]	6	240	—
1	5-7	49[6]	8	75	—
1	5-7	9	18	—	410
1	3-9	8	25	—	750
1	4-7	7	70	—	1000
9	5-9	9	17	—	380
9	5-9	8	18	—	440
50	8-20	6	6	—	340
50	9-20	5	18	—	390
50	8-20	6	19	—	500

^a Range of fluences for which growth of damage was observed.

^b Number of shots at stated fluence range. Bracketed values are numbers of shots during which damage grew prior to stabilization.

Table 6-5. Size of damage surrounding particles with various initial dimensions.

We do not yet have an explanation for the difference between the calculated and observed scalings of thresholds. Experimental errors could come only through having assigned too large a threshold during some measurements. That is, damage was observable at the three stated values of fluence, so thresholds cannot exceed those values. Therefore, the small slope of the experimental data could only be due to overestimation of the threshold at 1 ns. It could be argued that the threshold damage by 1-ns pulses was so slight that it was not observed, and that the assigned 1-ns threshold is actually the fluence needed to make the damage visible. However, that argument suggests that the measured values at 9 and 50 ns might also be too large, so it is not certain that visibility of damage can account for the small slope of experimental values. Two additional observations place some bounds on the 1-ns data. First, the 1-ns threshold is not less than $\sim 1 J/cm^2$, because all these inclusions had survived many irradiations at that fluence. Second, the level of damage found during

inspection of a selection of disks in Nova can be explained if the threshold at 1 ns is near $2.5 J/cm^2$. Additional experiments and modeling are required to clarify the scaling of thresholds with pulse duration.

All the data from our experiments are summarized in Table 6-5. Three particles with initial dimensions of 4, 5, and 8 μm caused damage that remained smaller than 250 μm under repeated irradiation by 1-ns pulses at fluences of 5 to 9 J/cm^2 . Except for one 6- μm particle that was irradiated by high-fluence (12 to 20 J/cm^2) 50-ns pulses, only particles with an initial dimension greater than 10 μm caused damage sites exceeding 250 μm . Therefore, there is reason to believe that inclusions smaller than 5 μm might be acceptable in disks for Nova, although the size of acceptable particles is probably a function of fluence.

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Modeling Laser Damage Caused by Platinum Inclusions in Laser Glass

Platinum inclusions in neodymium-doped phosphate laser glass on Nova and Novette cause damage when fluences are greater than $\sim 2 J/cm^2$. When laser light is absorbed by the opaque platinum, it vaporizes a thin skin of platinum on the front surface of the inclusion. A shock wave propagates through the inclusion and is transmitted to the glass. Glass, being brittle, cracks under the force of the shock wave and the pressure of the vaporized platinum. These cracks propagate until the pressure of the vaporized platinum drops below a critical value that depends on the material and the size of the crack. Our studies of inclusion damage, reported below, have given us information about the shape of fracture sites, and the relation between inclusion size and fluence damage limits.

Hopper and Uhlmann⁸ calculated that temperatures would exceed 10 000 K with fluences of 20 J/cm^2 over a 30-ns period. They theorized that very small inclusions were safe because heat losses to the glass kept the temperatures and, hence, stresses within allowable bounds. However, they

argued that inclusions larger than 1000 Å (0.1 μm) in diameter produced stresses that exceeded the strength of the glass and could cause fracturing. Their thermal modeling suggested that of such inclusions, large ones are less damaging than intermediate-sized ones. Surface-to-volume ratio decreases as the radius increases, so larger inclusions have more thermal capacity and hence lower average temperatures and stresses.

Sparks and Duthier⁹ also examined the inclusion problem and showed that glass damage was greater with high-intensity, short-pulse lasers than with low-intensity, long-pulse lasers. A number of authors¹⁰⁻¹² have studied heating of a sphere illuminated from one side. Their solutions all show that the high temperatures predicted by Hopper and Uhlmann will exist.

Our investigation of inclusion damage was broken into three parts: a fracture-mechanics theory, a crack-development hypothesis, and a thermal model. Our major conclusions, listed below, are followed by a detailed discussion of each approach.

- Fracture-mechanics theory suggests that crack propagation rates are proportional to the square root of the inclusion radius and the 3/4 power of the fluence. Inclusions over ~1 μm in diameter cause more damage than smaller inclusions.

- A crack-development hypothesis based on shock propagation successfully predicts the shape of many fracture sites: an annular crack around the inclusion equator with planar lobes at the poles. The lobe at the pole exposed to laser light is larger than the lobe at the unexposed pole.

- Thermal modeling predicts that glass damage will occur if the platinum vaporizes and that the fluence-damage threshold is both proportional to the square root of the laser pulse length and independent of the inclusion diameter above about 1 μm. Predicted fluence-damage limits of 1.4 J/cm² for a 1-ns pulse and 4.3 J/cm² for a 10-ns pulse agree with experimental data.

Fracture-Mechanics Theory. Consider glass to be a brittle material that initially has a penny-shaped crack with the same diameter as that of the platinum inclusion. The criterion for crack propagation is based on a material property, K_I , called the

fracture-toughness (stress-intensity) factor¹³

$$K_I = \sigma(\pi a)^{1/2}, \quad (1)$$

where

σ = applied stress [pressure (Pa)],

a = crack radius (initially equal to the inclusion radius m).

If K_I exceeds a critical value, then the crack will propagate. This critical value, termed the critical-fracture-toughness or critical-stress-intensity factor, K_{Ic} , is a property of the material; its magnitude is normally obtained experimentally.

The Mie-Grüneisen equation of state^{14,15} is applicable for shock-wave/high-stress values, where energy is deposited isochorically; it relates stress and specific energy as

$$\sigma = \Gamma \rho e, \quad (2)$$

where

Γ = Grüneisen constant (dimensionless),

ρ = density (kg/m³),

e = specific energy deposition above that required for vaporization (J/kg).

Following Hopper and Uhlmann⁸ for an inclusion with nonzero thermal diffusivity and a 10-J/cm², 1-ns pulse, and using their material properties [see their Eqs. (6) and (7b)], we approximate the spatial temperature distribution T at 1 ns and 10 J/cm² as follows:

$$T \simeq (4 \times 10^4) \times \left(\frac{1 \text{ ns}}{30 \text{ ns}} \right)^{1/2} \times \frac{10 \text{ J} \cdot \text{cm}^{-2}}{0.7 \text{ J} \cdot \text{cm}^{-2}} \times \text{ierfc} \left[\frac{0.6 z}{(1 \text{ ns}/30 \text{ ns})^{1/2}} \right],$$

and thus

$$T = 1 \times 10^5 [\text{ierfc}(3.286 z)]. \quad (3)$$

Here z is the distance into the inclusion from the surface in microns, and T is temperature in degrees Kelvin. Values of ierfc , the integral of the complementary error function, are tabulated by Abramowitz and Stegun.¹⁶ The surface temperature drops below the vaporization temperature (4800 K) at $z = 0.307 \mu\text{m}$. If we double the fluence to 20 J/cm², the surface

temperature drops below the vaporization temperature at $z = 0.371 \mu\text{m}$. If we use these values to calculate the power dependence of the fluence J on the specific energy deposition e (recalling that specific energy deposition is approximately proportional to temperature), then

$$\left(\frac{J}{J_0}\right)^m = \frac{\rho e}{\rho e_0} \approx \frac{T}{T_0},$$

$$\left(\frac{20 \text{ J} \cdot \text{cm}^{-2}}{10 \text{ J} \cdot \text{cm}^{-2}}\right)^m = \frac{20 \text{ J} \cdot \text{cm}^{-2} / 0.371 \mu\text{m}}{10 \text{ J} \cdot \text{cm}^{-2} / 0.307 \mu\text{m}},$$

and

$$m = 3/4. \quad (4)$$

Combining this result with Eqs. (1) and (2), we have

$$K_1 \sim J^{3/4} a^{1/2}. \quad (5)$$

The crack will propagate as long as $K_1 > K_{1c}$. Hence, this model suggests that the extent of crack propagation in the glass per shot is proportional to the square root of the crack radius and the three-quarters power of the fluence. Initially, the crack radius is equal to that of the inclusion but it would increase after each shot. Hence, at constant fluence, the rate of crack propagation would increase after each shot.

The model does not account for loss of pressure or for convection heat transfer as the vaporized platinum flows into the crack. The predicted crack propagation rate in later shots is counteracted by the pressure loss and heat transfer effects, both of which oppose the square-root dependence on the crack radius. This could partially account for some experimental data¹⁷⁻²⁰ that show that the amount of crack propagation first increases and then decreases, or alternatively, that the damage propagates and then stops. Cracks could also stop propagating in later shots if all the platinum initially in the inclusion has been deposited in the crack area with a thickness that is now transparent to laser light. In any case, this model suggests that the larger the inclusion, the worse the damage.

Additional development of this fracture-mechanics theory could lead to a better prediction of how far a crack would propagate during a particular shot. In order to proceed, we need to estimate the thickness (width) of a typical crack so that we can calculate the pressure distribution in the crack as a function of time. Further, we may need to account for the change in opacity below unity of platinum deposited in the crack region if the platinum thickness is 100 Å or less.

Crack-Development Hypothesis. Many damage sites have an annular crack around the equator of the inclusion with an additional planar lobe-shaped crack at each pole^{20,21} [Fig. 6-12(a)]. The crack at the pole facing the laser light illumination is always larger than the crack at the pole hidden from the laser light by the inclusion itself.

Hovingh²² suggested that the phenomenon was a result of shock propagation. As energy from laser light is deposited on the exposed surface of the inclusion, a shock wave develops and propagates through the inclusion [Fig. 6-12(b)]. The inclusion deforms to an oval shape. This puts pressure at the sides of the inclusion and initiates an annular crack site at the equator. The pressure over the annulus is uniform because of symmetry, and a completely symmetric annular crack site is formed.

Even though the forces are compressive, pressure causes glass to crack because it

Fig. 6-12. (a) Typical fracture site around an inclusion; (b) shock-propagation analogy showing changes in deformation responsible for this geometry.

Fig. 6-13. (a) Hemispherical model of a platinum inclusion in glass illuminated from the top with laser light; (b) a typical non-spherical inclusion.

is applied not only perpendicular to the incoming laser light, but also parallel to it. This tends to spread the glass material in tension parallel to the incoming laser light. Analogous forces occur in the walls of a spherical pressure vessel, where the hoop stress is tension and yet there is pressure exerted on the inside wall surface. In addition, any compression generated at the crack site changes to tension as the shock passes through the glass material itself.

The shock continues through the inclusion, dissipating energy and becoming weaker with time. The inclusion returns to a circular shape, but then deforms to an oval with major axis parallel (rather than perpendicular) to the direction of the incoming laser light. However, this time the oval has more deformation at the poles because the energy deposited is concentrated there. Polar crack sites are initiated at the poles only and do not form a complete annulus around the inclusion.

The shock propagation through laser glass produces the same qualitative deformation that occurs when a golfer hits a good drive. The golf ball has energy deposited at the surface and deforms to an oval with the major axis perpendicular to the direction of the golfer's swing. During flight, the golf ball returns to a circular shape before it deforms to an oval with major axis parallel to the direction of the golfer's swing.

The crack-development hypothesis explains why many crack sites have the form shown in Fig. 6-12(a). The hypothesis suggests the location of crack initiation sites. However, it is the vaporized platinum and its associated pressure that actually propagate the cracks beyond the initial site. Because the platinum vapor must flow farther to reach the crack site at the pole hidden from the laser light, less pressure develops there. A smaller crack should occur at the hidden pole than at the exposed pole. Observations of damage sites in laser glass verify this size difference.

Thermal Modeling. Our thermal model [Fig. 6-13(a)] was a half-sphere of platinum surrounded by a half-spherical annulus of glass; the distribution or zoning included in the finite-element code we used for the analysis is also shown. The model represents the typical inclusion in Fig. 6-13(b). Vaporization of the platinum is concentrated at the top of the inclusion, which in both cases is normal to the incident laser light. Hence, the spherical model is a good approximation to the more generally shaped inclusion of Fig. 6-13(b). A half-sphere in the model is sufficient because of symmetry. In the model, laser light was absorbed on the upper surface of the platinum sphere with a cosine distribution to account for the curvature of the surface.

Thermal properties used in the model for platinum and glass are shown in

Table 6-6. Platinum properties were obtained from Touloukian and DeWitt²³ and Hampel.²⁴ The boiling point shown in Table 6-6 is higher than in some other references (e.g., Weast²⁵ shows 4100 K), but is used for consistency with other properties taken from Hampel.²⁴ The key platinum properties are the heat of vaporization and the normal spectral absorptance, so the difference in the documented boiling point has an insignificant effect. The normal spectral absorptance α (platinum, 0.27; glass, 20) is at room temperature. The absorptance does change with temperature, but it is more a function of the surface itself. Changes of an order of magnitude are possible depending on whether the surface is polished or not and if the surface is oxidized. We used the

Table 6-6. Properties of platinum and glass used in the thermal model.

Thermal properties	Platinum inclusion	LHG-8 phosphate glass
Density (kg/m^3)	21 400	2 850
Specific heat ($\text{J/kg} \cdot \text{K}$)	146	263
Thermal conductivity ($\text{W/m} \cdot \text{K}$)	80	0.52
Melting point (K)	2 040	—
Heat of fusion (J/kg)	101 000	—
Boiling point (K)	4 800	—
Heat of vaporization (J/kg)	2 260 000	—
Normal spectral absorptance at $1 \mu\text{m}$	0.27	—

Fig. 6-14. Theoretical results of laser pulse length on the fluence damage threshold are independent of inclusion size above $\sim 1 \mu\text{m}$ in diameter; experimental data from Refs. 31-34 are also shown.

room temperature values of α because no elevated temperature values were available. The equation of state for platinum above the boiling point was that used in the LASNEX code.²⁶ Properties of LHG-8 phosphate glass were obtained from Stokowski, Saroyan, and Weber.²⁷ The surface of the platinum was allowed to radiate to a 300-K surrounding. No absorption was permitted in the glass, but heat was conducted between the glass and the platinum with negligible contact resistance.

Modeling was performed using Shapiro's TOPAZ²⁸ code; Hailquist's MAZE²⁹ and ORION³⁰ codes were used for preparing the input and processing the output. TOPAZ is a heat-conduction code, which is applicable to this problem because inclusion sizes are so small that natural convection is negligible. Problems were run on the Cray computer with an average run time of ~ 30 min each.

Results are shown in Fig. 6-14, where the fluence damage threshold is plotted against laser pulse length. Agreement with experimental data³¹⁻³⁴ is excellent. The error bars used with data from Avizonis and Farrington³¹ and Yamanaka et al.³² are also shown. The theory we developed is represented by the solid line; the slope of $1/2$ corresponds to a square root dependence on laser pulse length. The slight difference between our model and the data of Gonzales and Milam³³ at short laser pulse lengths may result because crack propagation per shot is so small in this region. Hence, it takes a number of shots before any damage can be physically observed.

Temperature profiles produced by laser illumination of a $5\text{-}\mu\text{m}$ -radius platinum inclusion at two different times (1 and 10 ns) are shown in Fig. 6-15. Peak temperatures are high (electron volt range; $1 \text{ eV} \approx 10\,000 \text{ K}$) and so are the resulting pressures of the vaporized platinum. Note that the temperature contours are all concentrated in the platinum at the exposed surface. Above $\sim 1 \mu\text{m}$ in diameter, there is no change in the temperature profiles with a change in the size of the inclusion because the thin layer where the profiles are located is similar to a planar surface. Table 6-7 lists maximum temperatures, maximum pressures, and the resulting crack propagation per shot for various

Fig. 6-15. Computer-generated temperature contours for a 5- μ m-radius platinum inclusion surrounded by phosphate glass; illumination is 1- μ m laser light for 1 ns with an incident fluence of 3.7 J/cm².

Table 6-7. Calculated characteristics of crack propagation.

Incident fluence ^a (J/cm ²)	Absorbed fluence (J/cm ²)	Pulse time (ns)	Maximum temperature (K)	Maximum pressure (Pa)	Crack propagation per shot (μ m)
37	10	1	1.3×10^5	2.8×10^{11}	500
18	5	1	7.6×10^4	1.5×10^{11}	150
3.7	1	1	1.5×10^4	0.5×10^{11}	8.4
1.8	0.5	1	6.5×10^3	$<0.5 \times 10^{11}$	3.5
1.5	0.4	1	5.3×10^3	$<0.5 \times 10^{11}$	1.5
3.7	1	10	4.5×10^3	—	—

^a Incident fluence damage threshold is 1.4 J/cm² for a 1-ns laser pulse length.

fluences. The maximum pressures were calculated using Eq. (2). We used a Grüneisen constant Γ equal to unity; the value of ρe was calculated from the TOPAZ results at the time of the peak temperature by dividing the total energy in the vaporized platinum at the exposed pole by the volume of the vaporized platinum.

We have no doubt that the pressures are high enough to propagate a crack. In conjunction with experimental results,¹⁷⁻²¹ this suggests that damage will occur if the temperature of the platinum exceeds its boiling point. Experimental data (Fig. 6-14) suggest a damage limit of about 2 J/cm² for a 1-ns pulse and 5 J/cm² for a 10-ns pulse. This is in agreement with the calculated values.

The fracture-mechanics theory is not well enough developed to predict crack growth; however, we can make a conservative estimate by multiplying the time during which platinum is in a vaporized state by the speed of sound in glass. Predicted values of crack growth for fluences close to those used in the experiments cited here are reasonable (see Table 6-7).

The key conclusions of this thermal analysis are that damage (1) occurs if the platinum temperature exceeds the boiling point, (2) is independent of inclusion size above $\sim 1 \mu$ m, and (3) is worse for laser pulses of shorter duration.

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Development of a Rapid-Scan Laser Glass Inspection Facility

We have constructed a facility to inspect for the presence of microscopic inclusions in optical glass components. We can now inspect large optics quickly and reliably for inclusion volume concentrations in the parts per quadrillion (10^{15}) range.

Most glass optical components for high-power lasers are melted and homogenized in platinum-lined crucibles. This glass melting and casting process potentially leads to the presence of microscopic platinum inclusions. In situations where large optics are exposed to high laser fluences, such as the Nova laser, these inclusions can significantly damage optical components even when present in spatial densities as low as one inclusion per component. Previously, we used direct visual examination or optical microscopy to inspect for these particles; this limited reliable detection to inclusions larger than 10μ m in diameter. Unfortunately,

inclusions much smaller than $10\text{ }\mu\text{m}$ can initiate laser damage.

The size of the inclusions that must be located extend into the submicron region—a territory in which conventional optical microscopy is ineffective. Light-scattering methods are unwieldy because bubbles, surface imperfections, and dust are indistinguishable from inclusions even when the sample is of high optical quality. Signals from these imperfections swamp the signals of interest because the number of inclusions is quite low, on the order of $10^{-4}/\text{cm}^3$.

The detection of inclusions early in the production of optical components allows us to discard unsatisfactory casts before final cladding and polishing, and significantly reduces the cost of the next set of Nova glass. To achieve this, we require a system capable of rapid and reliable inspection of Nd:phosphate glass samples for inclusions that absorb at $1.06\text{ }\mu\text{m}$, such as platinum and carbon.

We adopted a destructive technique for locating inclusions as the most reliable method available. A test facility has been constructed that detects small opaque inclusions in large transparent components; it uses a commercial laser to deliver high energy pulses at moderate frequency to the test sample in a small diameter beam. A schematic diagram of the facility is shown in Fig. 6-16. When present, platinum inclusions are easily detected after testing by the characteristic macroscopic fracture site that is created in the glass surrounding the inclusion.

A commercial, Q-switched YAG laser (JK Model HY 750) is focused on the sample, which is scanned in the x and z directions by a commercial motor-driven stage (Daedal Models 47215 LH and MLH). The output energy of the laser is 800 mJ at 8-ns pulse width and 20-Hz pulse repetition frequency. The 8-mm diam output beam is focused by a 2-m positive lens to 2 mm in diameter where it strikes the test

Fig. 6-16. Schematic of the rapid-scan inspection equipment used to detect small opaque inclusions in laser glass.

piece mounted on the x-z stage. The con-focal parameter of this beam is large with respect to the thickness of the test piece.

To produce damage sites at each inclusion that are large enough to inspect optically, each site in the sample is irradiated at normal incidence with 10 shots, each above 10 J/cm^2 . This fluence corresponds to twice the damage threshold for platinum inclusions in glass at the 8-ns pulse widths used here.³⁵ We tested finished amplifier disks from Nova, ranging in volume up to about 10 liters, and numerous samples from test melts. In practice, the device is capable of scanning a 46-cm-aperture phosphate-glass amplifier disk that is 7000 cm^3 in volume in about 4 h using 10 shots per site. Reducing the number of shots proportionately reduces the scan time. The optical examination of the damaged disk takes an additional hour or so, depending on the number of damage sites that must be characterized and mapped for later evaluation.

The device is located in a class-1000 clean room to avoid dust contamination and consequent damage to the device optics or the test samples. A laboratory computer (Hewlett Packard 9816S), interfaced through a general-purpose, data-acquisition control unit (HP 3421A), controls the laser and the stage drive. If dust or laser damage in the optical train upstream degrades the laser energy, the computer that monitors the output readings from the laser calorimeter interrupts the test. Automatic shut-off of the laser and the stage drive also occurs if the sample reaches position limit triggers, or if high temperatures are reached in the laser cooling-water system.

This laser glass test facility offers a unique capability in materials damage assessment. Prior damage testing has generally fallen into one of two categories. In one method, beams of 1 mm or less were used at high fluences to test several spots on the sample,³⁶ and a statistical analysis was performed on the data. The limitations of that method in situations where the damaging entities have a very low spatial density are obvious. In the second approach, the Nova laser was used to give high fluences over an entire area in single shots, but at a very low repetition rate.³⁷

Aside from the expense of the latter method, our new device provides higher throughput and is less cumbersome. With our new device, we can now inspect large optics quickly and reliably for inclusion volume concentrations in the parts per quadrillion (10^{15}) range.

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Development of Melting Technology to Eliminate Metallic Platinum Inclusions

Staff members at LLNL are working with scientists at Hoya Optics, Inc. and Schott Glass Technologies, Inc., two Nova laser glass vendors, to eliminate metallic platinum inclusions in phosphate laser glass. Both Hoya and Schott have been able to modify their melting processes to eliminate external sources of contamination and to keep the remaining platinum in solution. Some details of the glass manufactured at these facilities are proprietary, so progress can only be reported in general terms. An excellent discussion of the technology used in preparing optical glasses has recently been published by Izumitani.³⁸

Preparation of Laser Glass: A Brief Description. Nova laser glass is prepared in a two-step batch melting process (Fig. 6-17). In the first step, a cullet glass is made by melting raw starting materials. This cullet is the feedstock for the final melting of the optical quality laser glass.

The cullet melting is usually done in a relatively inert refractory crucible. The product of the melt is poured onto a cold surface, where it fractures into smaller pieces as a result of the thermal stresses. The cullet is generally full of bubbles, striae (chemically inhomogeneous regions of glass with different indices of refraction), and possibly some small particles of starting material that were not completely dissolved in the glass.

The fractured glass produced from the cullet melt is used in the final melting step, performed in an ~40-L platinum crucible. The use of platinum vessels has long been the standard in the glass industry because platinum is erosion-resistant and

Fig. 6-17. Simplified view of the two-step melting process used in the preparation of Nova laser glass.

Fig. 6-18. Vapor pressure of pure Pt and that of PtO_2 in equilibrium with pure Pt metal and 1-atm O_2 .

essentially insoluble in the glass matrix. Striae from erosion or corrosion of the melting vessel are not generated in platinum pots. In contrast, most refractory materials are oxides (e.g., Al_2O_3 , SiO_2 , MgO) that are soluble in the oxide-glass matrix and cannot be used in the production of large pieces of striae-free optical glass.

During the final processing step, the glass melt is mechanically stirred to homogenize the glass and remove any striae that might have been produced by

volatilization.³⁴ The glass is also treated to remove bubbles. At the end of the final melt cycle, the glass is cast into a mold and annealed.

Effect of Platinum Chemistry on Inclusion Formation. Platinum metal is widely used as a container material for processing optical-quality glass because it is largely inert to most molten glasses. Nevertheless, some small (ppb to ppm) level of contamination by Pt particles is known to occur. Except for high-power laser glass, the platinum does not affect optical performance.

The source of platinum-particle contamination in glass has been the subject of several studies.³⁹⁻⁴³ It is generally agreed that the particles originate from (1) abrasion or wear, (2) direct dissolution of the crucible and subsequent precipitation of the dissolved platinum, and (3) vapor-phase transport as $\text{PtO}_2(\text{g})$ and redeposition as Pt. In principle, abrasion and wear can be eliminated by care in processing, and have usually been ruled out as a source of contamination.

Platinum is an unusual metal in that it forms the volatile oxide $\text{PtO}_2(\text{g})$ when heated above $\sim 700^\circ\text{C}$ in an oxygen-containing atmosphere.^{44,45} The vapor pressure of PtO_2 in equilibrium with Pt (at 1-atm O_2) is much greater than that of the Pt metal (Fig. 6-18). PtO_2 exists only as a

gas; upon contact with a solid or liquid surface, it redeposits Pt metal⁴⁵:

$$\log [P_{\text{PtO}_2}] = - \left(\frac{8585}{T} \right) + 0.204,$$

$$\Delta G = 39\,270 - 0.93T, \quad (6)$$

where P_{PtO_2} is the vapor pressure of PtO_2 (Torr), ΔG is the free energy (cal/mol) and T is the temperature (K).

Thus, in oxidizing systems, it is possible to transport Pt to the molten glass as PtO_2 and then redeposit it as Pt. This produces the lace-like deposits of Pt metal often seen on the surface of unstirred glass melts prepared in platinum crucibles in air. This transport mechanism is further confirmed by experiments in which the glass is melted in a ceramic crucible that is physically separated (but in close proximity) to Pt metal.³⁹ When the melt is carried out in O_2 , Pt particles are found in the glass. On the other hand, when an inert gas (e.g., N_2) is used, no Pt particles are observed. It would appear that preparing glass in an inert environment eliminates Pt particle formation, but this is not the case. As shown by Ginther,³⁹ oxygen from the decomposition of oxide-glass components may be sufficient to produce Pt transport and the subsequent formation of inclusions.

In a related study of ED-2 silicate laser glass, Hornyak and Abenroth⁴² report a correlation of Pt solubility with oxidizing conditions of the melt; the higher the O_2 content of the gas, the greater the Pt solubility. They propose that Pt dissolution (by the glass) and redeposition may be the mechanism that causes inclusion formation in melts carried out in inert or reducing conditions.

Because phosphate glasses have a Pt solubility limit that is at least 1000 times that of silicate laser glasses, it was generally thought that the Pt would remain in solution in the Nova phosphate glass. However, Pt inclusion formation remains a problem in the manufacture of phosphate glasses.

Effects of Key Process Variables on Pt Inclusion Formation. Four key variables that affect the formation of Pt inclusions in phosphate laser glasses are (1) the O_2

Fig. 6-19. Relation between oxygen partial pressure in processing gases and measured density (per cm^3) of Pt inclusions for 0.5-L melts at $\sim 1100^\circ\text{C}$ with and without a proprietary additive developed by Hoya.

content of the process gas, (2) the melt temperature, (3) the processing time, and (4) Pt solubility in the glass. Both Hoya and Schott examined the effects of these variables on the density of platinum inclusions in their glasses. Some of the results from these experiments are reported here.

Increasing the oxygen content of the processing gas greatly reduces inclusion content of the phosphate laser glass. This is illustrated in Fig. 6-19, where Pt inclusion count is plotted as a function of O_2 partial pressure in the process gas. These results are from small, 0.5-L test melts of LHG-8 carried out at Hoya; a similar trend was observed at Schott in their LG-750 glass. Hoya also found that a proprietary additive further reduced the inclusion count under the most oxidizing conditions. Based on the experiments at both Schott and Hoya, it is now clear that oxidizing conditions promote the dissolution of Pt particles in the melt;

This can be seen by the increase in Pt^{n+} content of the glass with increased P_{O_2} (Fig. 6-19). Using a simple thermodynamic argument, Eq. (7) can be used to relate the effects of oxygen pressure to ionic Pt content in the glass. The equilibrium constant (K_{eq}) for the reaction is given by

$$K_{\text{eq}} = \frac{[\text{Pt}^{n+}][\text{O}^{-2}]^{n/2}}{(a_{\text{Pt}})(P_{\text{O}_2})^{n/4}}, \quad (8)$$

Fig. 6-20. Relation between oxygen partial pressure in the processing gas and Pt ion concentration $[Pt^{n+}]$ as measured by the 400-nm absorption, where $[Pt^{n+}] = 637 \cdot k_{400 \text{ nm}}$.

Fig. 6-21. Inclusion density as a function of melt temperature.

The activity of the Pt metal (α_{Pt}) is unity for this system. We can further assume that the O^{2-} activity remains approximately constant (i.e., the glass composition remains nearly fixed). This allows us to reduce Eq. (8) to the simple relation

$$\log [K_{eq}] = \log [Pt^{n+}] - \frac{n}{4} \log [P_{O_2}] + C_1, \quad (9)$$

or by rearranging, to the more useful expression

$$\log [Pt^{n+}] = \frac{n}{4} \log (P_{O_2}) + C_2, \quad (10)$$

where C_1 and C_2 are constants.

Note that the results in Fig. 6-19 seem to show some nonlinear behavior between $\log [Pt^{n+}]$ and $\log [P_{O_2}]$, which may indicate the system has not reached full equilibrium. This could be particularly true at low oxygen pressures because of the release of O_2 by the melt itself. Using only the data for oxygen partial pressures between 0.03 and 1.0 atm, the measured slope from the data in Fig. 6-20 is 2 ± 0.2 . Therefore, these results seem to indicate Pt dissolution as the $+2$ ion.

There are several limitations that must be considered when using Pt dissolution as a means of eliminating Pt inclusions. First, of course, is the solubility limit of Pt in the glass. Fortunately, for phosphate glasses this is quite high (≥ 1000 ppm) under normal processing conditions. Thus, reprecipitation of Pt during cooling has generally not been a problem. It is important to note that most of the ionic platinum that enters the glass does *not* come from the dissolution of the Pt inclusions, but rather from the simultaneous dissolution of the platinum crucible.

A greater concern than the solubility limit of the glass is the level of absorption at 400 nm from the ionic Pt. We found that if the absorbance, which reduces the pumping efficiency of the xenon flash-lamps, remains below 0.2 cm^{-1} at 400 nm (i.e., ~ 130 ppm Pt), the ionic Pt will not adversely affect laser performance.

Schott found that processing temperature also strongly affects the inclusion density in their LG-750 glass. Figure 6-21 is a plot of inclusion density vs processing temperature; increasing the processing temperature from 1100 to 1450°C produces a drop of $\sim 10^3$ to 10^4 in inclusion density. These results are for a series of 0.5-L melts under oxidizing process-gas conditions. For this same series of melts, Schott also measured the 400-nm absorption (Fig. 6-22), and found that it remains below the 0.2 cm^{-1} specification.

The effect of processing time on reductions of inclusion density or 400-nm absorption is much less dramatic than that of temperature. This is not surprising. The rate of dissolution is governed by a typical Arrhenius relation, which depends linearly on t and exponentially on $-1/T$. Therefore, changes in process temperature would be expected to produce much larger effects than corresponding changes in processing time.

Using results from these and other experiments, both Hoya and Schott have been able to modify the melting process to eliminate most, if not all, Pt inclusions. In brief, both companies have sought to control the platinum problem by (1) eliminating outside sources of platinum contamination and (2) adjusting process conditions (e.g., P_{O_2} , temperature, and time) to put or keep all platinum in solution (i.e., ionic). Considerable progress has been made to date. In early December 1985, we tested seven large pieces of glass with a total volume of ~ 17 to 18 L (equivalent to about three large laser disks), and found an average of <0.2 inclusions per liter of glass. Four pieces of glass were completely inclusion-free. By comparison, the current Nova disks have Pt inclusion concentrations between 10 and 1000/L; thus, the new glass represents a 50- to 5000-fold improvement. These pieces of new laser glass were prepared under conditions designed to convert any metallic platinum inclusions to the ionic form. The samples were damage tested both on Nova and with our new rapid-scan inspection tool (see the preceding article) at 2 to 4 times the threshold-damage fluence.

Further changes in the melting process are being implemented by Hoya and Schott that should reduce the Pt inclusion level to less than 0.1/L, our design specification.

Fig. 6-22. Absorption at 400 nm due to dissolved Pt vs processing temperature.

Plans to Replace Nova Laser Glass

By the end of 1985, the glass companies were able to significantly reduce the number of platinum particles in optical glass. In particular, Hoya made four 40-L melts with as little as 1/4 particle of platinum per liter. Since the 31.5- and 46.0-cm disks each contain about 7 L of glass, this is equivalent to an average of two particles in large disks. Hoya and Schott have begun the construction of new melters that embody the results of their experiments in platinum-particle reduction. They are also incorporating other changes that may further reduce the number of platinum particles—the goal, of course, is to have none.

Early in 1986, we plan to initiate pilot runs of about 20 large disks from each vendor's new melting facilities. Simultaneously, we will issue a request for quotation on production quantities. The baseline plan is to obtain replacements from the company that originally supplied the disks. However, estimates on alternative quantities will be used to negotiate replacement scenarios in the best interest of LLNL that consider cost, technical ability, and schedule.

We have been concurrently working on a new approach to edge cladding that would permit the glass manufacturers to deliver the disks unclad. Cladding could be applied at LLNL by cementing or potting preshaped absorbing strips to the glass edge. Preliminary estimates indicate that the price to us from the manufacturers of unclad disks would be about 60% of that of the monolithically clad disks. The price difference includes both added

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quotations and the lower yields of the poured cladding. Quotations for both clad and unclad disks will be obtained from the glass manufacturers and a decision will be made by early summer of 1986.

The tentative replacement schedule in Fig. 6-23 gives baseline replacement quantities including spare disks. Not shown in the schedule are forty 20.8-cm disks (30 disks plus 10 spares) used in the highest fluence area of the 20.8-cm beam, which will also be replaced.

Polymer Edge Claddings

Monolithic glass edge claddings that absorb at $1\text{ }\mu\text{m}$ are currently used to suppress parasitic oscillations caused by amplified spontaneous emission (ASE) in the Nova laser disks. Unless ASE is absorbed at the disk edge, parasitic lasing in the disk can occur, seriously reducing the

gain. Although they function very well, the monolithic glass edge claddings are expensive to produce and they reduce the manufacturer's yield of laser disks.

As a part of our previous single-segment amplifier project, we sought to develop a more flexible and economic cladding design than the monolithic glass. Consequently, we developed a composite system that uses a silicone polymer to directly bond the cladding absorber glass to the glass disk. The polymer was index matched to the disk and absorber glass by varying the diphenylsiloxane content of this silicone system with excellent results. An amplifier comprising four of the test disks has been shot over 4000 times during gain measurements with no visible bulk damage to the siloxane or degradation of amplifier performance (see "Single-Segment Amplifier: Experiments," Sec. 7).

The success of these experiments has led us to consider utilizing polymer-based edge claddings on the platinum-free phosphate glass laser disks being purchased for Nova. Unfortunately, the silicone-rubber-based cladding developed thus far is not directly applicable to Nova use. The 1.52-index siloxane is easily abraded and

not amenable to the stringent cleaning procedures used for Nova hardware. Thus, we have begun an effort to identify alternate polymers and cladding configurations applicable for Nova use.

We considered three general cladding configurations that employ a polymer component; these are shown schematically in Fig. 6-24. The first two configurations [Fig. 6-24(a) and (b)] are composite polymer-glass systems. In both designs, a copper-doped glass serves as the absorber, and the polymer is simply a bonding component. Note that the polymer is cast around the absorbing glass and thus forms a relatively thick bond in the design in Fig. 6-24(a). In contrast, a very thin (~ 20 - to $50\text{-}\mu\text{m}$) bond holds the absorber glass to the laser disk in Fig. 6-24(b); this bond is formed by applying the polymer to the absorber and laser-glass faces and then pressing them together until the polymer cures. The thin-bond approach has the advantage of minimizing the amount of polymer; thus, the probability for contamination by damaging inclusions is less because the volume of material is small.

The third cladding configuration [shown in Fig. 6-24(c)] is the direct polymer analog of our current monolithic glass cladding. In this design, the polymer is the absorber matrix. The absorber can be some dye homogeneously distributed in the polymer or a heterogeneous mixture of absorbers (e.g., opaque particles) suspended in a polymer binder. The advantage of the monolithic polymer cladding is that it is a single component system and, conceptually, could be easier and less costly to apply. The disadvantage is that the volume of polymer is large and, thus, chances for failure due to inclusion contamination could be higher. Note also that in this third configuration, one must find an absorber that is compatible with the polymer matrix.

The optical specifications of the polymer for the designs in Fig. 6-24(a) and (b) are that it be clear and nonscattering at $1\text{ }\mu\text{m}$ and that its index of refraction matches both the laser glass and the glass absorber. To prevent parasitics, the refractive index in all cases must be matched within approximately 0.01.⁴⁶ The most difficult optical requirement for the polymer to meet is that it survive without damage the

Fig. 6-24. Three potential configurations for polymer-based edge cladding: (a) cast cladding with glass absorber, (b) thin bond-line approach with glass absorber, and (c) monolithic dyed-polymer absorber.

combination of ASE and the flashlamp light, which is incident on it at a fluence greater than 10 J/cm^2 in a 0.5-ms pulse. As discussed in the next two articles, we believe that optical damage to polymers arises from the thermal decomposition that results from the heating of small absorbing particles by the incident light.

The required mechanical specifications of the polymer are less clear. We prefer a material that cures at room temperature so we are not required to heat (and potentially stress) the large laser disks. In addition, we require that the polymer provide a sufficiently strong bond that the absorber does not separate from the disk under temperature-induced stresses or during normal assembly and cleaning. In general, the mechanical properties of the polymer must be compatible with the configuration and assembly techniques that we have chosen.

Since the beginning of this effort in mid-October 1985, we have begun several evaluations of commercially available materials. The more promising among these include a flexible polysiloxane (Petrarch PSW 2239), a urethane-acrylate optical adhesive (Norland 65), a low-modulus epoxy (Mereco 4501), several high-modulus epoxies (Epotek 301, Kodak HE104-2, Hysol EA9470, and others), an ethylene-acrylic acid (EAA) copolymer (Dow 483), and a copper-doped EAA ionomer. All of the polymers, except the last, are directed at

configurations (a) and (b) in Fig. 6-24. Absorption in the case of configuration (c) is produced by doping the ionomer with copper ions.

Based on mechanical, optical, and processing tests, we hope to narrow down our potential choices of polymers and configurations early in 1986.

Authors: H. T. Powell, J. H. Campbell, and M. O. Riley

Temperature Rise of Absorbing Inclusions in Polymers

We analyzed the transient temperature change of an inclusion illuminated by light energy from all sides and surrounded by polymer bonding agent. This analysis is an extension of previous work⁴⁷ on platinum inclusions in laser glass where the inclusion was heated from one side. We are interested in the peak transient temperature rise that occurs at the inclusion-polymer interface. Other authors, such as Lundberg and Nelson,⁴⁸ have predicted the bulk temperature rise of inclusions. However, their results are only partially applicable to our study of thermal decomposition in edge-cladding polymers because the important value to us is the peak temperature rise.

Consider a spherical inclusion of radius a in intimate thermal contact with polymer material. Carslaw and Jaeger⁴⁹ give the solution for the steady-state temperature rise ΔT in a hollow sphere of inner and outer radii a and b that is illuminated uniformly at its inner surface as

$$\Delta T_{\text{steady}} = \frac{Pa(b-a)}{\kappa b}, \quad (11)$$

where

P = surface flux (W/m^2),

κ = thermal conductivity of the polymer ($\text{W}/\text{m}^2 \cdot \text{K}$).

If we let the outer radius approach infinity, the solution becomes

$$\Delta T_{\text{steady}} = \frac{Pa}{\kappa}, \quad (12)$$

which is the solution for a solid sphere of radius a illuminated at its surface. An identical solution was worked out by Powell and Kulander.⁵⁰ It can be argued that this solution is conservative in predicting the peak temperature because the light that illuminates the inclusion comes mainly from the amplifier flashlamps on one side of the inclusion. However, this light refracts and reflects as it passes into and through the laser disks, so the actual illumination could be closer to uniform than one-sided. Regardless, we expect that an analysis using one-sided illumination would produce only a slightly lower peak temperature rise than one using uniform illumination.

We also analyzed this problem numerically using the TOPAZ finite-element heat conduction code.⁵¹ We selected a carbon inclusion as a reference case and assumed a constant flux of $10 \text{ kW}/\text{cm}^2$ in the analytical solution. In our numerical solutions, we used a maximum flux of $10 \text{ kW}/\text{cm}^2$ with the half sine waveform extending over a time span of 0.5 ms. This results in an integrated absorbed energy of $3.2 \text{ J}/\text{cm}^2$ with the transient solution.

We used the property information shown in Table 6-8. The uncertainty in the properties, particularly the thermal conductivity, is large. Further, we do not know the exact chemical composition of the inclusions,

Table 6-8. Thermal properties of materials used in polymer inclusion heating analysis.

	Thermal conductivity κ ($\text{W}/\text{m} \cdot \text{K}$)	Specific heat c ($\text{J}/\text{kg} \cdot \text{K}$)	Density ρ (kg/m^3)	Surface absorptance (at $1 \mu\text{m}$)	Thermal diffusivity $\alpha = \kappa/\rho c$ (m^2/s)	$\alpha\tau$ (m^2)
Polymer ^a	0.19	1250	1030	~ 0	1.5×10^{-7}	8.6×10^{-6}
Carbon ^b	1.0	500	2250	~ 1	8.9×10^{-7}	21.0×10^{-6}
Dust ^c	0.5	1180	1640	0.07	2.6×10^{-7}	11.0×10^{-6}
Iron oxide ^d	7.0	800	5000	0.7	18.0×10^{-7}	29.0×10^{-6}

^a GE RTV 655; see Ref. 52.

^b Graphite; see Ref. 52.

^c 67 wt% SiO_2 , 33 wt% Al_2O_3 ; see Ref. 53.

^d Fe_2O_3 or Fe_3O_4 ; see Ref. 53.

which adds to the uncertainty (e.g., we assume the nominal "dust" particle is an aluminosilicate). In Table 6-8, the values of absorptance are for 1- μm light, which is a longer wavelength than the visible light primarily emitted from the flashlamps in the Nova amplifiers. However, the absorptance is about constant for the inclusion materials we considered between 1 and 3 μm , and we assumed that the absorptance does not change appreciably below 1 μm . Results are only as good as the property information included in the analysis. Until we have actual measurements to reduce the uncertainty, it is reasonable to apply the results to arbitrarily shaped inclusions as well as spherically shaped inclusions.

Results. If the radius of the inclusion is small, the analytical solution and the numerical solution agree because the time for the inclusion to reach steady state is short, and the temperature within the inclusion is essentially constant. For small inclusions, conduction of heat away from the inclusion surface depends on the properties of the polymer only, and the peak temperature rise is independent of the inclusion properties other than the absorptance.

Conversely, if the inclusion radius is large, only the numerical solution is accurate because the time for the inclusion to reach steady state exceeds the time for energy deposition. There is a temperature gradient within the inclusion and the peak temperature depends mainly on the thermal conductivity of the material.

Results for carbon inclusions are shown in Fig. 6-25, where peak temperature rise is plotted against the inclusion diameter. Carbon thermal conductivity is 5 times that of the polymer, so the results are independent of the polymer properties for large inclusion diameters. Also with large diameter inclusions, the temperature changes during times of interest occur in the outer portion of the inclusion, so results are independent of diameter. This independence can be seen at the right side of Fig. 6-25 where the curve is flat. At the left side, the curve agrees with the steady-state solution of Eq. (8) as expected. The curve should have a maximum at a diameter approximately equal to the square root of the product of the inclusion thermal diffusivity α and the flux heating time τ . For a carbon inclusion, the maximum in the curve occurs at an inclusion diameter d of

$$d_{\Delta T_{\max}} \simeq (\alpha\tau)^{1/2} = [8.9 \times 10^{-7} \text{ m}^2/\text{s} \times 0.5 \times 10^{-3} \text{ s}]^{1/2} = 21 \mu\text{m}. \quad (13)$$

The peak temperature rise scales linearly with the absorbed flux.

Results for the aluminosilicate dust and iron oxide inclusions are similar to those for carbon inclusions at both small and large diameters with one important difference. The absorptance of a carbon inclusion is ~ 14 times greater than that of dust (see Table 6-8). Hence, the maximum temperature for this particular dust inclusion is lower by a factor of 14 than for a carbon inclusion. With iron oxide inclusions, the

Fig. 6-25. Variation with peak temperature rise of a carbon inclusion vs particle absorber diameter for peak absorbed power of 10 kW/cm² in a 0.5-ms pulse.

predicted temperature rises are ~ 0.7 times those in Fig. 6-25.

The maximum temperature rise for carbon as shown in Fig. 6-25 is greater than 1500 K for the worst-size particles. The location of the maximum will be different for materials with different thermal diffusivities. Table 6-8 also shows the inclusion diameter corresponding to this maximum temperature as calculated using Eq. (13). The predicted peak temperatures for carbon are high compared to conventional polymer decomposition thresholds. However, the allowable peak temperature rise (~ 200 K) is known only for steady-state conditions. Because of the transient nature of the temperature rise, a higher instantaneous peak temperature may be allowable. Nonetheless, it appears there is a high premium on carefully filtering our polymers to remove inclusions, thus reducing potential peak temperatures and the consequent optical damage from these inclusions.

Author: J. H. Pitts

Damage Testing of Polymer Candidates

Bulk Temperatures in the Edge Cladding. Polymeric solids possess many properties that make them ideally suited for edge-cladding use, including low moduli of elasticity, tailorable refractive

Fig. 6-26. Instantaneous temperature distribution in the edge cladding for various absorption coefficients assuming an input of 10 J/cm^2 .

indices, and good optical clarity. In addition, the low-temperatures of two-part reactive liquid systems make it possible to remove particulate contamination by ultra-filtration. The relatively low degradation temperature (~ 300 to 400°C)⁵⁴ of optically clear polymeric solids, however, places an upper limit of typically 60 to 200°C on their continuous service temperature,⁵⁵ depending on the chemical type of the polymer. Consequently, it is important to characterize the thermal environment at the disk edge where ASE is converted to heat in the absorber glass.

The instantaneous one-dimensional temperature distribution in the absorber glass following a flashlamp pulse can be simply calculated using an adiabatic analysis. This is a good approximation since the time for thermal diffusion through the glass absorber plate of thickness D and thermal diffusivity α , following a step change in surface temperature, is⁵⁶

$$t = \frac{D^2}{12\alpha} = 1 \text{ s}, \quad (14)$$

which is orders of magnitude longer than the flashlamp pulse width of $\tau = 0.5$ ms. In other words, essentially no ASE-generated heat is conducted out of the absorber glass during the flashlamp pulse.

According to Lambert's law, the incident energy flux $I(x)$, in a medium of absorptivity k , gives a volumetric heat generation rate⁵⁷

$$\dot{q}(x) = \frac{dI(x)}{dx} = -Ik \exp(-kx) \quad (15)$$

from which the adiabatic temperature rise profile in the absorber immediately following a pulse is

$$\Delta T(t, x) = \frac{\dot{q}(x)\tau}{C_p} = \frac{I_0 k \tau}{C_p} \exp(-kx), \quad (16)$$

where C_p is the volume heat capacity of the absorber. Using $I_0 \tau = 10 \text{ J/cm}^2$ at the amplifier disk edge, and with $C_p = 2 \text{ J/cm}^3 \cdot \text{K}$ for the absorber glass,⁵⁸ one finds

$$\Delta T(x) = 50 k \exp(-kx), \quad (17)$$

with k in mm^{-1} and ΔT in K. Figure 6-26 shows the instantaneous temperature

distributions calculated from Eq. (17) for three different Schott absorber glasses and the Nova specification: KG-3, $k = 3.39 \text{ mm}^{-1}$; KG-1, $k = 2.02 \text{ mm}^{-1}$; KG-2, $k = 1.05 \text{ mm}^{-1}$; and the Nova specification, $k = 0.75 \text{ mm}^{-1}$.

It is clear from Fig. 6-26 that the maximum temperature rise occurs at the polymer/absorber interface. For the KG-3 and KG-1 absorbers, this temperature exceeds the maximum recommended continuous service temperature of many polymeric materials. The Nova specification falls well below the typical maximum temperature range, thus allowing greater design flexibility in choosing a polymer cladding material.

Laser and Flashlamp Damage Testing.

Beyond the temperature rise of the absorber discussed above, the viability of polymeric solids as Nova laser disk edge-cladding components depends primarily on the ability of these materials to survive direct beam and flashlamp fluence in the amplifier housing. The screening of polymers with a wide variety of chemical compositions and physical properties for their susceptibility to laser- and flashlamp-induced damage is discussed below.

Candidate polymeric materials were laminated between two 12- × 12-cm window-glass plates in a press, applying heat and pressure as required to obtain a sandwich with a thickness of 25 to 100 μm of cured polymer approximately in the center. The glass plates were each 3 mm thick with a transmission of 88%. Care was taken to obtain samples that were as clean and dust-free as possible within the limitations of the standard laboratory environment.

Table 6-9 summarizes the types and properties of the polymers tested in this sample configuration.

A Raytheon laser welder operating at 1.06 μm , characterized by beam imaging on photographic film and calorimeter measurements, produced 32 J/cm² with a full-width-to-10%-of-maximum pulse duration of 1.1 ms (average power, 30 kW/cm²). The intensity was judged to be uniform over the beam diameter of 0.5 cm at the sample surface. A 5- × 5-cm grid was drawn on the glass surface of the laminated sample; 20 of the 25 sites were initially tested with a single shot. One site on each sample that showed no visible damage ($>100 \mu\text{m}$) after the first shot was subjected to about 100 repetitive shots (reps) at 10 Hz to determine if damage could be initiated and propagated under these conditions. Typically, the sample would become hot to the touch after 100 reps, probably due to absorption by the glass plates.

A cerium-doped quartz flashlamp of 2.0-cm o.d. and 1.5-cm i.d. was operated at an explosion fraction of 0.2 using a full-width-to-10%-of-maximum pulse duration of 0.95 ms for the flashlamp tests. Average output energy at the surface of the cylindrical flashlamp was approximately 20 J/cm². The sample laminates were masked with an aluminum plate except for a circle 3.2 cm in diameter, and were placed as close as possible to the surface of the lamp. A white scatter plate was placed on the opposite side of the exposed area. Because no damage was induced in any of the polymers with this configuration (except 1.53 EC, the experimental composition silicone discussed in

Table 6-9. Types and properties of candidate edge-cladding polymers with summary of damage test results.

Sample ^a (chemical type)	Refractive index	$-\log_{10}$ (A/A_0)	Damage results	
			Laser welder	Flashlamp
Kodak optical cement (epoxy)	1.535	6.1	No damage	No damage
Hysol Glasshesive (epoxy)	1.53	5.7	No damage	No damage
Chemco resin (polyester)	1.545	5.6	No damage	Small discoloration; 1st shot
Petrarch PSW 2239 (silicone)	1.529	4.6	No damage	3 small sites
Epo-Tek 301 (epoxy)	1.544	4.4	1 site, single shot; 0 on reps	3-4 sites 1st shot; to 50 μm , shots 2, 3
1.53 EC (experimental composition silicone)	1.53	4.1	1 site singles; recoverable thermal separation on reps	3 sites enlarge to 1-2 cm on 2nd shot; different area chars completely after 1 shot
Saflex PVD (plasticized polyvinyl-butylal)	1.49	3.7	1-2 sites; no damage 100 reps	Numerous sites on 1st shot (size, several 100s of μm)
Mareco 4501 (epoxy)	1.52	3.7 (bubbles)	1 site; no damage on reps	Single sites enlarge to 2 cm in 4 shots

^a Ranked by suitability with the best material listed first.

^b At Sodium-D line.

Fig. 6-27. Flashlamp test configuration for polymer edge-cladding samples.

the introduction to "Polymer Edge Claddings", an aluminum foil reflector was placed over the lamp to direct energy at the exposed sample area (Fig. 6-27).

Observable ($>100\ \mu\text{m}$) damage usually occurred during three consecutive shots on a single surface using this procedure. We estimate that this configuration produced an energy density of $30\ \text{J}/\text{cm}^2$ at the sample surface. This number is comparable to the intensity of the laser welder, in agreement with the observation that the number density of damage sites initiated in a single shot was of the same order (i.e., 2 to 3 sites per $10\ \text{cm}^2$ of irradiated surface). Given the 1-ms pulse width of the flashlamp, its power density should also be similar to that of the laser welder.

Contaminant particle size and number density were estimated prior to testing using a Zeiss light microscope. From these data, a surface area fraction (total cross-sectional area of inclusions, A divided by area of sample A_0) of particulate inclusions was calculated for each sample, assuming a spherical inclusion of equivalent diameter. Particles less than $\sim 5\ \mu\text{m}$ in diameter were unobservable on the microscope, and were ignored in area fraction calculations. The majority of particles were in the range 10- to $100\text{-}\mu\text{m}$ -diam, with $20\ \mu\text{m}$ the most common size.

Damage was ascertained visually after each experiment on the laser welder or flashlamp, and examined subsequently with the microscope. Consequently, damage sites less than $\sim 100\ \mu\text{m}$ in

diameter may have been neglected in the visual screening.

Table 6-9 summarizes the damage to the eight polymer samples tested. Table 6-9 also lists the area fraction of particulate contamination expressed as $-\log_{10}(A/A_0)$. The qualitative ranking of candidate materials with regard to resistance to laser and flashlamp damage seems to correlate better with particulate concentration than with any other property. For example, epoxies are ranked both high and low, and flexible materials are both high (PSW 2239) and low (Mareco 4501).

The only pre-damaged sample to fail completely in a single flashlamp pulse (with charring and delamination of the entire exposed area) was one section of the 1.53 EC silicone polymer. Another rubbery polymer, Mareco 4501, also showed poor resistance to damage propagation in flashlamp tests, yielding 1- to 2-cm-diam delaminated areas with charred centers after four shots.

Assuming that delamination is a precursor to catastrophic failure and that it occurs at the polymer/glass interface due to pressure generated by local heating of dissolved gases in the polymer or actual volatilization of the sample at inclusions, a fracture-mechanics approach gives the critical pressure P for propagation of a circular debond of radius r at the interface:

$$P = \sqrt{\frac{2\pi E \gamma_b}{3r}}, \quad (18)$$

where E is the tensile modulus of the polymer and γ_b is the adhesive fracture energy (bond strength). This equation suggests that materials with high modulus and good adhesion should be better at controlling damage propagation that results from locally generated interfacial pressures at hot inclusions. This may account in part for the high damage propagation resistance of the Hysol and Kodak epoxies and the poor resistance of the rubbery Mareco 4501 and 1.53 EC, the latter being judged to have the poorest adhesion.

At roughly equivalent energy flux ($30\ \text{J}/\text{cm}^2$), the broad spectrum

cerium-doped flashlamp appears to be a more severe environment than the 1- μm laser light, producing gross (catastrophic) failure in a single pulse in certain pre-damaged materials. In contrast, the 1- μm laser light would only enlarge existing damage sites in a single pulse mode, although reps at 10 Hz could cause explosive failure due to the local accumulation of thermal energy at damaged inclusions.

The surface characteristics of dirt particles appear to play a primary role in damage initiation and propagation, more so in fact than particle size. Dark particles observed under a microscope would preferentially enlarge over bright reflecting particles when shot at 32 J/cm² with the laser. In addition, damage sites could initiate from unobservable particles (<5 μm) while larger sites (10 to 100 μm) would remain unchanged. These observations are consistent with our proposed thermal damage mechanisms and the calculations discussed in the preceding article when one takes into account the effective absorption of the particle.

Conclusions. Eight candidate polymers for Nova laser amplifier-disk edge cladding were tested for laser- and flashlamp-induced damage at fluence levels approximately 3 times those expected in the amplifier housing. The chemical type of the commercial polymers ranged from two-part reactive silicone, epoxy, and polyester compounds to thermoplastic polyvinyl-butyl sheet. Resistance to damage from both sources correlated with the microscopic (particle diameters >5 μm) cleanliness of the samples, implying a thermal damage mechanism that originates at particle sites and appears to be roughly independent of the chemical composition of the polymer. Damage propagation, however, may depend on the physical characteristics of the material.

Final choices of materials for edge-cladding use will be guided by physical requirements, including processability, refractive index, spectral characteristics, adhesive strength, and mechanical properties. The relation between polymer modulus, adhesive strength, and damage propagation resistance will be studied

on filtered samples prepared in a clean room.

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Prototype Cladding Design

Introduction. We are using a 46-cm amplifier to test a prototype polymer-based edge cladding with glass absorber plates cast or glued in place around the periphery of the laser disk. Monolithic glass edge claddings, such as those presently used in the Nova amplifiers, contribute significantly to disk cost and delivery time. We expect to decide in mid-1986 between the polymer-based option or the conventional cladding for Nova replacement disks.

Our current thinking is that the assembled edge cladding will be applied to polygonal laser disks as shown in Fig. 6-24(a) and (b). We plan to test for possible damage to the polymer edge cladding and measure the amplifier gain and the birefringence in the laser glass.

A spare Nova 46-cm amplifier will be used for prototype tests. The existing 46-cm amplifier uses elliptical disks with monolithic glass edge cladding in their holders, a pump cavity with top and bottom reflecting surfaces, and vertical flashlamps backed with shaped reflectors.⁵⁹ There are forty 46-cm amplifiers on Nova and two disk assemblies per amplifier. Also, there are fifty 31.5-cm amplifiers with two disk assemblies per amplifier. To use the polygonal shape shown in Fig. 6-28, the only modifications necessary on either amplifier are to the disk holder and the top and bottom reflectors.

Edge Cladding Assembly. The primary function of the edge cladding is as an ASE absorber. Secondly, it can act as a soft mount for the laser glass to minimize birefringence. Soft mounts have been used routinely on Nova optics, e.g., the large turning mirrors that operate at 1/10-wave distortion. However, in the mirror application, there is no flashlamp fluence present. The mirror mounts have significantly less birefringence than do the present split

Fig. 6-28. Subassembly of a polygonal disk assembled in a bezel with a polymer-based edge cladding.

laser disks in the 46-cm amplifier, which are only marginally acceptable.

We will first investigate the "cast" approach, which was previously used in the single-segment amplifier work,⁶⁰ to put absorber on the outside of the disk. Absorbing glass is placed at the peripheral edge of each piece of laser glass using a polymer whose index of refraction matches that of the laser glass. These pieces of absorbing glass are incorporated into a restraining structural frame that forms the final assembly: two sections of laser glass, absorbing glass, and the compliant mounting material (Fig. 6-28). The polygonal laser disks have straight edges, that allows us to use flat absorbing plates; the disks also have a flat bottom, that supports the disk weight without creating stress. The optical beam fits this shape closely with little wasted glass area.

Critical to this assembly is the 0.010-in. thickness tolerance of the two pieces of laser glass and the twelve pieces of absorbing glass. Ridges formed by variations in thickness would hinder cleaning and also make it difficult to control the gap between the disk and the mask. A large gap would allow light to reach the polymer

and damage it. During assembly, the polymer must be handled in a way that does not contaminate the laser glass transmissive surfaces. The investigation of these issues, and others associated with manufacturing and assembly, will depend heavily on our prototype experience. One of the dominant issues is the cleanliness of the total process. Clean, particulate-free materials are required to minimize damage to optical surfaces. We must control bubble formation and particulates in the polymer. To meet these requirements, the assembly and processing of materials will be done in a class-100 clean room. The laser glass subassembly will be washed with soap, rinsed with distilled water, and dried by centrifugal spinning. Finally, metal masks will be placed over the polymer edge-cladding material for protection from the flashlamp light. Before this assembly can be installed into an amplifier, it must be inspected and validated to the very strict cleanliness specification used for Nova amplifiers.

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Nonlinear Propagation and Frequency Conversion

Introduction

When construction of the Nova laser system was officially completed in December of 1984, basic system performance and laser diagnostics operation at the fundamental and harmonic frequencies were demonstrated in a series of impressive ten-beam milestone shots (see "Nova Laser" in Sec. 5). Our efforts during 1985 have included extensive laser system experiments in the high-power, high-energy-pulse regime to characterize the system performance, and to establish an important data base for future solid state ICF laser systems. In early tests, energy output at the harmonic wavelengths was somewhat less than expected, and this alerted us to the possibility that beam propagation problems at high power might be occurring in the laser chain and/or the air path to the frequency conversion crystals. Full-aperture experiments with 74-cm crystal arrays located either at the target chamber or (for test purposes) mounted in the switchyard showed an unexplained degradation in conversion performance. We observed that third-harmonic energy conversion reached a maximum of 50% at intensities near 1.75 GW/cm^2 and sharply decreased for intensities above 2 GW/cm^2 . Second-harmonic conversion would typically reach 50% efficiency at 1.5 GW/cm^2 and saturate at $\sim 55\%$ at higher intensities. Code calculations in both cases gave energy conversion efficiencies in excess of 70% for temporally square pulses with no conversion efficiency decrease for intensities below 3.5 GW/cm^2 . The Novette frequency conversion experiments with prototype Nova arrays gave qualitatively the same results as the early Nova experiments.

To explain these discrepancies, we proposed that certain nonlinear optical effects such as stimulated light scattering or intensity-dependent polarization rotation might be degrading the fundamental beam quality and the polarization, thereby adversely affecting frequency conversion.

A collaborative effort with the Nova engineering staff and with members of the Advanced Laser Development Group was therefore undertaken to investigate these issues. We report here our early observations of these effects. We are continuing efforts to obtain an understanding of the physical phenomena and the engineering implications for future Nova improvements and for next generation ICF laser designs.

At the high end of Nova performance, typically with square laser pulse shapes, unfocused Nova output beams can exceed average intensities of 2.5 GW/cm^2 and peak intensities can approach 4.0 GW/cm^2 . Laser optics in the 46-cm amplifier chain can experience intensities exceeding 6 GW/cm^2 and energy densities above 15 J/cm^2 for longer square pulses. Much of the historic work in nonlinear optics was performed with Q-switched, ruby-rod laser systems at these intensities with tightly focused beams. On Nova, such intensities can occur at the full system aperture and therefore a long, coherent interaction for stimulated light scattering or other nonlinear processes are possible either within the laser system or in the 45- to 65-m air path between the output aperture and the frequency conversion crystals.

We have encountered the following nonlinear optical processes and related phenomena⁶¹:

1. Beam-induced damage to optical components of the system due to large intensity modulation.
2. Nonlinear ripple growth and low-spatial-frequency beam breakup at high power due to nonlinear refractive index effects in the final amplifier stages.
3. Temporal pulse distortion due to gain saturation in the final amplifier stages.
4. An intensity- or B -integral-dependent rotation of the electric field in partially depolarized portions of the Nova beam.
5. Stimulated rotational Raman scattering from the nitrogen gas in the long air paths at the output end of the system.

The first three nonlinear effects are well known to us and have been important design issues on Nova and previous large lasers such as Cyclops, Argus, Shiva, and Novette.⁶² For example, optical damage to spatial filter lenses and amplifier disks due to intensity spikes caused by beam breakup sets the upper limit on the system

performance at the fundamental 1ω wavelength. Beam breakup is caused by diffraction around previously damaged sites in laser components, and by enhanced growth of small-scale self focusing from nonlinear refractive index effects. This nonlinear diffraction phenomenon leads to severe intensity modulation and to subsequent damage to the optical train. On Nova, an additional high-power effect occurs that is characterized by a periodic array of hot spots with dimensions of ~ 1 cm. This large-scale beam-breakup mode in multiple-stage amplifier chains cannot in general be suppressed by spatial filtering and image relaying. For square temporal pulse shapes, the nonlinear diffraction effects combine to set the upper power limit per Nova beamline at approximately 9.5 TW. The Nova energy performance curves at the fundamental and third-harmonic wavelengths for square temporal pulses are shown in Fig. 6-29. The high-end Nova 1ω performance per beam will be 9.5 kJ at 1 ns, 15 kJ at 2.5 ns,

and 18 kJ at 5 ns or longer. The lower 1ω performance curve represents what we would consider to be routine system operation. We have discovered, however, that metallic platinum inclusions in the Nova laser glass cause damage at energy fluences and beam energies significantly lower than the intrinsic bulk damage threshold, and therefore we have had to temporarily lower operating energies on Nova for routine experiments. The platinum damage issues and large aperture optics testing are discussed in "Nova Laser Glass" earlier in this section.

Item number three on our list is temporal pulse distortion. A square output pulse requires a monotonically increasing pulse in all amplifier stages in order to counteract the effects of gain saturation. Generating a square output pulse requires careful pulse shaping of the oscillator input, as discussed in "Pulse Generation and Shaping" later in this section. As a result of this pulse shaping, the B -integral for a square output pulse increases monotonically in

Fig. 6-29. Projected 1ω and 3ω energy performance per beamline. Upper dashed curve is stimulated rotational Raman threshold at 1ω . Lower dashed curves are improvements to 3ω performance that would be observed by polarizing the Nova output 1ω beam.

time, with the last photon seeing a significantly higher δB per amplifier stage than the first photon. Nonlinear optical effects that are B -integral dependent will be dynamic in time and therefore the last photon will set the performance limits when nonlinear diffraction dominates.

The currently projected third-harmonic performance to the Nova ten-beam target chamber is shown in the lower set of curves in Fig. 6-29. A nominal efficiency of 85% is factored into these curves to account for transport losses from the crystal array output to the target chamber center. In the high-power regime, that is for square pulses <2.5 ns, the third-harmonic conversion performance is degraded by the occurrence of B -integral-dependent polarization rotation effects that originate in the high-power end of the Nova chain and continue in the air path.⁶³ Second-harmonic conversion efficiency is similarly degraded and the performance curves resemble the third-harmonic curves. The Nova system is designed to be linearly polarized at the output aperture. The present Nova frequency conversion design with Type II KDP crystals is specifically optimized for a uniform linearly polarized fundamental wavelength beam. We discovered this year, however, that the Nova beams have a small orthogonally polarized electric field component, and that this depolarization is nonuniformly distributed across the aperture and appears to be present at all power levels. We estimate that 2 to 3% of the beam energy is orthogonally polarized, with the worst regions of depolarization occurring in the four corners of the beam where the depolarized energy can range from 6 to 8% of the local beam fluence. An orthogonal electric field component implies in the most general case that these regions of the beam are elliptically polarized, and that this must occur by propagation through one or more optical components in the laser chain that are birefringent. At high power, the nonlinear refractive index changes in the laser glass, focusing optics, and air path through the system conspire to cause a rotation in the orientation of the polarization ellipse. In the simplest case, this orientational change is proportional to the accumulated

B -integral over the air path and the initial degree of ellipticity. A slight amount of ellipticity (<0.2) by itself does not adversely effect Type II frequency conversion processes, but the rotation of the polarization orientation as a function of laser output intensity has a severe degrading effect on conversion efficiency. In "Polarization Changes at High Intensity" we summarize the details of this polarization reorientation effect and describe our small-aperture, time-integrated, and time-resolved measurements on beam 10 using the classical techniques of ellipsometry. In "Beam Depolarization Effects on Nova Frequency Conversion" we describe how the B -integral-dependent reorientation and beam depolarization affect the third-harmonic conversion process. We conducted third-harmonic experiments in the Nova switchyard on beam 10 at the 2.5-, 25-, and 70-cm apertures to investigate the conversion efficiency rollover phenomenon at high power, and to investigate the possible improvements to conversion efficiency when the Nova output beam is repolarized. At the 25-cm aperture, we achieved greater than 75% energy conversion with repolarization of the drive beam. At the full Nova aperture, we fielded a two-segment, 74-cm-aperture, thin-film polarizer at the output of beam 10 and achieved $>62\%$ energy conversion to the third harmonic with no apparent performance degradation due to B -integral effects. The dashed 3ω performance curves in Fig. 6-29 represent the improvements to Nova that would be expected with a high-rejection-ratio (better than 200:1), high-damage-threshold polarizer located at the output of each Nova beamline. With repolarization of Nova, we expect to achieve short-wavelength energy on target >40 kJ for 1-ns pulses and >60 kJ for 2.5-ns pulses. Temporal reshaping of the frequency converted pulse due to the dynamic nature of the elliptical rotation process is also discussed in "Beam Depolarization Effects on Nova Frequency Conversion" and is a serious issue for target irradiation experiments. We have demonstrated, however, that beam repolarization does maintain pulse fidelity in the frequency conversion step.

The fifth nonlinear effect primarily involves the air path between the system output and the location of the frequency conversion crystals at the target chamber. The two longest beam paths going into the Nova ten-beam chamber are beams 3 and 7 with air-path lengths of ~ 65 m. At the high end of Nova performance, at pulse durations from 0.5 to 2 ns, the possibility exists for a stimulated light scattering process to exceed threshold for these two beams. Beamlines 6 and 7 when directed into the two-beam target chamber both travel ~ 105 m from the system output, and even mid-energy pulses can exceed the threshold for stimulated light scattering. Stimulated rotational Raman scattering in hydrogen or deuterium gases has been known since the early days of nonlinear optics. It has only been relatively recent that SRRS has been observed to occur in nitrogen gas at atmospheric conditions.⁶⁴ The occurrence of SRRS on Nova at energy levels exceeding threshold by 10 to 15% would severely impact the system performance, because SRRS causes severe beam breakup and a substantial shift in beam energy from the fundamental laser wavelength to rotational Raman frequencies.⁶⁵ To investigate the effect of SRRS on Nova performance and to obtain a better understanding of the SRRS physics issues, we fielded experiments in the Nova switchyard and laser bay areas utilizing a subaperture of beam 10. The experiments and physics issues are discussed in "Stimulated Raman Scattering Effects on Beam Propagation." Air-path lengths from 65 to 150 m were investigated with either 1- or

2.5-ns square temporal pulses. The SRRS process is transient in the buildup of frequency shifted and scattered light, and therefore the energy threshold depends on pulse duration, as well as path length. Beam quality is also an issue here since a modulated beam will have spatial intensity spikes that will give a lower threshold for generating SRRS than would be calculated from a spatially averaged beam intensity. In Fig. 6-29 the threshold energy for SRRS in air for a 65-m air-path length is plotted over the projected 1ω Nova performance curves. We see that SRRS threshold can potentially be exceeded for 0.5- to 2-ns pulses. Spatial intensity modulation, particularly near the edges of the beam, will derate the observed threshold below the plotted curve by an amount corresponding to the degree of intensity modulation. The SRRS issues on Nova remain currently under investigation. It is clear however that the two-beam target chamber air path will require a partial evacuation and filling with an inert gas to avoid severe SRRS. The design and implementation of such a system has been completed on beamlines 6 and 7.

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Stimulated Raman Scattering Effects on Beam Propagation

Introduction. The beam tubes between Nova's output spatial filter and the target chambers are filled with air, so stimulated scattering and other nonlinear processes eventually limit the laser intensity that can

Fig. 6-30. Layout of SRRS experiments on beam 10 of Nova.

be propagated to the target chamber.

Averbakh and coworkers⁶⁶ have shown that the stimulated scattering process of lowest threshold in air at or below atmospheric pressure is stimulated rotational Raman scattering in nitrogen.

During the design phase of Nova, we estimated the stimulated Raman gain in the beam tube and decided that stimulated Raman processes would probably not impact normal operation of the laser. However, the uncertainties in Raman gain were large enough to prompt an investigation of the Raman conversion that might be seen at very high intensities. This experimental study of SRRS conversion on Nova was especially important in planning for the two-beam target chamber (see Sec. 5), which has a propagation path almost twice as long as the propagation path to the ten-beam chamber. Study results confirmed that the SRRS threshold does not affect normal operation with the 10-beam target chamber. However, we found that significantly less power would be delivered to the two-beam chamber if the path were entirely in air. Thus, the beam tubes for the two-beam chamber were designed to be evacuated or filled with an inert gas, which reduces beam degradation from SRRS to acceptable levels.

Observations of SRRS at Several Distances and Pulse Lengths. We diverted a 20-cm-diam section of beam 10 during the activation of Nova early in 1985 to study the growth of SRRS. As shown in Fig. 6-30, the test beam was propagated from the switchyard through an open-air path in the laser bay to a film spectrograph and a near-field film camera that recorded the SRRS conversion and its effect on the beam profile. A high reflector folded the beam once for propagation paths longer than 85 m. We studied both a region of uniform intensity near the center of the beam and a less uniform, somewhat depolarized section near the disk split and beam edge (Fig. 6-31).

The SRRS components observed are shown in Fig. 6-32 as a function of the average intensity $\times$ propagation distance $I_p L$ for several distances and two pulse lengths. For a 1-ns full-width-at-half-maximum (FWHM) flat-top (~ 0.1 -ns rise and fall) pulse, we begin to see first Stokes components of nitrogen SRRS at

Fig. 6-31. Location of the center and edge samples of beam 10 used in Raman experiments.

intensity-length products of 12 TW/cm at the edge of the Nova beam and 16 TW/cm in the center of the beam. The edge of the beam has more intensity modulation than the center. We estimate that the 1 ± 0.33 intensity modulation is consistent with peak $I_p L = 16$ TW/cm, which agrees with the threshold in the center of the beam. With a flat-top 2.5-ns FWHM pulse, we observed a threshold of 12 TW/cm in the clean central part of the beam; longer pulses have a lower threshold because of transient effects discussed elsewhere.⁶⁶⁻⁶⁸

Fig. 6-33. Spectrum of SRRS in air for a 1-ns pulse propagating 150 m at an $I_p L$ of 30 TW/cm.

Fig. 6-34. Intensity profile of the test beam below Raman threshold at an $I_p L$ of 14 TW/cm and a propagation distance of 150 m with a 1-ns pulse. Note diffraction fringes from the hard aperture and disk split.

Typical models predict a 30 to 40% higher SRRS threshold for 1-ns pulses than for 2.5-ns pulses, which is consistent with our data. The data also agree with the expected $\gamma I_p L$ dependence, where γ is a Raman gain coefficient. Diffraction and nonlinear index effects change the details of the beam

intensity profile over the propagation path, so it is difficult to be more quantitative about intensity-length products. The final spatial filter forms an image of the last laser amplifier at the 65-m point in the propagation path, so some diffractive effects are suppressed by image relaying.

Figure 6-33 shows a characteristic spectrum of nitrogen SRRS, well above threshold, at the 30-TW/cm point on Fig. 6-32. The spectrum shows significant first Stokes intensities on the S(6), S(8), and S(10) rotational Raman lines at the 59, 75, and 91 cm^{-1} shifts from the laser frequency. There are numerous second and third Stokes components that correspond to various intercombinations of these first Stokes components. Weak fourth Stokes and first anti-Stokes components were observed on some shots. For this spectrum, we estimate that 62% of the original pump energy remains in the pump and that 20, 15, and 3% of the energy is shifted to first, second, and third Stokes components, respectively. Although the S(6) line is the most intense, we find that the S(8) and S(10) lines are the first to go over threshold and that they do so at the same value of $I_p L$ (within the precision of this experiment).

Multiple rotational line and four-wave mixing effects far above threshold lead to a very complex spectrum, as previously noted by Averbakh.⁶⁶

The profile of the 20-cm test beam at the end of the 150-m propagation path shown in Fig. 6-34 is below the SRRS threshold; in Fig. 6-35, it is well above the SRRS threshold. The SRRS-shifted components have a speckled intensity profile with a characteristic scale of order 1 to 3 mm, and they occupy an area 2 to 3 times that of the collimated pump beam at the end of the propagation path. These photographs confirm that SRRS well above threshold will cause unacceptable degradation of Nova beam quality.

The SRRS components grow from spontaneous Raman fluorescence in this experiment. In this geometry, we have shown⁶⁷ that the small-signal gain along the propagation path that is required to reach a threshold of 1% conversion of the pump (approximately the minimum conversion detectable in the experiment) is about 33 nepers (e^{33}). Amplified spontaneous emission from nitrogen contained within the Nova disk amplifiers reduces this required gain to about 30 Np. Since the gain in nepers is just $\gamma I_p L$, we can estimate the gain coefficient γ for nontransient (2.5 ns) pulses in air as $\sim 2.5 \text{ cm/TW}$. This is $\sim 40\%$ higher than the gain Averbakh⁶⁶ calculated for these transitions, and it is consistent with recent preliminary gain measurements made at SRI International.⁶⁹

Figures 6-36 and 6-37 are near-field pictures of sections of a 76-cm beam propagating to the two-beam target chamber after installation of new beam tube number 27. The diffraction caused by the hard 20-cm aperture that is so evident in Figs. 6-34 and 6-35 does not occur; the intensity structure is also much cleaner, which may reflect the higher turbulence in the open-air path used in the earlier experiments. SRRS threshold values are nevertheless consistent with the earlier data and show the characteristic speckle of SRRS components.

These experiments show that the quality of the most intense part of the beam is seriously degraded for any $I_p L$ more than ~ 10 to 15% higher than the threshold defined in Fig. 6-30. These limitations are

Fig. 6-35. Intensity profile of the test beam in the preceding figure at an $I_p L$ of 26 TW/cm, well over Raman threshold. Note the divergent, speckled stimulated Raman components.

Fig. 6-36. Intensity profile of a section of a 1-ns, 76-cm beam propagating 105 m to the two-beam target chamber at an $I_p L$ of 13.5 TW/cm. The disk split and its diffraction pattern are visible in the center; however, no Raman components can be seen.

summarized in Table 6-10 for various pulse lengths and propagation paths to the two target chambers. The table shows that SRRS will not be a significant issue for normal operation with the ten-beam chamber, but that the power deliverable to the two-beam chamber is significantly reduced when the propagation path is wholly in air.

Fig. 6-37. Intensity profile similar to Fig. 6-36, but at an $I_p L$ of 16.5 TW/cm. Characteristic Raman speckle components are now visible.

Table 6-10. Propagation limits set by nitrogen SRRS in air in the Nova beam tubes.

Pulse length (ns)	Path	Beam power (TW)	Beam energy (kJ)	Intensity (GW/cm ²)
0.5	10-beam	10.3	5.1	2.9
1.0	10-beam	7.5	7.5	2.1
2.5	10-beam	5.8	14.4	1.6
0.5	2-beam	6.3	3.2	1.8
1.0	2-beam	4.6	4.6	1.3
2.5	2-beam	3.6	8.9	1.0

Consequently, we designed the beam tubes to the two-beam chamber so they can be evacuated or filled with an inert gas with no significant Raman gain.

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Polarization Changes at High Intensity

Introduction. In experiments to study the cause of lower than expected frequency conversion efficiency on Nova, we observed a small vertically polarized component that increases with intensity in some regions of the output beam, a beam that nominally has horizontal polarization. At high laser output intensities, the vertical component significantly reduces the efficiency of frequency conversion to the second and third harmonics. Having identified the cause for the low conversion efficiency, we can now search for methods to correct it and obtain the full inherent output of Nova at the second and third harmonic.

In the Nova laser system, the electric field vector of the beam is oriented in the horizontal plane at the exit of the final spatial filter on each amplifier chain by the thin-film polarizer sets in each of the four Faraday rotator assemblies up to the 31.5-cm amplifiers, and by polarizers in Pockels cell assemblies in smaller-diameter components at the front end of the amplifier system. In addition, all amplifier disks are set at Brewster's angle to the direction of propagation, thus ensuring a high net reflection loss for the vertical polarization. Nova's type-II harmonic crystals require a linearly polarized beam at the target chamber for optimum second- and third-harmonic conversion.

Sources of Vertical Polarization. The small vertically polarized component is distributed in a nonuniform pattern across the full aperture and exists at all power levels. Approximately 2 to 3% of the full-beam energy is vertically polarized on a typical shot (Fig. 6-38). Regions of the beam have apparently become elliptically polarized by propagation through one or more birefringent optical elements in the laser chain. Full-beam photos of the vertical polarization, obtained by reflection of the beam off a large Brewster-angle glass plate at various positions along the chain, indicate that the source of the birefringence is distributed among many components in the output end of each chain. Stress in the amplifier disks, caused by mounting and edge-cladding treatments, and vacuum loading of the spatial filter lenses are thought to be the source of the birefringence. At high power, the depolarization increases because of a rotation of the principal plane of polarization of the beam. This change in orientation of the polarization ellipse is an intensity-dependent refractive index or *B*-integral effect (defined below). It occurs in the laser glass, spatial filter lenses, and air path; and, to the first order, it is proportional to the ellipticity and the *B*-integral accumulated over the path length beyond the birefringent components.

Ellipsometry Measurements of Depolarization. To understand the details of the depolarization and the changes with intensity, we performed small-aperture time-integrated and time-resolved

ellipsometry measurements on beamline 10 in the regions of worst depolarization. Over the Nova operating range of 4 to 8 kJ for 1-ns pulses, the worst-case vertical-polarization fluence measured was 6 to 8% of the local beam fluence. Time-integrated measurements indicate a polarization ellipticity of 0.20 to 0.28, and an ellipse rotation of 3.9 to 5.1 deg for every GW/cm^2 of beam intensity for square temporal pulses. Time-resolved measurements, limited by streak camera performance, indicate that as much as 4 to 8 deg of rotation occur dynamically during a high-power shot. The time-averaged data are in excellent agreement with a theory that relates ellipse rotation to the B -integral, and some of the time resolved data exhibit behavior that is expected theoretically. In addition, we briefly investigated ellipse rotation in the air path. For an equivalent 65-m Nova air path (beamlines 3 and 7), we deduced a rotation rate of 2.2 deg per GW/cm^2 , which is in good agreement with calculations using the best available values for the nonlinear index of refraction n_2 for air.

The techniques used for our polarization measurements⁷⁰ take advantage of the well known fact that the polarization state of any monochromatic beam is fully specified by determining the four Stokes parameters from intensity measurements. Figure 6-39 shows the experimental arrangement of polarizers and calorimeters with which we

Fig. 6-38. Nova output beam in vertically polarized light. The beam was reflected from a glass plate at Brewster's angle.

Fig. 6-39. Arrangement of polarizers, calorimeters, and streak cameras used to measure the Stokes polarization parameters of a 2.5-cm-diam segment of Nova beam-10 output.

Fig. 6-40. The characteristics of the polarization ellipse.

Fig. 6-41. Measured polarization states for a 1-in.-diam depolarized portion of Nova beam 10; tilt angle ψ is shown on the polarization ellipses.

measured the energy in the four polarization directions (0, 45, 90, and 135 deg) relative to horizontal. This allowed us to determine the four Stokes parameters except for the sign of s_3 , which determines the handedness of the elliptically polarized beam. The streak camera pickoffs were used to determine the pulse shape for each of the four polarization directions. In addition to measuring the beam directly out of beamline 10, we propagated the beam through various lengths of the air

path before directing it into the measuring apparatus. A large aluminum plate blocked the beam except for a 2.5-cm-diam portion in either the right or left lower corner of the D-shaped half beams. The first time-integrated measurements used the left-hand corner. Later time-integrated and time-resolved measurements used the right-hand corner; the approximate location is shown in Fig. 6-39.

Figure 6-40 is a typical polarization ellipse. The equations used to determine the Stokes parameters are:

$$\begin{aligned} s_0 &= I_0 \text{ deg} + I_{90} \text{ deg} \\ s_1 &= I_0 \text{ deg} - I_{90} \text{ deg} \\ s_2 &= I_{45} \text{ deg} - I_{135} \text{ deg} \\ s_3 &= \pm [s_0^2 - s_1^2 - s_2^2]^{1/2}, \end{aligned} \quad (19)$$

where I is intensity. The ellipticity is defined by $\epsilon = \pm E_b/E_a = \tan \chi$, and $\eta = E_y/E_x = \tan \alpha$ (Fig. 6-40). The ellipticity angle $\chi = \tan^{-1} \epsilon = 0.5 \sin^{-1} s_3/s_0$. The tilt angle of the ellipse from the horizontal $\psi = 0.5 \tan^{-1} [\cos \delta \times \tan 2\alpha] = 0.5 \tan^{-1} s_2/s_1$. Figure 6-41 shows the effect of intensity on time-averaged polarization ellipses for three shots. As is the case for all the data collected, the ellipticity ϵ was nearly constant at values between 0.24 and 0.27. However, the tilt angle ψ increased with intensity. Both the constant ellipticity and intensity dependence of ψ are expected in observations of nonlinear-polarization ellipse rotation. Figure 6-42 shows tilt angle plotted against intensity

for three different sets of data with time-integrated ellipse rotation rates from 3.9 to 5.1 deg per GW/cm² of average beam intensity. The ellipticity ranges from 0.21 to 0.27 and the depolarization from ~5.5% at low intensity to ~8% at high intensity (>1.6 GW/cm²). The 3.9-deg data in Fig. 6-42 are the most accurate because the polarizers and the calorimeter channels were very carefully realigned. Some variation in tilt angle and ellipticity may arise because the three different sets of data were taken from different locations in the beam. Ellipticity is determined accurately because small polarizer misalignments have only a secondary effect on computed ellipticity. However, small ellipse tilts from the horizontal have greater uncertainty because measurement of the intensity difference in the 45- and 135-deg channels is subject to systematic error from polarizer misalignment and because $I_{45} \approx I_{135}$ for small tilt angles. Also, note that our apparatus is located in beamline 10 after two 45-deg turning mirrors in the switchyard that normally direct the beam through the port in the target bay wall. It has not been determined yet if these large mirrors, one of which is partially transmissive for the output sensor package, are biasing our data.

During measurement of ellipse rotation in the long air paths, we attempted to keep the full-beam laser energies between 5.6 and 7.0 kJ (i.e., 1.7 to 2.1 GW/cm²) while taking data at the three path lengths. Figure 6-43 plots the average ellipse tilt at each path length as a function of path length. The three data points lie approximately on a straight line, indicating that increasing the air path has a linear effect on ellipse rotation. A slope of 0.064 deg/m at an average intensity of 1.9 GW/cm² yields a rotation rate in air of 2.2 deg per GW/cm² for a 65-m air path. This is at the low end of our estimate for ellipse rotation in air, which is based on published values for the nonlinear refractive index of air.⁷¹

Simple ellipse rotation is ~1/3 times the product of beam ellipticity and the *B*-integral accumulated in the high-power end of the laser chain. When self-focusing from nonlinear refractive index effects can be neglected, a constant ellipticity is

Fig. 6-42. Tilt angle of the measured time-averaged polarization ellipse vs output beam intensity from three different sets of measurements.

Fig. 6-43. Measured ellipse rotation angle vs path length in air; shots were approximately equal in intensity.

maintained regardless of the amount of ellipse rotation. In other words, the right and left circular components of the elliptically polarized beam do not exchange energy. Only the linear birefringence in the optical components and differential Fresnel reflection losses cause changes in the beam ellipticity during propagation through the laser chain. Again, taking the simple view that the beam ellipticity ϵ is constant during propagation, the spatial rate of change of the ellipse tilt ψ is given by

$$\frac{d\psi}{dz} = \rho \frac{\epsilon}{1 + \epsilon^2} \frac{\pi n_2 E^2}{\lambda}, \quad (20)$$

where ρ is the depolarization of the nonlinear refractive index, n_2 is the nonlinear

refractive index, and E is the electric field amplitude. For laser glass, $\rho = 1/3$ when the electronic contribution to the nonlinear refractive index is substantially greater than the nuclear contribution. The B -integral for a narrow time-slice of the laser pulse is defined by

$$B = \int \frac{\pi n_2 E^2}{\lambda} dz, \quad (21)$$

where the spatial integral follows the propagation of the time-slice through the relevant portion of the laser chain. Experimentally, the Nova beam becomes depolarized in the 46-cm amplifier stages, and so the relevant B -integral increment for ellipse rotation is approximately from the input to the first 46-cm disk amplifier to the output of the 46- × 74-cm spatial filter. Formally integrating Eq. (21) over the path length and averaging over the pulse length, the time-averaged ellipse rotation $\bar{\psi}$ during the pulse is

$$\bar{\psi} = \rho \frac{\epsilon}{1 + \epsilon^2} \bar{B}. \quad (22)$$

From code calculations of Nova performance with 1-ns square output pulses,⁷² the time-averaged B -integral increment from the 46-cm amplifiers to the system output may be approximated by

$$\bar{B} = 0.74 I^{1.243}. \quad (23)$$

With an average ellipticity of 0.25, Eqs. (22) and (23) give a time-averaged ellipse rotation rate of 4.1 deg per GW/cm² at a shot intensity of 1 GW/cm², and a rotation rate of 5.4 deg per GW/cm² at 3 GW/cm². These rotation rates are in agreement with the range of measured time-averaged rotation rates.

These detailed measurements on small-aperture segments of the depolarized regions are consistent with the interpretation that a small degree of elliptical polarization is induced into the beam by birefringence in components in the output stages. The intensity-dependent rotation of the principal axis of the polarization ellipse is caused by the intensity-dependent, nonlinear index of refraction n_2 in the glass and

in air in the beam path. This increases the degree of elliptical polarization, which significantly reduces the conversion efficiency of the Nova beam.

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Beam Depolarization Effects on Nova Frequency Conversion

During 1985, we actively investigated the effect of nonuniform laser polarization in degrading Nova system performance. In particular, experiments were concentrated on locating the source of the birefringence in the high-power end of the chain, characterizing the dynamics of the polarization rotation with laser power, and determining the effects on third-harmonic (3ω) frequency conversion. With repolarization of the full-aperture beam (using a two-segment large-aperture thin-film polarizer at the output of beamline 10), we achieved >62% energy conversion to 3ω at ~ 2 GW/cm² with 1-ns square laser pulses. Third-harmonic, 25-cm aperture experiments with repolarization achieved >75% energy conversion, and with improved thin-film polarizers, we expect to achieve at least a 70% energy conversion to 3ω at full Nova aperture.

Frequency-conversion experiments during Nova activation indicated a 15 to 30% degradation in second- and third-harmonic performance at intensities beyond 2 GW/cm². Full-beam experiments with 74-cm arrays mounted in the Nova switchyard and at the target chamber showed an unexplained degradation in performance. Internal crystal conversion typically peaked at 50% at the third harmonic for intensities near 1.75 GW/cm², and sharply decreased for intensities beyond 2 GW/cm². Second-harmonic conversion typically reached 50% at 1.5 GW/cm², and appeared to reach a limit at $\sim 55\%$ for intensities beyond 2 GW/cm². Recent 3ω experiments conducted in the Nova

Fig. 6-44. Apparatus for 2.5-cm-aperture measurements used to study polarization dependence of 3ω conversion efficiency.

switchyard indicate that the major nonlinear effect responsible for Nova frequency conversion problems is the B -integral-dependent rotation of the principal polarization plane that affects the elliptically polarized regions of the Nova beam. Contributions to ellipse rotation from the air path, and any ellipticity change or rotation following reflection from the switchyard and target-bay turning mirrors, also complicate frequency conversion at the target chamber.

The beam depolarization at low intensity and effects at high intensity, such as beam divergence and beam breakup, are insufficient by themselves to account for the very pronounced and early rollover in conversion efficiency at moderate intensities (1.5 to 1.75 GW/cm^2). Schleiren tests of the Nova beam using the output spatial filter of beamline 10 show that the beam is well collimated (within $\pm 100 \mu\text{rad}$) at high intensity.

Two series of 3ω conversion experiments were conducted on beamline 10 in the Nova switchyard to verify the degradation of Nova harmonic conversion caused by ellipse rotation. Experiments were performed with 2.5- and 25-cm portions of the full beam with nominally 1-ns square pulses. Two $27 \times 27\text{-cm}$ -square antireflective-coated Nova production crystals were mounted in quadrature to a specially designed precision aluminum eggcrate. The eggcrate was positioned inside a large-diameter mirror mount, driven by a stepper motor. The mount was rotated so that the first crystal ordinary axis would be at 35.3° to the horizontal for optimum coupling to the nominally

horizontal Nova polarization. Precision Sony transducer probes were secured to the mirror mount to monitor the eggcrate rotation about the orthogonal ordinary and extraordinary crystal axes. Crystal orientation for phase matching was performed with the standard Nova procedure using the output sensor package on beamline 10.

The layout for the 2.5-cm-aperture 3ω conversion experiments is shown in Fig. 6-44. A badly depolarized portion of the full beam in the upper right-hand corner (looking into the beam) was chosen for these experiments. A polarizer was inserted before the 1-cm-thick (I_1 , I_2) doubling and mixing crystals in the 1ω beam to transmit horizontally polarized light. The light rejected by the polarizer was measured in a 1ω calorimeter.

The 2.5-cm third-harmonic generation (THG) conversion data closely match the theoretical curves (Fig. 6-45). The 1ω light fraction rejected by the polarizer is also shown. The peak conversion with repolarization exceeds 63% for intensities from 2.0 to 2.4 GW/cm^2 . The 1ω beam fraction rejected by the polarizer varies between 6 and 10% and is nearly flat over the 1.25 to 2.4 GW/cm^2 range. Without repolarization, the conversion efficiency peaks at $\sim 52\%$ to 54% at a drive intensity of 1.25 GW/cm^2 , and falls off for increasing intensity.

The 2.5-cm aperture 3ω conversion data are accurately modeled when ellipse rotation and beam repolarization effects are incorporated into harmonic conversion codes.^{73,74} This code includes the variation of the B -integral and ellipse orientation as a function of time in the pulse. Approximate analytical expressions are used for

Fig. 6-45. Third-harmonic conversion for 2.5-cm-aperture experiments vs model computations.

Fig. 6-46. Apparatus for 25-cm-aperture measurements used to study polarization dependence of 3ω conversion efficiency.

the "first" and "last" photon B -integrals for 1-ns pulses, and the B -integral is assumed to vary linearly through the pulse. The only variable parameters in this computation are the beam ellipticity ϵ , the initial tilt angle, and the angular deviation (θ_1, θ_2) from exact phase-matching in the doubling and mixing crystals.

Both data sets in Fig. 6-45 are modeled with an internal angular detuning $\theta_1 = \theta_2 = 175 \mu\text{rad}$ for the doubling and mixing crystals and $\epsilon = 0.25$. The initial

ellipse tilt angle is taken at the 35.26 deg polarization angle in the crystal reference frame, corresponding to a nominal horizontal Nova polarization. The $175\text{-}\mu\text{rad}$ detuning ($\pm 25 \mu\text{rad}$) is on the high side of what we would expect from our alignment procedures. The ellipticity of 0.25 ± 0.03 agrees with the time-averaged results discussed in the preceding article "Polarization Changes at High Intensity." An assumed polarizer transmission $T_p = 98.5\%$ (extinction ratio $\sim 2000:1$) agrees

well with the experimental conversion and loss data, and also with the measured transmissions for 2.5-cm Nova polarizers. With less angular detuning error, this experiment would have achieved the $\sim 72\%$ 3ω conversion at 2 GW/cm^2 that is expected with repolarization of a beam with an ellipticity of 0.25.

Third-harmonic conversion experiments were also performed at the 25-cm aperture. The layout of the apparatus is shown in Fig. 6-46. A sample of the Nova beam was taken on the right-hand side (looking into the beam) at approximately one-third the distance from the top of the beam. The 27×27 -cm Nova crystal pair, precision eggcrate, and mounting arrangement used in the 2.5-cm experiments were also used here. The 1ω drive energy was determined from the beam-10 output sensor package after recalibration to account for obscuration loss from placement of the 25-cm aperture ahead of the polarizer. The estimated drive intensity is based on the 1ω energy measured ahead of the polarizer. A removable large-aperture thin-film polarizer could be positioned before the crystals to repolarize the 1ω beam.

The experimental conversion data and the theoretical best-fit curves are plotted in Fig. 6-47 vs the estimated 1ω drive intensity. Without beam repolarization, the 3ω conversion peaks at ~ 70 to 72% at 1.8 GW/cm^2 , and falls to $\sim 52\%$ conversion at 2.7 GW/cm^2 . With repolarization, the conversion continues to rise for intensities beyond 1.8 GW/cm^2 and reaches $\sim 76\%$ conversion at 2.6 GW/cm^2 . When polarizer loss is subtracted, the estimate of peak conversion is $\sim 80\%$ at 2.4 GW/cm^2 , which is the upper limit reached so far on 3ω conversion in Nova production crystals.⁷⁵ An angular detuning of $125 \pm 25 \mu\text{rad}$ and an ellipticity of 0.0625 ± 0.0125 gives theoretical curves that are in good agreement with the data. The polarizer 1ω loss data is fit well, if we assume a polarizer transmission $T_p = 97.5\%$. This is close to the 98.5% measured for this polarizer. Compared to the 2.5-cm results in Fig. 6-45, beam repolarization at the 25-cm aperture has a less significant effect on conversion above 1.5 GW/cm^2 , because we expect a smaller fraction of depolarized light in this larger region of the Nova beam. This is confirmed by a value of ellipticity four

times smaller than that required to fit the 2.5-cm data, and by the low amount of 1ω light rejected by the polarizer.

To obtain information on the dynamical effects of ellipse rotation on harmonic conversion, streak-camera records of the 1ω input pulse and 3ω output pulse were obtained in the small-aperture conversion experiments. The two streak-camera channels were set up with fiber-optic links similar to those used in the ellipse rotation experiments discussed earlier. Figure 6-48 shows the 1ω and the corresponding 3ω streak records for the experiments with and without beam repolarization at an intensity of 1.85 GW/cm^2 . With the polarizer removed, the 1ω input pulse is substantially flat in time with a slight ($\sim 20\%$) increase in intensity at the tail end compared to the front end of the pulse. The resulting 3ω pulse (bottom left) exhibits a dramatic decrease in harmonic conversion during the pulse with the tail end showing some 55% lower conversion than the front end. With the polarizer in place, the input 1ω pulse exhibits a $\sim 20\%$ drop-off in intensity at the tail end, corresponding to the increasing fraction of 1ω light rejected from the polarizer during the pulse. The resulting 3ω pulse (bottom right), indicates a slight increase in conversion efficiency during the pulse.

Time-resolved ellipsometry data⁷³ obtained via the Stokes parameter measurements discussed in the preceding article indicate that the 1ω polarization ellipse

Fig. 6-47. Third-harmonic conversion for 25-cm-aperture experiments vs model computations.

Fig. 6-48. Polarizing the drive beam improves the pulse shape of the third harmonic in small-aperture tests.

can rotate a substantial amount (~ 10 deg) during the pulse. With the polarizer in place, the rejected light fraction should increase during the pulse and the transmitted 1ω should decrease. This is consistent with the 1ω streak data shown in Fig. 6-48. With repolarization, the 3ω conversion is substantially flat, despite the small loss in drive intensity, because the repolarized beam is able to maintain the optimum ratio of 1ω light along the ordinary to extraordinary crystal axes during the conversion process. Without the polarizer, the drive intensity does not suffer any loss, but the dynamical effects of ellipse rotation change the ratio of 1ω light along the crystal axes during the pulse. As the input

drive intensity is increased, the amount of ellipse rotation increases as the magnitude of the B -integral increases from the front end to the tail end of the pulse. Therefore, as a function of average 1ω drive intensity, the pulse-averaged 3ω conversion curves will roll over as the effect of ellipse rotation counteracts the normal improvement in conversion efficiency that should be seen with increasing drive intensity up to ~ 2.5 GW/cm² with 1-cm-thick harmonic crystals. This is consistent with the time-averaged conversion data shown in Fig. 6-45.

To verify the expected improvement in third-harmonic conversion efficiency with beam repolarization at full aperture, we

conducted experiments in beamline 10 in the Nova switchyard with a 74-cm harmonic array.⁷⁶ A prototype two-segment thin-film polarizer mounted in a special frame was secured to the optical table in the switchyard before the harmonic array. After transmission through the UG-11 filter array, the converted 3ω light was measured with the beamline-10 calorimeter on the switchyard wall. The intensity of the 1ω input drive was estimated from the full beam energy measured by the beamline-10 output sensor photodiode. The 1ω light rejected by the polarizer was also measured. The most important data points are shown in Fig. 6-49. These conversion results were obtained with nominal 1-ns square laser pulses with beam energies from 1.6 to 9.0 kJ. A peak 3ω conversion of $\sim 62\%$ was obtained with repolarization at intensities greater than $\sim 1.6 \text{ GW/cm}^2$ (conversion efficiency and drive intensity, $\pm 8\%$). Without repolarization, the peak conversion is $\sim 54\%$ from ~ 1.3 to 2.0 GW/cm^2 and drops to 36% at $\sim 2.5 \text{ GW/cm}^2$. The theoretical curves in Fig. 6-49 are for 200- μrad angular detuning of the mixing crystal and 100- μrad angular detuning of the doubling crystal. This is consistent with the experiment, since a difference of $\sim 160 \mu\text{rad}$ in wedge angle between the polarizer segments prevented optimum alignment in the sensitive axis of the mixing crystal. A beam-averaged ellipticity of 0.075 was assumed for the theoretical curves, which is consistent with the moderate value of ellipticity (0.0625) used to fit the 25-cm-aperture data.

The full-aperture conversion data shown in Fig. 6-49 is very encouraging for Nova. With repolarization, the highest energy shot (8.8 kJ) generated a 3ω energy of $\sim 5.4 \text{ kJ}$ (with the UG-11 filter array transmission and obscuration loss backed-out); this is the highest 3ω power per beam and highest 3ω energy at 1 ns generated on Nova. The prototype thin-film polarizer sustained optical damage over the course of these experiments. With a wedge-free, high-damage-threshold polarizer assembly, we expect $>70\%$ energy conversion to the third harmonic. At this time, it remains an open question as to whether thin-film dielectric polarizers will be able to survive the output fluence of the Nova laser. An

Fig. 6-49. Third-harmonic conversion results at full Nova aperture with drive beam repolarization.

alternate frequency-conversion scheme for THG that does not require beam repolarization is being considered,⁷⁷ and will soon be tested at the 25-cm aperture on beamline 10 in the Nova switchyard.

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Pulse Generation and Shaping

Nova's large disk amplifiers are operated at several times their saturation fluence to extract maximum energy. The resulting gain saturation causes pulse distortion, which is compensated for by controlling the pulse shape injected by the master oscillator. The master oscillator pulse-shaping system⁷⁸ also supplies the complex pulse shapes required by target experiments.

During 1985, the pulse-shaping system for the Nova master oscillator has been operating on a regular basis. Standard operating procedures were developed, and the system is workable and reliable. We measured the gain saturation of Nova's amplifier chains in order to calculate the

Fig. 6-50. Computer simulation of square pulse distortion for a Nova chain vs output energy.

pulse shapes needed at the laser chain input. We can now inject input pulse shapes that have a square profile at the chain output for energies up to 15 kJ per arm. We can also generate the "picket fence" pulses required for certain ICF target designs.

It is convenient to define the square pulse distortion (SPD) as the ratio of the gains seen by the leading and trailing edges of a square input pulse after propagation through the system. Using a computer simulation, we calculated the SPD for a Nova chain as a function of output energy (Fig. 6-50). This curve also shows the pulse shape that must be injected into a chain to produce a square output pulse. The contrast ratio shown in Fig. 6-49 between the leading and trailing edge of the pulse means that late portions of the pulse accumulate a larger B -integral than earlier ones. It also means that for a fixed output power, beam breakup and intensity-dependent polarization rotation will be greater for longer pulse lengths. Finally, because the slope of the SPD becomes very large for beam energies exceeding 18 kJ, it will be impractical to operate the Nova system much above this level and still provide well-controlled temporal pulse structure.

The emphasis of our pulse-shaping work in 1985 was to compensate for gain saturation so that we can produce a flat-top temporal profile at 1, 2, and 3ω . Good control of the input-pulse contrast ratio is required to produce satisfactory output pulse shapes, especially at higher energies (≥ 7 kJ per arm).

Arm-to-arm fluctuations in small-signal gain G_0 imply different input pulse shapes for the same output energy. Fortunately, the variations in G_0 are small enough such that no significant differences in the output pulse shapes are observed.

To predict the SPD, we begin with the small-signal gain,

$$G_0 = \exp\left(\frac{F_{st}}{F_{sat}}\right), \quad (24)$$

where

F_{st} = energy per unit area stored in amplifier,

F_{sat} = saturation fluence, and the trailing-edge gain,

$$G_t = \exp\left(\frac{F_{st} - \Delta F}{F_{sat}}\right), \quad (25)$$

where

ΔF = energy per unit area extracted from the amplifier.

Then, SPD is defined as

$$\frac{G_0}{G_t} = \exp\left(\frac{\Delta F}{F_{sat}}\right). \quad (26)$$

Defining the extraction efficiency η as $\Delta F/F_{st}$ gives $SPD = G_0^{\eta}$ for a given amplifier (for amplifiers at Brewster's angle, $SPD = G_0^{\eta/2}$). The system SPD is the product of the SPD from all amplifiers.

We used a computer code that predicts the system SPD to correlate Nova data with our saturation model. The code propagates the pulse from the spatial filter preceding the last rod amplifier to the frequency-conversion arrays. The procedure for predicting the system SPD for a given input streak is a four-step process:

1. Digitize the input streak.
2. Adjust amplifier gains to reflect measured energies on a specific shot (and chain) and calculate the input fluence necessary for a measured output energy.
3. Calculate system intensity gain as a function of input fluence.
4. Apply the gain function to the input streak.

The programs that perform the last three steps are written in the MATHSY language and run on the Octopus computer system. They include a

Fig. 6-51. (a) Streak camera lineout of a typical MOR-shaped pulse and the same pulse digitized. (b) Output streak at 1ω , 10 kJ and the same pulse as predicted by pulse-generation computer code.

two-gain cross-section model that calculates the saturation fluence F_{sat} fairly accurately as a function of output saturating fluence.⁷⁸ This process is similar to that used in MALAPROP to predict the input-contrast ratios required for flat-top output pulses. In this code, however, the gains are adjusted according to real energy measurements, and the pulse shape is defined experimentally.

Figure 6-51(a) shows the streak camera lineout of a typical shaped pulse from the master oscillator room (MOR) that was propagated through Nova; a digitized version of the pulse was also used by the code. Figure 6-51(b) gives both the output streak (1ω) at a measured energy of 10.5 kJ on beamline 7 and the code prediction for that energy. Note that the depth of modulation seen on the computer plot of the output may be present in the real pulse, but limited by the temporal resolution of the streak camera system.

Square pulse distortion is not the only measure of saturation. If the contrast ratio at the input is correct for the intended output energy but the shape is wrong, then the leading- and trailing-edge amplitudes at the output may be equal but the shape will not be flat. Consider that the gain seen by a particular portion of the pulse is a function of the energy extracted to that point in time $\Delta E(t')$. The extracted

energy is a function of the pulse shape since

$$\Delta E(t') = \int_{t_0}^{t'} I(t) [G(t) - 1] dt, \quad (27)$$

where $I(t)$ is the input pulse shape. Variations in $I(t)$ cause variations in $\Delta E(t')$ that affect the gain at t' . Therefore, shape and contrast ratios are equally important in producing the desired output shapes.

A temporal profile to be generated by the pulse-shaping system is translated into a plot of micrometer position vs distance along the variable impedance transmission line (VITL)⁷³ as shown in Fig. 6-52. The relation between the position of each micrometer and the optical power at a corresponding point in the pulse is determined experimentally. An approximate pulse shape is adjusted on the VITL and the resultant optical pulse is observed on the streak camera. The plot of the micrometer position vs optical power is compared with the required pulse shape, and iterations are made until the pulse is satisfactory. Micrometer positions are recorded and used when the same pulse is required again.

Early pulse profiles exhibited unwanted modulation that adversely affected the final output from the amplifiers. We found that much of this modulation came from poor impedance matching in the

Fig. 6-52. A graph of micrometer position vs distance showing a typical setting.

Fig. 6-53. A picket-fence pulse observed on the MOR streak camera.

connectors for the slicer Pockels cells. These Pockels cells receive a square electrical pulse with a rise time of <200 ps; any reflections of the rising edge from impedance discontinuities will result in modulation of the flat top of the square pulse at the Pockels cell. We lengthened the cables to move the reflections in time but some modulations still existed. At our request, Lasermetrics Inc. designed an improved Pockels cell that has wider bandwidth connectors and shorter connections to the

cell. Time-domain reflectometer tests of the prototype cell show smoother transitions from cable to cell. This device will soon be tested in the pulse-shaping system.

Another pulse shape that we are now able to produce is the aptly named picket fence shape. The picket-fence generator designed last year⁷⁸ was activated and tested in the MOR. A 100-ps pulse was split into three of the four legs and aligned onto the output beamline with ~ 1 -ns delay between the pulses. The alignment sensor was used to check that the three pulses were colinear. The optical system was designed to reflect the beams so that they emerge parallel, but in practice, small errors in the components cause variations in pointing between the beams that are resolved by this sensor. The beams can be manually adjusted so that these variations are eliminated. Figure 6-53 shows a streak-camera plot of the 3-pulse picket-fence waveform. The amplitudes were adjusted for equal heights, but any desired relative amplitudes can be set by simple wave

plate adjustments. Further tests will demonstrate the propagation of these pulses down the amplifier chains.

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Spectral Bandwidth of Nova Pulses

Concerns regarding harmonic conversion efficiency initially prompted us to determine the output spectral line width of Nova laser pulses. This measurement of line width, which had not previously been made, also provided additional characterization of the overall aspects of the output beam.

Although the pulse leaves the master oscillator room with a relatively narrow bandwidth (calculated at $<0.033 \text{ cm}^{-1}$), nonlinear mechanisms are capable of altering the width and character of the frequency spectrum. In particular, the nonlinear index n_2 can cause a variable Doppler shift in a pulse whose temporal intensity dependence is not linear (e.g.,

Gaussian and shaped pulses). The Doppler-shifted bandwidth produced in this way is directly proportional to n_2 .

The measured value of the line width of the MOR output pulse was instrument-limited at $0.056 \text{ \AA}$ because of the large free spectral range and low interferometer finesse of ~ 20 . The free spectral range was set to a value ($1.1 \text{ \AA}$) that would show only significant changes in bandwidth (i.e., widening of the rings); absolute wavelength measurements are not possible without a reference signal.

Figure 6-54 is a computer-enhanced picture of a small slice of the pattern produced by imaging the interferometer rings on a piece of Royal-X Pan film for a 7-kJ shot at 1 ns (7 TW). The adjacent plot is the microdensitometer trace of the slice. With the free spectral range at $1.1 \text{ \AA}$, the measured line width from this data is $\Delta\lambda = 0.08 \text{ \AA}$, which is very close to the instrument's resolution. It appears that no significant spreading takes place, although a higher resolution instrument might show interesting structure in the frequency domain.

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Fig. 6-54. Computer-enhanced slice of a Fabry-Perot interferometer pattern at 7 kJ, 1 ns and the corresponding microdensitometer trace.

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Interior of single-segment amplifier.

The development of an affordable multimegajoule laser as a fusion driver is an enormously challenging task, one that will require many innovations in laser technology. In our planning for achieving these goals, we break down the cost areas of fusion lasers into three major components: optical, electrical, and mechanical. The photos shown below on the left and opposite above, illustrate our work in nonlinear crystal development, a major optical cost center. The large photo below, left, shows the depolarization pattern produced by potassium dihydrogen phosphate (KDP), a biaxial crystal, when observed with laser light between crossed polarizers. We currently use KDP as a nonlinear frequency-conversion material. In studying new nonlinear materials, we use a spin-
dole-stage microscope to determine the refractive indices of samples as small as 25 μm in size as shown in the opposite-page, top photo.

Laser Research and Development

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Laser Research and Development

Introduction

J. F. Holzhrichter

The Advanced Laser Development (ALD) program has four principal objectives:

1. To support the activation of the Nova laser system for inertial confinement fusion (ICF) experimentation.
2. To develop a low-cost, 5- to 10-MJ laser for ICF research.
3. To develop efficient, high-repetition-rate (1 Hz to 1 kHz), high-quality lasers for ICF reactor applications, for the medium-power solid state laser program of the Defense Advanced Research Projects Agency (DARPA), and for the free-electron laser program of the Strategic Defense Initiative Office (SDIO).
4. To conduct basic laser research for energy (Office of Basic Energy Science) and for special tasks such as the development of UV laser coatings for the Air Force Weapons Laboratory.

The basis for our assumption of responsibilities beyond those of the ICF application arises from our intense R&D effort in solid state laser technology since 1971, when the ICF Program was formalized. These efforts, together with recent contributions from the Department of Defense (DOD)-sponsored programs described above, have led solid state laser technology into new operating regimes.¹ Our two reimbursable programs, described in Sec. 9, have as their objectives the demonstration of high-average-power solid state laser operating conditions for DOD applications. The DARPA program began late in fiscal year 1984, and the SDIO program began in fiscal year 1986. We now project solid state laser PRFs exceeding 1 kHz, efficiencies exceeding 10%, and beam quality approaching 1/10 wave rms, resulting in compact laser systems with great operational flexibility.

A major goal of our Nova support activities, described in Sec. 6, is to raise the output power limitation of the Nova system at 1.05 μm from its present 5 kJ per arm to at least 10 kJ per arm. This requires raising the damage threshold of laser glass and quantifying the rotational Raman scattering that occurs in the air path at high laser power. The production of platinum-free laser glass for Nova (the major step towards raising the damage threshold) appears to be compatible with low-cost, continuous-melting technology that strongly supports the cost and performance projections of our multimegajoule laser concepts, described below.

Our major goal continues to be meeting the ICF driver objectives. We have made substantial progress toward that goal: the technology developments described in this section form the basis for a high-energy, low-cost driver for high-gain ICF demonstrations, and for an efficient, moderate-cost laser driver for reactor applications.

Our primary near-term ICF laser design objective is to provide a basis for the construction of a versatile, 5- to 10-MJ, single-shot facility (the so-called high-gain test facility, or HGTF) with which to study high-gain ICF targets. In 1985, we continued to study megajoule-size solid state laser architectures, based on target-plane holographic injection and phase conjugation. Phase-conjugation measurements have begun, and a variety of systems using phase conjugation have been conceived. Several system studies, taking

account of modeling results, anticipated materials developments, and packaging considerations, indicate that a single-shot 5- to 10-MJ laser system could cost as little as 2 to 3 times as much as Nova.

We have made technical advances that affect all of our ALD work. These include the development of laser glasses and coatings, amplifiers, and optical switches, and the thorough study of nonlinear crystals for all applications and of laser crystals for advanced applications. Most of our effort has been in optical materials studies, which have the following goals:

- Acquiring a detailed understanding of damage in laser glasses, crystals, and optical coatings.
- Developing techniques for the rapid growth of nonlinear crystals.
- Gaining basic understanding of nonlinear optical crystals.
- Developing a methodology for the development of laser ions and hosts capable of efficient ($>10\%$), variable-PRF (1 Hz to 1 kHz) laser performance.

We have an essentially complete understanding of a high-efficiency glass-amplifier module (the single-segment amplifier, or SSA). Using the SSA we have demonstrated a storage efficiency of over 7%. We have completed studies of flashlamp efficiency, spectral output, and radiation pattern (including reflector contributions), and of the radiant efficiency of optical feedback.

We are developing a firm technical basis for the design of more powerful, energetic, and efficient solid state laser systems that are applicable to a wider variety of problems than have been considered in the past.

Solid State Amplifiers

Introduction

During 1984 and 1985 we performed experiments on a versatile Nd:glass amplifier that we call the single-segment amplifier (SSA). The SSA was designed to simulate a single segment of a large-array, rectangular-disk Nd:glass amplifier. The primary purpose of our experiments was to investigate the limits on the efficiency of flashlamp-pumped Nd:glass amplifiers. However many of the results apply to flashlamp pumping of solid-state disk lasers in general. Flashlamp pumping efficiency must be improved to reduce the cost of a multimegajoule laser for a high-gain test facility. Moreover, we must determine whether flashlamp pumping holds promise for reaching a laser system efficiency of 10%, which is a fundamental requirement for the driver in an ICF power plant.

Using the SSA we have demonstrated significant improvements in pumping efficiency over past designs. By optimizing the flashlamp packing fraction, the reflector configuration, and the pump pulse length, and by using flashlamp preionization, we have obtained peak

storage efficiencies as high as 7%, which occurs at an energy storage density of 0.05 J/cm^3 . At an energy storage density about equal to that obtained in the large Nova amplifiers (0.25 J/cm^3), we have demonstrated 4% efficiency. These results represent roughly a factor-of-3 improvement over the best Nova amplifier efficiencies at those energy densities. The improvement arises, in part, from the fact that past designs were directed at 1-ns output pulses and were thus optimized to maximize gain, not storage efficiency. This is explained in the next article. As the pulse length required for target implosion becomes longer, laser efficiency becomes more important than peak gain.

We have developed calculational methods to model our SSA results. These are discussed in "Single-Segment Amplifier: Modeling." We have used ray-tracing to analyze reflector designs, and we have used our fundamental measurements on the radiative properties of flashlamps (see the next subsection, "Flashlamps") to predict the time-dependent pumping efficiency of Nd:glass. The agreement between modeling and experiment is quite satisfactory.

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Single-Segment Amplifier: Experiments

Introduction. The optimum design of amplifiers for laser fusion drivers has changed as output energies have increased. Higher output energies drive larger fusion targets, which in turn require longer laser pulses. From Shiva through Nova to the concepts for a multimegajoule laser, the laser pulse length desired for target irradiation has increased by two orders of magnitude, from 0.1 to 10 ns.

The change to longer laser pulses has reversed the relative importance of gain and efficiency in amplifier design. In all cases we have attempted to minimize the system costs at a given output, and each system design has been constrained by damage and nonlinear self-focusing, the so-called ΔB limit.² However, the 0.1-ns pulse length of Shiva was short enough that the ΔB limit was reached at average energy fluences well below the saturation fluence of the laser glass. Shiva amplifiers were therefore designed for maximum gain, to minimize the beam path length in glass and thus minimize ΔB , with little regard for energy storage efficiency. For Nova, with pulse lengths of ~ 1 ns, the ΔB limit allows output fluences about equal to the saturation fluence, and energy storage efficiency is relatively more important. The pulse length for a multimegajoule facility

would be ~ 10 ns, allowing operation in the heavily saturated regime and making storage efficiency still more important. Efforts to improve amplifier performance are therefore concentrating on improving energy-storage efficiency.

As discussed in previous annual reports,^{3,4} we have constructed a versatile Nd:glass disk amplifier called the SSA. Figure 7-1 shows the rectangular-disk geometry of the SSA. Its square aperture is like that envisioned for each segment of a multimegajoule amplifier; it therefore tests the pumping configuration of such a multi-segment amplifier. Its primary purpose, however, is to provide a test bed for producing fundamental data against which to check modeling results and for demonstrating improved amplifier performance. This article presents results from experiments aimed at optimizing the flashlamp, drive circuits, and pumping configuration of the SSA with regard to storage efficiency and stored energy density.

Flashlamp and Drive-Circuit Optimization. To optimize the flashlamps and drive circuits, we varied the type of gas in the flashlamps, the fill pressure, and the input pulse length. We also operated the lamps with and without a preionization pulse, which preconditions the lamps for the main discharge. The flashlamps had cerium-doped-quartz envelopes with a 1.5-cm bore diameter and 234-cm arc

Fig. 7-1. SSA assembly showing $+x$ and $+y$ directions, one of two black end caps, and an absorbing panel that can be used in place of flashlamps and reflectors for one-pass pumping.

length. Most data were taken with xenon lamps, factory-filled to 300 Torr, but refillable lamps were used for this flashlamp and drive circuit optimization. The factory-filled and lab-filled lamps gave identical performance, within experimental error.

We compared SSA performance for different flashlamp pulse lengths and fill pressures, choosing the circuit capacitance C , inductance L , and initial bank voltage V to maintain a fixed value of the circuit damping parameter α (Ref. 5) at constant lamp explosion fraction f_x . The explosion fraction is the ratio of the total energy supplied to a lamp to the single-shot explosion energy of the lamp, and is therefore proportional to the total energy delivered to the lamp. The shape of the power pulse applied to the lamp depends primarily on the circuit damping parameter α . Matching α at constant f_x for circuits that deliver different pulse lengths $3T = 3(LC)^{1/2}$ requires that C , L , and V vary as $T^{5/6}$, $T^{7/6}$, and $T^{-1/6}$, respectively.⁶

The changes in C , L , and V required to optimize the lamp and drive circuit also required a new definition of storage efficiency. Previously, we defined storage efficiency as the ratio of energy stored in the laser glass to the energy stored in the

capacitor bank. However, changes in the gas type, the fill pressure, the circuit parameters, or the number of circuits fired caused small changes in the transfer efficiency of electrical energy from the bank to the flashlamps. To remove this effect, which obscures the performance differences in which we were interested, we re-defined storage efficiency as the ratio of stored energy in the laser glass to the energy (including the preionization energy) delivered to the lamps.

Xenon-filled flashlamps outperformed krypton and argon-filled lamps, as was the case in the open-flashlamp measurements.⁷ However, the storage efficiencies of krypton and argon were somewhat closer to those of xenon in the SSA than in the open-flashlamp measurements: the storage efficiencies for the SSA obtained with xenon, krypton, and argon were in the ratio 100/96/88 (for $f_x = 0.18$ and $3T = 670 \mu\text{s}$), while the corresponding neodymium-pumping efficiencies with an open flashlamp under similar operating conditions were in the ratio 100/93/83. We believe that this difference results from the reduced opacities of krypton and argon relative to xenon. The lower opacities produce lower radiant efficiencies, but they

Fig. 7-2. SSA storage efficiency vs explosion fraction, with and without preionization pulses to the flashlamps, for three flashlamp pulse lengths.

also give a lower absorption of pump light. In an amplifier, the laser glass competes more favorably for pump light using krypton- or argon-filled lamps and gets a larger fraction of the total pump light. In the SSA, the lower radiant efficiencies of krypton and argon are partially compensated by the correspondingly reduced flashlamp absorptions.

Varying the fill pressure of xenon in the SSA lamps gave very flat maxima in efficiency between 150 and 300 Torr, independent of $3T$ and f_x . The efficiency dropped less than 2% for factor-of-2 changes in pressure about the maxima. This behavior is consistent with the open-flashlamp results.

Preionization improved flashlamp performance under all operating conditions, but did so most markedly for short input pulses and relatively low input energies. The preionization pulse delivered 145 J/lamp (0.15 J/cm^2) 200 μs ahead of the main pulse. Figure 7-2 shows storage efficiency vs f_x for three pulse lengths, each with and without preionization. The data were taken with 5-lamp Rabl-shaped reflectors,⁸ with blast shields between the disks and flashlamps and with black end-plates on the SSA. We determined error bars as described previously.³ Figure 7-2 shows

that efficiency without preionization is a maximum for $3T \approx 290 \mu\text{s}$, but that with preionization the peak efficiency continues to increase for shorter pulse lengths. Preionization gives a 23% increase in peak efficiency for the 170- μs pulse length.

Figure 7-2 also shows that preionization slightly reduces the transfer efficiency to the lamps. The with-preionization data points are all shifted slightly to the left with respect to those without preionization for the same initial conditions, indicating a slightly lower delivered energy. This presumably occurs because preionization reduces the average lamp impedance.

Figure 7-3 shows the effect of shorter pulse lengths on SSA performance. The figure shows efficiency vs stored energy density for six pulse lengths, with preionization. We characterize the radiative lifetime of the SSA disks by the first e -folding time of the laser glass, since LG-660 glass does not give a single-exponential radiative decay. The average of the first e -folding times for the $5 \times 5 \times 0.5\text{-cm}^3$ witness samples made with the SSA disks was 350 μs . The six values of pulse length shown in Fig. 7-3 therefore represent ratios of pulse length to lifetime ranging from 0.5 to 2.7.

Fig. 7-3. SSA storage efficiency vs stored energy density for six flashlamp pulse lengths, all with preionization.

Figure 7-3 shows that reducing the input pulse length greatly improves storage efficiency at low stored energy density, but that the optimum pulse length depends on the desired operating point for the amplifier. The curves define a maximum-performance envelope, with different pulse lengths required for maximum efficiency at different stored energy densities. The individual curves are terminated at their high stored-energy-density ends by the maximum safe voltage for our circuit components.

Pumping Configuration Optimization.

We tested the seven reflector geometries shown in Fig. 7-4. The labels in the figure give the number of lamps per side and indicate the shape (cylindrical, flat, or Rabl). (What we now call the 5C reflector was referred to as the 5-lamp Hagen reflector in previous *Laser Annuals*). All reflectors were silver-coated, some by electroplating and some by vacuum deposition, and all were overcoated with MgF_2 or SiO_2 . We could not obtain accurate spatially averaged reflectivities in some cases because of the short and/or variable radii of curvature, but the measured reflectivity differences

Fig. 7-4. Cross sections of the seven reflector geometries tested.

do not significantly change the conclusions presented here.

Figure 7-4 shows reflectors for which both the shape and the flashlamp packing fraction vary. Flashlamp packing fraction is the projected area of the lamps divided by the total area of the output plane of the flashlamp panel. Our definition of projected lamp area is based on the 1.5-cm bore diameter of the flashlamps. The packing fractions were 0.14, 0.19, 0.23, and 0.38 for the 3-, 4-, 5-, and 8-lamp reflectors, respectively. Packing fraction and reflector shape are not entirely independent of one another, since a specific shape cannot be preserved for a change in lamp packing fraction. However, these quantities affect different aspects of the pumping, and therefore might be expected to optimize differently.

Flashlamp packing fraction determines the total pumping power inside the amplifier, and greater packing fractions give higher gains and stored fluences. However, a simple physical argument⁹ based on the blackbody-like character of the lamp output suggests an optimum packing fraction of $1/\pi$ when efficiency is considered, even for the highest possible gains and stored fluences. The flashlamps are in fact somewhat gray rather than black, so the optimum lamp packing fraction of $1/\pi$ is only approximate.

The reflector shape directly determines the three principal factors that decrease the transfer efficiency of pump light on the first pass from lamps to disks. These factors are the number of reflections from the silver-coated surfaces, the amount of pump light lost out the amplifier ends, and the amount of direct feedback to the lamps. Reflector shape also strongly influences the angular distribution of pump light, which in turn determines the amount of Fresnel reflection from the disks. However, Fresnel reflection contributes loss only by increasing one or more of these factors.

An accurate ranking of the seven reflector geometries with regard to direct feedback, average number of reflections, and end loss requires a self-consistent three-dimensional analysis of each geometry, which we did not attempt. However, ray-tracing calculations, described in the next article, indicate very similar numbers of reflections for one-pass pumping for all seven reflectors. A direct measurement of end

loss for one-pass pumping (described below) shows relatively little difference among the seven cases. This implies that direct feedback is likely to be the principal difference among the reflector geometries, at least for one-pass pumping.

We can estimate the ordering with regard to direct feedback among similar shapes and for different shapes with the same packing fraction. None of the light emitted from a lamp in a Rabl-shaped reflector can go directly back into that lamp, so direct feedback is due entirely to Fresnel reflection from the disks. Furthermore, the width of the angular distributions in the plane transverse to the lamp axis increases with lamp packing fraction for the Rabl reflectors, going from $\pm 30^\circ$ for the 3R reflector to $\pm 45^\circ$ and $\pm 80^\circ$ for the 4R and 5R reflectors, respectively. As a result, Fresnel reflection (and therefore direct feedback) increases with lamp packing fraction for the Rabl reflectors. Direct feedback is greater for the cylindrical and flat reflectors than for the Rabl reflectors, because they have the additional component from the lamp reflector. For the cylindrical reflectors the reflector cross sections in Fig. 7-4 make it clear that the total direct feedback increases

with lamp packing fraction. In summary, direct feedback increases with lamp packing fraction for similar shapes, and it is greater for flat and cylindrical shapes than for the Rabl shape at a given packing fraction.

We tested each reflector geometry in one-pass pumping configurations as well as in the usual two-sided configuration. One-pass pumping is easier to interpret because it eliminates the effects of the multipassing of pump light not absorbed on the first pass and absorption and recycling of light in the flashlamps. To obtain one-pass pumping data we replaced one flashlamp panel and reflector by the black absorbing panel shown in Fig. 7-1, so as to absorb all pump light after one pass through the disks. Two one-pass pumping configurations are possible, corresponding to lamps on the $+x$ and $-x$ sides of the amplifier. (Figure 7-1 shows the $+x$ and $+y$ directions.) Gain and efficiency differ in the two configurations, because lamps on the $+x$ side radiate some light directly out the ends of the amplifier, whereas all pump light from the $-x$ side passes through at least one disk.

Figure 7-5 shows one-pass pumping storage efficiencies (defined as for Fig. 7-3) vs

Fig. 7-5. One-pass storage efficiency vs explosion fraction for the seven reflector geometries with ASE effects removed.

Fig. 7-6. End-loss transmission factors vs explosion fraction with ASE effects removed (see legend for Fig. 7-5).

explosion fraction for the seven reflector geometries, with the lamps on the $-x$ side of the amplifier. This presentation hides a major influence of lamp packing fraction, namely the different amounts of total pump energy available at different packing fractions. It compares performance at the same lamp conditions, and thus shows only differences in the one-pass lamp-to-disk transfer efficiency. These data were taken without blast shields, with black end caps, and with a flashlamp pulse length of $660 \mu\text{s}$. The data have been corrected slightly to remove the effects of amplified stimulated emission (ASE).³ Under these conditions, ASE led to a linear reduction in efficiency with peak stored energy density, amounting to an 11% decrease at 0.20 J/cm^3 .

The general shape of the curves in Fig. 7-5 reflects the change in radiative efficiency and spectrum of an open flashlamp, as is explained in "Single-Pass Pumping," later in this section. The ordering of the curves seem to scale most strongly with lamp packing fraction. The higher packing fractions have lower efficiencies that fall off slightly faster with f_x . At a given packing fraction, the cylindrical and flat shapes give slightly lower efficiencies and fall off faster

than the Rabl shape. A likely explanation for these differences is the variation in direct feedback among the reflector geometries. As described above, the amount of direct feedback increases with lamp packing fraction, and it is greater for cylindrical and flat shapes than for the Rabl shape. Direct feedback leads to a faster falloff in efficiency with f_x because lamp absorption increases with f_x . The changes in direct feedback are thus consistent with the changes in ordering and falloff of the curves.

The differences in one-pass storage efficiencies between pumping from the $+x$ and $-x$ sides give a direct measure of end loss. We define an end-loss transmission factor T_{el} from the one-pass pumping efficiencies from the two sides $\eta(+x)$ and $\eta(-x)$ as follows:

$$T_{el} = \frac{\eta(+x) + \eta(-x)}{2\eta(-x)} \quad (1)$$

Figure 7-6 shows T_{el} vs f_x for all the reflector geometries. Clearly, SSA end losses are relatively small: only 5 to 8% of the useful pump light is lost on the first pass of the light across the amplifier. However, end loss is proportional to the number of disks

Fig. 7-7. Two-sided storage efficiency vs explosion fraction with ASE effects removed (see legend for Fig. 7-5).

and would be twice as great for a two-disk amplifier. Figure 7-6 also shows that end loss has no significant dependence on f_x for any of the reflector geometries, and that all but the 5C case give very similar T_{el} values. Both results are surprising, considering that the angular distribution of light varies with f_x (Ref. 6) and with reflector geometry. The 5C case gives distinctly higher T_{el} values, indicating an anomalously strong collimation of flashlamp light.³

Figure 7-7 shows two-sided pumping data for the seven reflector geometries, with a 660- μ s pulse length, with black end caps, and no blast shields, and with ASE effects removed as before. These data show somewhat higher efficiencies and a faster falloff with increasing f_x than the one-pass results in Fig. 7-5, but the ordering of the curves is essentially the same. There is surprisingly little difference among the reflector geometries at a given packing fraction, considering the substantial differences in reflector shape. Furthermore, our reflectivity measurements indicate reflectivities in the high-80% range for the 5C and 5R cases, as against the mid-90% range for the others. Consequently, some of the efficiency differences between 5C

and 3C and between 5R and 4R can be attributed to reflectivity differences.

Figure 7-8 shows the same data as Fig. 7-7, but without the ASE correction and with storage efficiency plotted against stored energy density. This presentation shows the total amplifier performance for each pumping configuration, and it substantially changes the relative positions of the curves. The lower packing fractions reach higher efficiencies than the higher packing fractions only at low stored energies, because their efficiencies decrease as the lamps are loaded more heavily. Higher stored energies require more lamps to avoid driving the lamps too hard. The 4R reflector performs the best over its entire operating range, although the difference between the 4R and 5R at maximum stored energy density might be attributed to reflectivity differences. The best lamp packing fraction depends on the desired operating point of the amplifier. For stored energy densities below 0.24 J/cm³, the 4R case gives the best performance. Above that the 5R would be best, and a 6-lamp configuration might even win near 0.30 J/cm³; but 8 lamps per side is too many. The optimum packing fraction for

Fig. 7-8. Two-sided storage efficiency vs stored energy density (see legend for Fig. 7-5).

maximum stored energy density appears to be between 0.23 and 0.38, in reasonable agreement with the $1/\pi$ figure obtained assuming blackbody lamp behavior.

Two-sided and one-pass pumping differ only because of end losses and the effect of multipassing and recycling of pump light unused after one pass through the disks. Both of these effects are absent in the one-pass case of Fig. 7-5, and both are present in the two-sided case of Fig. 7-7. To see the effect of multipassing and recycling alone, we averaged the one-pass efficiencies for pumping from the $+x$ and $-x$ sides to introduce the same relative amount of end loss in the average as is present in the two-sided data. Dividing the average into the two-sided data (with ASE effects removed) eliminates end loss and leaves only the multipassing and recycling effect. Figure 7-9 shows these ratios.

Figure 7-9 shows that the combined benefit of multipassing and recycling decreases with f_x for all seven reflector geometries. Multipassing and recycling compete: pump light not absorbed on the first pass through the disks is either reflected back through the disks (in multipassing) or is absorbed by the flashlamps and re-emitted as broadband pump light (in recycling). We know that flashlamp

absorption increases with increasing f_x , whereas multipassing is independent of f_x ; this means that recycling increases at the expense of multipassing with increasing f_x . The decrease in the efficiency ratio with f_x therefore implies that multipassing improves efficiency more than recycling does. This conclusion presumably depends on the doping-thickness product of the disks, because the efficiency improvement from multipassing decreases with increasing disk absorption, making recycling relatively more important. With a sufficiently high disk absorption, the efficiency improvement from recycling should exceed that of multipassing, but the $17.2\% \cdot \text{cm}$ doping-thickness product of the SSA disks is apparently not high enough for this to be the case.

Figure 7-9 also shows that the combined effect of multipassing and recycling gives only a small efficiency improvement over one-pass pumping. At best the improvement is only 25% at low lamp loading, and only 15% at high loading. A lower doping-thickness product should increase the efficiency ratio, because a greater fraction of the pump light would be available for multipassing and recycling.

Spatial Variations. Gain uniformity was a serious problem for those reflector

Fig. 7-9. Ratios of two-sided to one-pass pumping efficiencies vs explosion fraction with ASE effects removed (see legend for Fig. 7-5).

geometries in Fig. 7-4 with only 3 or 4 lamps per side. Figure 7-10 shows vertical (y -direction) gain scans for the 3R, 4R, and 5R reflectors. Large spatial oscillations are evident for the 3R and 4R cases. They change phase with distance from the lamps, creating local maxima and minima throughout the aperture. Overlapping radiation patterns from adjacent reflector segments presumably cause these oscillations, because Rabl reflectors produce a two-lobed pattern (see the next article). The 4R data show that vertical scans at symmetric positions with respect to the centerline have the same shape, as would be expected. The amplitudes of the spatial oscillations are lower and decrease more quickly with distance from the lamps for the 4R case than for the 3R case. The 3C case (not shown) also produces spatial oscillations, but only near the lamps. They damp out completely for the central 60% of the amplifier aperture. All reflectors with 5 or more lamps per side have smooth profiles, with none of the spatial oscillations seen with the 3- and 4-lamp reflectors.

Figure 7-10 also shows a gain drop at the top and bottom of the clear aperture, caused by the blockage of pump light by the disk holder. Although

the effect is most obvious here for the 5R case, which is otherwise very flat, all the reflectors tested show a similar effect.

Averaging the vertical gain profiles shown in Fig. 7-10 and plotting the averages vs horizontal position gives the averaged horizontal gain profiles shown by the dashed curves in Fig. 7-11. The solid curves in the figure show horizontal gain scans through the amplifier centerline for comparison. The slopes of the averaged horizontal gain profiles result from end loss, which leads to a variation in pumping across the aperture at the ends of the amplifier. The outside surface of the end disks sees less pump light on the $-x$ side than on the $+x$ side of the amplifier. All disk amplifiers with an even number of disks have this sloped gain profile. Figure 7-11 shows similar slopes for the four cases displayed; in fact, all cases tested had the same slope to within $\pm 3\%$. This implies little variation in end loss, in agreement with our earlier conclusion based on the more direct measure of end loss shown in Fig. 7-6.

We needed spatially averaged gains for all the one-pass and two-sided pumping configurations to obtain the performance data shown in Figs. 7-5 through 7-9. Since

Fig. 7-10. Vertical gain profiles for the three Rabl-shaped reflectors.

Fig. 7-11. As-measured profiles (solid curves) and vertically averaged horizontal gain profiles (dashed curves) for four reflector geometries.

we measured gain vs bank energy at the amplifier centerline, we defined a uniformity correction factor to account for the difference between centerline gain and the spatially averaged gain. We obtained the spatially averaged gain from an area-weighted average over measured spatial gain data. The resulting uniformity correction factors for the 3R and 4R cases were 0.880 and 0.945, respectively. The other five reflector shapes gave uniformity correction factors between 0.975 and 0.995. Gain profiles for one-pass pumping were similar to the corresponding profiles for two-sided pumping, and they gave similar uniformity correction factors.

We experienced shot-to-shot reproducibility problems with the 3- and 4-lamp Rabl reflectors that we have not experienced with other reflectors. We suspect that variation in the arc position in the lamp bores was the cause. For other reflector configurations the pump light from a given lamp spreads rapidly with distance from the lamp, so that the gain at any point is produced by light from several lamps.

Pump light from the 3R and 4R reflectors spreads more slowly, making gain at a given point more strongly dependent on individual lamp/reflector elements and correspondingly more sensitive to arc position.

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Single-Segment Amplifier: Modeling

Ray Tracing Calculations.

Introduction. In 1985 we extended the radiation transport calculations reported in the 1984 Laser Annual¹⁰ to other reflector shapes and found effective reflectivities for them. The calculated differences in pumping performance between the different shapes are small enough that they can

easily be outweighed in practice by differences in reflectivity. We calculated the spatial variations in the pump distribution transverse to the lamps. Agreement with experiment was within 2% for the 3-lamp Rabl reflector; the observed gain variations for the 4-lamp Rabl reflector were about 5% greater than calculated.

The purpose of this work is to understand and improve amplifier performance. In general, the performance of various reflector shapes is quite similar for most designs with the same lamp-packing fraction. Increasing the lamp-packing fraction to obtain higher gain reduces pump efficiency because of greater lamp absorption of pump light (see the preceding article and Ref. 11). Minimizing the cost of the amplifier per joule of laser energy may require higher gains than obtainable at peak pump efficiencies. Thus it may be necessary to compromise on pump efficiency to reduce the cost of hardware and optics, at an increase in bank cost.

Fig. 7-12. Angular emission profiles for Rabl (R) and cylindrical (C) reflectors.

The most noticeable differences between the various reflector designs are their 2-D emission profiles, which we calculated by ray tracing (Fig. 7-12). Despite appearances, these curves are all normalized to have the same integral over azimuth angle. (The area around the origin has a large weighting.) The Rabl and cylindrical reflectors (designated 3R, 4R, 5R, 3C, 5C, and 8C, indicating the number of lamps per side) were studied in the SSA experiments described in the preceding article. The Rabl reflectors avoid reabsorption of pump light on its initial path, but have wider emission profiles than cylindrical reflectors. Emissions at large angles lead to greater reflection losses on the side mirrors and blastshields. The cylindrical reflectors provide better collimation, but the flashlamps block some of the reflected light. The profiles in Fig. 7-12 are for both transparent lamps and opaque lamps (neglecting reemission). The intensities decrease for the cylindrical reflector, but not for the Rabl reflectors, as the lamps become opaque. As explained below, the trade-off is higher reflection loss for the Rabl shapes against higher lamp absorption for the cylindrical shapes.

The 3R and 4R reflectors provide better collimation than the 5R reflectors, but require larger lamp spacing and hence limit the amount of pump light. The 3R and 4R reflectors have cutoff angles of about 30° and 45°, respectively, and they emit more strongly near the cutoff angles, as shown in Fig. 7-12.

The 5C reflectors provide a good balance between lamp absorption and collimation. The emission profiles for the 3C and 8C reflectors are also shown in Fig. 7-12. Neither of these designs is optimal: the reflector curvature of the 3C is too flat, and that of the 8C is too steep. The 3C has double-lobed emission, while a well-collimated emission profile could be obtained with a reflector approximating a

parabola with the lamp at the focus. The emission profile of the 8C is wider than necessary, because the angles at the apex of the reflector approach 90°.

Effective Numbers of Reflections. The various reflectors can be characterized by intensity-weighted effective numbers of reflections for emission and reemission, which are given in Table 7-1 for initial and subsequent passes within the pump cavity. For example, it takes an average of 1.05 reflections for light to reach the amplifier disk from a 3R reflector ("emission"), and 2.7 reflections for radiation entering the 3R reflector at low lamp opacities to reach the disk ("reemission"). In the reemission calculations, we used the output angular distribution of each reflector as the input distribution entering that reflector. The more-collimating Rabl reflectors generally trade off the number of reflections in emission with reductions in the number of side-wall reflections. For one-sided pumping, the total effective numbers of reflections are within 2% of one another for the Rabl reflectors; the number for the best cylindrical reflector, the 5C, is about 20% lower than this. The 5C's advantage is reduced at high lamp loading, for which reabsorption of pump light by the flashlamps becomes significant.

Pump Distributions. Figure 7-13 shows the calculated transverse pump profiles with one-sided pumping for various SSA reflectors, for distances from the lamp panels corresponding to the experimental scan positions. The calculations neglect lamp absorption. As expected, the 3-lamp reflectors exhibit the greatest nonuniformities; otherwise, the pump distributions appear quite uniform. The distributions generally become smoother farther from the lamps. Small ripples are primarily caused by statistical fluctuations in the calculations themselves, and should be ignored. Different lamp opacities have little effect on the transverse pump profiles except to reduce the pump intensity. With two-sided pumping, some of the strong features close to the lamp (for example, those observed with the 3C reflector) are washed out by emission from the opposite side.

The measured and calculated pump distributions in the central vertical plane of the SSA are shown in Fig. 7-14. The

Table 7-1. Effective number of reflections from lamp to amplifier disk for Rabl (R) and cylindrical (C) reflectors.

Number of lamps/reflector	3R	4R	5R	3C	5C	8C
Initial emission	1.05	0.93	0.70	0.57	0.61	0.67
Initial side reflection	0.24	0.35	0.64	0.57	0.49	0.63
One-sided pumping ^a	1.29	1.28	1.34	1.14	1.10	1.30
Reemission (transparent lamps)	2.7	2.2	1.4	1.1	1.1	1.5

^a Sum of two values immediately above.

Fig. 7-13. Transverse pump profiles for various distances from the lamp panels, with the reflector shapes of Fig. 7-12.

Fig. 7-14. Transverse pump distribution in the central vertical plane of the SSA for 3-lamp Rabi reflectors.

agreement is very reasonable despite the 2-D simplifications used for the calculations. 3-D corrections should not alter these distributions significantly because the projected emission profiles¹² from the lamp/reflector combinations are essentially identical for a particular receiver plane pumped by different axial sections of the flashlamps.

The relative performance and pump distributions of various reflectors can be approximated quite well with 2-D simulations. The predicted pumping performances of all 3-, 4-, and 5-lamp reflectors are well within 10% of one another, so that actual performance is likely to be dominated by variations in reflectivity. It is more difficult to obtain uniform pump distributions with fewer lamps. The agreement between experiments and simulations is good enough that we can scale amplifier performance using the simulations.

Author: W. F. Hagen

Single-Pass Pumping. In the 1984 Laser Annual we described¹³ a simple model for calculating the single-pass pumping efficiency of the SSA. In that model we assumed that the flashlamp excitation of the laser glass was an instantaneous function of the electrical input power to the lamps. Our studies of xenon flashlamps have yielded more accurate spectra and radiative efficiencies, which significantly affect our estimates for the pumping efficiency of Nd:glass. In addition, as we describe in "Flashlamp Modeling," later in this section, we now believe that the radiative characteristics of xenon flashlamps are better described as a function of their output optical power. The output power may in turn be calculated from the electrical input power using a simple nonlinear differential equation that reduces, for long pulse widths, to the conventional quasi-stationary description of flashlamp behavior. We have used our updated flashlamp

data and model to obtain an improved prediction of the storage efficiency of the SSA for single-pass pumping.

We calculated the Nd³⁺-pumping efficiency by convolving the quasi-stationary output spectrum of the 1.5-cm-i.d. lamp, at a xenon fill pressure of 300 Torr, with the wavelength-dependent absorption of the laser glass. This procedure is described in detail in Ref. 13. For comparison with our earlier results,¹² Fig. 7-15 shows the pumping efficiency as a function of the electrical input power per unit area. The rollover in the efficiency that was present in our previous curves does not appear here, mostly because the lamps used in the present work were sufficiently preionized that the instantaneous radiative efficiency fell off less at low input power. The remaining falloff is made up for by the better match, at lower input powers, between the lamp spectrum and glass absorption.

Fig. 7-15. Calculated efficiency for xenon-lamp pumping (1.5-cm bore, 300-Torr xenon, 400-nm cutoff) of LG-660 laser glass (doping-thickness product 1.64×10^{21} Nd³⁺ ions/cm³) relative to electrical input and to lamp output power per unit area.

Fig. 7-16. Predicted and measured SSA Nd³⁺ storage efficiencies: (a) actual values, (b) ratio.

To account for the transient behavior of the lamp radiation (see "Flashlamp Modeling"), we used the pumping efficiency relative to the output power of the lamp in our calculation of the SSA storage efficiency. This quantity is plotted as a function of the observed output power in Fig. 7-15. The solid line, a fourth-order polynomial fit, was used to model the amplifier performance.

In the updated calculation we used the observed values of the electrical power input to the lamp, rather than attempting to predict the lamp's electrical behavior using the circuit equations of Markiewicz and Emmett.¹⁴ This eliminated small complications caused by the nonlinear behavior of the electrical switch in the lamp circuit and by changes induced in the electrical characteristics of the lamp by feedback of radiation. The time-dependent power output from the lamp was then determined according to the method described in "Flashlamp Modeling."

Once the time-dependent output power P_L and the instantaneous Nd:glass pumping efficiency $\eta(P_L)$ were known, the amount of energy E_s stored in the Nd³⁺ inversion was determined by solution of the equation

$$\dot{E}_s = \eta(P_L)P_L(t) - \Gamma_{\text{decay}}(t); \quad (2)$$

where $\Gamma_{\text{decay}}(t)$ is the effective decay rate for the stored energy, which was modeled by a three-component exponential fit as in Ref. 13.

The results of this calculation are presented in Fig. 7-16 for data obtained using the 3-lamp Rabl reflector, which gave the highest measured single-pass efficiencies, implying the most efficient transport of energy from flashlamps to disk. The solid curves are smoothed spline fits to the data, suitably weighted to account for experimental error. Aside from a slight increase in the calculated efficiencies relative to those observed at the lowest values of the explosion fraction, the overall flatness of the curve in Fig. 7-16(b) is an indication that the model correctly treats the variations in storage efficiency due to the blue shift of the lamp output with increasing input power.

The ratio of observed to predicted efficiency, about 75%, provides an estimate of

the transport efficiency of the pump radiation from lamp to disk, since the model does not take any losses into account. We can estimate the effects of the two primary sources of loss in the amplifier: the absorption of pump light by the reflectors and the blockage of the pump light by the disk holders. The reflectivity of the 3-lamp Rabl reflector, spatially averaged and spectrally averaged (over the five main pump bands) is $96 \pm 2\%$; two-dimensional ray-tracing (see the preceding article, "Ray Tracing Calculations") indicates that the average number of reflections from lamps to disk is approximately 1.3. These data indicate that the transport factor due to reflectivity considerations could be as low as 92%. Similarly, we estimate that as much as 8% of the light may be prevented from reaching the glass by the disk holders. Combining these two effects yields an expected transport efficiency of approximately 85%, somewhat higher than the observed 75%. This discrepancy may be resolved by the development of a more refined three-dimensional amplifier model that accounts for Fresnel reflections at the glass surfaces and the absorption of pump radiation by the flashlamps themselves. But we believe that the agreement is already close enough to support our quantitative measurements of the spectral output intensities of flashlamps.

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Flashlamps

Introduction

Since 1982 we have been measuring the absolute, spectrally and temporally resolved output of pulsed flashlamps under a wide variety of input powers, gas types, fill pressures, bore diameters, pulse lengths, and other parameters. This information will allow us to choose optimum flashlamp parameters for future solid state lasers. It is also necessary in the construction of accurate models of flashlamp pumping of solid state lasers, and it has been particularly useful in making predictions that can

be compared with the performance of the SSA (see "Single-Segment Amplifier: Modeling," earlier in this section).

The experimental portions of this study, which are now complete, have yielded surprisingly fundamental results. As described in the next article, we have discovered that the output spectra of pulsed lamps of different bore diameters are, at steady state conditions, nearly identical when compared at equal input power per unit area of the bore wall. This is contrary to the assumption, made in connection with previous models, that spectral equivalency occurs either at equal values of current density squared times bore diameter or at equal input power per unit volume. Although our result is to be expected for optically thick radiators, it is not clear why it applies to flashlamp plasmas whose opacities are intermediate between thick and thin.

"Diameter Scaling" describes our work on the time dependence of flashlamp emission. For pulse lengths of interest ($>100 \mu\text{s}$) and with appropriate preionization of the flashlamp gas, we find that, for practical purposes, the output spectral distributions measured as a function of time are well described by the steady-state results. The best agreement occurs at equal output optical power per unit area, rather than equal input electrical power per unit area; this modification takes the finite response time of the flashlamp plasma into account.

The results described in these two articles make possible a relatively simple model of flashlamp output, as described in "Flashlamp Modeling." Flashlamp spectral distributions have been analyzed to obtain a plasma temperature and a spectrally dependent emissivity as a function of input (or output) power. A simple method, using a single time constant and taking into account the reduction in efficiency caused by arc formation and expansion, has been found to predict the total time-dependent output power as a function of input power. Knowing the input vs time, we can then predict the spectrally resolved lamp output.

"Rapid Scan Analyzer" describes the design and performance of a device that can obtain snapshots (at $18\text{-}\mu\text{s}$ intervals) of flashlamp spectra with moderate spectral

resolution from a single shot. This device has been particularly useful in allowing us to see the changes in spectral output caused by lamp preionization.

Finally, in "Reflective Coated Flashlamps," we describe our work on multilayer dielectric-coated lamps that recycle radiation that is normally unused in laser pumping back into the plasma. Improvements of up to 12% in peak Nd* signals have been demonstrated using UV-reflective coatings. These studies have set a lower bound of 0.75 for the relative efficiency of optical (in particular, UV) input to electrical input in producing lamp output.

Author: H. T. Powell

Pulse Length Scaling

Introduction. Fluorescence decay is a major loss mechanism in neodymium:glass lasers when the flashlamp pulse length is long compared with the fluorescence lifetime of the laser glass. The longest fluorescence lifetimes obtainable at neodymium-doping concentrations appropriate for megajoule-class fusion lasers and for high-average-power lasers are approximately 300 μ s and 200 μ s, respectively.¹⁵ Flashlamps for fusion lasers are usually operated at pulse lengths of $\sim$ 600 μ s to maximize the peak gain.

Higher efficiency can be achieved at short pulse length (see "Single-Segment Amplifier: Experiments," earlier in this section), but short-pulse operation introduces several difficulties. These include^{16,17} reduced flashlamp lifetime, reduced radiant (electrical-to-optical conversion) efficiency, blue-shifting of the lamp emission away from the neodymium-pump bands, and increased lamp absorption. This article describes the changes in the radiant efficiency and spectral output of flashlamps with pulse length, and the remarkable effects produced by lamp preionization.

Time-resolved spectra at various pulse lengths are needed to test models of amplifier and flashlamp behavior.¹⁸ We generally depend on the quasi-stationary state (QSS) approximation, which is a fundamental assumption of the Trenholme-Emmett model.¹⁹ This approximation

assumes that the instantaneous lamp output is described as a function only of the instantaneous input power to the lamp, regardless of whether that power is rising or falling. As pulse length is reduced, this approximation should begin to break down because of the nonzero radiative decay time of the flashlamp plasma²⁰ and because of the nonzero time required for the arc expansion to occur.

Overview of Experiments and Results.

We studied xenon-filled flashlamps with 1.5-cm-diam bore and 112-cm arc length at pulse lengths of $3\tau = 3(LC)^{1/2} = 122, 195, 315,$ and 635μ s. (Throughout this and the next article, "pulse length" means 3τ ; for a critically damped pulse, 3τ is approximately the full width of the input power at 10% of maximum.) Instantaneous measurements of overall (wavelength-integrated) lamp radiant output power and of the fractions of radiant output in the UV, the neodymium-pump region, and the IR were made throughout each pulse. Emission spectra in the pump region were obtained at peak radiant output. We determined overall radiant efficiencies (ratio of optical energy to electrical energy) and single-pass neodymium-pumping efficiencies (ratio of peak neodymium-fluorescence signal to electrical pulse energy), which are useful figures of merit for comparing flashlamps operating under different conditions. Tests were carried out over a wide range of input energies, with and without preionization pulses. Framing-camera photographs of the arc development were made, for possible correlation with lamp operating conditions.

Preionization significantly reduces the falloff in the radiant efficiency of flashlamps at short pulse length. Its major effect is to increase the initial arc size. We believe that preionization increases radiant efficiency largely by reducing arc contact with the lamp wall. Because the arc has less time to expand away from the wall for short pulses and for low-energy pulses, the effects of preionization are much stronger for those conditions.

Without preionization, pronounced xenon-ion line emission and (during the rise of the flashlamp pulse) greater UV emission occur. These differences, which are greatest for short pulses, are probably

of secondary significance with respect to neodymium-pumping efficiency. It is interesting to note, however, that they are greatly reduced by preionization.

By characterizing flashlamp emission spectra by radiant output power, rather than electrical input power, we can modify the QSS approximation to account for the nonzero radiation decay time of the plasma. The radiant output power is a better measure of the state of the plasma than the electrical input power. With this modification, the QSS approximation provides a good description of lamp spectra for pulse lengths longer than 100 μs when preionization is used, because of the effects of preionization on arc expansion. The modified-QSS results described in the next article may therefore be used to model flashlamp pumping of solid state lasers for all conditions currently of interest to us.

Experiments. The basic experimental apparatus has been described previously.^{16,21,22} The time-dependent voltage across the flashlamp was measured differentially with two Tektronix 6105 voltage probes. The lamp current was measured with a Pearson 301X Rogowski coil. All optical instruments were apertured to detect only light from a scatter plate packed with white BaSO_4 powder, which was used to angle-average the flashlamp output.¹⁶ The scatter plate was placed 23 cm from the centerline of the flashlamp, so that the maximum viewing angle was 68° from the lamp normal. The optical instrumentation included a Moletron P1-35 pyroelectric detector, a neodymium-fluorescence detector,²³ and a Tektronix J-20 spectrometer. The instrument outputs were attached to digitizing oscilloscopes, which were interfaced with an LSI-11/23 computer.

The pyroelectric detector, which has a flat spectral response, was calibrated absolutely to give total flashlamp radiant output power vs time. Calibration was carried out by comparing time-integrated signals from the detector with total radiant output energies measured absolutely using black-can calorimetry.²² Time-integrated pyroelectric and black-can measurements agreed with each other, within experimental errors, over wide ranges of conditions. We calculated radiant efficiencies by dividing the integrated pyroelectric detector signal by

the total electrical energy deposited in the lamp, including preionization pulse energy.

The pyroelectric detector was also used with Hoya L-42 and Corning 7-56 filters to transmit wavelengths longer than 0.4 and 1.0 μm , respectively, to determine the fractions of radiant output in the UV ($\lambda < 0.4 \mu\text{m}$), the neodymium-pump region ($0.4 < \lambda < 1.0 \mu\text{m}$), and the IR ($\lambda > 1.0 \mu\text{m}$). To calculate the radiant fractions, we divided the filtered pyroelectric detector signals by effective filter transmissions, which are the average transmissions of the lamp spectra throughout the filter passbands. The effective filter transmissions and radiant fractions at peak radiant output were calculated by convolving the measured filter transmissions with measured flashlamp spectra from 0.4 to 1.0 μm and with blackbody extensions of the spectra for $\lambda > 1.0 \mu\text{m}$.

To calculate the radiant fractions at times during the pulse for which there were no detailed spectral measurements, we used our modified QSS approximation and assumed that the effective filter transmissions, like the flashlamp spectra, are functions only of instantaneous radiant output power. The effective transmissions were then obtained from measurements made at peak radiant output power.

We measured flashlamp spectra with the Tektronix J-20 spectrum analyzer, which uses a silicon Vidicon detector attached to the output plane of a 0.25-m spectrometer. Appropriate filters eliminated second-order contributions. The J-20 was gated (25 μs FWHM) with a rotating-wheel shutter to obtain time-resolved spectra. Electronic circuitry attached to the shutter controlled the timing of lamp pulses so that "instantaneous" spectra could be obtained, typically at peak radiant output. The flashlamp spectra were corrected²² for the nonlinearity of the Vidicon response to light and for the wavelength-dependent sensitivity of the instrument and filters. The corrected spectra were scaled absolutely by using the pyroelectric-detector measurements of total radiant power and the radiant fraction in the $0.4 < \lambda < 1.0\text{-}\mu\text{m}$ band at the time of the shutter opening.

The electrical energy E delivered to a flashlamp is conveniently expressed²⁴ in terms of the explosion fraction $f_x = E/E_x$,

Flashlamps

where E_x is the input energy at which the flashlamp explodes. Empirically, E_x varies with pulse length as²⁵

$$E_x = 20 l d \tau^{1/2} \text{ kJ}, \quad (3)$$

where l and d are the flashlamp length and diameter in centimeters and $\tau = (LC)^{1/2}$. Although E_x may be considerably lower than given by Eq. (3) under abnormal conditions (e.g., for large-bore lamps at high fill pressures), Eq. (3) is reasonably accurate under normal conditions, and we use it to calculate f_x in any case. For fixed τ , Eq. (3) shows that f_x is simply a measure of input energy per unit area of the flashlamp bore.

As we varied the pulse length, we attempted to maintain critical damping in the flashlamp discharge circuit at approximately $f_x = 0.25$ by choosing L , C , and the charging voltage V according to the prescription²

$$L \sim \tau^{7/6}, \quad C \sim \tau^{5/6}, \quad V \sim \tau^{-1/6}. \quad (4)$$

We followed Eq. (4) for the three longest pulse lengths, starting with $C = 173 \mu\text{F}$ and $L = 257 \mu\text{H}$ at 635 μs . At 122 μs , a much larger capacitor was used to provide higher input energies at accessible voltages, and in this case critical damping occurred at $f_x = 0.07$.

The preionization pulse was provided by discharging a small, independently charged capacitor C_p through the main circuit elements L and C (see "Single-Segment Amplifier: Experiments," earlier in this section). Given the bank and preionization charging voltages V and V_p for capacitors C and C_p , the maximum

energy available in the preionization pulse is $(1/2)C_p(V + V_p)^2$. The efficiency of energy transfer from the preionization circuit to the lamp was typically 50%. The width of the preionization pulse power was approximately 50 μs FWHM. Most tests were conducted with $C_p = 1.32 \mu\text{F}$, but C_p was varied for the measurements of the effect of preionization energy on lamp radiant efficiency.

Results. We determined the overall radiant efficiency of the lamp as a function of the time delay between the initiation of preionization and the main pulse. For these measurements, the lamp was filled to 300 Torr and was fired at $f_x = 0.2$ using a 315- μs pulse. Preionization and main pulse energies per unit area were 0.15 and 12.1 J/cm^2 , respectively. The results, shown in Fig. 7-17(a), are normalized to the 68.5% radiant efficiency obtained under these conditions without preionization.

Radiant efficiency initially increases with time delay because more time is allowed for preionization electrons and ions to diffuse from the lamp wall. This is strongly suggested by framing-camera photographs, which show that the initial arc size (as indicated by the apparent fraction of the bore filled by the arc) increases as time delay is increased. The larger arc size and reduced contact of the arc with the wall increases lamp efficiency by reducing plasma recombination and thermal-conduction wall losses.²⁶

The decline in radiant efficiency at longer delay times probably reflects the decrease in preionization electron and ion densities because of volume recombination. The optimum time delay of 200 to 400 μs may depend on gas-fill pressure

Fig. 7-17. Xenon flashlamp radiant efficiency (a) vs delay time between preionization and main pulses and (b) vs preionization pulse energy.

and preionization energy, but this dependence has not been explored.

Similar experiments were carried out to determine the dependence of flashlamp radiant efficiency on preionization pulse energy. The results, again normalized to the efficiency obtained without preionization, are shown in Fig. 7-17(b). These data were taken at a time delay of 200 μ s using the same pulse length, input energy, and fill pressure as in Fig. 7-17(a). Most of the improvement in radiant efficiency is obtained with a preionization energy per unit area of 0.1 J/cm². Radiant efficiency saturates with preionization pulse energy, in part because the preionization energy is

included in the denominator for the radiant efficiency calculation. At the highest preionization energy plotted, 1.5% of the total input energy is contributed by preionization.

Radiant efficiency varies strongly with pulse length and explosion fraction. This is shown in Fig. 7-18, in which radiant efficiency is plotted against pulse length at explosion fractions of 0.1 and 0.3, with and without preionization. Preionization energy per unit area was usually 0.15 J/cm²; delay time was 200 μ s. The data in Fig. 7-18 are taken from measurements made over the range $0.02 < f_x < 0.4$ for each pulse length. Fill pressures were 300 Torr for the

Fig. 7-18. Radiant efficiency vs flashlamp pulse length $3\tau = 3(LC)^{1/2}$ for $f_x = 0.1$ and 0.3, with and without preionization. Arrows indicate improvement resulting from preionization.

Fig. 7-19. Neodymium-fluorescence pumping efficiency vs flashlamp pulse length $3(LC)^{1/2}$ for $f_x = 0.1$ and 0.3, with and without preionization. Arrows indicate improvement resulting from preionization.

635- and 315- μ s pulses and 150 Torr for the 195- and 122- μ s pulses.

Without preionization, radiant efficiency drops sharply as pulse length is reduced, especially at low explosion fractions. This is not surprising, since framing-camera photographs show that the arc does not expand very far from the lamp wall under those conditions. Preionization increases the radiant efficiency, especially at short pulse lengths. Framing-camera photographs show that preionized arcs have a much greater initial size, and thus reduced wall contact. Preionization gives the greatest efficiency improvement for low

Fig. 7-20. Radiant-fraction history for $3(LC)^{1/2} = 635 \mu$ s, with preionization, for UV, pump, and IR regions and for $f_x = 0.1, 0.2$, and 0.3 : (a) vs electrical input power per unit area and (b) vs radiant output power per unit area.

explosion fractions, even when the preionization pulse energy is included in the efficiency calculation. At higher explosion fractions, preionization has less effect on efficiency because the arc expands away from the lamp wall even without preionization. At $f_x = 0.1$, preionization increased radiant efficiencies by 4, 11, 18, and 24% for 635-, 315-, 195-, and 122- μ s pulses, respectively.

Figure 7-19 shows neodymium-pumping efficiency (ratio of peak neodymium-fluorescence signal to electrical pulse energy) vs pulse length for $f_x = 0.1$ and 0.3 . The efficiencies are normalized to the efficiency for 635- μ s pulses at $f_x = 0.2$, without preionization. The experimental conditions were the same as for the preceding figure.

With preionization, neodymium-pumping efficiencies increase strongly as pulse length is reduced. This improvement reflects the reduced fluorescence decay that takes place during pumping. Figure 7-18 shows that there are only slight differences in the radiant efficiencies at $f_x = 0.1$ and 0.3 with preionization at any of the tested pulse lengths. The differences in pumping efficiencies shown in Fig. 7-19 therefore arise almost entirely from blue-shifting of the lamp spectrum as lamp input is increased. As might be expected, blue-shifting becomes more severe for short pulse lengths, which correspond to higher peak powers. Without preionization, the benefit obtained from reduced fluorescence decay as pulse length is reduced is offset by a reduction in flashlamp radiant efficiency.

Radiant-Fraction Results. Here we report the results of our broad-band, time-resolved spectral measurements, which we made to test the QSS approximation. If the approximation were valid, the flashlamp spectrum would be a single-valued function of electrical input power. As an example of the deviation of actual spectra from this approximation, Fig. 7-20(a) shows the radiant fractions in the UV, pump, and IR regions vs electrical input power per unit bore area. The data were taken with 635- μ s pulses at $f_x = 0.1, 0.2$, and 0.3 , with preionization. Arrows on the curves indicate the direction of time. The curves show considerable hysteresis even at this relatively long pulse length.

Figure 7-20(b) shows the result of plotting radiant fractions vs radiant output power for data taken under the same conditions as in Fig. 7-20(a). Hysteresis is reduced, and there is good overlap of the spectral distributions for shots at different input energies.

We tested this modified QSS approximation at our four pulse lengths, with and without preionization, by plotting the time-resolved radiant fraction vs output radiant power per unit area. To the extent that the radiant fractions are single-valued functions of P_a , the plots will show no hysteresis, and the flashlamp spectrum will at all times be that predicted by the QSS approximation.

The differences between spectra on the rise and fall of the flashlamp pulse are considerable for short pulses without preionization, even if plotted against output power. Figure 7-21(a) shows radiant fraction vs output power for 122- μ s pulses at $f_x = 0.1$ and 0.2, without preionization. The hysteresis is considerable; the greatest departures from the results in Fig. 7-20(b) occur on the rise of the power curves.

Figure 7-21(b) shows that preionization greatly reduces hysteresis. The curves closely follow those obtained at peak output for long pulse lengths, described in the next article (see Fig. 7-27), particularly on the downside.

We used the long-pulse-length QSS results to estimate the error in predicting laser pumping efficiencies at various pulse lengths. The single-pass pumping rate scales approximately as the product of the neodymium-pump band radiant fraction and the total radiant output power. This product can be time-integrated over the pulse to give the total optical energy delivered in the neodymium-pump band. We calculated these pump-band energies using the measured time-dependent radiant fractions in the pump region and the QSS results at peak output for 635- μ s pulses. The energies agree within 3% for shots with preionization at all observed pulse lengths, and within 5% for shots without preionization.

This agreement between the observations and the QSS approximation for time-integrated pumping is better than one might expect. One reason for this is that

more energy is deposited during the fall of the pulse (for which the agreement with the QSS prediction is good) than during the rise (for which the agreement is poor). Another reason is that, for shots in which the fractional pump energy is below the QSS curve on the rise, it tends to be slightly higher on the fall. Given this close agreement, we expect the modified QSS results to provide a good description of laser

Fig. 7-21. Time-dependent radiant fraction vs radiant output power per unit area for $3(LC)^{1/2} = 122 \mu$ s, for UV, pump, and IR regions and for $f_x = 0.1$ and 0.2: (a) without preionization and (b) with preionization.

pumping for pulse lengths greater than $100\text{ }\mu\text{s}$. However, if the modified QSS description is to be used effectively, the flashlamp model must predict the radiant output power from the electrical input power. A method for doing this is described in "Flashlamp Modeling," later in this section.

More detailed J-20 flashlamp spectra, with and without preionization, are shown in Fig. 7-22. These spectra were taken with a $25\text{-}\mu\text{s}$ window centered at peak optical output. Two of the spectra were taken at a pulse length of $122\text{ }\mu\text{s}$ at $f_x = 0.2$, with and without preionization; the third was taken at a pulse length of $635\text{ }\mu\text{s}$ at $f_x = 0.29$, with preionization. These spectra are from shots in which the radiant output powers per unit area were as close to one another as we could find in our data.

Without preionization, the $122\text{-}\mu\text{s}$ spectrum shows many strong emission lines between 0.4 and $0.7\text{ }\mu\text{m}$. These are all Xe^+ lines, in contrast with the neutral-xenon lines observed between 0.7 and $1.0\text{ }\mu\text{m}$ at low input power. The Xe^+ lines become much weaker or disappear with preionization. Moreover, with preionization we see

good agreement with the spectrum taken at long pulse length at essentially the same output power per unit area. At the $635\text{-}\mu\text{s}$ pulse length there is very little difference between spectra with and without preionization. We conclude that detailed spectra at peak radiant output are characterized by total radiant output power, provided that long pulse lengths are used or that preionization is used at short pulse lengths.

Measurements of flashlamp spectra with the rapid-scan analyzer (see "Rapid Scan Analyzer," later in this section) showed that there are also UV emission lines on the rising side of flashlamp pulses without preionization. The UV emission features were measured using a Princeton Applied Research optical spectrum analyzer with 0.1-nm spectral resolution. The uncorrected spectrum is shown in Fig. 7-23. The UV lines emanate from high Rydberg states of Xe^+ ($n \approx 15$ to 35), which are $\sim 32\text{ eV}$ above the ground state of neutral xenon. Since the plasma temperature is $\sim 1\text{ eV}$ after the QSS is established, these states are initially far from thermal equilibrium. Figure 7-23 also shows a flashlamp

Fig. 7-22. Flashlamp emission spectrum at $3(LC)^{1/2} = 122\text{ }\mu\text{s}$ and $f_x = 0.2$, with and without preionization; and spectrum at $3(LC)^{1/2} = 635\text{ }\mu\text{s}$ and $f_x = 0.29$, with preionization.

Fig. 7-23. Flashlamp emission spectrum early in the pulse, with and without preionization.

spectrum taken at similar conditions but with preionization: the Xe^+ lines have disappeared.

We have a hypothesis to explain the appearance of the visible and UV Xe^+ lines. First we note that the momentum-transfer cross section for several-volt electrons is orders of magnitude greater for xenon ions than for xenon neutrals.²⁷ To a large degree the momentum-transfer cross section controls the electron energy, since the electric field between the lamp electrodes can decelerate an electron after a collision. During arc expansion one might expect that there would be a diffuse sheath on the arc boundary, where the ion density is low. The electron temperature would be high in this sheath, and high-energy states of the sort observed would be produced. The temperature of the heavy particles would remain low. When the arc expansion is complete, the plasma boundary between the arc and the unexcited gas presumably becomes much sharper, eliminating these hot electrons. In the arc the electrons and ions communicate strongly, as do the ions and neutrals, leading to local thermodynamic equilibrium (LTE) and

disappearance of the ion lines. Preionization suppresses the hot spectral lines because the sheath is no longer well defined and LTE is attained sooner. We have no explanation of why these specific Xe^+ lines are produced.

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Diameter Scaling

Introduction. The ability to predict laser amplifier performance will facilitate the design of ICF laser drivers. An important element of this ability will be a model from which xenon flashlamp emission and absorption spectra can be obtained. This model will be useful in many applications if it can be applied over broad ranges of flashlamp bore diameter, arc length, xenon fill pressure, input power, and pulse length. We are especially interested in extending the range of our existing model to include large bore diameters, because

substantial cost savings should result from using a (relatively) small number of large flashlamps rather than a larger number of small flashlamps in a 10-MJ laser.

To develop such a model, we require time-resolved, quantitative measurements of flashlamp emission spectra and of instantaneous radiant efficiencies (ratio of radiant output power to electrical input power). We assume that we can characterize flashlamp spectra using radiant output power. That is, we use the modified QSS approximation described in the preceding article.

Measurements of emission spectra and radiant efficiencies must be made at many different output powers. These measurements are made most accurately when the output power and the plasma state are changing slowly. Furthermore, as described in the next article, the instantaneous radiant efficiencies are determined when the time derivative of radiant output power is zero. These considerations dictate that our time-resolved measurements be made at the time of peak radiant output power.

The data should be as free as possible from transient effects or effects dependent on arc history, so the arc should not be expanding when the data are taken. Data should therefore be taken using long pulse lengths and optimized preionization. Under these conditions, radiant efficiencies and spectra at peak radiant output power will be representative of the QSS.

We studied the electrical and optical characteristics of xenon-filled flashlamps with diameters of 0.5, 1.5, and 4.2 cm. Measurements were made at pulse lengths of approximately $3\tau = 3(LC)^{1/2} = 635 \mu\text{s}$, and preionization pulses were applied before each main pulse. (As in the preceding article, we note that 3τ is identical, within a few percent, with the full width of the input power pulse at 10% of maximum, and that we always mean 3τ when we say pulse length.) Several xenon fill pressures were used: 300 and 750 Torr for the 0.5-cm-bore lamp; 100, 300, and 750 Torr for the 1.5-cm lamp; and 100 Torr for the 4.2-cm lamp.

For each pulse we measured voltage, current, and radiant output power vs time. By combining results from three pulses, we determined the lamp emission spectrum from

0.4 to $1.0 \mu\text{m}$ at peak radiant output power and the fractions of radiant output in the UV ($\lambda < 0.4 \mu\text{m}$), the neodymium-pump region ($0.4 < \lambda < 1.0 \mu\text{m}$), and the IR ($\lambda > 1.0 \mu\text{m}$). We calculated the overall radiant efficiencies (ratio of radiant energy to electrical energy), and measured Nd⁺ fluorescence efficiencies (ratio of peak Nd⁺ fluorescence to integrated electrical power).

Our results show that QSS flashlamp emission spectra are well characterized by the output power per unit bore area, as is characteristic of a surface radiator. However, as discussed in the next article, the emissivities inferred from these spectra are less than unity, so emission must come from within the lamp. We therefore describe flashlamps as gray bodies with emissivities and temperatures that somehow adjust with power and diameter to give this power-per-unit-area dependence. The integrated radiant efficiencies and neodymium-pumping efficiencies of the lamps are very similar when they are compared at equal input energy per unit area. The radiant efficiencies change slightly with pressure: best performance was obtained with xenon fill pressures of 750, 300, and 100 Torr for the 0.5-, 1.5-, and 4.2-cm-bore lamps, respectively. Line-emission features broaden slightly as pressure is increased.

Experimental. The experimental apparatus is described in the preceding article and elsewhere.²⁸⁻³⁰ We attempted to maintain the same pulse lengths and pulse shapes at common explosion fractions as lamp diameters were changed. This is particularly important for comparing neodymium-pumping efficiencies, because they are sensitive to changes in pulse length and pulse shape because of fluorescence decay during pumping. The prescriptions for varying circuit capacitance C , inductance L , and charging voltage V to maintain constant pulse length and pulse shape are³¹:

$$C \propto \frac{d^{5/3}}{lK_0^{4/3}},$$

$$L \propto C^{-1},$$

$$V \propto \frac{K_0^{2/3} l}{d^{1/3}},$$

(5)

where d is bore diameter, l is arc length, and $K_0 = 1.3(P/450)^{0.2}$; P is the xenon fill pressure in Torr. These prescriptions achieve their purpose reasonably well.

The lamps we tested were refillable and were constructed with envelopes of clear fused quartz and electrodes of sintered tungsten. The 0.5- and 1.5-cm lamps were supplied by ILC Technology, and had arc lengths of 112 cm. The 4.2-cm lamp was constructed in-house, using O-ring seals around the electrode fittings. Its arc length was 115 cm.

To optimize lamp efficiency, greater time delays between preionization and main pulses are required for larger bore lamps. We used delays of 100, 200, and 300 μ s for the 0.5-, 1.5-, and 4.2-cm lamps, respectively. Larger lamps also require more preionization pulse energy. We used preionization pulse energies *per unit bore area* of 0.19, 0.15, and 0.4 J/cm² for the 0.5-, 1.5-, and 4.2-cm lamps, respectively.

A scatter plate packed with highly reflective BaSO₄ powder was placed 23 cm from the centerline of the lamps to angle-average the flashlamp spectra. The principal optical instruments—a pyroelectric detector, a neodymium-fluorescence detector,³² and a Tektronix J-20 spectrum analyzer—were apertured to detect only light from this plate. The analysis of these data to obtain radiant efficiencies, broadband spectral distributions, and detailed spectral distributions in the pump region is discussed in the previous article.

The distance between the scatter plate and the lamps must be short to provide good angle-averaging of the emission spectra, but it was too short to permit accurate comparisons of the optical output of the different lamps. For such lamp-to-lamp comparisons, we used a second scatter plate 2.3 m from the lamp centerline, a distance many times greater than the largest bore diameter. Signals from the pyroelectric detector, apertured onto this scatter plate, were used to scale the radiant output of the different lamps relatively. The pyroelectric detector was calibrated absolutely by comparing its integrated signal with the radiant output energy measured using black-can calorimetry³⁰ at identical conditions of the 1.5-cm-bore lamp.

Results. Figure 7-24 shows some of the results for overall (wavelength-integrated) radiant efficiency, plotted vs explosion fraction f_x (which is proportional to input energy per unit area of lamp bore). The explosion energies at the 635- μ s pulse length correspond to electrical inputs of 16, 49, and 137 kJ for the 0.5-, 1.5-, and 4.2-cm lamps, respectively.

These data, which were taken with preionization, confirm the observation²⁹ that there is little difference in the overall performance of flashlamps at different diameters. The radiant efficiencies are low at low f_x for several reasons. The two most important reasons are that (1) as f_x is reduced, the preionization pulse energy, which is included in the radiant efficiency

Fig. 7-24. Time-integrated radiant efficiency vs f_x for 0.5-, 1.5-, and 4.2-cm-bore xenon-filled flashlamps at fill pressures of 750, 300, and 100 Torr, respectively.

Fig. 7-25. Neodymium-pumping efficiencies vs f_x for the lamps and xenon fill pressures of Fig. 7-24.

Fig. 7-26. Instantaneous radiant efficiency at peak radiant power vs radiant output power per unit area for the lamps and xenon fill pressures of Fig. 7-24.

calculation, becomes a larger fraction of total input energy, and that (2) the energy required to establish the arc becomes relatively more important as f_x is reduced (this is discussed in the next article). At high f_x , radiant efficiency falls slightly, presumably because a larger fraction of the plasma's radiant output is in the UV, where the lamp envelope absorbs ($\lambda < 0.2 \mu\text{m}$).

Neodymium-pumping efficiencies (ratio of peak neodymium-fluorescence signal to

electrical input energy) for the same lamp conditions are plotted in Fig. 7-25. These efficiencies are normalized to the value obtained at $f_x = 0.2$ using the 1.5-cm lamp with 300-Torr fill pressure, without preionization. The vertical scales in the figure are greatly expanded to permit detailed comparison of the different lamps; the change in neodymium pumping efficiency with loading is in fact quite small. These measurements confirm that there is no reason

Fig. 7-27. Radiant fractions at peak radiant power vs radiant output power per unit area for the lamps and xenon fill pressures of Fig. 7-24.

to prefer one lamp bore diameter over another when considering performance alone. The neodymium-pumping efficiencies depend on overall radiant efficiencies, pulse length, pulse shape, and the spectral distribution of lamp output. Some of the differences between lamps are due to a failure to exactly match circuit pulse lengths and pulse shapes. The fall in neodymium-pumping efficiency at high f_x is mostly due to blue-shifting of the lamp emission spectrum away from the neodymium-absorption bands.

Figure 7-26 shows instantaneous radiant efficiency at peak radiant output vs output power per unit area for the three bore diameters. To compare these efficiencies with the time-integrated radiant efficiencies in Fig. 7-24, we note that a peak output power of 35 kW/cm^2 corresponds to an explosion fraction of 0.2 at this pulse length. Although the curves in Figs. 7-24 and 7-26 have similar shapes, the instantaneous radiant efficiencies are higher at low input energies than the corresponding integrated radiant efficiencies. This difference arises from the inclusion of the preionization pulse energy in the integrated radiant efficiency, and from the investment of a relatively constant amount

of energy to establish the arc. The latter effect causes the lamp to radiate less efficiently early in the pulse, lowering its time-integrated radiant efficiency. At high input energies these effects are of less relative importance, and the instantaneous and integrated radiant efficiencies are in close agreement.

We now compare the spectral outputs of the 0.5-, 1.5-, and 4.2-cm lamps. A rough measure of emission spectra is provided by the fractions of radiant output in the UV, the neodymium-pump region, and the IR (as already described, the respective boundaries are at 0.4 and $1.0 \mu\text{m}$). As Fig. 7-27 shows, the results are essentially identical for all three lamps when the radiant fractions are plotted vs radiant output power per unit area; the curves would be quite different from one another if the data were plotted vs input power per unit volume. Figure 7-27 clearly shows the blue-shift of lamp emission with increasing power.

Figure 7-28 shows detailed emission spectra of the 0.5- and 4.2-cm lamps at two values of radiant output power per unit area. The J-20 spectrometer was used to obtain these spectra at peak radiant output power. The spectra are normalized to the

Flashlamps

Fig. 7-28. Spectral output distribution at radiant output powers of 19 and 44 kW/cm^2 for 0.5- and 4.2-cm-bore xenon-filled flashlamps at fill pressures of 750 and 100 Torr, respectively.

Fig. 7-29. Emission spectra for the 1.5-cm-bore lamp with 300-Torr fill pressure at three radiant output powers.

neodymium-pump band radiant fraction. The spectra are quite similar to one another, especially at high powers; the 1.5-cm-bore spectra (not shown) agree very well with those shown.

Figures 7-27 and 7-28 show that the flashlamp emission spectra are characterized by the radiant output power per unit bore area, which is consistent with the behavior of blackbodies. Lamps operating as blackbodies would have unit emissivities, and would act as surface radiators. Flashlamp spectra, however, are clearly not

given by blackbody functions, especially at short wavelengths. We do not have an explanation for this fact. However, it appears that we can describe flashlamps as gray bodies whose emissivities and temperatures change with power and diameter to give this power-per-unit-area dependence.

Figure 7-29 shows instantaneous emission spectra of the 1.5-cm bore lamp at three output powers. These plots show that the blue-shift of the emission spectra with increasing power is due to greatly increased emission at short wavelengths and

to saturation at long wavelengths. The line emission, which is prominent at low input energies, saturates to the smooth continuum emission. The prominent emission lines between 0.75 and 1.0 μm show no reversal at low input powers. At high input powers, for which the continuum emission dominates, these lines are sometimes observed as weak absorption features. These results suggest that the lamp plasma is a hot arc of relatively constant temperature surrounded by a colder region of gas with a relatively sharp boundary.

Conclusions. We have presented a sample of the large body of data we have accumulated for a wide range of lamp conditions. The instantaneous measurements of radiant efficiency and spectra, taken with preionization at long pulse lengths, are representative of the lamp emission at quasi-steady state conditions. These data will be useful for developing a model of flashlamp emission and amplifier performance. Instantaneous lamp spectra are well characterized by the radiant output energy per unit area of flashlamp bore. The integrated radiant efficiency and neodymium-pumping efficiency are described by f_x , or (equivalently) by the input energy per unit bore area.

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Flashlamp Modeling

Analysis of Quasi-Stationary Flashlamp Spectra. Spectrally resolved measurements of the peak radiative output from xenon flashlamps have been presented in past Laser Annuals.^{33,34} In 1985 we analyzed flashlamp spectra taken at an input-power pulse width of 600 μs (FWHM) (which gives quasi-stationary flashlamp behavior) to determine effective lamp temperatures and plasma absorption coefficients for bore diameters d from 0.55 to 4.2 cm, fill pressures p from 75 to 750 Torr, and input power loadings P/A from 5 to 100 kW/cm^2 . The temperatures and absorption coefficients inferred from our data exhibit a simple dependence on

Fig. 7-30. Xenon flashlamp output spectrum ($d = 1.5$ cm, $p = 300$ Torr, $P_{\text{in}} = 22$ kW/cm^2) and blackbody spectrum at the inferred plasma temperature T_{eff}

current density, diameter, and fill pressure that confirms our experimentally observed scaling of flashlamp optical power output with input power loading.

The first step in determining the emissivity was to determine the effective plasma temperature T_{eff} . In principle, this can be easily done using our absolute intensity measurements if the emission reaches its blackbody limit at some known wavelength (e.g., at the center of a strong, un-reversed spectral line). Rather than guessing the wavelength at which this occurs, we define an effective temperature T_λ at every observed wavelength through

$$I(\lambda) = I_{\text{bb}}(\lambda, T_\lambda), \quad (6)$$

where $I(\lambda)$ is the output intensity from the lamp and I_{bb} is the blackbody intensity distribution function.³⁵ We then assume that T_{eff} is given by the highest value of T_λ . Figure 7-30 shows a lamp spectrum $I(\lambda)$ at intermediate loading and the blackbody distribution $I_{\text{bb}}(\lambda, T_{\text{eff}})$, where T_{eff} was determined using this procedure. The observed intensity is below I_{bb} over most of the wavelength interval of interest, showing that the lamp's emissivity is generally less than unity.

At all but the highest lamp loadings, T_{eff} turns out to correspond to one of the intense neutral-xenon lines at 0.823, 0.828, and 0.882 μm ; in many cases (that of Fig. 7-30 is an example), T_λ peaks at approximately the same value at several wavelengths. For input energies greater than 30% of the explosion energy (explosion fraction $f_x \approx 0.3$), some degree of line reversal occurs where a colder region of the plasma outside the arc blocks the hotter emission from the inside. For these

spectra, the wavelength at which the maximum T_λ is found shifts to as low as $0.7 \mu\text{m}$.

Figure 7-31 shows values of T_{eff} for xenon flashlamps with bore diameters d of 0.55, 1.5, and 4.2 cm, fill pressures p of 75, 150, 300, and 750 Torr, and peak current densities j from 0.4 to 6 kA/cm^2 . T_{eff} depends very weakly on p . The straight lines are the result of a Livariate nonlinear least-squares fit of the data to a power-law dependence on j and d . The best-fit

Fig. 7-31. Xenon flash-lamp plasma temperature vs current density for various fill pressures and lamp diameters.

Fig. 7-32. Spectral absorption of xenon flashlamp ($d = 1.5 \text{ cm}$, $p = 300 \text{ Torr}$) at loadings differing by a factor of 2.

dependence was

$$T_{\text{eff}} \propto j^{(0.21 \pm 0.01)} d^{(0.12 \pm 0.01)} \quad (7)$$

As we explain in detail below, this result indicates that T_{eff} scales with output optical power per unit area P_{out}/A , and not with power per unit length or per unit volume.

With this estimate for the plasma temperature, we calculated the effective angle-averaged emissivity at each wavelength by dividing the observed $I(\lambda)$ by $I_{\text{bb}}(\lambda, T_{\text{eff}})$. We then obtained the plasma absorption coefficient $\alpha(\lambda)$ by inverting the analytically derived relation between α and the average emissivity for cylindrical radiators.³⁶ Figure 7-32 shows the results of this calculation for a 1.5-cm-i.d., 300-Torr xenon flashlamp at input loadings of 22 and 44 kW/cm^2 . Although the absorption coefficient at a given wavelength varies strongly with lamp loading, the dependence of α on wavelength is relatively insensitive to loading.

To determine the dependence of the inferred value of $\alpha(\lambda)$ on current density, lamp diameter, and fill pressure, we examined the variations of α at 0.4, 0.6, and $0.8 \mu\text{m}$ in greater detail. We expect the

behavior of the absorption at these wavelengths to be characteristic of that of the continuum absorption in general, since they are far from any prominent line-emission features. At each wavelength, to minimize the effect of small-scale noise in the experimental spectra, we determined an effective absorption coefficient by averaging values in an $0.02\text{-}\mu\text{m}$ range about the central point.

The absorption coefficients thus determined for 0.55- and 1.5-cm -bore flashlamps at a fixed xenon fill pressure of 750 Torr are essentially independent of lamp diameter. (Recall that the temperature depends on both lamp diameter and current density.) At all three wavelengths, the absorption coefficient varies approximately as $j^{1.5}$.

The variation of α with p and j for $d = 1.5$ cm is plotted in Fig. 7-33. At each wavelength the data were least-squares fitted to a simple power-law dependence on j and p . The best-fit exponents and the resulting curves are shown in Fig. 7-33. In all three cases the absorption coefficients scale roughly according to the power law

$$\alpha(\lambda) \propto j^{1.5} p^{0.45}. \quad (8)$$

The dependence of α on j in Eq. (8) is somewhat weaker than the j^2 variation in αd reported by Emmett, Schawlow, and Weinberg.³⁷ In the present empirical model, our estimate of the plasma temperature, and hence our estimate of the product αd , depends on the value assumed for the effective arc diameter; we have assumed that it is equal to lamp's inner wall diameter. As noted in the introduction to this subsection, however, the arc never completely fills the bore, and the extent to which the lamp volume is filled increases with lamp loading. Because of this effect, one would expect αd to vary faster than our $j^{1.5}$, and to be closer to the j^2 of Ref. 37. The expansion of the arc may also be the cause of the threshold behavior of the absorption coefficient reported by Linford³⁸, the focused probe beam used in that experiment may have missed the smaller arcs formed at the lowest levels of lamp excitation. Our work exhibited no such threshold behavior.

The functional dependences of the temperature and absorption of the flashlamp plasma demonstrated above confirm the experimental observation that quasi-stationary lamp spectra scale with lamp input power per unit area P/A . Using the

Fig. 7-33. Xenon flashlamp absorption coefficient (lamp i.d. = 1.5 cm) vs current density for various xenon fill pressures.

empirically determined variation³⁹ of the lamp impedance parameter with j , p , and d , we find that

$$(P/A) \propto j^{1.5} p^{0.2} d. \quad (9)$$

Combining this with our power-law scalings for the temperature and absorption coefficient, we obtain the relations

$$T_{\text{eff}} \simeq 2240(P/A)^{0.14}, \quad (10)$$

where T_{eff} is in kelvins and P/A is in kW/cm^2 , and

$$\alpha(\lambda)d \propto j^{1.5} p^{0.45} d \propto (P/A)p^{0.25}. \quad (11)$$

Elementary emission physics for a semi-opaque radiator such as a flashlamp requires that the radiative output at each wavelength be given by the product of the blackbody intensity at the appropriate temperature and the angle-averaged emissivity, which is completely determined by $\alpha(\lambda)d$. Equations (10) and (11) indicate that T_{eff} and αd , and therefore the emission spectrum itself, may be represented approximately by a function of the lamp power per unit area, with a weak dependence on the fill pressure.

A true blackbody exhibits a similar scaling with P/A . However, in that case the temperature varies as $(P/A)^{0.25}$. The exponent for the flashlamp temperature is lower because the lamp emissivity increases with loading. It is clear from Fig. 7-30 that the lamps are far from totally opaque when the P/A scaling holds. We do not have a satisfactory physical explanation for this flashlamp behavior.

Transient Effects. As described earlier,^{40,41} the simple quasi-stationary model that we previously used for describing the

time dependence of the radiative output from xenon flashlamps was limited to the assumption of a constant time delay, generally about $30 \mu\text{s}$, between the input and output powers. While this prescription resulted in a good match between calculation and experiment under our standard operating conditions⁴⁰ (pulse width $\simeq 600 \mu\text{s}$, explosion fraction $f_x \simeq 0.2$), it fails to predict accurately the delay and shape of output pulses at shorter input pulse widths or at smaller loadings.

We have recently found (see "Pulse Length Scaling," earlier in this section) that the output spectra of flashlamps are better described as functions of the instantaneous optical output power P_{out} rather than the electrical input power P_{in} . We have developed a method to predict P_{out} from P_{in} , thus allowing us to modify the quasi-stationary model to predict the output history of a 1.5-cm-i.d. xenon lamp output over the entire range of lamp loadings and for pulse widths down to $100 \mu\text{s}$ (FWHM).

We begin by relating $P_{\text{out}}(t)$ and $P_{\text{in}}(t)$ through the differential equation

$$\dot{P}_{\text{out}} = \frac{1}{\tau(P_{\text{out}})} [\eta(P_{\text{out}})P_{\text{in}}(t) - P_{\text{out}}(t)], \quad (12)$$

where η is the instantaneous lamp radiative efficiency and the response time τ is the time required for the flashlamp to lose (primarily by radiation) its stored energy. This equation follows from the conservation of energy. We assume that the efficiency and the response time are both functions of $P_{\text{out}}(t)$ rather than of $P_{\text{in}}(t)$. This is in agreement with the experimental results described above, where it was found that the radiative characteristics of the lamp describe, to first order, a single-valued function when plotted vs P_{out} .

The instantaneous radiative efficiencies for 1.5-cm-i.d. flashlamps at a xenon fill pressure of 300 Torr were determined using Eq. (12) and the observed values of P_{in} and P_{out} at the time of peak output ($P_{\text{out}} = 0$). Figure 7-34 shows the results of this calculation and a fourth-order polynomial fit to the data. We then integrated Eq. (12), assuming a fixed response time of $30 \mu\text{s}$ and using the experimental input-power curves for the 1.5-cm-i.d. lamp, and

Fig. 7-34. Xenon flashlamp radiative efficiency ($d = 1.5 \text{ cm}$, $p = 300 \text{ Torr}$) vs output optical power per unit area.

Fig. 7-35. (a) Xenon flashlamp output ($f_x = 0.073$) and prediction of old flashlamp model. (b) Xenon flashlamp output and prediction of new flashlamp model, which includes an efficiency term that depends on lamp input energy.

compared the resulting predictions for $P_{out}(t)$ with the observed values.

For long pulse widths and for explosion fractions above 0.1, the agreement between the model and experiment is excellent, as one might expect. However, as illustrated in Fig. 7-35(a), the predictions at lower energies overestimate the output power at early times. This disagreement indicates the need to account for the energy investment necessary to establish the arc, which reduces the effective radiative efficiency early in the pulse.

We corrected for this effect empirically by replacing the term $\eta(P_{out})$ in Eq. (12) with an instantaneous efficiency η' that accounts for the energy deposited in the lamp up to a given time:

$$\eta'[P_{out}, E(t)] = [1 - C \exp\{-BE(t)\}] \eta(P_{out}), \quad (13)$$

where $E(t)$ is the integral of $P_{in}(t)/A$ up to time t . The agreement between the model for long times and large input energies is maintained with this formulation, since at late times (for which the deposited energy is much greater than the saturation energy B^{-1}) the instantaneous efficiency η' goes to its quasi-stationary value η .

We varied the coefficients B and C to obtain a best fit of the ratio of predicted and observed output powers for pulse widths from 200 to 600 μ s and for explosion fractions from 0.04 to 0.37. Figure 7-35(b) gives a typical example of the results of this calculation for the same conditions as in Fig. 7-35(a). With $B^{-1} = 0.354$ J/cm² and $C = 0.636$, agreement between the calculation and experiment is extremely good at all times and input energies, for pulse widths as low as 100 μ s. We have not yet

determined how B and C vary with bore diameter or preionization energy.

One use of this model has been to determine an improved estimate for the single-pass pumping efficiency of the single-sided amplifier, as is described in "Single-Pass Pumping," earlier in this section.

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Rapid Scan Analyzer

Our previous measurements of flashlamp optical emission have been either spectrally integrated and time-resolved or spectrally resolved and time-sliced.⁴² We have designed and built a system that provides both spectral and temporal resolution of the entire flashlamp emission on a single shot. This system, called the rapid scan analyzer (RSA), operates as a stand-alone instrument requiring only a zero-time trigger to synchronize it with an external event. Data are digitized and stored locally during the shot, and are retrieved, calibrated, interpolated, and merged to produce the spectral-temporal history.

The RSA, shown in block diagram in Fig. 7-36, achieves a 0.2- to 1.72- μ m spectral range (previous measurements covered 0.4 to 1.0 μ m) by utilizing four independent spectrometer subsystems. Spectral overlap between successive spectrometers allows us to merge the data to produce a composite spectrum with a uniform relative intensity scale. Table 7-2 summarizes the optical and temporal characteristics of each spectrometer subsystem.

As in previous flashlamp experiments, we use a BaSO₄ scattering plate to obtain

Fig. 7-36. Block diagram of the RSA.

Flashlamps

Table 7-2. Optical and temporal characteristics of the spectrometers in the rapid scan analyzer system.

Characteristic	Spectrometer number			
	1	2	3	4
Wavelength range (μm)	0.2-0.4	0.3-0.6	0.5-1.0	0.85-1.72
Number of useful detectors	37	37	37	11
Wavelength resolution (nm/el)	5.4	5.4	10.8	65
Total recording time (ms)	2	2	2	1.1
Temporal resolution (μs /temporal sample point)	18	18	18	10.5
Amplitude resolution	1:512	1:512	1:512	1:512

an optical angle-average of the flashlamp emission. Each spectrometer is independently aligned to view the center of the scattering plate. Long-pass filters prevent second-order light from reaching the longer-wavelength detectors. We control intensity by adjusting an iris in front of the spectrometer and by inserting neutral-density filters.

Three Reticon 64-element, self-scanned, linear silicon photodiode arrays record the spectra over the 0.2- to 1.0- μm range. Each array is scanned (read out) more than 100 times during an optical event. A single scan takes 18 μs . The array output is a train of pulses, each proportional to the integrated optical intensity during the scan period. Clear-time markers separate individual scans and allow us to associate a signal pulse with a specific detector. We integrate each pulse electronically before digitizing at 4 MHz.

For the 0.85- to 1.72- μm range, we designed a 16-element, discrete, linear germanium photodiode array. The germanium signals are individually amplified and then time-multiplexed onto a single line before being digitized. This digitizer operates at 1.8 MHz, giving a single-scan time for the array of 10.5 μs .

After the data are acquired, they are calibrated to account for the spectral response of the instruments, spliced together over the four spectral regions, and interpolated to give the instantaneous spectra at 10- μs intervals. We acquire more than 112 000 bytes of raw data per shot. We archive two data files on disk: raw and time-interpolated (calibrated) data.

Our CAMAC-interfaced transient digitizers, along with the analog signal-conditioning circuitry and spectrometers, are in the experimental laboratory near the flashlamp. The control system, consisting of an LSI-11/23 microprocessor and standard peripherals, is in a separate room. The two subsystems communicate through a fiber-optic CAMAC serial highway, chosen for its noise immunity and reliability.

The RT-11 operating system and the Fortran language form the backbone of our

application code development. We have written user-friendly, menu-driven codes for system calibration, data acquisition, and data analysis and display. The codes are large, so they are chained rather than overlaid to assure reliable operation. The codes allow us to select options for control, parameter setting, and analysis. We are able, for example, to plot an individual detector's output vs time or to plot the spectrum vs wavelength for an individual array or for all arrays. Both raw and calibrated data may be viewed.

Two spectral calibration techniques are used. This is dictated by differences in sensitivity and sampling rate between arrays and by the limited output intensity of standard calibration lamps. Between 0.5 and 1.72 μm , tungsten-filament sources provide adequate signal intensities to calibrate two of the spectrometers (Nos. 3 and 4 in Fig. 7-36).

Deuterium lamps do not have enough intensity to calibrate the system in the ultraviolet (0.2 to 0.4 μm), so we calibrate the two short-wavelength spectrometers (Nos. 1 and 2 in Fig. 7-36) using the flashlamp itself as the light source and using a gated optical multichannel analyzer (OMA) as a secondary standard. The OMA is a commercial device that utilizes a sensitive, intensified photodiode array but that provides only a single, time-sliced spectrum on each shot. First, we calibrate the OMA in the 0.2- to 0.8- μm range using the deuterium and tungsten lamps. The raw calibration data (typically an average of many scans), and the manufacturer's data on the lamp's spectral output, are analyzed on an off-line computer. We fit the manufacturer's data using a spline routine and from this curve calculate the output of the standard lamp at the spectral points available in the raw OMA data. These two sets of data are then combined to produce an OMA system calibration factor. Both instruments then simultaneously record the flashlamp emission in the same spectral region. We gate the OMA on for 2 μs at the peak of the light output. The raw flashlamp spectrum from the OMA is transmitted over a serial communications link to the RSA, where it is corrected using the OMA system calibration factor to give the spectral output of the flashlamp. The measured output of the flashlamp, and RSA

data sampled at the same instant, are combined to produce an RSA system calibration factor in the 0.2- to 0.6- μm range.

Preparation of the final data set of flashlamp output vs time requires several steps. First, the RSA system measures and subtracts electronic noise from the raw data. Scattered light in each spectrometer represents a significant error, which is corrected in the raw data with the following technique. The spectral overlap between arrays results in duplicated detectors, some of which lie below the cut-on wavelengths of the filters. Signals measured in these detectors below the filter cut-on represent scattered light. These signals are averaged on a scan-by-scan basis and subtracted from all detectors in the raw array. At this point we apply the calibration factor to each array element at each time increment.

Because of the differences in calibration, data from each array must be scaled to produce a composite spectrum. We reference each array to the known zero time and then choose a time (typically the temporal peak) at which to calculate a set of scale factors. Linear interpolation is used to obtain values at the selected time. We then choose a wavelength in the overlap region of each spectrometer pair and linearly interpolate to the chosen wavelength, if required. Scale factors computed at these wavelengths are then applied to all elements in the appropriate arrays. Once scaling is complete, we merge the data from the four arrays to obtain a composite spectrum at the scaling wavelengths.

Each digitizer runs independently of the others, even though accurate timing is distributed to all subsystems. Consequently, data samples from each array are not acquired simultaneously. Differences in array lengths and clock rates further time-skew the data between arrays 1, 2, and 3 on the one hand, and array 4 on the other. Using zero time as our absolute reference between arrays, we calculate, via linear interpolation, a value for all detectors at each selected time increment, typically every 10 μs . The final, time-synchronized data file consists of a 1.1-ms time history at 10- μs intervals, consisting of approximately 120 spectral points at each time increment.

Figure 7-37 shows the spectrally and temporally resolved emission history of a

Fig. 7-37. Spectral and temporal history of a xenon flashlamp driven at 0.28 explosion fraction without preionization, as measured by the RSA on a single shot.

Fig. 7-38. Three spectra from the shot of Fig. 7-37.

xenon flashlamp from a single shot, obtained with the RSA. The data are plotted at 20- μ s intervals. The flashlamp was operated at an explosion fraction of 0.28 and a 37-pulse length of 600 μ s. Figure 7-38 shows the spectra at peak light output and at the half-intensity points on the rise and fall. Figures 7-37 and 7-38 represent but two of the several ways in which we can display the data.

Figure 7-38 shows strong line emission near 0.25 and 0.5 μ m on the rise of the light flash but not on the fall. More detailed measurements using the OMA show that these emissions are xenon ion lines as opposed to xenon neutral lines, which are commonly observed in the near-infrared. The presence of the ion lines has been shown to be a function of whether preionization of the flashlamp is used: see "Pulse

Length Scaling," earlier in this section. Such differences are of interest for modeling the time-dependent pumping rate of lasers. We have used the RSA in studying reflectively coated flashlamps: see the next article.

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Reflective-Coated Flashlamps

Introduction. To the degree that xenon flashlamps are simply hot radiating gases in local thermodynamic equilibrium, one might expect that the absorption of optical power would increase light output exactly like a similar increment in the input electrical power. Despite the simplicity of this argument, there has been considerable uncertainty for many years over whether such optical recycling in flashlamps takes place. Only in the USSR has this issue been carefully investigated, and the results⁴³ of these experimental and theoretical investigations have not been confirmed in the West.

It is important to understand this issue for modeling the pumping of solid state lasers by flashlamps, in which (by design

or by accident) some optical output is usually reabsorbed by the lamps, which are often quite black for conditions of interest. Moreover, it is clear that the laser-pumping efficiency of flashlamps could be improved by reflective coatings that directly feed back to the plasma some of the energy that is normally unused in pumping. Estimates have been made⁴⁴ of the Nd:glass pumping efficiency that might be achieved with such an optically conserving system, assuming that the effect of the recycled optical energy is equivalent to an identical increase in the electrical input.

Figure 7-39 illustrates the regions of spectral output of a clear fused-quartz xenon flashlamp that are amenable to recycling when pumping Nd:glass. The flashlamp spectrum, obtained with the rapid scan analyzer (see the preceding article) represents the instantaneous spectrum at peak output. Radiation below 350 nm does not couple to the Nd³⁺ absorption bands. In addition, because the reflectivity of the silver platings in our laser cavities falls off strongly below 400 nm, and because the generally used cerium-doped quartz envelopes also absorb below 370 nm, pump radiation at wavelengths shorter than 400 nm is generally unused and lost. These facts suggest that recycling flashlamp radiation below 400 to 450 nm could

Fig. 7-39. Typical flashlamp output spectrum (1.5 cm bore, 300 Torr xenon) and Nd:glass absorption, illustrating spectral regions suitable for optical recycling.

significantly increase the pumping efficiency. There is radiation above 900 nm that could be beneficially returned to the lamp plasma as well.

During 1985, we primarily investigated the effects of recycling UV radiation. This has the greatest benefit when the lamps are operated at relatively high input powers, as is the case in pumping Nd:glass fusion lasers. We demonstrated increases of 12 to 13% in peak neodymium fluorescence for reflective-coated lamps for input explosion fractions between 0.1 and 0.2 at a pulse length $3(LC)^{1/2}$ of 600 μ s (Ref. 45). We used pyroelectric detector data to determine a lower bound on the equivalency factor between optical and electrical input to a pulsed xenon flashlamp operating at quasi-stationary conditions (see "Pulse Length Scaling," earlier in this section). We find that UV input is at least 0.75 times as effective as electrical input, and it is likely that this factor is in fact unity.

Predicted Pumping Enhancement with Recycling. The measured output spectra of xenon flashlamps discussed in previous articles in this subsection were used to predict the possible improvement in pumping efficiency using optical recycling. We assumed that the exterior surface of the lamp is covered with an ideal coating that reflects perfectly all wavelengths below some cutoff wavelength λ_c and transmits perfectly all longer wavelengths. We also assumed that the optical power at wavelengths below λ_c is absorbed by the plasma and is equivalent to electrical power.

We denote the normalized spectral distribution of lamp output by $f(P_T, \lambda)$, where P_T is the total power into the lamp, and the instantaneous radiant efficiency by $\eta_{rad}(P_T)$. The total lamp input P_T is then the sum of electrical and optical inputs,

$$P_T = P_e + P_T \eta_{rad}(P_T) f_{rec}(P_T), \quad (14)$$

where P_e is the electrical input, and f_{rec} is the fraction of light output that is optically recycled, given by

$$f_{rec}(P_T) = \int_0^{\lambda_c} f(P_T, \lambda) d\lambda. \quad (15)$$

The instantaneous pumping efficiency with recycling is then given by

$$\eta'_{pump}(P_T) = \frac{P_T}{P_e} \eta_{rad}(P_T) \times \int_{\lambda_c}^{\infty} f(P_T, \lambda) [1 - \exp[-\sigma(\lambda) \rho w]] \frac{\lambda}{\lambda_0} d\lambda, \quad (16)$$

where $\sigma(\lambda)$ is the absorption cross section, ρ is the doping density, w is the sample thickness, and λ_0 is the operating wavelength of the laser.

The simultaneous Eqs. (14) to (16) are solved straightforwardly by calculating the electrical power P_e associated with an assumed total power P_T for a given lamp coating. We did this using our experimentally determined radiant efficiencies and spectral distributions. Convolving the lamp spectrum with the Nd:glass absorption at $\rho w = 4 \times 10^{20}$ ions/cm² yields the results shown in Fig. 7-40 for the enhancement in Nd* pumping efficiency over the case with no lamp feedback. The lowest cutoff wavelength used for the calculation was 400 nm, because we do not have detailed spectra at shorter wavelengths. The curves in Fig. 7-40 have the same general shape as the pumping efficiency curves without feedback, but they roll off more and more quickly with input as λ_c is increased because of the greater lamp blue shifts. The increase in the instantaneous pumping efficiency is predicted to be as high as 30% for an ideal, optimally chosen coating.

Coating Specifications. Multilayer dielectric coatings were deposited on prefabricated flashlamps and on tubes for which O-ring-sealed electrodes could be installed to make a flashlamp. All the envelopes had a bore diameter of 1.9 cm and an outside diameter of 2.4 cm. The coatings covered the full arc length of 47 cm. The coatings were fabricated by Optical Coating Laboratory, Inc. (OCLI) and by Spectra Physics under contract to LLNL. To obtain uniformly reflecting coatings over the entire lamp surface, the vendors had to pay careful attention to source masking and to tube alignment as the lamps were rotated about their axes behind a narrow slit during coating. For the second batch of

Fig. 7-40. Calculated instantaneous enhancement in Nd^{3+} pumping efficiency of LG-660 laser glass due to UV recycling in the flashlamp.

Fig. 7-41. Normal-incidence transmission vs wavelength for the two UV-reflective flashlamp coatings.

coatings we utilized precision ground, precision bore tubes to minimize wobble and help obtain a more uniform coating. Reflective stacks were obtained using alternating hafnia and silica layers. Hafnia was chosen as the high-index material because of its low absorption in the UV down to about 220 nm.

Figure 7-41 shows examples of the transmission at normal incidence of two types

of coatings made by Spectra Physics. The curves were obtained with a conventional spectrometer using a sliced section of a coated tube, with a similar slice from an uncoated section in the reference arm. Difficulty with balance between the two arms permitted us to make measurements only at normal incidence using this technique.

The specifications for coating 1 were that it have a reflectivity greater than 90%

between 250 and 400 nm, and an average transmission greater than 92% between 450 and 900 nm. The coating substantially met this specification, as did similar coatings made by OCLI. The specifications for coating 2 (fabricated only by Spectra Physics) were similar, except that the transition region was shifted to longer wavelength by 50 nm. Hence, the cutoff wavelengths for coatings 1 and 2 were approximately 400 and 450 nm, respectively.

We obtained five samples of each coating, although for various reasons (described below) not all samples could be used. It is important to recognize that the reflectance edge shifts to shorter wavelengths as the angle of incidence of the light is increased (e.g., there is a 20-nm shift at 45°). There are also differences in reflectivity between the s and p polarizations. Because flashlamps emit over a broad angular distribution, there is considerable averaging of the coating properties.

Testing Procedure. Identical coated and uncoated flashlamps were tested for electrical and optical performance. We used a bank circuit consisting of a 750- μ F capacitor, charged to 2 to 5 kV and discharged through a 56- μ H inductor by a size A ignitron. A preionization pulse was used to initiate discharge in the xenon. Voltage across and current through the lamp were recorded digitally.

We used the scatter-plate (powdered BaSO_4) method to obtain an angle-averaged sample of the light emitted by each lamp. Angle-averaging is clearly required, given that both the coating and lamp properties vary with angle. The

optical diagnostics were apertured to view only the scatter plate. A peak fluorescence detector (PFD) was constructed of a 2-cm-thick piece of 2%-neodymium-doped LG-750 glass. To simulate the pumping conditions in our laser cavities, a 400-nm long-pass filter was generally placed in front of the PFD to prevent pumping by UV light. A spectrally flat pyroelectric detector, with various filters, also viewed the lamp emission. We also used the RSA, and in some cases we used a gated optical multichannel analyzer (Princeton Applied Research 1420 Reticon system).

Because the differences in output between the coated and uncoated lamps were relatively small, it was necessary to pay attention to small details to obtain quantitative results repeatable to about 1%. For example, the position of the grounded trigger wire relative to the scatter plate affected signal levels slightly and was held constant for the measurements. Azimuthal asymmetry in the lamp coating was particularly troublesome, since the scatter-plate method implicitly assumes that the lamp emission is independent of tube rotation. Some samples of coating 1 had sides that were reflective well into the visible, as noted by visual observation. If we placed these sides opposite the scatter plate, we observed apparent PFD enhancements ($\sim 10\%$) that were caused by a simple mirror effect. The lamps were always rotated so that the average emission was observed. Coating 2 was sufficiently uniform that careful positioning was unnecessary.

Results. Figure 7-42 shows RSA spectra for a coated (coating 1) and an uncoated

Fig. 7-42. Flashlamp output spectra for coated and uncoated lamps at moderate and high input powers.

lamp for two values of the instantaneous peak input power. The UV emission from the coated lamp is depressed with respect to that from the uncoated lamp because of the reflection by the coating. The coated lamp shows greater output at longer wavelength because of this optical feedback, clearly demonstrating that significant optical reemission occurs. The increase in output at long wavelengths is much greater at high input powers, presumably because of greater UV emission and greater plasma absorption.

Figure 7-43 shows the enhancement in peak neodymium fluorescence between coated and uncoated flashlamps, as indicated by the PFD signals. The data for coating 1 represent an average, over several runs, of results for five uncoated and five coated lamps fabricated by the two vendors. The data for coating 2 were obtained with a single uncoated and a single coated tube, which were precision-ground and identical in size; the error bars for this case represent only shot-to-shot statistics. None of the lamps with coating 1 was optically damaged after many shots at explosion fractions of up to 0.3. Because of its greater number of layers, coating 2 was slowly damaged at the highest inputs, limiting data collection.

The differences in performance between coatings 1 and 2 follow the general trends shown in Fig. 7-40. The 450-nm coating (coating 2) produces a greater enhancement, which peaks at lower input. The PFD enhancement factor is not as great as the predicted instantaneous enhancement in the Nd* pumping efficiency (at $f_x = 0.2$, the peak input is 40 kW/cm^2). This presumably occurs because there is averaging over time, because the coatings do not reflect over the full wavelength range (so that some light is lost to coating absorption), and because some light is lost by reflection out the ends of the lamp.

Recycling Efficiency. The comparisons between coated and uncoated lamps just described were obtained at fixed electrical inputs. A different scheme, using differing electrical inputs for the coated and uncoated lamps, was utilized to determine the optical recycling efficiency. One might expect that the same long-wavelength

Fig. 7-43. Measured enhancement in peak Nd* fluorescence produced by flashlamps with UV-reflective coatings.

output could be obtained from an uncoated lamp as from a coated lamp if the electrical input to the uncoated lamp were slightly higher than that of the coated lamp. Presumably, the plasma conditions would then be identical. We define the optical/electrical equivalency factor as

$$\chi = \frac{\Delta P_{\text{elec}}}{\Delta P_{\text{opt}}}, \quad (17)$$

where ΔP_{opt} is absorbed optical power and ΔP_{elec} is the increment in electrical power required to obtain the same plasma conditions, as indicated by long-wavelength output.

This argument provides a simple strategy for determining the recycling efficiency. Using a spectrally flat pyroelectric detector, filtered with a long-pass filter to see only long-wavelength output, we increased the electrical input to the uncoated lamp so as to produce the same instantaneous signal at peak output as the coated lamp. (Since it was difficult to obtain exactly the desired output with the uncoated lamp, interpolation was generally used to determine the value of ΔP_{elec} required to achieve equivalent conditions.) An upper limit on ΔP_{opt} was then obtained by comparing the total outputs of the coated and uncoated lamps for the equivalent condition. A signal S , proportional to this total optical output, was obtained by removing the long-pass filter on the pyroelectric detector. It is easy to show that a lower bound on the recycling efficiency is given by

$$\chi = \frac{1 - P_{\text{coated}}/P_{\text{unc}}}{1 - S_{\text{coated}}/S_{\text{unc}}} \frac{1}{\eta(P_{\text{unc}})}, \quad (18)$$

Fig. 7-44. Optical/electrical equivalency factor inferred from coated-flashlamp experiments.

where P_{coated} and P_{unc} are the electrical input powers at equivalent conditions for the coated and uncoated lamps and S_{coated} and S_{unc} are the corresponding unfiltered pyroelectric detector signals. Measurements of $\eta(P_{\text{unc}})$, the instantaneous radiant efficiency at the equivalent electrical power, are discussed in "Pulse Length Scaling," earlier in this section. Equation (18) clearly gives a lower limit on the recycling efficiency, since some of the missing power may be lost to sinks other than plasma absorption.

Figure 7-44 shows the results of this analysis as applied to lamp data with coatings 1 and 2. Both sets of data behave similarly. At low input power the apparent recycling efficiency is low; at high input powers, for which the plasma becomes more opaque, the apparent recycling efficiency becomes as high as 0.8. We presume that at low inputs, for which the plasma is quite transparent in the UV, the missing light is lost mainly to sinks other than plasma absorption (presumably end leakage and coating absorption). The recycling efficiency could vary with loading as indicated. However, since our method gives a lower limit on the optical/electrical equivalency factor, our results are not inconsistent with the view that χ is unity for all conditions.

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Optical Materials and Coatings

Introduction

Our work on optical materials and coatings in 1985 had four main objectives:

- To complete experimental measurements and analyses on the effect of transition-metal impurities on absorption and on neodymium fluorescence lifetime in silicate and phosphate laser glasses.
- To improve our understanding of the influence of neodymium concentration on fluorescence lifetime.
- To develop new materials and applications for sol-gel coatings.

- To continue measurements on the damage thresholds of laser materials, nonlinear optical materials, and coatings.

Transition-metal impurities in silicate and phosphate laser glasses have broad absorption bands whose absorption at the neodymium laser wavelength decreases extraction efficiency. Such impurities may also induce nonradiative transitions from the neodymium excited state, reducing the excited-state lifetime and thereby reducing laser efficiency. Our experiments show that control of absorption requires a more stringent purity requirement than does control of fluorescence lifetime. To keep the absorption loss coefficient below 0.1 m^{-1} , transition-metal impurity must be limited to 0.5 wppm.

Excited-state lifetime also varies with neodymium ion concentration N_A ; observed dependences range from N_A^{-1} to N_A^{-3} . In an attempt to understand this relation, we carried out numerical simulations of energy diffusion among ions. A function for the fluorescence intensity vs time was then determined that gave a best fit to the numerical results. Application of these results to laser glasses has not yet been undertaken.

We began promising work on highly reflective (HR) sol-gel coatings. We produced multilayer coatings of SiO_2 and TiO_2 with reflectivities up to 99% and damage thresholds about the same as those of conventional coatings (produced by e-beam evaporation). In work funded by

the Air Force Weapons Laboratory, we produced porous antireflection (AR) coatings of MgF_2 and CaF_2 with damage thresholds 2 times higher than those of conventional coatings. These coatings, the first non-oxide sol-gel coatings ever made, may have application in the chemically corrosive atmosphere of XeF lasers.

We measured thresholds for laser-induced damage arising from iridium inclusions in the laser material gadolinium scandium gallium garnet (GSGG). This information will be used to improve the GSGG growth process to eliminate iridium inclusions. We also measured the damage thresholds of several nonlinear optical crystals and of CaF_2 ; the CaF_2 work was funded by the Air Force Weapons Laboratory.

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Impurity Studies in Laser Glasses

In collaboration with Schott Glass Technologies, we have measured the spectra of transition-metal ions (Ti, V, Cr, Mn, Fe, Co, Ni, and Cu) and the effects of these ions on neodymium fluorescence quenching in neodymium-doped phosphate and silicate glasses. Our purpose was to determine the maximum allowable impurity content, given limits on the absorption loss at $1.053 \mu\text{m}$ and on the neodymium fluorescence quenching rate. To keep the absorption loss below 0.1 m^{-1} , the transition-metal impurity content should be kept below 0.5 wppm (weight parts per million). To keep the increase in the neodymium fluorescence decay rate below 1%, the impurity content should be less than 3 wppm .

Introduction. Transition-metal (TM) impurities are detrimental to laser-glass performance in two ways. First, their absorption bands may overlap the laser transitions between 1.050 and $1.060 \mu\text{m}$. This absorptive loss will decrease the amount of stored energy that can be extracted from the pumped laser glass. As energy is extracted, the gain coefficient G drops until it becomes equal to the loss coefficient L ; at this point no further

energy extraction is possible. The maximum possible extraction efficiency η is related to the initial gain G_0 and the loss L through the expression⁴⁶

$$\eta = \frac{1}{\ln G_0} \left\{ \sqrt{1 - L \ln \left(\frac{G_0 - 1}{1/L - 1} \right)} - \frac{1}{\sqrt{1 - L}} \ln \left[\frac{(1 - 1/G_0)}{L} \right] \right\}. \quad (19)$$

To obtain 75% extraction efficiency at $G_0 = 7.4$ requires a gain-to-loss ratio of 20. Because typical laser glasses in amplifiers have gain coefficients of 2.5 to 5 m^{-1} , this implies (if we allow for surface scattering) that the loss should be less than 0.1 m^{-1} .

Transition-metal impurities also quench the neodymium fluorescence; that is, they induce nonradiative transitions from the $^4\text{F}_{3/2}$ excited state of the neodymium ion.⁴⁷ The correspondingly shorter lifetimes τ of the $^4\text{F}_{3/2}$ state lead to lower amplifier efficiency in flashlamp-pumped systems, because the flashlamp pulse width is about the same as the neodymium fluorescence lifetime. As the lifetime decreases, the pump efficiency decreases as τ^x , where x is between 0.5 and 0.7. Trying to compensate for the shorter neodymium lifetime by decreasing the flashlamp pump pulse width does not work, because the radiative efficiency of the lamp decreases and its emission shifts toward the ultraviolet, which increases its spectral mismatch with neodymium. We want less than a 1% decrease in the neodymium fluorescence lifetime due to impurities.

We measured the spectra and quenching effects of TMs in a silicate and two phosphate glasses that may have applications in future fusion research lasers.

The metastable $^4\text{F}_{3/2}$ state of Nd^{3+} has a $400\text{-}\mu\text{s}$ fluorescence lifetime. The radiative quantum efficiency of this state is near unity in most glasses.^{48,49} The fluorescence can be quenched, however, by nonradiative energy transfer to vibronic or electronic energy states of the host material or of ionic or molecular impurities. Forster⁵⁰ and Dexter⁵¹ analyzed the dipole-dipole energy transfer process between localized energy states. Their analysis is now known

as the Forster-Dexter theory. In brief, the probability W_{DA} of energy transfer from a donor D to an acceptor A is given by

$$W_{DA} = \frac{9\chi^2}{128\pi^5 N_A n^4 \tau_{0D} R_{DA}^6} \times \int g_D(\nu) K_A(\nu) \nu^{-4} d\nu, \quad (20)$$

where the factor χ^2 takes into account the orientational averaging of the dipole, N_A is the acceptor concentration, n is the refractive index, R_{DA} is the donor-acceptor distance, τ_{0D} is the donor radiative lifetime, $g_D(\nu)$ is the shape factor of the fluorescence band or bands, $K_A(\nu)$ is the attenuation coefficient of the acceptor, and ν is the frequency (in wave numbers). This expression indicates that four important physical parameters determine energy transfer: the donor-acceptor distance, the donor fluorescence and acceptor absorption strengths, and the spectral overlap of fluorescence and absorption. Many authors have applied, extended, and

modified the Forster-Dexter theory for rare-earth energy transfer in crystals and glasses; reviews can be found in Refs. 52 and 53. The theory is also discussed in the following article, "A Model for Fluorescence Quenching."

Experiments. We investigated two phosphate glasses and one silicate glass.

The phosphate glasses, designated UP-91 and UP-16, had the following compositions (in mole %):

UP-91: 68.2 P_2O_5 , 11.8 Al_2O_3 , 14.0 K_2O , 6.0 (La + Nd) $_2O_3$.

UP-16: 70 P_2O_5 , 20 K_2O , 10 (La + Nd) $_2O_3$.

UP-91 may be a candidate for future fusion research lasers. It has good spectroscopic properties⁵⁴ and is easy to melt in production. It was developed primarily for its relatively low neodymium self-quenching and its good thermomechanical properties. UP-16 has even lower neodymium self-quenching than UP-91, but it is impractical because of its poor thermomechanical properties and poor environmental stability.

The silicate glass was LG-660, a commercial alkali-zinc-silicate laser glass. Its neodymium transition strength and fluorescence lifetime are near the median values of these quantities for silicate laser glasses. LG-660 has been tested in large amplifiers and shows good storage efficiency.⁵⁵

We chose neodymium doping concentrations for which we could observe quenching effects with and without energy diffusion in the neodymium system. We melted glasses with 0.5×10^{20} ions cm^{-3} ; at these concentrations, the glass exhibits no energy diffusion and little neodymium self-quenching. To see the effects of diffusion at two neodymium concentrations, we melted glasses with 5 and 10×10^{20} neodymium ions cm^{-3} . At these concentrations, neodymium self-quenching is present.

Transition metals generally appear as impurities in laser glasses at less than 10 wppm. Therefore, we kept the acceptor concentration considerably below that of neodymium, and added TM concentrations of 30, 100, 300, or 1000 wppm, depending on the magnitude of the absorption or quenching effect.

Fig. 7-45. Vanadium spectra in UP-91 and LG-660 laser glasses. Sample thickness: 1 cm. Nominal V concentration: 300 wppm.

Fig. 7-46. Manganese spectra in UP-91 and LG-660 laser glasses. Sample thickness: 1 cm. Nominal Mn concentration: 300 wppm.

Schott Glass Technologies inductively melted the silicate glass in a platinum crucible and the phosphates in quartz with an atmosphere of dry nitrogen, and homogenized the melted glasses with a platinum stirrer. The melted glass was cast in a graphite-coated steel mold, and then annealed at a temperature drop rate of 30°C/h. The platinum crucible and stirrers were cleaned in HCl solution between each TM series. Samples of the resultant glasses were cut and polished into 20- × 20- × 10-mm pieces for absorption measurements and into 20- × 20- × 1-mm pieces for fluorescence decay measurements. We kept the glasses as free of water as possible, because OH⁻ is a well-known quencher of neodymium fluorescence.⁵⁶ Dry nitrogen was bubbled through the melted phosphate glasses. Most of the glasses had water absorption coefficients <1 cm⁻¹ at the infrared band peak (~2900 cm⁻¹). For glasses with a water absorption coefficient of 1 cm⁻¹, the quenching rate is about 60, 160, and 300 s⁻¹ for neodymium concentrations of 0.5, 5, and 10 × 10²⁰ cm⁻³, respectively. When necessary we corrected the fluorescence decay rates for variations in the water content.

Because many of our decays were exponential or nearly so, we took the inverse of the first *e*-folding time as the decay rate. This procedure is accurate enough for our analyses, and is appropriate because the efficiency in a laser amplifier depends more strongly on the decay rate during the first *e*-folding time than on the rate during later *e*-folding times. For silicates, even in the absence of quenching acceptors, the decay function is nonexponential. Phosphate glasses, on the other hand, have nearly exponential decays.

Results.

Absorption Spectra. Figures 7-45 through 7-48 show the spectra of V, Mn, Co, and Ni in various glasses. The spectra of the titanium-doped glasses (not shown) exhibit only an increased ultraviolet absorption. The chromium-doped glasses have the usual Cr³⁺ spectra,⁵⁷ and Cr⁶⁺ is evident in the silicate glasses.⁵⁸ The iron and copper spectra are similar to those in previous measurements.⁵⁷ After analysis of the spectra and comparison with previous measurements and analyses, we have

assigned charge states and coordination for each TM: Table 7-3 gives the results.

In general the crystal field, or Dq value, for ions in octahedral coordination is higher for the silicate glass LG-660 than for the phosphate glasses UP-91 and UP-16. This probably results from the truly three-dimensional cross-linked nature of the silicate; the phosphate glasses are primarily chain structures with some chain cross-linking for these ultraphosphates. Thus the silicate structure is tighter, with smaller "holes" for the TM.

As Table 7-3 shows, we find higher oxidation states in the silicate than in the phosphates under our melting conditions. The one exception could be Cu-doped LG-660, in which the Cu²⁺ band is

Fig. 7-47. Cobalt spectra in UP-91 and LG-660 laser glasses. Sample thickness: 1 cm. Nominal Co concentration: 300 wppm.

Fig. 7-48. Nickel spectra in UP-91 and LG-660 laser glasses. Sample thickness: 1 cm. Nominal Ni concentration: 300 wppm.

Table 7-3. Observed ionic states and coordinations of transition metals in two laser glass types.

Phosphate (UP-91)	Silicate (LG-660)
Cu ²⁺ , octahedral	Cu ²⁺ , octahedral Cu ⁺ (implied)
Fe ²⁺ , octahedral (tetrahedral)	Fe ³⁺ , Fe ²⁺ octahedral (tetrahedral)
Co ²⁺ , tetrahedral (octahedral)	Co ²⁺ , tetrahedral
Ni ²⁺ , octahedral	Ni ²⁺ , octahedral (tetrahedral)
Cr ³⁺ , octahedral	Cr ⁶⁺ , Cr ³⁺ , octahedral
Mn ²⁺ (doubtful)	Mn ³⁺ , octahedral
V ³⁺ , octahedral	V ⁴⁺ (doubtful)

Table 7-4. Absorption coefficients of transition metals at the neodymium laser wavelength (m^{-1} per wppm metal added).

Ion	Phosphate glasses ($\lambda = 1.053 \mu\text{m}$)		Silicate glass ($\lambda = 1.057 \mu\text{m}$)
	UP-91	UP-16	LG-660
Cu	0.23	0.21	0.038
Fe	0.07	0.06	0.110
Co	0.035	—	0.084
Ni	0.028	—	0.064
V	0.017	—	0.013
Mn	(0.002)	—	0.009

Fig. 7-49. Quenching rate of neodymium fluorescence for 300 wppm added transition-metal impurity.

considerably less intense than expected, suggesting that some of the copper may be in the monovalent state.

We determined the TM absorption coefficients at $1.053 \mu\text{m}$, per wppm of metal added, using the nominal added concentrations. Table 7-4 gives the results. In phosphate glasses, Cu^{2+} is the strongest absorber; in silicates, Fe^{2+} is the strongest. To keep the absorption coefficient below 0.1 m^{-1} in laser glasses, the concentrations of these ions must be kept below 0.5 to 1 wppm.

Fluorescence Quenching. The quenching rates (i.e., the differences in decay rates with and without the added TM impurity) of neodymium fluorescence for 300 wppm added impurity are shown in Fig. 7-49 for glasses with 5×10^{19} neodymium ions cm^{-3} . Copper is the most effective quencher in phosphate glass; cobalt is most effective in LG-660 silicate glass. Arbuzov et al.⁴⁷ found the following order for quenching effectiveness of TMs in GLS-22 barium-phosphate glass: $\text{Cu} > \text{V} > \text{Ni}$, and $\text{Co} > \text{Fe}$. We are surprised by the large effect of cobalt in both the phosphate and silicate glasses. We cannot explain the quenching effectiveness of the various TMs by the variation in the spectral overlap of the neodymium fluorescence and the TM absorption [i.e., by

variation in the overlap integral

$$M = \int g_D K_A v^{-1} dv \quad (21)$$

in Eq. (20)]. We found the quenching rate in these glasses to be approximately directly proportional to the impurity content as determined from the absorption band intensities. For a true Forster-Dexter static decay,^{50,51} the inverse of the first e -folding time does not depend linearly on the acceptor concentration. In fact, for high quenching rates the dependence is quadratic. In addition, the Forster-Dexter theory does not correctly describe the initial decay, which is an exponential whose rate is determined by the smallest neodymium-acceptor distance. We sidestep this issue by accepting the measured linearity and asking how the slopes are related to the elementary dipole-dipole quenching process as described by the overlap integral M . The slopes should be proportional as we change the TM impurity for a given glass, assuming that the neodymium-TM distance distribution is the same.

Table 7-5 gives the normalized quenching rates for the various TMs in UP-91, UP-16, and LG-660. We calculated the overlap integral from the branching ratios of the $^4\text{F}_{3/2}$ transitions determined by using the Judd-Ofelt theory (which has been confirmed experimentally⁵⁹) and the measured TM spectra. We included the fluorescence transition from the $^4\text{F}_{5/2}$ state, which has a small population at room temperature.

Straightforward application of the dipole-dipole quenching theory would lead us to expect that all the ratios in Table 7-5 should be nearly the same for a given glass. The differences observable in the table could arise from deviations from a random distribution of acceptors; however, we suggest that the deviations probably result from phonon-assisted energy transfer, which is not included in the overlap-integral description. For instance, both cobalt and nickel have low energy states to which neodymium ions can "back-transfer" energy with the emission of phonons. Because the theory of phonon-assisted energy transfer is not well developed, we are unable to be quantitative in our analysis.

Table 7-5 indicates that the quenching rates of LG-660 are a factor of 2 lower than those of UP-91. Part of this is due to the lower neodymium radiative rate of LG-660, which is 72% of that for UP-91.

Note that energy transfer from neodymium to chromium is observed in UP-91, even though the Cr^{3+} absorption band overlaps only the weak neodymium fluorescence bands at 0.874 and 0.802 μm . Several unsuccessful attempts have been made⁶⁰⁻⁶² to use Cr^{3+} as a sensitizer for neodymium in glass. The low quantum efficiency of Cr^{3+} in glass,⁶³ coupled with the neodymium-to-chromium back-transfer, appears to explain this failure.

At higher neodymium concentrations, the quenching rates increase as energy diffusion takes place within the neodymium system. Figure 7-50 shows the ratio of the neodymium quenching rate for TMs at 5 and 10×10^{20} neodymium ions cm^{-3} to the rate at $0.5 \times 10^{20} \text{ cm}^{-3}$. We have not yet attempted to explain these results in terms of the quenching theories.

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A Model for Fluorescence Quenching

Introduction. Fluorescence quenching has a significant impact on the performance of Nd:glass lasers, and a better understanding of this phenomenon is therefore essential in the search for improved laser materials. Experiments show that fluorescence quenching can vary from a linear to more than a cubic dependence on neodymium ion concentration. Theories that attempt to explain quenching are incomplete, and the results of different workers are inconsistent because different theories are applied for different donor concentrations. With increasing donor concentration, fluorescence quenching usually is broken down into static quenching (SQ), diffusion-limited relaxation (DQ), migration quenching (MQ), and kinetic quenching (KQ). This complexity obviously leads to difficulties in transition regimes and to

Ion	UP-91	UP-16	LG-660
V^{3+}	92	—	—
Cr^{3+}	53	—	—
Fe^{2+}	66	98	34
Co^{2+}	107	—	47
Ni^{2+}	111	—	60
Cu^{2+}	38	47	51
τ_{0D} (μs)	390	—	540

^a Quenching rate (s^{-1}) divided by 100
 $\int g(\nu)\alpha(\nu)(\nu/\nu_0)^{-4} d\nu$, where ν is in cm^{-1} , α is the absorption coefficient, and $\nu_0 = 1.053 \mu\text{m}$.

Table 7-5. Normalized quenching rates^a of neodymium by transition-metal ions for 5×10^{19} neodymium ions cm^{-3} .

Fig. 7-50. Increase of quenching rate (expressed as ratio to quenching rate at 5×10^{19} neodymium ions cm^{-3}) with neodymium concentration.

uncertainty about which theory is applicable. We have developed a more general, unified theoretical approach that provided smooth transitions from one regime to the next and that uses a single set of material constants instead of four.

Theoretical Review. The effective fluorescence decay rate $\phi(t)$, including energy transfer from donors D (e.g., Nd^* ions) to acceptors A (e.g., impurities) can be approximated for electric dipole interactions and for uniform distributions of randomly oriented excited and quenching ions by⁶⁴:

$$\phi(t) = \bar{N}_D / N_{D0} \approx \exp\left(-\frac{t}{\tau_0} - \gamma t^{1/2}\right), \quad (22)$$

where τ_0 is the intrinsic donor lifetime in the absence of acceptors. If there is no energy migration between donors, a condition that corresponds to static quenching and that occurs at low donor concentrations,

γ is given by⁶⁴

$$\gamma = \gamma_0 = \left(\frac{\pi}{\tau_0} \right)^{1/2} \left(\frac{R_0}{R_A} \right)^3 = \frac{4}{3} \pi^{3/2} N_A (C_{DA})^{1/2}, \quad (23)$$

where $R_0 = (C_{DA} \tau_0)^{1/6}$ is the critical transfer distance, $R_A = (3/4 \pi N_A)^{1/3}$ is an effective radius of the sphere of influence around each quenching center, and N_A is the acceptor concentration. The microparameter C_{DA} for resonant donor-to-acceptor interactions depends on the oscillator strengths and the spectral overlap integrals of the interacting transitions, and can be obtained from spectroscopic luminescence and absorption data with⁶⁵

$$C_{DA} = W_{DA} R_{DA}^6 = \frac{3 C S_A}{6^4 \pi^5 n^4 \tau_i} \times \int g_D(\nu) g_A(\nu) \frac{d\nu}{\nu^4} \quad \text{cm}^6/\text{s}, \quad (24)$$

where W_{DA} is the transfer probability, R_{DA} is the donor-to-acceptor distance, C is a local field correction factor close to unity, and $S_A = \int \sigma_A(\nu) d\nu$ (cm) is the absorption band strength, with $\sigma_A(\nu)$ the absorption cross section and ν the wave number; $g(\nu)$ represents the line shape function, normalized to $\int g(\nu) d\nu = 1$. n is the index of refraction, and $\tau_i = \tau_R/\beta_i$ is the radiative lifetime of the transition, with τ_R the net radiative lifetime and β_i the branching ratio.

The effects of diffusion on energy transfer are described by a diffusion coefficient D for energy transfer between donors, which depends on the donor concentration N_D and can be correlated with the donor-donor interaction characteristic C_{DD} by⁶⁴

$$D = s N_D^{4/3} C_{DD} \quad \text{cm}^2/\text{s}, \quad (25)$$

where $s \simeq 2.8$ is a dimensionless constant. The parameter⁶⁶

$$x = D(t^2/C_{DA})^{1/3} \quad (26)$$

provides a measure for the significance of diffusion-limited relaxation. Initially, when $x \ll 1$, diffusion is not important, but later ($x > 1$) diffusion leads to exponential diffusion-limited relaxation⁶⁶ with

$$W_D = k_D N_A N_D (C_{DA} C_{DD}^3)^{1/4} \quad \text{s}^{-1}, \quad (27)$$

where k_D is a constant whose value depends on the models and constraints used by different authors.

There are other regimes, in which energy transfer is described by migration-controlled quenching, using random-walk models, or by kinetic quenching, in which the diffusion rate between donors is much greater than the acceptor quenching rate. The decay rates for these cases are in general described by⁶⁷

$$W_M = k_M N_A N_D (C_{DA} C_{DD})^{1/2} \quad (28)$$

and

$$W_K = k_K N_A N_C C_{DA}, \quad (29)$$

where N_C is the critical donor concentration for the transition from MQ to KQ. The values of the dimensionless coefficients $k_M = 24 \pm 4$ and $k_K = 2\pi^2/3$ depend to some extent on the models used by different authors.

Cross-relaxation⁶⁸ is a significant quenching process for neodymium-doped glasses, because excited Nd^{3+} ions can interact with unexcited neighbors by exchange transitions to produce two intermediate-state ions ($^4I_{15/2}$) that then both decay nonradiatively to the ground state. In this process the neodymium ions also act as acceptors, leading to the usually observed proportionality of the decay rate to the square of the neodymium concentration N for diffusion and migration-controlled quenching.

Unified Approach. To provide a general description of quenching, including all regimes from static to kinetic quenching, the diffusion equation

$$\dot{N}_D = D \nabla^2 N_D - \frac{N_D C_{DA}}{R^6} - \frac{N_D}{\tau_0} \quad (30)$$

was evaluated numerically within the sphere of influence of an acceptor for various values of D , C_{DA} , and R_A . The results were then approximated with a more refined decay function of the form

$$\phi(t) = \exp \left(-\frac{t}{\tau_0} - Wt - 0.2\gamma_0^2 t - \gamma \frac{t}{\sqrt{\tau_1 + \nu t}} - \delta t^{1.5} \right), \quad (31a)$$

where

$$W = W_D \left\{ 1 + k_1 \frac{C_{DD}}{C_{DA}} \right. \\ \times \left[1 + k_2 N_D^2 N_C^4 \left(\frac{C_{DD}}{C_{DA}} \right)^2 \right] \left. \right\}^{0.25}, \quad (31b)$$

$$\gamma = \gamma_0 \left[1 + k_3 N_D \left(\frac{C_{DD}^{0.75}}{C_{DA}^{0.25}} \right) \right]^{-1}, \quad (31c)$$

$$\delta = k_4 N_A^2 N_D (C_{DA} C_{DD})^{0.75}, \quad (31d)$$

and where $k_D \simeq 22$, $k_1 = (k_D/k_M)^4 \simeq 0.7$, $k_2 = (k_K/k_M)^4 \simeq 0.0056$, $k_3 \simeq 1.3$, and $k_4 \simeq 170$. These approximations give results within 5% of those given by the numerical simulations. The values of the coefficients k_M and k_K are taken from Ref. 69.

The decay function of Eq. (31a) provides smooth transitions from diffusion to migration and from migration to kinetic quenching as donor concentration is increased. The fourth-power transition applied in Eq. (31b) from diffusion [Eq. (25)] to migration [Eq. (26)] was obtained from numerical simulations, while the fourth-power transition from MQ to KQ [Eq. (27)] was obtained by matching the data in Refs. 69 and 70 on (La-Nd)F₃ and Li-La-Nd phosphate glasses.

To match the numerical simulations, Eq. (31b) required several terms. For $\gamma_0 t^{1/2} \ll 1$, which corresponds to low quenching rates, the method of Förster [Eq. (22)] is adequate. For $\gamma_0 t^{1/2} \simeq 1$, the next higher term, which can be approximated by $0.2\gamma_0^2 t$, must be included.

The initial infinite decay rate $-0.5\gamma/t^{1/2}$ of Eq. (22) arises because of the integration of the quenching probability over the spherical volume from 0 to R_A ; it is more of a computational annoyance than a serious error. A more reasonable approach is to begin the integration at some nonzero R_{min} , recognizing that it is impossible to find a donor and acceptor in the same place. This correction can be expressed quite well by replacing $\gamma t^{1/2}$ with $\gamma t/(\tau_1^{1/2} + t^{1/2})$, where $\tau_1 = R_{min}^6/\pi C_{DA}$.

The minimum distance between acceptor and donor for neodymium ions is about $R_{min} \simeq 4$ Å, which yields $\tau_1 = 13 \mu s$ for typical values of $C_{DA} = 10^{20} \text{ cm}^{-3}$. The initial decay rate for the quenching term is then limited to $\gamma/\tau_1^{1/2}$.

The γ term of Eq. (31c) vanishes for high diffusion coefficients, while at high acceptor concentrations an additional δ term is required to match the numerical simulations. This is only noticeable at the later stages of decay and hence is not particularly significant.

The numerical simulations based on diffusion lead to a correction term [the k_1 term in Eq. (31b)] that provides a smooth transition from diffusion to migration. Even though the motion of the excitation differs between these regimes, the character of the motion seems to be irrelevant. That is, it appears not to matter whether the excitation falls into the quenching sphere by making many short jumps along the "excitation" gradient or by one of the many random hops.

This can be better understood by examining the growth of the quenching sphere and the excitation density within it. The growth of the quenching sphere is shown in Fig. 7-51. Without diffusion [Fig. 7-51(a)], the sphere grows with time, while with diffusion [Fig. 7-51(b)] the void in excitations near the acceptor does not grow, but the gradients are reduced. Figures 7-51(c) and 7-51(d) show the radial distributions of excitations around a quenching center for various values of D and R_0 at fixed $t = \tau_0$. For high diffusion coefficients, excitation gradients of the quenching sphere become steeper, and a larger volume is affected by the quenching center than is the case without diffusion. For very high diffusion coefficients, the transition to the random walk model with localized quenching spheres surrounded by a uniformly depleted volume is quite natural.

Determination of Material Constants.

Despite the complex appearance of Eq. (31), concentration quenching is primarily determined by the energy transfer coefficients C_{DA} and C_{DD} . Only for high donor concentrations is it necessary to find the appropriate value of N_C , and the value of τ_1 is not critical because it has little impact on the decay function.

At low donor concentrations, Eq. (31a) reduces to

$$\phi_A(t) \approx \exp\left(-\frac{t}{\tau_0} - 0.2\gamma_0^2 t - \gamma_0 \frac{t}{\sqrt{\tau_1 + \sqrt{t}}}\right), \quad (32)$$

which allows the determination of τ_0 , τ_1 , and γ_0 from the measured fluorescence decay curve. At low acceptor concentrations, $\phi_0(t)$ is dominated by τ_0 and becomes exponential.

The ratio of the initial and final decay rates is

$$K_0 = \frac{\dot{\phi}_0(0)}{\dot{\phi}_0(\infty)} \approx 1 + \frac{\gamma_0}{\sqrt{\tau_1} \left(\frac{1}{\tau_0} + 0.2\gamma_0^2\right)}, \quad (33)$$

and this expression can be used to determine γ_0 and τ_1 . Since it is experimentally difficult to determine the initial decay rate,

Fig. 7-51. Radial excitation distributions (a) without and (b) with diffusion for various times t/τ_0 . The diffusion coefficient is normalized by $D^* = D\tau_0/R_A^2$. (c) and (d) The influence of diffusion is reduced for lower R_0 values.

the ratio t_1/t_3 of the first to third e -folding times provides a more reliable measure for γ_0 . For $\tau_1 \ll \tau_0$ and $\delta \rightarrow 0$, the nonexponential decay ratio is

$$K_{31} = \frac{t_3 - t_2}{t_1}$$

$$= \frac{1 + \sqrt{2\beta + \frac{\beta^2}{4}} - \sqrt{3\beta + \frac{\beta^2}{4}}}{1 + \frac{\beta}{2} - \sqrt{\beta + \frac{\beta^2}{4}}}, \quad (34a)$$

where

$$\beta = \frac{\gamma_0^2}{\left(\frac{1}{\tau_0} + W + 0.2\gamma_0^2\right)}$$

$$\simeq 1 - 2.2K_{31} + 1.2K_{31}^2. \quad (34b)$$

Using Eq. (34), one can determine the value of β from the measured nonexponential decay ratio K_{31} . For dry phosphate laser glasses the fluorescence decay is almost exponential, so K_{31} is close to unity. This implies $\beta \ll 1$ or $\gamma_0^2 \ll W_0 = 1/\tau_0 + W + 0.2\gamma_0^2$, which in turn requires either very low acceptor concentrations N_A or small energy-transfer coefficients C_{DA} . Phosphate glasses with high water content exhibit nonexponential decays as expected, however. Silicate laser glasses exhibit significant nonexponential decay with $K_{31} \simeq 1.5$ and $K_0 \simeq 5$, which implies $\beta \simeq 0.4$, $\beta/\gamma_0\tau_1^{1/2} = K_0 - 1 \simeq 4$, and hence $\gamma_0\tau_1^{1/2} \simeq 0.1$ or $\gamma_0 \simeq 30 \text{ s}^{-0.5}$ for $\tau_1 \simeq 10 \mu\text{s}$. These values are of the correct magnitude.⁶⁷ Evaluations of specific laser materials are required to determine their material constants and correlate these to overlap integrals and finally to compositions.

The change in the nonexponential decay with neodymium concentration N is quite complex, but to first order it can be scaled with $\beta^{1/2} \sim N/(N_0^2 + N^2)^{1/2}$ for $N_A \sim N_D \sim N$. This implies that the nonexponentiality first increases with N but then stabilizes for $N > N_0$, where N_0 is a complex function of τ_0 , C_{DA} , and C_{DD} , combining Eqs. (27) and (31). This behavior agrees qualitatively with experimental observations on various laser glasses. However, it should also be recognized

from Eq. (34) that long-lifetime glasses will exhibit stronger nonexponential decay than short-lifetime glasses, as observed experimentally. Site-to-site variations will obviously modify such general conclusions, but the trend is correct.

Concentration Quenching. We must determine whether the unified approximations provide the correct scaling with neodymium concentration. The scaling is nearly linear with concentration for $N > N_c$, or is independent of N for cases in which impurities dominate; the scaling is nearly quadratic for processes such as cross-relaxation, where $N_A = N_D = N$ with the additional constraint $N < N_c$. The scaling becomes near-cubic for $N \ll N_c$ and $N_A \sim N^2$ as a result of formation of neodymium pairs or more complex associations. Phonon-assisted transitions could also increase the dependence of quenching on neodymium concentration and are more likely. Other considerations can influence concentration quenching, such as irregular R_{DA} distributions due to segregation, structural changes of the glass with increasing neodymium doping, impurities introduced along with neodymium doping, and changes in the homogeneous and inhomogeneous line-width ratios. Short-pulse excitation will emphasize nonexponential decay, while the effective decay rate is more exponential for typical pump pulses in actual laser systems.

Unifying the diffusion, migration, and kinetic quenching regimes eliminates the possibility of large errors in the determination of energy transfer coefficients. This method can give us better understanding of fluorescence quenching.

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Sol-Gel Coatings

Highly Reflective (HR) Coatings. We investigated HR coatings prepared from colloidal oxide suspensions. It was hoped that high reflectors could be made in a solution process by depositing alternating layers of porous high- and low-index oxides on a suitable substrate. SiO_2 was selected as the low-index component because we had already developed a high-damage-threshold coating prepared from a

SiO₂ suspension. This coating, which has an index of refraction of about 1.22, is used on certain optics in the Nova laser to reduce Fresnel reflection. Among a number of possibilities for the high-index component, TiO₂ was selected for initial investigation. After considerable effort we found that a colloidal suspension of TiO₂ could be obtained from the hydrolysis of titanium isopropoxide with a strong organic base catalyst in water:

After solvent exchange, this suspension could be used to prepare porous TiO₂ coatings with an index of about 1.9 by application to suitable substrates at room temperature.

HR coating samples were prepared on silica substrates by depositing alternating layers of TiO₂ and SiO₂ at room temperature with a few minutes of drying time between coats. Coatings with as many as 31 layers showed no sign of peeling or

crazing and were probably therefore strain-free. Figure 7-52 shows the transmission spectrum of a 17-layer coating whose reflectance is ~99% at 0.6 μm. By changing the thickness of the individual oxide layers, HR coatings with reflection maxima at wavelengths from 0.4 μm to several microns can be prepared.

Laser-damage thresholds of 16 single-layer films of TiO₂ were measured with 1-ns, 1.064-μm pulses. The thresholds ranged from 5 to 12 J/cm²; the median was 8 J/cm². These films were deposited from colloidal suspensions of TiO₂ made from Ti(OC₃H₇)₄ using (CH₃)₄NOH or (C₂H₅)₂NH as a catalyst, and were dried at various temperatures between 20 and 150°C. Variations of the deposition conditions did not systematically alter the damage thresholds.

Four five-layer reflectors made of alternating quarter-wave layers of TiO₂ and SiO₂ deposited from colloidal suspensions were also tested. The thresholds of these films ranged from 3 to 7 J/cm², which is within the typical damage threshold range now obtained for conventional e-beam-deposited coatings. We hope that these thresholds can be improved when coatings are prepared under clean-room conditions.

Porous Antireflective (AR) Coatings for the USAF. AR coatings for the internal optics of fluoride lasers must obviously be resistant to the reactive environment in the lasing cavity; metal fluorides are therefore a natural choice as coating materials. The use of fluoride layers in dielectric coatings is also a possibility for ICF applications. Under a U.S. Air Force contract,⁷¹ we are investigating the possibility of preparing high-damage-threshold, porous (hence, low-index) AR coatings of MgF₂ or CaF₂ for use in fluoride lasers. Our experience with silica showed that colloidal suspensions could be used to prepare excellent porous coatings at room temperature. These coatings consist of layers of silica particles and are used on the fused silica and potassium dihydrogen phosphate (KDP) optical components on the Nova laser. It was hoped that colloidal fluoride suspensions could be used similarly to prepare porous fluoride coatings.

Calcium and magnesium fluorides may be prepared by several methods and from

Fig. 7-52. Transmission of a 17-layer TiO₂-SiO₂ HR coating on a SiO₂ substrate.

many different starting materials. A number of combinations were investigated before we obtained useful colloidal suspensions of the desired particle size.

Two preparations were developed. The first uses the reaction of metal acetates with hydrofluoric acid in anhydrous methanol solution:

Colloidal suspensions of both calcium and magnesium fluorides were obtained by this method. The particles were substantially monodispersed and were 100 to 200 Å in diameter.

The second method involved the reaction of metal alkoxides with hydrofluoric acid in anhydrous methanol solution:

A colloidal suspension of magnesium fluoride was obtained by this method; the particles had properties similar to those obtained from the acetate. The insolubility of calcium alkoxides precluded their use in this preparation.

Colloidal suspensions prepared as described above were used to coat polished fused silica and calcium fluoride substrates. All coatings were applied using a spin coater and were air dried at room temperature. The two-surface transmission spectra of MgF_2 on CaF_2 and CaF_2 on SiO_2 are shown in Figs. 7-53 and 7-54, respectively. Similar spectra were obtained for colloidal CaF_2 on CaF_2 and MgF_2 on SiO_2 . These spectra are characteristic of quarter-wave AR coatings in which the refractive index (1.22) corresponds to about the square root of the index of the substrate, thereby giving near-100% transmission.

Laser damage thresholds for eight MgF_2 coatings on fused silica substrates were measured with 0.351- μm , 25-ns pulses from an XeF laser. Each test site received 1000 pulses at a repetition rate of 25 Hz. Figure 7-55 shows thresholds for these coatings. A few measurements on CaF_2 coatings gave similar results. All

Fig. 7-53. Two-surface transmission of a MgF_2 coating on a CaF_2 substrate.

Fig. 7-54. Two-surface transmission of a CaF_2 coating on a SiO_2 substrate.

Fig. 7-55. Damage thresholds of MgF_2 coatings on silica substrates for $0.351\text{-}\mu\text{m}$, 25-ns pulses (1000 shots per site at 25 Hz).

Fig. 7-56. Damage thresholds of SiO_2 coatings on silica substrates for $0.351\text{-}\mu\text{m}$, 25-ns pulses (1000 shots per site at 25 Hz).

of the coatings tested were prepared for maximum transmittance at wavelengths near $0.350\text{ }\mu\text{m}$.

For the USAF, we made similar damage tests of SiO_2 AR coatings deposited from our standard colloidal silica suspension onto silica substrates. These thresholds are given in Fig. 7-56. Thresholds of the fluoride and silica coatings for 25-ns pulses are both more than twice as high as those of the best e-beam-deposited coatings under these test conditions.

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Laser Damage in Optical Crystals

Crystals play several vital roles in fusion lasers: they are candidate hosts for lasing ions and are useful as transmissive

windows, and optically nonlinear crystals are used as harmonic converters. This article describes the results of laser damage tests on a number of optical crystals. Additional damage studies are discussed in the articles "Measurement of Damage Initiation and Growth due to Metallic Platinum Inclusions" in Sec. 6, and "Sol-Gel Coatings," and "Optical Damage of KDP," later in this section.

Iridium Inclusions in GSGG Crystals.

Crystals of GSGG, which are grown in iridium crucibles, sometimes contain iridium inclusions that can lead to laser-induced damage. The volume density of inclusions varies greatly. Some laser rods with volumes of a few cubic centimeters are apparently free of inclusions, while isolated crystalline segments sometimes contain as many as 6×10^3 inclusions per cubic centimeter. The inclusions appear as opaque objects when backlighted and observed by optical microscopy. The largest of these objects are triangular or

Fig. 7-57. (a) Damage growth around a triangular iridium inclusion and two smaller, initially circular iridium inclusions in GSGG. Pulse duration was 9 ns. (b) Fluences in successive pulses.

hexagonal, suggesting that they were formed by condensation of iridium from the melt. In the investigations reported here, we could not resolve the smallest particles well enough to determine whether they had crystalline shapes.

A crystalline GSGG segment containing many inclusions was irradiated with 9- and 50-ns, 1.064- μm pulses. In the volume irradiated with 9-ns pulses, we observed nine inclusions in a single plane, ranging in size from 2.4 to 18 μm . All of these inclusions survived six shots at fluences below 1.6 J/cm^2 . Five survived two additional shots at 2.7 and 2.9 J/cm^2 , but were damaged by an additional shot at 4.5 J/cm^2 . Growth of the fractured volume surrounding the inclusions stopped after seven to

nine additional shots at 4 to 6 J/cm^2 . The evolution of the size of damage surrounding three of these inclusions is shown in Fig. 7-57. (Small variations in the measured sizes of damage in Fig. 7-57 are due to subtle differences in focusing and illumination of the photomicrographs.) Irradiation at higher fluences would probably cause additional growth of these damage sites, but use of GSGG above 4 to 5 J/cm^2 is unlikely, because the saturation fluence of this laser material is about 1 J/cm^2 .

In the 50-ns damage experiment we observed 27 inclusions in 6 planes at different depths in the 1-mm-diam inspection volume. The sizes of these inclusions varied from 2.8 to 42 μm . Two were equilateral triangles with side lengths of 16 and

Fig. 7-58. (a) Damage growth around a triangular iridium crystal, a 4.6- $\times$ 29- μm iridium particle, and a 3- μm -diam particle in GSGG during 41 50-ns pulses. (b) Fluences in successive pulses.

42 μm . Two were rectangular bars with dimensions of $6 \times 17 \mu\text{m}$ and $4.6 \times 29 \mu\text{m}$, respectively. The remainder were circular and had diameters from 2.8 to 4.8 μm . We believe that slight cracking of the crystal around the smaller of the triangular inclusions occurred during the third shot at 2.6 J/cm² and that these cracks expanded slightly during the next two shots at 3.3 and 3.4 J/cm². Fractures were obvious around this inclusion, the larger of the bar-shaped inclusions, and one of the small circular inclusions after a sixth shot at 4.3 J/cm². Except for slight growth of the damage surrounding the triangular crystal, there was no accumulation of damage during the last 35 pulses. Data from this test are shown in Fig. 7-58. During this sequence of irradiations, 24 of the 27 inclusions did not cause damage.

These data indicate that thresholds for damage at iridium inclusions in GSGG are $\sim 2 \text{ J/cm}^2$ for 9-ns pulses and not more than 4 J/cm² for 50-ns pulses; most of the inclusions survived the 50-ns irradiation at 4 J/cm². These thresholds are about half as high as those observed in similar tests of platinum inclusions in laser glass. (See "Measurement of Damage Initiation and Growth due to Metallic Platinum Inclusions," in Sec. 6.)

Linear Optical Crystals. We used 1- and 9-ns, 1.064- μm pulses to test three samples of the high-purity ("fluorescence-free") CaF₂ now being produced by Harshaw Chemical Corporation. The growth process was designed to eliminate the fluorescence sometimes observed in CaF₂ crystals irradiated by UV light. The failure mode in all three crystals was damage at isolated fracture sites in the bulk. Two were tested with 1-ns pulses, and their thresholds were 5.0 ± 0.7 and $7.9 \pm 0.8 \text{ J/cm}^2$, respectively. One, tested with 9-ns pulses, had a threshold of $10.0 \pm 1.5 \text{ J/cm}^2$. Two Harshaw crystals of high-purity MgF₂ were also tested with 9-ns, 1.064- μm pulses. Both failed at isolated bulk sites. Their thresholds were 14 ± 2 and $20 \pm 3 \text{ J/cm}^2$, respectively.

The failure of these high-purity crystals by bulk damage is disappointing, because they are prototypical of materials with strong electronic bonds and low refractive indices that should have high resistance to damage. We test a few fluoride crystals

each year to maintain a current evaluation of these materials, but have not yet found a crystal whose short-pulse, 1.064- μm bulk damage threshold exceeds its surface damage threshold. Therefore, applications requiring use of fluorides at high fluence may require identification and elimination of the defect structure responsible for the bulk damage.

Nonlinear Optical Crystals. Since the problem of isolated bulk damage in crystals of KDP has not been solved, we tested samples of several nonlinear materials in hopes of finding something with much higher damage resistance than KDP. In single-shot tests with 1-ns, 1.064- μm pulses, we observed isolated bulk damage at 11 J/cm² in a crystal of MgO:LiNbO₃, at 4 J/cm² in a crystal of LiIO₃, and at 8 J/cm² in a crystal of L-arginine phosphate (LAP). A second LAP crystal did not suffer bulk damage in tests at up to 15 J/cm², the threshold for surface damage. A crystal of strontium barium niobate (Sr_xBa_{1-x}Nb₂O₃) exhibited surface damage at $8 \pm 2 \text{ J/cm}^2$ and was free of bulk damage at that fluence. Multiple-shot tests involving large numbers of shots were not performed on any of these materials. The bulk damage threshold for the better of the two LAP crystals is equal to the highest threshold measured in 265 KDP crystals. Because of this encouraging result, we are continuing to investigate LAP as a potential harmonic conversion material for fusion applications.

Tests of CaF₂ for the USAF.⁷² In preparation for evaluation of AR coatings deposited from solutions onto polished surfaces of CaF₂ crystals, we measured surface and bulk damage thresholds for 11 bare polished crystals of CaF₂. The thresholds were measured with 25-ns, 0.351- μm XeF laser pulses. Each test site received 1000 shots delivered at 25 or 100 Hz. Three of the samples were made from standard UV-grade CaF₂ and eight from Harshaw fluorescence-free CaF₂.

The eight crystals of fluorescence-free material had similar damage resistance, and the small differences observed were probably caused by insufficient sampling. Therefore, we believe the union of all the data provides the best representation of this material. The damage that occurred during tests of 139 sites in these samples

Fig. 7-59. Damage observed in tests of 139 sites in 8 fluorescence-free CaF_2 crystals. Test conditions: 0.351- μm , 25-ns pulses; 1000 shots per site at 25 or 100 Hz.

Fig. 7-60. Damage observed in tests of 46 sites in 3 UV-grade CaF_2 crystals. Test conditions as in Fig. 7-59.

is summarized in Fig. 7-59. Although some sites survived irradiation for 1000 shots with fluences of 20 J/cm^2 , other sites exhibited bulk or surface damage at fluences as low as 10 J/cm^2 . The randomness of the data indicates that defects were responsible for the damage and that the median separation of these defects was similar to the beam dimension, $\sim 1.2 \times 0.3 \text{ mm}$. In tests with large-area beams this material

would not survive at fluences greater than 10 J/cm^2 .

Damage in the three UV-grade CaF_2 crystals is summarized in Fig. 7-60. These crystals exhibited surface damage at single-pulse fluences of about 10 J/cm^2 , but the damage at that threshold was a spatially uniform alteration of the surface. Bulk damage was not observed. We believe that the surface damage resulted from contaminants

introduced during surface fabrication, so it should be possible to raise the thresholds by etching the surfaces.

For all 11 crystals, thresholds were the same at 25 and 100 Hz. This indicates that the crystals were damaged by the high fluence in single pulses, not by average-power loading. This is not surprising, however, because the average-power loading for 10-J/cm² pulses delivered at 100 Hz was only 1 kW/cm².

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Nonlinear Optical Materials

Introduction

Advanced lasers for fusion research will benefit from the development of higher-

damage-fluence materials for nonlinear frequency conversion and of rapid and inexpensive methods for the production of large-aperture crystals. In 1985 we continued our work on rapid growth of KDP by growing crystals in rectangular plates with near-phase-matched orientation at growth rates 10 to 20 times higher than commercial production rates. The key to achieving such high growth rates is providing adequate mass transport to the growth surface. The damage thresholds of these KDP crystals are comparable to those of commercial crystals. Because of our interest in higher damage thresholds, we conducted KDP growth experiments in which key growth parameters were systematically varied; we will determine whether correlations exist between the damage thresholds and the growth parameters.

Variable	Current industry	LLNL goals	Cost reduction factor
Growth rate	1-1.5 mm/d	15 mm/d	10
Utilization	0.2-0.4	0.7-0.8	2.5
Yield	<0.55	>0.8	1.5

Table 7-6. Variables influencing crystal-growth costs, and potential cost reductions.

Fig. 7-61. Positions of type-II second harmonic generation plates in monopyramidal and sector crystal forms.

In our search for improved nonlinear crystals, we made improvements on standard powder tests for the efficiency of second-harmonic generation (SHG), particularly for fluoride materials. We measured quantities whose values are needed for calculating the nonlinear coefficients of *L*-arginine phosphate (LAP), whose nonlinear coefficient (about 3 times higher than that of KDP) and damage threshold (2 times higher) make it attractive as a possible replacement for KDP.

We developed a new interferometric method to measure optical dispersion.

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Fig. 7-62. Axial-centrifugal pump for conversion of rotational fluid flow into a downward-directed jet.

Rapid KDP Growth

Introduction. Our primary objective is to develop a process for growing KDP crystals at a producer's cost of about \$1/cm³ for unfinished plates of dimensions 35 × 35 × 1 cm, one-tenth the cost of the current industrial process. We have identified the technical limitations of the current process and the operating conditions required for faster growth.^{73,74} Our approach to obtaining cost reductions is to seek increases in growth rate, utilization, and yield, as summarized in Table 7-6.

To increase the growth rate we increase the rate of transport of solute to the crystal surface by allowing a high-speed jet of KDP solution to impinge on the crystal surface under controlled hydrodynamic conditions.

Utilization, which refers to the number of plates that can be cut from a mature crystal, depends on growth habit and on total quantity of growth. Most industrial suppliers grow KDP in the monopyramidal or bipyramidal form. Alternatively, net growth may be constrained to a single crystallographic plane or sector with mechanical barriers and directed solution flow. Such "sectorial growth" yields roughly 2 or 3 times as many type-II plates as does the monopyramidal growth shown in Fig. 7-61.

Yield refers to the fraction of growth starts that result in useful, mature crystals. While not all factors influencing yield are known, we believe that yields could be greatly improved through continuous measurement and control of solution supersaturation. Computerized process control should prevent the inadvertent temperature or concentration fluctuations that result in void defects (veils) or in spontaneous nucleation.

In 1985 we grew KDP crystals in the form of (101) sector blocks of dimensions 0.1 × 0.1 m at rates of 11 to 23 mm/d. We identified spontaneous nucleation and solution contamination as the important factors limiting long-duration, high-rate growth.⁷⁵

We derived the relation between the hydrodynamic conditions, the average mass-transport rate, and the solute concentration distribution at the surface of

KDP crystal sectors and pyramids.⁷⁶

We obtained this relation by solving the convective-diffusion problem of the wedge and cone (models for sector and pyramid growth, respectively) under axially impinging flow, with the constraint of uniform surface flux. We determined the flow speeds required for achieving high growth rates at full scale, and derived the maximum length (in the direction of flow from leading to trailing edge) for coherent crystal growth. This maximum length corresponds to the point at which depletion reduces the supersaturation of the solute

in the boundary layer at the crystal surface to zero, so that growth ceases.

Growth of Large-Dimension. (101) Sectors under Axially Impinging Flow. An axial-centrifugal pump and nozzle (Fig. 7-62) were designed to efficiently convert centrifugal fluid motion into a downward-directed jet. The pump was fabricated from polypropylene and polysulfone and was used in high-rate growth experiments on the 0.1-m scale. A three-chambered growth vessel (Fig. 7-63) was designed to provide the required heat and solute transport rates.⁷⁴ The vessel consisted of a double-

Fig. 7-63. Three-chambered system for growth of crystals from aqueous solutions at constant temperature and supersaturation.

Fig. 7-64. KDP crystal grown as a (101) sector block at 11 to 23 mm/d under impinging flow at 1 m/s.

walled stainless-steel tank whose interior was coated with polytetrafluoroethylene (PTFE) to prevent corrosion of the metal. The pump was used with the vessel in four large-scale growth experiments.

The sector block shown in Fig. 7-64 was grown by the temperature-drop method (with temperatures in the range 73 to 67°C) at rates of 11 to 23 mm/d. The linear solution speed at the pump outlet was 1 m/s. The crystal contained no obvious inclusions or cracks, and its surface showed a series of macrosteps lying parallel to the leading (upstream) edge. Growth was terminated because of uncontrolled spontaneous nucleation of KDP crystallites and growth of parasitic crystals on the walls of the vessel or the pump.

In some experiments, little growth was observed on the unshielded (100) surfaces, and the (101) growth contained fractures and voids. Subsequent analysis of the growth solutions by inductively coupled plasma electron spectroscopy found 1- to 2-ppm contaminations of Fe, Cr, and Mn. These cation impurities are preferentially adsorbed on the (100) surfaces, where they impede crystal growth. Since this amount of contamination is compatible with flaw-free growth, the deterioration of crystal morphology may be due to the buildup of an unidentified organic or anionic impurity, or to surface damage

caused by KDP crystallites entrained in the flow.

These experiments show that high-rate growth on the 0.1-m scale is feasible, and that it is sustainable at 70 to 80°C, at which the tendency toward spontaneous precipitation is less pronounced. We intend to control spontaneous precipitation through more efficient pump designs, reduced pump throughput, additional chambers for redissolving the crystallites, and careful control of supersaturation.

Convective Transport to Crystals Growing as Single-Sector Blocks or Pyramids

In our flow configuration, the leading edge of the sector (or apex of the pyramid) is most accessible to diffusion from the bulk solution. The steady-state concentration at the interface decreases from the leading edge to the trailing edge because of solute depletion in the diffusion boundary layer. Far enough from the leading edge, the surface concentration may approach the saturation value, so that supersaturation goes to zero and growth ceases. We calculate the solution flow rates required to sustain a desired growth rate over a crystal of a specified size and determine the maximum length of the crystal (in the flow direction) for coherent growth.

Growth of KDP crystals from aqueous solutions generally occurs at relative supersaturations $\sigma = (C_b - C_{sat})/C_{sat}$ of less than 0.05. (C_b and C_{sat} are the bulk and saturation molar concentrations, respectively.) Under these conditions, the diffusion coefficient D and the kinematic viscosity ν may be taken as constants, and the results of transport theory in two-component dilute solutions are applicable. Steady-state coherent growth requires a uniform flux of solute at the surface over the entire crystal plane. Newman and Smyrl⁷⁷ give a general expression for concentration distribution over a two-dimensional body (or a three-dimensional body with an axis of symmetry) in the form of a Stieltjes integral:

$$C_b - C_{surf}(x) = \frac{(D/3)^{1/3}}{\Gamma(2/3)} \times \int_0^x \frac{f_{int} R(\tau) d\tau}{\left[\int_\tau^x R \setminus R\beta(x) dx \right]^{2/3}}, \quad (38)$$

where $C_{\text{surf}}(x)$ is the molar concentration at the surface, Γ is the gamma function, j_{mt} is the mass transport flux (here, independent of distance x from the leading edge), and R is the distance of a point on the surface of the cone from the axis of symmetry ($R = 1$ for a two-dimensional body such as a sector). The function $\beta(x)$ is the partial derivative of the tangential velocity:

$$\beta(x) = \left. \frac{\partial V_x}{\partial y} \right|_{y=0} \quad (39)$$

The velocity derivative can be found from the velocity distribution near the surface of a wedge in two-dimensional streaming flow⁷⁸:

$$V_x = U(x)f'(\eta), \quad (40a)$$

where $U(x)$ is the potential flow velocity and f is a tabulated function⁷ of the similarity variable

$$\eta = y \left[\frac{m+1}{2} \frac{U(x)}{vx} \right]^{1/2}, \quad (40b)$$

where

$$m = \frac{\alpha}{2 - \alpha} \quad (40c)$$

and

$$U(x) = u_1 x^m; \quad (40d)$$

α is the wedge angle, expressed as a multiple of π radians. Differentiation of Eq. (40) at $y = 0$ yields

$$\beta(x) = u_1^{3/2} f''(0) \left(\frac{m+1}{2v} \right)^{1/2} x^{\frac{3m-1}{2}}. \quad (41)$$

Solution of Eqs. (38) to (41) yields

$$j_{\text{mt}} = KD[C_b - C_s(x)] \left[\frac{U(x)}{vx} \right]^{1/2} \left(\frac{v}{D} \right)^{1/3}, \quad (42)$$

where K is a dimensionless constant that depends on α . The (constant) flux j_{mt} may be expressed in terms of the growth rate g_z measured in the [001] direction:

$$j_{\text{mt}} = \frac{g_z d_{\text{KDP}}}{M_{\text{KDP}} \lambda^2} \quad (43)$$

Equation (42) has been written with dependent and independent variables on the right to show its similarity to the maximum transport rate obtained when the flux is allowed to vary and the surface concentration has the fixed value C_{sat} (Ref. 76):

$$j_{\text{mt}}(x) = K' D (C_b - C_{\text{sat}}) \times \left[\frac{U(x)}{vx} \right]^{1/2} \left(\frac{v}{D} \right)^{1/3}. \quad (44)$$

K' also depends on α ; Fig. 7-65 shows the variation of K and K' with $I_0 \alpha$.

These results may be used to predict the solution flow rate necessary to sustain the required average mass transport for a sector block and to predict the maximum length of crystal (in the direction of flow) for which coherent growth is possible. For each flow rate/growth rate combination, Eqs. (42) and (43) imply a critical distance L at which $C_{\text{surf}}(L) = C_{\text{sat}}$ and growth ceases. Growth at rates exceeding the critical growth rate may be expected to produce rough, flawed crystals, because the growth rate in the [101] direction cannot remain constant in the vicinity of the trailing edge. This expectation has been verified (see below).

This approach may also be used to predict the flow required for coherent pyramidal growth, although the cone is at best a rough approximation of a pyramid.

Laser Damage and Growth Kinetics in KDP.

Aqueous Crystal Growth Facility. We have established a versatile facility for the study of aqueous crystal growth at constant solute concentration under

Fig. 7-65. Coefficients K and K' vs included angle for maximum mass transport rate and uniform flux, respectively.

well-defined hydrodynamic conditions. We have used this facility to investigate KDP growth kinetics and to determine the effects of growth conditions on the resistance of the crystals to laser damage.⁴⁰

Four stations of the design shown in Fig. 7-66 were constructed for natural or sectorial crystal growth. Seeds are fixed to a pedestal that coincides with the axis of the 6-cm-diam rotating pump, which provides a steady stream of solution about the crystal at flow speeds up to 1 m/s. The tank volume (21 L) is much greater than the maximum crystal volume (a few cubic centimeters), so the concentration remains nearly constant, and supersaturation is a single-valued function of temperature. Temperature and supersaturation are programmed using an HP9816s computer and a Wahl controller.

Fig. 7-66. System for investigation of crystal-growth kinetics under strong impinging flow at fixed concentration.

Fig. 7-67. Growth rate vs relative supersaturation for $U/L = 15 \text{ s}^{-1}$.

Kinetic Studies. We studied the relations between crystal topography, growth rates, supersaturation, and flow conditions to determine the dependence of damage resistance on growth morphology and kinetics. Growth rates were measured on (101) sectors held at 45° to the direction of flow. Solutions were prepared from reagent-grade KDP, sterilized at 70 to 90°C , and passed through PTFE filters with pore sizes of $0.2 \mu\text{m}$ or 1 to $2 \mu\text{m}$. The position of the leading edge was monitored during growth using a telescope with a vernier reticle. Figure 7-67 shows growth rates vs relative supersaturation at fixed flow rates ($U/L = 15 \text{ s}^{-1}$) for temperatures corresponding to saturation temperatures of 57 to 76°C and for supersaturations of up to 0.1 . Veil-free growth was obtained at rates as high as 60 mm/d on samples of characteristic length of 21 mm .

The effect of flow rate was determined on sectors having a fixed length (21 mm) under constant supersaturation by stepwise variation of the flow rate. Figure 7-68 shows that the dependence of growth rate on flow rate was nearly linear. This indicates that even at values of U/L as high as 50 s^{-1} , conditions were still far from those under which growth is controlled by kinetics.

The assumption of a spatially uniform growth rate leads to the conclusion that concentration decreases monotonically from the leading edge of the sector (where it is equal to the bulk concentration) in the direction of fluid flow. For each critical supersaturation and flow condition (indicated by U/L), there exists a maximum growth rate at which the surface concentration at the downstream edge falls to the saturation concentration. The conditions are found by entering $C_{\text{surf}}(L) = C_{\text{sat}}$ into Eq. (42) and expressing the mass transport rate in terms of the corresponding [001] growth rate using Eq. (43):

$$g_{z,\text{max}} = \left(\frac{M_{\text{KDP}} \sqrt{2}}{d_{\text{KDP}}} \right) K D (C_b - C_{\text{sat}}) \times \left[\frac{U(L)}{\nu L} \right]^{1/2} \left(\frac{\nu}{D} \right)^{1/3} \quad (45)$$

Growth rates are plotted against supersaturation in Fig. 7-69, and the locus of points of $g_{z,\text{max}}$ from Eq. (45) is shown. Crystals grown at rates exceeding $g_{z,\text{max}}$

were highly flawed and cracked, while clear growth was obtained otherwise. Note that Eq. (45) depends only on U/L and is independent of any kinetic assumptions. Figure 7-69 also shows the predicted rate and the kinetics-limited rate (derived from our earlier model based on low-temperature data and extrapolated to these conditions).⁷³ Actual rates were scattered but were generally higher than predicted by our hydrodynamic/kinetic model. The scatter may indicate differences in surface topography or in macrostep density, which were noted when numerous crystals were grown at the same supersaturation.

Crystals grown in the sectorial configuration showed growth macrosteps lying parallel to the leading (upstream) edge. When part of the crystal face was dissolved in water, etch pits characteristic of screw dislocations appeared. At high growth rates, individual spiral hillocks were not observed. Although nucleation may occur by an edge mechanism, it is more likely that the asymmetric mass transport conditions upstream and downstream of the spiral hillock tend to cause trailing microsteps to coalesce into a macrostep, while upstream steps tend to vanish into the macrostep pattern.

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Optical Damage of KDP

Relation Between Laser-Damage Resistance, Growth Conditions, and Crystal Morphology. A total of 22 KDP crystals were grown under rigorously controlled conditions as part of a parameter study to relate damage resistance to growth conditions and crystal morphology. This work extended earlier studies to temperatures, convective flow rates, growth rates, and growth morphologies that are inaccessible to industry. Our growth systems differed from those in industry in the exclusion of glass, metals, and acrylics (all of which are unstable in KDP solutions) and used only saturated hydrocarbons, fluorocarbons, or polysulfone. Solids were removed by filtration (to 0.2- or 1- to 2- μm

Fig. 7-68. Growth rate vs U/L for fixed relative supersaturation and high temperatures.

pore sizes). The parameters varied were as follows:

- U/L (3 to 50 s^{-1}).
- Growth rate (3 to 30 mm/d).
- Temperature (55 to 70°C).
- pH (3.9 to 5.5).
- Degree of filtration.

The following additional treatments were applied:

- Cation complexation by ethylenediamine tetraacetic acid (EDTA).
- Anion complexation by tetramethyl ammonium cation.
- Defect refinement (by growth of crystals from comparatively small seeds).

The number of damage sites per cubic centimeter (at 10 J/cm^2) depended on flow speed, defect refinement, anion

Fig. 7-69. Maximum rate for coherent growth at the downstream portion of the sector vs relative supersaturation, and curves for kinetics-limited growth and for growth under mixed kinetic and transport control, according to our model.⁷³

complexation, and thermal cycling after growth. These dependences will be discussed elsewhere.

Biological Material in Crystals and Growth Solutions. Biological assays of Cleveland Crystal growth solutions (at natural pH) showed algae, fungi, and bacteria.⁸¹ Yokotani et al.⁸² reported "microbes" growing in KDP solutions between room temperature and 55°C, and found improvement in damage thresholds in samples grown from UV-irradiated solutions. UV irradiation also reduced absorption bands at 270 nm. We have identified viable *Bacillus* bacteria in double-distilled water used to prepare growth solutions, and *Bacillus* cells and spores in solutions from which our crystals were grown. We have found no evidence of viable bacteria in KDP crystals grown from contaminated solutions, however. Crystals grown from sterilized solutions did not show improved damage resistance. At the end of this article, we speculate on how semitransparent bacteria (visible only under phase-contrast microscopy) or partially carbonized organic particulates could absorb enough energy to cause crystal damage.

General Conclusions. Sectorial growth results in parallel growth macrosteps that originate at the leading edge. These steps do not collide, so entrapment of solution between imperfectly closed steps should not occur. While the expected step morphology was observed, no pronounced effect on damage incidence was evident. The use of EDTA did not refine the step dimensions: this may indicate that step-forming adsorbates in our case are not restricted to

transition-metal cations, but that they may be anions or molecules with electron-rich groups. Some samples, grown from relatively fresh solutions, showed unusually high n -on-1 damage thresholds (14 to 20 J/cm²). We will report on this research when damage-testing is complete.

Modeling of Damage Associated with Inclusions of Biological Material. We have developed a model for laser damage of KDP by hydraulic fracture resulting from pressure increases in a constant-volume cavity filled with solution and light-absorbing debris. We use the model to estimate the debris absorption characteristics required to achieve pressures capable of fracturing the crystal.

Laser light of incident intensity I_0 propagating through a solution in a cavity of characteristic length S is attenuated by an amount

$$\Delta I = I_0 \alpha S, \quad (46)$$

where α is the absorption coefficient of the solution. If the cross-sectional area of the cavity is S^2 , then the increase ΔU in the internal energy of the fluid in a cavity of constant volume $V (= S^3)$ is

$$\Delta U = I_0 \alpha V. \quad (47)$$

From the identity

$$dH = dU + PdV + VdP \quad (48)$$

and the constraint of constant volume, we can calculate the pressure increase from

$$\Delta P = \int_{P_0}^P dP = \rho(\Delta H)_\rho - I_0 \alpha, \quad (49)$$

where P_0 is a reference pressure (1 atm) and where $(\Delta H)_\rho$ represents standard tabulated values of specific enthalpy, interpolated for fixed density ρ . Figure 7-70 shows ΔP vs $I_0 \alpha$ according to Eq. (49).

The tangential stress σ_t in a spherical cavity at internal pressure P_i in an infinite body is given by $\sigma_t = P_i/2$. The fracture toughness K_{Ic} relates crack propagation to tangential stress through

$$\sigma_t = \frac{K_{Ic}}{b^{1/2}}, \quad (50)$$

Fig. 7-70. Internal pressure increase vs energy absorbed for inclusion filled with solution having absorption coefficient α . Critical pressures are indicated for fracture lengths b of 0.5 and 5 μ m.

where b is the fracture length. Hence the critical pressure for hydraulic fracture is

$$P_c = \frac{2K_c}{b^{1/2}} \quad (51)$$

For ammonium dihydrogen phosphate (ADP, which is similar to KDP, for which K_c is not known), $K_c = 2.3 \times 10^5 \text{ Pa} \cdot \text{m}^{1/2}$; although b could be of atomic dimensions (0.5 nm), it is more probably between 0.5 and 5 μm . This range corresponds to critical pressures between 200 and 650 MPa. For a damage fluence of 3 to 10 J/cm^2 , this pressure range corresponds to absorption coefficients of 60 to 400 cm^{-1} . With such low absorption coefficients, bacteria of dimensions 1 μm would absorb only 0.5 to 4% of the incident light and might be difficult to detect with conventional optical microscopy.

These results indicate that pure KDP solution (having an absorption coefficient of $\sim 10^{-1} \text{ cm}^{-1}$) cannot absorb enough energy at 1.06 μm to increase the pressure in voids to critical values. Hence voids are not sites of laser damage in crystals. Absorption by an inclusion at 1.06 μm could promote damage if the absorption coefficient were about 60 to 400 cm^{-1} , as might be the case with pigmented or partially carbonized biological material.

The hypothesis of damage by hydraulic fracture suggests a possible mechanism for laser-hardening (the increase of damage resistance following a series of sub-threshold laser pulses). Such pulses subject the crystal in the vicinity of an inclusion to temperature cycles above and below the saturation temperature, during which the cavity walls dissolve and recrystallize. Since dissolution occurs preferentially at flaws or dislocations in the crystal, such alternating phases of solution and dissolution should tend to increase the fracture resistance of the crystal surrounding the cavity.

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New Nonlinear Materials for Frequency Conversion in High-Power Laser Systems

Introduction. In 1985 we made significant progress in developing a systematic basis for engineering nonlinear materials for high-power laser systems. Our goal is to understand the relations between the structure, chemistry, and optical properties of solids, and especially to understand the correlations between the various susceptibilities and the intrinsic limits on them.

Our experimental program involves the study of known materials and a systematic search for new crystals with favorable characteristics. Therefore, in addition to developing the capabilities for making high-precision linear and nonlinear optical measurements, we have also begun a program of high-quality powder surveys. From powder data we have already discovered a pattern of structure/property relations in a large class of new materials.

One approach to material optimization we are studying is "index engineering," in which the noncritical wavelength of a nonlinear crystal is varied by chemical substitution. We are therefore looking for large classes of widely substitutable structure types and are determining their stability boundaries (or "structure field"). We have identified a new crystal, LAP, that has potential as a replacement for KDP in a number of applications.

The Powder Survey. Our survey is designed to obtain the most accurate information that can be obtained from powder samples about the harmonic-generating properties of crystals. The survey has three aspects.

1. Compounds chosen for synthesis are drawn from those known to belong to noncentrosymmetric point groups, since this is a necessary condition for SHG. The x-ray diffraction pattern of each synthesized powder is checked to see that it matches the pattern reported for that compound. The structure field can then be explored by comparing the measured powder patterns of substitutional analogs with simulated patterns calculated assuming the same structure. X-ray powder patterns also serve as a check on sample phase purity.

2. The dispersion and birefringence of microscopic (50 to 100 μm) grains are measured using a spindle-stage refractometer and an Ehringhaus compensator, instruments developed for optical mineralogy. The refractive index can be determined to the third decimal place and the birefringence to about 10%.

3. Using an apparatus designed by J. P. Dougherty of Advanced Material Technologies, Inc., we can measure the angle-averaged single-particle SHG efficiency with an accuracy of 10%, a considerable improvement from the order-of-magnitude accuracy previously common. Documenting the precision and accuracy of this measurement was an important part of our effort in 1985.

Angle-Averaged Single-Particle SHG Measurements. Assaying nonlinear materials for SHG by measuring the efficiency with which powdered samples produce second-harmonic light is a standard technique. Aside from the original work of Kurtz and Perry,⁸³ however, quantitative aspects of this technique have been largely ignored. Crude, qualitative tests have sufficed in most surveys of materials for possible development as nonlinear crystals.

We wanted to explore structure/property relations using our survey data, so we required a more refined experimental technique. We measure the intensity I of second-harmonic light generated by a

single, sparsely packed layer of graded crystalline material, and compare it with the intensity generated by some standard material. We disperse the sample in a refractive-index-matching fluid and support it on a glass slide. Microscopic characterization of the particle density and size distribution allows us to extract a per-particle, angle-averaged SHG efficiency having a precision of about 10%.

A significant advantage of the single-layer method is its relative immunity to large scattering-loss corrections arising from refractive-index mismatch. Mismatch is inescapable in birefringent materials and can lead to severe bias in measurements on multilayer samples. For example, Fig. 7-71 shows that with a mismatch of 0.1, the intensity drops by a factor of 2 in our single-layer measurements, while it drops by an order of magnitude for multilayer samples. Our apparatus was designed for maximum sensitivity, and we can detect signals 10^{-4} times as intense as those from quartz, which is a moderate harmonic generator.

As in all powder measurements, the fidelity with which the measured signal represents the true angle-averaged SHG efficiency of the sample is compromised by powder morphology. Many workers have commented on orientational bias, which occurs when preferential cleavage of grains leads to partially oriented samples.

Fig. 7-71. Second-harmonic intensity from a single-layer quartz sample in fluids of various refractive indices. Lobes indicate angular distribution of second-harmonic light scattered because of index mismatch. Dotted line shows rapid intensity falloff with mismatch for conventional multilayer measurements.

We tested for this in some cases by carefully air-tumbling quasispherical particles, which have less tendency to prefer special orientations on the slide. We demonstrated that the crystallinity of the sample, which can affect the apparent SHG efficiency by orders of magnitude, is a more severe source of inaccuracy. We have therefore emphasized synthesis techniques that produce samples with large (100- μm) single crystalline domains.

The per-particle SHG data can be analyzed in two ways. The most satisfactory use of the data is to test models of nonlinear susceptibility by comparing measured and calculated SHG efficiencies. Figure 7-72 shows such a comparison for a number of materials for which the known nonlinear coupling tensor elements⁸⁴ were used to calculate relative powder signals. In effect, this analysis tests the accuracy of our measurements. When good sample morphology is maintained as described above, the calculated and measured signals agree within a factor of 2, which is quite sufficient for identifying materials whose SHG efficiencies make them interesting.

A somewhat less satisfactory way of analyzing the SHG efficiency data is to extract "average" or "effective" d -coefficients, using the known linear index properties and approximate formulas for the powder signal. These coefficients are difficult to interpret when the symmetry of the material is such that several independent d -coefficients exist. Nonetheless, considerable insight into structure/property relations can be obtained by studying the variation of such averaged coefficients among large classes of compounds. Table 7-7 lists the effective d -coefficients derived from powder measurements on several standard materials, and known values of specific d -coefficients.

Fluorides as Nonlinear Crystals. The first large new class of materials we have studied are the fluorides, which are combinations of one or more cations with fluoride anions. For several reasons, one might expect the fluorides to yield useful materials for high-power applications. Compared with oxides, for example, the fluorides have much larger transparency regions and hence smaller intrinsic absorptions and smaller susceptibilities for parasitic Raman gain. Melt growth of

Fig. 7-72. Measured vs calculated powder signal intensities for materials with known d -coefficients and refractive indices. A least-squares fit of the data to $I_{\text{meas}} = aI_{\text{calc}}$ gives $a \approx 1$, as shown.

Table 7-7. Estimated and known nonlinear parameters.

Material	d_{est} (pm/V)	d_{lit} (pm/V)
NaClO_4	0.61	$d_{36} = 0.71$
LiKSO_4	0.85	$d_{33} = 1$
		$d_{31} = 0.6$
Urea	1.17	$d_{14} = 1.4$
$\text{LiClO}_4 \cdot 3\text{H}_2\text{O}$	0.32	$d_{33} = 0.40$
		$d_{31} = 0.35$
LiIO_3	5.76	$d_{33} = 7.0$
		$d_{31} = 7.1$

fluoride crystals is facilitated by their generally low melting points (<1000°C), and there are many water-soluble fluorides. Two phasematchable fluorides, BaMgF_4 and BaZnF_4 , have been characterized.⁸⁵ They have small nonlinearities but are noncritically phasematchable near 1.06 μm . It is reasonable to look for other, more nonlinear fluorides with similar phasematching properties.

The low dispersion of the fluorides implies that only a small birefringence is necessary to obtain phasematching in the visible or IR. This in turn implies that the fluorides will generally have large acceptance angles, even for critical phasematching. In fact, the dispersion of many fluorides is so low that even the minimum coherence lengths are expected to be macroscopic—50 to 100 μm .

Long coherence lengths are apparent in the results we obtained in phasematching tests in which the per-particle SHG was measured as a function of particle size. When the characteristic dimension r exceeds the minimum coherence length, we expect⁸³ to find $I \propto r^2$ for nonphasematchable materials, and $I \propto r^3$ for materials that have phasematching directions at the incident frequency. Figure 7-73 exhibits this difference in power-law dependence among several materials whose phasematching properties are known. The

Fig. 7-73. Signatures of phase-matchable and non-phase-matchable materials. The materials represented in the three upper curves (average slope 3.3 ± 0.2) are phase-matchable at or near $1.06 \mu\text{m}$; those in the three bottom curves (average slope 2.2 ± 0.1) are not.

Fig. 7-74. SHG intensity vs particle size for several fluorides. All have slopes of ≥ 3 , but KLiBeF_4 almost certainly does not have enough birefringence to compensate for its dispersion between 1.064 and $0.532 \mu\text{m}$ and hence is not phase-matchable.

mean slopes are 3.3 ± 0.2 for the phase-matchable materials and 2.2 ± 0.1 for the non-phase-matchable materials. By contrast, Fig. 7-74 shows that several fluorides, both phase-matchable and not, have a high (≥ 3) power dependence, presumably because the particle sizes do not sufficiently exceed the minimum coherence length. Hence this standard test for phase-matching is not useful for the fluorides, and one must rely more on the measured dispersion and birefringence to decide (at the powder level) if a material is potentially phase-matchable.

On the other hand, the low refractive indices of the fluorides place a stringent limitation on the nonlinear strengths expected for these substances.⁸⁶ Miller's rule states that for most substances a typical nonlinear coupling coefficient d is expected to vary as the cube of the linear susceptibility:

$$d \propto (n^2 - 1)^3 \delta. \quad (52)$$

The intrinsic nonlinearity parameter δ , sometimes called the "Miller δ ," varies between 1 and $10 \times 10^{-2} \text{ m}^2/\text{C}$. In fact, Kurtz⁸⁷ has shown that most nonlinear crystals have values of δ near $5 \times 10^{-2} \text{ m}^2/\text{C}$; we call this a "normal" intrinsic nonlinearity. (KDP, for example, has $\delta = 4 \times 10^{-2} \text{ m}^2/\text{C}$). Thus a typical low-index fluoride ($n = 1.4$) with a normal intrinsic nonlinearity would have a nonlinearity only half that of KDP ($n = 1.5$). Conversely, a material with a moderate d -coefficient and low refractive index would represent a structurally optimized substance (in the sense that it has low n) with a high δ .

Figure 7-75 shows the properties of the fluorides in our fairly thorough survey. The figure shows that NaLaF_4 has a normal nonlinearity. The striking fact shown by this figure is that, even beyond the small values of d expected from Miller's relation, the majority of fluorides also have small intrinsic nonlinearities, with δ -values below $1 \times 10^{-2} \text{ m}^2/\text{C}$. Only NH_4F and KLiBeF_4 have high intrinsic nonlinearities.

We believe that there is a clear structural basis for this dichotomy in the nonlinear responses of the fluorides. Both NH_4F and KLiBeF_4 are derived from a parent structure known as wurtzite, and are composed of a linked network of tetrahedral units. In NH_4F , this network is stabilized by hydrogen bonding. In $(\text{K,Rb})\text{LiBeF}_4$, it is stabilized by the small radii of lithium and beryllium, which prefer tetrahedral coordination by fluorine. None of the low- δ compounds has this type of structure. Instead, their nonlinear response is often determined by small distortions of coordination polyhedra of six or more fluoride ions.

Another feature of the optical properties of fluorides can be similarly explained. The birefringence of both of the high- δ crystal types is very low (~ 0.001), while higher birefringences (0.005 to 0.01) are

observed among some of the low- δ types. This is consistent with the general observation that tetrahedral networks with undistorted tetrahedra have very isotropic linear dielectric properties. From a practical point of view, it is disappointing that the materials that have the highest intrinsic nonlinearity are probably not phase-matchable because their birefringence is so small.

Table 7-8 lists the three possibly useful fluoride crystals we have found. Two of them appear to have noncritical phase-matching wavelengths in the visible or near-IR.

Sodium lanthanum fluoride (NaLaF_4) has the strongest nonlinear response, but suffers from the disadvantage of belonging to a point group ($\bar{6}$) for which the nonlinear coupling vanishes in the noncritical configuration.

A very large class of water-soluble compounds with small nonlinearities was discovered in the A_2BF_6 structure family. Figure 7-76 shows the structure field for this group. We have demonstrated that the substitution of Rb and NH_4 for K increases the birefringence in the hexagonal phases, and hence that it is possible to tune the noncritical wavelength in these compounds. The structures of the zirconium compounds labeled with an asterisk in the figure were not determined, but they are optically biaxial and have large positive birefringences. The nonlinearity of all these fluorides is about the same as that of BaMgF_4 and BaZnF_4 , which were the only fluoride harmonic generators known before our study.

Finally, while its small birefringence precludes its use as a harmonic generator at visible or near-IR frequencies, KLiBeF_4 may be a useful electro-optic material for applications in which very low dispersion or very low loss is required.

Other Materials. Although our most complete study in 1985 concentrated on fluorides, we obtained preliminary information on many other interesting materials. We surveyed 25 compounds with the benitoite structure type, and found them to be characterized by a low intrinsic nonlinearity. In spite of high refractive indices, no powder efficiency exceeding that of quartz was observed. The high natural birefringence of benitoite itself (~ 0.04) made this structure type interesting as a

Table 7-8. New nonlinear fluoride crystals.

Compound	Space group	d (pm/V)	Birefringence
NaLaF_4	$\bar{6}$	0.5	0.02 (+)
A_2BF_6	$\text{P6}_3\text{mc}$	0.06	0.005-0.001 (+)
KLiBeF_4	P6_3	0.5	0.001 (+)

▲ Fig. 7-75. "Surprisal plot" showing the nonlinear properties of the fluorides. Only two structure types, NH_4F and KLiBeF_4 , have very high values of d ; most of the rest have very low values.

Fig. 7-76. Structure field of the A_2BF_6 compounds.

possible source of phasematchable materials. However, our preliminary conclusion from measurements of SHG vs size is that the dispersion of our natural benitoite sample is too high to allow phasematching in the visible or near-IR. Because the morphology of these samples was not uniformly satisfactory in our initial survey, and because of the wide range of chemical substitution allowed by this structure, we plan to pursue this study more carefully.

We made the further interesting observation that the crystallographic space group P6_3 comprises a number of materials that differ structurally and yet have a

high nonlinear response. Two previously characterized crystals in this group are lithium iodate and KLiSO_4 ; KLiSO_4 is isostructural with KLiBeF_4 , which we have already identified as one of the most nonlinear fluorides. Another structure in this group is Ti_3PO_4 , whose powder signal is about 6 times stronger than that of quartz, and which has a large positive birefringence. As part of this study, A. Zalkin of LBL undertook a crystal structure determination of thallous phosphate.

Properties of L-Arginine Phosphate.

Part of our effort in 1985 was devoted to developing techniques for precise measurement of the refractive indices and nonlinear coupling tensor elements (d -coefficients) of crystals. Our purpose is to assess in detail the properties of new nonlinear materials, or of those whose properties have only been partially determined.

A material with high potential for frequency conversion in future high-power lasers is LAP, the phosphate salt of the amino acid L-arginine. Xu et al.⁸⁸ report that large crystals of LAP can be grown from solution, and that both type I and type II phasematching is possible at

$1.06\text{ }\mu\text{m}$. They observed phasematched SHG efficiencies 3 to 4 times greater than that of KDP. Tests at LLNL showed that the damage threshold of LAP was about twice that of KDP. These properties implied that LAP might be a candidate for high-power applications, and experiments were undertaken to characterize it more carefully. In particular, because the individual nonlinear tensor elements had not been determined, it was not certain that the phasematching direction chosen by Xu et al. was optimum.

The LAP crystal is monoclinic, belonging to the space group $P2_1$. It is therefore biaxial and has three unequal refractive indices. Table 7-9 shows the values we obtained. Our values (and those of Xu et al.) were obtained by fitting data obtained at visible wavelengths to a Sellmeier formula; the two sets of values agree only to the third decimal place. Since the two larger indices are nearly equal, the phasematching properties of LAP approximate those of a negatively birefringent crystal such as KDP. In particular, there is both type I and type II phasematching for SHG from $1.064\text{ }\mu\text{m}$ in LAP. Figure 7-77 shows the positions of these phasematching directions in the α - γ plane, and the crystallographic and dielectric axes. It should be noted that, because of symmetry, the nonlinear coupling for the type I position in this plane is exactly zero, but in the analogous position in the α - β plane it is nonzero. Because d_{eff} varies for both type I and type II phasematching conditions as one moves along the phase-matching curve from one plane to the other, there will be an optimum position for each type that maximizes d_{eff} . These positions depend on the values of the elements of the nonlinear coupling tensor.

Symmetry dictates that the coupling tensor for point group 2 has the form

$$d = \begin{pmatrix} 0 & 0 & 0 & d_{a\gamma\gamma} & 0 & d_{aa\beta} \\ d_{\rho\alpha\alpha} & d_{\rho\beta\beta} & d_{\rho\gamma\gamma} & 0 & d_{\beta\alpha\gamma} & 0 \\ 0 & 0 & 0 & d_{\gamma\beta\gamma} & 0 & d_{\gamma\alpha\beta} \end{pmatrix} \quad (53)$$

If the material is not too dispersive, coefficients that are related by a permutation

Table 7-9. Refractive indices of LAP.

Index	Wavelength	
	$1.06\text{ }\mu\text{m}$	$0.532\text{ }\mu\text{m}$
n_α	1.499	1.510
n_β	1.563	1.579
n_γ	1.570	1.587

Fig. 7-77. Crystallographic axes (a , b , c), principal dielectric axes (α , β , γ), and type-I and type-II phasematching directions in the α - γ plane for LAP.

of indices are nearly equal, a condition known as Kleinman symmetry:

$$\begin{aligned}d_{\alpha\eta\beta} &\approx d_{\beta\alpha\eta} \approx d_{\gamma\alpha\beta}, \\d_{\alpha\alpha\beta} &\approx d_{\beta\alpha\alpha}, \\d_{\gamma\beta\gamma} &\approx d_{\beta\gamma\gamma}.\end{aligned}\tag{54}$$

The optimum phasematching directions in LAP can therefore be determined, to a good approximation, from a measurement of the four coefficients $d_{\beta\alpha\alpha}$, $d_{\beta\beta\beta}$, $d_{\beta\gamma\gamma}$ and $d_{\alpha\beta\gamma}$.

We have made these measurements using a standard technique⁸³ in which the second-harmonic intensity generated by a wedge-shaped crystal is monitored as the crystal is translated perpendicular to an incident laser beam. The crystal orientation and the laser polarization are selected to measure a single coefficient, or a known linear combination of a few coefficients.

Figure 7-78 shows typical data from such an experiment, in which we measured $d_{\beta\beta\beta}$ (d_{22}) for LAP. The second-harmonic intensity varies periodically with wedge translation, and the period gives the coherence length for the fundamental and second harmonic for the chosen polarization. The maximum intensity represents the second harmonic generated over one coherence length in the crystal. The precision of our amplitude measurement is about 15%, while coherence lengths can be obtained to a few parts per thousand. Table 7-10 gives the four d -coefficients, and the corresponding coherence lengths, obtained in this way. These d -coefficients can be used to calculate d_{eff} for any phasematching direction in LAP.

Table 7-10 indicates a striking discrepancy between the measured coherence lengths and those calculated from the data in Table 7-9. Evidently it is not valid to extrapolate the LAP refractive indices to the IR, as we and Xu et al. have done. The discrepancy may be due to absorption bands we have found in the near-IR. These are of moderate strength (10 to 20% per cm) and arise from N-H stretching overtones. Our latest results indicate that these bands are substantially shifted by deuteration and

▲ Fig. 7-78. Second-harmonic intensity (normalized for small variations in input laser intensity) vs transverse displacement of an LAP wedge.

Nonlinear coefficient (pm/V)	Coherence length l (μm)
$d_{\beta\alpha\alpha} = 0.8$	$l_{\beta\alpha} = 3.3$
$d_{\beta\beta\beta} = 1.5$	$l_{\beta\beta} = 13.0$ (16) ^b
$d_{\beta\gamma\gamma} = 1.35$	$l_{\beta\gamma} = 21.7$ (27) ^b
$d_{\beta\alpha\gamma} = 0.85$	—

^a Experimental precision for d is $\pm 20\%$, for l , $\pm 1\%$.

^b Calculated from data in Table 7-9.

◀ Table 7-10. Nonlinear coefficients and coherence lengths l of LAP.^a

Parameter	KDP	LAP
d_{eff} (pm/V)	0.5	1
β ($\text{cm}^{-1}/\text{rad}$)	2500	4000
J_D (MJ/cm^2)	5–10	10–15
α ($1.06\ \mu\text{m}$) (%/cm)	5	10–20

Table 7-11. Frequency conversion parameters for LAP and KDP, both of which are capable of large-crystal growth.

that the absorption at $1.06\ \mu\text{m}$ can be reduced by at least a factor of 5. We plan to characterize the deuterated form of LAP more carefully. Table 7-11 gives the parameters relevant to high-power frequency conversion for LAP and KDP.

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Dispersion Measurements in Gases and Solids Using Second-Harmonic Generation

Nonlinear optics involves phenomena in which energy is transferred between light

waves of different optical frequencies. The efficiency of this energy transfer depends sensitively on the relative velocities of the interacting waves (phase matching). Accurate dispersion data are therefore essential in evaluating the efficiency of these nonlinear optical processes. In particular, when arrays of several crystals are used to increase the efficiency of optical frequency conversion, the dispersion of intervening centric (i.e., linear) materials becomes an important design parameter.⁹⁰

Typically, dispersion is determined from absolute measurements of the refractive indices at several wavelengths, so that the absolute error in the index difference Δn is the sum of the errors in the measured refractive indices. For materials with very small index variation, such as gases, this can lead to large relative errors in the dispersion.

Clearly, greater relative accuracy would be possible if the dispersion could be determined directly. We have constructed a nonlinear interferometer that directly measures the index difference of a material between a given frequency and its second harmonic,

$$\Delta n = n(2\omega) - n(\omega). \quad (55)$$

The technique is simple and gives Δn with a relative error of less than 0.1%. While the principles and utility of interferometers that make use of SHG have been discussed by Hopf and co-workers,⁹¹ this is the first systematic use of such a device to measure dispersion. We used the technique to determine the dispersions of gases of interest in high-average-power frequency conversion, and demonstrated that it can also be used for centric crystals or isotropic solids.

Theory of the Nonlinear Interferometer. The essential features of the interferometer are illustrated in Fig. 7-79, which shows the setup for measurements on gases. A laser beam passes consecutively through two

crystals of an optically nonlinear material, generating light at the second harmonic in each crystal. The intensity of the incident light is measured using a photodiode and edge filter, and the intensity of the second harmonic is measured similarly after each crystal. The relative phase between these second harmonics depends on the Δn given in Eq. (55), namely, the difference between the refractive indices, at the first and second harmonic frequencies, of the material between the two crystals. By varying the optical path length of a sample material placed between the crystals, and observing the corresponding variation in the intensity of the second harmonic after the second crystal, we may determine Δn for the material. If the material is a gas, the path length can be varied by changing its pressure in a cell of fixed length. For solids (and liquids), the path length can be varied by rotating a plane-parallel slab relative to the beam direction or by translating a wedge-shaped slab.

Application to Gases. The theory of harmonic generation shows that the second-harmonic intensity I at the final detector is given by

$$I = I_2 + I_2' + 2(I_2 I_2')^{1/2} \cos \phi, \quad (56)$$

where

$$\phi = \Delta n \frac{4\pi l}{\lambda} + \phi_0; \quad (57)$$

ϕ_0 is a constant. For low-pressure gases, we can write

$$\Delta n = \Delta n_0 \rho, \quad (58)$$

where Δn_0 is the index difference evaluated at a density of 1 amagat and ρ is the gas density in amagats. Since we do not measure ρ itself, but rather the gas temperature T and pressure P , we rewrite Eq. (58) as

$$\Delta n = \Delta n_0 \rho_0 \frac{P}{P_0} \frac{T_0}{T}, \quad (59)$$

where $P_0 = 14.696$ psia and $T_0 = 273.15$ K. For the simple gases we have measured, Eq. (59) holds to within 1 part in 10^3 for pressures up to 100 psia.

Fig. 7-79. Schematic diagram of the nonlinear interferometer. Crystals 1 and 2 are second-harmonic generators separated by a distance l . The second-harmonic field E_2 generated in crystal 1 interferes at the detector with field E_2' generated in crystal 2.

Using Eq. (59), we rewrite Eq. (57) as

$$\phi = \Delta n_0 \frac{4\pi T_0}{\lambda P_0 T} Pl + \phi_0. \quad (60)$$

If we now fit observed values of $I(P)$ to the expression

$$I(P) = a + b \cos(cPl + d), \quad (61)$$

we see from Eqs. (56) and (60) that the fitting parameter c so obtained will represent the quantity

$$c = \Delta n_0 \frac{4\pi T_0}{\lambda P_0 T}, \quad (62)$$

thus yielding Δn_0 . (The fitting parameter d represents a constant phase difference due to fixed material, such as air gaps and cell windows, between the two crystals.)

We used the technique described above to determine the dispersions of possible coolant gases for a proposed multicrystal high-average-power frequency conversion scheme. Our results are given in Table 7-12; we estimate an overall 95% confidence interval of $\pm 2\%$ for these results. Typical data and the corresponding fits to Eq. (61) are shown in Fig. 7-80 for helium and nitrogen.

Application to Solids. This method can be applied directly to centric crystals or isotropic solids as a variant of the Maker-fringe technique⁹⁴ in which the gas cell is replaced by a plane-parallel slab of the solid sample. In this case we vary the optical path l in Eq. (57) by rotating the slab, so that l is given by

$$l = \frac{l_0}{\cos \theta_{\text{int}}}, \quad (63)$$

where l_0 is the slab thickness and θ_{int} is the angle between the light beam inside the slab and the surface normal.

In this case we rewrite Eq. (57) as

$$\phi = \Delta n \frac{4\pi l_0}{\lambda \cos \theta_{\text{int}}} + \phi_0, \quad (64)$$

fit $I(\theta_{\text{int}})$ to the expression

$$I(\theta_{\text{int}}) = a + b \cos\left(\frac{cl_0}{\cos \theta_{\text{int}}} + d\right) \quad (65)$$

Gas	Δn (in 10^{-7})	
	Measured	Calculated ^a
He	2.160	3.6
Ne	4.030	4.7
Ar	40.04	40.6
Kr	82.05	79.8
Xe	192.1	189.5
H ₂	28.06	28.6
N ₂	42.84	42.9
O ₂	51.62	51.2
CO	73.61	71.6
CO ₂	90.55	78.7
CH ₄	94.40	—
Air	44.29	44.30

^a Except for air, data for calculated values are those of Cuthbertson and Cuthbertson in corresponding volumes of Ref. 92; data for air are from Ref. 93.

Fig. 7-80. (a) Second-harmonic intensity vs pressure for helium in a 112.48-cm cell. The solid line is the best fit to Eq. (60). (b) Second-harmonic intensity vs pressure for nitrogen in a 10.41-cm cell, with the best-fit curve. The noise in (a) is typical, and is largely due to laser-pointing instability.

and interpret the fitting parameter c as

$$c = \Delta n \frac{4\pi l_0}{\lambda}. \quad (66)$$

Figure 7-81 shows results of applying this version of the technique to a plate of crystalline lithium fluoride ($n = 1.39$, $l_0 = 0.198$ cm) and to a glass slide ($n = 1.512$,

Table 7-12. Dispersion $\Delta n = n(2\omega) - n(\omega)$ of possible coolant gases for a proposed multicrystal high-average-power frequency conversion scheme.

Fig. 7-81. (a) Maker fringes for a plate of lithium fluoride. The solid line is the best fit to Eq. (65) and corresponds to $\Delta n = 0.0069$. The discrepancy between data and fit near $\theta = 0$ may be due to neglect of reflection losses. (b) Maker fringes for a glass plate. The least-squares fit gives $\Delta n = 0.0133$.

$l_0 = 0.0940$ cm). Although the quality of the nonlinear least-squares fit is not very good, we deduce a value of $\Delta n = 0.0069$ for lithium fluoride, in good agreement with the value 0.0072 calculated from published refractive index data. For the glass slide we find $\Delta n = 0.0133$, which is consistent with values reported for silicate glasses.

Summary. We have demonstrated a technique that is capable of measuring the difference in refractive index of a material between a given frequency and its second harmonic with an unprecedented precision of about 0.1%. We have used this technique to measure this quantity for possible coolant gases for a proposed multicrystal high-average-power frequency-conversion scheme, and have shown that the technique can be applied to centric crystals or isotropic solids.

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Advanced Crystalline Laser Materials

Introduction

The development of a solid state laser as an ICF power plant driver will require new laser media capable of high average power (HAP) output at high efficiency. Recent analyses^{95,96} suggest that it may be possible to achieve overall laser efficiency η greater than 10% with appropriately sensitized crystalline hosts and optimized pump sources. The laser's thermomechanical design can be considerably simplified by use of crystalline materials, which conduct heat more rapidly than glasses.

Much of the current interest in solid state laser materials for the ICF application results from the discovery⁹⁷⁻¹⁰⁶ of efficient sensitization of neodymium laser emission by the co-doping of neodymium and chromium in various crystalline hosts. While the lasing efficiency of these materials might meet laser fusion system requirements, the nonlinear refractive indices of present hosts are too high to allow efficient extraction of stored optical energy in pulses short enough (~ 10 ns) to be effective in igniting thermonuclear reactions. (The maximum acceptable value of n_2 is about 2×10^{-13} esu.) It appears likely, however, that efficient chromium-to-neodymium energy transfer can be achieved in a number of other crystalline media, some of which have suitably low values of n_2 . To explore this possibility, and to develop a more fundamental understanding, we have concentrated a large part of our near-term research on such systems, regardless of whether they suit our ultimate needs for fusion.

In 1985 our work on two oxide garnets, gadolinium gallium garnet (GGG) and scandium-substituted GGG (GSGG), taught us much about the need for low-loss laser materials and the importance of high-temperature chemistry in the growth of good-quality crystals. For example, the valence of the chromium ion in these crystals was found to be unstable in the presence of certain

impurities. This produces undesirable spectroscopic characteristics, which can only be avoided through careful control of impurity concentrations. In addition, a number of defects commonly produced in Czochralski crystal growth have adverse effects on laser performance. As a result, a project to develop improved crystal-growth technology has been established at LLNL. The next article gives a more detailed discussion of the results of work on oxide garnet laser crystals.

Crystals with octahedral cation sites and with low electrostatic field strength at these sites offer the best prospects for efficient chromium-to-neodymium energy transfer. We have continued our search for new materials of this type, and have performed more intensive studies on some such materials that were already known. Studies of the spectroscopic properties of $\text{LaMgAl}_{11}\text{O}_{19}$ (lanthanum magnesium hexaaluminate, or LMAO) doped with chromium and neodymium, and of strontium aluminum pentafluoride (SrAlF_5) doped with chromium, showed that chromium enters a number of different crystallographic sites in each crystal. Detailed knowledge of the radiative and nonradiative properties of the different sites is needed for quantitative assessment of the potential of these materials for laser applications. The final article in this subsection summarizes the results of our work on LMAO, SrAlF_5 , and other crystal laser candidates, and of our work in the synthesis of new crystalline materials for use in these studies.

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Oxide Garnet Crystalline Laser Materials

Spectroscopic Studies on GSGG:Cr,Nd.

Although GSGG has a high nonlinear refractive index n_2 (estimated at 8×10^{-13} esu), which makes it unsuitable as a fusion-laser medium, it may be used as a model system for such an application. Therefore, we have determined several of the spectroscopic properties of GSGG that

are of prime importance in designing a laser, including the following:

1. The absolute spectral absorption cross section vs wavelength, which characterizes the opacity of the material to pump radiation.
2. The stimulated-emission cross section, which determines the saturation fluence and the wavelength dependence of the small-signal gain coefficient (for a specified stored energy density).
3. The dependence of the fluorescence decay time on neodymium-ion concentration, which is needed to determine the concentration that yields the highest efficiency.

Properties (1) and (2) were also studied for chromium-doped GSGG, to maximize the benefits of co-doping GSGG for higher lasing efficiency. A more detailed discussion of the experimental procedures and of the theoretical expressions derived to interpret the data have been published elsewhere.¹⁰⁷

The distribution coefficient k , defined as the ratio of the concentration of an element in the crystal to the concentration in the melt, has been determined for Cr^{3+} and Nd^{3+} in GSGG to be 1.00 ± 0.05 and 0.65 ± 0.05 , respectively.¹⁰⁷ A nonunity k results in a concentration gradient in the crystal. Fortunately, for Nd^{3+} in GSGG, k is high enough that the concentration gradient will not be too great in the crystals that will be used in HAP laser systems.

The absorption spectrum of GSGG:Nd^{3+} , in terms of absolute spectral cross section, is shown in Fig. 7-82. These data can be used to predict the radiative lifetime of the $^4\text{F}_{3/2}$ excited J manifold and the branching ratios of fluorescence transitions to the lower-lying $^4\text{I}_J$ manifolds by applying the Judd-Ofelt theory.¹⁰⁸⁻¹¹⁰ The radiative lifetime so obtained is $284 \mu\text{s}$, quite close to the $278\text{-}\mu\text{s}$ value we found for GSGG doped with trace amounts of Nd^{3+} . The predicted branching ratio for the laser transition ($^4\text{F}_{3/2} \rightarrow ^4\text{I}_{11/2}$) is 0.51.

The stimulated-emission cross section may be determined from the fluorescence spectrum if the radiative lifetime and intermanifold branching ratios are known.¹⁰⁷ The gain spectrum, in terms of normalized cross section, is shown in Fig. 7-83. The peak cross section is

Fig. 7-82. Spectral absorption cross section of Nd^{3+} in GSGG at 296 K.

Fig. 7-83. Normalized $^4\text{F}_{3/2} \rightarrow ^4\text{I}_{11/2}$ spectral emission cross section of Nd^{3+} in GSGG at 295 K.

Fig. 7-84. Concentration dependence of the exponential fluorescence lifetime of Nd^{3+} in GSGG at 295 K.

$1.3 \times 10^{-19} \text{ cm}^2$, close to the value we determined.¹¹¹ The homogeneous saturation fluence for 1.06- μm laser emission is $hc/\lambda\sigma = 1.4 \text{ J/cm}^2$. To extract the energy stored in the laser medium efficiently, the beam fluence must be about twice this

value. Since the damage fluence of GSGG ($\sim 20 \text{ J/cm}^2$) (Ref. 112) is much higher than the saturation fluence, efficient energy extraction should be achievable.

The concentration dependence of the exponential fluorescence lifetime of the $^4\text{F}_{3/2}$ manifold is shown in Fig. 7-84. Since the quantum efficiency $\eta = \tau_i/\tau_r(^4\text{F}_{3/2})$ of neodymium in GSGG is essentially unity for Nd^{3+} concentrations below about $5 \times 10^{19} \text{ cm}^{-3}$, the decrease in lifetime with increasing concentration is entirely due to nonradiative relaxation. The fluorescence lifetime τ_i is given by

$$\tau_i^{-1} = p_r + p_{nr}, \quad (67)$$

where $p_r = \tau_r(^4\text{F}_{3/2})^{-1}$ is the radiative decay rate and p_{nr} is the nonradiative decay rate. In Dexter's theory of nonradiative ion-pair relaxation,¹¹³ p_{nr} depends on the mean separation between the energy donor and the energy acceptor. In the case considered here the donor and acceptor are equivalent, and the separation distance may be replaced straightforwardly by the concentration. For nonradiative ion-pair relaxation via the electric dipole-dipole mechanism, the rate may be written as

$$p_{nr} = p_r \left(\frac{\rho}{\rho_0} \right)^2, \quad (68)$$

where ρ_0 is the concentration at which $p_r/p_{nr} = 1$. Substitution of Eq. (68) in Eq. (67) yields, after some algebra,

$$\tau_l = \frac{\tau_r}{1 + (\rho/\rho_0)^2}. \quad (69)$$

The solid line in Fig. 7-84 is a fit of Eq. (69) to the data, with $\tau_r = 280.8 \mu\text{s}$ and $\rho_0 = 4.97 \times 10^{20} \text{ cm}^{-3}$. Thus the concentration quenching is due to electric dipole-dipole ion-pair relaxation. The same conclusion was reached by Denker et al.¹¹⁴ by a different method.

The absorption and emission spectra of Cr^{3+} in GSGG are shown in Fig. 7-85.

The room-temperature fluorescence of samples having concentrations up to $1.5 \times 10^{20} \text{ cm}^{-3}$ decayed in single-exponential fashion with a lifetime of $112 \pm 5 \mu\text{s}$, which is consistent with the results of earlier studies.^{115,116} We determined the peak emission cross section of Cr^{3+} in GSGG to be $9.2 \times 10^{-21} \text{ cm}^2$ from the spectroscopic data.

GSGG:Cr,Nd Laser Testing. As part of an effort to assess the state of the art in crystal growth, and to establish a point of reference for evaluation of other laser materials in which chromium sensitization of neodymium laser action is employed, we tested a large number of GSGG:Cr,Nd laser rods. The testing^{117,118} consisted primarily of measuring the internal losses present in the bulk material and the efficiency of the free-running (long-pulse) laser. The results are summarized in Sec. 9, in the subsection "GSGG:Nd,Cr Crystals," since this work was done primarily for the High Average Power Laser project.

The GSGG crystals tested generally had significantly higher internal losses than commercial YAG laser rods, which is understandable in view of the fact that GSGG crystal-growth technology has had relatively little time to develop. When losses in the material were accounted for, the intrinsic lasing efficiencies of a number of co-doped GSGG laser rods exceeded that of the best YAG:Nd rod by more than a factor of 2.

Spectroscopic properties of GGG(Ca,Mg,Zr):Cr,Nd. A major problem inhibiting the use of co-doped GGG in

Fig. 7-85. Absorption cross section and normalized emission cross section of Cr^{3+} in GSGG at 295 K.

large-aperture laser systems is the high distribution coefficient (~ 3) of Cr^{3+} in the material. Under such circumstances, laser slabs grown by the conventional Czochralski method would have to be cut in sections transverse to the boule axis, thus significantly limiting the maximum achievable dimensions. Another problem that limits the size of the laser slab is amplified spontaneous emission. The stimulated emission cross section of GGG:Nd is comparable to that of GSGG:Nd and is larger than we desire.¹¹⁹ One possible solution to this problem¹¹¹ is to modify the structure of GGG by substituting charge-compensating ions such as Ca^{2+} or Zr^{4+} . These ions cause line-width broadening by increasing lattice disorder. We found the emission cross section in GGG(Ca,Zr):Nd to be about half that for GGG:Nd.¹²¹

We have since studied a related compound, GGG(Ca,Mg,Zr):Cr,Nd, which appeared to have several advantages over GGG(Ca,Zr):Cr,Nd. One of these was that the segregation coefficients of the charge-compensated pairs (in this case Ca-Zr and Mg-Zr) might be closer to unity over a wider composition space than in GGG(Ca,Zr):Cr,Nd.¹²⁰ It was also believed that this composition might have a Cr^{3+} distribution coefficient closer to unity.¹²¹

A boule of GGG(Ca,Mg,Zr):Cr,Nd was loaned to us by the Airtron division of Litton Industries. The chemical composition of the top of the boule, determined by the inductively coupled plasma technique, was within 10% of the initial melt

composition. The distribution coefficients of Nd^{3+} and Cr^{3+} were 0.5 and 3.0, respectively, and their concentrations at the top of the boule were $6.2 \times 10^{19} \text{ cm}^{-3}$ and $1.2 \times 10^{20} \text{ cm}^{-3}$. These concentrations (and the corresponding distribution coefficients) are accurate to within 10%.

Table 7-13 summarizes the spectroscopic properties determined from this boule. The lower stimulated emission cross section relative to that of $\text{GGG}(\text{Ca,Zr})$ shows that the addition of magnesium increased the disorder of the lattice, as expected. Unfortunately, there was evidence of banding of the Cr^{3+} parallel to the growth interface, as occurs in $\text{GGG}:\text{Cr}$. The material studied here also solarized during UV

irradiation, as was reported earlier for material from the same vendor.¹⁰⁴ This may be a resolvable impurity problem.

In summary, $\text{GGG}(\text{Ca,Mg,Zr}):\text{Cr,Nd}$ does not appear to be superior to the simpler compound $\text{GGG}(\text{Ca,Zr}):\text{Cr,Nd}$, based on spectroscopic evidence.

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Other Crystalline Laser Candidates

LMAO. The crystals discussed in the preceding article are unsuitable for ICF applications because of their high values of the nonlinear refractive index n_2 . This is caused by the presence of cations with highly polarizable oxygen bonds. LMAO is an oxide crystal that is mechanically strong, provides a lanthanide ion site for Nd^{3+} , and yet has a manageable n_2 ($\sim 3 \times 10^{-13}$).¹²² LMAO crystallizes in the magnetoplumbite structure, with a hexagonal unit cell consisting of a spinel-like central block bounded by mirror planes containing the La^{3+} ions.¹²³ As its formula suggests, the host is primarily Al_2O_3 , which keeps the (estimated) value of n_2 low. There are three sixfold-coordinated

Table 7-13. Spectroscopic properties of $\text{GGG}(\text{Ca,Mg,Zr}):\text{Cr,Nd}$.

Property	Value
Measured neodymium lifetime	260 μs
$^4\text{F}_{3/2}$ Stark splitting	74 cm^{-1}
$^4\text{F}_{3/2} - ^4\text{I}_{11/2}$ peak wavelength	1064.1 nm
$^4\text{F}_{3/2} - ^4\text{I}_{11/2}$ peak cross section	$5 \times 10^{-20} \text{ cm}^2/\text{a}$
$\text{Cr}^{3+} \ ^4\text{A}_2 - ^4\text{T}_2$ peak wavelength	640 nm
$\text{Cr}^{3+} \ ^4\text{T}_2 - ^4\text{A}_2$ fluorescence lifetime	120 μs
Cr — neodymium transfer efficiency	60% ^b
Nd^{3+} distribution coefficient	0.5
Cr^{3+} distribution coefficient	3

^a Assumes a $^4\text{F}_{3/2} - ^4\text{I}_{11/2}$ branching ratio of 0.5.

^b For concentrations quoted in text.

Fig. 7-86. Absorption and emission spectra of Cr^{3+} in $\text{LaMgAl}_{11}\text{O}_{19}$.

Fig. 7-87. Absorption and emission spectra of Cr^{3+} and Nd^{3+} in $\text{LaMgAl}_{11}\text{O}_{19}$.

Al^{3+} sites, suggesting that they will accept Cr^{3+} substitutionally. We suggested co-doping LMAO earlier,¹¹¹ but considered it an unlikely candidate for efficient energy transfer because spectroscopy indicated that the crystal field at the Cr^{3+} site(s) was too strong to allow rapid (or efficient) energy transfer. Nonetheless, since there were no reports of Cr,Nd co-doping of LMAO in the literature, we carried out a preliminary study.¹²⁴

The particular LMAO:Cr,Nd crystal on which most of our studies were performed contained $9 \times 10^{20} \text{ cm}^{-3} \text{ Cr}^{3+}$ and $3 \times 10^{19} \text{ cm}^{-3} \text{ Nd}^{3+}$, as determined by analytical chemistry. A crystal of LMAO:Cr with a Cr^{3+} concentration of about $2.3 \times 10^{20} \text{ cm}^{-3}$ was also studied. The distribution coefficients for Cr^{3+} and Nd^{3+} were 1.2 ± 0.2 and 0.5 ± 0.2 , respectively.

The polarized absorption and emission spectra of the LMAO:Cr crystal are shown in Fig. 7-86. The positions of the absorption bands and the presence of ${}^2\text{E} \rightarrow {}^4\text{A}_2$ line emission near 700 nm indicate an intermediate-strength crystal field at the Cr^{3+} site(s), similar to the cases of alexandrite and YAG:Cr. The observed fluorescence decay was the sum of three single exponentials, with lifetimes of 0.28, 1.0, and 2.7 ms. It thus appears that Cr^{3+} is incorporated in all sixfold-coordinated Al^{3+}

sites. It does not appear, based on the absorption spectrum, that Cr^{3+} enters non-sixfold-coordinated Al^{3+} sites in this crystal.

When the Cr^{3+} ions are excited in LMAO:Cr,Nd we observe Nd^{3+} emission, as shown in Fig. 7-87. This indicates Cr^{3+} -to- Nd^{3+} energy transfer. The Cr^{3+} decay after pulsed excitation was faster in the co-doped sample than in the LMAO:Cr sample. There is also evidence that the energy is transferred rapidly, as the Nd^{3+} fluorescence peaked about 40 μs after the pulse. However, the Cr^{3+} concentration was too high for all the Cr^{3+} quenching to have been caused by nonradiative energy transfer. We have since grown another crystal having the same Cr^{3+} concentration as the LMAO:Cr crystal, and about $2 \times 10^{20} \text{ cm}^{-3} \text{ Nd}^{3+}$. By studying the decay kinetics of the Cr^{3+} and Nd^{3+} ions, we found a total energy transfer efficiency of about 45%, with 15% of the transfer occurring in one neodymium lifetime (340 μs). It appears that a LMAO:Cr,Nd laser operated in the short-pulse mode would only be slightly more efficient than the same laser without chromium sensitization.

We intend to explore other ways of sensitizing neodymium emission in LMAO. One possibility is decreasing the crystal field at the Cr^{3+} site by increasing the

lattice parameter. A mixed crystal in which Sc^{3+} or Ga^{3+} partially substitutes for Al^{3+} might accomplish this without greatly increasing n_2 . Other sensitizer ions could be tried. Ce^{3+} is readily accepted in LMAO, and emits in the UV.¹²⁵ It should be possible to dope LMAO with Ti^{3+} : this should increase the amount of light absorbed in the green, and Ti^{3+} should have a high transfer rate to Nd^{3+} . The transfer efficiency will depend on how much higher that rate is than the radiative rate.

$\text{SrAlF}_5\text{:Cr}$. Although the Nd^{3+} ion was used as the laser-active ion in most of our studies, other ions, particularly the first-row transition-metal ions, may also be used.⁹⁵ We studied chromium-doped crystals in particular, because of chromium's broad absorption bands in the near-UV and visible, and because of its emission in the near-IR in crystals with low crystalline field environments.¹¹¹ Generally, the chromium-doped crystals studied are unsuitable for ICF applications because the emission cross section or the system efficiency is too low. We have reported preliminary results on the spectroscopic properties of $\text{SrGaF}_5\text{:Cr}$.¹¹¹ More work, with better samples, revealed that this material has a serious problem, in that the fluorescence is quenched at room

temperature. This work is discussed below. We have had more success with the aluminum analog, $\text{SrAlF}_5\text{:Cr}$, which also emits in the near-IR and which produces laser emission that is tunable from 825 to 1010 nm (Ref. 126). The known spectroscopic and laser properties of this material are summarized below, and further areas of study are discussed.

Crystals of SrAlF_5 , nominally doped with 1 or 2 at.% Cr^{3+} were grown by H. Guggenheim of AT&T Bell Laboratories using the horizontal zone-melting technique. The tetragonal unit cell of SrAlF_5 has four inequivalent, sixfold-coordinated aluminum-ion sites on which Cr^{3+} can substitute.¹²⁷ The fluorescence decay exhibits a double-exponential time dependence, indicating that Cr^{3+} occupies more than one Al^{3+} site. About 80% of the time-integrated fluorescence had a lifetime of $\sim 95 \mu\text{s}$ at room temperature in the 1-at.%-doped sample. The other 20% of the luminescence had a lifetime of $\sim 40 \mu\text{s}$. Both lifetimes increased by about 10% when the crystal was cooled to 15 K; such a slight increase indicates that the decays are primarily radiative at room temperature.

The absorption and emission spectra of the 2-at.%-doped sample are shown in Fig. 7-58. The absorption spectrum exhibits

Fig. 7-58. Absorption and emission spectra of Cr^{3+} in SrAlF_5 .

a long-wavelength tail that extends beyond 900 nm and that is much stronger in the σ -polarization ($\mathbf{E} \perp \mathbf{c}$). Analysis of the spectra indicates that there is a lack of reciprocity between the absorption and emission for wavelengths longer than about 750 nm (Ref. 128), implying that absorption in this wavelength range may have a low (perhaps zero) quantum yield. Low quantum yield is also indicated by a comparison of the absorption and excitation spectra.¹²⁹ Assuming that the measured fluorescence lifetime is purely radiative, and using the fluorescence line shape, which is approximately the same for all emitting ions, we deduced a peak cross section of $2.6 \times 10^{-20} \text{ cm}^2$ for the primary emitting center ($\tau = 95 \mu\text{s}$) at 860 nm. At about 910 nm, which is a useful wavelength region for lasers for submarine communications, the cross section falls to $2.2 \times 10^{-20} \text{ cm}^2$.

SrAlF₅:Cr Laser-Pumped Laser Experiments. A series of laser-pumped laser experiments were performed to

evaluate the potential of the SrAlF₅:Cr crystal as a tunable laser medium. We were also interested in examining some of the properties of this crystal (in particular, the quantum efficiency of the chromium luminescence) that would affect its usefulness as a laser material when co-doped with chromium and neodymium.

A 1.93-mm-thick crystal doped with 2 at.% Cr³⁺ was pumped longitudinally in a nearly concentric laser resonator by a krypton-ion laser, as shown in Fig. 7-89. The pump laser light was chopped to yield a 2% duty factor to limit thermal loading of the sample. A single-plate birefringent tuner was used with two sets of high-reflectivity mirrors to vary the output wavelength. Laser emission was σ -polarized with output polarization parallel to the c-axis only (π -polarization). The variation of threshold pump power with output wavelength is shown in Fig. 7-90. In these experiments the waist diameter of the pump beam was about 100 μm ; the diameter increased somewhat with increasing

Fig. 7-89. Experimental arrangement for production of laser action in SrAlF₅:Cr.

Focal length

$L_1 = L_2 = 15 \text{ cm}$
 $L_3 = 10 \text{ or } 20 \text{ cm}$
 $L_4 = 10 \text{ cm}$

Mirror radius

$M_1 = 5 \text{ cm}$
 $M_2 = 10 \text{ cm}$

Labels

T = crystal quartz birefringent tuner
 P = thermopile power meter
 B = beam dump
 S = spectrometer/optical multichannel analyzer
 A = photodiode/signal averager
 θ = output coupling angle

Fig. 7-90. Threshold pump power vs laser output wavelength for $\text{SrAlF}_5:\text{Cr}$.

pump power. The rise in threshold at the longer wavelengths is due to decreasing emission cross section, while the rise at shorter wavelengths is primarily due to increasing absorption loss.

Output-power measurements at 910 nm were made using a silicon photodiode calibrated against a thermopile. With 1% output coupling, a slope efficiency of 14.7% was obtained. In the standard mode!¹³⁰ the slope efficiency is given by $\eta_s = \eta_0 T / (L + T)$, where T is the output coupling fraction, L is the round-trip resonator loss (including both passive losses and excited-state absorption), and η_0 is the quantum-defect-limited efficiency (71% in this case). Using this equation to estimate the loss coefficient in the laser material, we obtained a value of $10\% \text{ cm}^{-1}$.

To see if the loss coefficient was actually this high, we carried out a series of experiments to measure the linear loss coefficient and the gain per unit pump power from 890 to 970 nm. The experiments used a modified Findlay-Clay technique,¹³¹ in which the pump power threshold for laser oscillation was determined as a function of the resonator output coupling fraction. This fraction was varied by tilting the intra-cavity birefringent tuner plate as

much as a few degrees away from its usual orientation at the Brewster angle. In this situation the threshold pump power P_{th} is related to the total resonator output coupling $\sum_i L_i$ by

$$G_p P_{th} = \sum_i L_i + L, \quad (70)$$

where L_i is the output coupling fraction of the i th optical element in the laser resonator (exclusive of the laser active crystal), L is the passive insertion loss of the laser active crystal (compare this with the definition of L in connection with the equation for η_s , above), and G_p is the round-trip gain produced in the crystal per unit pump power.

The results are shown in Fig. 7-91, which plots pump power absorbed by the crystal at threshold vs output coupling fraction for output wavelengths from 890 to 970 nm. The intercepts at the vertical axis represent the insertion loss L due to the laser crystal at each output wavelength. The crystal insertion loss so deduced is plotted vs laser wavelength in Fig. 7-92. The figure also shows the π -polarized absorption spectrum of the laser crystal. The uncertainty in the baseline of this spectrum,

Fig. 7-91. Pump power absorbed at threshold for $\text{SrAlF}_5\text{:Cr}$ laser action vs resonator output coupling for selected output wavelengths.

Fig. 7-92. Insertion loss determined from the preceding figure and absorption coefficient measured with a spectrophotometer for the SrAlF_5 laser crystal.

estimated to be $\pm 0.17\% \text{ cm}^{-1}$, is indicated by the error bar attached to the curve. The measured insertion losses agree very well with the more standard absorption measurement, indicating that other contributions to the insertion loss are negligible. The apparent loss coefficient at 910 nm is about $2\% \text{ cm}^{-1}$. A loss coefficient this high would make it difficult to operate this low-gain laser by flashlamp pumping, but it is not high enough to explain the low lasing efficiency achieved by Kr^+ laser pumping.

The slopes of the lines in Fig. 7-91 determine the gains per unit pump power, G_p , achieved in the $\text{SrAlF}_5\text{:Cr}$ laser crystal at each output wavelength. These values are plotted vs wavelength in Fig. 7-93. The decrease in gain with increasing wavelength is to be expected, because of the decreasing stimulated-emission cross section. However, the gain achieved is considerably lower than expected.

In the absence of nonradiative losses and excited-state absorption, G_p should be given by

$$G_p = \frac{4\lambda_p \tau_n \sigma(\lambda_L)}{\pi \omega_0^2 h c} \quad (71)$$

Fig. 7-93. Gain per unit absorbed pump power, determined from Fig. 7-91.

where λ_p is the pump wavelength, τ_{11} is the fluorescence lifetime of the excited state, $\sigma(\lambda_L)$ is the emission cross section at laser wavelength λ_L , and ω_0 is the Gaussian radius of the pump beam at the laser crystal. The emission cross section was obtained from Fig. 7-88. Gains deduced from the Findlay-Clay measurements (Fig. 7-93) were lower than those calculated using Eq. (71) by factors of 3 to 5, depending on wavelength.

Two mechanisms that could explain the reduction of gain from the value predicted by Eq. (71) are intense excited-state absorption and absorption of pump radiation by nonradiative chromium sites. Since there are four inequivalent sites on which chromium can reside in this crystal, some chromium sites could be relaxing nonradiatively at room temperature. As mentioned above, the existence of at least one nonradiative site is implied by the lack of reciprocity between the absorption and emission spectra¹²⁸ and by comparison of the absorption and excitation spectra of a $\text{SrAlF}_5\text{:Cr}$ crystal at cryogenic temperatures.¹²⁹ The low slope efficiency achieved in the laser-pumped laser experiments may be due to the same factors responsible for reducing the gain.

These experiments indicate that $\text{SrAlF}_5\text{:Cr}$ has a system efficiency that makes it, at present, unsuitable for high power laser applications. This low efficiency is the result of either low quantum efficiency of the absorbing ions or strong excited-state absorption, together with high linear absorption losses at the laser wavelength. It might be possible, through some growth strategy, to reduce the linear losses, but the low quantum efficiency or strong excited-state absorption might still remain. The low quantum efficiency would affect the feasibility of co-doping with chromium and neodymium, since the transfer efficiency would be reduced by the competing nonradiative processes, as in co-doped glasses. Cr^{3+} excited-state absorption would not affect chromium-to-neodymium energy transfer, however, so co-doping may yet produce a viable laser material.

Further characterization of the $\text{SrGaF}_5\text{:Cr}$ samples obtained in 1985 revealed that the crystals were a solid solution of two phases of the $\text{SrF}_2\text{-GaF}_3$ system, SrGaF_5

and $\text{Sr}_3\text{Ga}_3\text{F}_{19}$. Crystals of $\text{SrGaF}_5\text{:Cr}$ exhibited the same long absorption tail as was exhibited by Cr^{3+} in SrAlF_5 ; the tail extended to about 1300 nm. The fluorescence intensity was also quenched above about 200 K. The reasons for this greatly different temperature behavior from SrAlF_5 are not known.

Other Crystals. During 1985 we contracted with Allied Signal Corporation's Material Research Center to grow oxide crystals so we could assess their potential as laser materials for ICF and HAP applications. Crystals that have possible HAP applications are discussed in Sec. 9. Three naturally occurring crystals with potential application in ICF systems were chosen for study; $\text{Ba}_2\text{TiSi}_2\text{O}_8$, CaTiSiO_5 , and $\text{Ca}_2\text{MgSi}_2\text{O}_7$. $\text{Ba}_2\text{TiSi}_2\text{O}_8$ (fresnoite) was co-doped with chromium and neodymium in the hope that Nd^{3+} would go on the Ba^{2+} site and Cr^{3+} on the Ti^{4+} site, so that the two dopants would charge-compensate one another. Bluish crystals of about 1 cm^3 volume were obtained from a slowly cooled melt. An absorption spectrum showed bands peaking at $\sim 380\text{ nm}$ and a weaker set of overlapping bands centered at $\sim 700\text{ nm}$. There was no evidence that neodymium entered the crystal. No fluorescence was observed at room temperature. It was later determined that the Ca^{2+} and Ti^{4+} sites were tenfold- and fivefold-coordinated, respectively,¹³² which may explain the lack of useful spectroscopic properties.

CaTiSiO_5 (titanite) was also co-doped with Cr^{3+} and Nd^{3+} to attempt mutual charge compensation. The absorption of crystals grown by the Czochralski technique indicated that very little of the Nd^{3+} entered the crystal, and the coloration arose from broad bands, apparently due to Cr^{3+} , peaking at about 440 and 640 nm. There was an additional, overlapping series of bands to the long-wavelength side of the 640-nm band. Excitation and emission spectra showed Nd^{3+} ion fluorescence ($\tau = 190\text{ }\mu\text{s}$), but gave no indication of energy transfer from the broad bands. Crystals grown later have shown differences in intensity between the 440- and 640-nm bands, making it doubtful that they are due to Cr^{3+} alone. The lack of Nd^{3+} sensitization, after several growth attempts, made further study unwarranted.

Table 7-14. Crystals synthesized in 1985.

One crystal of $\text{Ca}_2\text{MgSi}_2\text{O}_7:\text{Nd}$ (akermanite) was prepared. This crystal has an estimated n_2 low enough for a fusion driver ($\sim 2 \times 10^{-13}$ esu). Fluorescence has been observed, but more work on the cation-site symmetries and on possible sensitization schemes will be carried out before we proceed further with the development of this material.

Growth of about 80 crystal compositions was attempted at LLNL under the direction of H. W. Newkirk. Table 7-14 lists the crystals synthesized. Some of the crystals were considered for HAP applications and are discussed in Sec. 9.

Several themes were followed in our search for an ICF driver:

i. Continued study of chromium-neodymium sensitization in low- n_2 systems.

2. Study of other sensitization schemes such as $\text{Sm}^{2+}-\text{Nd}^{3+}$ and $\text{Mn}^{4+}-\text{Nd}^{3+}$.

3. Study of ion-pair transitions, in which the initial excited state is on one ion and the transition terminates on the other ion's excited state.

A brief discussion on findings in each of these areas follows.

Chromium-Neodymium Sensitization.

Besides the work on LMAO, already discussed, attempts were made to chromium-dope hosts that could easily be doped with Nd^{3+} . Crystals of $\text{K}_5\text{CeLi}_2\text{F}_{10}$ singly and co-doped with Cr^{3+} and Nd^{3+} were grown by the horizontal Bridgman technique. The CrF_3 appeared to have condensed in some regions of the chromium-doped crystals; the rest of the crystal was transparent. Attempts to grow chromium-doped KY_3F_{10} and $\text{KTb}_3\text{F}_{10}$ from a flux have not been successful.

Other Sensitizer-Activator Systems.

Krupke¹³³ has discussed the possibility that Sm^{2+} can be an effective sensitizer for Nd^{3+} under appropriate circumstances. This has been demonstrated in CaF_2 by Payne and Chase¹³⁴ (see "Sensitization of Nd^{3+} -doped Crystals by Sm^{2+} ," later in this section). We have attempted to co-dope SrY_2F_8 and BaY_2F_8 with Sm^{2+} and Nd^{3+} , but have only obtained samarium in the trivalent state. We believe that the divalent state can be stabilized, and plan to pursue this system further.

Since Mn^{4+} is isoelectronic with Cr^{3+} , Mn^{4+} might also be used to sensitize

Crystals grown from a flux

$\text{GdAlO}_3:\text{Nd}$

$\text{GdAlO}_3:\text{Cr}$

Crystals grown from a melt

$\text{LaMgAl}_{11}\text{O}_{19}:\text{Cr}$

$\text{LaMgAl}_{11}\text{O}_{19}:\text{Cr},\text{Nd}$

$\text{K}_5\text{CeLi}_2\text{F}_{10}:\text{Cr}$

$\text{K}_5\text{CeLi}_2\text{F}_{10}:\text{Nd}$

$\text{K}_5\text{CeLi}_2\text{F}_{10}:\text{Cr},\text{Nd}$

BaY_2F_8

$\text{BaY}_2\text{F}_8:\text{Eu},\text{Ti}$

$\text{BaY}_2\text{F}_8:\text{Nd}$

$\text{BaY}_2\text{F}_8:\text{Sm},\text{Nd}$

$\text{SrF}_2:\text{Na}$

$\text{Gd}_3\text{Sc}_2\text{Ga}_3\text{O}_{12}:\text{Nd}$

$\text{LaTa}_7\text{O}_{19}:\text{Nd}$

SrY_2F_8

$\text{SrY}_2\text{F}_8:\text{Eu},\text{Ti}$

$\text{SrY}_2\text{F}_8:\text{Sm}$

$\text{K}_5\text{CeLi}_2\text{F}_{10}:\text{Sm},\text{Nd}$

$\text{RbCdF}_3:\text{Cr}$

RbMgF_3

$\text{BaLaGa}_3\text{O}_7:\text{Nd}$

$\text{CeMgAl}_{11}\text{O}_{19}$

$\text{CeMgAl}_{11}\text{O}_{19}:\text{Cr}$

$\text{EuAlO}_3:\text{Cr}$

Nd^{3+} . Manganese is not easy to obtain in this valence state. We have attempted—unsuccessfully, so far—to grow crystals of the form LnMeF_7 , where Ln represents a lanthanide ion such as La or Ce, and where Me represents Zr or Hf. We are trying other techniques.

Ion-Pair Transition Systems. Ion-pair transition systems create a situation in which Cr^{3+} , in a crystal with a strong-field environment, together with a rare-earth ion such as cerium, which has a thermally unoccupied level, could operate as a four-level laser system. The phenomenon of ion-pair transitions was first demonstrated in chromium-doped EuAlO_3 at low temperatures.¹³⁵ No mention of the temperature dependence of this effect was made, and measurements on our sample showed only weak R-line emission. Measurements of the emission of $\text{CeMgAl}_{11}\text{O}_{19}:\text{Cr}^{3+}$ at room temperature indicated that the Cr^{3+} emission was very similar to that of LMAO:Cr. Low-temperature experiments to attempt to observe this effect are planned.

Authors: J. A. Caird and M. D. Shinn

Major Contributor: H. W. Newkirk

Multimegajoule Laser Studies

Introduction

The prospect of virtually inexhaustible controlled fusion energy has motivated the construction of high-power lasers since 1974. A sequence of ever more powerful lasers has evolved at LLNL, each

Fig. 7-94. Design of ICF research lasers has progressed rapidly at LLNL. High target gain is thought to be achievable with a 10-MJ, short-wavelength laser driver.

embodying information gained with its predecessor. The largest and most recent laser of this group is the 100-kJ Nova system, which is now irradiating targets.

Ever since the 1972 review paper¹³⁶ by Nuckolls, Wood, Thiessen, and Zimmerman, it has been clear that laser pulse energies of 1 to 10 MJ may be needed to demonstrate high gain in an inertially confined fusion plasma. Figure 7-94 indicates that a laser able to burn fusion fuel efficiently would need to deliver 3 to 10 MJ to the fusion target, as described in Ref. 137. Our current goal is to provide a conceptual design for such a multimegajoule laser.

ICF research advances since the early 1970s have been paced by driver technology.^{138,139} Laser powers and energies thought to be needed to reach significant thermonuclear gains have remained in the neighborhood of 10¹⁴ to 10¹⁵ W and 3 to 10 MJ in 10-ns shaped pulses for several years. At LLNL, we chose the neodymium-doped glass laser as the major tool for research into the laser-plasma interaction in 1971. New technologies have recently emerged that make Nd:glass a candidate for an economically feasible high-gain facility that could be operational in the 1990s. Potential improvements in solid state lasers make it possible that related

laser technology could apply to the 1- to 5-Hz driver ultimately required for commercial energy production.

The following article presents two closely related conceptual designs for a multimegajoule laser. These designs assume the use of laser components now under development. They also rely on raising the damage fluence of selected optical components by about a factor of 2, which is a reasonable goal given our active research in this field. Within the next two or three years, when our design, component, and materials work is completed, we expect to demonstrate a small prototype laser to prove the viability of an advanced solid state laser design such as those presented here.

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Multimegajoule Laser Designs

Many high-power laser architectures exist that embody what we have learned over the past ten years or so concerning saturation of the medium, parasitic oscillation, and limits on nonlinear intensity-driven phase distortion. A sampling of such architectures is shown in Fig. 7-95. The Nova laser is of the master oscillator/power amplifier (MOPA) design,^{140,141} in which the laser pulse makes a single pass through the amplifying medium. Since the largest-aperture gain medium employed in a MOPA laser is generally its most expensive component, it becomes attractive to send the beam through that medium more than once (a process called multipassing) to avoid the cost of earlier stages while efficiently extracting its stored energy.

Multipass designs vary from angle-multiplexed systems to various regenerative amplifiers. The most cost-effective designs will be those that avoid expensive components and that occupy the least volume. We are thus led to the last two designs in Fig. 7-95, the harmonic-switched regenerative amplifier and the Cassegrainian design, both of which can be front-fed (that is, the laser pulse to be amplified starts near the target). The 1984 Laser

Fig. 7-95. We have invented many new laser architectures. In 1985 we studied front-fed stackable designs such as the two harmonic-switched systems in the lower right.

Annual¹⁴² introduced the concept of stackable front-fed Cassegrainian architectures. In 1985, we added cost-saving ideas and carried out related calculations with the MALAPROP propagation code. We considered designs with arrayed multipass laser architectures using harmonic switchout, target-plane holographic injection, phase conjugation (see "Phase Conjugation for Application to Fusion Lasers," later in this section), continuous apodization, and higher amplifier storage efficiencies than in Nova. To achieve an affordable laser system design we must exploit manufacturing economies of scale, including higher-volume glass production, rapid harmonic-crystal growth, capacitor sizing and packing to handle more energy, and standardization of parts, all of which reduce costs with respect to Nova.

The amplitude, phase, and polarization of the laser field must be carefully managed in any high-power laser. Nova and its predecessors were designed to maximize output consistent with optical damage limits, and this will presumably be the case for future lasers. In these MOPA designs, beam expansion was used to keep the average flux within tolerable limits. Image

relaying ensured uniform-intensity beams on laser components, and spatial filtering removed noise of high spatial frequency. Random phase gradients in the beam were reduced by a factor equal to the magnification of each beam expander as it was traversed. Automatic alignment and centering was used to avoid clipping, which would have led to damaging fringe patterns in high-power beams. Spatial coherence was thereby maintained, permitting the beam to be focused onto targets. We seek less expensive and more efficient scalable technologies that can perform these functions.

We envision a 10-MJ laser, focusable to a 1-mm-diam spot and using a large number of almost independent beams that are clustered to reduce systems costs but between which spatial coherence need not be maintained. The system has 24 beams (12 beams per side), each of which is made up of a 6×6 array of amplifier modules with a projected element cross section of 35×35 cm. Altogether there are 864 elements or cells in 24 groups of 36. Such a system should be capable of achieving the necessary power and energy, given reasonable technological advances during the next few years. The magnitude

Fig. 7-96. Cassegrainian cell capable of delivering about 12 kJ of 0.54- μm light to a target.

of this enterprise can be appreciated by noting that it will require a capacitor bank capable of storing about 10^9 J.

The building block of the design described here is a single cell with a 35- $\times$ 35-cm aperture. Our analysis indicates that this cell, in a Cassegrainian optical geometry, is capable of achieving our aims. A potentially superior regenerative amplifier design, discussed briefly, has been studied in less detail.

Figure 7-96 is an optical schematic of a single Cassegrainian cell. It features target-plane injection using a sacrificial holographic lens, multipass extraction of a 20-disk amplifier section, phase conjugation halfway through the optical train to remove passive aberrations, a resonant acousto-optic modulator and a Pockels cell for dynamic isolation, and second-harmonic switchout.

Figure 7-97 follows the laser pulse through this cell. In Fig. 7-97(a), the input wave from the oscillator first reaches a holographic lens, which is integral with the target. This pulse illuminates a region around the target, leaving the delicate target untouched. The target could be mounted within a 3- $\times$ 3-cm square in the center of a standard 4- $\times$ 5-in. photographic plate, for example. About 100 cm^2 of film area would be left, which should be capable of directing about 0.1 J/cm^2 for a total of 10 J in 10 ns. We will need about 2 to 3 mJ input to each cell for a total of 2.6 J in 10 ns or roughly 25% of the energy

directed by the hologram toward each lens in the array. At the holographic lens, each cell will subtend an 0.76-cm-square spot whose location will be determined by the wedge-shaped final turning mirror shown in Fig. 7-96. Typical offsets between the 1ω and 2ω focal spots might be 3 to 4 cm, requiring a 3- to 4-mrad wedge angle if a 10-m focal length focusing lens is used. Since we must specify the amplitude (uniform), phase (uniform), and polarization (linear and horizontal) of the field entering each cell, we must start with a well-polarized wave with a plane phase front and establish its amplitude and phase with the hologram.

Figure 7-97(b) shows a typical IR pulse on its way to the focusing lens. The pulse passes through the front surface of the final turning mirror (which is coated for high reflectivity at 0.53 μm but for high transmission at 1.05 μm), is reflected by the rear surface, and is directed into the cell by the first turning mirror. When this several-millijoule beam reaches the dichroic mirror in Fig. 7-97(c), it encounters a 12- $\times$ 12-cm mask that occults the central third of the aperture (an area fraction of 0.12). The edges of this shadowed region, like the edges of most of the optical components in this device, are apc-ized by bead blasting. Here a Gaussian stripe 1.5 cm wide with a Gaussian radius of 0.7 cm would keep diffraction ripple to less than 3% of a plane-wave input at 60 m. A spherical dome with a radius of curvature

of 24 m and almost-parallel inner and outer surfaces, this dichroic mirror (the harmonic switchout) is designed to transmit green light but to reflect 90% of the IR energy incident on its inner surface, so that only 10% of the input field reaches the amplifier.

Once inside the cell, the beam passes through the gain region twice at large aperture, as indicated in Fig. 7-97(d). Here two options, which give roughly equal laser output, have been considered: (1) 20 laser disks of aluminophosphate glass, with a cross section of $3 \times 10^{-20} \text{ cm}^2$ and a small-signal gain of 30 per pass; (2) 20 disks of LG-660 silicate glass with a cross section of $2 \times 10^{-20} \text{ cm}^2$ and a small-signal gain of 17 per pass. After being reflected by the rear mirror, the pulse is amplified by a second pass through the gain medium. The field is still much too weak to cause frequency conversion as it passes through the type-I converter crystal. (Type-I conversion is required so as to leave the polarization state unaltered by passage through the crystal.) The beam is not altered by the acousto-optic shutter modulator, since the beam reaches this component at a null in its cycle. This time

as the pulse reaches the inner surface of the dichroic mirror, 90% of it is reflected and focused while 10% proceeds back toward the holographic lens, with an intensity no more than 10 times greater than that of the input wave.

The beam is focused through an aperture about 4 mm in diameter, as shown in Fig. 7-97(e), which spatially filters the field presented to the central portion of the amplifier on the third pass. Figure 7-97(f) shows the beam being amplified, filtered, and further demagnified. In the 4-cm-diam part of the system, the pulse traverses the Pockels cell, is amplified by a rod with a gain of 15, and is reflected by the phase-conjugating mirror, as shown in Fig. 7-97(g). Retracing its path, the beam is once again filtered twice as it extracts about 40% of the energy stored in the amplifier. Arriving once more at the harmonic-conversion crystal, which is in a relay plane, the field is uniform enough and (now) intense enough to be converted to its second harmonic [Fig. 7-97(h)], with an efficiency of 70 to 80%. The resulting green beam passes through the dichroic mirror to the target, as shown in Fig. 7-97(i). This time the beam is reflected

Fig. 7-97. A laser pulse moving through a Csesegrainian amplifier cell. The pulse is frequency-doubled when it reaches the harmonic converter on its final pass, exits the cavity through the dichroic mirror, and is focused onto the target.

Fig. 7-98. A holographic lens fixed to the target can provide a means for efficient self-aligning injection of the master oscillator pulse into each beamline.

by the front surface of the final turning mirror and is focused at the target by the lens, whose focal length is shorter for green light than for IR. Unconverted IR energy must not reenter the central portion of the amplifier. The planar rear mirror can be tilted by about 0.5 mrad to prevent this pulse from getting through the spatial filter on this pass. Since the region about 1 cm from the spatial filter aperture must now intercept one to several kilojoules, that region must be designed carefully so that debris is not ejected onto the optics.

Targeting is greatly simplified in this design: it is done off line and is totally under the control of the target fabricator who makes and attaches the hologram. Figure 7-98 shows how a holographic lens might be positioned so as to provide an efficient, self-aligning method for injection of the oscillator pulse into the beamlines. The injected beam need not map perfectly onto each lens: rather, the beams can illuminate the entire cluster of lenses. Each lens would be bead-blasted to apodize the beam passing through it so as to avoid severe diffractive ripples in the laser field presented to each cell.

After passing through the amplifier several times, the pulse is retroreflected by a phase-conjugating mirror. A simple stimulated-Brillouin-scattering (SBS) cell

may be adequate to automatically point each sub-aperture beam and produce adequate phase-conjugated energy (~ 3 J) (Ref. 143). By establishing the trajectories of the IR input rays, the hologram automatically and uniquely determines the trajectory of the green output beam. The phase conjugator ensures that the reflected IR light retraces the path followed by the input beam. Our system designs all incorporate relay spatial filters, so the phase-conjugating wave front corrector must always be located in a relay plane. Figure 7-96 shows how component spacings have been chosen to meet the relay conditions and thus to map the uniform input near-field beam, as it is first presented to the harmonic converter, into the plane of the phase-conjugate mirror and then back into the plane of the harmonic converter on the final pass. The phase conjugator is therefore only required to correct pointing within the $\pm 200\text{-}\mu\text{rad}$ field of view of the spatial filters. On the final transit through the harmonic generator, about 70% of the 2-GW/cm^2 IR beam should be converted to green, and the final turning mirror and focusing lens then deliver the green pulse to the target.¹⁴⁴ We have set a nominal tolerance of $\pm 50\text{ }\mu\text{rad}$ on the pointing accuracy of the laser beam to ensure efficient

second-harmonic conversion in the crystal. This corresponds to a transverse alignment tolerance of ± 0.5 mm at the target chamber center, assuming 10-m focusing optics. Target-diagnostic alignment tolerances on Nova are often about five times smaller than this, so we assume that target diagnostics will probably call for more accurate target placement than will the multimegajoule laser system itself.

An analogous method for target alignment could be used in a fusion reactor. A reactor target mounted in the center of a holographic lens could be spin-stabilized as it is launched toward the center of the reactor vessel. Once the target was near the center and within the laser's field of view, an oscillator pulse could launch the beam from the target. In the few microseconds needed for the pulse to transit the amplifier and be frequency-converted, the target would have moved less than 1 mm in an easily predicted direction. Of course, the alignment of this spinning target to the required precision to direct the laser beam would be a challenging problem.

Beam-centering tolerances are not as strict as those in Nova, because of the use of continuous apodization. There will be almost 100 m between the target focusing lens and the dichroic mirror, so a relay telescope (without a filtering focal-plane

aperture) will be needed. The single 10- to 20-m focusing optic shown in Fig. 7-96 can be made up of three lenses or of a combination of lenses and curved mirrors. Only simple lenses have been considered so far, and centering tolerances of ± 0.5 cm on all apertures appear adequate. Should this tolerance be exceeded, the result would be a gradual reduction of power delivered to the target, but no significant increase in risk of laser damage.

Figure 7-99 shows the acousto-optic isolator, an essential component in this laser system design, and gives a sample of the performance of a subscale prototype at 10 MHz. An acoustic standing wave in fused quartz can provide a fast beam shutter to suppress amplified spontaneous emission (ASE) and retroreflected pulses. The range of frequencies over which this test shutter can provide dynamic isolation is determined primarily by the lithium niobate transducers. Cassegrainian design dimensions call for 6.25-MHz resonant dynamic isolation, well within the demonstrated capability of this shutter. A pulse that arrives at the modulator at some time other than one of its nulls will find a phase plate that adds a sinusoidal phase to the field. The pulse will then be diffracted into higher orders and will not pass through the spatial filter pinhole. A

Fig. 7-99. An acoustic standing wave can provide a fast beam shutter to suppress ASE and retroreflected pulses.

Pockels cell in the small-aperture portion of the beamline on the right hand side of Fig. 7-96 could prevent the buildup of oscillation that circulates at the modulation frequency (as in a mode-locked oscillator). Since the phase-conjugating mirror should provide very little feedback unless presented with an intense beam, the Pockels cell may not be necessary.

Laser system simulation codes such as MALAPROP¹⁴⁵ have been used to study the beam profile to be expected from a

Fig. 7-100. MALAPROP prediction of the intensity structure in a 12-kJ laser pulse from a Cassegrainian cell with Nova-quality optical components.

Fig. 7-101. Second-harmonic energy output vs first-harmonic energy into a Cassegrainian cell. With 2-cm-thick type-I crystals of potassium di-deuterium phosphate (KD*P), cell output is virtually independent of input energy.

single cell. Imperfections in the optical components are modeled by introducing randomly placed obscurations with sizes and densities typical of dirt and damage sites in operating glass lasers. The diffraction ripples produced by these imperfections are amplified by small-scale nonlinear self-focusing of the high-power beam. Starting with an ideal plane-wave input to the holographic lens, MALAPROP predicts that an alumino-phosphate or silicate Cassegrainian amplifier will produce a 12-kJ green beam with an intensity profile similar to the one shown in Fig. 7-100. The figure shows a calculated horizontal line scan from the center to the right edge of a typical toroidal beam as it would appear after conversion to green. This beam must be projected to the target focusing lens without degradation. The amplitude fluctuations shown depend on the power, the density of scatterers in the optical train, and the phase error built into each part. The Nova baseline obscuration fraction (5×10^{-5} of the area of each optical component) was assumed for this calculation. As the calculation proceeds by expressing the electric field seen by each component on a 512×512 grid, single pixels (about 1/500th of the beam diameter) are blanked at random until the required obscuration fraction is achieved on each surface. Reference 146 contains a sampling of beam profiles calculated for 15 locations in the cell. Preliminary calculations including random phase errors suggest that this design could tolerate optical components with 6 to 12 times more rms phase error than the Nova optics.

The energy delivered by each cell depends on the input energy and on the gain and loss in the cell. Energy on target should vary slowly with input energy, as shown in Fig. 7-101. The 2-cm-thick type-I crystals assumed for this simulation make each Cassegrainian cell insensitive to the input energy. Temporal pulse shaping is most easily accomplished by timing the arrival of the 864 individual pulses at the target. Even if the phase-conjugating reflector were replaced by a plane mirror and all components were assumed to have the maximum amount of passive aberration allowed for a Nova optic, each cell's beam could still be focused to a $1/e$ intensity radius of 1 mm. Centering of individual

Fig. 7-102. The harmonically switched multipass cell, another candidate architecture for a multimegajoule laser system.

beams on the target will depend on the hologram's tolerance, but alignment accuracies of 100 to 200 μm (demonstrated in Novette and Nova target tests) should be routinely achievable. This laser system should be able to deliver 10 MJ in a temporally shaped, 10-ns laser pulse to a focal spot about 1 mm in diameter.

The collimated Cassegrainian geometry outlined above has at least two potentially costly disadvantages. Square holes are needed in the doubling crystal and the rear mirror, both of which must incorporate mounts for smaller-aperture optics. A rod amplifier is used to boost the gain before and after phase conjugation, violating the principle that small-aperture components should be eliminated by multipassing the large amplifier.

Both of these disadvantages may be overcome by a multipass cell such as that shown in Fig. 7-102. The phase conjugator is now located near the focus of the relay spatial filter, and the beam traverses the disk amplifier more times to make up for the lack of a rod amplifier. Only on the final pass is the beam properly aligned for efficient frequency conversion to provide harmonic switchout. Five to seven passes before and after phase conjugation would be required to achieve enough gain to get 15 kJ out of such a cell.

Since it allows more passes, the multipass architecture seems to optimize the deliverable optical energy per dollar

using longer-lifetime laser media with lower gain coefficients. Such an amplifier has smaller cross section and therefore needs less pulse shaping to offset saturation, but it requires lower loss in the gain medium. More passes also mean greater passive phase errors to be corrected by the phase conjugator. The phase conjugator does not correct obscuration noise, so a question remains as to how many passes are tolerable with readily obtainable optics.

The multipass design needs better optical material than the Cassegrainian design, but just how much better remains to be seen. Both cell designs are optimized to deliver an average flux of 20 J/cm^2 to the frequency-doubling crystal. Because the multipass beam has 11% more area at its output than does the Cassegrainian beam of the same aperture dimensions, it delivers that much more energy per cell.

Figure 7-103(a) shows the beam path for a parallel-beam, 2×7 -pass option in more detail. (Transverse dimensions are expanded by a factor of 100 with respect to longitudinal dimensions for better visualization.) Seven passes take place before and seven passes after phase conjugation. The injected beam proceeds from pass 1 to pass 7 by reflection from the tilted mirrors, and is then deflected near the focal area toward the phase-conjugating mirror. Figure 7-103(b) shows an expanded view of the focal region. (The transverse scale

is expanded by a factor of 10.) The conjugate beam retraces its path and exits the amplifier cavity after harmonic conversion as in the Cassegrainian design. The arrangement shown in Fig. 7-103(a) places the seventh pass on the outside for easier insertion of a beam-deflecting element, provides maximum angular separation of passes 13 and 14 to reduce harmonic conversion on the next-to-last pass, and reduces vignetting by overlapping six passes on three trajectories. Vignetting can be eliminated by using slightly converging beam paths (± 1 mrad). In this option, the

harmonic converter would also serve as a weak lens with a focal length of about 100 m. The beam within the harmonic converter would be collimated, while in the amplifier it would be slightly converging. The relay lens would have to be moved by about 1 m to compensate for the converging beam and maintain a reimaging cavity. The phase conjugator would have to cope with the astigmatism accumulated during seven converging passes through the disk amplifier.

A wide variety of related designs exist that appear to meet the needs of a 10-MJ

Fig. 7-103. (a) Ray-tracing schematic for a multipass cell, showing how the phase-conjugating SBS cell may be located near the focus of the large spatial filter. Transverse scale expanded by a factor of 100. (b) Expanded view of the focal region of the cell. Transverse scale expanded by a factor of 10.

fusion driver. Although these designs must be sorted and refined in further analysis and experimentation, we can make some general statements about them and can indicate how they can substantially reduce the cost per delivered joule of laser light.

Mechanical and electrical costs per cell scale favorably with amplifier size up to

some limit determined by handling procedures. An amplifier module containing a 6×6 array of four slabs appears to be most cost-effective. Such a device would resemble the compact, high-efficiency amplifier structure sketched in Fig. 7-104 (see "Single-Segment Amplifier: Experiments," earlier in this section, for detailed results on

Fig. 7-104. A compact, high-efficiency glass amplifier with flash-lamp banks between rectangular glass slabs.

Fig. 7-105. Laser components for a 400-kJ, 10-ns Cassegrainian amplifier occupy much less space than Nova.

a single subsection). Since 20 slabs are used in each cell, this amplifier would require five such modules in series. Adding the two $m = 3$ telescopes, the rod amplifiers, and the phase-conjugating mirrors, we end up with a beamline 48 m long, such as the one sketched in Fig. 7-105. With a length of 60 m, the multipass cell shown in Fig. 7-103 would fit comfortably within the 67-m-long Shiva high bay. An array of 36 Cassegrainian multipass cells would produce about 400 kJ of IR light in 10 ns and would fit into a significantly smaller volume than that occupied by Nova. A 10-MJ system might look like the concept shown in Fig. 7-106.

UV ($0.26\text{-}\mu\text{m}$) target irradiation can be most easily provided by frequency-doubling the green laser pulse just before it reaches the focusing lens. This avoids the severe damage risks associated with reflecting 1- to 2-GW/cm² UV laser beams. New UV lens materials will probably be required, since the sol-gel-coated fused silica lenses used in fourth-harmonic tests on Novette experienced damage at an average flux of 0.5 to 0.7 GW/cm² in 1-ns pulses.

The large number of quasi-independent beams involved in these architectures makes it essential that alignment and diagnostics be simple. Since the laser beam is injected near the target and is retroreflected by the phase conjugator, alignment becomes essentially a beam-pointing task. Continuous apodization by

bead-blasted edges on most optical components makes these cells self-centering and ensures high fill factors. The holographic optics appear straightforward; in fact, related designs are used in grocery store scanners.

New technologies make multimegajoule glass lasers economically feasible. New laser architectures using harmonic switchout, target-plane holographic injection, phase conjugation, continuous apodization, and higher amplifier efficiencies have been devised. A plan for a multimegajoule laser that could be built for a recurring cost under \$300 million relies on manufacturing economies of scale. These include continuous-pour glass production, rapid harmonic-crystal growth, switching of large blocks of power using larger capacitors packed more economically, and using a large number of identical parts.

Authors: K. R. Manes, W. F. Hagen, and J. F. Holzrichter

Phase Conjugation for Application to Fusion Lasers

Introduction. Phase conjugation is a process in which nonlinear optical interactions cause a laser beam to retrace exactly the path it followed through a distorting

Fig. 7-106. The units shown in the preceding figure can be stacked to produce a system capable of delivering 5 MJ of green light in 10 ns.

medium, correcting for the phase distortions produced by the first pass as it does so. Although phase conjugation has been studied for more than a decade,¹⁴⁷⁻¹⁵¹ it has not yet found practical application in large laser systems.

It has long been recognized that phase-conjugate correction of optical irregularities, and of pointing and alignment errors, might significantly reduce the cost of fabrication, alignment, and control of very large laser systems for fusion or other applications. To explore these possibilities, it is necessary to go beyond the impressive but poorly diagnosed demonstration experiments that make up a major fraction of the literature on phase conjugation. Phase conjugators must be developed into engineering components with well-diagnosed properties. For fusion lasers, in particular, it would also be useful if conjugators could be operated with high reflection efficiency at energies of several tens of joules. This is a much higher energy than anyone has studied in the past, although a 90%-efficient conjugator with pulse energy up to 3 J has

been demonstrated.¹⁵² In 1985 we began a study of phase conjugators directed toward these goals.

Figure 7-107 shows how a phase conjugator can be used to correct aberrations and pointing errors in a large array of laser amplifiers. A pilot beam that is exactly counterpropagating (i.e., exactly phase-conjugate) to the desired laser output beam enters an array of laser amplifiers and reads all of the optical distortions in them. The beam then enters an array of phase conjugators that exactly reverse its direction. On the second pass through the amplifiers the first-pass distortions are cancelled, yielding an output beam that is a conjugate of the input pilot beam and therefore exactly what we desire.

The phase conjugator can correct only for linear distortions. Nonlinear distortions, such as distortion due to self-focusing in the amplifiers, are different for the low-intensity first pass of the pilot beam and the amplified second pass, so phase conjugation will be significantly less useful if self-focusing is important in a laser system. Current short-pulse fusion lasers are

Fig. 7-107. Conceptual design of a large laser system using phase-conjugate correction of aberrations and misalignments. Phase conjugators allow target illumination by parallel laser chains that do not require precision alignment.

severely limited by self-focusing. Larger fusion laser systems with higher output energy and longer pulses will be much less affected by self-focusing, and at the same time must be constructed for much lower cost per joule than current lasers. Phase conjugators could therefore be very significant in these future systems. Amplifier saturation is another nonlinear effect that can affect phase-conjugate corrections.¹⁵¹

Phase-Conjugate Reflection by Stimulated Brillouin Scattering. Many nonlinear optical processes can generate phase-conjugate reflections, but the simplicity and efficiency of conjugators using backward SBS in gases or liquids makes these devices the leading candidates for practical application to large laser systems. Figure 7-108 shows a typical SBS conjugator. A lens system focuses a distorted pilot wave to an intensity at which high SBS gain occurs. The wave conjugate to the distorted pilot wave has the highest

possible gain of any SBS signal that could grow from noise in the SBS cell, so it grows and establishes in the cell a sound wave that behaves as a phase-conjugating mirror. The sound wave is moving, of course, so the reflected wave is Doppler-shifted, but the shift is only a few gigahertz and is insignificant for most purposes.

A simple SBS cell such as that shown in Fig. 7-108 does not give a perfect phase-conjugate reflection, although simple, qualitative experiments often miss this point. Careful studies at the Lebedev Institute^{151,153} show that SBS conjugation is not quite as simple as one might infer from a cursory examination of far-field patterns of the pilot and reflected waves, which may appear to be very similar. Theoretical calculations at the Naval Research Laboratory¹⁵⁴ confirm that the pilot wave and its reflection from an SBS cell can have a rather different appearance in the near field, although their far-field distributions are similar. It is very important to understand these effects in detail if we are to use SBS conjugators in practical systems.

Experiments with SBS Phase-Conjugation Cells. In 1985 we began an investigation of conjugation fidelity in a simple phase-conjugation experiment, shown in Fig. 7-109. The experiment used a frequency-doubled, single-mode Nd:YAG

Fig. 7-108. Typical phase conjugator using backward SBS in a gas cell.

Fig. 7-109. Arrangement used for preliminary phase-conjugation experiments.

Fig. 7-110. Near- and far-field intensity distributions of a weakly aberrated pilot beam and its phase-conjugate reflection. The pilot beam is intensity modulated by a cross hair in the beam, as shown in Fig. 7-109.

laser with a pulse energy of about 20 mJ in 30 ns. For convenience, some of this first series of experiments used SBS conjugation in liquid methanol rather than in a gas.

We measured incident and reflected pulse intensities with calibrated photodiodes. With an unaberrated beam, the SBS cell reflected about 80% of the incident laser intensity into a phase-conjugate return beam; for a severely aberrated beam, it reflected about 30%. These reflectivities are typical of those observed in similar experiments described in the literature.

We measured near- and far-field distributions of the incident and reflected beams simultaneously, using a wedged beam splitter to reflect all four beams to ground-glass diffusing screens and recording the four intensity distributions on standard photographic film in cameras focused on these screens. The wedged beam splitter was used so that the incident and reflected near-field images could be recorded with one camera, and the far-field images on another.

Resolution is slightly degraded by the granularity of the diffuser when this technique is used, but this loss is greatly outweighed by the advantages of the arrangement. The cameras cover a large dynamic range with great convenience. We can place bandpass and attenuating filters at the position of the camera lens rather than directly in front of the film, as would be necessary if we were to record the intensity distributions directly on film. If these filters are in an image plane, their inevitable flaws and local distortions complicate the analysis of the image. If they are far from an image plane, near the camera lens, the effect of these flaws is uniformly distributed throughout the image and causes less confusion in interpreting the results.

Problems with High Spatial Frequencies. Figure 7-110 shows sample near- and far-field distributions of a weakly aberrated pilot beam and its phase-conjugate reflection as observed in this apparatus. We inserted a cross hair pattern in the pilot beam, as shown in Fig. 7-109, to introduce a well-controlled intensity modulation in the near field of the pilot beam. The distance from the cross hair to the SBS cell was set equal to the distance from the near-field diagnostics plane to the SBS cell, so a "lensless image"¹⁵⁵ of the cross hair was formed in this plane in the phase-conjugate reflected beam. The incident pilot beam shows the usual effects of near-field Fresnel diffraction in this plane. The far-field distributions of the pilot and conjugate waves show the expected far-field diffraction pattern of the cross hair, but the diffraction fringes in the conjugate beam do not extend nearly as far from the central peak as they do in the pilot beam. This "spatial filtering"¹⁵³ of the conjugate wave is real; it manifests the imperfection of SBS conjugation.

Figure 7-111 shows qualitatively how spatial filtering arises in an SBS conjugator. The pilot beam is very weakly aberrated, so most of it comes to a sharp focus in the

Fig. 7-111. Spatial filtering in a simple SBS conjugator. High-spatial-frequency information in the pilot beam, focused outside the central focal region in the SBS cell, is not reflected if the conjugator is near threshold.

SBS cell. The focal plane of the lens is a Fourier-transform plane of the input beam, so the high-spatial-frequency (HSF) information imposed on the pilot wave by the cross hair comes to a focus well outside the focal region of the main part of the pilot beam. An SBS reflection builds up in the sharp focus for the main beam, but the HSF information that lies outside this focus is not reflected and is therefore lost, as illustrated in the far-field patterns of Fig. 7-110.

SBS conjugators have a further complexity that can in fact lead to the recovery of some of the lost HSF information if they are operated well above SBS threshold. The rays carrying the HSF information cross the main part of the beam and its conjugate just in front of the focal volume in which SBS is occurring. SBS gain occurs in this region, although it is lower because of the lower laser intensity. In this region the main part of the pilot beam and its conjugate can act as the pump beams for four-wave mixing via SBS, which generates reflected rays conjugate to the rays that carry some of the HSF information. The degree to which this occurs obviously depends on the gain in this region, and therefore on the degree to which the SBS threshold is exceeded in the main focal volume.

The spatial-filtering effect of an SBS conjugator is not necessarily a disadvantage in fusion laser systems, since we normally suppress HSF information on fusion laser beams, using spatial filters, to avoid nonlinear processes. But it is important that we know the dynamic range over which HSF suppression occurs, and that we take account of the reintroduction of HSF information for beams well over the SBS threshold.

Another effect of poor HSF information transfer in SBS conjugators can be seen in

severely aberrated beams. In Fig. 7-112, for example, a pilot beam has been severely distorted by passing it through a glass plate of poor optical quality. The divergence of the pilot beam in the far field is perhaps 0.5 mrad after this distortion. The conjugate beam is much less divergent, as expected, since it cancels these aberrations. In fact, the conjugate beam is even less divergent than it appears in the figure, since the photograph was deliberately overexposed to emphasize the weak noise components that occupy a large area around the central far-field spot. These weak nonconjugate components contain only about 10% of the far-field energy and occupy roughly two orders of magnitude more area in the far-field plane than the central phase-conjugate return beam. The local intensity in these components is therefore lower than that of the central peak by about three orders of magnitude, and it is easy to miss them entirely if the film is exposed to give an unsaturated image of the central peak. The noise in these components in the far field corresponds to random HSF modulation of the intensity of the conjugate beam in the near field; the near-field photographs of Fig. 7-112 show that this modulation is indeed present. This agrees with computer simulations¹⁵⁴ of SBS conjugators, which show considerable near-field noise due to small imperfections in the conjugate signal. Severe near-field noise will affect laser extraction and will drive the growth of nonlinear effects such as self-focusing in fusion lasers, so it is important that we understand how to design a conjugator so that this noise is minimized.

As a first step towards quantifying the conditions under which SBS conjugators begin to introduce HSF near-field noise in the conjugate reflection, we studied the reflection of a weakly aberrated pilot beam

Fig. 7-112. Near- and far-field intensity distributions of a severely aberrated pilot beam and its phase-conjugate reflection. A small nonconjugate component in the far field of the phase-conjugate beam corresponds to severe intensity modulation in the near field.

to which we added controlled amounts of astigmatism as a representative distortion. We introduced the astigmatism by tilting the focusing lens in front of the SBS cell (Fig. 7-109). The far-field distribution of the conjugate wave showed correction of the astigmatism in the central peak of the reflection for all conditions that we investigated. For distortions less than about five waves of astigmatism across the beam, the near-field distribution of the conjugate beam was quite uniform, but for more severe distortions the near field began to develop intensity striations, which correspond to the spurious HSF modulations seen in the highly aberrated beam. This gives a qualitative picture of the operating regime over which a conjugator of this sort can give a clean near-field return beam free of spurious HSF modulation.

Considerable further theoretical and experimental study will be needed before we

understand the behavior of SBS conjugators well enough that we can confidently design them into systems. These studies will proceed in 1986.

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Modeling Large-Aperture Electro-Optical Switches with Plasma Electrodes

It is possible to design longitudinal Pockels cells with large apertures and high aspect ratios using a low-pressure glow discharge as a transparent electrode medium.^{158,157} We have constructed a Pockels cell with a clear aperture of 26×26 cm using such plasma electrodes.^{158,159} We present here a detailed engineering model¹⁶⁰ for scaling the device to the large

Fig. 7-113. Four-electrode geometry for Pockels cell with plasma electrodes.

apertures (1 to 3 m) of interest in high-energy fusion laser systems.

Figure 7-113 shows the design of the plasma-electrode Pockels cell. The cell consists of two Lucite plasma-discharge chambers and a central section containing a crystal holder and internal optical windows. Outside windows on the discharge chambers allow a laser beam to pass unattenuated through the two discharge regions, internal windows, and electro-optic crystal. Typical fills are 5 to 20 Torr of He, Ne, or Ar in each chamber. For our large-aperture cell, a z-cut KDP crystal of dimensions $27 \times 27 \times 1$ cm provides a clear aperture of 26×26 cm. A silicone elastomer bonds the crystal to the Lucite holder plate and provides electrical insulation. Two sets of Viton O-rings on the holder plate allow operation of the cell with a bare KDP crystal or with thin internal optical windows adjacent to each crystal face. The 1-mm gap between each window and the crystal can be filled with an index-matching fluid to minimize optical reflection losses. The need for index matching can be eliminated by applying sol-gel AR coatings on the polished KDP

crystal surfaces. Index matching might be an advantage in a plasma-electrode Pockels cell designed for operation at high repetition rates. Use of electro-optic crystals exhibiting low optical absorption (e.g., deuterated KDP) and flowing and cooling the fluid will allow high-average-power switching.

A preionizing electrode (cathode) and a switching electrode (anode) are mounted inside each discharge chamber. The switching electrode is a stainless steel bar with dimensions of $30 \times 4 \times 1$ cm. Its corners and edges are rounded, and its surfaces are lightly polished. The preionizing electrode consists of six stainless steel buttons, each 2.5 cm in diameter and with a 3-mm slot, and operates as a segmented hollow cathode.

Electrical Performance. Figure 7-114 shows the electrical circuits for the 26×26 -cm-aperture plasma-electrode Pockels cell. These circuits create a region of uniform, high-conductivity plasma on either side of the crystal, and after ~ 290 -ns delay they apply a nominally square 100-ns voltage pulse across the plasma sheets to switch on the Pockels cell. Typical pulse

Fig. 7-114. Electrical circuit for 26×26 -cm Pockels cell with plasma electrodes.

Fig. 7-115. Switching-voltage and charging-current wave forms applied to the 26×26 -cm Pockels cell (inset), and voltage-current characteristics of one plasma electrode at a peak preionization current of 1200 A.

shapes for the switching voltage and charging current are shown in Fig. 7-115 for a charged-line voltage of 24 kV. The charging current follows the derivative of the switching voltage, as expected for a capacitive circuit. The current midway through the pulse is fixed by the termination (typically 50Ω) on the preionizing side of the cell. This termination is necessary to obtain uniform voltage and optical transmission across the crystal aperture.

Figure 7-115 also shows the anode-to-cathode voltage drop and the preionization current. The preionization voltage rises to 12.5 kV before gas breakdown. After the discharge is established at ~ 100 ns, the preionization current pulse grows to a peak value of ~ 1200 A at ~ 220 ns. The current pulse swings negative after 340 ns, because of inductive effects in the preionization circuit and the plasma-electrode current loop. At the peak of the current pulse the inductive voltage drop is zero, and we calculate a bulk plasma resistance of 3.2 to 4.0Ω per plasma electrode. As desired, the series plasma electrode resistance of 6.4 to 8.0Ω is roughly equal to the effective driver impedance of 5Ω

Fig. 7-116. Voltage-current characteristic of one plasma electrode at a peak preionization current of 150 A.

(arising from a $10\text{-}\Omega$ termination in parallel with a $10\text{-}\Omega$ transmission line).

The transverse resistance of a plasma electrode can be determined from the voltage-current characteristic of the preionization circuit when the switching pulser circuit is deactivated. Typical pulse shapes at high preionization are shown in Fig. 7-115. To minimize inductive voltage drops, we limited the preionization current to 150 A by adding a resistance in the preionization circuit. The voltage and current pulses for this case are shown in Fig. 7-116. The discharge is fully established at 200 ns, and the current pulse is fairly flat for 300 ns after that.

Fig. 7-117. Measured plasma electrode conductivity vs time at low preionization current (150 A), and estimated conductivity vs time at high preionization (1200 A).

Fig. 7-118. Plasma electrode conductivity vs transported charge determined from voltage-current characteristic in Fig. 7-116.

Fig. 7-119. Optical transmission at extreme points A and B on 26×26 -cm Pockels cell with high preionization current and bare crystal configuration.

to the charge. The slope is $2.7 \pm 0.5 \times 10^3 \text{ mho/C}$. Multiplying the transported charge by $2.7 \times 10^3 \text{ mho/C}$ gives the dashed curve in Fig. 7-117 for the low-preionization (150 A) data. If we assume that the same linear conductivity-charge relation is valid at high current, then the conductivity vs time at high preionization (1200 A) is given by the upper curve in Fig. 7-117. When the main switching voltage is applied at $\sim 290 \text{ ns}$, we deduce that the transverse plasma resistance is $\sim 17.8 \Omega$ at low preionization and $\sim 2.5 \Omega$ at high preionization. This is consistent with an inverse dependence of transverse resistance (at switching time) to peak preionization current.

Optical Performance. Quantitative optical diagnostics were performed with a 1-W argon-ion laser operated at 514 nm in an arrangement described previously.^{158,159} Typical high-preionization optical transmission is shown in Fig. 7-119 for the bare-crystal Pockels cell. A switching voltage of 6.8 kV was applied to the cell; this is somewhat lower than the half-wave voltage of 8.0 kV at 514.5 nm. The optical transmission at extreme points A and B on the aperture (the sides nearest the main and preionization electrodes, respectively) varied from 94 to 97.5%.

High voltage gradients across the crystals disappear quickly with plasma electrodes, especially with high preionization. However, small gradients of several hundred volts persist for long times in the vicinity of the hollow-cathode electrodes. We believe that this results in nonuniformities in transmission through the crystal faces near the electrodes. These nonuniformities are hard to observe for half-wave voltage operation of the Pockels cell. In our case, these effects are of about the same magnitude as those arising from transmission problems from birefringence along the KDP optical axis due to internal crystal strain. We noted this by observing the isogyre pattern with divergent laser light at specific points on the KDP crystal face. In operation, transmission uniformity of $\sim 5\%$ was achieved over most of the crystal aperture. This figure might be improved with a strain-free crystal.

Model. To simulate the performance of a large-aperture Pockels cell with plasma electrodes, we have adopted a

Fig. 7-120. Transmission-line circuit for modeling performance of Pockels cell with plasma electrodes.

one-dimensional, nonlinear transmission-line model with discrete circuit components (Fig. 7-120). We mathematically divide the electro-optic crystal surface between the switching and preionization electrodes into N_{seg} segments of equal area. The model assigns a crystal capacitance C , a plasma-electrode inductance L , a transverse plasma resistance R_t , and a longitudinal plasma resistance R_l to each segment. (Here "transverse" and "longitudinal" denote the directions parallel to and normal to the crystal face, respectively.) The loop currents are designated I_1 through I_n , with $n = N_{\text{seg}}$. The plasma resistances R_t and R_l are functions of the net charge transported through the segment, and the coupled voltage and current equations are therefore nonlinear in time. For modeling, we take both plasma electrodes as equivalent, despite the fact that one acts as a cathode and one as an anode when switching pulser voltage is applied. We assume equal preionization current through each plasma electrode, as required by the symmetry of the preionization circuits.

To develop the circuit equations we identify the node voltages $V_{n,0}$ through $V_{n,n}$ for the positively charged plasma electrode and $V'_{n,1}$ through $V'_{n,n}$ for the negatively charged electrode. The node voltage $V_{n,0}$, the voltage drop across the resistive termination R_{load} , is determined by the difference between the driver current I_0 and the first-segment loop current I_1 :

$$V_{n,0} = (I_0 - I_1) R_{\text{load}} \quad (72)$$

In terms of the driver voltage $V(t)$, with the functional form

$$V(t) = V_0 [1 - \exp(-t/\tau)] \quad (73)$$

the driver current is given by

$$I_0 = \frac{2V(t) - V_{n,0}}{Z_{\text{driver}}} \quad (74)$$

Eliminating the driver current I_0 gives the first node voltage equation,

$$V_{n,0} = [2V(t) - I_1 Z_{\text{driver}}] Z_{\text{ratio}} \quad (75)$$

where the impedance ratio

$$Z_{\text{ratio}} = \frac{R_{\text{load}}}{Z_{\text{driver}} + R_{\text{load}}} \quad (76)$$

has a value of 0.5 for an impedance-matched driver circuit. The voltage drop V_i across the capacitance of segment i is determined by the difference in loop currents I_i and I_{i+1} :

$$\frac{dV_i}{dt} = \frac{1}{C} (I_i - I_{i+1}), \quad (77)$$

and the voltage drop between adjacent nodes is

$$\begin{aligned} V_{n,i-1} - V_{n,i} \\ = V'_{n,i} - V'_{n,i-1} = L \frac{dI_i}{dt} + R_{t,i} I_{t,i} \end{aligned} \quad (78)$$

where $R_{1,i}$ is the transverse plasma resistance for segment i . The node voltages $V_{n,i}$ with respect to ground are given by the voltage drop through the crystal capacitance and the longitudinal plasma resistance:

$$V_{n,i} = 2R_{1,i}(I_i - I_{i-1}) + V_i + V'_{n,i} \quad (79)$$

It is easily seen that $V'_{n,i} = V_{n,0} - V_{n,i}$, and therefore that

$$V_{n,i} = R_{1,i}(I_i - I_{i-1}) + (1/2)(V_i + V_{n,0}) \quad (80)$$

The crystal capacitance and plasma inductance (both constants) are calculated from the cell geometry. For a 1-cm thick KDP crystal (with a dielectric constant of ~ 20), the capacitance C per segment is

$$C = \frac{1.3 \text{ nF}}{N_{\text{seg}}} \quad (81)$$

at the 26- × 26-cm crystal aperture, and an approximate value for the plasma inductance L per segment is

$$L = \frac{15 \text{ to } 30 \text{ nH}}{N_{\text{seg}}} \quad (82)$$

The plasma resistances $R_{1,i}$ and $R_{1,i}$ vary according to the net charge transported through the plasma. That is,

$$R_{1,i}^{-1} = R_{1,0}^{-1} + A_1 \int_0^t dt' |I_i(t')| \quad (83)$$

and

$$R_{1,i}^{-1} = R_{1,0}^{-1} + A_2 \int_0^t dt' |I_{i+1}(t') - I_i(t')| \quad (84)$$

where A_1 and A_2 are avalanche coefficients, and $R_{1,0}$ and $R_{1,0}$ are initial values determined by the preionization conditions.

To model $R_{1,0}$ at preionization currents and plasma electrode apertures other than those experimentally investigated, we assume that the resistance is inversely proportional to the peak preionization current:

$$R_{1,0} = \frac{(100 \pm 10)L_p}{I_{\text{pre}}N_{\text{seg}}} \Omega, \quad (85)$$

where L_p is the plasma electrode aperture in centimeters and I_{pre} is the peak preionization current per electrode. For a single segment, and with $L_p = 28$ cm, Eq. (85) is consistent with transverse resistances of 17.8 and 2.5 Ω at low and high preionization, respectively. The aperture scaling in Eq. (85) assumes the worst case, a linear increase in resistance with aperture (for a square aperture).

The transverse avalanche coefficient A_1 is determined from the slope of the conductivity vs transported charge data from Fig. 7-118. Including aperture scaling that is consistent with Eq. (85), we obtain

$$A_1 = \frac{(7.6 \pm 1.4)}{L_p} \times 10^4 N_{\text{seg}} \text{ mho/C.} \quad (86)$$

The values of $R_{1,0}$ and A_2 must be determined empirically by comparing Pockels cell performance predicted by the model with experimental data at low and high preionization.

We used a fourth-order Runge-Kutta routine to numerically integrate the coupled voltage and current equations. A time step of 0.1 ns was required for integration stability; a three- or five-segment model was sufficiently accurate, given the uncertainty in coefficients.

Figures 7-121(a) and 7-121(b) show measured current charging times at several switching voltages for low and high preionization, respectively. The current charging time is the width of the current pulse at the 10% points relative to the peak value. The charging time so defined is within 3% of the voltage rise time, defined as the time for the last crystal segment to reach 95% of the driver voltage. The optical rise time of the Pockels cell from 0 to 95% transmission is typically 75% of the voltage rise time. The dashed curves in Fig. 7-121 are model predictions with the transverse and longitudinal plasma resistances both set to zero. The charging time of 26.3 ns, which is independent of driver voltage and preionization current, indicates an ultimate optical switching time of less than 20 ns. At low preionization, the longitudinal resistance varies from 7.5 to 10 Ω , and at high preionization the best value is $\sim 12.5 \Omega$. The same value of A_2 was optimum for both low- and high-preionization data.

Fig. 7-121. Measured charging current pulse duration vs main switching pulser voltage at high and low preionization current, and corresponding model calculations.

Fig. 7-122. Model predictions of optical performance (0 to 95% transmission) for plasma-electrode Pockels cells with apertures from 25 cm to 3 m. Preionization current linearly scaled with aperture at 40 A/cm for each discharge chamber.

From our data we cannot determine the dependence of the longitudinal resistance on preionization current, and we therefore take

$$R_{l,0} = (10 \pm 2.5 \Omega) N_{\text{seg}}, \quad (87)$$

assuming that $R_{l,0}$ scales linearly with the number of crystal segments. The best value of A_2 for our conditions (helium at 10 Torr) was

$$A_2 = \frac{11.4 \pm 3.8}{N_{\text{seg}}} \times 10^4 \text{ mho/C.} \quad (88)$$

Depending on the plasma discharge, the particular inert gas, and the (unknown) concentration of trace impurities such as hydrocarbons, the transverse and longitudinal resistances and coefficients might vary significantly from the values given above, so those values should be used with care.

There is a minor discrepancy between the model and experimental observations. The model predicts that the optical transmission at point A will be ahead of point B by a few nanoseconds. The most likely explanation for the observation (Fig. 7-119) that the transmission of point B is *ahead* of point A is that the values $R_{l,0}$ for each segment are not the same. The glow from the plasma near the hollow cathodes is much brighter than in the rest of the

aperture, suggesting a higher degree of ionization in the vicinity of point B. If we assume a nonuniform distribution of $R_{l,0}$ values in the circuit, our computer model easily predicts that point B will rise before point A. None of this is of any consequence, however, since the overall optical performance of the Pockels cell is unchanged by this effect.

We can use our model to predict the performance of Pockels cells of various size operated with plasma under typical experimental conditions. Figure 7-122

shows results for Pockels cells of apertures up to 3 m. In all cases the KDP crystal is assumed to be 3 cm thick. To maintain a constant preionization current density, the preionization current per discharge chamber is scaled linearly with aperture at 40 A per centimeter of aperture. At the larger apertures, an array of crystals glued to an insulating frame might be fabricated from 25- x 25-cm segments. Assuming that the driving circuit has an impedance no greater than 5 Ω , we predict rise times of ~ 100 ns per meter of aperture. To obtain performance approaching that of a single 25-cm segment, we might also consider segmenting the main switching electrodes. However, the prospect of routing low-impedance, high-voltage cables to the central segments of an array is not very appealing. For most applications the turn-on (or turn-off) time of the optical switch must be less than the round-trip transit time of light through the laser system. For large, megajoule-class ICF lasers with regenerative multipass architectures, the transit times are about 1 μ s or greater; thus the predicted performance with unsegmented switching electrodes is probably adequate.

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Basic Materials Studies

Introduction

The optical materials research program, which is supported by the Office of Basic Energy Sciences of the Department of Energy, involves basic research on topics of importance to the Laser Program. Three research areas of intermediate- or long-range impact on laser development provide the framework for this research effort:

1. Spectroscopy and theory of optically active ions in solid state laser media.
2. Basic research on the mechanisms of optical surface damage in optical materials.

3. Measurement and modeling of the magnitude and dispersion of the nonlinear refractive indices of optical materials.

During 1985, our research activities emphasized the first two of these areas. In addition, we developed plans to facilitate rapid and flexible measurements of the nonlinear refractive index for various laser wavelengths and polarizations. Our major accomplishments are summarized below, and further details are given in the remaining articles in this subsection.

Divalent samarium was investigated as a possible alternative to Cr^{+3} as a sensitizer for neodymium and other rare earth ions in fluoride hosts. It was concluded that a somewhat restrictive set of conditions involving the crystal-field environment and valence stability of the samarium must be satisfied by a crystalline host material to make this sensitization scheme useful.

Ab initio calculations of the energy levels and wave functions of impurity ions in crystals have been made and tested against optical measurements. The results offer the possibility of predicting and accounting for a wide range of optical properties of ion-host combinations of possible utility for solid state laser systems.

Calculations of two-photon absorption (TPA) in ZnS using the pseudopotential method have been shown to account quite well for the observed spectra. This is the first application of pseudopotentials to calculations of the frequency dependence of TPA.

The emission of electrons and ions produced by laser excitation of the surfaces of optical materials is being employed as a probe of optical damage processes. Results obtained during 1985 show that this approach has good potential for investigating the basic causes of optical surface damage.

Author: L. L. Chase

Sensitization of Nd^{3+} -Doped Crystals by Sm^{2+}

The Nd^{3+} ion has long been recognized as having many desirable properties for laser applications. But its pump bands arise from the forbidden 4f-4f transitions, which are weak and sharp in crystalline

hosts, so that broadband flashlamp light is inefficiently used. A possible solution to this problem involves a suitable sensitizer ion co-doped into the laser host with the Nd^{3+} ions. Such an ion would absorb in regions of the spectrum where Nd^{3+} does not, and would then efficiently transfer its energy to the Nd^{3+} ions. The Cr^{3+} ion has been used effectively for this purpose in $\text{Gd}_2\text{Sc}_2\text{Ga}_3\text{O}_{12}$ (GSGG), but Cr^{3+} is not soluble in many crystals, and the energy transfer rate is too slow in many other hosts. As a result, Cr^{3+} is often unsuitable for Nd^{3+} sensitization. It is therefore of interest to identify other sensitizers.¹⁶¹ To this end we have examined the optical properties of a variety of crystals doped with Sm^{2+} and Nd^{3+} .

Spectroscopy of $\text{CaF}_2:\text{Sm}^{2+}, \text{Nd}^{3+}$.

Figure 7-123 shows the absorption and emission spectra of a CaF_2 crystal co-doped with Sm^{2+} and Nd^{3+} . The sharp absorption features are due to the Nd^{3+} ions; the broad Sm^{2+} absorption bands effectively fill the gaps between the Nd^{3+}

lines, so flashlamp light will be absorbed much better by the co-doped crystal than by $\text{CaF}_2:\text{Nd}^{3+}$. The Nd^{3+} lines at 740 and 790 nm overlap the Sm^{2+} emission band, as is required for efficient energy transfer from Sm^{2+} to Nd^{3+} .

The $\text{Sm}^{2+} \rightarrow \text{Nd}^{3+}$ transfer efficiency η is defined as the ratio of the energy transfer rate to the overall rate of Sm^{2+} emission decay. It can be described within the Forster-Dexter model of a dipole-dipole interaction^{162,163} with

$$\eta = \pi^{1/2} \alpha \exp(\alpha^2) [1 - \text{erf}(\alpha)], \quad (89)$$

where $\alpha = (\pi^{1/2}/2)(C_{\text{Nd}}/C_0)$, and C_{Nd} is the concentration of Nd^{3+} . The critical concentration C_0 is defined by

$$C_0^{-1} = \left(\frac{\hbar c}{n} \right)^2 \left[\frac{4\pi}{3} \int \frac{g_{\text{Sm}}(E) \sigma_{\text{Nd}}(E)}{E^4} dE \right]^{1/2} \quad (90)$$

where n is the refractive index, g_{Sm} is the area-normalized emission spectrum of the

Fig. 7-123. Absorption and emission spectra at 200 K for CaF_2 containing 0.39 mol % Nd^{3+} and 0.0087 mol % Sm^{2+} . The broad bands are due to Sm^{2+} and the sharp features to Nd^{3+} .

Sm^{2+} sensitizer, σ_{Nd} is the absorption line shape of the activator expressed in terms of the absolute value of the cross section, and E is the photon energy. The integral therefore represents the spectral overlap

Fig. 7-124. Absorption and (vertically offset) excitation spectra at 200 K of CaF_2 containing 0.39 mol % Nd^{3+} and 0.0087 mol % Sm^{2+} . The Nd^{3+} emission was monitored for the excitation spectrum. The calculated transfer efficiency is $\eta = 0.86$.

Fig. 7-125. Sm^{2+} emission transients at 100 K, detected at 720 nm, following an excitation pulse near 616 nm. ▼

defined qualitatively above. The transfer efficiency η increases monotonically with C_0^{-1} , so to achieve the highest possible energy transfer rate we must maximize the overlap integral. Using the $g(E)$ and $\sigma(E)$ data in Fig. 7-123, and using $\sigma_{\text{Nd}}(790 \text{ nm}) = 2.3 \times 10^{-20} \text{ cm}^2$ in Eq. (90), we find¹⁶⁴ that $C_0 = 0.43 \text{ mol \%}$. According to Eq. (89), a sample doped with $C_{\text{Nd}} = C_0 = 0.43 \text{ mol \%}$ of Nd^{3+} will have a transfer efficiency of 72%.

We can verify that energy transfer is occurring in $\text{CaF}_2:\text{Sm}^{2+}, \text{Nd}^{3+}$ by monitoring the Nd^{3+} luminescence at 1040 to 1055 nm and scanning the exciting wavelength. Figure 7-124 shows such a data set. This method can be understood by considering two limiting situations: for slow energy transfer, only the sharp Nd^{3+} line would be observed in the excitation spectrum, while for rapid energy transfer the excitation and absorption spectra would be identical. In between these limiting cases, the efficiency can be calculated with

$$\eta = \frac{A_{\text{Sm,exc}}/A_{\text{Sm,abs}}}{A_{\text{Nd,exc}}/A_{\text{Nd,abs}}} \quad (91)$$

where A is the area of a selected samarium band (at 625 nm) or a neodymium band (at 575 nm), derived from an absorption or an excitation spectrum, as indicated in the subscript. For the case shown in Fig. 7-124, the sample contained 0.39 mol % Nd^{3+} and the transfer efficiency was $\eta = 0.86 \pm 0.3$. Although this method is not very accurate, it provides important confirmation that the energy transferred from Sm^{2+} is converted into Nd^{3+} luminescence, and not nonradiatively depleted in some way. Foerster-Dexter analysis of the Sm^{2+} emission decays, described below, provides the needed accuracy.

Emission Lifetimes of $\text{CaF}_2:\text{Sm}^{2+}, \text{Nd}^{3+}$ and $\text{SrCl}_2:\text{Sm}^{2+}, \text{Nd}^{3+}$. The energy transfer rate can be deduced by time-resolving the emission decays of Sm^{2+} resulting from pulsed excitation. Such data are shown in Fig. 7-125 for five samples of CaF_2 with different Nd^{3+} concentrations.¹⁶⁴ Increasing the Nd^{3+} concentration results in more rapid energy transfer to Nd^{3+} , which competes directly with Sm^{2+} emission. Hence a shortening of the emission decay is observed. The intensity

$I_{\text{Sm}}(t)$ can be fitted to the Foerster-Dexter equation to yield C_0 :

$$I_{\text{Sm}} = \exp\left(\frac{-t}{\tau_{\text{Sm}}}\right) \exp\left[-\pi^{1/2} \left(\frac{C_{\text{Nd}}}{C_0}\right) \left(\frac{t}{\tau_{\text{Sm}}}\right)^{1/2}\right] \quad (92)$$

where τ_{Sm} is the Sm^{2+} emission lifetime. The average value of $C_0 = 0.57$ mol % determined from Fig. 7-125 is in good agreement with the value of 0.43 calculated from the overlap of the Nd^{3+} absorptions with the Sm^{2+} emission [see Fig. 7-123 and Eq. (90)].

The values of C_{r}/C_0 vs temperature can be used to calculate the transfer efficiency from Eq. (89). These results are shown in Fig. 7-126. The decrease above 200 K is due to nonradiative emission quenching of the Sm^{2+} ions. Thus Sm^{2+} is not a viable room-temperature sensitizer in CaF_2 . We consider the mechanism of Sm^{2+} nonradiative quenching below.

The $\text{SrCl}_2:\text{Sm}^{2+}$ crystal luminesces efficiently at room temperature. The Sm^{2+} emission band in SrCl_2 is 30 nm higher in energy than it is in the CaF_2 host. The resulting poor spectral overlap with the Nd^{3+} absorption lines makes the energy transfer rate too low to be useful in SrCl_2 . In fact, we find that $C_0 = 6$ mol % Nd^{3+} for SrCl_2 , compared to 0.6 in CaF_2 . The case of $\text{SrCl}_2:\text{Sm}^{2+}, \text{Nd}^{3+}$ illustrates the complexity of finding a host in which the Sm^{2+} ion luminesces efficiently and has the desired spectral overlap.

Emission Quenching of Sm^{2+} in the Alkaline Earth Halides.¹⁶⁵ The luminescent properties of Sm^{2+} -doped materials are sensitively affected by the relative position of the lowest of the $4f^{5d}$ and the $^5\text{D}_0$ ($4f^6$) excited states. For $\text{CaF}_2:\text{Sm}^{2+}$, the lowest excited state is derived from the $4f^{5d}$ manifold. For the other hosts we investigated, SrF_2 , BaF_2 , and SrCl_2 , the $^5\text{D}_0$ ($4f^6$) level is lower than the $4f^{5d}$ state; since the $^5\text{D}_0$ lifetime is ~ 10 ms (at least at low temperature), while the $4f^{5d}$ emission lifetime is ~ 1.4 μs , the radiative rate k_{rad} must include the rate from the thermally populated $4f^{5d}$ level. Experimentally one measures a lifetime on the order of 10 ms at low temperature, and the lifetime decreases as the temperature rises. Although the SrCl_2 data in Fig. 7-127 can

Fig. 7-126. Calculated energy transfer efficiency vs temperature for samples A-D of Fig. 7-125.

be fitted to this two-level emission model, the situation is quite different in the alkaline earth fluorides because of nonradiative emission quenching. At high temperatures the lifetime data for these materials diverge from the dashed curves in Fig. 7-127, which represent the calculated radiative lifetime. These nonradiative effects can be included with an equation of the form

$$\tau^{-1} = k_{\text{rad}} + A_{\text{nr}} \exp\left(\frac{-E_{\text{nr}}}{kT}\right), \quad (93)$$

where A_{nr} and E_{nr} are called the nonradiative pre-exponential factor and activation energy, respectively. The solid lines in Fig. 7-127 are fits to Eq. (93).

We considered the possibility that the nonradiative processes might be due to competition of Sm^{2+} ionization from the $4f^{5d}$ level with the emission process. To test this possibility, we calculated the "optical" energy separation E_{opt} of the $4f^{5d}$ level from the conduction band using the photoionization theory developed by Pedrini, McClure, and Anderson.¹⁶⁶ This "optical" energy involves a vertical transition that does not allow for relaxation of the surrounding lattice, and is therefore expected to be somewhat greater than the "thermal" energy separations that we obtain from our data. Table 7-15 gives the calculated values of E_{opt} . It also gives the thermal energy separations that we obtained experimentally using $E_{\text{therm}} = E_{\text{nr}} - E_{\text{ld}}$, where E_{ld} is the spectroscopically determined energy gap between the $4f^{5d}$ and $^5\text{D}_0$ ($4f^6$) states. Despite the different magnitudes of E_{opt} and E_{therm} , the relative values of these two quantities follow

Fig. 7-127. Emission lifetime vs temperature for Sm^{2+} in alkaline earth halides.

Table 7-15. Experimental thermal energies and calculated optical energies from the $4f^{5/5d}$ state of Sm^{2+} to the conduction band.

Crystal	E_{therm} (eV)	E_{opt} (eV)
CaF_2	0.28	1.9
SrF_2	0.15	1.2
BaF_2	0.03	0.1
SrCl_2	>0.8	~3.8

the same trend for the various alkaline earth halides. Thus the data in Table 7-15 provide evidence that ionization is the source of non-radiative quenching in the alkaline earth fluorides, and this explains why the quenching temperature increases in the BaF_2 - SrF_2 - CaF_2 - SrCl_2 series.

Properties of Sm^{2+} in the Mixed Fluoride Crystals.¹⁶⁵ The physical properties of the mixed fluorides of the general formula $\text{Me}_{1-x}\text{Ln}_x\text{F}_{2+x}$, where Me may be the metals Ca, Sr, or Ba, and Ln is Sc, Y, or a trivalent rare earth, have been of great interest for some time because of their potential application for high-energy lasers activated by rare earth ions.¹⁶⁷ We have found that the optical properties of Sm^{2+} in these materials differ considerably from those of Sm^{2+} in CaF_2 . The absorption and emission spectra of $\text{Ca}_{1-x}\text{Y}_x\text{F}_{2+x}$: Sm^{2+} and $\text{Ca}_{1-x}\text{La}_x\text{F}_{2+x}$: Sm^{2+} show a

crystal-field splitting of 9500 cm^{-1} , compared to 17000 cm^{-1} found in CaF_2 : Sm^{2+} . Since charge neutrality requires that each Ln^{3+} ion incorporate an extra interstitial fluoride ion into the CaF_2 lattice, it is reasonable to attribute the reduction in the crystal field to the presence of additional fluorines in the Sm^{2+} environment. The extra negative potential due to these interstitial fluorines is also expected to raise the energy levels of the Sm^{2+} ion enough that the $4f^{5/5d}$ excited state may become degenerate with the conduction band. This would result in nonradiative quenching as described by our ionization model.

Figure 7-128 shows the Sm^{2+} emission lifetime and intensity vs temperature for the mixed fluoride hosts $\text{Ca}_{1-x}\text{Y}_x\text{F}_{2+x}$ and $\text{Ca}_{1-x}\text{La}_x\text{F}_{2+x}$. As the temperature increases, the emission lifetime remains constant, but the emission intensity decreases. This shows that there is competition between relaxation into the emitting 5D_0 ($4f^6$) state and some other nonradiative process from the $4f^{5/5d}$ level. The two possibilities for this competing process are ionization of the $5d$ electron and a nonradiative multiphonon transition direct to

the ground state. The direct transition to the ground state from the $4f^{5d}$ level would have a negligible effect on the quantum efficiency, however, because nonradiative transition rates decrease very rapidly as the energy gap between the initial and final states increases. Therefore the nonradiative decay rate to the emitting 3D_0 excited state must be orders of magnitude greater than that to the ground state at all temperatures, and no decrease in emission intensity should be observed. This leaves ionization as the only alternative. Thus the data in Fig. 7-128 provide additional evidence that the quenching of Sm^{2+} is determined by the relative positions of the conduction band and the $4f^{5d}$ level.

Requirements for Sm^{2+} Sensitization of Nd^{3+} . The $\text{CaF}_2:\text{Sm}^{2+},\text{Nd}^{3+}$ system exhibits many of the properties that will ultimately be required to successfully utilize Sm^{2+} sensitization for a Nd^{3+} -doped laser host. The absorption bands are dipole-allowed and broad, and they fall approximately between the Nd^{3+} lines, assuring good absorption of flashlamp light. The good overlap of the Sm^{2+} emission band with the Nd^{3+} absorption bands provides for efficient energy transfer. Also, the short Sm^{2+} lifetime diminishes potential problems associated with excited-state absorption of Sm^{2+} at the Nd^{3+} laser wavelength. Since the Sm^{2+} luminescence quenches above 220 K, CaF_2 is not the optimal host. We must select a host in which the photoionization threshold of the Sm^{2+} ion is greater, and this must be done without altering the other desirable host-sensitive optical properties of $\text{CaF}_2:\text{Sm}^{2+}$.

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Theoretical Studies of Impurity Ions in Crystals

The characterization of the electronic structure of impurities in solids is particularly important for the systematic development of new solid state laser media. The large number of ion-host combinations available precludes a case-by-case experimental study of materials that may be important in the search for tunable solid

Fig. 7-128. Emission intensity and lifetime vs temperature for Sm^{2+} in the mixed fluorides, excited at 355 nm.

state lasers, so we are developing and testing methods with which to study the optical properties of such materials.

Unlike the Bloch states of a perfect crystal, the electronic states associated with impurities are localized on a very few atoms and are usually well separated in energy from those of the host crystal. An impurity ion and the neighboring host ions constitute a cluster that can be studied with *ab initio* molecular structure methods. We have modeled the Cu^+ ion impurity in NaF and NaCl hosts by carrying out Hartree-Fock (H-F) calculations for the ground $^1A_{1g}$ and excited $^{168,169}E_g$ and $^{1,3}T_{2g}$ states of such a cluster. The excited states correspond to the promotion of an electron from the $3d^{10}$ ground-state configuration to a $4s$ orbital. The resulting energy level structure was compared with the experimental spectra.^{168,169} Potential energy curves corresponding to symmetric displacement of the nearest-neighbor ions were calculated for each state. Spin-orbit coupling was included, and the coupling parameter was varied to reproduce the observed splitting of the T_{1g} and T_{2g}

Fig. 7-129. The 33-atom molecular cluster for NaF and NaCl. The Cu^+ impurity is at the center of the cluster.

Fig. 7-130. Potential energy curves for symmetric stretch motion of the nearest F^- ions of $\text{NaF}:\text{Cu}^+$.

fine-structure levels arising from the 3E_g state.

Hartree-Fock Calculations for $\text{NaF}:\text{Cu}^+$ and $\text{NaCl}:\text{Cu}^+$. The cluster was chosen so that the impurity ion and the nearest halide-ion neighbors were completely coordinated. The resulting 33-atom cluster is shown in Fig. 7-129. All electrons on Cu^+ were explicitly included in the calculation. The Gaussian basis set for Cu^+ was taken from the work of Hood et al.¹⁷⁰ The effective core potential (ECP) and basis set of Kahn¹⁷¹ were used on the nearest-neighbor F^- ions. The $\text{NaCl}:\text{Cu}^+$ calculations were carried out using the ECP and basis set of Christiansen¹⁷² for Cl^- . The halide basis sets were augmented with a diffuse p-function, as suggested by Dunning and Hay.¹⁷³ The Na^+ ions were represented by total ion potentials.^{174,175}

We fixed all the ions except the nearest neighbors at the experimental lattice positions, and determined the potential energy curves for the symmetric stretch motion of the six ions nearest the Cu^+ in NaF . The

results are shown in Fig. 7-130. The calculated equilibrium distances for the ground and excited states for NaF:Cu⁺ and NaCl:Cu⁺, and the values for the undoped crystals, are given in Table 7-16. The lattice constants for NaF and NaCl were calculated using the same method, except that the Cu⁺ impurity was replaced by a Na⁺ ion. The agreement in row 1 of the table between nearest-neighbor distances in crystals with and without the Cu⁺ impurity is consistent with the similar ionic radii of Cu⁺ (1.814a₀) and Na⁺ (1.795a₀). The increase in the lattice constant from the ground state to the excited states is very nearly the same for both crystals. The experimental lattice constants for the undoped crystals are 4.38a₀ for NaF and 5.44a₀ for NaCl.¹⁷⁶ The calculated lattice constants are slightly higher in both cases, because of our neglect of the crystal field of the rest of the lattice. Our calculated equilibrium distances for the doped and undoped crystals indicate that the experimental nearest-neighbor distance for the undoped crystal should be a very good approximation to the value in the doped crystal.

Discussion of Results for NaF:Cu⁺. In Table 7-17 the vertical excitation energies for NaF:Cu⁺ at the calculated CuF and experimental NaF separations are compared with the positions of the absorption bands measured at pumped-helium temperature by Payne et al.¹⁶⁸ Both theory and experiment predict that the excitation energies are greater for NaF:Cu⁺ than for the free ion. The differences between the calculated and measured excitation energies most probably arise from our neglect of the contribution from electron correlation and the much smaller contributions from relativistic effects and spin-orbit coupling. For the free ion the correlation energy difference (the difference between the calculated and measured H-F excitation energies) is ~1.5 eV. For NaF:Cu⁺, the correlation energy difference is less: for example, at the experimental nearest-neighbor distance R_{CuF} = 4.38a₀, the difference is ~0.7 eV, while at our calculated value R_{CuF} = 4.61a₀, the difference is ~1.2 eV. Changing R_{CuF} has a significant effect on the calculated excitation energies, but (as Fig. 7-129 and Table 7-17 show) the positions of the excited levels with

State	NaF:Cu ⁺	NaCl:Cu ⁺	NaF ^b	NaCl ^c
3d ¹⁰ 1A _{1g}	4.61	5.45	4.57	5.46
3d ⁹ 4s 3E _g	4.71	5.57	—	—
3d ⁹ 4s 3T _{2g}	4.72	5.59	—	—
3d ⁹ 4s 1E _g	4.71	5.57	—	—
3d ⁹ 4s 1T _{2g}	4.72	5.79	—	—

^a All distances are in atomic units.

^b The experimental lattice constant for NaF is

4.38a₀.

^c The experimental lattice constant for NaCl is

5.33a₀.

State	ΔE (eV)	ΔE (eV)	ΔE (eV)
	R _{CuF} = 4.61a ₀	R _{CuF} = 4.38a ₀	Experiment ^a
3d ¹⁰ 1A _{1g}	0.0	0.0	0.0
3d ⁹ 4s 3E _g	2.61	3.14	3.87
3d ⁹ 4s 3T _{2g}	2.77	3.37	4.04
3d ⁹ 4s 1E _g	3.06	3.58	3.99, 4.29 ^b
3d ⁹ 4s 1T _{2g}	3.26	3.86	4.51

^a For NaF:Cu⁺ the experimental energies are from Ref. 168.

^b Jahn-Teller components from Ref. 168.

respect to one another are less affected by a change in R_{CuF}.

Interpretation of the 1E_g excitation energy is complicated by the Jahn-Teller distortion of the upper state. For NaF:Cu⁺, Payne et al.¹⁶⁸ observed two E_g peaks split by 2400 cm⁻¹, corresponding to transitions to the two branches of the upper surface. Our calculations did not include this effect and should correspond to the energy difference between the ground state and the conical intersection of the two excited Jahn-Teller surfaces.

We included the spin-orbit interaction among the 3d⁹4s states of NaF:Cu⁺ using perturbation theory.¹⁷⁷ Table 7-18 gives the spin-orbit excitation energies and wave functions for each of the states in terms of the LS states for the experimental lattice constant, as determined using the spin-coupling constant of the free ion (−828 cm⁻¹). By increasing the calculated excitation energies by 0.63 eV, we can bring the A_{1g} — 3T_{2g} transition into agreement with experiment. This corrects the remaining calculated levels (observed and unobserved) for electron correlation errors and allows a quantitative comparison with experiment. The shifted excitation energies are listed in the last column of Table 7-18 and are shown alongside the low-temperature absorption spectrum in Fig. 7-131. The energy levels are shown for several values of the spin-orbit coupling constant ζ, and are shifted so that the highest T_{2g} level agrees with experiment in

Table 7-16. Comparison of the calculated nearest-neighbor distance for the ground and excited states of Cu⁺-doped NaF and NaCl to the calculated lattice constants for the undoped crystals.^a

Table 7-17. Comparison of the theoretical excitation energies of NaF:Cu⁺ with experiment, for the calculated equilibrium CuF distance and for a value equal to the experimental lattice constant of NaF.

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Table 7-18. Spin-orbit energies and wave functions for NaF:Cu^{2+} using the experimental NaF lattice constant and the Cu^{2+} spin-coupling constant.

State	Calculated ΔE (eV) $R_{\text{CuF}} = 4.38a_0$	Wave function	Shifted ΔE (eV)
T_{1g}	3.09	$0.935 ^3E_g\rangle + 0.356 ^3T_{2g}\rangle$	3.72
T_{2g}	3.11	$0.952 ^3E_g\rangle + 0.290 ^3T_{2g}\rangle + 0.120 ^1T_{2g}\rangle$	3.74
A_{2g}	3.27	$1.0 ^3T_{2g}\rangle$	3.90
T_{2g}	3.32	$0.268 ^3E_g\rangle - 0.946 ^3T_{2g}\rangle + 0.182 ^1T_{2g}\rangle$	3.95
E_g	3.35	$0.877 ^3T_{2g}\rangle - 0.480 ^1E_g\rangle$	3.98
T_{1g}	3.47	$-0.356 ^3E_g\rangle + 0.935 ^3T_{2g}\rangle$	4.10
E_g	3.65	$0.480 ^3T_{2g}\rangle + 0.877 ^1E_g\rangle$	4.28
T_{2g}	3.88	$-0.149 ^3E_g\rangle + 0.145 ^3T_{2g}\rangle + 0.978 ^1T_{2g}\rangle$	4.51

Fig. 7-131. Shifted Hartree-Fock energy levels for three values of the spin-orbit interaction parameter ζ , and corresponding experimental observation.

each case. The spectrum clearly shows the two spin-allowed but parity-forbidden transitions, including the two bands corresponding to the Jahn-Teller components of the E_g state. The relative position of the Jahn-Teller peaks depends on the Franck-Condon distribution and on the splitting of the two branches of the distorted state. On the basis of qualitative arguments regarding the Jahn-Teller splitting of the 1E_g state, we estimate that the high-energy E_g band lies slightly above the energy of the conical intersection of the two surfaces. This puts our calculated position for the 1E_g state in good agreement with experiment. Payne et al.¹⁶⁸ eliminated the E_g contribution from the two-photon room-temperature spectra using polarized light and found a weak T_{2g} transition at 4.04 eV; this transition could correspond to the

3.95-eV $T_{2g}(^3T_{2g})$ level in Table 7-18. The calculations predict that the lower-energy Jahn-Teller band is also overlapped by the lowest E_g state. Table 7-18 shows that the lowest E_g and the second T_{2g} state have 23 and 3.3% singlet character, respectively, and should contribute to the absorption in the 4-eV region of the spectrum. However, the shift from room temperature to pumped-helium temperature may change the apparent agreement between theory and experiment in this case.

Further information on the spin-forbidden levels can be obtained from the emission spectrum. Goldberg et al.¹⁷⁶ measured emission bands at 25 600 and 26 600 cm^{-1} for NaF:Cu^{2+} at 300 K. At 77 K, they saw only one band, with a peak at 27 000 cm^{-1} . They suggested that the two peaks might be the Jahn-Teller levels of the 3E_g

state. We have used the calculated potential energy curves and the results of Table 7-18 to estimate the emission wavelengths. The two lowest T_{2g} states have singlet character, as shown in Table 7-18. The calculated (and unshifted) emission energies to the A_{1g} state are $\Delta E[T_{2g}({}^3E_g)] = 2.55$ eV ($20\,575\text{ cm}^{-1}$) and $\Delta E[T_{2g}({}^3T_{2g})] = 2.69$ eV ($21\,680\text{ cm}^{-1}$). The difference is 1105 cm^{-1} , in excellent agreement with the measured 1000 cm^{-1} . In addition, the energy shift of the T_{2g} levels required to bring the emission wavelengths into agreement with experiment is $+0.62$ eV, which agrees with the $+0.63$ -eV shift we applied to the vertical energy levels determined from the absorption spectrum. Sturge¹⁷⁸ has pointed out that the emission band due to a transition from a distorted E state to an A state will have only one peak. This gives further support to the argument that the two T_{2g} states, and not the Jahn-Teller components of the 3E_g state, are responsible for the observations of Goldberg et al.¹⁷⁶

In their study of the effect of a magnetic field on the luminescent lifetime of Cu^{+} in alkali-halide hosts, Payne et al.¹⁶⁸ found the splitting of the lowest T_{1g} and T_{2g} states derived from the 3E_g state to be 26 cm^{-1} . Using the free-ion spin-orbit coupling constant, we obtain a calculated splitting of $\sim 160\text{ cm}^{-1}$. If we reduce the spin-orbit coupling constant to $\sim 250\text{ cm}^{-1}$, we obtain a calculated splitting of 30 cm^{-1} , roughly the value observed. The energy level patterns determined with the two values for the coupling constant are compared with the energies without spin-orbit coupling Fig. 7-131. The pattern is qualitatively the same for each value of the coupling constant, but the excitation energies with reduced spin-orbit interaction should be more accurate. The reduced spin-orbit splitting relative to that of the free ion may be a consequence of the quenching of the orbital angular momentum by the Jahn-Teller distortion.¹⁷⁸

Summary. Although the theoretical results presented here do not include electron correlation or relativistic effects, the relative positions of the excited levels are in good agreement with experiment. The predicted locations of the spin-forbidden states and their splittings are also in

agreement with the measured spectra. On the basis of our calculations, we assign the emission bands observed for $\text{NaF}:\text{Cu}^{+}$ at $25\,600$ and $26\,600\text{ cm}^{-1}$ to the lowest two T_{2g} states.

Our study of the $\text{NaCl}:\text{Cu}^{+}$ system gave results very similar to those obtained for $\text{NaF}:\text{Cu}^{+}$. After we corrected the H-F levels for electron correlation and relativistic effects, the electronic transitions from the ground state overlapped the observed absorption spectra, although the width of the bands makes quantitative comparison difficult. In both hosts, observations indicate a significant reduction in the spin-orbit splitting, which we attribute to quenching of the angular momentum by the Jahn-Teller distortion. This effect requires further theoretical study. For $\text{NaCl}:\text{Cu}^{+}$ the shifted emission energy from the $T_{1g}({}^3E_g)$ state was 3.69 eV, in excellent agreement with the observed fluorescence at 3.6 eV.¹⁷⁹

Chermette and Pedrini¹⁸⁰ have carried out muffin-tin X_α calculations for $\text{NaCl}:\text{Cu}^{+}$. A comparison of our theoretical results with theirs shows considerable differences in the energy level patterns. This can be traced to the inability of the X_α method to predict singlet-triplet splittings correctly. Consequently when spin-orbit coupling is included, the energy level pattern is essentially inverted with respect to our results. Since the H-F method is relatively accurate in calculating singlet-triplet ordering and separation, a great deal of caution should be exercised in using X_α results to interpret or predict electronic spectra of impurities in crystals.

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Mechanisms of Optical Surface Damage

Although optical surface damage by high-energy laser pulses has been studied experimentally and theoretically for over two decades, the fundamental damage mechanisms are not yet well understood. It is important to achieve a basic understanding of the surface properties and interactions that lead to optical damage so that surface preparation procedures can be

devised that will lead to higher damage thresholds. The fact that optical damage on surfaces is usually initiated in small, isolated regions suggests that this goal can be achieved by understanding and controlling the conditions that lead to damage in those regions at fluences well below the damage threshold of the rest of the surface.

Our approach to investigating surface-damage mechanisms involves the use of surface-analytical techniques that have not heretofore been exploited systematically in such investigations. We have measured electron and ion emission resulting from laser excitation of polished surfaces of several amorphous and crystalline optical materials. The objectives of this research are to identify the basic mechanisms of optical damage at photon energies much less than the optical band gap and to identify the surface regions that are most susceptible to damage so that they can be characterized by surface-analytical techniques. Although the emission of electrons, ions, and neutrals resulting from surface bombardment by ions and UV photons has been an active field of research for over two decades,¹⁸¹ little work has been done involving surface excitation at photon energies in the transparency region, where the mechanisms for surface emission are not generally known.

We used two quite different approaches to characterize laser-induced particle emission from optical materials. In one approach, measurements are made in ultrahigh vacuum utilizing optical excitation of large surface areas at calibrated fluences.¹⁸² These measurements provide information about surface processes that occur at fluences below the threshold for observable damage. The second method involves laser-ionization mass spectroscopy (LIMS), a relatively new technique of chemical analysis in which excitation of a sample surface with a well-focused laser produces ions that are analyzed with a time-of-flight mass spectrometer, providing high sensitivity and mass resolution. This approach is promising as a probe of surface damage mechanisms initiated by impurities or microscopic inclusions of foreign matter near the surface.

Electron and Ion Emission. We performed our measurements at a residual pressure of 10^{-9} to 10^{-10} Torr. The peak fluence in the 1-mm-diam irradiated area was measured with a Vidicon and calorimeter.¹⁸³ Neodymium-laser pulses with nearly rectangular temporal profiles and variable widths from 1 to 50 ns were employed. Data were obtained for several materials at the laser's first-harmonic ($1.06\ \mu\text{m}$) and third-harmonic ($0.355\ \mu\text{m}$) wavelengths.

Figure 7-132 shows the total emitted negative charge Q vs peak laser fluence F for several materials. These data fit the power law

$$Q = kF^N, \quad (94)$$

where N is approximately equal to the order of the multiphoton absorption required to excite electrons across the band gap of the material.¹⁸² Table 7-19 gives values of N_{calc} , obtained from the ratio of the optical band gap energy to the photon energy, and N_{obs} , obtained by fits to Eq. (94). Although it might be possible to fit the data with a function representing an exponentially activated thermal process with a temperature proportional to the fluence,¹⁸⁴ it would be difficult to account for the change in slope of the data with laser wavelength with a temperature model alone.

One mechanism that could account for these results is ordinary photoelectric emission resulting from the multiphoton excitation of the photoexcited electrons up to the vacuum level, after which they may leave the sample surface. The total charge emission would be expected to depend on F and the temporal pulse width t_p as

$$Q = \frac{KF^N}{t_p^{N-1}} \quad (95)$$

The ratio of the fluences required to produce a given charge emission Q for pulse durations t_{p1} and t_{p2} is then given by

$$\frac{F(t_{p2})}{F(t_{p1})} = \left(\frac{t_{p2}}{t_{p1}} \right)^{1-1/N} \quad (96)$$

Fig. 7-132. Total emitted negative charge Q vs laser fluence for two photon energies. Values of N are fitted slopes from Eq. (94).

Table 7-19. Calculated and observed power-law exponents N_{calc} and N_{obs} for electron emission from several materials with a range of optical band gaps E_g and work functions E_w . $h\nu$ is the laser photon energy.

Material	E_g (eV)	E_w (eV)	$h\nu$ (eV)	N_{calc}	N_{obs}
CdTe	1.5	6.0	1.16	2	2.1 ± 0.3
NaCl	8.75	4.2	1.16	8	7.2 ± 1.0
SiO ₂	8.0-8.4	5.0	1.16	7-8	11.0 ± 1.3
ZnS	3.88	5.5	3.5	3	5.0 ± 0.5
			1.16	4	4.4 ± 0.4
			3.5	2	2.1 ± 0.3

Charge emission data for CdTe, ZnS, and NaCl at several pulse durations are shown in Fig. 7-133. The logarithm of the ratio in Eq. (96) is given by the horizontal displacement of the lines for the different pulse lengths. This ratio scales approximately as $t_p^{1/2}$ for all three materials. This agrees with Eq. (96) only for CdS, for which $N_{\text{calc}} = 2$. For ZnS ($N_{\text{calc}} = 4$ to 5) and NaCl ($N_{\text{calc}} = 7$) the behavior expected for the simultaneous multiphoton photoemission model is clearly not observed.¹⁸² Since optical-damage thresholds for bare polished surfaces often scale as $t_p^{1/2}$, however, the data of Fig. 7-133 suggest a connection between the charge emission and the damage mechanism.

A second indication that something more complex than photoemission is occurring is that, at least for SiO₂ and ZnS, roughly equal amounts of positive and negative charge were emitted. We investigated the nature of these two emission

components for ZnS by measuring the time of flight of the emitted charge, which yields a mass distribution. These measurements employed a microchannel intensifier detector positioned at right angles to the incident laser beam. The negative emission from ZnS is shown in the upper trace in Fig. 7-134 for a 1-ns laser pulse. The time delay and duration of this emission are too short for anything but electrons to reach the detector, and there is clearly no detectable negative-ion component.

Time-of-flight data for positive ions emitted from ZnS are shown in the lower trace

Fig. 7-133. Laser-pulse-width dependence of negative charge emission for CdTe, ZnS, and NaCl (irradiation wavelength $1.06 \mu\text{m}$).

Fig. 7-134. Time dependence of negative and positive charge emission from ZnS following a 1-ns, $1.06 \mu\text{m}$ laser pulse.

of Fig. 7-134. The ion mass associated with a particular sample-to-detector transit time can be obtained from the calibration based on the time delay of the sharp leading edge of the electron emission. The two strongest peaks in Fig. 7-134 result from emission of a single mass species from the front and back surfaces of the sample, which was rotated so that the back surface emission had a shorter path to the detector. If the species is singly charged, its molecular or atomic weight is about 100. This is approximately the mass of both $(\text{ZnS})^+$ and $(\text{S}_3)^+$, the only products with this mass that might reasonably be expected from a ZnS surface. The only other significant peaks in Fig. 7-134 are two weak signals from the front and back surfaces at a mass probably corresponding to desorption of some hydrocarbon from the surface.

The absence of any prominent peaks in the emission spectra corresponding to ionized zinc or sulfur indicates a notable selectivity in the ion-emission mechanism. A plausible explanation for these results is that neutral atoms and molecules are emitted from the surface, possibly because of heating by the laser beam, and that ZnS or S_n is then preferentially photoionized by the laser light. It is, in fact, well established for a number of insulating materials that neutral-atom and molecule emission is far more probable than ion emission for surface bombardment by ions or UV photons.¹⁸¹ The proposed ejection-photoionization mechanism would also account for the fact that only electrons and positive ions are observed. Measurements are planned incorporating an ionizing mass spectrometer in place of the microchannel detector to measure the relative abundance of neutral zinc and sulfur atoms and the molecular products. The dependence of this neutral emission on excitation fluence and on the mass distribution of the emission may be of great utility in assessing the roles of thermal evaporation and multiphoton absorption in the solid and in the emitted vapor.

An important conclusion that can be drawn from our observations on ZnS is that microscopic damage in the form of surface ablation occurs at fluences well below the damage threshold. This may also be true of the other materials for which we have measured electron and ion emission.¹⁸² The charged- and neutral-particle emission appears to be a sensitive way to locate damage-susceptible surface regions and to study the mechanisms causing the emission, which may be the initial stage of the processes leading to observable damage.

LIMS Measurements. In a LIMS, the fundamental and harmonics of a Q-switched YAG:Nd laser are focused onto a micron-size region of the sample surface. Ions (and electrons, if desired) are accelerated in a time-of-flight drift tube with a velocity-reversing region that compensates for initial velocity at the sample surface. The detector signal goes to a transient recorder that records a complete mass spectrum on a single laser shot. Although the absolute fluence at the center of the focal spot on the sample surface is not known, relative fluences can be varied by inserting filters in the laser beam. As the surface is moved laterally by an x-y translator, one can find spots at which damage, identified by an abrupt increase in the ion emission, occurs at relatively low fluences. These spots can be chemically characterized by the mass spectrum of the ion emission. Spots with a higher damage threshold can be characterized by increasing the laser fluence until damage occurs. For research into optical surface damage, LIMS has the obvious advantage over other methods of surface analysis that only those constituents at the damage site are observed in the mass spectrum.

We have made LIMS measurements on glass samples polished with iron oxide, silicon carbide, or cerium oxide. Optical damage thresholds are substantially higher for surfaces polished with cerium oxide than for those polished with the other two materials. Typical results for the cerium-polished sample are shown in Fig. 7-135. The upper trace shows the mass spectrum obtained at a low-threshold spot on the sample surface. Only Ce and CeO are observed in the mass spectrum at masses 140 and 156, respectively. For many regions of the sample no ion emission is observed at this fluence. As the fluence is increased, ion emission is more probable everywhere, and mass spectra typical of the constituents of the sample surface are observed in addition to the Ce and CeO. The lower trace in Fig. 7-135 shows, for example, strong peaks at 28 to 30 from silicon isotopes and at 44 from SiO₂, in addition to peaks from cerium and other impurities and surface contaminants. Similar results were obtained for the surfaces polished with silicon carbide and iron oxide. The fluence required to produce

detectable ion emission is lowest for the iron oxide, highest for the cerium oxide, and intermediate for the silicon carbide polished surfaces. These results correlate with the optical damage thresholds.

Those experiments suggest that incorporation of polishing compound into the surface has a measurable effect on the optical damage threshold. The fact that the fraction of ions from the polishing material relative to those of the intrinsic surface material increases as the threshold for damage decreases implies that the polishing material does not act simply as an inert impurity distributed more or less uniformly over the surface.

Although the LIMS experiments elucidate the chemical composition of damage-sensitive surface regions, they yield no information on the possible chemical incorporation or depth of incorporation of polishing material and other foreign

Fig. 7-135. LIMS spectra for two laser intensities on different regions of a polished fused silica surface.

matter. We have investigated these questions by utilizing SIMS (secondary-ion mass spectroscopy) and ESCA (electron spectroscopy for chemical analysis) measurements on glass surfaces.

A SIMS survey was made of a cerium-oxide-polished glass surface. The depth profile is obtained by sputtering away the surface while monitoring the emitted ions. In addition to the chemical constituents of the glass and bulk impurities, such as iron, one observes cerium at decreasing relative concentration as the surface layer is sputtered away. The depth of incorporation of the cerium estimated from these results is roughly the diameter of the polishing grains. Lanthanum and neodymium were also observed, presumably because they are present as impurities in the polishing compound. The ESCA measurements, which give the relative concentrations of various elements, show that the lanthanum is incorporated into the surface layer at a concentration relative to cerium that is much greater than in the original polishing materials. This suggests a selective chemical mechanism for the incorporation of polishing materials into the surface layer of the glass.

The results we obtained in 1985 indicate that laser-induced charged- and neutral-particle emission and some of the common surface characterization techniques are promising methods for investigating the basic mechanisms of optical surface damage.

Authors: L. L. Chase and W. J. Siekhaus

Gas Laser Studies

Introduction

The Advanced Laser Program has continued its interest in the development of short-wavelength gas laser media. A small-scale experimental effort has been funded to investigate electronic-excitation transfer from metastable nitrogen in two possible high-power storage media, SO and PF. The weakly allowed A-X transitions in these radicals may provide attractive fusion laser candidates if efficient means to pump them can be found.

The xenon fluoride excimer can be pumped by means of photolysis of XeF₂.

This is a classic rare-gas halide laser system. An extensive configuration-interaction calculation of the ground and low-lying excited-state potential energy surfaces and dipole transition moments for XeF₂ provides a detailed description of the photodissociation process. These projects maintain our interest in current gas laser research.

Author: H. T. Powell

Energy-Transfer Laser Possibilities Based on N₂ A-State Molecules

Our investigations of possible gaseous fusion-laser media have focused on utilizing the excitation energy stored in N₂ ($A^3\Sigma_u^+$) molecules. By finding collision partners with appropriate fluorescence lifetimes to which the excitation energy can be rapidly transferred, we expect to be able to develop an efficient, high-power, short-wavelength laser.

It has been found that the SO molecule quenches the N₂ A-state, apparently producing a high fraction in excited electronic states. This molecule emits photons with wavelengths around 260 nm and has a radiative lifetime of about 10 μ s, so it is an excellent candidate for an ICF laser. In 1985, the fluorescence spectrum of this molecule was measured and assigned. The PF molecule was also found to produce emission due to excitation transfer from the N₂ A-state. This work was carried out by D. Setser and his group at Kansas State University, under contract to the LLNL Laser Research and Development Program.

The SO system was studied in a flowing-afterglow reactor. In assigning the emission spectrum, the distribution of population among the vibrational levels of the upper $A^3\Pi_0$ state was determined. Vibrational relaxation rates were evaluated on the basis of line strength variation with pressure. Cao and Setser¹⁸⁵ made a careful study of these line strengths and found discrepancies between their results and earlier work of Clyne and coworkers.^{186,187}

Because of the uncertainty in the emission spectrum and the importance of vibrational relaxation rates in determining the utility of this system for a fusion laser,

Cao and Setser began a new experiment to measure these quantities directly using time-resolved, laser-induced fluorescence. Exploratory and calibration experiments were carried out in 1985 on the known $A^2\Pi$ state of NO, using one- and two-photon excitation to specific vibrational levels of the upper state. Measured lifetimes and self-quenching rates for this system agreed well with published values. Application of these methods to SO will provide needed information about the line strengths.

A study of PF in the flow reactor also produced emission due to excitation transfer from the N_2 A-state. However, the observed fluorescence was not that of the expected $A^3\Pi \rightarrow X^3\Sigma^+$ transition and is not yet understood. The A-X emission can be produced by excitation transfer from helium metastables. A spectrum obtained in this way is shown in Fig. 7-136. Assignment of the peaks indicates that the ground state equilibrium bond distance differs significantly from that in the excited state, so that a population inversion can easily be achieved. The lifetime of the excited state is expected to be greater than $1 \mu s$, so this state should also be an excellent storage medium.

Further work on characterizing both these potential laser systems is in progress.

Author: K. C. Kulander

Major Contributor: D. Setser
(Kansas State University)

Theoretical Studies of Xenon Difluoride

The spectroscopic and kinetic properties of XeF_2 are of considerable interest to the laser-development community because XeF_2 is an efficient source of XeF^* , through collision with electrons or metastable rare-gas atoms or through vacuum-UV photolysis. XeF_2 is the simplest known stable molecule containing xenon. Spectroscopic analysis of gas-phase XeF_2 has established that its structure is linear and symmetric, with an Xe-F bond distance of $1.977 \text{ \AA}$. The dissociation energy of XeF_2 into free atoms has been measured at 2.78 eV . If we assume that the bond energy of F_2 is

1.69 eV , we obtain a value of 1.09 eV for the molecular dissociation energy for the reaction $XeF_2 \rightarrow Xe + F_2$.

The ground state of XeF_2 has been well characterized, experimentally and theoretically. Large-scale theoretical calculations by Bagus et al.¹⁸⁸ on the $^1\Sigma_g$ ground state give a dissociation energy of 2.14 eV and show the bonding to be strongly ionic ($F^{-1/2}Xe^+F^{-1/2}$) near the equilibrium geometry; there is a sharp transition to covalent character at longer Xe-F bond distances.

Vacuum-UV photolysis of XeF_2 at 7 to 9 eV produces XeF^* with a high quantum yield.¹⁸⁹ Simple arguments based on molecular orbital theory¹⁹⁰ identify this absorption band as an intramolecular charge transfer that decreases the ionic contribution to the bonding and leads to dissociation.

We have carried out theoretical investigations to extend the work of Bagus et al., to characterize the low-lying excited states of XeF_2 , to assign the observed electronic transitions, and to calculate dipole transition moments between the various states. In addition to the symmetric dissociation

Fig. 7-136. Fluorescence spectrum of PF, excited by the $He^+ + PF$ reaction. Transition assignments are indicated.

of XeF_2 studied by Bagus et al., we have also obtained slices through the potential surface corresponding to bending and to the removal of a single fluorine atom.

In our calculations, we treated XeF_2 as an effective 22-electron system: we used effective core potentials (ECPs) to replace the 46 inner-shell electrons of xenon and the 2 inner-shell electrons of fluorine. The reliability of properly constructed ECPs in calculations of molecular potential energy curves is well documented. However, the ability of an ECP to accurately predict molecular properties such as dipole moments and transition moments is not as well established. We therefore carried out preliminary calculations on Xe_2^+ and XeF , for which comparison with previous all-electron calculations^{191,192} could be made.

Fig. 7-137. Asymmetric dissociation of XeF_2 : potential energy curves for the two lowest $^1\Sigma$ and $^1\Pi$ states.

Table 7-20. XeF_2 electronic transition moments (atomic units). One Xe-F bond distance is frozen at 1.98 Å; the other, R , varies as shown.

R (a.u.)	$1^1\Sigma-2^1\Sigma$	$1^1\Sigma-1^1\Pi$	$1^1\Sigma-2^1\Pi$
3.0	1.68	0.21	0.11
3.25	1.85	0.14	0.16
3.5	2.05	0.08	0.20
3.74	2.23	0.0	0.21
4.0	2.31	0.08	0.17
4.5	2.33	0.11	0.09
5.0	2.47	0.10	0.06
5.5	2.59	0.07	0.05
6.0	2.44	0.05	0.05
6.5	1.99	0.03	0.05
7.0	1.46	0.02	0.05
8.0	0.73	0.008	0.05
9.0	0.36	0.005	0.05
100.0	0.27	0.0	0.05

The ECPs gave transition moments that agreed to within $\sim 5\%$ with the values obtained from the all-electron calculations.

A molecular-orbital basis for each geometry considered was obtained by first carrying out a self-consistent-field (SCF) calculation on the $^1\Sigma$ ground state of XeF_2 . The configuration of this ground state at the SCF level (treating only the fluorine 2p and xenon 5s and 5p electrons explicitly) is

$$1\sigma^2 2\sigma^2 3\sigma^2 1\pi^4 2\pi^4 3\pi^4. \quad (97)$$

Although this single configuration is dominant for Xe-F distances near the equilibrium value, it fails to describe the molecule at large separations and does not dissociate properly to Hartree-Fock atoms. For proper dissociation, we must add the configuration

$$1\sigma^2 2\sigma^2 4\sigma^2 1\pi^4 2\pi^4 3\pi^4. \quad (98)$$

We obtained a two-configuration multicomponent self-consistent-field (MCSCF) wave function and a new set of orthonormal molecular orbitals by an iterative natural-orbital procedure. This wave function, while dissociating properly to Hartree-Fock atoms and expressing the transition from ionic to covalent bonding in a qualitatively correct manner, does not give a quantitatively accurate description of the bonding in XeF_2 . This fact was noted by Bagus et al. The binding energy in XeF_2 is a correlation effect. In our most extensive calculations, we employed a first-order wave function that contains a full set of valence configurations and all single and double excitations in which no more than one electron occupies an orbital beyond the valence shell. "Valence" here refers to those orbitals occupied in the two-configuration MCSCF wave function. This prescription gave ~ 1600 to 1800 configurations, depending on the symmetry considered.

We consider first the asymmetric stretch ($C_{\infty v}$) case. Keeping one fluorine atom fixed at an Xe-F distance of 1.977 Å, we calculated the potential energies of the two lowest Σ and Π states as a function of the other Xe-F bond separation. The results are shown in Fig. 7-137. A broad continuous absorption, centered at ~ 7.9 eV, has been observed in XeF_2 . Our calculations

identify this absorption with transitions from the ground state to the first excited $^1\Sigma[3\sigma_{(g)} - 4\sigma_{(u)}]$ state at ~ 8.5 eV and to the second $^1\Pi[2\pi_{(g)} - 4\sigma_{(u)}]$ state at ~ 7.5 eV, correlating with the Σ and Π ionic states of XeF^+ , respectively. Although both transitions are symmetry-allowed from the ground state, the transition moment for the Σ - Σ transition is an order of magnitude greater than the Σ - Π transition moment. This explains why UV photolysis of XeF_2 produces only XeF (B-X) emission; $\text{D}(^2\Pi_{1/2})\text{-X}$ and $\text{C}(^2\Pi_{3/2})\text{-A}(^2\Pi_{1,2,3/2})$ emissions are not seen. The calculated transition moments are given in Table 7-20.

For the case of symmetric dissociation, we calculated the two lowest Σ potential curves ($^1\Sigma_g^+$ and $^1\Sigma_u^+$) and the associated dipole transition moment. The results are shown in Fig. 7-138.

The calculations give a dissociation energy of 1.95 eV for the symmetric process $\text{XeF}_2 \rightarrow \text{Xe} + 2\text{F}$, in agreement with the results of Bagus et al.; as already noted, the experimental value is 2.78 eV. The calculations show that the ground state undergoes a transition from ionic to covalent structure between 5 and 5.5 a.u. Thus in

tracing the energy path from the equilibrium structure to atomic fragments, one must in effect calculate the electron affinity of atomic fluorine, which is notoriously difficult to do; this is the major source of error in the calculations.

The optical absorption spectrum of XeF_2 shows a weak absorption that peaks at 5.4 eV. Referring again to Fig. 7-137, we can identify this as the transition to the first

Fig. 7-138. Symmetric dissociation of XeF_2 : potential energy curves and transition moment.

Fig. 7-139. Bending of XeF_2 : potential energy curves and transition moment.

$^1\Pi$ state. The calculated vertical energy for this transition is 5.9 eV. This transition is vibronically allowed, but symmetry forbidden, since at the equilibrium geometry the first excited state has $^1\Pi_g$ symmetry. [The transition is essentially $3\pi_{(u)} \rightarrow 4\sigma_{(u)}$.] The ground state and $^1\Pi$ excited state both dissociate to neutral atoms. Note in Fig. 7-137 that the asymptotes for these two states are split by ~ 1 eV, because one Xe-F bond distance is frozen at 1.98 Å. Thus, the 1-eV asymptotic splitting simply reflects the $1^2\Sigma^+ - 1^2\Pi$ splitting of XeF at this distance.

Further information about the 5.4-eV vibronic band is obtained by allowing the molecule to bend. Under bending, the ground state A_1 remains nondegenerate, while the $^1\Pi$ state splits into A_2 and B_1 states, the latter being dipole-connected to the ground state. Figure 7-139 shows the ground, B_1 , and A_2 bending curves and the $A_1 \rightarrow B_1$ dipole transition moment. The electronic transition energy decreases with increasing bond angle, so we expect the 5.4-eV absorption band to broaden and shift to lower energies as the temperature is increased.

It is important to bear in mind that the equilibrium XeF₂ bond distance $R_e = 1.98$ Å is considerably shorter than the R_e for the XeF ion-pair states. Thus, although Fig. 7-137 implies that a vertical transition to the excited $^1\Sigma$ state produces XeF* below the XeF*($B^2\Sigma$) + F asymptote, this is true only for a constrained XeF*. The asymptote for relaxed XeF* + F lies some 2 eV lower. Thus photodissociation will also involve a rearrangement of the diatomic fragments and may produce a substantial fraction of vibrationally excited XeF*. We plan further studies of the photodissociation dynamics to examine this question.

Author: T. N. Rescigno

Theoretical Atomic and Molecular Physics

Multiphoton Ionization Theory

Anomalous strong laser-atom interactions for which no adequate theoretical

treatment exists have been observed. Collective motion of atomic electrons (in effect, an "atomic plasmon") induced by a strong laser field is the most probable explanation for the phenomena. Such collective behavior would be of fundamental physical significance, and it could have profound consequences for the development of UV and x-ray lasers, because it might permit population inversion on an atomic scale rather than requiring large-scale recombination.

Theoretical approaches to this problem must necessarily include strongly nonlinear photon-electron effects that have not been addressed previously. We are developing and applying nonanalytic numerical methods for solving this problem; here we briefly describe recent results that indicate the power and scope of the method in simple atomic systems.

Unusual nonlinear effects have been observed in several multiphoton ionization experiments.¹⁹³⁻¹⁹⁵ These investigations showed greater production of highly ionized species than is predicted by linear theory, thus suggesting a resonant phenomenon. For example, the relative abundance of Xe^{q+} ($q = 1, \dots, 6$) produced in the multiphoton ionization of xenon is predicted by standard linear (single-electron) theory to drop strongly with q : the populations should vary as $1:10^{-4}:10^{-8}$ for $q = 1, 2, 3$, and they should become undetectable for $q > 4$. Experimentally, the relative abundances of the higher charge states are many orders of magnitude greater than predicted, decreasing monotonically from 44 for $q = 1$ to 1 for $q = 6$. This behavior depends on the atomic number, with an apparent change in ionization mechanism occurring between argon ($Z = 18$) and krypton ($Z = 36$). The effect depends strongly on the number of inner-shell electrons.

Collective electronic effects have been suggested in simpler single-ionization experiments on barium.¹⁹⁶ These experiments were stimulated by the theoretical prediction of resonant phenomena.¹⁹⁷ Although the details of this particular theoretical model are still in dispute, the suggestion that a dipole resonance is involved appears viable. More refined calculations than those of Ref. 197, using a time-dependent density functional

approach with the local-density approximation for exchange effects,¹⁹⁸ yield quantitative agreement with the observed barium spectra and demonstrate the collective nature of the excitation. It is thus highly plausible that similar phenomena exist in multiphoton ionization.

Since the problem is strongly nonlinear, no analytic technique (almost always perturbative or linear) is likely to provide a complete description of the intense-field phenomena. Thus, computational methods must be developed for this purpose. In particular, both finite-differencing schemes (the Peaceman-Rachford method) and pseudo-spectral (the Fast Fourier Transform) techniques have been implemented to solve the time-dependent Schrödinger equation. Experience with these techniques, which includes applications to quite complicated systems,^{199,200} suggests that they might be suitable here.

Since it is too difficult to obtain the full solution of the time-dependent Schrödinger equation for intense-field effects in atoms, some physically reasonable nonperturbative approach is required. One relatively simple but qualitatively accurate method is the time-dependent Hartree-Fock (TDHF) approximation.²⁰¹ The method's simplicity arises from its restriction of the entire many-body wave function to single determinantal form; its accuracy arises from the self-consistency demanded of the resulting wave function and the potential it generates. In essence, the TDHF approximation successfully incorporates the basic nonlinearities of the problem and then avoids the limitations of perturbative (linearized) theories.

Within this scheme, the TDHF equations for an atom in a laser field are (in atomic units)

$$i\frac{\partial}{\partial t}\Phi_j(r,t) = \{h + V(r,t) + ezA(t)\cos\omega t\}\Phi_j(r,t), j = 1, \dots, n \quad (99)$$

for each electron's wave function ϕ_j . Here h is the zero-field Coulomb Hamiltonian, $V(r,t)$ is the self-consistent potential, and $A(t)$ is the amplitude of the applied field, which is switched on linearly over 10 cycles of the photon field oscillation.

Fig. 7-140. Time dependence of the probability P_{00} of a hydrogen atom remaining in its ground state in a 10^{14} -W/cm² laser field of 15-eV photons.

As a first test of the method, we applied these equations to the hydrogen atom: they were solved numerically on a cylindrically symmetric grid using the Peaceman-Rachford propagation scheme. The energy of the impinging photons was taken to be 15 eV, and the intensity of the photon beam was 10^{14} W/cm². One 15-eV photon is sufficient to ionize the atom. Figure 7-140 shows the time dependence of the ground state probability P_{00} . After the initial adiabatic turn-on time of the laser, P_{00} shows a steady decline. From the slope of the curve, we obtain a photoionization cross section of 55×10^{-18} cm², in good agreement with the exact value of 50×10^{-18} cm² for 15-eV photons. These calculations encourage the hope that the TDHF method can be successfully applied to multielectron systems in intense laser fields.

Authors: C. J. Cerjan and K. C. Kulander

Angular Dependence of Photoejection

Summary. We have used complex basis function methods to calculate the asymmetry parameter $\beta(E)$ that determines the angular dependence of photoejection in the dipole approximation. Heretofore the basis-set approach has been able to produce only the total photoionization cross section.

Discussion. Photoionization processes play an important role in the operation of short-wavelength lasers. Over the last decade, several approaches to the calculation of atomic and molecular photoionization cross sections have been developed that rely entirely on expansions in terms of

square-integrable basis functions and that extract the ionization probability from a diagonalization of the resulting matrix representation of the electronic Hamiltonian. These techniques are practical because they rely on bound-state computational methods and thereby allow the treatment of polyatomic molecules and the use of higher levels of approximation in both the bound (initial-state) and continuum (final-state) portions of the calculations.

The Stieltjes moment theory method and its variants²⁰² constitute one of the two most successful approaches. This theory has been applied with remarkable success to a large number of systems, mostly within the separated-channel static-exchange approximation. More recently an attractive alternative, the complex basis function method,²⁰³ has been developed; it provides reliable results at computational effort similar to that of moment-theory calculations even when narrow resonance features dominate the cross section. Several important formal properties of the complex basis function method allow it to be used as a general approach for computing collision amplitudes; we have exploited those properties in the theory presented here.

However, both of these approaches have produced only total photoabsorption cross sections, and have provided no information about partial cross sections (except through the use of approximate schemes for separating different channel contributions) or about the angular dependence of photoejection. The origin of this limitation is perhaps illustrated most succinctly by the expression that provides the starting point of the complex basis function approach. The cross section $\sigma(\omega)$ is written as the imaginary part of a matrix element of the full Green's function or resolvent,

$$\sigma(\omega) = -\frac{4\pi\omega}{c} \operatorname{Im} \lim_{\eta \rightarrow 0} \times \langle \Psi_0 | \mu [E_0 + \omega - H + i\eta]^{-1} \mu | \Psi_0 \rangle, \quad (100)$$

where ω is the frequency of the incident radiation, c is the speed of light, Ψ_0 is the initial-state wave function of energy E_0 , μ is the component of the electric dipole operator in the direction of the applied field,

and H is the target Hamiltonian. This expression produces only the modulus squared of the integral over ejection directions and the sum over ionization channels, and thus yields only the total cross section.

On the other hand, the differential cross section for photoejection (in the dipole approximation) has the form²⁰⁴

$$\frac{d\sigma}{d\Omega} = \frac{\sigma_{\text{tot}}}{4\pi} [1 + \beta(E) P_2(\cos \theta)], \quad (101)$$

where $P_2(\cos \theta)$ denotes the Legendre polynomial of order 2. The asymmetry parameter $\beta(E)$ is thus central in determining the angular dependence of photoejection. For a one-electron atom initially in a state of angular momentum l , for example, the asymmetry parameter is given by

$$\begin{aligned} \beta(E) = & \frac{1}{D} \left[l(l-1) d_{l-1}^2 + (l+1)(l+2) d_{l+1}^2 \right. \\ & \left. - 6l(l+1) d_{l+1} d_{l-1} \cos(\delta_{l-1} - \delta_{l+1}) \right], \quad (102) \end{aligned}$$

$$D = (2l+1) \left[l d_{l-1}^2 + (l+1) d_{l+1}^2 \right], \quad (103)$$

where $d_{l\pm 1}$ denotes real-valued radial dipole matrix elements between the initial state and continuum states of angular momentum $l \pm 1$. The cross term in Eq. (102) contains the cosine of the difference of the phase shifts in the $l \pm 1$ partial waves. To evaluate this cross term we need the phases of the complex dipole transition amplitudes, which Eq. (100) does not provide. This deficiency also afflicts the moment-theory approach to the problem.

We have overcome this difficulty through use of the complex basis function method. The power of this method is that it can be used to provide an approximation to essentially any matrix element of the full Green's function. We have exploited this property to develop a procedure for computing the differential cross section $d\sigma/d\Omega$ or the asymmetry parameter $\beta(E)$.

In the dipole approximation, the time-dependent wave function that contains

all the information about the ionization process is

$$\Phi(t) = e^{-iHt} \mu \Psi_0. \quad (104)$$

The total cross section, for example, is related to $\Phi(t)$ by

$$\sigma(\omega) = \frac{4\pi\omega}{c} \text{Im} \int_{-\infty}^0 e^{i(E - i\eta)t} \langle \Psi_0 | \mu | \Phi(t) \rangle dt. \quad (105)$$

To extract the angular dependence of photoionization or photodetachment from Eq. (100), we must find the flux (or probability current density) in a particular direction due to $\Phi(t)$. Let $F(\mathbf{r}_0)$ be the symmetrized radial flux

$$F(\mathbf{r}_0) = \delta(\Omega - \Omega_0) \frac{1}{2m} \times [\delta(r - r_0) \hat{\mathbf{p}} + \hat{\mathbf{p}} \delta(r - r_0)], \quad (106)$$

where $\hat{\mathbf{p}}$ is the radial momentum operator. This operator gives the flux in the direction Ω_0 evaluated at the distance r_0 from the center of the coordinate system; it is simply the polar-coordinate representation of the flux or current density operator of familiar textbook discussions. The radial flux $f(t, \mathbf{r}_0)$, evaluated at time t at distance r_0 and in direction Ω_0 due to the wave function in Eq. (104), is just

$$f(t, \mathbf{r}_0) = \langle \Phi(t) | F | \Phi(t) \rangle \\ = \langle \Psi_0 | \mu | e^{iHt} F e^{-iHt} | \mu \Psi_0 \rangle. \quad (107)$$

The total flux for detachment in a given direction can be written either as an integral over this time-dependent quantity or as an integral over all ejection energies of the flux at a fixed value of the energy, $P(E, \mathbf{r}_0)$:

$$f_{\text{tot}} = \int_0^\infty f(t, \mathbf{r}_0) dt = \int_0^\infty P(E, \mathbf{r}_0) dE. \quad (108)$$

Here $P(E, \mathbf{r}_0)$ is given by

$$P(E, \mathbf{r}_0) = \frac{1}{2\pi} \langle \Psi_0 | \mu \| E - H - i\eta \|^{-1} \\ \times F \| E - H + i\eta \|^{-1} | \mu \Psi_0 \rangle. \quad (109)$$

The differential cross section for photodetachment can be easily related to the flux: in atomic units, it is given by

$$\frac{d\sigma}{d\Omega} = \frac{2\pi\omega r_0^2}{c(2l+1)} \sum_{m_l} \langle \Psi_{m_l} | \mu \| E - H - i\eta \|^{-1} \\ \times F \| E - H + i\eta \|^{-1} | \mu \Psi_{m_l} \rangle. \quad (110)$$

The sum over m_l is over m_l sublevels of the initial state of the atom.

Our objective has been to write the differential cross section in terms of a matrix element between bound functions of an operator expression involving only the Green's function: Eq. (110) is just such an expression. Equation (110) can be evaluated with complex basis function techniques in the same way in which those techniques allow us to approximate bound-bound matrix elements of the Green's function.

This whole analysis works because the complex basis function procedure effectively allows us to compute the quantity $\langle \mathbf{r} | (E - H + i\eta)^{-1} | \mu | \Psi \rangle$ for real values of r , as opposed to complex values ($re^{i\theta}$), which would result from complex coordinate rotation. Furthermore, after partial-wave analysis, this amplitude behaves asymptotically as

$$\langle \mathbf{r} | E - H + i\eta |^{-1} | \mu | \Psi_0 \rangle \\ \xrightarrow{r \rightarrow \infty} -\frac{2}{k} \sum_{l,m} h_l^+(k, r) d_l \exp(i\delta_l) C_{lm} Y_{lm}(\hat{\mathbf{r}}), \quad (111)$$

where $h_l^+(kr)$ is an irregular Coulomb function (or $\mathcal{P}$ Hankel function, in the case of photodetachment of an anion) with *no phase shift*, and C_{lm} is a collection of constant factors.

Equation (111) provides an alternative to the flux expression in Eq. (110) for computing the asymmetry parameter. Projecting on Eq. (111) with a spherical harmonic and dividing by the irregular Coulomb function (containing no phase shift), we immediately obtain the radial dipole matrix element d_l and its (Coulomb plus short-range) phase δ_l .

We used both the flux expression in Eq. (110) and the matching procedure

Fig. 7-141. Asymmetry parameter $\beta(E)$ for photolization of 2p state of atomic hydrogen. Matching procedure based on Eq. (111) used 60 and 70 Laguerre functions in s and d channels, respectively; extrapolation of flux expression, Eq. (110), used 70 Laguerre functions.

based on Eq. (111) to compute the asymmetry parameter $\beta(E)$ for ionization of the 2p state of atomic hydrogen. We used overly large basis sets of complex Laguerre functions so as to demonstrate conclusively the validity of this approach.

Figure 7-141 shows the results of these calculations. Not surprisingly, the interference terms in Eq. (110) approach their asymptotic values only as $1/r$, as do the interference terms in the outgoing flux of any scattering wave function. Thus, the flux expression must be extrapolated to very large values of r to yield the correct results, and this extrapolation is less stable at higher energies. The matching procedure, on the other hand, can be performed at points near the target as long as they are outside the range of $\mu\Phi_0$.

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Wave-Packet Theory of Molecular Photodissociation

Molecular photodissociation can play a role in exciting gas laser systems and in initiating chemical reactions. Efficient methods for modeling the dynamics of dissociation help us understand the characteristics of the potential energy surfaces, which yield particular fragment angular, velocity, and internal-state energy distributions. This understanding improves our ability to develop laser photochemistry or to produce photon-pumped laser systems. Because of the high energy required, the photodissociation of polyatomic molecules can often lead to

final states with three fragments. The dynamics of such a dissociation cannot be handled using standard, time-independent molecular scattering techniques. Under these circumstances time-dependent wave-packet methods are required.

Pattengill²⁰⁵ has studied the three-body direct dissociation of cadmium dimethyl [$\text{Cd}(\text{CH}_3)_2$] using classical trajectory methods; the relative weights of the trajectories were determined from a Wigner distribution function for the initial state. Classical mechanics should be very accurate for the dynamics of this system because of the high energies and large masses involved. Only the uncertainty in selecting the initial conditions makes the classical results suspect.

In Pattengill's model, cadmium dimethyl is treated as pseudo-triatomic, with each methyl (Me) group treated as an "atom" of mass 15 amu:

It is assumed that the C-Cd-C bond angle remains fixed in the collinear geometry of the initial, ground electronic state. Thus, the problem is reduced to the photodissociation of a collinear triatomic molecule: the system has just two degrees of freedom, corresponding to symmetric and asymmetric stretch motions. The ground electronic state is bound; the observed Cd-Me bond length is $3.99a_0$. The ground vibrational state is assumed to be harmonic, with frequencies of 465 and 538 cm^{-1} for symmetric and asymmetric stretching, respectively. The excited-state surface is completely repulsive, and has a saddle point at a Cd-Me bond distance of $5a_0$. The two-body dissociation valleys lead to CdCH_3 in its ground electronic state.

We have carried out a time-dependent wave-packet calculation of the dissociation dynamics for this model system. The initial-state wave packet is localized around the ground-state equilibrium separation. During dissociation, the final-state wave packet quickly traverses the saddle-point region, so that only a limited amount of spreading occurs in the asymmetric stretch direction. The photodissociation therefore results in three fragments.

Fig. 7-142. Normalized cadmium atom velocity distributions for the photodissociation of $\text{Cd}(\text{CH}_3)_2$ at three photon wavelengths.

When the propagation on the excited-state surface is complete, the final wave packet is projected onto plane waves and the distributions of fragment velocities are determined. It has been shown²⁰⁶ that the partial photodissociation cross sections are proportional to the product of the photon frequency and the square of the projection onto the final state. The total absorption cross section is the sum (or, in this case, the integral) over the final states at the given total energy.

We compared our calculated fragment velocity and energy distributions with those of Pattengill for the ground vibrational level of the initial electronic state. The continuous distributions obtained from the wave-packet calculations were binned to permit comparison with the classical results.

Figure 7-142 shows the classical and quantum-mechanical cadmium velocity distributions (normalized so that the integrals over the distributions are all unity) for three photon wavelengths. Velocities are expressed as fractions of the maximum velocity energetically allowed, v_0 . The maximum velocity for the cadmium is that for completely asymmetric dissociation. The agreement between these results is very good.

These calculations show the power of the time-dependent wave-packet method. To obtain the fragment distributions classically, roughly 1000 trajectories must be run for each photon energy of interest. The wave-packet propagation requires a

computational effort equivalent to only five trajectories and provides results for all photon wavelengths from a single calculation. Furthermore, the evolution of the wave packet gives detailed insight into the dynamics of the dissociation.

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Time-Dependent Theory of Gas/Surface Accommodation

Despite several attempts to model the phenomenon,²⁰⁷⁻²¹¹ the mechanism by which a gas particle transfers its energy to a crystalline surface is imperfectly understood. A low-mass atom scattered or adsorbed by a low-temperature surface composed of high-mass atoms must be treated quantum-mechanically: The propagation of the scattering particle must be given by the Schrödinger equation, and the bulk atomic motion must be described by quantum mechanics as well.

A practical instance in which these concerns arise is the flow of a gas along a surface that may be at a different temperature than the gas. The macroscopic transfer of energy to or from the surface has macroscopic consequences for the flow of the gas. In the usual continuum description of the flow, the temperature jump at the surface alters the boundary conditions of the

hydrodynamic equations, which then modify the macroscopic properties of the moving particles. A convenient dimensionless parameter that links the microscopic and macroscopic descriptions is the thermal accommodation coefficient $\alpha(T_s, T_g)$, which is usually a function of both the gas temperature T_g and the surface temperature T_s . The conventional definition of α is²¹²

$$\alpha(T_s, T_g) = \frac{\bar{E}_i - \bar{E}_f}{2k(T_g - T_s)}, \quad (113)$$

where $\bar{E}_i$ and $\bar{E}_f$ are the initial and final energies of the scattering particle and k is the Boltzmann constant. If $T_g = T_s = T$, then α is known as the equilibrium accommodation coefficient (EAC) (Ref. 212):

$$\alpha(T) = \lim_{T_g \rightarrow T_s} \alpha(T_s, T_g). \quad (114)$$

The EAC is the quantity measured in several experimental situations, especially in flash desorption,²¹³ in which the gas and surface are thermally equilibrated and then the vacuum chamber is suddenly pumped to reduce the gas pressure. The desorbing particles are collected and the surface coverage is then monitored as a function of time.

Although the first theoretical calculations²¹⁴ date from 1935, there is still qualitative disagreement between theory and experiment as the temperature of the gas is decreased.²¹⁵ For helium on tungsten, the theoretical prediction is that $\alpha(0 \text{ K}) = 0$, but the experimental data indicate that $\alpha(0) \neq 0$. This behavior persists even with the inclusion of a binding potential well in one dimension²¹⁵; it has therefore been speculated that one-dimensional theories are inherently deficient.

On the other hand, the theoretical predictions^{214,215} were perturbative results that did not incorporate higher phonon-gas couplings, as is done in Refs. 216 and 217. More extensive, but still perturbative, computations²¹⁸⁻²²³ indicate that one-dimensional models can be used successfully, at least for very weakly bound gas/surface systems. The extension of this reasoning requires numerical rather than analytical techniques for the solution of the Schrödinger equation. To bypass the

limitations of simple time-dependent or time-independent perturbation theory, the Schrödinger equation must be solved for a gas particle interacting with a quantum-mechanical reservoir of phonons.

To solve the time-dependent Schrödinger equation, we use a Fast Fourier Transform (FFT) technique.²²⁴ The phonon bath is treated by a semigroup expansion technique^{225,226} similar to that used in studies of vibrational relaxation of liquids.^{227,228} We consider couplings between the gas particle and phonons only to second order, but this is not an essential limitation of the technique. The model potential we use supports 11 bound states, indicating that the previous restriction²¹⁸⁻²²³ to a few shallow bound states can be lifted. The main qualitative result of this work is that one-dimensional quantum-mechanical models can be usefully employed in the examination of the EAC without many of the limitations, or complicated expressions, arising from the use of perturbation theory.

We employed a consistent one-dimensional model to study the dynamic (time-dependent) scattering of helium atoms by tungsten as a function of temperature. The initial wave packet was chosen to have a Maxwell-Boltzmann distribution in momentum space. That is, the width and normalization of the wave packet varied with temperature as

$$\begin{aligned} \psi(t=0, z=+\infty) &= \left(\frac{mkT}{\pi} \right)^{1/4} \\ &\times \exp(-z^2/2mkT) \exp(ip_0 z), \end{aligned} \quad (115)$$

where z is the normal direction to the surface. This choice mimics the conditions of most desorption and accommodation experiments. As the temperature approaches 0 K, a Bose-Einstein distribution must be used for the incident helium atoms. For the temperatures we considered, the Bose-Einstein and Maxwell-Boltzmann distributions are indistinguishable.

The wave packet was propagated by the time-dependent Schrödinger equation

$$\begin{aligned} i \frac{\partial \Psi}{\partial t} &= -\frac{1}{2m} \frac{\partial^2 \Psi}{\partial z^2} \\ &+ [V_0(z) + V_s(z, x(t))]\Psi; \end{aligned} \quad (116)$$

the FFT algorithm was used to evaluate the kinetic energy operator, and second-order differencing was used for the time dependence.²²⁴ The time dependence is implicit, arising from the phonon motion $x(t)$. Time evolution was determined by a harmonic model with initial conditions dictated by the known equilibrium values.

Calculations were performed for $T = 35$, 75, and 115 K. Initial runs at these temperatures were made with the time dependence removed—that is, for a completely static surface. To the accuracy of the calculations, no bound-state trapping occurs. In all cases, substantial average energy change occurs, as shown in Table 7-21. The amount of energy transferred to the surface remains essentially constant as the

$T(K)$	$\Delta \bar{E}_{\text{final}}$	$\Delta \bar{E}/E_{\text{init}}$
35	1.50×10^{-5}	0.112
75	1.95×10^{-5}	0.066
115	2.55×10^{-5}	0.056

temperature is decreased. These results are in accord with those of similar time-dependent perturbation approaches,²¹⁸⁻²²³ which indicate that it is not the one-dimensional aspect of the previous models that is deficient but rather the limitations of the perturbative treatments.

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Table 7-21. Average and fractional energy changes (Hartree units) after 16 ps for three temperatures.

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The Pulse*Star ICF reactor (right) is one innovative approach to electricity generation that employs a fusion chamber submerged in a liquid-metal coolant and a wetted screen for first-wall x-ray protection.

The Cascade ICF reactor concept, shown below and on the opposite page, is one of the most safe, promising, and economically competitive systems for future electricity generation. An artist's rendition demonstrates on these two pages a possible application of a Cascade reactor in a power plant. Multiple laser beams generated in the building at the bottom, right, reflect off mirrors shown at the bottom, left, and enter the ceramic reactor, which is displayed in detail below.

ICF Applications

ICF Applications

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ICF Applications

Introduction

W. J. Hogan

Inertial fusion power reactors will be safer and potentially less harmful to the environment than fission or coal plants. The major thrust of ICF Applications this year has been to demonstrate that ICF power reactors can be economically competitive as well. Our design for the Cascade reactor has established a new benchmark and standard for ICF reactors: A future design will receive attention only if it is likely to outperform Cascade. This year we also examined trends in the power industry, related them to fusion reactor design, and developed the methodology for future comparisons. Finally, we began to examine new reactor ideas and to consider how to develop ICF for multiple applications.

The Cascade reactor is the first ICF design that is economically competitive with fission and coal plants while retaining ICF's desirable characteristics—safety and minimal environmental impact. Cascade's design, which converts energy with about 55% efficiency and at reasonable cost, has withstood much scrutiny as we and others have looked, unsuccessfully, for fatal flaws in the concept. GA Technologies, in a just-completed major study of the reaction chamber concept and the balance of plant, concluded that the design is feasible, and Cascade's cost of the electricity would compare favorably with future fission and coal power plants.

This section includes articles that describe the current version of Cascade and analyze specific issues studied this year. Specifically, we added an inner carbon layer to the blanket to absorb and transport the surface heat flux. Carbon was chosen because the alternatives—diatomic ceramics—would give rise to partially dissociated gases that will not readily recondense under the conditions found in Cascade. Activation in this new blanket is acceptably low. New heat exchanger designs will allow higher energy-conversion efficiency at reasonable cost. The tritium inventory of a Cascade plant would be extremely low, about 260 g (only 140 g would be liberated from the granules over a 30-day period in a worst-case release scenario), and the routine tritium permeation would be negligible (less than 25 Ci or 2.5 g per day). Our final cost analysis indicates that electricity from a Cascade power plant would cost the same as—or somewhat less than—power from future fission and coal plants. With this work we completed our study of Cascade; we are now preparing a final report.

The earlier Pulse*Star reactor was conceived as a way to reduce the costs of a liquid-metal-based fusion reactor. Our current Pulse*Star design increased the pulse rate, eliminated liquid lithium (a fire hazard), eliminated the intermediate heat exchange loop, and adopted a pool-type design (similar to that of Superphenix, the French LMFBF) to reduce the number of pipes and connections required.¹ This year we pursued design solutions to three issues: structural support for the LiPb pool, design of the inner screen that defines the LiPb droplet region, and containment of tritium. Although design solutions seem possible in every case, Pulse*Star would not be competitive with Cascade because Pulse*Star is still based on the steam cycle. Therefore, even though the Pulse*Star design study is far from complete, we have decided to stop work on it.

Several general studies contributed to our understanding of what is required to be economically competitive. We commissioned a study of trends in the power industry (e.g., use of electricity and power plant size). The results demonstrate that while the current emphasis is on smaller plants, the definition of "smaller" is changing. Fission-reactor designers now specify plants that produce 100 to 300 MW_e, but by the time fusion plants are ordered, plants that produce 500 to 1500 MW_e are likely to be considered small. In this year's economic studies, we concluded that a laser fusion power plant can compete with fission and coal plants if the recirculating power fraction is kept below 25% and the driver costs between \$0.4 and \$0.6 billion (\$130 to \$200/J) for a 1-GW_e plant. Starlite, a new reactor design concept, represents one of the new breed in which we attempt to exploit features unique to the ICF process to design a balance of plant much cheaper than the steam cycle. In the Starlite reactor, helium absorbs fusion neutrons and converts their energy into excimer states that will radiate with high efficiency in a very narrow bandwidth. Specially designed PN junctions could then convert this radiation to electricity with reasonable efficiency. The net efficiency of this several-step process is much less than that of Cascade. However, this process might be considerably less costly than Cascade, so the idea is worth pursuing. Unfortunately, the use of helium presents formidable problems in the design of the containment vessel.

The final article in this section deals with issues for developing ICF facilities for use by a number of possible applications including weapons physics, energy and materials production, and space propulsion. If this strategy is followed, it will be important to design the development facilities (especially the high-gain test facility) with multiple test chambers and very flexible drivers.

Cascade Reactor

Introduction

Cascade is a reactor concept that offers the promise of efficient, economical, and safe conversion of inertial confinement fusion energy into electrical power. This year, we completed the Cascade design and began our final report; the following ten articles are a summary of our progress. We are continuing technical work in two areas that remain as issues with Cascade: recondensation of vaporized blanket material and granule fragmentation—see "Condensation of Vaporized Cascade Blanket Material" and "Strain Energy Effects on Dynamic Fragmentation of Blanket Granules," later in this section.

The Cascade reactor (see Fig. 8-1) is double-cone shaped and contains a 1-m-thick flowing granular blanket that serves three functions:

- Absorbs and transfers fusion energy to heat exchangers.

- Breeds tritium to replace fuel burned in the fusion reactions.

- Protects the reactor wall from pressure pulses caused by fusion reactions (otherwise, high stresses would result).

The blanket is made of three high-temperature materials, and its profile of velocity vs radius is sufficiently similar to the radial heating-rate profile to permit high efficiencies to be obtained. Indeed, the velocity profile has been experimentally verified. Further, we believe that material vaporized from the inside layer of granules after each fusion pulse will recondense within the 200-ms interpulse time, so that a 5-Hz repetition rate of the reactor can be achieved.

Granule fragmentation may affect economics if excessive remanufacturing of comminuted granules is necessary. However, our work indicates that only the inner layer of granules—which are exposed to x rays and fusion-fuel pellet debris—could fragment.

The average blanket outlet temperature is high (1440 K), but ceramic materials

Fig. 8-1. Cascade: A rotating, ceramic-granule-blanket reactor.

used in the reactor walls can withstand the high temperatures in the adjacent blanket. Accordingly, no active cooling is required. We use a three-zone ceramic-tube heat exchanger to transfer energy from the blanket granules to high-pressure (5 MPa) helium gas. The power-conversion system is a simple once-through helium gas turbine (Brayton) cycle without reheats.

Activation is dominated by ^{24}Na that is produced in the blanket by (n, α) reactions with the aluminum in the lithium aluminate (LiAlO_2) tritium-breeder zone. The 1-GCi total radioactivity in Cascade at shutdown is lower by a factor of 4 than that in recent conceptual magnetic-fusion power plants operating at the same net electrical power.

The power-conversion system efficiency of 55% is made possible by the high temperatures present in the blanket granules and by an efficient heat exchanger that transfers energy directly from the granules to high-pressure helium gas without an intermediate loop. The net power-plant efficiency is 49%. The balance of plant capitalizes on the inherent safety of Cascade, which allows most of the buildings

and equipment to be built as nonnuclear components; this leads to lower capital costs. Because the Cascade design includes simplicity, high efficiency, and many safety features, its cost is competitive with those of high-temperature gas-cooled reactors, pressurized-water reactors, and coal-fired power plants. In addition, tritium inventory, permeation, and recovery offer no restrictions for the Cascade concept.

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Final Reactor Design

The final design of the Cascade reactor²⁻¹² (see Fig. 8-1) takes the form of a double cone with a maximum radius of 5 m; each cone has a half-angle of 35° and a half-length of 4.3 m. The reactor is evacuated to permit efficient illumination of the fusion-fuel pellets by laser or ion beams (to 13 Pa if laser beams are used). The reactor walls are 20-mm-thick silicon carbide (SiC) panels held in compression by SiC-fiber/aluminum-composite tendons

that gird the chamber circumferentially and axially. As the reactor rotates at 50 rpm, a flowing blanket is held against the inside wall by centrifugal force.

The blanket absorbs 95% of the fusion energy, breeds tritium by neutron interaction with lithium, reduces neutron activation in the reactor wall, acts as a heat-transfer medium, and ameliorates wall stresses. The blanket is made of three granular materials (see the following article, "Flowing Granular Blanket," for a complete description), with the granule density lowest in the interior and greatest in the region adjacent to the reactor wall. Each blanket layer or zone floats on the next one radially outward, and mixing between the three zones is minimized. The granules enter the reactor through ports at the smaller-radius ends and flow along the reactor wall to exit shelves that rotate with the reactor at the larger-radius center. As the granules flow through the reactor, they absorb energy from 300-MJ fusion fuel pellets that are injected into the center of the reactor at 5 Hz.

One-third of the fusion energy produced in the pellets leaves in the form of x rays and pellet debris and is deposited in the first few microns of the surface layer. The rate of energy deposition in this thin region is high enough to vaporize about 1 kg of material, which in turn fills the center of the reactor, cools by radiative heat transfer, and then recondenses on the remaining granules at the inner surface on the blanket. In the final design, we chose elemental pyrolytic carbon for the surface-layer material because it can withstand temperatures of 1600 K without adversely interacting with the adjacent material, and because an elemental material will recondense in a similar form, whereas diatomic substances will dissociate (see "Condensation of Vaporized Cascade Blanket Material," later in this section). Although it takes up to 100 ms for the pyrolytic carbon vaporized from the inner blanket surface to recondense from the central volume of the reactor, calculations⁷ show that once the vapor reenters the surface layer, it recondenses in less than 10^{-5} s.

The remaining two-thirds of the fusion energy takes the form of energetic 14-MeV neutrons and is deposited throughout the blanket with a profile that decreases exponentially with radius. The BeO front zone acts as a neutron multiplier for the outer LiAlO_2 tritium breeder zone.

Final design parameters are listed in Table 8-1. Note that we obtained a power conversion efficiency of 55%, which results in a net plant efficiency of 49%. These high efficiencies are possible because of the high blanket temperatures permitted when ceramic materials are used, and they compare favorably with the 35% net plant efficiency normally found in coal-fired power plants producing the same net electrical power. This relative efficiency, plus the inherent safety of Cascade (see "Balance of Plant," later in this section), reduces the cost of the Cascade power plant to a level competitive with light-water reactors and fossil-fuel-fired power plants (see "Cost Analysis," later in this section).

We have completed additional calculations⁸ concerning recondensation of material contained in the fusion fuel pellets. Suppose we consider that the majority of the mass in each fusion fuel pellet is a

Table 8-1. Final design parameters for Cascade. (See the following article, "Flowing Granular Blanket," for blanket parameters.)

System parameters	Values
Fusion power (MW)	1500
Thermal power (MW)	
Blanket thermal power	1670
Turbine thermal power input	1645
Electrical power	
Net turbine electrical output (MW)	905
Total station electrical demand (MW)	90
Net electrical output (MW)	815
Power conversion efficiency (%)	55
Net plant efficiency (%)	49
Tritium-breeding ratio	1.05
Blanket energy multiplication	1.11
Laser and fuel pellet parameters	
Beam energy (MJ)	1.5
Fuel pellet gain	200
Fuel pellet yield (MJ)	300
Repetition rate (Hz)	5
Laser efficiency (%)	10
Laser electrical power consumption (MW)	75
Reactor	
Wall radius	
Minimum (m)	2
Maximum (m)	5
Wall thickness (mm of SiC)	20
Wall half angle (°)	35
Half length (m)	4.3
Maximum laser beam half angle (°)	10
Clear volume inside blanket inner radius (m^3)	160
Wall surface area (m^2)	230
Rotational speed (rpm)	50
Rotational electrical power drive (MW)	5

heavy material such as lead. Lead will condense if its pressure exceeds the vapor pressure of molten lead at the minimum temperature anywhere in the reactor or pump duct system. Knowing the variation in vapor pressure of lead with temperature, we can calculate the mass of lead (see Fig. 8-2) that would be in the reactor before a liquid phase would condense. At 1800 K (200 K above the maximum blanket temperature), for example, about 100 kg of lead (mostly Pb₁) would be present in the reactor. The amount of lead (or other heavy material) contained in each fusion fuel pellet is unknown, but should certainly be less than this. In any case, in calculating vacuum-pumping requirements, we considered the three cases shown in Table 8-2. We also considered the amount of helium, tritium, and deuterium-tritium that is unburned, generated in the fuel pellet and blanket, or both. A comparison of the results in Table 8-2 with the maximum vacuum pumping of about 200 000 L/s through the ends of the reactor shows that vacuum-pumping requirements can be met for all cases except that of 10 g/s lead at 10⁻⁵ atm. The chamber background pressure is 13 Pa (1.3 × 10⁻⁴ atm), so we have ample margin for vacuum pumping.

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Flowing Granular Blanket

Deuterium-tritium ICF reactors are unique in that a large fraction of the energy yield, typically about 30%, is in the form of a surface heat flux from x rays and the debris of fusion-fuel pellets. The balance of the energy yield is in the form of volumetric heating from neutron absorption.

We redesigned the Cascade blanket this year by changing the surface layer material to carbon. The blanket is now ideally suited to absorb energy depositions: The thin, fast-moving carbon surface layers can absorb and transport the surface heat flux.

Fig. 8-2. Minimum quantity of lead required to establish a liquid-lead phase in the Cascade reactor (volume = 160 m³).

Table 8-2. Pumping requirements for Pb + He + DT + T₂ at 1800 K.

Required vacuum	Added Pb (g/s)	Pb vapor requirement (L/s)	He + DT + T ₂ requirement (L/s)	Total (L/s)
10 ⁻⁵ atm	0.1	7 180	62 000	69 000
	1.0	71 800	62 000	133 800
	10.0	718 000	62 000	780 000
10 ⁻⁴ atm	0.1	718	6 200	6 900
	1.0	7 180	6 200	13 400
	10.0	71 800	6 200	78 000

and the thicker, slower-moving, high-temperature and tritium-breeding layers can absorb and transport the neutron energy. The average blanket exit temperature is 1440 K, which simultaneously permits high power-conversion efficiency (55%) and adequate tritium breeding.

The blanket,^{3,13} shown in Fig. 8-3 and characterized in Table 8-3, is divided into three regions called (from outside in) the breeder zone, the front zone, and the surface layer. Each zone is composed of granules of high-temperature materials; all granules are 1 mm in diameter—about the size of sand grains.

The breeder zone, the outer 900 mm of the blanket, is composed of flowing lithium aluminate (LiAlO₂) granules that absorb neutrons, capturing their energy in the form of heat and producing tritium to replace that burned in the fusion reactions. We require that the tritium breeding ratio (tritium bred to tritium burned) be greater than unity, so that no additional tritium need be supplied to the reactor. Granules from the breeder zone leave the reactor at 1355 K and are used in the first section

Fig. 8-3. The Cascade blanket.

Cascade Reactor

Table 8-3. Characteristics of the Cascade blanket.

	Surface layer	Front zone	Breeder zone
Material	C	BeO	LiAlO ₂
Thickness (mm)	10	90	900
Inlet temperature (K)	1380	1410	1285
Outlet temperature (K)	1600	1510	1355
Total flow rate (m ³ /s)	1.0	1.3	8.4
Power fraction	0.28	0.13	0.59
Peak exit velocity (m/s)	1.9	0.22	0.22

of the heat exchanger for initial heating of cold helium gas returning from the turbine-generators.

The front zone, a 90-mm-thick middle region, is made up of beryllium oxide (BeO) granules. BeO acts as a neutron multiplier and makes it possible to obtain tritium-breeding ratios greater than unity in the breeder zone. The BeO granules are slightly less dense than the LiAlO₂ granules, so that they "float" on the surface of, and do not mix with, the LiAlO₂. The front-zone granules, which leave the reactor at 1510 K, are used for intermediate heating of the helium gas in the heat exchanger. The velocity distribution of the front- and breeder-zone granules in the reactor is continuous and decreases with radius. There is about an order of magnitude change from the highest to lowest velocity.

The 10-mm-thick surface layer is composed of pyrolytic carbon granules that absorb the x rays and fuel-pellet debris. These granules move with a velocity about an order of magnitude greater than that of granules in the front zone, and leave the reactor at 1600 K. These granules are used for final heating of the helium gas in the heat exchanger. The high velocity of the surface layer—discontinuous from the velocity in the front zone—has been verified by experiments¹⁴ of sand flow on inclined

chutes and is necessary to keep the exit temperature of the surface-layer granules within reasonable bounds.

Our ideas about blanket materials have evolved along with the Cascade concept. Originally we proposed Li₂O granules to achieve a tritium-breeding ratio greater than unity without a neutron multiplier. A detailed study³ of material properties and allowable temperatures showed Li₂O could not exceed 1200 K without excessive LiOT mass transport out of the blanket. This temperature limit was excessively restrictive, and we finally selected the three-material blanket described above. Table 8-4 lists the maximum material compatibility limits for the blanket materials.

We selected pyrolytic carbon for the surface layer because the alternatives, diatomic ceramics such as BeO or SiC, will give rise to partially dissociated gases (Be + 1/2 O₂ and Si + C) that will not readily recondense under the conditions found in Cascade. If recondensation is not completed between fusion pulses, the vacuum pumping capacity required would be prohibitive.

The likelihood of dissociated gases with diatomic ceramic materials can be deduced from thermodynamic considerations alone. The surface layer temperature in Cascade is 1600 K; the corresponding vapor pressures of Be and Si are of the order of 1 000 to 10 000 Pa, and O₂ is, of course, above the critical point. All these vapor pressures are far higher than the 13-Pa maximum pressure inside the reactor necessary to permit laser illumination of the fusion fuel pellets. The gas would, therefore, have to be pumped out. However, the maximum pumping capacity of Cascade is of the order of 4 × 10⁵ L/s, which at 13 Pa and 1600 K corresponds to production of about 10⁻¹ g moles of dissociated gas per pulse, if the repetition rate is 5 Hz. The dissociated material remaining in the reactor just before the next pulse is about 15% of the initially vaporized material, or about 150 g (17 g moles of Be or 5 g moles of O₂). This is far greater than the 10⁻¹ g moles of dissociated gas per pulse that can be pumped out of the reactor. Thus, diatomic or polyatomic surface layer materials would, in general, not be useful in Cascade.

The problem of dissociation, recondensation, and vacuum pumping can be avoided by using pyrolytic graphite as the surface layer. Because the condensation coefficient of monatomic carbon vapor (~ 0.5) is greater than that of the polyatomic vapor (< 0.1), recondensation is assisted by the expected high degree of dissociation.

The proposed 10-mm-thick surface layer of carbon granules is sufficient to condense all the carbon vapor before it penetrates to the BeO front zone below. The granular bed also increases the effective condensation coefficient over that of a smooth wall by about 50%. Details of the recondensation phenomena that reduce recondensation times are discussed later in this section; see "Condensation of Vaporized Cascade Blanket Material."

The carbon/BeO interface limits the maximum permissible temperature in the surface layer. The surface layer has good mixing, and the reaction

prevails at the interface. The temperature must be kept low enough so that the combined Be and CO vapor pressure is a small fraction of the reactor background pressure (13 Pa). We chose to keep this vapor pressure an order of magnitude lower, thereby limiting the interface temperature to 1600 K (see Table 8-4). This temperature is well below the 3000 K bulk temperature permitted in the carbon, and hence limits the peak temperature permitted in the blanket.

Another characteristic of this blanket is a step in the heating rate at the BeO/LiAlO₂ interface. The flow rate must be high enough to keep the LiAlO₂ exit temperature below 1500 K (see Table 8-4). Because the velocity profile is continuous in the front and breeder zones, this flow rate, along with the heating-rate profile, results in average outlet temperatures below the maximum values permitted. The average exit temperature for the BeO front zone is only 1510 K, but the BeO can withstand temperatures up to 2300 K. Similarly, the average exit temperature from the

Material	Temperature limit (K)	Basis
SiC, bulk temperature	2100	Thermal decomposition overpressure of 1 Pa ^a
SiC/LiAlO ₂ interface	1400	Overpressure of 1 Pa
	1500	Stability of SiO ₂ layer on SiC at overpressure of 4 Pa
LiAlO ₂ , bulk temperature	1500	0.8 × melting temperatures (Ref. 4)
BeO/LiAlO ₂ interface	No limit	LiAlO ₂ has much lower (more negative) free energy of formation than Be·Al ₂ O ₃
BeO, bulk temperature	2300	Material phase transition
C/BeO interface	1600	Thermal decomposition of BeO for overpressure limit of 1 Pa
C, bulk temperature	3000	Overpressure limit of 1 Pa

^a Overpressure is the total pressure of gaseous products resulting from the material interaction.

LiAlO₂ breeder zone is 1360 K, 140 K below the allowable 1500 K.

The temperature of the helium in the heat exchanger is reduced by two phenomena: the average exit temperatures for all three zones are lower than maximum, and 59% of the fusion energy is deposited in the lowest-temperature breeder zone. Helium behaves like a perfect gas with constant specific heat, so that the temperature rise in each region of the heat exchanger is proportional to the energy transferred. Because 59% of the energy is associated with the breeder zone, 59% of the helium temperature rise is associated with the corresponding region of the heat exchanger. Further, the temperature of the helium leaving the breeder-zone region of the heat exchanger is only 1160 K (the 1360 K average exit temperature of the breeder zone less 200 K, which we used as a minimum temperature drop across the heat exchanger). Helium will leave the heat exchanger at 1300 K, well below the 2800 K (3000 K less 200 K) possible in the surface-layer region of the heat exchanger. Efficiency of the conversion system increases with temperature. A gross plant efficiency of 55% is feasible with helium at 1300 K. Should we find more acceptable blanket materials with higher allowable temperatures, then a higher helium exit temperature, and a correspondingly increased conversion system efficiency, could be attained.

A zoned blanket has the further advantage of reducing cracking of the surface-layer granules as they absorb x-ray and

Table 8-4. Material compatibility temperature limits for Cascade.

pellet-debris energy. If the surface layer is made of Li_2O , as in the original Cascade design, thermal stresses exceed those allowed and some granule cracking would be anticipated. If the surface layer is composed of carbon, however, little or no cracking is expected.

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Condensation of Vaporized Cascade Blanket Material

The maximum fusion pulse rate that can be achieved in Cascade will be limited by the time needed to remove gases introduced or generated with each pulse. Previous research^{16,17} indicated that vaporization and condensation would not be an issue with a beryllium-oxide surface layer: the condensation time for a laser driver would be about 20 to 25 ms, shorter by an order of magnitude than that required for the desired 5-Hz pulse rate. However, several phenomena that were identified but not included in the earlier model might extend the condensation time. Of particular concern are dissociated gases that might arise from the surface layer of the blanket and inhibit recondensation. In this article we discuss BeO, SiC, and C granules as candidates for this surface layer, evaluating dissociation, ionization, recombination, and sticking coefficients of these materials and examining the interaction of the vapors of these materials with the solid granules to produce incondensable gases. (Thermodynamic properties are shown in Table 8-5.) We chose a 1-cm-thick layer of

carbon granules for the surface layer; the reasons for our choice are summarized in this article.

Amount of Material Vaporized. We estimated, in a fashion similar to that described in Ref. 16, the amount of material vaporized from the inner granule layer by a 300-MJ fusion pulse. We used the BUCKL code to calculate the energy density resulting from x-ray deposition for each candidate material. We assumed that any material whose energy density exceeds the cohesive energy (heat of fusion plus heat to raise material to vaporization temperature plus heat of vaporization) would be vaporized; additional material would also be vaporized because more than enough energy is deposited in it to reach the boiling point. For an upper limit we integrated the energy-density profile between the point at which the energy density just equals the cohesive energy and the point at which the energy density is just sufficient to raise the temperature to the boiling point and assumed that that amount of energy is available to vaporize additional mass. Using these two methods, we calculate that between 0.4 and 1.0 kg of BeO or SiC would vaporize, while for carbon the total mass vaporized ranges from 0.06 to 0.4 kg. The vaporized material has an inward velocity of approximately 10^4 m/s and thus a momentum of about 10^4 kg·m/s. The fusion debris moves outward at a velocity of about 10^7 m/s, but its mass is much less than 1 g. We therefore assume that the debris is completely stopped by the vaporized material. Thus, about 100 MJ of energy (the sum of x-ray and debris energies) less the cohesive energy for each material is deposited in the vaporized material. Pressures of about 10^5 Pa and temperatures of 1 to 3 eV result.

Vapor Flow Into the Blanket. The vaporized blanket material will expand toward the center of the chamber and also into the granular bed. At a jump-off velocity of 10^4 m/s (estimated from the energy density), this material takes about 0.6 ms to rebound from the center and reach the blanket again (for a blanket at 3 m). On the other hand, the mean velocity of a carbon atom in a 1-eV gas is about 4500 m/s, and the 1-cm-thick blanket surface layer may be penetrated faster. The mean free path in the gas after it has expanded to fill

Table 8-5. Material properties of BeO, SiC, and C.

Material property	BeO	SiC	C
Melting temperature, T_m (K)	2800	—	—
Vaporization temperature, T_v (K)	4200	3000	3925
Heat of fusion, H_f (J/kg)	2.6×10^6	—	—
Heat of vaporization, H_v (J/kg)	19.8×10^6	13×10^6	60×10^6
Density, ρ (kg/m ³)	3000	3200	2100
Specific heat at 1270 K, C_p (J/kg·K)	2100	1300	2000
Thermal conductivity at 1270 K, k (W/m·K)	28	50	280

the chamber is about 0.1 mm, while before expansion it is 10^6 times shorter than this. The mean free path for a gas molecule to collide with a granule in the bed is

$$\lambda_g = \frac{4}{\pi d^2 n_g}, \quad (2)$$

where d is the granule diameter (~ 1 mm) and n_g is the number density of the granules. For a blanket with porosity p ,

$$n_g = \frac{6p}{\pi d^3}, \quad (3)$$

so that the mean free path for a porosity of 0.5 and a granule diameter of 1 mm is

$$\lambda_g = \frac{2d}{3p} = 1.2 \text{ mm}. \quad (4)$$

Finally, the gas present in the blanket at the instant of the fusion pulse is at the background chamber density of 10^{20} to 10^{21} m^{-3} —the density range necessary for laser propagation to the target. Because the mean free path of a gas molecule at this density is 6 to 60 mm, the gas expands into the blanket as it is expanding toward the chamber center. The rate of expansion into the blanket can be estimated by the kinetics of gas molecules interacting with the granules. The velocity of the expansion front v_g is given by

$$v_g = \frac{1}{\sqrt{3}} \omega_g \lambda_g, \quad (5)$$

where the $1/\sqrt{3}$ comes from geometric effects and ω_g , the frequency of collisions between gas molecules and granules, is given by

$$\omega_g = n_g \pi \left(\frac{d}{2} \right)^2 v_a, \quad (6)$$

and v_a is the average velocity of a gas molecule. Inserting Eqs. (4) and (6) into Eq. (5) results in $v_g = 0.3 v_a = 1200 \text{ m/s}$ for carbon gas at 1 eV flowing into a granule bed with 0.5 porosity.

For another estimate of the velocity of the gas into the granule bed, we applied Darcy's law:

$$v_g = -\frac{k}{\mu} \frac{\partial p}{\partial x} \quad \text{m/s}, \quad (7)$$

where k is the permeability, μ is the viscosity, and $\partial p/\partial x$ is the pressure gradient. For spherical, 1-mm-diam granules and a porosity of 0.5, the blanket permeability k has been estimated to be $1.5 \times 10^{-9} \text{ m}^2$ (see Ref. 16). For carbon vapor at 1 eV, the viscosity μ has been estimated (in a manner analogous to the method in Ref. 16) to be $2.2 \times 10^{-4} \text{ Pa}\cdot\text{s}$. The initial pressure gradient across the 1-cm surface layer is 10^7 Pa/m ; therefore, from Eq. (7), the initial gas-flow velocity is about 70 m/s.

From these estimates we conclude that the time necessary to penetrate the 1-cm surface layer would be from 0.008 to 0.14 ms, which is much shorter than the time for the gas to expand to the center of the chamber and return to the blanket (0.6 ms). Therefore, we conclude that the gas will be exposed to the entire 1-cm surface layer after 0.1 ms or less. The inner surface of the granules is assumed to be at the vaporization temperature of the material (3925 K for carbon), but the temperature of the outer surfaces of these granules and the surfaces of other granules in the surface layer is determined from neutron-deposition calculations to be about 1600 K. Thus a large, relatively cool surface is available for condensation in 0.1 ms or less.

Condensation. Neglecting dissociation and ionization for the moment, we can estimate the condensation time simply from kinetic theory. If a gas is in equilibrium, the number of molecules hitting a unit area of wall per unit time can be found by integrating the Maxwell-Boltzmann distribution function over velocity space

$$\begin{aligned} & \int_0^\infty \int_{-\infty}^\infty dv_2 \int_{-\infty}^\infty dv_3 n f(\mathbf{v}) v_1 dv_1 \\ &= n \left(\frac{kT}{2\pi m} \right)^{1/2}, \end{aligned} \quad (8)$$

where v_1 is the velocity perpendicular to the wall, v_2 and v_3 are velocities parallel to it, $n = \rho m^{-1}$ is the number density, $f(\mathbf{v})$ is the Maxwell-Boltzmann function, T is the gas temperature, k is Boltzmann's constant, and m is the molecular mass. If we assume that all molecules that hit the wall are condensed (a sticking coefficient of

unity), then the rate of change of mass in the chamber is

$$\dot{M} = -4\pi R^2 m n \left(\frac{kT}{2\pi m} \right)^{1/2} = \frac{4}{3} \pi R^3 \rho. \quad (9)$$

where R is the chamber radius and ρ is the density. Solving this differential equation leads to

$$\rho(t) = \rho(0) e^{-t/\tau}, \quad (10)$$

where the decay time τ is given by

$$\tau = \frac{R}{3} \left(\frac{2\pi m}{kT} \right)^{1/2}. \quad (11)$$

For carbon gas at 1 eV in a 3-m chamber, τ is 1.3×10^{-3} s. The initial number density for carbon is $1.4 \times 10^{23} \text{ m}^{-3}$, and a number density below about 10^{21} m^{-3} is needed to propagate laser pulses. Thus, Eq. (10) suggests a condensation time of about 6 ms.

The hot plasma will radiate energy to the wall and vaporize additional mass. To include this effect we ran the CONRAD code for a carbon wall, just as we had previously done for a BeO wall.¹⁶ CONRAD simulates the behavior of a solid or liquid layer that is exchanging mass and energy with a vapor. One-dimensional Lagrangian hydrodynamics and multifrequency radiation transport are used to model the vapor. The Lagrangian zones are automatically rezoned when mass transfer between the vapor and the solid or liquid significantly changes the mass in any zone. The

maximum permitted mass change in any zone is an input variable. Non-ideal equations of state for use in CONRAD are obtained from the MIXERG code, which calculates ionization level, opacity, etc., to create tables. CONRAD uses a one-dimensional, finite-difference method to determine the time-dependent temperature profiles in the solid or liquid layer.

In CONRAD, heat transfer across the interface between the vapor and liquid (or solid) is the result of mass transfer (either condensation or vaporization), thermal radiation from hot vapor, and x-ray deposition. Each of these phenomena deposits energy either on the interface or within the solid—in the case of x-ray deposition, with a deposition profile of a multigroup sum of exponentials. Vapor within the 1-mfp-thick boundary region exchanges heat and mass with the surface layer of a solid or liquid. The mass transfer is determined by the temperature of the solid or liquid layer at the interface and by the conditions of the vapor in the boundary layer. Since the Lagrangian method is valid only within the hydrodynamic approximation, CONRAD stops executing when the thickness of the boundary layer equals the thickness of the entire vapor region.

To estimate the effect of the granular nature of the blanket, we considered two CONRAD cases, just as was done previously for the BeO calculation.¹⁶ Figure 8-4 shows the results for a carbon blanket with the properties of Table 8-5. Figure 8-4(a) depicts the case with the correct average heat flux into the granular

Fig. 8-4. Condensation of C on the first granule layer for two cases: (a) correct average heat flux, heat capacity $\times 3$, and (b) correct heat capacity, heat flux $\times 9$.

Table 8-6. Equilibrium dissociation fractions α for BeO, SiC, and C.

ρ (kg/m ³)	T (eV)	α_{BeO}	α_{SiC}	α_{C}
3×10^{-3}	1.0	0.2	0.2	0.0
3×10^{-3}	4.0	0.4	0.5	0.5
5×10^{-3}	0.5	1.0	1.0	1.0
5×10^{-3}	1.0	1.0	1.0	1.0
5×10^{-6}	0.1	0.0	0.0	0.0
5×10^{-6}	0.5	1.0	1.0	1.0

and for SiC it is

$$\frac{\alpha^2}{1 - \alpha} = 2.2 \times 10^6 \left(\frac{T_0}{T} \right)^{1/2} \times \left(\frac{\rho_0}{\rho} \right) e^{-4.6 T_0/T}, \quad (14)$$

where the subscripted quantities are the values of the variable at the point of maximum expansion of the gas (i.e., when it reaches the center of the chamber). These approximations are reasonable for values of T less than a few electron volts. Shown in Table 8-6 are the values of α found for BeO and SiC for three values of the gas density corresponding to the initial density at vaporization ($\rho = 3 \times 10^{-3}$ kg/m³), the value at the time of maximum expansion ($\rho = 5 \times 10^{-3}$ kg/m³), and the value at the end of condensation ($\rho = 5 \times 10^{-6}$ kg/m³). For each density in Table 8-6, minimum and maximum gas temperatures (taken from the CONRAD calculations) are used to derive values of α . Table 8-6 clearly shows that the dissociation is complete by the time of maximum expansion. Recalling that the dynamics drives the reaction toward dissociation even more than equilibrium calculations would indicate, we can be assured that dissociation is complete. The products of dissociation for BeO and SiC (with the exception of the carbon) will not condense on the blanket. Therefore, dissociation will effectively prevent condensation until the gas cools sufficiently to first allow recombination. Radiation effectively reduces the gas temperature to 0.5 eV, but becomes very ineffective below that value (recall that radiative power varies as T^4). With neither radiation nor condensation taking place, the gas will cool only very slowly by conduction into the blanket granules. It must cool to less than 0.2 eV for recombination to become significant. In this case the condensation time can be quite long.

The dissociation of carbon is, in general, complex. At high pressure, the gas exists

blanket but with a heat capacity 3 times too large, while Fig. 8-4(b) depicts the case with the correct heat capacity but a heat flux 9 times too large. In either case the number density drops to 10^{21} m⁻³ in about 6 ms. These calculations agree remarkably well with the simple kinetic-theory estimate above, particularly in light of the fact that in CONRAD a great deal of additional carbon is vaporized by reradiation from the gas (see Fig. 8-4).

Dissociation and Ionization. Given that the dissociation energies of BeO and SiC are both approximately 5 eV and that the available specific energy per molecule is 20 to 40 eV, we should suspect that some dissociation will occur. The relevant dissociation reactions are

The dissociation reactions are two-body interactions, while the recombination reactions must involve three bodies. Detailed kinetic calculations of the two- and three-body collision rates¹⁸ show that immediately after vaporization these rates are $\nu_2 = 10^{11}$ s⁻¹ and $\nu_3 = 10^8$ s⁻¹, and at the end of the expansion into the chamber they are $\nu_2 = 10^8$ s⁻¹ and $\nu_3 = 10^2$ s⁻¹. Because this expansion takes about 0.3 ms, it is safe to assume equilibrium between dissociation and recombination initially. However, as the gas expands, the recombination rate will slow much faster than does the dissociation rate, and the system will thus be driven toward greater dissociation than even equilibrium kinetics predicts. As a lower limit on dissociation, therefore, we estimated the equilibrium dissociation expected at various times.

The equilibrium degree of dissociation α was estimated by constructing the partition functions Q for the various mixtures created by the dissociation/recombination reactions [Eq. (12)] and then minimizing the Helmholtz free energy $F = -kT \ln Q$ with respect to α (Ref. 18). The result for BeO is

$$\frac{\alpha^{3/2}}{1 - \alpha} = 8.7 \times 10^2 \left(\frac{T_0}{T} \right)^{1/4} \times \left(\frac{\rho_0}{\rho} \right)^{1/2} e^{-2.0 T_0/T}, \quad (13)$$

as a mixture of carbon chains C_n where n varies from 1 to about 7 (Ref. 19). At sufficiently low pressure the monatomic form C_1 dominates, and the next most important species is C_3 (rather than C_2) because of the anomalous stability of C_3 . For simplicity, we have considered the equilibrium between C_1 and C_3 only. Using the same methods as were used for BeO and SiC and the data from Ref. 19, we calculated the dissociation fraction for C_1 that is shown in Table 8-6. Similarly to BeO and SiC, the carbon is likely to be completely dissociated by the end of the expansion phase. This is fortunate, since, as we show below, the sticking coefficient for C_3 is much smaller than that for C_1 and the condensation time would be correspondingly longer (the time varies inversely as the sticking coefficient).

Because a highly ionized gas can also hinder condensation, we estimated the degree of ionization that might exist in the carbon vapor at the time of maximum expansion. The first thing to note is that the collision diameter of the electron-ion interaction—a Coulomb interaction—is about 10 times that of the chemically reacting species. Furthermore, the collision diameter varies inversely with T ; the collision diameter grows as the temperature drops. Since the three-body collision frequency is proportional to σ^5 , the ion-electron recombination rate will initially be 10^5 times faster than chemical recombination and will grow as the temperature drops. That gives us the chance that the recombination will proceed rapidly even if the initial ionization is significant. It also suggests that an equilibrium calculation of ionization is reasonable and conservative. In a manner analogous to that used for molecular dissociation, we found that the ionization fraction α could be found from

$$\frac{\alpha^2}{1 - \alpha} = 2.4 \times 10^4 T^{3/2} e^{-11.2/T} \quad (15)$$

at the time of maximum expansion. For $T = 0.5$ eV, $\alpha = 2 \times 10^{-3}$, and we conclude that ionization can be ignored during condensation.

Condensation Coefficients. The condensation coefficient β is defined as the probability of a gas molecule sticking to a solid surface with which it collides.

Ordinarily, one would think the condensation coefficient to be unity for a cold surface, unless the energy of the incident molecule is greater than about twice the latent heat of sublimation (10 to 15 eV per molecule). Carbon is an exception. The reflection coefficient $(1 - \beta)$ of C_1 and C_3 vapor have been measured directly by Chupka et al.²⁰ The values of the inferred condensation coefficients for two different temperatures were

$$\left. \begin{aligned} \beta_1 &= 0.4 \pm 0.2, \\ \beta_3 &= 0.1 \pm 0.2, \end{aligned} \right\} \text{ at } T = 2300 \text{ K}$$

$$\left. \begin{aligned} \beta_1 &= 0.70 \pm 0.15, \\ \beta_3 &= -0.1 \pm 0.2, \end{aligned} \right\} \text{ at } T = 800 \text{ K} \quad (16)$$

where T is the solid-surface temperature. In the case of the carbon blanket the inside surface could be at 3925 K, but the bulk of the 1-cm-thick surface layer would have a granule temperature of 1600 K. At this temperature we estimate the condensation coefficient for C_1 to be about 0.5. This value would double or triple the previously calculated condensation time to about 12 to 18 ms. As can be seen from the above data, the condensation coefficient for C_3 is quite small and very uncertain. It is therefore fortunate that the above estimate of dissociation indicates that there will be no C_3 .

Reactions of Atomic Carbon Vapor with BeO. Hot carbon vapor can react rapidly with solid BeO to produce gaseous Be and CO, both incondensable in the Cascade chamber. It is thus very important that the vaporized carbon recondense before it completely penetrates the carbon layer and reaches the BeO. We calculated the rates of gas-producing reactions as a function of the rate of penetration of C_1 through the carbon layer. From these rates we derived the chamber pumping rates necessary to remove these incondensable gases as a function of the same variable, as shown in Fig 8-5. Since the inherent Cascade pumping capability (through the chamber ends) is 4×10^5 L/s, Fig. 8-5 indicates that we should be able to tolerate ~ 0.1 g/s of carbon penetrating the surface layer. At a pulse rate of 5 Hz, we can allow up to 0.02 g to penetrate on each pulse. Since about 30 to 460 g are vaporized with

each pulse, the penetration factor must be less than 4.3×10^{-5} . For a condensation coefficient β , the probability of an atom not condensing after s collisions is $(1 - \beta)^s$. For a β of 0.5, $s = 15$ collisions are necessary. From Eq. (4) recall that the mean free path for a C_1 molecule in a bed of porosity 0.5 is $4d/3$, where d is the granule diameter. The mean free path in the chamber's radial direction is $4d/3^{3/2}$. Thus, the thickness necessary to provide 15 collisions is

$$\Delta r = 15 \frac{4d}{3^{3/2}} = 12 \text{ mm.} \quad (17)$$

Even if the condensation coefficient were only 0.3, only 28 mm would be necessary to provide a sufficient barrier. The pumping rate necessary to handle the gaseous products of the interactions between carbon and BeO can thus be made negligible.

Conclusions. We examined the condensation times for surface layers made of three granular materials for the blanket in the Cascade reaction chamber. The high likelihood that BeO and SiC will dissociate and significantly inhibit condensation means that they would not be suitable as first layers. It appears that a carbon-granule layer 1- to 3-cm thick will suffice and result in a condensation time of less than 20 ms for a laser driver. If a heavy-ion driver requires 2 density 10^3 times lower to operate, then the condensation time would be about 7 times longer, or about 140 ms. It appears that either driver could operate at 5 Hz; the laser driver might be able to operate at a pulse rate of a few tens of hertz.

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Strain Energy Effects on Dynamic Fragmentation of Blanket Granules

We are studying whether, and under what conditions, the Cascade blanket would fragment and what would be the average fragment size and distribution. In this article we discuss how the blanket might fragment as a result of neutron

Fig. 8-5. Required chamber pumping speed to remove the gaseous products of reactions between carbon vapor and solid BeO. The chamber pumping capacity is about 10^5 L/s.

heating; our study builds on our previous work on liquid fragmentation.²¹ Fragmentation can also happen if the inner radius of the blanket absorbs energy from x rays or the debris of fusion-fuel pellets. We have left this latter case to future study.

Typically, neutrons penetrate the blanket quickly in comparison with the time it takes for release waves to move any significant distance from available free surfaces. The blanket material is therefore heated at constant volume, and high internal pressures may be generated that, when released, could fracture and fragment the blanket components. In some designs this disassembly is desirable: In the HYLIFE system, for example, disassembly of the liquid-lithium jets significantly speeds condensation of the lithium plasma.²²⁻²⁴ In other designs, however, blanket fragmentation is deleterious. The Cascade blanket consists of granular layers of carbon and ceramic oxide granules, and substantial comminution would necessitate a refabrication facility, significantly raising the plant's total cost.²⁵

Model. Grady^{26,27} derived a simple model of dynamic fragmentation in terms of the total free energy of the resulting fragments in an equilibrium configuration. Grady assumed the free energy to be the sum of fragment surface energy (which is created during the formation of new fragment surface area) and a local kinetic energy (which measures the intensity of expansion of the body and is responsible for driving the fragmentation process). A crucial point was the recognition that the total kinetic energy of the body is not

available for fragmentation. Rather, only that fraction relative to the center of mass will participate in the breakage, while the remainder stays in the rigid-body flight of fragments.

Although the model agrees qualitatively with results from experiments on such diverse materials as oil shale and steel, it neglects the effect of stored elastic (strain) energy. In addition, the mean fragment size is obtained by minimizing the free-energy density E with respect to the acquired surface area per unit volume A . The surface-energy density associated with the equilibrium fragment $\Gamma_e = \gamma A_e$, where γ is the surface energy per unit area, is then exactly twice the kinetic-energy density that was supposedly available for fragmentation.

A modest reformulation corrects this inconsistency and shows the effect of the potential energy as well. We assume, with Grady, that the body under load undergoes homogeneous dilatation, that the process is adiabatic, and that compressibility effects can be ignored. We add the condition that no fragmentation occurs until a critical stress σ_c is attained. The total available kinetic energy for driving the breakage is $\Delta T = 2\pi N \rho \epsilon^2 a^3/5$, where N is the overall number of fragments, ρ is the mass density, ϵ is the strain rate ($= \rho/3\rho$), and a is the fragment radius. If the initial radius of the dilating body is R , then $N = (R/a)^3$, so that $\Delta T = T_i(a/R)^2$; T_i is the initial kinetic energy resident in the body.

Since the body is expanding uniformly, the strain energy of distortion is zero, and only the hydrostatic strain energy is involved. The fragments are presumably stress-free when formed, so that all the stored elastic energy is available for fragmentation, i.e., $\Delta P = P_i = (4/3)\pi R^3(\sigma_c^2/2K)$, where K is the bulk modulus.

Finally, the surface energy associated with the fragments is simply $4\pi N a^2 \gamma = \Gamma_i R/a$, so that the incremental surface energy acquired is $-\Delta\Gamma = \Gamma_i(1 - R/a)$. Then, balancing the change in system energy that results from fragmentation, we have

$$\begin{aligned} \Delta T + \Delta P + \Delta\Gamma &= T_i(a/R)^2 \\ &+ P_i + \Gamma_i(1 - R/a) = 0. \end{aligned} \quad (18)$$

Equation (18) can be written in terms of the fragment radius:

$$a^3 + \alpha a - 2\beta = 0, \quad (19)$$

where

$$\alpha = \frac{2\beta}{R} + \frac{5}{3} \left(\frac{\sigma_c}{\rho c \epsilon} \right)^2,$$

$$\beta = \frac{5}{2} \left(\frac{K_i}{\rho c \epsilon} \right)^2,$$

$K = \rho c^2$, and the surface energy is expressed in terms of the fracture toughness (i.e., $\gamma = K_{Ic}/2\rho c^2$). When the surface energy of the unfragmented body can be ignored and the stress required to initiate breakup is insignificant (as would be the case, for example, with a fluid body), $\alpha \rightarrow 0$, and the solution to Eq. (19) is simply

$$\alpha = (2\beta)^{1/3}, \quad (20)$$

which is $2^{1/3}$ times the value obtained by Grady.²⁶ The discrepancy is exactly related to the kinetic energy that is presumed available for disassembly.

The general solution to Eq. (19) is:

$$a = 2 \left(\frac{\alpha}{3} \right)^{1/2} \sinh \left(\frac{\phi}{3} \right), \quad (21)$$

where $\phi = \sinh^{-1} \beta(3/\alpha)^{3/2}$.

When the parameter $\beta(3/\alpha)^{3/2} \gg 1$,

$$a \simeq (2\beta)^{1/3}, \quad (22)$$

as in Eq. (20). On the other hand, for a truly brittle material, with low fracture toughness and a high fracture-initiation stress, $\beta(3/\alpha)^{3/2} \ll 1$ and

$$a \simeq 2\beta/\alpha. \quad (23)$$

In this limit, the final fragment size should be essentially independent of the strain rate and proportional to the square of the ratio of the fracture toughness to the initiation stress.

We compared these results with the FS-01 explosive fragmentation experiments of Weimer and Rogers²⁸ cited in Ref. 26. In these experiments, $\epsilon \simeq 4 \times 10^4 \text{ s}^{-1}$, $c \simeq 5 \text{ km/s}$, and $\rho = 7840 \text{ kg/m}^3$. We

assume that the data on yield strength with 0.2% offset provide reasonable estimates of σ_y and that R is just twice the ratio of the specimen volume to surface area ($R = 12.8 \text{ mm}$). Table 8-7 lists the experiment variables, as well as the measured Mott size parameter μ (half the average fragment mass); we note that the Mott size data appear to be incorrectly plotted in Ref. 26. In Fig. 8-6 we plot the results of evaluating Eq. (21) with these data, as well as the asymptotic solutions represented by Eqs. (22) and (23). The measured fragment size, based on the Mott parameter, agrees well with Eq. (22) over the entire range $0.26 < \beta(3/\alpha)^{3/2} < 5.8$. If these data are representative, the implication is that the stored elastic energy is essentially uncoupled from the fragmentation process, at least under conditions in which the local kinetic energy is sufficient to produce fragment sizes small compared with the unfragmented body. Moreover, this seems to be the case even when the kinetic energy density is small compared with the (potential) strain energy density.

Summary. It is unclear whether carbon (graphite) and ceramics are well represented by the brittle-steel data. The above result is especially puzzling in light of the well-known Rayleigh-wave limit on crack(s) propagation speed; i.e., the time for body waves to propagate (and consequent release of strain energy) should be considerably less than the time for cracks to traverse the specimen. Moreover, for quasi-static loading (which more nearly characterizes the rear of the blanket), it is certain that fragmentation is controlled entirely by the balance between potential and surface energy, in accordance with the theory originally advanced by Griffith.²⁹ Work is in progress to resolve this apparent paradox, as well as to determine the size distribution of any fragments that might develop. Equation (22) is probably a lower bound on the average fragment size; applied to the Cascade blanket energy deposition, it predicts no fragmentation except very near the inner surface where x-ray heating is encountered. We will examine the effects of the strong temperature gradient expected in this region, as well as the consequences of thermal ratcheting (caused by reactor cycling),

Steel type	Critical stress (MPa)	Fracture toughness (MPa·m ^{1/2})	Mott size parameter μ (mg)	$\beta(3/\alpha)^{3/2}$
FS-01(A)	2570	21.7	19.0	0.256
FS-01(B)	2036	24.2	24.1	0.635
FS-01(C)	1777	31.4	34.1	1.497
FS-01(D)	1520	57.3	97.4	5.750

Table 8-7. Data from explosive fragmentation experiments of Welmer and Rogers (see Ref. 28).

Fig. 8-6. Fragment radius as a function of fracture toughness. The curves are predicted by Eqs. (21)–(23); data points are from explosive fragmentation experiments (see Ref. 8).

which may significantly lower the fracture toughness.

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Activation Analysis

We calculated neutron activation for the reference design of Cascade. The results, which are reported in detail in Ref. 30 and summarized in this article, indicate that the activity in Cascade is significantly lower than in magnetic fusion reactor designs.

Neutronics Model. For these calculations, Cascade was modeled in spherical geometry. A 14.1-MeV neutron source was uniformly distributed in a spherical region of compressed DT that had a density-radius product of 3 g/cm^2 . The inner surface layer of the blanket was represented by a 1-cm-thick region of carbon with an inner radius of 3.34 m and a packing

fraction of 50%. The BeO multiplier was a 10-cm-thick region (3.35 to 3.45 m from the center) at 40% of the normal material density. A 90-cm-thick region of LiAlO_2 extended from 3.45 to 4.35 m (measured from the center) and was at 50% normal density. Thus, the total thickness of the blanket was about 1 m, while the effective material thickness was about 0.5 m because of the porous nature of the flowing granule material. A 2-cm-thick SiC chamber wall, extending 4.35 to 4.37 m from the center, surrounded the blanket.

The entire chamber was surrounded by a borated-water radiation shield with an inner radius of 7.48 m. It consisted of a 2-cm-thick aluminum liner, a 1-m-thick water region, and a 1-m-thick borated-water region. Both water regions contained 10 vol% aluminum as the structural material. Boron was excluded from the inner half of the shield to make it a more effective neutron reflector for the chamber.

Impurity elements, which in some cases can dominate neutron activation, were included in the calculation. The impurity levels used in our calculation were based on the values given in Refs. 31–33 and are listed in Ref. 30.

Activation. We calculated one-group neutron-activation cross sections for each region with the ORLIB code,³⁴ which makes use of the energy-dependent fluxes from TART³⁵ and the ACTL cross-section library.³⁶ The resulting activation cross sections were then used in the FORIG code³⁷ to determine the time-dependent activity, afterheat, biological hazard potential (bhp), and shallow burial index (SBI) for each region. The irradiation time was taken as the assumed 30-y life of the plant.

The activity (by region and the total) as a function of time after shutdown is shown in Fig. 8-7. The total activity 1 s after shutdown—1 GCi—falls by an order of magnitude during the first day and by a factor of

Fig. 8-7. Activity as a function of time after shutdown.

Region	Time after shutdown					
	0	1 d	1 month	1 y	10 y	100 y
C surface layer	4.63×10^5 ¹² B	3.69×10^3 ³⁷ Ar	2.03×10^3 ³⁷ Ar	4.19×10^2 ⁵⁵ Fe	4.88×10 ⁵⁵ Fe	1.51×10 ¹⁴ C
BeO multiplier	3.08×10^4 ⁶ He	3.79×10^5 ²⁴ Na	2.44×10^5 ⁵⁵ Fe	1.43×10^5 ⁵⁵ Fe	2.58×10^4 ⁶⁰ Co	5.33×10^3 ¹⁴ C
LiAlO ₂ breeder	7.31×10^8 ²⁴ Na	8.64×10^7 ²⁴ Na	4.05×10^6 ⁶⁰ Co	1.87×10^6 ⁶⁰ Co	4.69×10^5 ⁶⁰ Co	5.12×10^4 ¹⁴ C
SiC wall	3.11×10^6 ²⁸ Al	4.35×10^3 ⁶⁰ Co	2.64×10^3 ⁶⁰ Co	1.44×10^3 ⁶⁰ Co	4.13×10^2 ⁶⁰ Co	2.54 ¹⁴ C
Al liner	1.03×10^8 ²⁸ Al	1.27×10^6 ⁷⁶ As	1.50×10^5 ⁵⁵ Fe	7.18×10^3 ⁵⁵ Fe	6.67×10^3 ⁵⁵ Fe	4.98×10 ⁶³ Ni
H ₂ O shield	6.76×10^7 ²⁸ Al	6.98×10^5 ⁷⁶ As	1.08×10^5 ⁵⁵ Fe	4.82×10^4 ⁵⁵ Fe	4.82×10^3 ⁵⁵ Fe	4.84×10^2 ¹⁴ C

Table 8-8. Activity (Ci) at selected times after shutdown and predominant isotope.

200 during the first month after shutdown. Thirty years after shutdown the activity is reduced by four orders of magnitude to 0.1 MCi.

As seen in Fig. 8-7, the total activity is strongly dominated by the activity in the LiAlO₂ breeder region. For the first day, ²⁴Na, which is produced by (n,α) reactions with ²⁷Al, dominates the activity. Sodium-24 beta-decays with a half-life of 15 h and emits hard gamma rays with energies of 2.75 and 1.37 MeV. From several days to about 20 y ⁶⁰Co, which is produced by neutron capture in the ⁵⁹Co impurities, is the primary contributor to the activity in the LiAlO₂. Cobalt-60 beta-decays with a half-life of 5.3 y and emits gamma rays at energies of 1.17 and 1.33 MeV. Beyond 20 y the activity in the LiAlO₂ is dominated by ³⁸Ar, which is primarily produced by (n,p) reactions with ³⁹K impurities. Argon-39 beta-decays with a half-life of 269 y and does not emit any gamma radiation. Table 8-8 lists dominant contributors and the activities at progressive times after shutdown for the various regions.

The activity levels in Cascade are significantly less than in other recent fusion reactor studies. The activities in the Mars³⁸ and Starfire³⁹ magnetic fusion reactor designs were compared with Cascade on the basis of curies per unit net electric power; the net electric power of Cascade is 815 MW_e, while Mars and Starfire both have net powers of 1200 MW_e. The activity in Cascade is lower by a factor of 3.5 to 4.8 at shutdown, a factor of 190 to 280 at 1 y after shutdown, and a factor of 1.2 to 3.1 at 100 y after shutdown. (These results do not include the activity of first walls or blanket materials that may require replacement during the life of the power plants.)

Fig. 8-8. Decay heat as a function of time after shutdown.

Neutronics and activation calculations for an earlier version of the Cascade chamber, which used a Li₂O breeder, were reported by Blink.⁴⁰ For this version, activity at shutdown with the Li₂O breeder was about an order of magnitude less than those we calculated for the LiAlO₂ case.

Decay Heat. Figure 8-8 shows the radioactive decay heat as a function of time after shutdown. The decay heat at shutdown is 21.8 MW, and it drops to 6 MW within an hour after shutdown. During this period decay heat is dominated by ²⁴Na in the LiAlO₂. The SiC chamber wall, which operates at 1200 K, can radiate ~30 MW to the surrounding water shield, which provides a large thermal sink. Thus we expect that the decay heat can be removed by thermal radiation without the need for active cooling systems. Even with absolutely no heat

removal from the LiAlO_2 , it would take about one week to reach its melting point.

Biological Hazard Potential. The inhalation bhp is the volume of air required to dilute vaporized material to breathing standards. Table 8-9 gives the bhp for the various regions in the Cascade chamber. If we assume a plume containing vaporized radioactive material that has a volume tens of cubic kilometers, then the allowable release fraction can be determined from Table 8-9. At one day after shutdown a release fraction of $\sim 10^{-4}$ of the BeO would be tolerable, while the release of LiAlO_2 would have to be limited to $\sim 10^{-6}$. These fractions are somewhat more restrictive than the Li_2O version of Cascade.⁴⁰ The probability of release is expected to be low with the use of high-temperature ceramic materials and the absence of volatile materials within the chamber.

Shallow Burial Index. The SBI is the ratio of the specific activity (Ci/m^3) to the legal limit; thus, a value below 1.0 allows legal shallow land burial. The activity for each region at 10 y after shutdown was divided by the volume of material corresponding to that region and then divided by the legal limit to give the SBI. The volume used in this calculation was the full-density volume of the total inventory of material. The results for each region are given in Table 8-10. The first column is based on guidelines given in 10CFR61.⁴¹ There are, however, several long-lived

isotopes produced in fusion reactors that are not currently limited. Maninger and Dorn have estimated limits for some of these isotopes.³³ The resulting SBI based on the full set of limits is given in column 2 of Table 8-10. Note that Refs. 4 and 12 do not give a limit for ^{39}Ar . Since ^{39}Ar and ^{14}C both beta-decay without emitting any gamma radiation, Chen¹³ suggested that the limit for ^{39}Ar be set equal to that for ^{14}C . The limit for ^{14}C is 8 Ci/m^3 if it is not in a metal and 80 Ci/m^3 if it is in a metal. The 8-Ci/m^3 limit is used in column 2 of Table 8-10.

As indicated, BeO exceeds the limit by a factor of nearly 24, virtually entirely because of the ^{14}C . About 60% of the ^{14}C is produced by (n,α) reactions with ^{17}O , while the remainder is produced by (n,p) reactions with ^{14}N impurities. The SBI of 90.6 for LiAlO_2 is dominated by ^{39}Ar (51.8) and ^{26}Al (32.7). Argon-39 is produced by (n,p) reactions with ^{39}K impurities, while ^{26}Al is produced from $(n,2n)$ reactions with ^{27}Al . Reducing the K impurity level in the LiAlO_2 and using a thicker BeO region would help to reduce the SBI for the LiAlO_2 . Even so, it appears that significant dilution will be required in order for the BeO and LiAlO_2 to qualify for shallow land burial. All other materials qualify for shallow burial without dilution.

Summary. Neutron-activation calculations have been carried out for the Cascade reference design. The total activity in the Cascade reactor at shutdown is less than in the magnetic fusion reactor designs Starfire and Mars—by factors of 3.5 and 4.8, respectively. Sodium-24, which is produced from $^{27}\text{Al}(n,\alpha)$ reactions in the LiAlO_2 , dominates the activity, decay heat, and biological hazard potential for the first couple of days after shutdown. With this high-temperature chamber design, the decay heat can be removed by thermal radiation to the surrounding water shield to prevent melting of blanket and chamber materials. Shallow land burial will be possible only if the BeO and LiAlO_2 are diluted by factors of 24 and 91, respectively. Carbon-14 is the limiting isotope in the BeO , while ^{39}Ar and ^{26}Al limit the burial of LiAlO_2 .

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Table 8-9. Inhalation biological hazard potential (km^3 of air) and predominant isotope.

Region	Time after shutdown	
	1 h	1 d
C surface layer	$7.8 \times 10^2 (^{22}\text{Na})$	$6.1 \times 10^2 (^{22}\text{Na})$
BeO multiplier	$3.0 \times 10^5 (^{60}\text{Co})$	$2.5 \times 10^5 (^{60}\text{Co})$
LiAlO_2 breeder	$5.8 \times 10^7 (^{24}\text{Na})$	$2.7 \times 10^7 (^{24}\text{Na})$
SiC wall	$1.7 \times 10^4 (^{31}\text{Si})$	$6.9 \times 10^3 (^{60}\text{Co})$
Al liner	$6.6 \times 10^5 (^{76}\text{As})$	$2.9 \times 10^5 (^{76}\text{As})$
H_2O shield	$3.4 \times 10^5 (^{76}\text{As})$	$1.4 \times 10^5 (^{76}\text{As})$
Total	5.9×10^7	2.8×10^7

Table 8-10. Shallow burial index and predominant isotopes for Cascade reactor components.

Region	SBI based on limits given in:	
	10CFR61 (Ref. 41)	Maninger and Dorn ³³
C surface layer	0.07 (^{14}C)	0.07 (^{14}C)
BeO multiplier	23.61 (^{14}C)	23.96 (^{14}C)
LiAlO_2 breeder	3.23 (^{14}C)	90.6 ^a (^{39}Ar)
SiC wall	0.06 (^{14}C)	0.06 (^{14}C)
Al liner	0.01 (^{63}Ni)	0.73 (^{26}Al)
H_2O shield	0.07 (^{14}C)	0.07 (^{14}C)

^a Assumes limit of 8 Ci/m^3 for ^{39}Ar .

Heat Exchanger Design

The Cascade power-conversion system uses a ceramic-tube heat exchanger to transfer fusion energy from the high-temperature granules to the helium working fluid in a closed-cycle gas turbine. This year, we refined the design of the heat exchanger to include a separate heat-exchange region for each blanket zone, which makes for high efficiency. Other design choices were made with an eye to reliability, ease of manufacture, and cost effectiveness.

During initial studies,² we selected a shell-and-tube concept that operates on the basis of vacuum heat transfer, with granules flow-mixing by gravity in the primary (shell) side of the heat exchanger. We selected a horizontal array of staggered coolant tubes configured in triangular pitch to enhance granule mixing. The once-through coolant tubes contain 5-MPa (secondary-side) helium gas and pass back and forth, ascending countercurrent to the downward granule flow.

In a later study,³ we surveyed heat-exchange parameters to determine ways to reduce costs. These parameters included primary-side gas pressure and temperature; granule size and temperature; granule/tube contact time; tube pitch, diameter, and length; heat-exchanger width; and secondary-side Reynolds number. Our design objectives were to minimize these parameters: heat transfer area, pressure drop and pumping power, the number of heat-exchanger tubes, and the active heat-exchange height. Our motivation was as follows:

- A decrease in heat-transfer area, pressure drop, or pumping power translates directly into reduced capital cost.
- Generally speaking, reliability decreases as the number of tubes increases.
- The active heat-exchange height dictates the minimum capability requirements of the granule throw system.

We developed an attractive mechanical design for a heat exchanger and adapted its design parameters to each blanket zone. Each blanket zone has a separate heat-exchanger region; the helium flows through each region in series. Figure 8-9 shows the temperature of the blanket granules and the helium as a function of

Fig. 8-9. Temperature profiles in the Cascade heat exchanger. The solid curves represent the granule side of each heat exchanger region. The dashed line represents the helium side of the heat exchanger.

normalized distance through the heat exchanger. Fifty-nine percent of the total energy is transferred to the helium in the tritium breeder zone, which has the lowest temperature. An additional 13% is transferred in the front zone. Finally, 27% is transferred in the surface layer, which has the highest temperature. The temperature discontinuities between regions on the granule side of the heat exchanger are caused by the difference in average temperature of the granules in each blanket zone as they leave the reactor. The helium behaves as a perfect gas, and its temperature rise through the heat exchanger is linear. Because the average temperature difference between the granules and the helium in each region is nearly constant, heat-transfer efficiency is improved where granule temperatures are kept below maximum.

Mechanical Arrangement and Thermal Performance. At the operating temperature of the blanket (up to 1600 K), the only functional materials for the heat-exchange surface are ceramics, such as SiC or Si₃N₄. We chose to incorporate SiC tubes in the structural layout shown in Fig. 8-10. The solid granules composing the blanket flow across the ceramic-tube array in a vacuum (1.3 Pa), cross-countercurrent to the multi-pass helium in the tubes. The tubes are single-walled, and heat is transferred directly from the primary-side granules to the 5-MPa secondary-side helium. Thus, though the tubes are in hoop tension because of the internal helium pressure, the tensile stresses are modest (~30 MPa in 2.5-cm-o.d. tubes with 2-mm-thick walls) and can be accommodated while allowing adequate reserve for thermal stresses.

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Proof-testing of ceramic components has advanced significantly in recent years,^{42,43} and the small size and low stresses called for in this design make this a suitable application for SiC tubes.

The design of the heat exchanger takes into account that the ceramic tubes must

eventually interface with metal structure. Thus, the high-pressure headers shown in Fig. 8-10 are coaxial, insulated, and actively cooled metal, and incorporate a practical metal-to-joint. A low-stress, ceramic, spacer wall separates the granules from the metallic wall and protects the

Fig. 8-10. Schematic of a portion of the Cascade heat exchanger.

structural members from abrasion and heat. This configuration does not include extremely complex or highly stressed ceramic components, making it easier to manufacture. The heat exchanger is also made more reliable by a modular design, which makes it possible to change tubes—or, more practically, a bank of tubes.

Table 8-11 lists mechanical and thermal design characteristics of the heat exchanger, which has a total heat-transfer area of $13\,200\text{ m}^2$. Total pressure in the granule side of the heat exchanger is 1.3 Pa—considerably lower than inside the reactor, where the pressure before each fusion energy pulse is 13 Pa if a laser driver is used. The lower pressure in the heat exchanger eases the removal of tritium bred inside the LiAlO_2 granules and reduces the total tritium inventory to 264 g. Of this inventory, 140 g is

characteristically released in less than 30 d. We considered two materials (carbon and beryllium oxide) for the surface layer because two designs were considered for the Cascade blanket.

Details of the thermal calculations are included in Ref. 3. The heat-transfer area depends largely on the granule/tube heat-transfer coefficient. Parametric analysis of primary-side gas pressure, granule size, and temperature demonstrated that temperature is the controlling variable. The main effect of temperature is through the fourth-power dependence of the radiative heat-transfer term in the granule/tube heat-transfer coefficient. Increasing temperature decreases the required area, with attendant reductions in the total number of tubes, pressure, pumping power, and heat-exchanger height. The secondary-side parameters (tube diameter and

Table 8-11. Mechanical and thermal characteristics of the Cascade heat exchanger.

Material Power level (MW _t)	Surface layer		Front zone	Breeder zone
	(C) 460	(BeO) 480	(BeO) 220	(LiAlO ₂) 990
Primary side				
Gas pressure (Pa)	1.3	1.3	1.3	1.3
Average gas temperature (K)	1550	2200	1460	1320
Granule diameter (mm)	1	1	1	1
Granule inlet temperature (K)	1600	2300	1510	1355
Granule outlet temperature (K)	1380	2100	1410	1285
Bed porosity	0.5	0.5	0.4	0.5
Granule flow rate (m ³ /s)	1.0	1.0	1.0	7.5
Overall granule/tube heat-transfer coefficient (W/m ² ·K)	790	2200	660	490
Log-mean temperature difference (K)	216	580	220	220
Secondary side				
Helium pressure (MPa)	5.0	5.0	5.0	5.0
Helium inlet temperature (K)	1200	1200	1155	945
Helium outlet temperature (K)	1300	1300	1200	1155
Helium velocity (m/s)	90	150	82	69
Total helium mass flow rate (kg/s)	910	940	920	910
Reynolds number	1.5×10^5	2.5×10^5	1.5×10^5	7.0×10^4
Fractional pressure drop	0.0045	0.0029	0.0022	0.0088
Pumping power fraction	0.023	0.015	0.022	0.018
Heat-transfer coefficient (W/m ² ·K)	2600	2600	2500	3000
Log-mean temperature difference (film drop) (K)	70	320	60	40
Overall heat-exchanger dimensions				
Heat exchange area (m ²)	2600	380	1600	9000
Active height (m)	5.1	4.3	5.5	5.7
Active width (m)	15	15	15	20
Active tube length (m)	1.0	0.25	0.5	1.5
Total number of tubes	1.8×10^4	9100	1.9×10^4	8.2×10^4
Number of working fluid passes	6	6	6	5
Tube pitch (cm)	7.0	9.0	7.0	4.0
Tube outside diameter (cm)	5.0	5.0	5.0	2.5
Tube inside diameter (cm)	4.6	4.6	4.6	2.1
Tube wall temperature difference (K)	7	50	6	4
Granule flow area (m ²)	4.3	1.7	2.1	11

Reynolds number) have relatively minor roles in establishing the heat-transfer area.

For a given heat-transfer area, the heat exchanger configuration—i.e., height, width, and length—is determined by the constraints imposed by the granule transport system, tube manufacturing capability, and minimum number of working-fluid passes. The heat exchanger cannot be wider than 20 m without exceeding the capability for lateral granule throw with the scoop system; in addition, it must be between 5 and 10 m high. A minimum of five passes of the working fluid is necessary to satisfy the configuration of the cross-counterflow coolant. These constraints are not severe and allow much flexibility for future design iterations.

Given the heat-transfer area and heat-exchanger configuration, the tube pitch and diameter are determined by the need to maintain adequate granule flow area and to minimize heat-exchanger height, number of tubes, pressure drop, and pumping power. The tube pitch, which affects only heat-exchanger height, was selected to keep heat exchangers near the minimum acceptable height, i.e., 5 to 6 m. The tube-gap spacing, determined from granule flow considerations, was chosen by approximating our rectangular flow area condition with the equation for mass flow

through a pipe (see Ref. 44):

$$Q = 1.055 D^{2.5}, \quad (24)$$

where Q is the volume flow rate (m^3/s) and D is the pipe diameter (m). For the three heat-exchanger-volume flow rates, we used Eq. (24) to calculate the necessary flow areas. We then calculated the equivalent hydraulic diameter (minimum tube spacing) for the selected tube lengths. The minimum gaps were thus calculated to be 1.5 cm for 1.5-m tubes in the LiAlO_2 heat exchanger and 2.0 cm for 0.5- and 1.0-m tubes in the BeO and C heat exchangers. The minimum gap would be 4.0 cm for 0.25-m tubes in the surface-layer heat exchanger. For a given pitch, we can minimize heat-exchanger height, number of tubes, pressure drop, and pumping power by maximizing the tube diameter. After maximizing the pitch from height considerations as described above, we obtained the maximum tube diameter by selecting the minimum acceptable gap.

At this point, the heat exchanger is completely designed except for the secondary-side Reynolds numbers. This variable strongly affects the pressure drop and pumping power. The Reynolds numbers shown in Table 8-11 were obtained by restricting the pressure drop and pumping power fractions to between 2 and 3%. These levels, which are acceptable by current thermal-hydraulics design practices, were used to calculate the efficiency of the power-conversion system.

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Power-Conversion System

The high blanket temperatures possible with Cascade mean that thermal conversion can be highly efficient: a net efficiency of 49% is possible with the once-through, regenerative, no-reheat Brayton cycle.

We evaluated possible power-conversion systems for Cascade mainly on the basis of net thermal efficiency, which is a function of the peak working-fluid temperature. We

Fig. 8-11. Power-conversion efficiency for various thermodynamic cycles as a function of peak working-fluid temperature. Ceramic heat-exchanger tubes must be used above 1200 K because of materials limitations for metallic tubes.

evaluated three primary power-cycle options: Rankine, with water as the working fluid; Brayton, with helium (an ideal gas) as the working fluid; and a Rankine/Brayton hybrid called the Field cycle, also with water as the working fluid. We also considered secondary power-cycle options, including combinations of these in the form of topping and bottoming cycles, different working fluids, non-ideal gas Brayton cycles such as supercritical CO_2 , and dissociating gas cycles. The results of this evaluation are shown in Fig. 8-11; details can be found in Refs. 3 and 45. [The steam-injected gas turbine (open Brayton cycle) is shown in Fig. 8-11 for comparison.] Below 900 K, the standard steam Rankine cycle used in essentially all fossil-fired and nuclear power plants gives the highest efficiency. Between 900 and 1200 K, the steam Field cycle appears promising. Above about 1200 K, the helium gas turbine (Brayton) cycle is most efficient.

Following this evaluation, we chose the once-through, regenerative, no-reheat Brayton cycle (see Fig. 8-12) for Cascade because of its simplicity, cost effectiveness, and high efficiency. The cycle includes three compressors with intercoolers, a regenerator, heat addition in the heat exchanger, and a turbine-generator to extract energy and produce electrical power. The regenerator transfers heat from the helium gas exiting the turbines to the compressed gas before it enters the heat exchanger as shown by the dashed arrows in Fig. 8-12. The maximum pressure ratio is 2.3. No intermediate loop is needed to isolate the

tritium present in the granules from the helium gas, because diffusion of tritium through the heat-exchanger tube walls is only 25 Ci/d. Excluding the intermediate loop improves thermal efficiency and eliminates costly equipment.

A power balance is shown in Fig. 8-13. Fusion power, which is initially 1500 MW, rises to 1670 MW by neutron capture in the blanket. About 25 MW are subsequently lost in the heat exchanger, but 1645 MW of thermal power are converted to 905 MW of gross electrical power in the turbine generator. Ten percent of the gross electric power is recycled for the laser (75 MW) and auxiliary systems (15 MW), resulting in a net electrical power of 815 MW.

The heat exchanger tubes must be ceramic because of the high granule and helium temperatures; however, enough expansion occurs in the turbine nozzles to lower the temperature of the helium

Fig. 8-12. Pressure-temperature plot of the Cascade power-conversion system.

Fig. 8-13. Cascade power balance.

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impinging on the turbine blades. This temperature allows the use of metallic turbine blades, which are less expensive and more reliable. Cascade's power-conversion efficiency is 55%, which results in a net plant efficiency of 49%. In comparison, a coal-fired power plant that can deliver the same net electrical power achieves a net plant efficiency of about 35%.

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Balance of Plant

The "balance of plant" includes all power-plant buildings and components other than the reactor, heat exchanger, and power-conversion system. In this article we describe the balance-of-plant buildings and components, explain the analyses that led to the choices of shielding materials and designs, and address safety issues. Details of these analyses and safety issues can be found in Ref. 3.

The building arrangement is shown in Fig. 8-14. We determined the building dimensions by considering the form and function of comparable buildings in other reactor studies. Structures for nuclear power plants are normally divided into

two categories: Seismic Category I and non-Seismic Category I. Category I structures house all safety-related equipment and are designed to withstand the Design Basis Safe Shutdown Earthquake, which is defined as having a peak ground acceleration of 0.25 g. These structures generally consist of exterior and interior walls, floor and roof slabs, and foundation mats—all of reinforced concrete. The floor and roof slabs are normally supported on heavy structural steel framing.

The majority of buildings at the Cascade site are designated non-Category I because they house equipment and components not required for plant safety. These are generally steel-frame structures with insulated metal sidings and insulated roofs. The buildings are designed for normal dead and live loads, including expected loads from natural events and additional loads for seismic conditions as provided in the Uniform Building Code. Thus, they are similar to the buildings of coal-fired power plants. Where required by proximity, these buildings are designed so that their failure would not lead to loss of function in a Category I structure. The specifications for the Cascade plant buildings are consistent with the inherent safety resulting from the low levels of radioactivity within the site.

Fig. 8-14. Cascade power plant building arrangement. Dimensions shown are in meters.

The reactor building houses the vacuum chamber, within which the reaction chamber, heat exchanger, and shielding are located. The arrangement of the former two is shown in the isometric perspective of Fig. 8-15. The integrated arrangement within the vacuum chamber is shown in Fig. 8-16. Other equipment within the reactor building includes transportation, lift, and component-handling fixtures and equipment; the chamber drive system; the blanket-granule scoop-drive system; the

shield system; the final turning and focusing mirrors; and the target injection system. Also included are portions of the secondary helium loop, shield cooling, and recirculating-gas cooling systems. The building is metal, $40 \times 30 \times 50$ m high, supported on concrete footings, and framed with structural steel. The ground floor is a concrete slab and the upper floors are grating.

The turbine building contains the turbine generator, helium compressors,

Fig. 8-15. Major component positions within the vacuum chamber.

Fig. 8-16. Vacuum chamber arrangement: (a) side view and (b) end view.

regenerative heat exchanger, intercoolers, and the support equipment necessary to generate electricity. This three-level, non-Category-I building is estimated, because of the size of the 905-MW_e turbine it will house, to be 40 × 80 × 40 m high.

The reactor-maintenance building, estimated to be 30 × 30 × 15 m high, provides working areas for radiation containment during maintenance, assembly, disassembly, and repair functions outside the reactor building. Specifically included in the reactor-maintenance building are a warm cell for maintaining radioactive components, an area for decontaminating activated components and removing activated target material, and a waste processing area.

The tritium-recovery building, estimated to be 30 × 30 × 15 m high, includes systems for fuel cleanup, atmosphere detritiation, and water detritiation. Other buildings have dimensions as indicated in Fig. 8-14.

Shielding. The shield system for Cascade, schematically represented in Fig. 8-17, consists of a 2.0-m-thick borated-water radiation shield, a 2.5-m-thick concrete biological shield, and a cooling system yet to be designed. The radiation shield attenuates neutron and gamma fluxes during operation to levels that keep damage and activation within safe bounds. The design allows maintenance workers to service the heat exchanger and vacuum chamber, which are located behind the radiation shield, shortly after shutdown.

Two materials were considered for the radiation shield: boronated graphite and

borated water. The shield structure is aluminum, 10% by volume, and the coolant for the solid shield is water, also 10% by volume. We chose borated water for the reference design after studying three arrangements:

1. Borated-water shield behind the blanket.
2. Boronated-graphite shield behind the blanket.
3. Boronated-graphite shield behind a 0.5-m SiC reflector, which itself is placed behind the blanket.

The SiC/boronated-graphite design was considered primarily in response to concern that the presence of neutron-absorbing materials immediately behind the blanket would reduce the tritium-breeding ratio in the reactor by about 5%. In fact, without a reflector the tritium-breeding ratio would be marginal for the reference blanket design. If either of the first two arrangements is used, an adequate tritium-breeding ratio can be restored by one of three methods:

- Thickening the blanket's front zone by 2 to 3 cm.
- Enriching the lithium in ⁶Li.
- Modifying the shields so that they reflect neutrons by removing the boron from the 30 to 50 cm closest to the reactor.

Figure 8-18 displays the calculated neutron flux distribution in the shield. From this figure, it appears that borated water has the best flux attenuation capability, followed by the SiC/boronated-graphite arrangements. The boronated-graphite shield shows the poorest attenuation capability.

Fig. 8-17. Schematic of the Cascade shield system.

We estimated the maximum neutron flux in metallic materials behind the shield at which direct access to these materials would be allowable 1 d after shutdown. If we assume that the plant operates for 30 y and that a work week is 40 h, continuous direct access is possible for a dose rate of 2.5 mrem/h.

The allowable neutron flux was estimated to be about 4×10^7 n/cm²·s for nickel- and iron-based alloys. We extrapolated the results of Fig. 8-18 to the required limiting flux, assumed a neutron attenuation coefficient of 0.078 cm^{-1} , and estimated the required shield thicknesses to be 1.58 and 2.04 m, respectively, for nickel- and iron-based alloys.

Selection of a 2.0-m borated-water shield maximized flexibility in selecting materials for components in this location. To ensure a tritium-breeding ratio greater than unity, two coolant loops will be required: one to carry unborated water to the portion of the shield nearest the reaction chamber, and the other to carry borated water to the rear of the shield.

The biological shield was designed to assure that the operating neutron flux outside it is attenuated to below $0.1 \text{ n/cm}^2 \cdot \text{s}$, thus keeping the dose rate below 1 mrem/h. As shown in Ref. 46, the attenuation coefficient of concrete on the operating biological dose is estimated to be about 0.0617 cm^{-1} . Using this attenuation coefficient and the allowable neutron flux for shutdown dose-rate considerations mentioned above, we estimated the required additional concrete thicknesses for biological shields to be about 2.5 and 3.1 m, respectively, if nickel- and iron-based alloy components are placed outside the radiation shield but inside the biological shield. In conjunction with the 2.0-m borated-water shield, a 2.5-m concrete biological shield was thus adopted for the baseline design.

Safety. The Cascade approach to fusion power has some attractive safety features. First, tritium inventory and routine permeation release are very low. Second, the use of low-activation materials brings the benefits of minimum radioactive inventory and minimum production of decay heat. Third, the use of a solid breeder avoids the potential for a large uncontrolled chemical energy release. Cascade thus

Fig. 8-18. Neutron flux distribution in three proposed radiation shields.

minimizes both the radioactive source term and the mechanism for its release, thereby approaching the inherent safety and environmental potential of fusion. These inherent advantages can translate into economic advantage because unnecessary safety systems are avoided and because more components and equipment are classified as conventional rather than nuclear. By planning for safety early in the design stage, we can minimize safety-related costs.

Safety-related equipment includes all systems and components that function to prevent events or mitigate the consequences of events that would result in offsite exposures that approach or exceed regulatory dose limits.^{47,48} Licensing criteria are adopted to protect the safety of the general public—power plants are designed so that safety equipment is sufficient to safely shut down reactor operations after any credible emergency. In the context of design practice and standards for nuclear power plants, safety means averting or not causing undue radiological effect on the health of the public.^{47,49} Thus, from a regulatory perspective, safety is measured strictly with respect to public dose criteria. Preventing economic damage to the reactor is the purview of the plant owner, but economic considerations are secondary when safety-related equipment is selected.

The general approach to evaluating safety and establishing safety requirements

is to analyze postulated accidents to determine consequences in terms of the release of radioactivity and doses to operators and the public. In these analyses it is customary to include such factors as the length of time for accident progression and the time available for corrective action. Damage to equipment is considered only insofar as it leads to a radioactive release. For Cascade, much of the safety classification system is inappropriate because there is no pressure-retaining primary coolant boundary. We satisfy the intent of a safety classification system by identifying the equipment necessary to satisfy the radiological safety requirements. Although the accident scenarios and classification systems developed by the fission industry may not apply, the dose guidelines used by the fission industry⁴⁹⁻⁵¹ will probably either be directly applicable or serve as useful references in defining the radiological safety requirements. By applying the established and applicable safety analysis and licensing methodology to Cascade, we can design equipment such that the numerical dose limits are inherently satisfied to the greatest extent possible, avoid nuclear classification of equipment where possible, and identify the minimum number of additional systems required for compliance.

We designed the Cascade blanket so that it would not melt, even with maximum decay heat, thereby eliminating the major energy source that could lead to the release of radioactivity. Cascade's ability to meet this guideline is a significant achievement, approaching the fusion goal of inherent passive safety.

The major contributor to decay heat in the blanket is ^{24}Na produced from (n,α) captures in ^{27}Al . This is a threshold-type reaction, and the fraction of ^{24}Na in the LiAlO_2 can be controlled by moderating the neutron flux entering the LiAlO_2 zone below the threshold energy. The blanket can absorb 710 GJ before melting. The blanket decay power Q_0 at shutdown generated by ^{24}Na is 7 MW, with a radioactive decay constant λ of $1.3 \times 10^{-5} \text{ s}^{-1}$; the total energy produced in the blanket is then Q_0/λ or 540 GJ, which is below that required to melt the blanket. Furthermore, for the reference blanket, the radiation shield is a massive ultimate heat sink that can passively absorb all the decay energy

produced and still experience a temperature rise of only 100 K. An active ultimate heat sink is not necessary.

This analysis assumes that heat generated within the LiAlO_2 will be uniformly distributed to the other blanket materials. In an accident, the chamber would slow to a stop and the chamber contents would be mixed—a chain of events that may not occur in the heat exchangers and granule-transport system. We did not conduct detailed transient thermal analyses, but it is possible to calculate the thermal diffusivity of the LiAlO_2 and thus a thermal time constant for heat conduction. We know that the time necessary for LiAlO_2 to heat adiabatically to the melting point as a result of its own activation is 13 h. The thermal time constant for 1-mm granules is 0.43 s; for a 25-cm-thick slab of LiAlO_2 granules the thermal time constant will be less than 13 h. This means that if the heat exchangers and granule transport system have a maximum LiAlO_2 thickness of no more than 25 cm, the blanket will not melt.

For passive heat transfer to the radiation shield via thermal radiation, the controlling resistance is the insulation. The thermal time constant is 30 h for a 15-cm-thick insulation with a thermal conductivity of $0.1 \text{ W/m}\cdot\text{K}$ and density of 250 kg/m^3 . In the longer time scale, this heat transfer mechanism would further limit the heat rise in the blanket.

The analysis above shows that the Cascade blanket is inherently safe from meltdown and thus from accidental radioactivity release. The vapor pressures of the activation products at high temperatures must be included in safety analyses of radioactivity release. Nonetheless, these results form the foundation for eliminating a dedicated, active, safety cooling system from the baseline reactor design, reducing the containment requirements imposed on the reactor building, and also reducing the costs of system components by permitting the specification of conventional rather than nuclear construction and qualification.

We also developed design guidance for the maximum tritium inventory and leakage allowable without nuclear qualification—i.e., make it unnecessary that the reactor building either be seismically leak-tight or that it include dedicated safety and tritium-control systems. We can

determine the maximum routine tritium leakage rate that will satisfy 10CFR50 (Ref. 50) and the maximum accidental tritium release that will satisfy the 10CFR100 (Ref. 51). The former calculation, detailed in Ref. 52, indicates that a tritium leakage rate from all sources of 100 Ci/d from a 100-m stack, neglecting nontritium contribution to dose, will meet the present 5 mrem/y limit. The allowed leakages from a 10-m stack and at ground level would be 10 and 1 Ci/d, respectively. Cascade's leakage rate is calculated to be 25 Ci/d (see the next article, "Tritium Inventory, Permeation, and Recovery.") We then calculated the maximum tritium release that will meet current accident guidelines. Results show that a tritium inventory of 210 g released over 12 h will meet the 25-rem regulatory guideline of 10CFR100 (Ref. 51) without requiring an effluent release stack or containment leak-rate or seismic qualification of the reactor building. Cascade has a total inventory of 142 g that can be released in less than 30 d, which more than satisfies this requirement (see Table 8-12).

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Tritium Inventory, Permeation, and Recovery

This year we calculated the tritium inventory for the blanket, chamber, granule distribution and transport components, heat exchangers and helium manifolds, shielding, and vacuum chamber. We also examined tritium permeation (permeation is the product of solubility and diffusion) and transport into the vacuum system, vacuum-chamber cooling system, shield cooling system, helium working fluid, and atmosphere of the reactor building. Finally, we designed a tritium-recovery system that will incorporate the flow from all tritium-bearing streams and maintain tritium concentrations well below levels considered safe for public exposure. This work supplements that reported in the 1983 *Laser Program Annual Report*.⁵³

Component	Tritium inventory (g)	Characteristic release time
BeO (based on Al_2O_3)	4	0.4 h
LiAlO_2	100	10 h
SiC chamber	8	130 y
SiC granule distribution and transport components	39	100 y
SiC in primary heat exchangers	75	120 y
Steel manifolds in primary heat exchangers	2	74 h
Steel helium ducting	6	12 d
Al/ H_2O /C/ B_4C shielding (based on steel)	1	28 d
Vacuum-chamber atmosphere	9	Instantaneous
Aluminum vacuum chamber (based on steel)	9	19 d
Fuel cleanup system	11	Instantaneous
Total releasable in less than 30 d	142	
Total secure	122	
Total	264	

To calculate inventory and solubility, we adopted the assumptions of Stickney⁵⁴ about solubility (defined as mass of T_2 per mass of solid) and diffusivity (cm^2/s):

$$S = S_0 e^{(-Q_s/RT)} \times p^{1/2}, \quad (25)$$

and

$$D = D_0 e^{(-Q_d/RT)}, \quad (26)$$

where S is solubility, the pre-exponential S_0 is the solubility coefficient (g of T_2 per g of solid/ $\text{atm}^{1/2}$), Q_s is the heat of solution (kcal/mole), R is the universal gas constant (1.9872×10^{-3} kcal/mole $\cdot$ K), T is the temperature of the solid (K), p is absolute pressure (atm), the pre-exponential D_0 is the diffusivity coefficient (cm^2/s), and Q_d is the activation energy for diffusion (kcal/mole). Characteristic release times τ were conservatively calculated with the solution for diffusion through a plane sheet⁵⁵:

$$\tau = \frac{L^2}{15 D}, \quad (27)$$

where L is a characteristic dimension for escape of the tritium contained in each component. The characteristic release time—the time required for 1/e of the inventory to permeate out of the component when the surface concentration is suddenly reduced to zero—indicates the vulnerability of the tritium inventory to accidental release. [Reference 38 contains an excellent discussion of tritium migration and solubility mechanisms in metals and ceramics. The tritium-recovery system is based on the tritium system test assembly (TSTA) design.⁵⁶]

Table 8-12. Tritium inventories and characteristic release times in Cascade components.

Inventory. We calculated tritium inventories by determining the tritium solubility as a function of temperature and tritium overpressure for each material and component exposed to tritium in the Cascade reactor. The total inventory was calculated by multiplying solubility by the mass, then summing for all materials and components. The tritium overpressure is set by the rate at which DT pellets are injected into the reaction chamber, the tritium breeding rate, and specifications of the vacuum pumping system. We sized the vacuum system to maintain a total pressure of 0.1 Torr throughout the reaction chamber and 0.01 Torr in the heat exchanger. The fraction of tritium is slightly below 40%, giving a tritium overpressure

of 0.04 Torr in the reaction chamber. (The tritium inventories in Cascade materials and components are shown in Fig. 8-19 and Table 8-12. The solubility and diffusion data used are summarized in Table 8-13.)

Our results show that Cascade has a low total inventory of 264 g (excluding tritium contained in the fuel pellet fabrication plant). The vulnerable portion of the tritium inventory (the amount capable of release within 30 h) is 142 g, of which 100 g are contained in the solid LiAlO_2 tritium breeder with a characteristic release time of 10 h.

All of the secure portion of the tritium inventory is contained within the SiC components, with characteristic release times of hundreds of years. These components

Fig. 8-19. Schematic illustration of the tritium inventory in and permeation from Cascade components.

Table 8-13. Solubility and diffusion data used in tritium analysis.

Material	Solubility coefficient (g of T ₂ per g solid · atm ^{1/2})	Heat of solution (kcal/mole)	Diffusivity coefficient (cm ² /s)	Activation energy for diffusion (kcal/mole)	Permeation coefficient ^a	Reference
Al ₂ O ₃	2.9×10^{-3}	18.1	NR ^b	NR	NR	6
BeO	—	—	1.3×10^{-6}	30.8	NR	7
LiAlO ₂	—	—	1.12×10^{-6}	35.8	NR	8
SiC ^c	6.6×10^{-10}	-37.0	1.58	73.6	SD ^d	9
Stainless	4.8×10^{-5}	2.54	2×10^{-3}	12.5	6.4×10^{-2}	10 and 11
Aluminum	—	—	—	—	6.4×10^{-5}	11 (700 K)
					6.4×10^{-7}	12 (500 K)

^a Includes oxide-barrier permeation reduction factor of 100.

^b NR: not required.

^c For temperatures below 1300 K, use results of transport calculations of Ref. 13 (see text).

^d SD: permeation coefficient obtained from product of solubility and diffusivity coefficients.

exhibit very low tritium diffusivity, and the tritium inventory will not approach the solubility limits within the 40-y lifetime of a Cascade plant. Results⁵⁷ show that after 30 y at 1300 K with a conservative tritium overpressure of 0.1 Torr (2.5 times the actual value), the tritium inventory per unit of exposed area in SiC components is 3.3×10^{-6} g/cm². In 30 y, tritium atoms would have diffused just 2 mm into the SiC. Thus, equilibrium calculations on inventory at temperatures below about 1300 K can be overly conservative because the characteristic diffusion times are long compared with the plant lifetime. Since the diffusion time is shorter for components above 1300 K, we conservatively calculated the inventory in these components with the equilibrium solubility data of Ref. 59. The tritium inventory in SiC components below 1300 K (see Table 8-12) was calculated on a non-equilibrium basis.

We computed the tritium inventory for BeO from data in Refs. 58 and 59; because no solubility data were available for BeO, we based our calculations on a comparable ceramic, Al₂O₃. We did not consider tritium generation in the beryllium because this source is expected to be even smaller than the diffusive inventory. Similarly, we based our calculations for aluminum solubility and diffusivity on data available for steel—which results in conservatively estimated inventory and release rates, since aluminum permeation is several orders of magnitude lower than that of stainless steel.

The tritium inventory in the fuel cleanup system is scaled from data presented in Ref. 38. We used a similar fuel cleanup

Component	Permeation (Ci/d)	Comment
Heat exchanger manifolds and helium piping	23.4	Based on steel at 700 K
SiC tubes in surface layer heat exchanger	1.2	Based on SiC at 1600 K
SiC tubes in all other heat exchangers	0.0	Based on SiC below 1300 K
Vacuum chamber	0.03	Based on aluminum at 700 K
Shield coolant system	0.008	Based on aluminum at 500 K
Total	24.6	

system for Cascade and found³ that the tritium inventory in the molecular sieves is 1 g and the tritium inventory in the helium separator is 10 g. The total inventory in the Cascade fuel cleanup system is thus 11 g.

Permeation. The coefficient for tritium permeation across the primary system boundary was calculated as the product of the solubility and diffusivity coefficients⁵⁴ as functions of temperature. This product must be multiplied by the density of the solid, the exposed surface area, the square root of the tritium overpressure, and a conversion curies per unit mass of tritium; the results for all components are summed to obtain the total permeation rate.

The tritium permeation rates of Cascade systems and components—presented in Table 8-14—are extremely low compared with permeation rates of other nuclear power plants. As shown in Fig. 8-19, tritium permeates out of the primary tritium containment boundary at 25 Ci/d, a rate that compares quite favorably with:

- The record for tritium handling in Canadian power plants, which shows tritium release rates from fission-reactor stations of 115 to 170 Ci/d (average from 1978 to 1982), with a low of 50 Ci/d.⁶⁵

Table 8-14. Tritium permeation results.

- The U.S. regulatory guideline of 100 Ci/d, which would mean an individual dose of 5 mrem/y at a site with a 100-m stack and a boundary 500 m from the plant.^{hi} It is noted that the major contributor to Cascade's permeation rate is leakage into the helium secondary coolant, which has a tritium-recovery system within its own containment. Because the primary leak rate is very low, actual losses to the environment were not calculated. Virtually negligible permeation rates into the environment are expected.

The primary permeation path for tritium is through the heat-exchanger manifolds and helium ducting. Since permeation is a strong function of temperature, the calculations were performed with the assumption of a conservative (700 K) operating temperature. We also assumed that the presence of a natural oxide barrier on the helium side would reduce the permeation rate by a factor of 100. (This factor has been observed to vary from 10 to 10^4 ; see Ref. 65.) The 2-mm heat exchanger tubes do not permit tritium breakthrough in 30 y of full-power operation at temperatures below 1300 K and predict only 1.2 Ci/d at 1600 K. Tritium would permeate through the aluminum vacuum chamber and shield coolant tubes at a very low rate because of the effects of low operating temperatures and the natural Al_2O_3 oxide surface that forms.

Tritium-Recovery System. The tritium-recovery system is designed to prepare a highly pure DT fuel stream for delivery to the fuel pellet fabrication plant, maintain tritium concentrations below occupational limits in all fluid streams, and provide for rapid cleanup of the building atmosphere in the event of an accidental release. Figure 8-19 shows a flow schematic of the tritium-recovery system.

Primary tritium recovery for Cascade takes place in the vacuum system. Tritium is transported through the vacuum system to a fuel-cleanup system, which is based on the design developed for recovering fuel from the plasma exhaust of the MARS tandem mirror reactor.³⁸ The walls and ducts of the vacuum chamber provide primary containment of tritium. However, as shown in Fig. 8-19, tritium permeates out of its primary containment and into several other systems, including:

- Secondary coolant (helium working fluid) loop.
- Shield coolant loop.
- Vacuum chamber coolant loop.
- Reactor building atmosphere.

The tritium-bearing stream from any of these systems can be directed to a tritium-recovery unit as necessary. Below, we discuss the requirements for each unit.

The chamber exhaust consists principally of unburned DT fuel, tritium produced in the blanket, helium produced in fusion reactions and in the blanket, pellet debris, and vaporized species from the blanket surface layer. In similar fashion to the fuel cleanup system described in Ref. 38, this gas stream is routed to columns of molecular sieves (type 5A) cooled to 77 K, which absorb all impurities except helium. When a column is saturated with impurities, it is taken off-line and regenerated at 700 K while the stream is routed to another column. The system is based on a design now being developed at the TSTA.⁵⁷ This design will remove C, N, O, heavy metals, and some He, as well as decompose tritium-bearing hydrocarbons. The impurity gases contain tritium in low concentrations and are prepared for direct disposal. If the level is too high, the gases are routed to an oxidizing system and the tritium absorbed on desiccants. The water is then routed to the tritiated-water treatment system. After passing through the molecular sieves, the chamber exhaust is composed of only D, T, and He. Researchers at TSTA studied a technique for separating helium from hydrogen that uses helium evaporation from a condensed falling film of hydrogen isotopes. This method was adopted by the authors of Ref. 38 and will also serve for Cascade.

The largest fraction of tritium permeation from the primary tritium-recovery system is leakage into the secondary helium loop. If this tritium were allowed to build up over the full life of the plant, the tritium inventory in the secondary loop would be 27 g, or about one-fifth of the 129-g vulnerable inventory in the reactor building. Even if the entire contents of this loop and the reactor building were released and transported to the site boundary, public exposure would be within present regulatory guidelines. Thus, there is no explicit need to recover tritium from

the secondary loop. However, a clean-up system for the secondary loop is probably needed to maintain helium purity at proper levels for controlling corrosion and maintaining the coolant chemistry. If it becomes necessary to reduce tritium inventory to an absolute minimum, methodology to remove tritium from the secondary loop can be added to the cleanup system with little additional expense.

The requirements for removing tritium from the building atmosphere are dictated by the maximum allowable occupational exposure to airborne tritium, 5×10^{-6} Ci/m³ air, as set forth in regulatory guide 10CFR20 (Ref. 67). The building atmosphere is maintained below occupational exposure limits with the atmosphere detritiation system (ADS). The unit is sized by the more restrictive of two processes:

- Remove routine vacuum-chamber leakage from the reactor building atmosphere.
- Bring tritium concentrations in the atmosphere of the reactor and tritium buildings to the level specified in the guidelines on occupational exposure following an accidental release.

For Cascade, the more restrictive requirement is set by a hypothetical accidental tritium release into the reactor building. The ADS processing rate is governed by economic considerations in that, following a release, the sooner the tritium concentration is below regulatory limits, the sooner workers can return to the area and restart the reactor. The atmosphere processing rate P is calculated from

$$C_t = C_0 e^{-\left(\frac{P}{V}t\right)}, \quad (28)$$

where C_t is the desired concentration at time t (5×10^{-6} Ci/m³), C_0 is the initial concentration following the release, V is the building volume, and t is the time to achieve C_t .

If we assume that the entire 129-g vulnerable inventory in the reactor building is instantly released to the 6×10^4 m³ of available building atmosphere, and that 2 d is an economically acceptable time for tritium recovery, the processing rate required is 150 m³/min. The atmosphere-processing system proposed in Ref. 38 consists of individual units of 140 m³/min capacity. A single unit would thus serve the needs of Cascade, requiring 2.2 d for

full cleanup of the reactor building. The same unit would require 8.7 h to process a maximum release in the tritium building. (These scenarios assume that no tritium is released to the public. Because all of the Cascade inventory could be released to the public without exceeding regulatory guidelines, this approach is consistent with a philosophy of assuring that any external releases will be as low as possible.)

During routine operation, the ADS would process the secondary coolant at a rate of 136 m³/min to remove the 25 Ci/d in-leakage and maintain the tritium inventory in the helium to 8 Ci. The reactor building atmosphere would be routinely processed at a rate of 4 m³/min to remove the 0.03 Ci/d leakage from the vacuum chamber and thus maintain tritium concentrations below 5×10^{-6} Ci/m³.

The tritium concentrations in water coolants are maintained below 10CFR20 regulatory guidelines¹⁷ with a water-detritiation system. The low tritium leakage into the shield and vacuum chamber coolants eases the requirements on this system. A state-of-the-art review of tritium recovery from water coolants³⁸ concludes that technologies for removing tritium from water are commercially available. The system selected in Ref. 38 and proposed for Cascade uses electrolysis of the tritiated water followed by separation of tritium by cryogenic distillation. Electrolysis of water is well established in industry and is routinely used for hydrogen production. In the present application, the water-detritiation system maintains tritium concentration in the water at the 10CFR20 requirement of 0.1 Ci/m³ with a water-processing rate of 0.4 m³/d, or 16 kg/h. For comparison, the system proposed for Ref. 38 processes water at 68 kg/h.

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Cost Analysis

Analysts at GA Technologies, Inc.³ estimated the capital cost of the Cascade power plant on the basis of the estimated cost of high-temperature gas-cooled reactors (HTGR)⁶⁸⁻⁷⁰ and supplemental information on the cost of unique Cascade

Cascade Reactor

components and systems. With a method consistent with our other economic studies at LLNL,⁷¹ we used GA's estimate to calculate the cost of electricity (COE). In this article we summarize these estimates.

Capital Cost. The fission power industry is considering lowering costs by dividing construction activities into two grades: nuclear and conventional. Table 8-15 shows capital-cost estimates for Cascade for two possible types of construction—nuclear plus conventional and all conventional. By nuclear-plus-conventional construction, Cascade would be constructed according to standards and methods for fossil-fuel plants except where the more costly methods and standards for fission-reactor plants are appropriate. It is assumed that Cascade would include a combination of nuclear and non-nuclear systems similar to that found in modern fission-reactor plants.

Table 8-15. Capital cost summary for the Cascade power plant. All figures are in millions of 1985-constant dollars.

Account No.	Description	Capital cost	
		Nuclear plus conventional (\$ million)	All conventional (\$ million)
20	Land and land rights	5	5
21	Structures and improvements	84	37
22	Reactor plant equipment	283	254
23	Turbine plant equipment	184	156
24	Electric plant equipment	71	47
25	Miscellaneous plant equipment	16	12
26	Heat rejection system	22	22
27	Driver equipment	133	133
28	Target factory equipment	100	100
	Total direct cost	898	766
91	Construction services and equipment	180	110
92	Home office engineering and services	135	43
93	Field office engineering and services	90	34
94	Owner's cost	63	43
95	Project contingency	137	100
96	Escalation during construction	0	0
97	Interest during construction	140	102
	Total capital cost	1643	1198

It has been suggested, for a variety of reasons, that an ICF power plant may have less nuclear-grade construction than does a fission plant. Therefore, to determine the lowest possible construction cost, we calculated the cost of all-conventional construction, assuming no nuclear-grade components.

The driver cost (\$133 million direct, shown as account 27) represents a goal for advanced short-wavelength lasers rather than an actual cost estimate. For the Cascade reference design, the driver energy is 1.5 MJ. The target factory cost of \$100 million (account 28) is a rough estimate based on the cost of current facilities that mass-produce precision products such as semiconductor microchips.

The total direct capital costs for the nuclear-plus-conventional case and the all-conventional case are \$898 million and \$766 million, respectively. The fractions used to determine the indirect costs (accounts 91 through 94) for the nuclear-plus-conventional case are averages of the fractions for nuclear and fossil plants (see Ref. 72). The fractions used for the all-conventional case are those used for the single-unit fossil plant in Ref. 72. A project contingency of 10% and interest during construction (IDC) are added in both cases. The IDC factor of 5.3% is based on a real, after-tax cost of money of 2.83% and a construction period of 8 y. The escalation during construction is zero since we are using a constant-dollar analysis and have assumed that the escalation rate is equal to the general inflation rate.

The total capital costs for nuclear-plus-conventional and all-conventional construction are \$1643 million and \$1198 million,

Table 8-16. Comparison of costs for Cascade and other nuclear- and coal-powered plants. (Levelized busbar cost in 1985-constant dollars.)

	Cascade			Other plants		
	Nuclear plus conventional	All conventional	HTGR	PWR	Coal	
Thermal power (MW _t)	1670	1670	2240	2400	2285	
Net electric power (MW _e)	815	815	855	800	800	
Availability factor (%)	70	70	75	70	70	
Capital cost (million)						
Direct	898	766	901	625	650	
Total	1643	1198	1542	1181	996	
Cost of electricity (¢/kW _e h)						
Capital	2.73	1.99	2.28	2.00	1.69	
Operation and maintenance	0.96	0.96	0.66	0.71	0.56	
Fuel	0.01	0.01	0.96	0.80	2.27	
Total	3.70	2.96	3.90	3.51	4.52	

respectively. Since the net electric power of Cascade is 815 MW_e, these capital costs translate to \$2016/kW_e and \$1470/kW_e, respectively. Note that most of the savings in the all-conventional case result from the lower indirect costs.

Cost of Electricity. Table 8-16 compares the COE for Cascade with that for several other plants. All direct capital costs were estimated by GA Technologies. The indirect costs and costs of operation and maintenance and fuel are based on information from GA and from Ref. 72. The annual fixed charge rate (the annual cost per dollar of capital investment) for the constant dollar analysis is 8.3%. As indicated, the COE for the nuclear-plus-conventional version of Cascade is midway between the COE for the HTGR and the pressurized water reactor (PWR). The all-conventional case is 20% lower than nuclear-plus-conventional construction. Note that the COE from the coal plant is sensitive to the assumed real escalation rate of coal, which is 1.5%/y in this case. If we assume no real escalation in coal prices, the COE for the coal plant drops to 3.83¢/kW_eh.

Major uncertainties that can significantly affect the Cascade results are target performance and cost of the driver. For Cascade we used the gain-vs-energy curve corresponding to advanced target concepts (see "Economic Requirements for Competitive ICF Power Production," later in this section). This resulted in a driver energy of 1.5 MJ for a target yield of 300 MJ. If the advanced target technology is not developed and 3 MJ is required to achieve this same yield, the direct capital cost of the driver increases by \$83 million, and the COE for the nuclear-plus-conventional case increases by 6.8% to 3.95¢/kW_eh. If the 1.5-MJ driver cost doubles to \$266 million, the COE increases by 10.9% to 4.10¢/kW_eh. If both the required driver energy and the cost per MJ double, the direct capital cost of the driver increases by \$299 million, and the COE increases by 24.6% to 4.61¢/kW_eh. Even so, this is still within the range of economic competitiveness with future nuclear- and coal-powered plants.

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Pulse*Star Reactor

Introduction

Pulse*Star is a pool-type reactor in which the fusion chamber, pumps, and heat exchangers are submerged in a pool of liquid-metal coolant (see Fig. 8-20). The baseline design uses a liquid lead/lithium eutectic (Li₁₇Pb₈₃) to transfer heat from the fusion chamber to the steam generators, to breed tritium, and to protect the bell jar, which is the structural element of the chamber.

Our work on Pulse*Star this year is presented in the next three articles:

- The first article discusses the pool support structure and the pool, pump, and heat exchanger connections. These are patterned after the Superphenix reactor, a 1200-MW_e liquid-metal fusion breeder reactor (LMFBR) located near Lyon, France.
- The inner screen of Pulse*Star must intercept the LiPb spray and keep it from entering the driver beam lines. It is wetted on both sides and—when the fusion pulse occurs—some liquid metal is ablated, which results in a large impulsive load on the screen. The second article discusses the design of the Pulse*Star screen necessary to survive this pulsed load.
- Containing tritium in LiPb systems has proven to be a difficult problem for magnetic fusion and ICF reactors. The last

Fig. 8-20. Pulse*Star reactor concept.

article in this subsection describes progress made in this area.

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Pulse*Star Reactor and Coolant System

Pulse*Star's liquid-metal pool and reactor vessels are supported from the top (as shown in Fig. 8-21); their total weight is transmitted by the lid to the reinforced concrete cylinder surrounding the pool. Suspending the pool above the ground in this fashion eliminates any possible interference with surrounding structures that might result from thermal expansion as the reactor temperature changes during startup and shutdown.

In addition, the reactor vessel is insulated to minimize energy loss and reduce the temperature of the surrounding concrete. Cooling coils between the liquid level and the lid (on the outside of the reactor vessel) and thermal baffles (inside) are necessary to keep the lid cool enough to prevent damage to the concrete, to minimize thermal expansion and bowing, and to enable access to the head area by maintenance personnel. The reinforced concrete cylinder is 2 m thick at the bottom and 1.5 m thick at the top. The 20-m span across the top of the vessel requires a 2.5-m-thick lid made of steel webbing and concrete.

The thickness of the reactor-vessel wall is dictated by two requirements: sufficient strength to contain the pressure of the liquid-metal pool and sufficient

Fig. 8-21. Pulse*Star reactor tank and support system layout.

strength to support the weight of the pool and reactor vessel. Figure 8-22 shows how these requirements affect reactor-vessel thickness. The wall would most likely be stepped to minimize manufacturing costs.

The coolant enters the reactor chamber from above (see Fig. 8-21) and is discharged at the bottom. Because the center of the reactor is under vacuum, the coolant leaving the reactor is not pressurized. The pressure must be increased before the coolant can be forced through the steam generators. Thus, the reactor outlet is connected to four pump inlets at the bottom of the pool. The pumps discharge liquid metal into a hot plenum at the top of the pool, after which the liquid metal flows circumferentially around to the heat exchangers. Heads of the hot and cold plenums differ enough that the coolant is under sufficient pressure to flow downward through the heat exchangers and exit into the cold plenum.

In designing the Pulse*Star heat transport loop, we wish to eliminate piping connections between the primary pumps and heat exchangers inside the pool so that components can be singly lifted out of the pool for maintenance without disconnecting lines. The Pulse*Star heat exchanger/pump connections are based on those used at the Superphenix LMFBF plant,⁷³ which is also a pool-type reactor. Each pump fits into a well that has a non-leaktight connection with the reactor outlet piping. This connection permits easy removal of the pumps for maintenance. Each heat exchanger has an overlapping seal at the plenum interface. Both the pumps and heat exchangers are supported from the lid.

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Screen Design

A key structural component in the Pulse*Star design is the screen that is exposed to pulsed loads from fusion

Fig. 8-22. Reactor vessel wall thickness.

reactions. Designs proposed for the screen include metallic honeycomb and perforated plates. We analyzed the stress in one candidate design: a perforated plate with a triangular hole pattern.

Consider the screen to be cylindrical and made of 2-1/4 Cr-1 Mo steel with an impulsive load applied at the inside surface following each fusion pulse. Glenn and Young⁷⁴ show that the peak wall stress is related to the impulse (or momentum) by

$$\sigma = \frac{M}{At} \sqrt{\frac{E^*}{\rho}}, \quad (29)$$

where σ is the peak stress (Pa),

M is the applied momentum (kg·m/s),

A is the wall surface area (m²),

t is the wall thickness (m),

E^* is the effective modulus of elasticity (Pa),

ρ is the wall material density (kg/m³).

The applied momentum is just

$$M = mv = \sqrt{2m\epsilon}, \quad (30)$$

where m is the mass (kg) of material ablated as a result of energy absorbed from x rays and fusion-fuel-pellet debris, v is the velocity of the ablated material (m/s), and ϵ is the energy absorbed in the ablated material (J).

The energy per unit area absorbed at the horizontal centerline of the wall is approximately

$$e/A \approx \frac{Y/3}{4\pi R^2} \quad (31)$$

where Y is the yield of the fusion reaction (J) and R is the radius of the wall (m). If we assume a 600-MJ yield, a 4-m wall radius, and a 20- μ m-thick ablated material, and if the density of the ablated material (assuming LiPb eutectic) is 10^4 kg/m³, then the momentum per unit area M/A is 600 Pa·s.

There is some question about what magnitude of E^* to use in Eq. (29). Slot and O'Donnell⁷⁵ give values of E^* vs λ for steady plane-strain loading in both triangular- and square-hole configurations. Here, λ is defined as unity minus the ratio of hole diameter to hole spacing. For want of better values, we use the E^* values of Slot and O'Donnell for our dynamic loading case. We used the methodology of Pitts⁷⁶ to calculate the wall thickness required to withstand fatigue loading on the screen exposed to fusion pulses at 5 Hz for 5 y, and incorporated strain range data $\Delta\epsilon$ of 6×10^{-4} (Ref. 77). The allowable strain range is half that of the data, which results in: an allowable stress range: $\Delta\sigma = 3 \times 10^{-4} E^*$. The stress range is defined as the peak-to-peak variation of stress. For our loading conditions, the peak stress is half the stress range, so the allowable peak stress is $\sigma = 1.5 \times 10^{-4} E^*$ as defined in Eq. (29). The density ρ , modulus of elasticity E , and Poisson's ratio ν of 2-1/4 Cr-1 Mo steel are 7800 kg/m³, 1.6×10^{11} Pa, and 0.3 respectively.

Using the numerical values given above, we can solve Eq. (29) for the

required screen thickness as a function of ligament efficiency λ . Results are shown in Table 8-17. Thermal stresses are considered negligible because the screen is cooled by liquid metal.

A representative ligament efficiency of $\lambda = 0.5$ (hole diameter equals half the hole spacing) requires a screen thickness of only 153 mm. This is a reasonable thickness if the flow characteristics required for the screen can be obtained. Required screen thickness decreases away from the horizontal plane because the peak momentum per unit area that we used occurs only there. In the actual screen design, the thickness could be reduced above and below the horizontal centerline.

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Containment of Tritium

Previous analysis⁷⁸ suggests that about 46 g/d of tritium will diffuse into the steam supply from the Li₁₇Pb₈₃ eutectic in Pulse*Star unless special precautions are taken. If all tritium entering the steam system escapes into the environment, the leakage will exceed acceptable limits by three to four orders of magnitude. In 1984 we considered three possible ways to decrease the amount of tritium entering the steam system: hydrogen dilution, cold trapping in a sodium intermediate loop, and the use of pure liquid lithium instead of Li₁₇Pb₈₃ as a coolant. In 1985 we examined other means of tritium containment: stripping tritium from the Li₁₇Pb₈₃ with counterflowing helium, removing tritium from the water in the steam cycle, and using double-walled heat exchanger tubes and oxide layers with a tritium scavenging gas. We also explored hydrogen dilution further. Our studies show that tritium containment will be neither simple nor inexpensive.

Helium Countercurrent-Stripping System. It was suggested that since tritium is so insoluble in LiPb, a carrier gas such as helium could be bubbled through the LiPb, allowing the tritium to diffuse into the gas and thus be stripped from the

Table 8-17. Required screen thickness (last column) for a triangular hole pattern, found by applying Eq. (29).^a

Ligament efficiency λ (Pa)	Ratio of actual to effective modulus of elasticity E^*/E	Effective modulus of elasticity E^* (Pa)	Allowable peak stress σ (Pa)	Required thickness t (mm)
0.05	0.0172	2.75×10^9	4.12×10^5	363
0.1	0.0520	8.32×10^9	1.25×10^6	497
0.2	0.155	2.48×10^{10}	3.72×10^6	287
0.333	0.325	5.20×10^{10}	7.80×10^6	199
0.5	0.545	8.72×10^{10}	1.31×10^7	153
0.7	0.800	1.28×10^{11}	1.92×10^7	127
1.0	1.0	1.60×10^{11}	2.40×10^7	112

^a Momentum per unit area M/A is 600 Pa·s, modulus of elasticity E is 1.6×10^{11} Pa, and density ρ is 7800 kg/m³.

eutectic. Several types of stripping towers were examined. Figure 8-23 shows the principles of a shower-tray arrangement. LiPb is poured in the top and cascades from tray to tray. Each tray breaks the LiPb into small droplets to provide a large surface area. The helium gas is forced to flow in a zigzag fashion to provide a long path. Each space between trays represents a stage, and in each stage the path is long enough that the gas and liquid come to equilibrium with respect to tritium concentrations. The mole fraction of hydrogen isotopes in the gas is y_i as it enters and y_{i+1} as it leaves a stage, while the mole fraction in the liquid is x_{i+1} as it enters and x_i as it leaves. By mass conservation,

$$L(x_{i+1} - x_i) = V(y_{i+1} - y_i), \quad (32)$$

where L is the flow rate of the liquid (in moles/s) and V is the flow rate of the gas. For n stages, the result is:

$$L(x_{n+1} - x_i) = V(y_{n+1} - y_i), \quad (33)$$

which simplifies to

$$y_{n+1} = \frac{L}{V}(x_{n+1} - x_i) + y_i, \quad (34)$$

the equation for a straight line with a slope L/V in the x,y plane (where we can use the continuous variables x and y in place of the discrete ones x_{n+1} and y_{n+1}). This line is called the operating line for the tower. Clearly, if we want to operate with the minimum helium-flow velocity, we want the slope of this line to be as great as possible. The slope can be increased until the line is tangent to the curve representing Sievert's law in the x,y plane. That is, hydrogen isotopes will continue to diffuse out of the liquid until the concentration in the gas is too high. Sievert's law for the mole fraction is

$$x = \sqrt{\frac{P_H}{K}}, \quad (35)$$

where P_H is the pressure of the hydrogen isotopes in the gas and K is Sievert's constant. If we assume that

$$y = \frac{P_H}{P_{\text{tot}}} \quad (36)$$

Fig. 8-23. Shower-tray countercurrent tower for tritium removal.

as is the case for ideal gases at low partial pressures, then Eq. (35) can be written as

$$y = (9.28 \times 10^8) x^2, \quad (37)$$

where we have inserted Sievert's constant for hydrogen at 770 K and have assumed the helium is at 1 atm. Solving Eqs. (34) and (37) simultaneously and finding the slope (dy/dx) of Eq. (37) at that point will give us the maximum slope of the operating line. The minimum gas-flow velocity $V_{\min}$ found in this way is just

$$V_{\min} = 2.7 \times 10^{-10} \frac{L}{x_i}, \quad (38)$$

if we assume that $y_i \ll (4 \times 10^9)x_i^2$ (i.e., that the partial pressure of hydrogen initially in the helium is far less than the partial pressure that would be created by the LiPb as it exits the stripping tower).

$L = 1200$ kg-mole/s is the flow rate of LiPb required through the heat exchangers; x_1 , the mole fraction allowable in the liquid as it leaves the stripping tower (and enters the heat exchangers), depends on the hydrogen-to-tritium dilution ratio in the LiPb. For no dilution, $V_{\min} \sim 150\,000$ kg-mole/s = $610\,000$ kg/s, which is clearly an impractically high helium-flow rate. For comparison, consider that the helium flow rate in a 1000-MW_t HTGR is approximately 1200 kg/s. If hydrogen is added to the LiPb (see below) and dilutes the tritium (and deuterium) by a factor of 1000, then the required helium flow rate is 610 kg/s. For this case we estimate that a shower-tray stripping tower similar to the one in Fig. 8-23 would be roughly the same size as the entire LiPb pool (about 20 m × 20 m). This does not look like an attractive option.

Removal of Tritium from Water in the Steam Cycle. A variety of systems have been proposed for removing tritium from the water in the steam system. These systems use various combinations of liquid- or vapor-phase catalytic exchange, water distillation, and electrolysis. Once the gaseous hydrogen isotopes are removed by these methods, cryogenic distillation is used to separate the various isotopes. Mound Laboratories compared costs of four tritium-removal systems for the International Tokamak Reactor (INTOR) study.⁷⁹ They found that the capital cost depended on the process flow rate because the flow rate determines the size, and therefore the cost, of the equipment. The volume of tritiated water to be processed in any system is given in liters per day by

$$V_w = \frac{D_1}{C_1 \epsilon}, \quad (39)$$

where D_1 (Ci/d) is the tritium permeation rate, C_1 (Ci/L) is the desired tritium concentration in the water (to keep escape to the environment within limits), and ϵ is the efficiency of the process. The four systems were sized for $C_1 = 1$ Ci/L, $D_1 = 1000$ Ci/d, and $\epsilon = 0.9$ corresponding to $V_w = 1100$ L/d. Capital costs between \$15 and \$25 million were found, and it was estimated that the costs would scale as $V_w^{0.55}$. If we simply use these relations with the

maximum Pulse*Star tritium permeation rate (46 g/d = 460 000 Ci/d), we can see that the water processing rate would be 510,000 L/d, and the tritium extraction system would cost 29 times the Mound estimate, or about \$600 million. [Caution should be used with this extrapolation, however, since it is outside the limits recommended by Mound (<10 000 L/d)]. To stay within reasonable monetary bounds, the tritium permeation rate should be kept below about 1 g/d, which is lower than Pulse*Star's uninhibited level by a factor of 50. It appears that systems to remove the tritium from water alone cannot do the job economically.

Oxide Layers and Double-Walled Heat Exchanger Tubes. There has been a great deal of interest in the use of oxide layers on heat exchanger tube walls to inhibit tritium permeation.^{80,81} Experiments have shown barrier effects of factors of 10 to 1000, depending on conditions and materials. The extent to which the efficacy of an oxide layer can be depended on is, however, controversial because the layer is often rather fragile and may crack or even flake off, particularly under thermal cycling. On the other hand, operating at high temperature and in an oxidizing atmosphere (as on the steam side of the Pulse*Star heat exchanger tubes) produces a rather stable environment for the oxide layer, and the layer even seems to heal itself after thermal-cycle-induced degradation. In such conditions it appears reasonable to assume a 100-fold reduction in tritium permeation over the bare-metal value. Before relying on this barrier for Pulse*Star, however, we would need to thoroughly analyze the conditions under which thermal cycling will occur. For example, thermal cycling would occur whenever a heat exchanger was removed for servicing. The maintenance procedure would have to be designed to include re-establishment of the oxide layer before re-introducing the heat exchanger to the tritium environment. The maintenance procedures used on the Superphenix LMFBR seem to support these concerns.

Another suggestion for controlling tritium permeation in LiPb systems is to use double-walled heat exchanger tubes and flow oxygen between the walls as a sweep gas (see Fig. 8-24). Ideally the tritium

would diffuse through the first wall as atomic tritium, leave the wall as T_2 , and react with the O_2 to form T_2O that would not be likely to diffuse through the second wall. The tube walls would have to have oxidized surfaces to keep the T_2O from reducing. In one experiment at Idaho National Engineering Laboratory⁸² it was observed that more than 98% of the tritium diffused through a single barrier was converted to T_2O . Questions have been raised about the importance of the nature and required integrity of the oxide layer on the walls. For example, thermal cycling may crack the oxides, allowing T_2O to come into contact with ferrous material, where it will be reduced. More experimental work will be required to determine the true value and cost of this proposed method for controlling tritium permeation.

Hydrogen Dilution. In last year's discussion of hydrogen dilution,⁷⁸ we assumed that the dilution occurred in the reaction chamber. It could, of course, also occur in the LiPb. For example, high-pressure hydrogen could be injected into an array of metal tubes submerged in the LiPb pool. For 100-atm hydrogen and with tube walls similar to those in the heat exchangers, a tube area about the same as that of the heat exchanger would be required to obtain a hydrogen dilution ratio of about 30 in the liquid. The hydrogen so added to the liquid would be evolved in the reaction chamber and carry the dissolved tritium with it. About 30% of the tritium, 5 mg/s, is created in the LiPb. The factor-of-30 dilution in the liquid implies a factor-of-1000 dilution in the chamber [see Eq. (35)], or about 5 g/s of hydrogen entering the chamber. To pump this out of the chamber at a pressure of 5×10^{-3} Torr would require a surface area of 280 m². Since this is about the same as the surface area of the entire chamber, this is clearly an impractical idea. It would appear that only about a factor-of-10 dilution could be achieved in practice, and that would require a substantial investment in equipment.

Conclusions. An inexpensive solution to the problem of tritium containment in the LiPb version of the Pulse*Star reactor is not certain. The use of oxide layers in single- or double-walled heat exchanger tubes seems to be attractive, but questions are not resolved about the containment

Fig. 8-24. Double-walled heat exchanger tube with oxide layers.

factor obtained, the integrity of the oxide layers, and the system costs. A combination of oxide layers and double-walled tubes with tenfold hydrogen dilution may reduce tritium permeation by a factor of 1000, and that may be acceptable. The full costs of such a solution have not been estimated in detail, but they will not be negligible. The pure-lithium version of Pulse*Star⁸³ solves the tritium permeation problems but also incurs increased costs. The capital costs for the required intermediate sodium loop and for the lithium fire-suppression systems would be significant, even though they could be offset by decreased operating costs for pumping the lighter metals.

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General Studies

Implications of Electricity Projections and Utility Requirements for Future ICF Reactors

We contracted with Titan Systems, Inc., to review current and projected energy-use patterns in the United States, utility requirements for future power plants, and the implications for ICF reactors.⁸⁴ This article summarizes the key findings of their

Fig. 8-25. Electricity generation in quads ($Q = 10^{15}$ Btu) and as a percent of the total U.S. energy consumption to 2010.

study, focusing on electric power production and fusion electric power plants.

Electricity Growth Projections. The demand for electricity in the U.S. continues to grow, albeit at a slower rate than in years past. The rate of increase in electricity use has been greater than the rate of increase in total energy use for the last several decades. As a result, the fraction of the total primary energy consumption used for electricity production continues to increase.

Figure 8-25, taken from the 1983 edition of the National Energy Plan,⁸⁵ shows electricity generation as a percentage of total primary energy projected to the year 2010. Electricity generation, which grew at 6.9% a year in the 1960s and 4.0% a year in the 1970s, required about one-third of the nation's primary energy in 1980. Projected growth is 3.6% a year in the latter part of the 1980s and 1.9% a year between 1990 and 2010.

The electric utility industry also projects electricity demand. Figure 8-26, taken from the most recent National Electric Reliability Council (NERC) publication of 10-year forecasts,⁸⁶ shows summer peak electric demand. The NERC projections are compiled by adding individual projections received from each utility that submits a forecast. The forecasts are non-coincident hourly peaks, and, therefore, their summation does not represent an actual peak demand. However, the projected demand indicates the generating capacity that must be available to meet the load. The figure shows that the peak demand increased at a compounded rate of 7.5%/y from 1965 to

1973. The compounded rate of increase from 1973 to 1983 was much less—2.3%. The figure includes 10-y projections of peak demand starting with the projection released in 1974. The drop in the growth rate of peak demand during the past 10 y is evident. The actual demand in 1983 was about 40% less than what was forecast in 1974. It is clear that projections tend to reflect recent experience: during periods of high growth the projected growth rates are high and vice versa.

Net annual energy use is also projected by NERC, whose projections tend to be somewhat higher than those made by DOE. The NERC projects the use of about 35 Q of primary energy for electricity production in 1990; the DOE study predicts 32 Q. Extrapolated, the NERC estimates predict 43 Q in the year 2000, while DOE predicts 39.

The NERC also compiles information on types of fuel used to supply electricity. In 1983, coal was used to generate 54.8% of the 2.34×10^9 MW·h of electricity produced, nuclear 12.9%, hydro 14.2%, gas 11.6%, and oil 5.9%. By 1993, the coal fraction is projected to remain about the same (54.6%), while the nuclear fraction increases to 23.6% of generation. To meet the increase in demand by 1993, utility plans call for 107 new coal-fired plants (equivalent to 57.8 GW_e), 50 new nuclear units (55.1 GW_e), and 30 oil-to-coal conversions (6.0 GW_e), with 351 units retired (15.1 GW_e).

Titan Systems examined trends in energy use to determine the fraction of U.S.

Fig. 8-26. Diagram of actual summer peak electricity demand for the period 1960 to 1983 and projections for 1984 to 1993.

energy consumption that goes for electricity production. Data were gathered for four major industries representing about 7% of the total primary energy use in the U.S.: pulp and paper mills, petroleum refining, glass melting, and steel making. In all cases, the total energy consumption per unit of production has decreased during the 1970s and into the 1980s at rates ranging from 1.2 to 1.8%/y. The electricity consumption per unit of production, however, has increased at rates between 0.9 and 3.7%. The fraction of primary energy used for electricity production in the residential and commercial sectors is also increasing; it grew from 49 to 63% between 1973 and 1984. These trends are consistent with the information presented in Fig. 8-25, which shows that an increasing fraction of our total energy use is going for the production of electricity. Hence, it is important that we continue to study electric power production as a potential future application of ICF.

It should be pointed out, however, that the existence of diverse viewpoints and practices makes reliable projections of

energy and electricity use very difficult. Whereas the electricity growth rate for the remainder of the century is about 2.2%/y in Fig. 8-25, some low-growth projections predict an actual decline in electricity use from current levels.⁸⁷ We must also continue to explore alternative applications for ICF in the event that the future market for electric power plants is much smaller than currently anticipated.

Utility Requirements. The most extensive and quantitative set of utility requirements for future power plants was developed in 1981 in an Electric Power Research Institute (EPRI) project for fusion plants.⁸⁸ These requirements were developed from the responses of 32 utilities, which ranged in size from 800 to 17 800 MW_e and represented about 28% of the U.S. generating capacity. The 23 requirements along with measures of their importance (on a scale from vital to slightly important) are listed in Table 8-18. Each requirement also had associated with it one or more determining factors; the factors that describe capital cost, for example, are total cost and cost per kW. All factors

A. Utility planning and finance		
1. Plant capital cost		Vital
2. Plant operating, maintenance, and fuel costs		Important
3. Outage rates		Very important
4. Plant life		Important
5. Plant construction time		Very important
6. Financial liability		Vital
7. Unit rating		Moderately important
B. Safety, siting, and licensing		
1. Plant safety		Vital
2. Flexibility of siting		Very important
3. Waste handling and disposal		Very important
4. Decommissioning		Important
5. Licensability		Vital
6. Weapons proliferation		Important
C. Utility operations		
1. Plant operating requirements		Very important
2. Plant maintenance requirements		Very important
3. Electrical performance		Very important
4. Capability for load change		Moderately important
5. Part load efficiency		Moderately important
6. Minimum load		Moderately important
7. Startup power requirements		Important
D. Manufacturing and resources		
1. Hardware materials availability		Very important
2. Industrial base		Very important
3. Fuel and fertile material availability		Very important

Table 8-18. Summary of utility requirements and their overall weightings (from Ref. 88).

are specified with their parameters, the ranges of these parameters, and numerical ratings for various sub-ranges of the parameters. This description of the requirements allows for a numerical evaluation of candidate power plants.

Item 7 in Table 8-16, the unit rating—that is, the size of the plant in terms of its net electric power output—may be particularly important for ICF. ICF has an inherent economy of scale because a certain minimum investment must be made in the driver independent of how much electricity is produced. For a given driver cost, the more electric power produced, the lower the cost of electricity. Therefore, the size of future power plants has important economic implications for ICF.

Historical data for coal and nuclear power plants show that, in both cases, a number of 100-MW_e units were built early in the development of the technology. Once the technology was proven and utilities gained operating and economic experience, the size of plants gradually increased. Planned additions of power plants from 1984 through 1993 average 1100 MW_e (nuclear) and 540 MW_e (coal).

The benefits of building plants that are smaller than current standard designs have been widely discussed for several

years. Potential advantages of smaller plants include: reduction in the total capital cost; smaller fractions for indirect construction costs such as field construction, supervision, and engineering; and shorter lead times, which increase the utility's ability to follow the growth in demand. It will be many years before we can tell if this current thinking will lead to smaller plants.

In response to one survey, utilities' desire for smaller plants was quantified as equal to 1 to 3 y of growth. In the U.S. now, many large utilities have generating capacities between 10 and 20 GW_e. For a 10-GW_e utility, 1 to 3 y of growth at 2%/y is equivalent to a 200- to 600-MW_e plant. If the same utility experiences a 2% growth rate for the next 50 y, the desirable plant size then will be between 540 and 1600 MW_e. Hence, it is not unreasonable to assume that ICF will be able to take advantage of its economy of scale and still be considered a small plant. Even so, it is prudent for us to investigate small ICF power plants, especially for development and early commercialization.

While we have not yet evaluated any of our reactor concepts using the list of requirements in Table 8-16, it does serve as a check list of parameters to be specified for a concept to receive a full-utility evaluation of its potential as a generating unit.

Results and Conclusions. Titan's examination of recent energy and electricity growth projections, alternative futures, and utility requirements has led to the following results and conclusions:

- Recent projections of energy and electricity demand predict decreasing rates of growth into the next century. The currently projected growth rates for electricity production range from negative values to 2.2%/y. Summer peak demand is estimated to grow at an average rate of 2.5%/y from 1984 to 1993, requiring 119 GW_e of new capacity by 1993.

- Industrial conservation trends and new utility practices, such as load management, appear to contribute significantly to reductions in actual and projected capacity requirements. EPRI estimates that load management may be equivalent to 50 GW_e of new capacity by the year 2000. This represents 20 to 25% of the projected required new capacity.

- Futurists link fusion qualitatively with fission in the areas of complexity, capital intensiveness, and nuclear proliferation. In many cases, fusion is not perceived to have distinctive and strong motivating factors except the potential for utilizing a very large fuel supply.

- New electricity-generating technologies now being commercialized may result in a more difficult market entry for fusion energy systems. Modular generating units with capacities of a few hundred megawatts are under development for both coal and nuclear plants.

- Utilities are stating their preferences for small units (100 to 1000 MW_e) with known construction time and institutional factors, even at higher cost, to avoid uncertainties. A rule of thumb for sizing a future unit is to set its capacity equal to 1 to 3 y of load growth. At present this implies 200- to 600-MW_e plants for a 10-GW_e utility. If the 2% growth rate prevails, in the year 2030 the desirable plant size will be 540 to 1600 MW_e.

- Projections and forecasts are not meant to predict the future, hence planners "... should cultivate the capability of remaining flexible, preserving options, and remaining constantly on the alert for new developments and surprises."⁸⁹

Regarding fusion systems and fusion research and development, we recommend the following:

1. Identify small developmental and commercial reactor units whose designs assure ease of construction, licensing, and operation. Use developmental units to establish a construction and operating history for reactor vendors and utility personnel.

2. Identify and actively evaluate driving factors that lead to high cost, large size, or complex reactor systems with the goal of stimulating novel approaches for circumventing them.

3. Subject fusion reactors to extensive system analysis and critique to identify desirable scientific, technological, and engineering directions.

4. Assess the potential of ICF reactors to repower or retrofit existing utility plants.

5. Continue to explore applications other than electric power production.

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Economic Requirements for Competitive ICF Power Production

The desirability of an ICF power plant will be measured partly by how well it competes economically with other power sources. One figure of merit used to judge economic competitiveness is the projected COE from such a plant. To attempt, now, to estimate the COE of an ICF power plant that is not likely to be built for 20 to 40 years might seem the ultimate exercise in futility. Indeed, we caution against using any COE presented in this article in an absolute sense. There is value, however, in constructing a simple economic model of a future ICF plant that can be used to identify the technical developments that would lead to the greatest reduction in the COE. The model can also be used (if a great deal of care is taken) to make rough comparisons between fusion and other future power sources, provided the same economic assumptions are used in the comparison. It is for these two purposes that we have constructed such a model.

The Economic Model. The economic model we used for this study is described in detail in Ref. 90. In our model, COE is directly proportional to capital cost of the power plant and inversely proportional to net electric power, which is held constant. The capital cost is broken into three items: reactor, driver, and target factory. The cost of the reactor, which includes the fusion chamber and balance of plant, is based on a recent cost estimate for the Cascade reactor.⁹¹ The Cascade reference design has an estimated direct capital cost of \$0.659 billion at a plant thermal power of 1.670 GW_t and a gross electric power of 0.905 GW_e. We can scale the cost to other thermal powers by the usual relation

$$C_{rd} = C_r \left(\frac{P_t}{1.67} \right)^a, \quad (40)$$

where C_r is a reactor cost coefficient (\$ billion), P_t is the thermal power (GW_t), and a is the power scaling exponent. The fission industry is considering lowering costs by dividing construction activities into nuclear-grade and conventional-grade categories. In our base case we assume a combination of nuclear- and conventional-grade

construction with the same split as for a future fission plant. It has been suggested, however, that an ICF plant may have less nuclear-grade construction than does a fission plant. Therefore, to set an upper limit on possible cost reductions, for our second case we assumed all-conventional-grade construction. The base-case parameters for Eq. (40) are $C_t = \$0.66$ billion and $a = 0.5$, while the parameters for the all-conventional construction are $C_t = \$0.53$ billion and $a = 0.8$.

The direct capital cost of the driver scales with beam energy and pulse rate. In our model the direct capital cost is given by

$$C_{dd} = C_d E_d^b \omega^c, \quad (41)$$

where C_d is a driver cost coefficient (\$ billion), E_d is the driver energy (MJ), b is an energy scaling exponent, ω is the pulse rate (Hz), and c is a pulse-rate scaling exponent. The base-case parameters are $C_d = \$0.1$ billion, $b = 0.7$, and $c = 0$. The cost coefficient is a goal for advanced short-wavelength lasers, and the energy scaling exponent is in the range of values reported in Refs. 92 and 93. The pulse-rate scaling exponent is quite uncertain, and we have set it equal to 0 for our base case.

The direct capital cost of the target factory is assumed to be a constant \$0.1 billion. Target-factory costs have not been studied definitively; this assumed cost is simply a rough estimate based on analogies to current facilities that mass-produce precision products (e.g., semiconductor microchips).

Thermal power, which appears in Eq. (40), can be related to the net electric power and other system parameters by

$$P_t = \frac{P_n \eta GM}{(1 - f_a) \eta GM \epsilon - 1}, \quad (42)$$

where P_n is the net electric power, η is the driver efficiency, G is the target gain, M is the fusion energy multiplication factor that results from neutron reactions, f_a is the fraction of the gross electric power required for auxiliary power within the plant, and ϵ is the thermal-to-electric conversion efficiency of the plant. These parameters can be inserted into Eq. (40) to give the reactor cost in terms of the target gain when all other parameters are fixed.

The pulse rate, which appears in Eq. (41), can be written in terms of E_d and G because the pulse rate is related to the thermal power by

$$\omega = \frac{1000 P_t}{E_d GM}. \quad (43)$$

(The factor of 1000 is needed because we express E_d in MJ and P_t in GW.)

The final item needed to find the COE is a relation between target gain and driver energy. These gain curves are very uncertain, of course, since high gain has not yet been achieved. Complex target implosion and burn calculations are done to estimate the shape of these curves for a variety of target concepts. The most recent target-gain curves from the target physics group at LLNL have been fitted by the simple expressions in our model. Moving from the lowest gain curve to the higher curves represents a progression from simple to technologically advanced target designs. We refer to this process as increasing target technology.

We can draw several inferences by combining the cost, power balance, and gain relations. Our reference case parameters are $P_n = 1.0$ GW, $f_a = 0.017$, $\eta = 10\%$, $M = 1.11$, and $\epsilon = 54\%$. The net power is standard for baseload plants, the driver efficiency is a reasonable goal for advanced short-wavelength lasers, and the other parameters are based on the Cascade design. Note that the electric conversion efficiency is much higher than for other fusion reactor designs (see Ref. 94 for details).

Exercising the Model. It is instructive to compare the target gain projected from the physics with the gain required for economic competitiveness, the gain required to meet certain reactor-design constraints, and the gain required for both economic and design considerations. Such comparisons, which can tell us which improvements will produce the greatest benefits, are made by overlaying three curves of the types shown in Fig. 8-27. The advanced-concepts curve, which is based on target physics calculations, describes the projected target gain as a function of driver energy.

With our economics model we derived an expression for the COE as a function of driver energy and target gain alone; for the

other variables we used the base-case values. From this expression we can find the gain required to obtain a given COE as a function of driver energy; this gain is shown on the $4\text{¢/kW}_e \cdot \text{h}$ curve in Fig. 8-27.

Finally, one reactor-design constraint is the pulse rate achievable in the reaction chamber. The higher the desired pulse rate, the more difficult it is to design a chamber that will recondense the vaporized material in time for the next pulse. From Eqs. (42) and (43), we can derive an expression for the gain required to obtain the desired net power given an achievable chamber pulse rate. The $\omega = 5 \text{ Hz}$ curve in Fig. 8-27 is the result of this calculation. The shaded area bounded by the three curves defines a possible design and operating space. That is, if gains up to the advanced concepts curve and chamber pulse rates of 5 Hz prove to be achievable, then the cost of electricity will be below $4\text{¢/kW}_e \cdot \text{h}$.

We chose to plot the COE curve for $4\text{¢/kW}_e \cdot \text{h}$ because the COEs for future fission and coal plants are $3.5\text{¢/kW}_e \cdot \text{h}$ and $4.4\text{¢/kW}_e \cdot \text{h}$, respectively. These values were determined on the basis of capital, fuel cycle, and operating and maintenance costs given in Ref. 95; the same fixed charge rate (8.3%) and availability factor (70%) were used for all three plant types. [The COEs actually reported in Ref. 95 are somewhat higher, since they assume a lower availability factor (65%) and include escalation in excess of general inflation during construction for the nuclear plant.]

The shape of the constant-COE curves is instructive. As driver energy increases, so does driver cost. To keep the COE constant in this situation, it is necessary to increase gain so that the recirculating power will decrease. Just as in Fig. 8-27, at some point the necessary gain goes to infinity. This behavior is the opposite of that exhibited in the target-gain curves—curves that, for each target design type, are relatively flat at high driver energy and drop rapidly to zero below some threshold.

Each curve in Fig. 8-27, in fact, is only one of a family of curves that make up a grid in target-gain/driver-energy space. Figure 8-28 shows three target-physics curves (each representing a greater achievement in terms of target technology), three constant-COE curves that

bracket the future fission and coal COEs, and three chamber pulse-rate curves. It is interesting to note that, for any given target technology, the minimum COE will be achieved where that target-gain curve is just tangent to a constant-COE curve. If the tangent point lies beyond the chamber pulse-rate limit, then the minimum COE will occur at the pulse-rate limit.

Examination of Fig. 8-28 leads to several conclusions:

- Gain improvements at low driver energy are more beneficial than those at high driver energy. Unfortunately, the rapidly declining gain curves at low driver energy mean that gain improvements at low driver energy are more difficult to achieve.
- The COE/target-technology tangent point moves to the left as we move to the higher target-technology curves. That is, the optimal driver energy decreases with improving target technology. This result, of course, is highly dependent on the shape of the gain curves, which is very uncertain. Once high gain is achieved, it will be important to determine the actual shape of the gain curves at low driver energies.
- The higher the chamber pulse rate that can be achieved, the lower the gain required to obtain a given COE. On the other hand, for any given target technology (e.g., for the advanced concepts in

Fig. 8-27. Possible design and operating space for a 1-GW_e ICF power plant competitive with future fission and coal plants.

Fig. 8-28. The families of COE curves, target design curves, and pulse-rate curves identify beneficial technological improvements.

Fig. 8-28), the target gain and COE curves are almost parallel between 5 and 10 Hz. This implies that the COE does not change rapidly as a function of chamber pulse rate in that interval; i.e., we are near the minimum anyway. Therefore, we conclude that pulse rates between 5 and 10 Hz are sufficient if target gain curves are similar in shape to those we show here. It should be noted, however, that if the gain curves turn out to be different from these, other rates may be more desirable.

● Finally, Fig. 8-28 tells us that if pulse rates of 5 Hz or more are achievable, gains of 70 to 150 are sufficient to make ICF plants competitive with future fission and coal plants.

In Fig. 8-29 we show the effect of decreasing the cost of the balance of plant. The two constant-COE curves are for the base case and for the case for which

construction is all conventional grade. In the latter, both the cost coefficient and the scaling exponent in Eq. (40) were changed. Reducing the coefficient flattens the curve; increasing the scaling exponent steepens the curve. Clearly, in this case, the 20% decrease in C_r in Eq. (40) far outweighs the 60% increase in α . This again points out the relative importance of reducing capital costs in increasing the available design and operating space.

We examined several other parameter variations in the target-gain/driver-energy space. We also constructed several cross plots of the data to develop the sensitivity study summarized in Table 8-19. An examination of the table supports the conclusions reached earlier and leads to several new ones: In particular, it appears that decreasing the plant net power makes economic competitiveness more difficult. It is

most important to reduce the capital cost. A final inference we draw from the table is that the sensitivity of the COE to variations in cost and performance parameters decreases with increasing target technology.

Conclusions. From studies of this simple model, we conclude that an ICF power plant can be economically competitive with future fission and coal plants. The following inferences were drawn from the analysis:

1. Gain improvements at low driver energy are more valuable. Target-physics calculations, however, show that high gain at low driver energy is also more difficult to achieve.

2. The optimal driver energy and the sensitivity of the COE to variations in cost and performance parameters both decrease with increasing target technology.

3. Raising the achievable chamber pulse rate lowers the gain required for a given COE. The COE, however, is relatively

Fig. 8-29. Reducing the balance-of-plant costs increases the available operating space greatly.

Parameter	Cost of electricity (¢/kW _e h)		
	Advanced concepts	Upper edge, best estimate	Lower edge, best estimate
Base case	3.27	3.67	4.55
Reactor scaling coefficient			
[0.49]			
0.70	3.37 (+3%)	3.82 (+4%)	4.83 (+6%)
0.90	3.48 (+6%)	3.97 (+8%)	5.11 (+12%)
Conversion efficiency			
[0.54]			
0.35 (-35%)	4.00 (+22%)	4.53 (+23%)	5.83 (+28%)
0.45 (-17%)	3.56 (+9%)	4.01 (+9%)	5.04 (+11%)
Net power (GW _e)			
[1.0]			
0.5 (-50%)	4.94 (+51%)	5.61 (+53%)	7.09 (+56%)
1.5 (+50%)	2.59 (-21%)	2.89 (-21%)	3.56 (-22%)
Chamber pulse rate (Hz)			
[5]			
2	3.47 (+6%)	3.97 (+8%)	5.08 (+12%)
10	3.18 (-3%)	3.56 (-3%)	4.39 (-4%)
Driver efficiency			
[0.10]			
0.05 (-50%)	3.39 (+4%)	3.95 (+8%)	5.34 (+17%)
0.15 (+50%)	3.23 (-1%)	3.59 (-2%)	4.34 (-5%)
Driver cost coefficient (\$ billion)			
[0.10]			
0.20 (+100%)	3.75 (+15%)	4.46 (+22%)	6.00 (+32%)
0.15 (+50%)	3.51 (+7%)	4.11 (+12%)	5.28 (+16%)
Driver energy exponent			
[0.7]			
0.9	3.32 (+2%)	3.89 (+6%)	5.30 (+16%)
0.5	3.22 (+2%)	3.50 (-5%)	4.07 (-11%)
Driver pulse rate exponent			
[0]			
0.2	3.45 (+6%)	3.97 (+8%)	5.10 (+12%)
0.1	3.35 (+2%)	3.81 (+4%)	4.81 (+6%)

Table 8-19. Changes in the COE resulting from changes in the base case parameters for three gain curves shown in Fig. 8-28. Base-case parameters are given in square brackets. Fractional changes (%) from the base case are given in parentheses.

sensitive to pulse rate between 5 and 10 Hz. If these pulse rates are achievable, then gains of 80 to 100 will be sufficient for ICF plants to be competitive.

4. Lowering the cost of the balance of plant has the greatest leverage in reducing the relative COE.

5. As the plant size (measured in terms of the net power distributed) decreases, relative economic competitiveness decreases.

Several of these conclusions are very dependent on the shape of the gain curves. Achieving high gain initially is more likely at high driver energy. Once high gain is achieved, it will be crucial to determine the true shape of the gain curves.

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Physics and Systems Issues for the Starlite Excimer Direct-Conversion Concept

Starlite is a recently proposed⁹⁶ concept for a reactor that will directly generate electricity with high efficiency. In this design, the reactor is filled with high-pressure helium gas everywhere except for a region that is filled with low-density gas through which driver beams propagate to the target. When the targets ignite, the helium absorbs the fusion neutrons and converts their energy into excimer states of the diatomic molecule. These states then decay and emit 15-eV photons, which can be intercepted and directly converted to electricity with positive/negative (PN) junctions around the walls of the reactor chamber. In this article we examine this proposal and discuss unresolved physics and systems issues, including the necessity for PN junctions that can withstand high current transients and for a containment vessel for the excimer fluid that can withstand damage from x rays and debris.

Excimer Concept. The radiative and kinetic mechanisms of excimer noble gases are familiar to scientists studying advanced laser fusion.⁹⁷ Figure 8-30 displays the interactions necessary to produce the excimer and its decay photons, where X is an atom of any noble gas.

In the first process, an ion with several MeV of kinetic energy forms when the energetic neutron undergoes elastic scattering. In the second process, the fast ion interacts with the background gas through Coulomb scattering to produce low-kinetic-energy (thermal) ions and a manifold of excited states. With a sufficient column density of background gas, the combined efficiency for these two processes usually exceeds 90%. In the third process, the slow ions rapidly form molecular ions at high pressures (X_2^+), and dissociative recombination of the molecular ions feeds into the atomic excited-state manifolds (X^*). The excited atomic states feed the upper excimer levels (X_2^*), which radiate to the repulsive ground state, yielding a photon. This photon, which is about 15 eV for helium, is not radiatively trapped because an excimer exists only in the excited state—there is no ground state. Along the excimer chain, other interactions may occur: e.g., the diatomic ion can form clusters ($X_2^+ + 2X \rightarrow X_3^+ + X$), and the excited atom can form an ion through electron collision ($X^* + e^- \rightarrow X^+ + 2e^-$) or Penning ionization ($2X^* \rightarrow X^+ + X + e^-$). Nevertheless, these alternative interactions produce species that reenter the excimer chain at other locations. The third process converts the deposited Coulomb energy at low excitation levels into subsequent excimer radiation at an efficiency between 60 and 80% (the theoretical upper bound). For doubly charged helium, the efficiency has been measured to be approximately 70%.⁹⁸

Physics Issues. Photodiodes cannot collect the excimer photons because they are limited by space charge. Photovoltaic PN junctions might be used, but their high-current capability at sub-microsecond time scales is unknown. That is, an inherent problem in all direct-conversion ICF schemes is the change in time scales from fusion burn, at tens of picoseconds, to the steady state of the required electricity. Because the excimer chain slows the time scale only to the sub-microsecond level, the junctions must be able to withstand extremely high current transients if PN junctions and capacitors are used to complete the time-scale conversion.

Fig. 8-30. Starlite excimer chain processes.

First process:

Second process:

Third process:

Future experiments must prove that such capability exists.

Systems Issues. First note that the helium in Starlite must be a gas because liquid helium has a very low heat of vaporization: A sphere of liquid helium 4 m in radius would be entirely vaporized in a few fusion pulses, and cryogenic refrigeration would be expensive and would consume a great amount of power. Unfortunately, however, extreme pressures are required to obtain enough column density to stop 14-MeV neutrons when the gas is at operating temperatures. For example, a pressure of 850 atm is required for a 4-m-radius system operating anywhere near 1000 K. Such pressures necessitate containment vessels with walls at least 0.5 m thick. In addition, the ionization stage,

which forms the input to the excimer chain of Fig. 8-30, has a formation efficiency that is highly sensitive to contaminants. Any collisional losses to impurities will circumvent the excimer chain processes. Helium must therefore be kept uncontaminated by target debris or reactor-wall ablation products. The necessity for both high-pressure gas and uncontaminated helium requires that the helium be separated from other reactor systems by a containment wall, which will be subject to serious erosion from target x rays and debris. We do not yet know how to contain the gas without continually replacing this wall. In addition, we have yet to investigate ways to add tritium breeding to the excimer reactor systems.

The maximum possible efficiency of the excimer concept is between 30 and 60%. This efficiency is the product of a measured excimer formation efficiency of 70%, an assumed photon collection efficiency of 100% (which is probably unattainable), a photovoltaic efficiency of 40 to 80%, and an assumed short-time-scale circuit efficiency of 100%. The maximum efficiency of 30 to 60% should be compared with the 55% efficiency thought to be attainable in the Cascade reactor concept.⁹⁹ The excimer concept is appealing because high electrical efficiency is possible with no moving parts, so that a lower COE is possible if the capital cost of the reactor can be lower than the cost for schemes using thermal cycles. In addition, the residual thermal heat in the gas can be removed with a thermal bottoming cycle, thereby allowing even higher efficiencies and perhaps even lower COE. The excimer concept thus has a potential for lowering electrical costs, but it imposes formidable mechanical constraints on the containment vessel for the excimer fluid.

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Developing ICF Facilities for Multiple Uses

Development facilities required for ICF must be clearly affordable if ICF is to be competitive. While costs are being reduced, we are considering a collateral route: make ICF development facilities attractive for multiple purposes, so that

several sponsors can split the costs. With ICF's many present and future applications, it is obvious that this strategy will be adopted; the question is how many such applications will be accommodated. In this article we

- Outline some possible applications of ICF technology.
- Sketch development paths for each.
- Examine the implications for design and operation of the ICF development facilities if multipurpose use is encouraged.

Our conclusion is that, although there are risks, a strategy for varied multiple uses will result in more stable funding for ICF facilities.

ICF Applications and their Development Requirements. Military applications of ICF include production of special nuclear materials (SNM) for weapons and studies of nuclear weapons physics, vulnerability, and lethality. Civilian applications include use of x-ray lasers for diagnosis, production of fissile fuel and electric power, and space propulsion. Although details of military applications of ICF technology are classified, general principles are discussed here in broad terms.

Fusion Power. The various proposals for development steps and facilities for commercial production of fusion power have similar characteristics (see Fig. 8-31). With present ICF facilities (which cannot achieve high gain), researchers are studying driver technology, driver/target coupling, and target implosion dynamics. It is generally agreed that the next step is to build a high-gain test facility (HGTf) capable of obtaining gains greater than 100. In such a facility the physics of burning targets can be studied in detail and the

Fig. 8-31. Typical commercialization scenario for an ICF power reactor.

dependence of target gain on driver energy can be fully resolved for a variety of target types. This facility needs a driver that not only can support energy and peak-power capabilities in excess of ultimate requirements for a reactor, but also one that is quite flexible in pulse duration, pulse shape, and wavelength. The facility will need to provide pulses only occasionally (a few per day) and may be needed for only about 10^4 shots.

The next development step is assumed to be an engineering test facility (ETF) in which several potential reaction-chamber concepts can be tested to identify how best to protect the chamber walls and clear the chamber. The ETF has been proposed as an add-on to the HGTF that uses the same driver.¹⁰⁰ In this case, it would be desirable to have the capability to pulse the driver several times in a short burst even if the driver were not fully rep-rated for continuous operation; that is, not capable of producing several laser pulses per second on a continuous basis. For example, perhaps a 10-MJ driver could also deliver five 2-MJ pulses in rapid succession (assuming that, in our HGTF experiments, we learn how to make a capsule burn with only 2 MJ). Alternatively, the ETF could be a new facility or a facility added to a rep-rated driver developed independently of the HGTF. In any case, the ETF would be the first facility with multi-pulse capability, but it would not need to maintain the pulse rate for long times. Driver requirements for the ETF are similar to those for the HGTF except that bursts of several or a few tens of shots are desirable, and the driver for the ETF should last longer than a HGTF driver by about an order of magnitude.

An experimental power reactor (EPR) would be the next facility needed. In the EPR, the best rep-rated driver would be integrated with the most desirable reaction chamber concept for thorough engineering and systems testing. It is possible that the ETF could evolve into the EPR.

Finally, a demonstration power reactor would be needed to prove the concept sufficiently reliable and economical for industry to begin construction. These last two reactors require fully qualified drivers in terms of pulse rate, but do not need the

pulse-shape flexibility of the HGTF and the ETF.

Nuclear Weapons Physics. When a nuclear weapon explodes, dense, highly ionized, hot plasmas form. As the materials move, instabilities at the interfaces become important, and the materials may not be in local thermodynamic equilibrium. The equations of state and opacities of such materials are difficult to calculate, and the conditions are difficult to reproduce in the laboratory. These phenomena can be thoroughly studied only in underground nuclear tests, which are expensive and infrequent.

Nuclear weapons physics problems can now be simulated in the laboratory with an exploding fusion capsule or with the laser driver alone. Novette was the first laser sufficiently powerful and flexible to produce conditions of interest to weapons designers. For example, when a 450-ps pulse was focused onto a thin selenium foil 200 μm wide by 1 cm long, Novette's power density was $5 \times 10^{13} \text{ W/cm}^2$ with 0.53- μm light. This power density produced a plasma with an electron density of $3 \times 10^{20} \text{ cm}^{-3}$ and a temperature high enough to strip 24 electrons from each selenium atom. With each increase in power density, additional regimes of weapons physics become accessible. When fusion capsules achieve high gain, they will provide energy densities greater than we presently achieve by at least a factor of 100, and their emissions can then be used to create rare conditions in matter.

Weapons physics experiments are already an important element in ICF facilities, and their impact on design and operation of those facilities is being felt in added requirements for diagnostic capabilities and target irradiation geometries. Expanded ranges are needed for pulse length, pulse shape, and wavelength flexibility.

Vulnerability, Lethality, and Effects. Studies of the interactions between the outputs of a weapon and either the environment or other military systems are called vulnerability, lethality, and effects (VLE) studies. The designer must determine the vulnerability of a system to an opponent's weapon as well as the lethality of weapons when used against the opponent's systems. The ICF facilities are

suitable for VLE studies for nuclear and laser weapons.

When a nuclear weapon explodes, it emits energy in the form of neutrons, x rays, γ rays, and debris. The weapon output may interact with the environment and create effects such as blast or electromagnetic pulse (EMP). The possible physical interactions among an exploding nuclear warhead, the environment, and another military system are numerous, and we will discuss only a few examples relevant to ICF in this article.

Neutrons from the explosion of a warhead can cause fissions and heat buildup in another warhead. X rays can create shock impulse in the skin of a reentry vehicle. X rays and γ rays can cause photoelectrons or Compton electrons to be emitted from the circuit elements inside a system. The rebounding electrons can create large pulses of electromagnetic energy in those circuits. Radiation can cause loss of bits of information in integrated circuits.

To study these effects, VLE researchers have built laboratory simulators and have conducted underground tests with nuclear explosive sources. To be realistic, a simulation must re-create the energy partition, output spectra, time scale, environmental effects, and interactions among all phenomena. Not surprisingly, the present simulation methods are inadequate to cover all cases of interest. Even underground nuclear tests do not have complete fidelity.

ICF technology could significantly increase this country's understanding of nuclear weapons effects only if high gain is achieved. Existing simulators are better suited for this purpose than are existing ICF drivers. Even small thermonuclear ICF yields (gain 1 to 10) will not exceed current simulator capability except in very limited cases; thermonuclear yields of more than a few hundred megajoules will be necessary to make ICF attractive as a simulator. In many cases an entire structure must be present to create the effect, and this capability exists only at high ICF yields.

A HGTF will be fully exploited only if there is sufficient space in the target chamber to expose large pieces of hardware. Because there is interest in a variety of output spectra and energy partitions, it will be necessary to alter the output of the fusion capsules. The explosion time scale

of an ICF capsule must be as long as possible to better simulate a nuclear weapon. It is also possible that the target chamber will need to handle fissionable material.

Some beam-weapons experiments can be done with current ICF drivers. Laser propagation studies are the most obvious, and—although the characteristics of ICF lasers and laser weapons are different—some VLE experiments are possible. To do useful VLE research on directed-energy weapons such as the x-ray laser, however, will require the HGTF.

X-Ray Laser Applications. The peak output brightness obtained in the laboratory x-ray laser experiments to date is 6×10^{10} W/cm²·Å·sr. These x-ray sources are 10^{10} times brighter than other x-ray generators (e.g., synchrotron sources). They are, of course, one-shot, short-pulse devices that must be developed substantially before they become useful laboratory x-ray sources. The potential exists, however, to make the x-ray laser an important tool.

It has been proposed to do x-ray holography on biological organisms. The spatial resolution would be comparable to that from an electron microscope (50 to 2000 Å). However, x-ray holography would have the advantage of exposing living tissues, thereby avoiding alteration of structures while preparing samples for microscopy. Suggested experiments include studies of the cell membrane structure, nerve fiber interconnections, and the morphological differences between normal and malignant cells.

Some simple (nonholographic) experiments are possible with the present x-ray laser capability. However, more experiments will become possible as the driver power increases and x-ray wavelength decreases. The high-brightness source mentioned above is at a wavelength of about 200 Å, and some lasing was detected at about 150 Å in yttrium. To do microholography, a bright source at perhaps 30 Å would be needed. Cylindrical focusing, high peak power, and pulse flexibility are very important laser characteristics for this application.

Production of Nuclear Materials. The neutrons in a fusion reactor can produce valuable nuclear materials. Fusion has a natural advantage over fission for this application because many more excess

neutrons are available per unit thermal power. The advantage is a factor of 9 or 10 compared with a fission reactor that obtains fuel from another source and a factor of 30 to 60 compared with a breeder reactor. In one fusion reactor design that uses a beryllium neutron multiplier and flowing lithium to produce tritium, a tritium-breeding ratio of 2.1 was calculated with a thermal power of 530 MW. This design could produce excess tritium (beyond that needed to fuel the reactor itself) at a cost two to three times lower than with a new fission reactor. Such a reactor could produce tritium and plutonium for the military, ^{235}U for fission reactor fuel, or both.

Materials production will come to fruition only when a high-gain pulse rate over 1 Hz becomes available. To develop this application will require a full-scale engineering-development program. Limited integrated neutronics tests and chamber phenomenology experiments will be possible on the HGTF. These first experiments would be similar to those for power reactor development, but would use different materials. If production testing begins on the ETF, the facility will need more extensive capabilities for processing radioactive materials than are needed for power reactor development alone.

Space Propulsion. In this application, fusion capsules would be injected into a thrust chamber at the rear of a space vehicle and burned at the rate of tens or even hundreds of pulses per second. Magnetic fields would divert the charged debris and produce thrust in the desired direction. While the propulsion system for such a vehicle would be very large (perhaps 500 tons), it would have unique properties: Specific impulses of at least 10^5 s and specific powers of 100 W/g are achievable. With such a rocket, payloads of 1500 tons could be moved between Earth and Mars in 23 days and between Earth and some nearby stars in less than 100 years.

The development path for such a capability is very long. Most proposals use advanced fuels such as $\text{D}-^3\text{He}$ or DD to increase the debris yield relative to the neutron yield. However, some thought is being given to developing a DT-fueled rocket for testing in the near future. If such a system appears attractive, then some development steps could be taken on the next few ICF

development facilities. Obviously, it would be desirable on the HGTF to test DT targets in which the amount of debris yield is maximized. Different methods of converting the capsule yield to thrust could also be tried on both the HGTF and the ETF. Considerable driver flexibility would be desirable for these different and less well-defined targets.

The Impact on ICF of Sharing Facilities. Each new ICF facility will be suitable for one or more of the applications described above, depending on whether the ICF facility is designed to be sufficiently flexible. Careful planning will be needed to fit the parts of each application's unique development path to the most suitable ICF research facilities.

Even today, facilities are shared between ICF technology development and weapons physics experiments. Thus the HGTF, the ETF, and perhaps the EPR should all be able to accommodate experiments for some of the applications discussed. On the other hand, at some point (e.g., a demonstration power reactor) the facility would be so narrowly designed that the usefulness to other applications would be limited.

Figure 8-32 shows the ICF development facilities and the experiments that can be done for each application. For example, in addition to ICF high-gain experiments, the HGTF could handle expanded weapons physics and x-ray laser applications experiments, VLE exposure and phenomenology tests, first-wall exposure tests for both material and power reactors, and target design and the first thrust-concept experiments for propulsion. It is clear that to accommodate so many different types of experiments, the HGTF should have multiple target chambers, as shown schematically in Fig. 8-33.

The use of multiple target chambers is not new; Nova now supports 2 chambers, with either 10 beams sent to the main chamber or 2 beams sent to the smaller chamber. Experiments in different chambers are done on successive pulses, which may be hours apart.

Several advantages accrue from the use of multiple chambers. The major time spent on an experiment is in setting up and taking down apparatus in or near the target chamber. Often, necessary laser

Fig. 8-32. Applications of ICF technology (other than power production) that could do experiments on multipurpose ICF development facilities.

	Present ICF facilities	HGTF	ETF
Weapon physics	Substantial use of first high-energy-density experiments	Extensive use of very high-energy-density experiments	N/A
Vulnerability, lethality, and effects	N/A	Source development, phenomenological studies	Multiple exposures
X-ray laser applications	Develop x-ray laser sources, first demonstration experiments	Substantial use of standard user sources	N/A
Nuclear materials production (military and civilian)	N/A	Neutronics experiments, chamber phenomena	First materials production demonstrations
Space propulsion	N/A	Maximization of target debris yield	Comparison of thrust chamber concepts

Fig. 8-33. The use of multiple target chambers will facilitate experiments for many applications on ICF development facilities such as the HGTF and the ETF.

configuration changes can be accomplished much faster. It is possible for different groups to work in parallel in separate areas, and several types of experiments can be handled during each run. Each applications group is able to do its work on its own schedule. Since the target chambers are inexpensive relative to the laser and the costs of the driver construction and operations are shared, total

costs for each application are lower. Some applications require large exposure areas, while others use materials that will contaminate a target chamber, requiring considerable clean-up time. Using separate chambers mitigates these problems. If the chambers can be adequately isolated from one another, it is even possible that classified and unclassified experiments could share the same driver but exist in separate

buildings. This suggests that there should be chambers for basic ICF studies, weapons physics, VLE studies, and perhaps industrial applications or reactor studies. Priorities could be established easily for each chamber because delays caused by other applications needs would be minimal. If adequate space is allowed and the driver characteristics are flexible enough, the chambers could be built sequentially, as the need arises.

Including the capability to support multiple chambers has some implications on the appropriateness of ICF technologies under consideration. Laser and heavy-ion drivers of some types could easily support multiple chambers. However, it would be difficult for light ion, free electron laser, and compact torus drivers to support multiple chambers. In addition, supporting multiple chambers would be considerably easier with indirect-drive targets, which use fewer beams, than with direct-drive targets. However, the issue of target classification would have to be addressed if one experimental area was to be unclassified and the others were classified.

The second major issue in multiple use of ICF development facilities is the flexibility of the driver to provide pulses appropriate to all applications. In Table 8-20 we compare the driver requirements for power reactors and those for the HGTF, assuming that multiple applications are contemplated. It is clear that a HGTF driver would be considerably more flexible than the reactor driver, but not necessarily capable of high average power. It would be desirable if the driver technology selected for the HGTF could be extrapolated to a reactor driver. However, a sufficiently flexible design is likely to differ greatly from a reactor driver. Furthermore, while the HGTF is being operated, research on high-average-power drivers will help to define the best driver for any particular reactor application. It is sufficient that a possible development path to a reactor driver is foreseen for the chosen HGTF driver. It is more important to emphasize performance parameters that can assure high gain and the flexibility necessary to support multiple applications.

	Research driver	Reactor driver
Pulse energy	5-10 MJ	3-5 MJ
Peak power	500-1500 TW	~500 TW
Pulse duration	0.1-100 ns	5-10 ns
Pulse shape	Flexible	Fixed, complex
Pulse rate	Few per day	5-10 Hz
Efficiency	Little importance	≥ 10%
Wavelength	Flexible, 250-1000 nm	Fixed, < 500 nm
Focusability	Flexible geometry	Fixed, 1-mm spot at several meters
Reliability	Unknown	> 70% availability
Durability	10 ³ -10 ⁴ shot life	~10 ¹⁰ shot life
Maintainability	Skilled, controlled	Routine, hands-on
Cost	25-50 \$/J	100-150 \$/J

A final factor to be considered is the effect of multipurpose development on support facilities:

- A variety of target types would need to be fabricated.

- Diagnostics support would have to be broader, although individual researchers could probably supply some of the diagnostics equipment for their experiments.

- Data would probably be collected separately for each target chamber, with only laser pulse characteristics collected in common and distributed to the appropriate area.

- A multipurpose ETF would require an expanded chemical-processing capability that could handle all the types of materials anticipated.

The operating costs for these expanded support facilities would be shared by the users. The impact on the ETF would be similar to the that on the HGTF: Multiple chambers, driver flexibility, and shared support facilities would still be important. It is likely, though, that when an EPR is built, driver flexibility will be a lower priority; by that time, the driver best suited for the primary mission will be known.

Conclusions. Because ICF already has multiple goals, it is clear that its development will include building facilities to accommodate a number of applications. The only question is the number of such applications that will be accommodated.

Proceeding along such a development path is advantageous for ICF because a number of groups will have an interest in seeing each ICF facility built. Each group may perceive its share of the cost as affordable, even if the total could not be supported by any one application. Such a

Table 8-20. Driver requirements for research and development facilities such as HGTF and ETF emphasize flexibility and broad parameter ranges, whereas those for a commercial reactor can be less flexible but need high average power and high durability and reliability.

development strategy may provide more stable support by being less sensitive to changes in national priorities (although it will be as sensitive as any other strategy to across-the-board cuts). It should be recognized that there are also risks to such a strategy. During times of general budget squeeze, each user group may cut back more than its share in the hope that the others will take up the slack.

The success of such a development strategy for ICF will require careful

planning. Before each new facility is approved, the needs of all possible user groups will have to be considered. It is clear that driver flexibility, use of multiple chambers, and shared support facilities must be included in the planning for the HGTF and the ETF. Done properly, the multipurpose facility strategy can accelerate ICF development.

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Measured optical fringes produced by silicate glass laser slab.

Gas-cooled-slab laser test bed.

Solid state lasers at a broad range of energy have been the focus of considerable development during the past few decades. Large laser systems (Shiva and Nova), which are used to study ICF, store large amounts of energy, which they deliver in single pulses of very high peak power. Conversely, low-power glass lasers are widely employed in many industrial, military, and medical applications. However, technical limitations on the mid-range power and efficiency achievable with solid state lasers have prevented them from being used in moderate-to-high power applications such as high-power welding and cutting, communications, and radar.

The High-Average-Power Laser Program's goals are to develop a broad technology base for solid state lasers and to demonstrate operation of these devices with higher power, repetition rate, efficiency, and beam quality than is currently possible. A well-integrated and well-balanced research and development effort spans the disciplines of laser system design, quantum electronics, materials science, and the engineering sciences. Major activities are the zig-zag laser testbed and the gas-cooled-slab laser test bed. The photo below shows a silicate glass laser slab before insertion into the test-bed structure. Photos at lower right and on the opposite page, top, show the measured optical fringes produced by this slab while operating in the test bed. The opposite-page photo, bottom, shows the gas-cooled-slab laser test bed.

9

High-Average-Power Lasers

High-Average-Power Lasers

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High-Average-Power Lasers

Introduction

W. L. Gagnon and J. F. Holzrichter

The goals of the High-Average-Power Laser Program at LLNL are to develop a broad technology base for solid state lasers and to demonstrate high-average-power laser operation with more efficiency and higher beam quality than has been possible with current technology. We are placing a dual emphasis on a broad science base and on technology advancement. Our objectives are oriented toward general military applications rather than to specific missions or platforms.

Solid state lasers have been employed in a wide variety of low-power industrial, military, and medical applications: x-ray lithography, laser fusion, laser isotope separation, and materials fabrication and processing. However, technical limitations have prevented solid state lasers from being used or considered in moderate-to-high-power applications; gas lasers were used instead. Solid state lasers could be considered for these applications if restrictions on average power could be removed and if their efficiency were improved because they have intrinsic features that make them attractive candidates for many demanding applications.

Low-power neodymium lasers have already been used in some military applications, most notably in range finders and designators. However, medium- and high-power applications (communications, countermeasures, radar, etc.) are for the future and are paced by the availability of the developing solid state laser technology. Unbounding the average power and increasing the efficiency of solid state lasers will put them in contention for these applications. Compactness for compatibility with military vehicles is an additional requirement, and at medium-to-high power, all laser concepts have been deficient in this respect.

Over the past decade, numerous new design concepts, materials, and techniques relevant to solid state lasers have emerged that, collectively, suggest that the traditional technical limitations on power (a few hundred watts or less) and efficiency (less than 1%) can be removed. The basic characteristics of solid state lasers make them highly attractive for high-average-power applications—they are operationally benign, logistically unobtrusive, and potentially very flexible in their wavelength and pulse width. Therefore, the Defense Advanced Research Projects Agency (DARPA), LLNL, and others¹⁻¹¹ recognized that an aggressive development program could greatly expand the utility of solid state lasers. Such a development program would boost solid state laser technology to a level where the following performance goals could be selectively demonstrated:

Power	Tens to thousands of watts per plate
Beam quality	$\lambda/10$ rms (Strehl ratio of 0.6)
Storage efficiency	>10%
Extraction	>60%
Lifetime	10^8 to 10^{10} shots
Wavelength	Many

Summary of Accomplishments

Development goals	Fiscal year 85	Fiscal year 86	Fiscal year 87	Fiscal year 88
Laser material				
Glass	Evaluate	Cast plates	Low loss	Composite
GSGG	Small-slab growth	Large-slab growth	Working slabs	
Test beds				
Zig-zag geometry	Construct	Experiments	Low PRF laser	Power PRF Efficiency Beam quality Multi-wavelength
Gas-cooled slab geometry	Facility	Experiment design	Measurements	
Nonlinear optics				
Q-switches	Evaluate	Low PRF	High PRF	
Frequency conversion	Evaluate materials	Develop materials	Low PRF	
Design codes	1-D	2-D	3-D	Anisotropic 3-D
Application designs	Free-electron-laser oscillator	Laser radar countermeasures Submarine laser communication		

Table 9-1. DARPA/LLNL advanced laser development schedule.

The full realization of the potential of solid state lasers is being achieved through a broadly based, well-integrated, and well-balanced research and development effort that spans the disciplines of laser system design, quantum electronics, materials science, and the engineering sciences (thermomechanics, optics, and pulsed power). The program schedule shown in Table 9-1 represents our broad-based approach to the problems involved. In the schedule, we show the technologies addressed and the fiscal years in which they will be undertaken.

Summary of Accomplishments

Experimental Test Beds

The two key factors in high-average-power operation are the control of thermal gradients in the laser medium and a laser configuration in which thermal gradients do not generate phase distortions. For both of these reasons, plate geometries of the laser medium, as contrasted with traditional rod geometries, are required. Two plate geometries appear to be useful: the water-cooled zig-zag geometry and the gas-cooled slab geometry. Figure 9-1 shows that in these plate-like gain configurations, phase and gain profiles are uniformly sampled by the beams. Further, these geometries make possible large surface areas for cooling and assure that deposited heat is relatively close to the cooling surfaces.

To test the validity of these concepts, we assembled test beds and support facilities for both geometries. In these test beds, we can perform rapid and flexible experiments that check and validate the basic design principles and the thermal/

optical design codes. Using the water-cooled zig-zag test bed, we characterize the system cooling requirements, investigate the effects of pump uniformity on beam quality, design and test reflector geometries that provide uniform pump intensity, and identify and test materials suitable for bonding edge absorbers and evanescent-wave coatings to the slabs. Using the gas-cooled slab test bed, we operate and characterize the gas flow system, operate the flashlamps, and measure the surface temperature distribution of the slab during pump conditions. We characterized beam quality and thermal loading over a wide range of operating power in the test beds. Our ultimate goals are to demonstrate that both of these geometries are capable of operation with high efficiency, high output power, and high beam quality.

Materials Development

The achievement of our long-term goals will ultimately be paced by the availability of suitable materials with appropriate laser and thermal-mechanical properties and with desirable scaling properties to reach

Fig. 9-1. Two geometries appear to be useful for high-average-power operation, the zig-zag and the gas-cooled slab.

the required sizes. In the near term, we will grow materials to reasonable sizes in order to demonstrate high-average-power designs and validate our design methodology and design codes. In the longer term, we are developing crystal-like laser media with properties responsive to the above described applications. In addition, we will develop a basic understanding of growth scaling and develop scaled growth techniques.

The major thrust of this effort is to develop and certify the techniques and processes needed collectively to produce fully functional laser quality plates made of chromium- and neodymium-doped gadolinium scandium gallium garnet (Cr:Nd:GSGG) of nominal dimensions $1 \times 10 \times 20 \text{ cm}^3$. The techniques and processes are being demonstrated by performing the following tasks:

- Demonstrate growth of 13-cm-diam, core-free boules of Cr:Nd:GSGG with adequate refractive index homogeneity, stress birefringence, and optical loss.
- Acquire and refine scandium oxide for scaled crystal growth efforts.
- Support crystal growth scaling experiments by off-line small-melt studies directed to assuring control of impurity-induced $1.06\text{-}\mu\text{m}$ losses in the crystals.

These tasks are currently being carried out under contracts with Allied Corporation, ATT Bell Laboratories, and Ames Research Laboratory.

To transform laser-quality crystal boules into functional laser plates, it is necessary to extend to large crystals the edge-cladding methods that have been successful on centimeter-sized crystals and to develop adequate optical fabrication methods. These tasks are currently being carried out under contracts with Hoya Optics, Inc. of Japan and Zygo Corporation. At LLNL, we are performing all the measurements necessary to evaluate the materials produced under these contracts (optical quality, optical absorption and scattering losses, solarization effects, cladding performance, coating effectiveness, etc.) so as to maximize the combined rate of technical progress. We are also developing methods for increasing the breaking strength of large glass and crystal plates. We have demonstrated that the strength of a plate or slab is a strong function of the size of the largest flaws in the areas of high tension. These flaws originate during preparation of the optical surfaces. We have demonstrated that by developing surface finishing techniques that minimize both the number and size of these

flaws, the breaking strength can be substantially increased.

In addition to these central efforts, we have conducted exploratory research experiments on spectroscopy for new laser materials, properties characterization, and laser demonstration. The most notable development has been the identification of the promethium ion as an effective laser ion for submarine communication.

Pump Sources

We are interested in both laser diode arrays and xenon flashlamps as pump sources for high-average-power lasers. Over the past several years, laser diodes have undergone considerable development. Efficiencies of about 40% have been demonstrated in single diodes and 1×1 -cm arrays have been tested at efficiencies exceeding 20%. This work has been carried by RCA, McDonnell Douglas, and Spectra Diodes under contract to DARPA. For diode arrays to be feasible in a systems application context, the manufacturing cost of producing large, efficient arrays must be reduced by about an order of magnitude; efforts are proceeding in this direction.

In addition, we have worked to develop xenon flashlamps for high-average-power operation. We tested water-cooled lamps at power loadings on the walls up to about 400 W/cm^2 . A number of lamps lasted for several million shots and one lamp lasted for 44 million shots before failing. Failure modes are usually fractures in the lamp envelope near the cathode, where the wall temperature is highest. These tests were all carried out in a water-cooled calorimeter with provisions for measuring the radiant efficiency and spectral characteristics of the output power. In addition, under the sponsorship of the ICF program, we developed reflective dielectric coatings that, when applied on the outside surfaces of the lamp walls, increase output in the neodymium pump bands by recycling otherwise unused portions of the output spectrum back into the lamp plasma. We are continuing to characterize the behavior of flashlamps in the high-average-power regime and to develop coatings that will survive for 10^8 to 10^{10} shots.

Nonlinear Materials

Our efforts in developing nonlinear materials have been directed at assessing materials for high-average-power operation. The current state of the art in high-average-power frequency conversion is a few tens of watts. The limiting process is linear optical absorption in the crystals causing temperature gradients that reduce the conversion efficiency by disturbing the phase matching. For this reason, we carried out a careful assessment of the thermodynamics in nonlinear crystals. This analysis showed that thermal gradients must be perpendicular to the crystal faces in order to successfully operate at high average power. We determined that there are two contributions to the temperature variation in the crystals: one is the longitudinal temperature rise associated with the flow of heat outwards to the crystal surfaces; the other is the transverse temperature rise associated with the temperature rise in the coolant as it traverses the crystal face. High average power is achieved by configuring the nonlinear material as a set of thin plates with a large ratio of surface area to volume and by cooling the exposed surfaces with a flowing gas. We project that this configuration, using KD^*P , can be extended to average powers approaching 100 kW. Furthermore, the design and material fabrication of such a harmonic generator is well within current technology.

Design Methodology

A major objective of the program is to lay a sound and broad technical foundation for high-average-power solid state lasers in a three-year time frame. An essential ingredient is a thorough design methodology that addresses all the relevant issues. We completed the formalism started by Emmett, Krupke, and Trenholme¹² and now have a capability to design a high-average-power laser system for almost any given application, starting either with given materials or by using the inverse approach of letting the application define the desired material properties. We concentrated on improving laser performance by increasing efficiency, pump

uniformity, repetition rate, and beam quality, and on simplifying design methodology without sacrificing accuracy. We analyzed the pump properties of existing test-bed units in detail and used the results to find the highest efficiency operating points. We also found reflector designs with excellent uniformity and high efficiency. We studied amplified spontaneous emission (ASE) effects that reduce storage lifetime, and calculated slab aspect ratios that minimize unwanted ASE. Extraction in the presence of losses has been determined for both Q-switched and amplifier operation, and limits on permitted losses have been set. Surface and edge heating, which limit repetition rate, have been quantified. We incorporated the results of these and previous analytic efforts into a number of fast but accurate design codes that allow rapid production of slab designs to meet any system requirement.

We have, in addition, directed considerable effort toward analyzing and understanding the thermal/optical effects in laser materials. The TECATE and TOPAZ/NIKE-2D codes have been used to evaluate thermal and stress distributions in the zig-zag test-bed experiments. These results are used to analyze both thermal engineering design issues and beam wave front and polarization deviations. We have also developed a birefringent ray-tracing code, BREW, which we are using to evaluate wave front distortions due to thermally induced refractive-index changes (dn/dT) and to displaced boundaries.

System Applications

During 1985, we developed conceptual designs for two potential applications of the technology. In the first of these system designs, advanced solid state laser concepts produce 10-J, 50-ns pulses at a repetition rate of 1 kHz for a 1- μm free-electron laser driver system. This high quality pulse train (nearly diffraction limited) is injected colinearly with a relativistic electron beam into the wiggler amplifier structure of the free-electron laser where it is amplified to very high power.

In addition, we were asked by Lincoln Laboratories to study the feasibility of developing a high-average-power solid

state radar. We concluded in this study that the required performance level can probably be reached, given reasonable development of the component technologies.

Conclusions

In conclusion, we have in place a broad-based program to develop solid state laser technology. We have operating test beds for the two laser geometries of interest and are taking data from them. We have in progress an aggressive program to grow crystals at large scale and to identify and characterize both laser and nonlinear materials that show promise for high-average-power applications. We have evolved a thorough design methodology that makes possible rapid evaluation of new concepts and techniques. Finally, we have identified a number of near-term applications for the technology that are both timely and useful.

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Laser Design

Introduction

The average power of solid state lasers is limited by two thermal effects. Beam quality is degraded as the power is raised, and the laser material fractures because of thermal stress at very high power.

Thermally induced laser beam distortion arises when different rays in the beam have different optical path lengths, or when thermally induced birefringence depolarizes the beam. Most solid state lasers employ rod geometries, in which thermal effects are large. However, the laser medium can be deployed in plates, in such a way that most rays have the same thermal history and so undergo the same path change. Far from the edges of the plates, the thermoelastic effects do not depolarize the beam. Thus appropriate plate geometries can greatly reduce beam degradation. In addition, if the plates are pumped through their faces, so that all rays have the same average gain, there will be no changes in beam profile arising from gain differences. Thus laser materials in plate

form can be run at powers near the fracture limit without reducing beam quality to the point of unusability.

Thermal fracture is a result of heat buildup, which causes thermal expansion and stress buildup. The stress eventually exceeds the strength of the material. Glasses have low thermal conductivity and low strength, so heat buildup is large and fracture occurs quickly. Crystalline materials have high thermal conductivity and high strength, so heat buildup is small and fracture occurs only at very high loading. Crystalline laser hosts are thus desirable for high-power laser operation.

We conclude that laser slabs can produce high average power without thermal fracture or excessive beam distortion (Refs. 13 and 14). Glass slabs can be designed for low distortion at lower power densities; we shall see that they are appropriate for higher energy at lower repetition rate. Crystalline slabs are appropriate for higher power densities. Their best application is for use at lower energy and higher repetition rate.

The Slab Concept

Operation of a laser medium in slab geometry greatly reduces four problems that limit high-power operation: heat buildup in the material, thermal distortion of the wave front, depolarization of the beam, and beam profile changes due to gain variations.

Heat buildup is reduced by making the thermal path to the coolant as short as possible. Because the temperature rise for constant volumetric thermal loading increases as the square of this thermal path length, the potential improvement from path shortening is large. The slab geometry puts all the material close to the coolant, and therefore does a good job of cooling.

Thermal distortion of the wave front is reduced if the laser beam passes through the slab in such a way that all the rays in the beam have the same thermal history. An arrangement that leads to differing thermal histories, for example, is for the beam to pass at right angles to the thickness in a slab (and thus at right angles to the thermal profile gradient); some rays will then travel entirely in warmer regions, while others will travel entirely in cooler regions [Fig. 9-2(a)]. This leads to wave front distortion due to the change of refractive index with temperature. But if the beam passes through the slab in the same direction as the thermal gradient [Fig. 9-2(b)], all rays spend the same time in warmer and cooler regions and all have the same optical path.

Beam depolarization occurs when the laser material's refractive index changes with beam polarization (birefringence) and when the beam polarization is not aligned with one of the medium's two principal axes of polarization. Polarization-dependent

Fig. 9-2. Effect of spatially nonuniform temperature distribution on laser output beam quality. (a) In rod amplifiers, a radial temperature gradient generates a radial gradient in the index of refraction, which causes focusing or defocusing of the output beam. (b) In disk amplifiers the beam propagates parallel to the temperature gradient, and no focusing occurs.

indices arise from intrinsic material birefringence and from stress-induced birefringence. In a slab whose lateral dimensions are large compared with its thickness, the stresses are oriented in the plane of the slab and induce birefringence that is independent of position in the plane of the slab. The stresses do change with position through the thickness of the slab, but the principal axes of the index indicatrix do not rotate with distance along the beam if the slab material is initially isotropic. Thus the polarization of a beam polarized in or perpendicular to the plane of incidence is not disturbed. In fact, this conclusion holds for 16 of the 32 possible crystal symmetry classes. For another 12 classes, it is possible to orient the slab faces with respect to the crystalline axes in such a way that the polarization is undisturbed. Within roughly one slab thickness of the slab edges, the stresses will be nonplanar and some depolarization will take place. The fraction of the beam that is depolarized will be small if the lateral dimensions of the slab are large compared with its thickness. In addition, if the material is a uniaxial or biaxial crystal it is possible to align the natural birefringence in such a way that thermally induced changes in the index indicatrix are minimized.

The intensity profile of the beam will be distorted if different portions of the beam have different gain histories [Fig. 9-3(a)]. Thus it is desirable to pump through the slab faces, so that the beam path will cross the gain gradient and each ray will have the same gain history. The gain will

then be the same for every ray in the beam, and the beam will be amplified uniformly rather than having different parts end up at different intensities because of gain differences [Fig. 9-3(b)].

Physics and Engineering of Slab Amplifiers

Introduction. The design of high-average-power solid state lasers requires analysis of the pumping and decay properties of laser materials, of the control of parasitics at the material edges, of the thermal management of bulk and edge heating and the resulting depolarization and optical path changes, and of the gain, extraction, and fill-factor issues involved in output-beam production.¹³ Systems are synthesized on the basis of these analytic results and of laboratory experience with slab lasers. The optimum system for a given application is then found by maximizing the appropriate figure of merit for that application, subject to constraints imposed by slab physics and engineering and by the application itself.

Pumping and decay issues have been studied intensively for more than a decade in support of the Laser Fusion Program and are well understood. Computer simulation models permit selection of the best doping level and pump pulse duration for many different materials. Absorption and decay information needed by the models can be acquired from a few small samples of a material. Amplified spontaneous emission (ASE) in slab geometries has been analyzed by Monte Carlo methods to

(a) Rod amplifier

(b) Disk amplifier

Fig. 9-3. Effect of spatially nonuniform pump energy deposition on laser output beam quality. (a) In rod amplifiers, spatially nonuniform power deposition produces a radial gradient in the gain. The beam propagates perpendicular to this gradient, leading to a distortion in the beam. (b) In disk amplifiers, the beam propagates parallel to the gradient, and is not distorted.

find the size and gain limits to efficient slab operation.

Parasitic control has also been intensively studied for fusion applications, and the results have been adapted to high-average-power operation. The difference between the two applications is in the thermal load on the parasitic-suppressing coating. Removal of heat without excessive temperature rise combined with good parasitic suppression requires a coating with the proper combination of optical absorption depth (enough to achieve one-way attenuation by a factor of 10 to 100) and thermal diffusion length between pulses (more than half the coating thickness). In addition, the coating must be able to handle the stress caused by sudden heating during the laser pump cycle. We have developed computer programs to analyze temperature and stress in the coating under cyclic volumetric thermal loading (Fig. 9-4).

Thermal management to avoid damaging stress and optical degradation is the most exacting problem with high-average-power solid state lasers. The schematic analysis below reveals the most important factors involved, but full treatment of two- and three-dimensional thermal deposition and removal and of related optical effects in anisotropic slabs requires large and sophisticated computer programs. We have developed and are using such programs to carry out our analysis and synthesis tasks (see "TECATE," later in this section).

Gain, storage, extraction, and fill-factor questions have also been intensively studied for fusion applications. Analyses of the performance of single- and multiple-pass

amplifier systems are available, as are computer programs for very detailed laser beam propagation calculations, including the effects of diffraction, saturation, optical noise sources, and spatial relaying.

Basic Thermal Analysis. If a rectangular slab of isotropic material with thickness t (with t small compared with slab width and length) is uniformly heated throughout its volume by a volumetric power density P_v and cooled in such a way that its large faces are both at the same constant temperature, it reaches equilibrium with a parabolic temperature distribution across its thickness. The peak temperature of the distribution is higher than the face temperature by an amount T_c given by

$$T_c = \frac{P_v t^2}{8\kappa}, \quad (1)$$

where κ is the thermal conductivity. The temperature distribution leads to a stress distribution that is also parabolic but that has zero mean value. The slab surface is in tension; the surface stress is

$$\sigma_t = \frac{\alpha_t E P_v t^2}{12(1 - \nu)\kappa}, \quad (2)$$

where α_t is the thermal expansion coefficient, E is Young's modulus, and ν is Poisson's ratio. If the material is operated at the maximum permissible surface yield tensile stress S , the maximum volumetric thermal load is

$$P_{v,max} = \frac{12R}{t^2}, \quad (3)$$

where the so-called thermal shock parameter R is a combination of material constants:

$$R = \frac{(1 - \nu)\kappa S}{\alpha_t E}. \quad (4)$$

The thermal load P_v is in general related to the laser operating point by

$$P_v = f\rho\chi, \quad (5)$$

where f is the operating pulse repetition frequency (prf), ρ is the peak stored energy density in the laser material during a pump pulse, and χ is the ratio of thermal

Fig. 9-4. Calculated ASE absorber stresses due to repetitive pumping.

energy deposited in the material to peak optical energy stored in the material. In turn, χ is given approximately by

$$\chi \approx \frac{1}{\delta} \left(\frac{1}{Q} - 0.8 \frac{\tau_d}{\tau_u} + Z \right), \quad (6)$$

where δ is the ratio of peak stored energy to total available energy, Q (the so-called laser/pump quantum ratio) is the ratio of the laser photon energy to the energy of an average pump photon, τ_d and τ_u are the doped and undoped lifetimes, and Z is an extra factor to account for any pump input that bypasses the upper laser level. Note that Q is about 0.6 for flashlamp pumping of neodymium ions in glasses or crystals, and about 0.76 for narrow-band diode pumping at 810 nm.

The maximum frequency at which the laser can be operated without thermal fracture at the surface is found from the equation for P_v to be

$$f_{\max} = \frac{12R}{\rho\chi t^2}. \quad (7)$$

Because this frequency varies inversely with t^2 , one is strongly motivated to use

thinner slabs to achieve higher repetition rates. However, such slabs will absorb light less efficiently than thicker slabs and will have lower laser efficiency. We therefore want to examine the pumping of slabs in detail to see how much efficiency we must give up to get high repetition rate.

Slab Pumping Analysis. We take the pump efficiency to be the product of six factors:

- The pump cavity transfer efficiency.
- The slab absorption efficiency.
- The laser/pump quantum ratio Q .
- The storage efficiency (compared with that of an infinite-lifetime material).
- The extraction efficiency.
- The aperture fill factor.

The cavity transfer efficiency is the most difficult of these factors to model, because it depends on the geometric, emitting, absorbing, and reflecting properties of all the materials in the cavity. Rather than calculate these factors for each pump cavity, we fitted the extensive data available from the single-segment amplifier (SSA) experiments. (See "Single-Segment Amplifier: Experiments," in Sec. 7.) Because experiment has shown that lamp properties depend mostly on instantaneous wall power

Fig. 9-5. SSA amplifier efficiency (prepulsed) and the approximation which results from use of the cavity transfer function shown. T is the pump base-to-base duration.

loading (see "Diameter Scaling," in Sec. 7), we assumed that the cavity properties depend only on wall loading and pump pulse duration. We made a fit to the peak gains observed in the SSA experiments (Fig. 9-5), resulting in a pulse-averaged cavity transfer efficiency (Fig. 9-6) that rises rapidly with loading (because some energy is needed to prepulse the arc) and then falls slowly with further loading (because of spectral blue shift and increased reabsorption). The cavity is less efficient at shorter pulse durations, presumably because the lamp arc is less efficient if it does not have time to form properly.

Slab absorption is obtained from VODEP-code absorption coefficients reported in the Laser Glass Handbook,¹⁵ after multiplying the Handbook values by a factor of 3.7 to convert them to fractional absorption. This formulation does not take into account multiple-pass absorption when the slab is optically thin, which makes thin slabs in tight pump cavities relatively more efficient.

The laser/pump quantum ratio Q is also obtained from VODEP calculations. For flashlamps pumping neodymium in glass, this ratio is always near 0.6.

The storage efficiency δ was found by modeling the inversion pumping rate as a half-sine shape in time, and the decay as a simple exponential. The fraction stored was found numerically and then fitted by the expression

$$\delta = \left[1 + \frac{1 + 96\epsilon^2}{\epsilon(1.57 + 192\epsilon^2)} \right]^{-1}, \quad (8)$$

where ϵ is the ratio of decay e -folding time to pump base-to-base time.

The extraction efficiency η_{ex} is found by use of the Frantz-Nodvik result for an amplifier of gain G ,

$$\eta_{ex} = \frac{1}{\ln G} \times \left[T \ln \left(\frac{G-1}{1/T-1} \right) - \ln \left(\frac{1-1/G}{1-T} \right) \right], \quad (9)$$

where T is the transmission of the unpumped laser cavity and slab.

The aperture fill factor depends on the size and apodization of apertures in the laser beam path, and on optical relaying and spatial filtering (if used). It can range from very low values (0.01 or less in

Fig. 9-6. SSA cavity efficiency (pulse average) which leads to the approximation shown in Fig. 9-5.

Fig. 9-7. (a) Efficiency of slab laser with Hoya LHG-5 glass for various values of the total (half-sine-wave) pump duration T . (b). Efficiency for various dopings and pump durations. (c). Efficiency for various pump energies, dopings, and pump durations. (d). Efficiency for various slab thicknesses, pump energies, dopings, and pump durations.

poorly designed systems) to values as high as 0.7 or 0.8 in large fusion lasers.

Design of Slab Lasers

Lasers are designed by plotting the power and efficiency available from various materials against repetition rate. To obtain data points for such a plot, the thermal and pumping models are used to calculate the power, efficiency, and repetition rate of a slab as its parameters are varied. The appropriate material for a specific application can then be selected and its dimensions specified.

For example, Fig. 9-7(a) shows the efficiency of a 2-cm-thick slab of LHG-5 glass as the pump pulse duration is varied, for fixed values of pump energy, slab doping, and internal and external losses. The figure shows that there is a pump duration for which maximum efficiency is realized, but that repetition rate can increase almost indefinitely at the expense of efficiency. Figure 9-7(b) shows curves for several dopings. At any frequency, there is an optimum combination of doping and pump duration for highest efficiency. For each doping in Fig. 9-7(b), Fig. 9-7(c) shows

families of curves for different pump energies. We now see good performance at a wide variety of repetition rates. Figure 9-7(d) adds the possibility of changing the slab thickness, which gives us an even broader operating region that maintains high efficiency.

For phosphate glass, Fig. 9-8 indicates that the achievable efficiency stays high up to surprisingly high repetition rates, although it drops steadily as repetition rate increases. Figure 9-8 also shows the potential of YAG and co-doped GSGG, which give high repetition rate and high efficiency, respectively. Figure 9-9 shows the output powers corresponding to the efficiencies of Fig. 9-8. All three materials can produce about 1 kW per plate, but phosphate glass does this by producing 1 kJ at 1 Hz, while the crystalline materials produce 1 J at 1 kHz. For all three materials, output power drops as repetition rate rises.

Some Design Examples

First consider a 2-cm-thick phosphate glass slab. Taking a beam aspect ratio of 10:1 for high fractional use of the slab

width, we have an active width of 20 cm. To get enough gain for good extraction efficiency, we choose a pumped length of approximately 50 cm. The length, thickness, and refractive index are interrelated because of the beam-bounce geometry.

With a glass index of 1.53 and a coolant (or evanescent-wave coating) index of 1.34, the internal bounce angle must be at least

$$\sin^{-1}(1.34/1.53) = 61^{\circ} \quad (10)$$

Fig. 9-8. Ratio of laser beam power to lamp power input for three slab laser media configured as shown.

Fig. 9-9. Power output for laser media shown in preceding figure.

for total internal reflection (TIR) to occur. If we require both straight-in entry and 100% fill of the volume, the bounce angle must be

$$\sin^{-1} \frac{1 + \sqrt{1 + 8(1.53)^2}}{4(1.53)} = 62.75^\circ. \quad (11)$$

This is far enough from TIR trouble that we choose it as the bounce angle. With 16 bounces inside, the slab must then be

$$\frac{L}{t} = (16 + 1)\tan 62.75 = 33.01 \quad (12)$$

times as long (tip to tip) as it is thick. Thus a 2-cm slab must be 66 cm long. We want the beam to overlap itself exactly inside the slab, so that each part of the slab is swept twice by the beam and no parts are missed. For this to occur, the tip wedge angle must be the complement of the bounce angle, or 27.25° . This is small enough to reduce our fear of end-to-end parasitic modes. To reduce beam distortion, we want to pump the slab over a length corresponding to an even number of bounces. The length per bounce is $t \tan 62.75 = 1.94t$; for $t = 2$ cm, 12 bounces thus take 46.6 cm and 14 bounces take 54.4 cm. We choose to pump over 12 bounces of length to keep the pumped region well away from the end regions and thus avoid beam distortion.

The slab doping gives maximum efficiency at an ion density of about 6 or $7 \times 10^{20} \text{ cm}^{-3}$ (Fig. 9-7).

The lamp size and spacing is chosen to give a somewhat lower pump density in the slab head than in the SSA test head, since the laser material has higher cross section and will give good gain at lower pump densities. Also the pumped material is in the pump plane, rather than tilted, which gives 20% more energy density for the same pump energy at the lamps. The SSA has 1.44 units of lamp area (both sides) per unit of pump plane area. We choose to operate with a lamp/slab area ratio near unity. This can be achieved by using three of our standard 15-mm lamps per side, for a lamp/slab area ratio of

$$\frac{3\pi \times 1.5 \times 2}{20} = 1.41. \quad (13)$$

However, this setup is electrically inconvenient. We therefore use two 19-mm lamps per side, giving a lamp/slab area ratio of

$$\frac{2\pi \times 1.9 \times 2}{20} = 1.19. \quad (14)$$

We assume that the efficiency gain resulting from the wider lamp spacing in the slab head (compared with that in the SSA reference) and from the placement of the slab directly below the lamps is compensated by the efficiency loss due to the cooling jackets around the lamps and the water channel over the slab, so that the cavity efficiency is the same as that of the SSA.

We now have in hand all the factors with which to approximate the slab performance as we vary pump energy and pump pulse duration. Figure 9-10 shows the resulting slab power vs repetition rate for various energy inputs to the flashlamps. We assume that the slab material has a loss coefficient of 0.1 m^{-1} , and that the external cavity transmission is 0.85. We use a fill factor of 80% to allow for apodization and diffraction effects. The highest power is achieved at 2 to 4 Hz, but the half-power frequency is near 10 Hz. Figure 9-11 shows efficiency vs repetition rate. This curve has the same general behavior as the power, peaking near 2 Hz. Note that at some low repetition rate (and at a correspondingly high energy per pulse), we are apt to run into parasitic problems and be unable to further increase the power and efficiency.

Fig. 9-10. Power output of a $2 \times 20 \times 50$ -cm phosphate glass slab laser at various prf's for various energy inputs to the flashlamps.

Fig. 9-11. Efficiency of the laser of the preceding figure.

Now consider a plate of $GdGG$ co-doped with chromium to increase its absorption efficiency. A pumped volume of $0.4 \times 5 \times 10$ cm can produce 200 mJ per pulse at 1500 Hz, for a power of 300 W per plate. Twenty such plates in a gas-cooled slab geometry would deliver 4 J per pulse, for a total power of 6 kW, at an overall efficiency of about 2%. The cooling gas must remove about 10 W/cm^2 from the plate surfaces. About 2 kW/cm^2 must be removed from the lamp walls; this will require improved technology in the form of thin sapphire envelopes or the use of liquid-vortex inside-wall-cooled lamps.

Author: J. B. Trenholme

Experiments

Zig-Zag Slab Architecture

Introduction. The mission of the experimental program is to evaluate the performance of medium- and high-average-power (HAP) solid state laser architectures. We have constructed facilities to test two architectures: the liquid-cooled zig-zag slab and the gas-cooled Brewster's-angle slab. Although our immediate goals are to characterize laser power, beam quality, and efficiency, our long-range objective is to understand the design limits of HAP solid state laser systems. This can only be accomplished by developing a computer modeling capability that accurately calculates

the prescribed performance parameters. Consequently, our experimental efforts have focused on evaluating laser designs that have well-defined optical, thermal, and mechanical boundary conditions, thereby providing an unambiguous set of data for validating our computer models.

We have characterized the optical performance of a flashlamp-pumped slab amplifier geometry using a helium-neon probe laser to measure the wave front distortion and depolarization. We measured the gain uniformity (in space) and compared it with measurements of the pump uniformity. Our gain data can also provide a measure of the storage efficiency, which, together with accurate thermal and electrical calorimetry, can provide a measure of the ratio of thermal to optical power for the laser material. This scaling parameter is essential for accurate laser performance modeling.

We can cross-check the accuracy of these measurements against the wave front and depolarization calculations, which are obtained by 3-D simulation, if we have accurate values for the slab's thermo-optic, stress-optic, and mechanical properties. To further check the modeling accuracy, we took optical data on two different geometries: the zig-zag and the "straight through" propagation.

Wave-Front Distortion. The 10-bounce zig-zag slab test bed consists of an LSG-91H silicate slab, $32 \times 10 \times 1.5$ cm, which is mounted in a flow channel where water flows over the slab's two flashlamp-pumped faces. With the thermal gradients predominantly across the thin dimension, we can average the optical aberrations usually observed when using the straight-through propagation by using a zig-zag propagation. We can adjust thermal boundary conditions at the edge of the slab to help minimize the wave front distortion. The slab wave front distortion was measured by using a Fizeau interferometer and a helium-neon probe laser. Figure 9-12(a) shows a double-pass interferogram of the zig-zag slab that is observed at 1600 W of flashlamp input power without adjusting the edge temperature and slab coolant temperature. Figure 9-12(b) displays an optically corrected (vertical direction only) version of Fig. 9-12(a). We can obtain similar interferograms without optical

correction by adjusting the slab edge temperature. This tunes out aberrations resulting from conduction through the slab

holder as well as some effects caused by pump nonuniformities. These thermal boundary condition adjustments influence

(a) Uncorrected

(b) Corrected

Fig. 9-12. Zig-zag wave front distortion with focus corrected optically ($P_T = 1600$ W).

(a) Rabl, $P_T = 750$ W

(b) Goodwin, $P_T = 750$ W

Fig. 9-13. Uniform pumping with a Goodwin reflector reduces the wave front distortion.

the wave front distortion in the vertical direction. The horizontal distortions, which are across the thin dimension of the aperture, are predominantly affected by the slab-surface dimensional changes that result from pump and cooling transitions. The ray tracing code indicates that these effects can be minimized by adjusting the relative location of beam "footprints" in the TIR surface and the thermally induced surface deformations.

Although we did not achieve complete optical compensation in the initial test geometry, the numerical models predict that small changes in pump and cooling geometry will give an improved beam. This is the focus of our next experimental series.

Reflectors. We illustrated the sensitivity of wave front distortion to pump uniformity by comparing two radically different

reflector designs. The highly nonuniform pump pattern of the Rabl reflector design produced a severely distorted wave front, as displayed in Fig. 9-13(a). In contrast, pumping with the Goodwin reflector design produced a significantly improved wave front. To more precisely measure the pump profile, we photographed the reflector/flashlamp assembly (see Fig. 9-14). The film was digitized, which resulted in the pump intensity plots displayed in Fig. 9-15. When the as-built reflector dimensions and reflectivity were included in the numerical simulations, the measured and calculated scans transverse to the lamp axis and in the orthogonal direction agreed closely.

We attempted to improve the pump uniformity by applying a diffuser stripe to the lamp cooling jacket, thereby reducing the intensity of the direct-shine component from the lamp. We obtained the desired effect [see Fig. 9-15(b)], but we believe that the pump uniformity is not sufficient to produce a high-quality wave front because of thermal gradients induced by the roll-off at the edges. We believe that this spatial behavior is responsible for the vertical focusing component observed in the interferograms of Fig. 9-12. Numerical simulation of the reflector indicates that improved pump uniformity can be obtained with higher reflectivity ($>85\%$) and small adjustments to the lamp location.

Gain. Gain measurements on the zig-zag test bed were completed to provide data for comparison with the

Fig. 9-14. Schematic of reflector performance measurement. The diffuse stripe was evaluated as a method of fine tuning the pump uniformity.

Fig. 9-15. Pump uniformity, measured photographically, agrees closely with our numerical model.

well-characterized single segment amplifier (SSA). We took measurements on a 10-bounce, 2.7% neodymium-doped zig-zag slab and a 1.9% neodymium-doped straight-through slab. We measured storage efficiencies of up to 5.3% (see Fig. 9-16), which compare favorably with the SSA results, despite the lower doping-thickness product. The measured gain profiles along the height (100 mm) of the slab compare well with the measured illumination profiles of the reflector and the calculated pump distributions (see Fig. 9-17). The measured gain profiles through the thickness (15 mm) of the slab compare well with simulations when recirculation of flashlamp light within the pump cavity is included.

The observation that recirculation of pump light can significantly reduce the nonuniform gain distribution within a slab or disk has not been made previously. This effect influences the thermal deposition and stress within a slab, the efficiency of an amplifier as a function of pump power, and our understanding about face parasitics.

In 1985 we constructed the facilities required to measure the performance of medium- and high-average-power solid state lasers. We also began the necessary experiments to evaluate numerical models needed for laser system design and scale-up. Our first data set is sufficient to verify the numerical models needed to fine tune the beam quality and efficiency.

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Gas-Cooled-Slab Architecture

Introduction. We installed the gas-cooled-slab (GCS) test bed in April 1985. Although program priorities dictated that most of the experimental resources of the HAP Program were directed to the zig-zag test bed, two significant experiments were performed on the GCS facility. With 3 kW of electrical power input to the flashlamps, we used an IR camera to look through the windows and coolant and to

Fig. 9-16. Probe beam gain implies a storage efficiency of $\pm 5\%$.

Fig. 9-17. Gain distribution reveals pump uniformity in one dimension.

record the surface temperature profile of the Nd:glass slab. As described in the first article of this subsection, "Preliminary Thermal Characterization," the temperature profiles were consistent with numerical calculations of thermal boundary-layer development.

We performed a series of experiments to demonstrate that the nitrogen gas flow across the slab is both rectilinear and of uniform velocity. The results, which met our design goals, are described below.

Author: J. A. Blink

Preliminary Thermal Characterization.

We performed a thermal test series as part of the overall qualifications of the GCS test bed. We measured slab surface temperatures at 0-, 0.72-, and 2.56-kW flashlamp input power levels, all at a fixed nitrogen-coolant flow rate of 0.84 kg/s. We used a Hughes IR camera to measure temperatures on one side of the slab as the test bed approached and

Table 9-2. Thermal test conditions and observed slab-surface temperatures.

Experiments

	Lamp input power (kW)		
	0	0.72	2.56
Nitrogen flow rate (kg/s)	0.84	0.84	0.84
Inlet stagnation pressure (kPa)	202	202	202
Mach number in channel	0.6	0.6	0.6
Surface temperature (°C)			
At channel inlet	18.3	19.4	25.4
Peak value	18.3	21.4	29.2
At channel exit	18.3	21.1	28.8

Fig. 9-18. Comparison of calculated and measured slab-surface temperature.

attained steady-state performance at each of these operating conditions.

Table 9-2 summarizes the specific test conditions and the observed steady-state surface-temperature distributions. For the 0.72- and 2.56-kW power conditions, surface temperatures near the channel inlet increased in the flow direction, as expected for developing convective conditions on the slab surface. In contrast, the surface temperature decreased from a peak value near the slab center to a slightly lower value at the channel exit; this behavior reflects a continued adjustment of the channel flow structure, an axial conduction at the slab edge, or a spatial variation in the incident flashlamp energy.

The zero-power condition provided a baseline measurement with a uniform surface temperature that by definition is the adiabatic wall temperature for the fixed coolant flow conditions. The observed lateral uniformity of surface temperature

in all three tests was particularly important for qualifying the test bed. About 90% of the slab width was at a nearly constant temperature, indicating that two-dimensional channel flow had been achieved except for corner regions of limited extent.

The test bed was not sufficiently instrumented to establish detailed thermal conditions for the slab, but we deduced additional information from a match of calculated and measured surface temperatures. Figure 9-18 compares measured and calculated slab-surface temperatures for the 2.56-kW flashlamp power. A uniform surface heat flux was assumed in formulating the two-dimensional turbulent channel flow model. We adjusted the surface heat flux to 0.9 W/cm^2 to achieve this match; this value we considered to be the “best estimate” of average surface heat flux. Furthermore, heating of a 5-mm-wide shim that covers the leading edge of the slab also had to be assumed for a good match at the channel inlet; it follows that the shim is absorbing a significant amount of incident flashlamp light.

Based on our observed and deduced results, we concluded that the convective cooling system in the test bed is uniform, predictable, and can facilitate meaningful slab-laser experiments. Specifically, the gas-flow and surface-convection characteristics are predominantly two-dimensional and the measured thermal boundary layer development along the slab surface is close to that calculated with the computer model. This supports previously calculated cooling performance estimates. However, more detailed experiments are needed to explain in detail the decrease of surface temperature near the channel exit.

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Flow Characterization. We experimentally evaluated the nitrogen cooling system for the gas-cooled-slab test bed and found that it provides a stable and uniform flow of gas.¹⁶ We measured dynamic and static pressures at various locations in the vertical ducting to characterize the system. We

assessed velocity variations on the basis of dynamic pressures and streamline curvature on the basis of static pressures. Flow development was assessed on the basis of boundary layer traverses at two positions along the duct wall.

Experimental Arrangement. We positioned the rectangular vertical duct sections at a 30° angle to the horizontal pipes in order to mate with the flow channels in the test bed, which are positioned at 30° (Brewster's angle) to the horizontal pipes. We substituted a venturi duct section for the test bed to directly measure flow maldistributions at the test bed elevation. Figure 9-19 is a top view of this arrangement; the rectangular duct cross-section at the venturi elevation is shown with a phantom line. The adopted north-south reference frame is also identified in Fig. 9-19.

The venturi duct section simulates the test bed: its throat height is 10 mm, which represents the combined cross section of the two 5-mm channels in the initial test bed configuration. Its overall flow characteristics are expected to be comparable to those of the initial test bed.

We provided wall static pressure taps and access ports for pressure probes at ten elevations in the vertical duct, including the venturi throat. To minimize the number of duct wall penetrations,

pressure taps and access ports were machined in only one side of most duct sections; sections can be rotated 180° if access from the opposite side is needed. We arranged the inlet sections with taps and ports on the north side. The exhaust diffuser was arranged with taps and ports on the south side. The venturi duct section is an exception; it has static pressure taps on both sides.

The following instruments were used to obtain the reported measurements.

- A Wallace and Tiernan gauge, with a full-scale reading of 1550 mmHg (207 kPa) for measuring pressure at a given location.

- A Merriam inclinometer, with a full-scale reading of 1 in. H_2O (248 Pa) for measuring differential pressure between two locations.

- United Sensor pitot-static pressure probes, which simultaneously indicate stagnation and static pressures.

- United Sensor Kiel-type probes, which measure stagnation pressure more accurately than pitot-static probes.

- United Sensor boundary layer probes, constructed from thin-wall, small-diameter tubing.

Experimental Results. Our static-pressure measurements indicate that the flow is rectilinear at all measurement elevations. The most important variation

Fig. 9-19. Relative orientation of the test bed, the venturi replacement, and horizontal pipe sections.

Experiments

occurs at the venturi throat (see Fig. 9-20), since a comparable variation is expected in the test bed flow channels. The implied radius of curvature is 14 m. This small deviation from rectilinear flow is considered acceptable with regard to convective cooling of the laser slab.

We measured dynamic pressures in the inlet and outlet ducts. At the venturi duct

inlet, the dynamic-pressure variation was within $\pm 1\%$ of the section's mean, a small variation that is comparable to the measurement uncertainty of $\pm 0.6\%$ (see Fig. 9-21). On the basis of these data, we concluded that the flow uniformity is acceptable.

We also used the dynamic-pressure measurements to calculate velocities for comparison with values obtained from the volume flowmeter. The comparison indicates an uneven split between the two sides of the diffuser. The volume flowmeter and its associated pressure and temperature sensors were used to determine the mass flow rate. The static pressure at the elevation, coupled with the elevation's flow area and the mass flow rate, allows calculation of the average velocity at that elevation. Table 9-3 displays the velocities calculated from the two measurements.

The velocities for the upstream sections differ by less than 25%. However, downstream of the venturi on the south side, the velocity based on measured dynamic pressure is more than twice the value based on total mass flow. This high value indicates a north-south flow maldistribution

Fig. 9-20. Static pressure variation across the duct at the venturi throat. Data at the north and south faces obtained at a mass flow of 0.38 kg/s are shown.

Fig. 9-21. Dynamic pressure uniformity in the downstream portion of the inlet duct, near the north and south walls. These data were taken at a flow rate of 0.38 kg/s.

downstream of the venturi, because reasonable variations were seen at upstream elevations.

Our final measurements were boundary-layer traverses in the inlet and exhaust ducts, which were used to characterize the channel flow development. We measured each profile on the duct wall opposite the midspan access port. We measured the profile on the south wall in the inlet duct and on the south side of the center vane in the exhaust duct.

The boundary-layer profile measured in the inlet duct is characteristic of a turbulent boundary layer extending approximately 7 mm out from the wall. Figure 9-22 displays a plot of velocity V , normalized by the free-stream velocity V_∞ , as function of the distance y from the wall, normalized by the boundary-layer thickness δ . Our calculated boundary-layer parameters are itemized in Fig. 9-22. The boundary-layer profile measured in the exhaust duct is also characteristic of a turbulent boundary layer, but it extends approximately 18 mm out from the wall. The free-stream velocity V_∞ is 160% higher than that expected for uniform flow based on the measured mass flow; this is similar to the finding in the dynamic-pressure surveys.

The boundary-layer profiles that we used to determine the friction factor highlight a difficulty we had in obtaining these data. Figure 9-23 displays a profile in the standard dimensionless coordinates V^+ and y^+ based on matching the profile to theoretical predictions in the turbulent region. Data in the laminar sublayer region do not match the theoretical sublayer correlation; this suggests that the probe position relative to the wall was not measured correctly. We initially assumed that the probe tip resided on the wall during the traverse; we now believe that it was displaced off the wall. (During insertion into the duct, flexing of the delicate probe was required, and it may not have returned to the correct orientation.) To compensate for the orientation, we applied a fixed correction to y (0.634 mm) to achieve a match in both the turbulent region and the laminar sublayer. Figure 9-23 also includes our corrected data. With the correction, the friction factors, which are based on a match in the turbulent region, changed by

Elevation	Average velocities at the elevation (m/s)		
	From flowmeter (V_{fm})	From dynamic pressure (V_p)	V_p/V_{fm} , %
1	3.15	2.4 to 3.1 ^a	77 to 99 ^a
1	9.3	7.5	80
4	12.1	13.3	110
10	3.0	6.5	215

^a The two values represent the averages near the south and north walls.

Fig. 9-22. Measured velocity distribution on the south wall in the inlet duct section.

Fig. 9-23. Measured velocity distribution near the wall in the inlet duct section. The data correction compensates for the probe position uncertainty. Solid line from Ref. 17.

Fig. 9-24. Comparison of "best estimate" friction factors with a flat plate correlation. Boundary layer separation, indicated for friction factors below 0.0001, is not expected. Solid line from Ref. 17.

less than 3%. This indicates that our method of determining friction factors is somewhat insensitive to probe position error. Hence, the values displayed in Fig. 9-23 were selected as "best estimates" for purposes of qualifying the system.

Figure 9-24 compares the results with the generally accepted correlation for friction factors in a turbulent boundary layer on a flat plate. The data fit the correlation remarkably well and imminent separation of the boundary layer in the exhaust diffuser, corresponding to a friction factor of 0.0001, is not indicated. We conclude from this comparison that the boundary layers are developing on the duct walls in a normal manner and that separation in the exhaust diffuser is not occurring.

After we removed the venturi duct, we installed the test bed and again surveyed the dynamic pressures in the exhaust duct. The previously observed high velocities in the south side of the diffuser were gone, and the mean velocity based on total mass flow and duct cross sectional area agreed with the probe values within 10%. We concluded that the test bed's configuration, which includes honeycomb flow straighteners and a center vane, had corrected the flow maldistribution in the exhaust. Consequently, we considered the system qualified for tests using the initial test bed configuration.

We must make design modifications to the exhaust system to isolate the flows in the vertical duct from downstream flow maldistributions. Subsequently, we will be

able to evaluate alternative test bed configurations with the assurance that the flows in the slab coolant channels are uniform. We have inserted flow-straightening screens in the square-to-round transition duct downstream of the diffuser duct as a means of achieving this goal.

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Typical System Designs

Introduction

Research in high-average-power solid-state laser technology, under DARPA sponsorship, has been directed toward removing the limitations on average power, beam quality, and wavelength variability. When developed, this technology will be used in Department of Defense applications such as communications, countermeasures, radar, and for specialized applications such as an oscillator for the free electron laser (FEL). An important part of this development is system design and analysis. This process gives the technologist hardware objectives and "directional derivatives" to guide research into areas with the highest systems paybacks.

In analyzing the degree to which average-power solid state laser technology can meet the system requirements in these applications, we studied several laser configurations and derived some important conclusions.

1. Single-plate crystalline zig-zag amplifiers can operate at a PRF well above 1 kHz. Laser cavities using single-plate amplifiers can provide over 1 kW of average power. These systems are capable of meeting the experimental FEL master oscillators goals and, in addition, may meet many industrial applications.

2. Flashlamp-pumped glass zig-zag lasers are capable of providing over 1 kJ of frequency converted light at ~ 1 Hz (average repetition rate) for an average power of > 1 kW. Such a laser can assume a rather compact package, < 1 m³ in volume and

less than 1000 kg in weight, with an overall efficiency exceeding 4%.

3. The solid state laser used as an optical radar transmitter can provide power levels exceeding 1 kW. It uses optical switches to multiplex many lasers into one output beam and it appears capable of maintaining phase coherence for pulse lengths exceeding tens of microseconds.

System studies show the range over which high-average-power solid state lasers can be used. Studies of this nature have provided, for example, the impetus to develop new laser materials (the promethium ion in a fluoride host) and to develop new components such as high-power multiplex switches. Such studies have oriented our prototyping demonstrations of the base technology program into important system areas. An example is our use of an injection-locked, relayed, phase-conjugated laser resonator for our first laser demonstrations planned for the end of FY 86.

Developing the average-power solid state laser could resolve many Department of Defense, Department of Energy, and industrial problems; the list grows rapidly. We are finding that the technical richness of this technology provides a wide range of solutions to system problems. The main limitations to average-power solid state laser applications are in the growth, fabrication, and purity of laser materials. They are consuming a major fraction of our development resources. Materials development activities are deserving of additional support by government funding agencies. We continue to be excited by the remarkable number of applications for this technology.

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Solid State Transmitter Laser for Submarine Laser Communication

Introduction. The objective of the DARPA/Navy submarine laser communication (SLC) program is to develop component and system technologies for a satellite-borne laser transmitter capable of providing significantly improved strategic

and tactical communication with submerged submarines as compared with existing non-optical communication systems. In 1985 we identified, analyzed, and ranked a number of solid state transmitter lasers (SSTL) in our search for an *advanced* laser transmitter that will meet SLC requirements.

The present design for a SLC system requires that the SSTL output wavelength match that of an atomic resonance absorption and fluorescence (ARAF) filter in the blue portion of the ocean-water transmission spectrum. The ARAF filter exhibits both the large field of view and the extremely narrow-band spectral responsivity required to adequately reject solar background radiation.

DARPA and the Navy are presently developing a *baseline* SLC system in which the output wavelength of the xenon-chloride excimer-discharge laser is Raman converted in lead vapor to match the ARAF filter.

While this baseline transmitter module based on XeCl + Pb technology is likely to achieve useful performance levels (efficiency, pulse energy, pulse repetition frequency, size, and weight), concerns remain about the operating life of the system and its effect on system cost.

Preliminary analyses of the SLC system (in terms of satellite and launch costs and their dependence on the size, weight, and life of the laser transmitter) suggest that a significant cost benefit can be realized with laser transmitters that can operate for at least seven years or that are light enough and small enough to permit several units to be placed on a single satellite. A SSTL might well result in a SLC system with improved life-cycle costs compared with the existing baseline design.

Solid State Transmitter Laser. Each SSTL consists of a primary solid state laser pumped by broad-area diode laser arrays operating at some discrete wavelength in the near infrared and of one or more wavelength converters used to shift the light from the primary output to a wavelength matching that of the ARAF filter. Out of 27 solid state transmitter module candidates identified (a half-dozen concepts and many material variants), we determined that 12 concepts were sufficiently simple and potentially efficient for further detailed analysis. These concepts are

listed in Table 9-4. We roughly ranked each concept quantitatively on the basis of estimated efficiencies and qualitatively on the basis of such optical design issues as the number of wavelength-stabilized beamlines and the presence of higher or lower beam fluences along a beamline. The projected efficiency of each SSTL was determined by taking into account material parameters, cavity losses, number and type of nonlinear conversion steps required, etc.

In addition to transmitter electrical efficiency, an accurate rank ordering requires data on size, weight, and operating life. These parameters are extremely complex and were outside the scope of the present study; however, their determination will be undertaken in the future.

The systems represented in Table 9-4 fall into two general categories. The first category employs moderately tunable wavelength shifters, which allow some flexibility in optimizing the solid state laser

host material. The second category employs host-specific neodymium laser materials with strong transitions at specific wavelengths, which can be converted to reach the filter wavelengths.

Laser System Characteristics. We evaluated the electrical efficiencies, nominal sizes, and other salient characteristics of each laser system. The purpose of this evaluation was to highlight the relative electrical efficiency, optical complexity, and materials-specific parameters of the candidate options. The three most efficient laser systems are summarized in Fig. 9-25. Not included in the efficiency estimates are the factors associated with power conditioning from the solar cells to the diode array, thermal management requirements, and all other auxiliary functions required to make the total system functional. Future work will explore the properties of configuration, size, weight, and the issues associated with integrating the spacecraft system and subsystems for some of the more promising designs.

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Table 9-4. Laser options matrix.

Tunable systems	Fixed-wavelength systems
Nd:() + 3 ω + OPA	Nd:LaMgAl ₁₁ O ₁₉ + 3 ω
Nd:() + 3 ω + dye	Nd:BEL + 3 ω
Nd:() + 2 ω + OPA + 2 ω	Nd:() + CD ₄ /SiH ₄ + 3 ω
Nd:() + 2 ω + TiAl ₂ O ₃ + mix	Nd:LuAlO ₃ + 3 ω + H ₂
Nd:() + 2 ω + TiAl ₂ O ₃ + 2 ω	Nd:LuAlO ₃ + LiNbO ₃ + 3 ω
Nd:YAP + 3 ω + D ₂	
Nd:YAP + 3 ω + CH ₄ (2)	

Legend (): any host material
OPA: optical parametric amplifier
BEL: lanthanum beryllium aluminate
YAP: yttrium aluminum perovskite

Nonlinear Optics

Introduction

Laser technology is moving towards increased average power. Potential

Fig. 9-25. The three most efficient neodymium laser systems calculated for a diode-array efficiency (light output to electrical power input) of 0.3.

applications include submarine laser communication, high resolution optical radars, high energy battlefield lasers, and laser fusion. Frequency conversion and optical switching are essential elements of high-power laser technology. In 1985 we completed three tasks necessary to the development of nonlinear optical devices at high average power.

- We completed a thorough analysis of the thermal behavior of harmonic generators.
 - We completed a thorough study of various methods of optical switching at high average power.
 - We completed the characterization of three useful nonlinear optical materials for harmonic generation, so that realistic designs for high average power devices can be drawn up.
- We summarize here the results of these studies.

High-Average-Power Harmonic Generation

Frequency conversion is a useful technique for extending the utility of high-power lasers. It is in common use for single-pulse and low-average-power applications, but there has been comparatively little interest in scaling to high average power. Broadly speaking, this is because the average power in the converted beam is limited primarily by the average power capability of the laser drive. It has been widely assumed that solid state materials do not handle thermal loads efficiently, and it has thus been argued that frequency conversion using solid state materials would not be practical even if diffraction-limited, high-average-power lasers were available.

Increasing the average power of a laser generally degrades the beam quality or damages the laser by thermal loading. Under these conditions the primary requirement on the harmonic generator is that it maintain high conversion efficiency for a laser with poor spatial quality. The highest conversion efficiency is then obtained in materials with large acceptance angles and high damage thresholds. Although the linear optical absorption in nonlinear crystals gives rise to thermal gradients that destroy phasematching, managing the heat

deposition has been important only if the acceptance angle is very large—that is, under conditions of noncritical phasematching. Typically, managing laser beam quality has been more important than managing thermal gradients in harmonic generators.

Thus, frequency conversion has not progressed beyond relatively modest average powers, in part because without a high-average-power Nd:YAG laser there is little or no imperative to develop the frequency conversion techniques to handle higher average power. Although the development of high-average-power solid state lasers languished during the last two decades (with CO₂ being the prime average power laser), these lasers have recently begun to enjoy renewed interest. One reason for this is the recent appearance of new applications for lasers, including those associated with military programs and the Strategic Defense Initiative.¹⁸ The average power required in these applications varies up to several orders of magnitude beyond the current state of the art in average-power frequency conversion.

As part of our program to develop a laser for the commercial application of laser fusion, we reexamined the potential of solid state lasers to achieve high average power.¹⁹ New design concepts and recent developments in materials and techniques provided a basis for that study. Particular attention was paid to thermal management and optical distortion, and especially to efficient utilization of pump energy. It was concluded that solid state lasers are capable of unconstrained scaling of average power while maintaining high beam quality. Moreover, it was found that efficiencies of over 10% should be achieved at high power and irradiance with careful attention to conservation of pump energy and to thermal management. In 1985, on the basis of experience gained in frequency conversion of a 100-kJ, 100-TW laser,^{20,21} the potential of solid state nonlinear materials to reach high average power while maintaining high conversion efficiency was thoroughly evaluated.

The state of the art in high-average-power lasers is a few tens of watts. We have shown that solid state nonlinear materials can surpass this level by at least three orders of magnitude with careful attention to thermal gradients and beam

shape and by balancing the various sources of phase mismatch. Moreover, configurations exist that permit essentially unconstrained scaling of average power.

In average-power operation, the linear optical absorption in the crystals causes temperature gradients that can reduce the conversion efficiency by disturbing phasematching. If nonlinear crystals were perfectly transparent, there would be no special aspects to average-power harmonic generation. Design would then proceed along standard lines, and average power would cease to be an issue. However, most nonlinear materials are not transparent enough that thermal effects can be

ignored. A linear absorption of 0.1 m^{-1} is often great enough to be significant.

Because the interest is in high average power, intrinsically unscalable geometries such as tight focusing were not considered. A careful assessment of heat flow in the nonlinear crystal (Fig. 9-26) showed that the thermal gradients must be perpendicular to the crystal faces for efficient, scalable heat removal. The optimum geometry uses plane waves and a number of crystal plates, between which a coolant gas flows to remove the heat deposited by optical absorption. The gas flow is high speed and turbulent for efficient heat removal. By dividing a large aperture into segments (a technique demonstrated on Nova²¹) and optimizing each segment, the total average power in the entire aperture can be made essentially unconstrained. Each segment consists of several plates of single crystals through which the beam passes consecutively, as shown in Fig. 9-27. The average power in each segment is constrained by thermal dephasing, surface fracture, surface temperature, and optical damage. Each of these effects was analyzed in detail. It was found (see Fig. 9-28) that a single segment of a generator using KD*P to frequency-double $1.06\text{-}\mu\text{m}$ light is ultimately capable of an average power approaching 100 kW. With a relatively modest number of segments, very high average powers can be obtained over the total aperture.

For doubling wavelengths below $1.06 \mu\text{m}$, the absorption of KD*P is considerably lower than at $1.06 \mu\text{m}$, where it is 0.6 m^{-1} , so that the ultimate average power per segment for doubling to the blue or ultraviolet is much higher than 100 kW. A precise value cannot be given without accurate absorption data; however, a value typical for transparent window materials, 0.03 m^{-1} , permits the ultimate average power per segment to be greater than 1 MW.

Temperature gradients cause thermal dephasing of the waves and disturb phasematching. The temperature variation in the crystal (Fig. 9-26) can be characterized by the longitudinal temperature rise T_l , associated with the flow of heat outwards to the crystal surfaces, and the transverse temperature rise T_{tr} , associated with the temperature rise of the coolant as

Fig. 9-26. The temperature profile in a thin plate gives rise to dephasing and surface stress.

Fig. 9-27. A gas-cooled harmonic generator maximizes average power.

Fig. 9-28. More than 50 kW of harmonic output is achievable from a single segment of KD*P at 1.06 μm .

it traverses the crystal face. Now, the total dephasing budget is constrained by the requirement of high conversion efficiency. This budget contains six terms, which arise from T_i , T_{tr} , the optical dispersion in the coolant, the angular divergence and line width of the drive, and the index inhomogeneity in the crystals. Of these parameters, only the optical dispersion in the coolant and T_{tr} depend on the plate separation. This separation is therefore uniquely determined by the trade-off between these two terms. With the plate separation fixed, the thermal part of the total dephasing then includes only two terms, one arising from T_i and the other from the optimized combination of the optical-dispersion and T_{tr} terms. Because the average power is related to T_i , highest average power is obtained by minimizing the transverse terms. Following the analysis and notation of Ref. 19, the transverse terms are given by the transverse figure of merit Em :

$$Em = \frac{C_p v l_c}{32 K \kappa_t} \quad (15)$$

and their relative size is given by $BN^{1/2}$, where

$$BN^{1/2} = \left(\frac{N}{Em \delta_{T0}} \right)^{1/2} \quad (16)$$

Here N is the number of plates that would be required if the transverse terms were zero, and

$$\delta_{T0} = \delta - \delta_\theta - \delta_\lambda - \delta_n \quad (17)$$

is that part of the total dephasing budget available to the longitudinal dephasing, assuming that the transverse terms are zero. The effect of the transverse terms is of the order of $BN^{1/2}$.

The precision with which the flow speed must be controlled was also addressed. As the speed of the coolant is reduced, both the cooling rate h and the transverse figure of merit Em decrease. The effect can be visualized in Fig. 9-28 as a downward translation of both the thermal dephasing and the boundary layer temperature rise lines. These constraints depend in different ways on the flow speed; in some situations their dependence on v can be described approximately with an effective power law. This power is generally not great; depending on which effect is limiting it, the average power is either independent of v , or drops moderately as v is reduced. At some point the Reynolds number becomes low enough for the flow to become laminar, and the average power then drops discontinuously with v by at least an order of magnitude.

Overall, the prospect for nonlinear optics at high average power is very

encouraging. The heat load in a nonlinear material with moderate optical absorption and nominal thermomechanical properties can be handled gracefully with turbulent-gas cooling. This conclusion is not limited to second-harmonic generation, but appears to generalize to any optical-energy-conserving process that involves frequency conversion. If a nonlinear material with nominal properties is capable of high efficiency on a pulse-by-pulse basis, then it will also scale to high average power.

High-Average-Power Electro-Optic Switching

Electro-optic switches rotate the plane of polarization of a light beam through an angle whose value depends on the voltage applied to the crystal. There are two realizations of this technique, the longitudinally excited Pockels cell (LEPC) and the transversely excited Pockels cell (TEPC). In the LEPC the electric field is parallel to the beam direction and is applied to the crystal using plasma electrodes that have high optical transmission. The TEPC uses the standard electrode design. The voltage requirements of the two cells are similar. Thermal management is virtually nonexistent in present-day devices, and no cooling is used. The LEPC has the advantage that it can switch large apertures in a few nanoseconds and that it has no discernible elasto-optic transients of the type found in large-aperture TEPCs. Electro-optic switches of either design have high transmissions ($>98\%$), impose no limits on pulse length or position, and allow precise, real-time polarization control through servomechanisms. They also have the advantages of zero chirp and zero beam steering or slewing.

At high average power, optical absorption causes thermal gradients that couple to the birefringence and reduce the uniformity of the output polarization over the

beam aperture. The temperature sensitivity of the polarization angle is high, and precise temperature control is required to maintain performance. A detailed study of the thermal gradients in the presence of absorptive heating and gas cooling has been carried out, using experience and results from the earlier study¹⁹ on high-average-power harmonic generation. Two designs for cooling were studied, in which the principal thermal gradients are either transverse to the beam direction or parallel to it. The longitudinal-gradient (LG) and transverse-gradient (TG) designs are illustrated in Figs. 9-27 and 9-29. The TG design avoids the degradation of beam quality associated with passing through the coolant. But the entire temperature variation across the plate, from one cooled side to the other, must remain within the temperature bandwidth of the device. In the LG design the beam passes through the coolant, which must be a thin layer of high-speed gas to avoid increasing the beam divergence. But the LG design allows the temperature variation from one cooled face to the other to be averaged, so that only this average must remain within the temperature bandwidth. Thus the LG design is ultimately capable of the highest average power.

The LG design can handle about 10 times more average power than the TG design can handle, but even LG Q switches are about an order of magnitude more sensitive to thermal gradients than harmonic generators. A thorough study of the thermal limits on Q switches, using the appropriate figures of merit for the various available materials, showed that this thermal sensitivity limits the average optical power to between 1 and 10 kW.

The dominant contribution to the temperature variation comes from the temperature rise of the coolant gas as it crosses the face of the crystal. (See Fig. 9-26.) This temperature rise can be reduced by increasing the thermal transport of the flow, going to higher pressure and turbulent flow. Ultimately, the heat transport is limited by turbulent dephasing in the coolant gas. We cannot state precisely what the turbulent dephasing will be, although it can be estimated readily. However, it would clearly be advantageous to find a design that allowed the coolant-gas

Fig. 9-29. The transverse-gradient Q switch or harmonic generator is more sensitive to thermal profiles in the crystal.

temperature rise to exceed the temperature bandwidth of the crystal. A technique for accomplishing this has been invented, and it allows the average power to increase by more than an order of magnitude. One realization of this approach is described here.

With current materials and technology, a Q switch can be constructed that can handle 100 kW of average power. It is an LG TEPC using plates made from KD*P, cut as shown in Fig. 9-30, followed by a phase-error compensator. The compensator is a mechanically and thermally identical device using KD*P plates oriented at 90° to those of the first cell. The compensator is constructed so that its temperature profile is identical to that of the switch. However, because of the cut of the crystals, the temperature coefficients of depolarization for the compensator and switch are exactly equal in magnitude and opposite in sign. In the absence of higher-order effects, the thermal depolarization in the switch is exactly reversed by repolarization in the compensator. This cancellation occurs for arbitrary temperature profiles in the switch and compensator, and therefore occurs regardless of the spatial variation in the beam intensity or the design of the flow cooling. It requires only that the two elements be mechanically and thermally identical. In the absence of other effects, the average optical power the Q switch can handle would be limited only by mechanical and thermal limits on precision engineering.

In practice, the average optical power is limited by stress-induced fracture or by the coupling of the temperature profile to the birefringence through the stress-optic effect. Stress-induced fracture occurs in KD*P at average powers between 50 and 150 kW, depending on the quality of the crystal surface. Ohmic heating of the crystals by electric-field cycling is slight compared with the absorptive heating. The stress-optic coupling effect has been evaluated: in KD*P, the birefringence should be an order of magnitude less sensitive to the thermally induced stresses than to the direct dependence of the refractive indices on the temperature.

Several further points are worth noting. (1) Only the anisotropy in the stress-optic effect causes the polarization rotation to

Fig. 9-30. The high-average-power Q switch uses a compensator for thermal depolarization.

be different in the switch and compensator. (2) Only the transverse temperature gradients and the anisotropic part of the stress-optic tensor contribute to this anisotropy. (3) Linear temperature gradients produce no stress, and therefore produce no stress-optic effects. From this final point we conclude that only a very weak stress-optic effect is associated with the temperature rise of the coolant gas, because the associated temperature gradient is almost constant.

All the transverse temperature gradients are associated with beam nonuniformity. We anticipate that the small-scale spatial variation in the intensity is weak, because the beam is assumed to be close to diffraction-limited. Then we can reduce the transverse temperature gradients by making the plate thin compared with the scale of the transverse variations in beam intensity. One expects that the greatest transverse gradients will occur at the spatial edges of the beam, where the intensity is dropping to zero. These gradients can be reduced by using thin plates or by using a small auxiliary heating wire to even out the temperature profile at the edge of the beam. The stress-optic coupling is already an order of magnitude smaller than the direct coupling of temperature to birefringence, and there are apparently several means of reducing the effects still further. A conservative approach to a high-average-power Q switch suitable for an optical radar uses five $15 \times 3 \times 0.3$ -cm plates in both the switch and the compensator. The coolant-gas temperature rise is

~1°C, and the switching voltage is then ~12 kV. Less conservative designs would have fewer plates, at the expense of a greater coolant-temperature rise.

We conclude that compensation will be inexact only to the extent that fabrication and construction tolerance limits are reached, or that transverse-gradient-induced anisotropy in the index profiles is present. It appears that, in KD*P, the limits on average power derived from each of these effects are comparable with the limit due to stress-induced fracture. Ultimate limits will be determined only after this design is reduced to an operating prototype, but it is clear that a compensated Q switch is capable, in principle, of average powers in the 100-kW range.

High-Average-Power Nonlinear Materials

Operational devices using nonlinear optics are constrained by the availability of suitable nonlinear optical materials. In many military applications the short engagement times demand high-average-power capability. Concepts currently under consideration envision utilizing frequency conversion and optical switching at physical sizes considerably greater than current technology has reached, and in environments that are potentially much more stressful. The effectiveness of these concepts ultimately rests on the availability of materials that can operate efficiently at scale in the required environments.

The paramount importance of materials in realizing any proposed laser system concept has been clearly recognized at LLNL. From the outset, the development

of effective optical materials has been an essential part of all laser programs, and it continues to be emphasized in the new initiatives we have undertaken.

Research on nonlinear materials at LLNL is a broad-based program directed towards identifying and developing materials and techniques for perceived future national needs. This work involves both basic research on the optical properties and crystal structure of nonlinear materials and applications engineering, in which the needs of specific missions are addressed. Here we describe our work under this program on high-average-power frequency conversion.

Our analysis of high-average-power frequency conversion showed that even very weak optical absorption can result in thermal gradients that can reduce conversion efficiency or can even fracture the crystal. There are four figures of merit for average-power materials:

1. The threshold power P_{th} . The conversion efficiency is high only if P_{th} is lower than the total laser power divided by the square of the etendu. (The etendu is the ratio of the beam divergence to the beam divergence of a diffraction-limited laser of the same aperture.)
2. The optical damage threshold J_d .
3. The fracture temperature ΔT_f , which depends linearly on the fracture toughness S_0 . If the temperature difference between the central plane and the surface planes of a thin-plate harmonic generator exceeds ΔT_f , the plate will crack.
4. The thermal dephasing parameter $Wm = 8K/\alpha\beta_t$, where K is the thermal conductivity, α is the linear optical absorption coefficient, and β_t is the temperature sensitivity of the phasematching.

Table 9-5. High-average-power figures of merit for nonlinear materials at 1.06 μ m.

Material	P_{th} (MW)	J_d (J/cm ²)	K (W m ⁻¹ K ⁻¹)	α (m ⁻¹)	β_t (m ⁻¹ K ⁻¹)	Wm (W·mm)	S_0 (MPa·m ^{1/2})	ΔT_f (°C)
KDP	74	7	2	3	61	87	0.09	7.0
KD*P	67	7	2	0.6	53	503	(0.09) ^a	—
CD*A	ncpm ^b	(7)	(2)	2	(200)	[40] ^c	(0.09)	—
MgO:LiNbO ₃	ncpm	10	5.6	0.1	1122	400	(0.14)	3.8
LiIO ₃	135	2-4	1.3	0.1	17	6120	0.3	—
KTiOPO ₄	0.05	(10)	(2)	0.1	14	11 430	—	—
D-LAP	30	15	(2)	0.1	(15)	[10 000]	0.25	—

^a Quantities in parentheses are estimates.

^b Noncritically phasematched.

^c Quantities in square brackets are calculated from estimates.

This parameter measures the extent to which thermal gradients degrade the conversion efficiency.

These four parameters determine the performance of a given material in a particular application. We have determined them for several common nonlinear materials; the data are summarized in Table 9-5.

These data indicate that KD^*P and lithium iodate are capable of very high average power (approaching 100 kW) in a single aperture. Other materials, such as lithium niobate and CD^*A , cannot handle so much optical power because their temperature sensitivities are greater.

In summary, we have begun to determine the figures of merit of nonlinear materials for high-average-power applications. The most common nonlinear materials appear to be capable of very high average power. This conclusion is not limited to frequency-doubling $1.06\text{-}\mu\text{m}$ light, but applies by a simple extension to any nonlinear optical process.

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Thermal-Stress Optics Modeling

Introduction

The articles in this subsection describe thermo-optical modeling of solid slab lasers capable of high-average power output and high beam quality. The physical effects that must be controlled to advance this technology are those that govern thermal-stress fracturing, thermal focusing, stress-induced birefringence, and depolarization. That is, thermal gradients, thermo-elastic stresses, and material deformations must be controlled in the flashlamp-pumped, convectively cooled slab systems under investigation. Numerical modeling is an effective research and design tool for this application because a multitude of physical parameters, boundary conditions, and material properties, which may be difficult to designate and/or control in the laboratory, can be specified

precisely, varied with ease, and diagnosed in depth in computations.

During 1985, we developed a suite of 2- and 3-D thermal stress codes (TECATE, TOPAZ, and NIKE) and an optical ray trace code (BREW), and applied them to zig-zag slab systems. We considered both isotropic and anisotropic glass and crystalline laser host media because we believed the thermo-optical properties of some candidate host crystals might prove to be better than those of laser glasses.

The zig-zag laser was originally conceived to eliminate those thermal and stress-induced focusing and depolarization effects that degrade both output power and beam quality in straight-through slab and rod lasers. The zig-zag beam path is designed to traverse the slab interior by a series of TIRs from the flashlamp-pumped, convectively cooled slab surfaces. The purpose of the zig-zag beam trajectories is to sample thermal and stress gradients and surface deformations in such a way that accumulated optical path distances are constant and that undistorted output polarizations are obtained for every ray in the exit beam. Thermo-optical modeling of such systems must therefore accurately describe all source-function and boundary effects that may generate significant wave front phase and polarization deviations in zig-zag laser beams.

The modeling work described below accounts in detail for variations in thermal source functions due to (1) flashlamp absorption distributions within slabs and (2) ASE absorbers that are in contact with either slab edges or contiguous elastomers. Optical ray tracing is used to evaluate wave front quality and thus to develop optimal slab/pumping dimensions and optimal coolant temperature profiles on TIR surfaces, at slab ends, and in chill bars. We find that optical path differences can be sensitive to all of these physical and geometrical factors, as well as to associated factors of material thermal/optical properties, incident beam angles, slab wedge angles, and exact numbers (odd vs even) of zig-zag beam traversals.

The first article below presents an overview of our thermo-optical modeling codes and results from baseline design work.

The second article describes the ray tracing code BREW and selected benchmark problems. The third article describes the 2-D thermal stress code TECATE, which was developed for efficiency of execution with large grid meshes. Dynamic alternating direction implicit (DADI) methods are used to solve stress potential equations in TECATE for both isotropic and anisotropic zig-zag slabs.

Author: R. J. Gellinas

Thermal-Stress Optics Modeling of Zig-Zag Slabs

Introduction. To quantitatively design zig-zag slab systems for operation at high average power, we have developed a suite of logically linked 2- and 3-D computational physics codes (Fig. 9-31) and applied them to benchmark designs. System designs capable of producing high beam quality require the full range of these codes to identify those flashlamp-pump intensity distributions, pump dimensions, slab dimensions, heat-flux distributions, and coolant flow boundary conditions over all portions of the slab that can minimize optical path differences (OPDs) in the output beam.

Flashlamp-pump intensity contributes two thermal source functions to the governing of beam quality at high average power: (1) the distribution of absorbed flashlamp radiation that is available as heat in the neodymium-doped glass slab and (2) the heat generated in the ASE absorbers as a result of transmission of both ASE and flashlamp radiation through the slab's side edges and elastomer into the highly absorptive edge cladding.

Physics. Beam phase distortions are measured in terms of OPDs that are proportional to path integrals of the refractive index $n(x, y, z)$ over beam trajectories. Such path integral quantities may vary significantly for different rays in the beam due to nonuniform deformations of slab surfaces that can be both mechanically and thermally induced and to non-uniform index of refraction variations Δn that occur as consequences of both thermal gradients ($\Delta n \sim (dn/dT) \nabla T$) and thermal and mechanical stresses [$\Delta n \sim f(\pi \cdot \sigma)$]. Beam depolarization is a function of principal stresses σ_p in both glass and crystalline slabs.

Computational Codes. Figure 9-31 presents an overview of our suite of codes, illustrating their logical connections.

SLAB. This is a collection of subroutines that defines slab doping, gain, losses,

Fig. 9-31. A suite of 2- and 3-D design/analysis codes is used to model wave front quality in high average-power-slab lasers.

lifetimes, storage and extraction parameters, flashlamp operating parameters, heat removal, system energy budget, and cavity properties. These factors are major determinants of thermal source functions and heat removal in thermal-stress evaluations.

SOURCE. This set of routines simulates flashlamp thermal source functions for the heat transfer codes TOPAZ and TECATE.^{22,23} Shading functions for radiation intensities near boundaries within the slab are computed for various slab, reflector, and mounting configurations.

TOPAZ, NIKE. These are multipurpose thermal and stress engineering codes^{22,24}

that have been adapted to advanced laser applications for both steady-state and transient operating conditions. Complicated system geometries with multiple regions, multiple materials, and a variety of loading boundary conditions are solved with these codes. Phase-change and slide-line options are available for both elastic and inelastic material models.

ORION. This graphics postprocessor was originally developed in conjunction with the TOPAZ and NIKE codes; adaptations that permit processing of TECATE data have been developed. Both 2- and 3-D versions of ORION are used to plot

Fig. 9-32. (a) Computed 3-D deformations (magnified by 10^4) in a pumped glass zig-zag slab. (b) Contours of axial stress on the slab surface show potential fracture sites in pump transition region.

contours, time histories, and deformed shapes in color and in black and white.

TECATE. This 2-D thermal stress code²³ is used for efficient, dedicated analysis of high-average-power slabs fabricated from solids with arbitrary material symmetry properties. Steady-state and transient, multi-region (e.g., slab/elastomer/ASE absorber combinations) problems are solved by a DADI method in the TECATE code. Stable, high-order-accurate solutions are obtained with high numerical grid resolution.

BREW. This 3-D geometrical optics code computes the OPDs and fringe patterns for slab output wave fronts. BREW uses ray tracing methods to evaluate beam propagation in both zig-zag and straight-through slabs as well as the effects of beam entrance and exit refraction on undeformed surfaces. (See the next article.) Thermal lensing and stress birefringence are computed for arbitrarily deformed slabs with arbitrary stress and temperature distributions. Means for treating depolarization effects are under development.

Results. We have obtained both steady-state and transient results in benchmark designs of LSG-91H flashlamp-pumped slabs. Figure 9-32(a) shows greatly magnified 3-D deformations that were computed by the TOPAZ and NIKE codes on an $11 \times 21 \times 47$ numerical grid (height, width, length). The deformations relative to the original unpumped slab are typically on the order of several microns. The intrinsically 3-D nature of the problem, prior to design optimizations, is clearly evident.

Figure 9-32(b) presents axial stress contours for this same glass slab, in which total pump power is $1600 W_e$ and all coolant and ambient temperatures are $20^\circ C$. The maximum tensile stresses occur near the edges in the transition region between the pumped and unpumped slab areas. These computed locations are consistent with morphologies of stress fracture events that have been observed in laboratory experiments. OPDs for this unoptimized pumped slab system appear in Fig. 9-33(a); they are clearly dominated by surface deformations. By optimizing the thermal distributions, beam quality can be improved greatly, as shown by the results in Fig. 9-33(b).

Certain applications may require highly controlled beam quality at all times of operation, from startup to steady state. Figure 9-34 presents transient computations of thermal stresses sampled over all time regimes of operation. In this example, stress gradients are small over 80% of the slab area, which is essential to minimizing the effects of depolarization at all times of interest.

Summary. In 1985 we made the following advances in thermo-optical modeling:

- Developed 2- and 3-D thermal-stress optics codes that include source physics to model and design zig-zag slabs with high beam quality.
- Computed the effects of 3-D stress fracture sites in zig-zag slabs.
- Improved beam quality in zig-zag slabs using optimized thermal designs.

Fig. 9-33. (a) Computed OPD in a ten-bounce zig-zag slab, with a total lamp power of $1600 W_e$. (b) Improved thermal/optical design, with a total lamp power of $500 W_e$.

Fig. 9-34. Transient calculation of horizontal stress component, σ_{xx} . Results are shown at selected pulses.

• Identified sensitive physical processes and assessed them in conjunction with laboratory experiments.

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BREW, An Optical Ray Tracing Code for the Design of Zig-Zag Slabs

Introduction. We have begun to develop a 3-D numerical code for birefringent ray tracing of electromagnetic waves (BREW) that computes wave-front phase and polarization distortions in slab lasers. This optical code uses thermal, stress, and displacement data from companion 2-D and 3-D thermoelastic codes (TECATE, TOPAZ, and NIKE²⁵⁻²⁷) for both anisotropic and isotropic media. The physical processes described in BREW include ray reflection,

bending, and refraction from deformed surfaces. We believe these effects and processes are the major determinants of beam quality in the thermo-optical design of zig-zag slabs. This article describes the present state of development of the BREW code.

Code Description. Figure 9-35 shows the geometry of the zig-zag laser slabs that are to be modeled. While such other materials as elastomer ASE cladding, evanescent coatings, and chill bars may be attached to various surfaces of the slab and are modeled thermoelastically, the BREW code traces rays only in the laser slab and in its ambient atmosphere. We consider slabs with arbitrary dimensions and wedge angles $\alpha \leq \pi/2$ in the Cartesian coordinate system shown in Fig. 9-35.

The height of the slab between TIR surfaces OPQR and O'P'Q'R' is denoted by h , slab width by w , and slab length by l . Slab temperature, stress, and displacement data used in BREW are abstracted from larger data sets obtained in thermal stress computations for the state of the slab. We

Fig. 9-35. Zig-zag laser slab in Cartesian coordinate system.

perform ray tracing computations by assuming that light rays incident upon the slab in the plane $x = \text{const.}$ will remain in that plane; thus, beam walk-off effects are temporarily neglected. Coding of depolarization is in early stages of development.

For a typical plane at $x = OO''$ (Fig. 9-35), a beam of light rays is defined as a collection of orthogonal trajectories to a traveling plane wave front, with a reference point at some negative value of z . After refraction at the left beveled face of the slab, the beam traverses the slab—making a total of $2m$ TIR bounces at the upper and lower slab surfaces—and exits at the right beveled face. Its phase front after exit from the slab is calculated at a plane parallel to the initial plane wave surface at a known distance from the right beveled face of the slab. Optical path lengths in multiple-pass interferometers are computed as the multiple of the optical path length obtained for a single pass. Phase-front distortions of emergent electric fields are analyzed by superposition upon the incident field, ignoring changes in field amplitudes caused by depolarization. Fringe patterns are plotted from loci of zero intensities, which are essentially isocontours of optical path differences in wavelength units.

Optical Path Length Calculations. A typical y - z cross section $O''P''Q''R''$ of the ideal undeformed slab is shown in Fig. 9-35. The condition that each ray incident

on $O''R''$ make exactly $2m$ TIR bounces before emerging at plane $P''Q''$ is

$$l = 2mh \tan \psi. \quad (18)$$

When provided with input values for all but one of the quantities l , m , h , or ψ , the BREW code determines the unknown quantity and then performs a set of internal consistency tests. For example, the TIR condition $\sin \psi > n_3/n_2$ must be satisfied, where n_1 , n_2 , and n_3 are indices of refraction for air, slab, and coolant, respectively. Incident beam angles θ , which satisfy $n_1 \sin \theta = n_2 \sin (\psi - \alpha)$, must be real. Moreover, incident rays that would reflect back into the deformed face or inside the slab from the bottom surface $O''P''$ are identified and rejected.

Unperturbed optical path lengths $\int n_0 ds_0$ between reference incident and exit planes are identical for all rays in the ideal undeformed slab. The optical path length for variable temperature, stress, and material deformations is partitioned for a given ray as

$$\text{OPL} = \text{OPLT} + \text{OPLS} + \text{OPLD}, \quad (19)$$

where $\text{OPLT} = (\partial n / \partial T) \int T ds$, $\text{OPLS} = 1/2 \int (n_{11} + n_{22}) ds$, and $\text{OPLD} = \int n_0 ds$. The integration path increment ds corresponds to slab surface deviations from s_0 . The coefficient $\partial n / \partial T$ is a material constant, the eigenvalues n_{11} and n_{22} are derived from the π - σ perturbation to the impermeability tensor, and the path integrations extend over the bent ray paths s according to the eikonal equation. The ray trajectories extend to and are reflected from or refracted by the deformed slab surfaces. Redirecting the beam at deformed surfaces is essential to producing results that are closer to the observed results than simpler approximations that neglect such effects.

Numerous code validations have been performed against benchmark problems with known solutions. For example, the optical path differences due to $\partial n / \partial T$ have been computed and verified in a flashlamp-pumped Nd:glass slab. Similarly, sample problems with prescribed π - σ perturbations and prescribed surface deformations have been validated

satisfactorily. Full-scale analyses of optical fringe data obtained for convectively cooled, flashlamp-pumped zig-zag slabs are in progress.

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TECATE, A Dynamic Alternating Direction Implicit Code for Thermal Stresses in Isotropic and Anisotropic Media

Introduction. Thermal gradients and stresses adversely affect both the phase and polarization quality of high-average-power laser output beams. Thermal gradients can cause significant focusing effects in materials with temperature-dependent indices of refraction, while thermally induced stresses can introduce optical birefringence into isotropic and anisotropic glass and crystalline laser media. Such thermo-optical effects are sensitive to many physical and system-geometry factors that, although difficult to isolate experimentally, can be analyzed effectively by numerical modeling. Modeling of multi-dimensional thermo-optical effects in high-average-power solid lasers requires numerical accuracy and high computing speed.

To solve the heat and stress partial differential equations, we have applied the

DADI method to the totally Eulerian code for anisotropic thermoelasticity (TECATE). Higher order spatial discretization methods have been used in the TECATE code. These numerical methods readily solve mixed Dirichlet, Neumann, and film boundary conditions.

This article describes the prominent features of the 2-D TECATE code and presents results of sample thermal test problems in both isotropic and anisotropic media. TECATE code results for practical experimental systems are also discussed in the two preceding articles.

Thermoelasticity in Anisotropic Media.

The TECATE code solves the equations of thermoelasticity for the slab geometries shown in Fig. 9-36. Zig-zag slab systems are considered exclusively in this article, although other slab systems can frequently be analyzed with the logical framework of zig-zag slab codes. Beam phase distortions are measured in terms of optical path lengths, $\int n ds$, where n is the refractive index and s is distance along the optical path. Refractive indices are sensitive to temperature variations δT within the slab and to thermoelastic stresses σ_{mn} (elements of the second-rank stress tensor σ), where $m, n = 1, 2, 3$. The time-dependent heat equation for anisotropic media is

$$\rho c \frac{\partial T}{\partial t} = \nabla \cdot (\mathbf{K} \cdot \nabla T) + Q, \quad (20)$$

where T = temperature (K); ρ = mass density (kg/m^3); c = heat capacity ($\text{J/kg} \cdot \text{K}$);

Fig. 9-36. (a) A 3-D slab that is practically infinite in the z -direction. (b) A cross sectional cut Ω perpendicular to the z -axis.

$\mathbf{K}$ = the symmetric, second-rank thermal conductivity tensor ($\text{W/m} \cdot \text{K}$); and Q = rate of internal heat generation (W/m^3).

The stationary 2-D problems considered in this article are always written in terms of the coordinates (x, y) , and ∇ denotes the reduced 2-D vector operator $(\partial/\partial x, \partial/\partial y)$. We use $\bar{\mathbf{K}}$ to denote the appropriate two-by-two submatrix of $\mathbf{K}$, namely

$$\bar{\mathbf{K}} = \begin{pmatrix} k_{11} & k_{12} \\ k_{12} & k_{22} \end{pmatrix}. \quad (21)$$

The product $\mathbf{K} \cdot \nabla T$ is the heat flux and is interpreted as the vector

$$\mathbf{K} \cdot \nabla T = \begin{pmatrix} k_{11} T_x + k_{12} T_y + k_{13} T_z \\ k_{12} T_x + k_{22} T_y + k_{23} T_z \\ k_{13} T_x + k_{23} T_y + k_{33} T_z \end{pmatrix}. \quad (22)$$

The heat flux vector in an anisotropic medium will not be normal to the isotherms $T = \text{const}$, as would be the case in an isotropic medium. We assume that ρ , c , and the k_{ij} are all constants, although this assumption can be relaxed. The stationary heat equation is simply given in divergence form as

$$\nabla \cdot (\mathbf{K} \cdot \nabla T) = -Q. \quad (23)$$

To compute thermal stresses and material deformations in advanced laser applications, the full stress tensor σ must be evaluated in solids with arbitrary symmetry properties. The stress potential formulation is applied because it yields high accuracy in the stress variables for the freely suspended laser slabs that are of interest in our work. In the compact vector notation discussed by Nye,²⁸ Hooke's law for small deformations is written as

$$\epsilon = \mathbf{S} \cdot \sigma + \alpha T, \quad (24)$$

where ϵ denotes material strains, σ denotes stresses, α represents thermal expansion coefficients, and $\mathbf{S}$ is the material compliance matrix (s_{ij}) . The stress potentials ϕ and ψ are defined in the laboratory $(x-y)$ frame by

$$\begin{aligned} \sigma_1 &= \frac{\partial^2 \phi}{\partial y^2} & \sigma_2 &= \frac{\partial^2 \phi}{\partial x^2} & \sigma_6 &= -\frac{\partial^2 \phi}{\partial x \partial y} \\ \sigma_5 &= \frac{\partial \psi}{\partial y} & \sigma_4 &= -\frac{\partial \psi}{\partial x} \end{aligned} \quad (25)$$

for the conditions of 2-D generalized plane strain, which are frequently applicable in examples of interest. Equation (25) satisfies the balance of force equation, $\nabla \cdot \sigma = 0$, identically. Using Hooke's law, the following system of fourth-order partial differential equations is obtained for ϕ and ψ :

$$L_4 \phi + L_3 \psi = M_2 T, \quad (26a)$$

$$L_2 \psi + L_3 \phi = M_1 T, \quad (26b)$$

where the operators L_i and M_j are linear combinations of terms containing mixed differential operators of fourth and lower order as well as material constants.²⁹ For efficient numerical solution, the $x-y$ plane is mapped onto the $\zeta-\eta$ plane so that the parallelogram Ω is mapped onto a unit square in $\zeta-\eta$. The free-surface boundary conditions for ϕ and ψ are

$$\begin{aligned} \phi &= 0, \\ \frac{\partial \phi}{\partial \hat{n}} &= 0, \quad \text{and} \\ \psi &= 0, \end{aligned} \quad (27)$$

where $\hat{n}$ is the outward unit normal to the boundary of the unit square. It can be shown that the operators L_4 and L_2 in Eq. (26) are elliptic and that the DADI method is therefore an attractive one for numerically solving both the stress potential equations and the heat equation. (These operators remain elliptic in the plane of the unit square.) A 5×5 finite-difference stencil centered at some grid node (i, j) is used for the differencing formulae.²⁹ Material stresses are obtained from Eq. (25), and displacements are obtained by integration of the material strains ϵ according to Hooke's law.

The DADI Solution Method. The DADI solution method closely follows the logic of standard ADI methods, but differs from them by seeking systematically to accelerate the convergence of the solution process. Briefly, ADI methods are used to find the solution of finite-difference or finite-element equations of the form

$$Lu = f, \quad (28)$$

where L is a linear elliptic operator and f is a given function. Assume that L can be

written as the sum $L = A + B$. The Peaceman-Rachford³⁰ approach for solving Eq. (28) is applied heuristically as follows: $Lu \equiv (A + B)u = f$ is converted into the evolution equation

$$\frac{\partial v}{\partial t} = Av + Bv - f, \quad t > 0, \quad (29)$$

where v is some initial guess u^0 at $t = 0$ whose steady state ($t = \infty$) solution is obviously a solution of $(A + B)u = f$. One discretizes Eq. (29) in time, solving odd-numbered steps (which we designate as $n + 1$, because $n = 0$ represents initial values) implicitly in A and explicitly in B (a process we call the x -sweep) according to

$$\frac{u^{n+2} - u^n}{\Delta t} = Au^{n+1} + Bu^n - f. \quad (30)$$

Then, reversing the process, even-numbered steps ($n + 2$) are solved explicitly in A and implicitly in B (a process we call the y -sweep) according to

$$\frac{u^{n+2} - u^{n+1}}{\Delta t} = Au^{n+1} + Bu^{n+2} - f. \quad (31)$$

It is unnecessary to find the true transient solution in ADI methods; the objective is simply to reach the steady state as quickly as possible.

The innovation of DADI as compared with standard ADI is the development of a strategy that automatically changes the step size Δt dynamically. In this way, DADI should be able to: (1) keep Δt within regions of fast convergence; (2) recognize and suppress instabilities by decreasing Δt as instabilities start to develop; and (3) avoid the burdensome and less efficient a priori choice of iteration parameters required in the standard ADI method. Such a DADI strategy has been applied in this work based on numerical convergence factors that are described in a separate report.²⁹ This scheme has been shown to be unconditionally stable for the heat equation.²⁹

Results of Test Problems. The DADI method, along with appropriate subroutines for thermal source functions in flashlamp-pumped, convectively cooled solid slabs, has been validated against analytic test problems,²⁹ other finite element codes,^{31,32} and benchmark experimental applications.³³ We consider here two

Fig. 9-37. Slab temperature contours ($^{\circ}\text{C}$) in pumped GSGG slab.

benchmark design calculations for pumped Nd:Cr:GSGG (cubic crystal symmetry) and Al_2O_3 (trigonal symmetry) slabs.

Nd:Cr:GSGG Zig-Zag Slab. In our first example, an $8 \times 18 \times 0.7$ -cm slab of gadolinium scandium gallium garnet (GSGG) is pumped with 270 W/cm^2 of flashlamp power. Dopant densities are $[\text{Nd}] = 2 \times 10^{20} \text{ cm}^{-3}$ and $[\text{Cr}] = 1 \times 10^{20} \text{ cm}^{-3}$. The convective cooling power removed by water flowing over the TIR faces is 11.5 W/cm^2 , and the Brewster faces at the slab ends are exposed to room-temperature air. Maximum slab deformation magnitudes were approximately $3 \mu\text{m}$ near the Brewster faces exposed to air outside of the pumped and convectively cooled regions in section B of Fig. 9-37. Computer slab temperatures in contour g of section A (Fig. 9-37) show the

effect of a 2°C temperature rise in the coolant, which flows across the slab surfaces from left to right. These temperature rises can also affect beam quality over the entire slab area.

Figure 9-38 presents computed principal stress vectors in the x - y plane. Heavy lines denote tension, and light lines denote compression. These vector orientations are of considerable interest, because beam depolarization tends to be greatest in regions where principal stress rotation angles are large. Such regions of depolarization occur only near the slab side edges in the present baseline design.

Material Deformations in Trigonal Al_2O_3 . The combined effects of crystal symmetry orientation in slabs and various flashlamp pumping directions are examined in our second design calculation.

Fig. 9-38. Principal plane stress vectors in pumped GSGG slab.

Fig. 9-39. Magnified deformations in Al_2O_3 (trigonal) slabs computed for alternative orientations of pump vs crystal symmetry axes. Crystalline axes are denoted by a_1 , a_2 , and a_3 .

When crystals are cut in shapes and orientations that do not correspond directly to crystal symmetry conventions,³⁴ then material property tensors for thermal conductivities, thermal expansion coefficients, and elastic compliances must be transformed from their conventional frames of reference into the appropriate laboratory frame. Figure 9-39 represents the conventional crystal-axis orientations²⁸ for the trigonal crystal Al_2O_3 by a_1 , a_2 , a_3 , and c , where c is normal to the plane of a_1 , a_2 , and a_3 . Computed deformations, with artistic rendering, are shown on the slab faces in the lower part of Fig. 9-39 for three alternative pumping orientations relative to the crystalline axes. Readily apparent in these results is a preferred orientation that maintains purely planar deformations in the x - y plane, in contrast to the generally nonplanar deformations in the other two orientations. These examples were computed with spatially uniform pumping of the top and bottom slab surfaces under conditions of generalized plane strain.³⁵

The observed orientation and symmetry-sensitive deformations are due strictly to the thermoelastic responses of the anisotropic Al_2O_3 medium. Apparently, such deformations can be significant determinants of optical beam quality in anisotropic media. A 51×81 computational mesh was used in the TECATE computations on a Cray-1S computer. The numerical solutions converged in approximately 30 DADI iterations.

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Thermal Design

Facilities

Introduction. In support of HAP project objectives, laboratory facilities were constructed to control the thermal environment

Thermal Design

of laser slabs and to characterize the thermal aspects of laser operation. Data from these facilities provided the basis to

- Evaluate stress-optics codes.
- Update models of energy distribution in a laser cavity.
- Refine the test bed designs.

In 1985, the following components and systems of the thermal test facilities were activated and brought into full operation. We fabricated and assembled the zig-zag and gas-cooled slab (GCS) test units, and both units were used for thermal experiments. The coolant systems, which were installed and operated to remove waste heat from the test units, include temperature-controlled, high-cleanliness water circuits and a high-volume (1 kg/s) open-loop

nitrogen gas system. We developed a computer-based data acquisition system to measure coolant temperatures, flow rates, and component temperatures. A Hughes Probeye thermal-imaging camera monitors the temperature distribution of the slab in the GCS test unit and the pump uniformity of the zig-zag flashlamp reflectors. Finally, we constructed a calorimeter to test the thermal life of flashlamps, coatings, and optical elastomers.

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Zig-Zag Test Unit. The zig-zag test unit shown in Fig. 9-40 was used extensively for thermal, optical quality, and gain experiments in 1985. It represents a significant

Fig. 9-40. Zig-zag test unit.

departure from previous LLNL solid-state lasers in its configuration, operation, and use of materials.

The test bed is composed of a 32- × 10- × 1.5-cm LSG-91H laser slab supported by a cast optical elastomer. In 1985, three of these slabs were used. Each slab is mounted in a vertical orientation to avoid gravity-induced sag. The slabs are face pumped with water-cooled flashlamps in a "birdhouse" reflector cavity designed by E. J. Goodwin. The Rabl reflectors used in the initial experiments were slightly more efficient, but produced a less uniform illumination than the birdhouse design. The laser slabs are face cooled by temperature-controlled laminar water flow in 0.5- to 1.75-mm parallel-walled channels. Two of the slabs have beveled ends; the third has square ends and is used for straight-through beam propagation. All components in the laser cavity are water cooled, and are monitored by the thermal data acquisition system. In addition, all cooling circuits that affect thermal conditions of the slab are temperature controlled. The test unit assembly is mounted on a dedicated air-leg optical table to prevent flow-induced vibration from disturbing the adjacent interferometer table.

A urethane developed at LLNL^{36,37} that has a high index of refraction was used to mount the slab in the support structure. This elastomer couples fluorescence, parasitics, and ASE to a KG-3 absorber. The energy deposited in the absorber is removed with chill bars. The adjustment of chill-bar temperature, used to control thermal boundary conditions of the slab, was found to have a significant effect on beam quality.

Casting elastomers is an important step in test unit assembly. We developed casting technology that provides essentially distortion-free mounting, leak-tight seals, and smooth flow surfaces. A number of complex mold geometries were used. The tolerance on dimensional uniformity of the cast parts is typically 0.001 in.

The test unit is modular and was configured in different ways. Several seal arrangements, placement of electric guard heaters near the beveled ends of the slab, and two chill-bar geometries were used in the experiments. In some configurations, a

fluorinated polyurethane (FPU) coating with an index of refraction of 1.3 was applied to the TIR surfaces to eliminate the interaction of the coolant with the evanescent wave. Since the FPU became cloudy after prolonged exposure to water, a Teflon overcoat was used to protect the FPU from the coolant.

A straight, heated section of the slab coolant channel was installed to induce thermal boundary layer development in the slab coolant. This section was composed of an aluminum vane with several water channels oriented transversely with respect to the slab-coolant flow. Temperature-controlled turbulent water flow was used to deposit heat in the slab coolant water.

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Gas-Cooled Slab Test Unit. The GCS test bed is shown in Fig. 9-41. A single slab of LSG-91H is mounted vertically in a Brewster's angle configuration, similar to that used for Nova disk amplifiers, and illuminated by four water-cooled flashlamps. Simple radius reflectors are used.

The GCS test unit is modular; various flashlamps, reflector designs, channel configurations, and mounting techniques may be used.

The slab is cooled by the turbulent flow of nitrogen gas in 5-mm parallel-walled channels. The test unit was designed according to ASME pressure vessel codes to operate at up to 2 atm. Coolant velocities as high as Mach 1.0 were used. The channels are defined by the slab surface and windows attached to the slab support frame. Pressure in the test bed is allowed to equilibrate with the static pressure in the nitrogen channel, minimizing distortions in the channel windows. The windows at the ends of the test unit support the pressure difference between the static pressure in the channel and ambient pressure. These windows are thicker and have a lower aspect ratio than the channel windows. Sapphire windows on the flow

Fig. 9-41. Gas-cooled slab test unit.

channels and at the ends of the housing permit thermal imaging of the slab surface with an infrared camera. Fused silica windows with superior optical figure are available for beam quality measurements.

The flashlamps, reflectors, and slab support frame are water cooled. Thermal conditions in the GCS unit are monitored by a data acquisition system similar to the one servicing the zig-zag test unit described in the preceding article. In addition, temperature and flow of the nitrogen gas are monitored.

The mounting arrangement is similar to the one used for the zig-zag test unit except that the slab in the GCS is potted on all four edges. A Kovar slab holder is used. This provides a good match in thermal index of expansion between the ASE

absorber and the slab holder. The slab holder is water cooled and functions as a chill bar. The holder cooling circuit is composed of a network of gun-drilled channels. Holders with one or two circuits have been fabricated and installed.

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Coolant Supply System. The experimental area of the HAP facility consists of two 1200-ft² laboratories flanking a common control room. One laboratory is nominally dedicated to zig-zag test bed experiments and the other to the GCS

experiments. The experiments require a capability to individually control the steady-state operating temperature of various test-bed components. In most cases, thermal energy is removed from the experiment, and calorimetric measurements on the coolant provide valuable data on the energy-balance relations of the operating test-bed components, such as flashlamps, reflectors, slab mounts (chill bars), and the laser slab.

The coolant supply system built into the laboratories is designed for flexibility and adequate capacity for experiments of increasing size. Open-loop city water circuits that ultimately discharge into a storm drainage system are the basic component cooling system. City water is a relatively cool, stable coolant that is economical for the limited duration of typical experimental runs. The LLNL closed-loop, low-conductivity water (LCW) system is used to cool components exposed to high voltage (primarily flashlamps).

The laser slabs, which currently are Nd:glass, but are projected to be replaced by a crystalline material, Nd:Cr:GSGG, have special cooling requirements. The zig-zag slab is optically pumped by flashlamps. The energy radiates through the liquid slab coolant to the slab at peak optical intensities as high as 6 kW/cm^2 in current experiments; the intensity will be higher in planned experiments. The slab must be cooled by very clean, pure water to avoid the surface damage caused by light absorption in contaminants. For the case of the gas-cooled slab, the intense energy of the laser beam propagating through the cooling medium dictates the use of a very clean, gaseous coolant.

Each of the two HAP laboratories has a permanent system of water-cooling circuits. A generous number of circuits provides individual temperature control of numerous components and allows for multiple experimental setups. Relatively high flow capacities are available if near-isothermal conditions are required.

There are ten city-water circuits in each lab, each with a maximum capacity of 30 gpm. A total of 200 gpm of city water is available, filtered to $50 \mu\text{m}$, pressure regulated for flow stability, branched, valved,

and measured at wall-mounted trees, and finally delivered to the experimental area through overhead piping. After passing through the experiment, city-water flows are directed to an elevated collection manifold that subsequently discharges to an exterior storm drain. The collection manifold is vented to atmosphere, which ensures that back pressures (and hence flow rates) in individual circuits are not influenced by flow changes in other circuits.

There are six LCW circuits in each lab, each with a maximum capacity of 15 gpm. LCW is filtered to $10 \mu\text{m}$, regulated, controlled, and delivered in a manner similar to that described for city water. After passing through the experiment, LCW is returned to the LLNL closed-loop system.

Portable, self-contained chiller units are used when it is necessary to provide temperature regulation of components at a specified level. One chiller with a closed-loop flow capacity of 18 gpm and three chillers with 8-gpm capacities are available.

Special high-cleanliness, temperature-controlled coolant systems have been devised for the critical slab coolants. The zig-zag liquid slab coolant is temperature controlled and circulated by a special chiller that has only stainless steel and selected plastic components in contact with the coolant water. High-purity distilled and deionized water is used; both feedwater and circulating filter systems can remove particles larger than $0.1 \mu\text{m}$ and remove ions to an electrical resistance of $18 \text{ M}\Omega$ between probes spaced 1 cm apart.

The slab coolant system for the gas-cooled slab test bed is an open-loop nitrogen circuit that draws up to 5000 scfm (2 kg/s) of gas at ambient temperatures and pressures up to 80 psi from the 28 000-gallon Nova nitrogen supply tank. Several hundred feet of 10-in.-diam stainless-steel pipe convey the gas to an exterior area just outside the east wall of the GCS lab. In this area, a system of valves, pressure regulators, filters, heaters, cooling coils, and pressure-relief devices controls the flow rate, pressure, and temperature of the nitrogen gas delivered to the experiment. The 10-in. piping enters the GCS lab above ceiling level and extends to a point

directly above the GCS test bed. Nitrogen flow is straight downward through the experiment via ducts that gradually change in cross section to match that of the test-bed channels. The exhaust piping continues downward through the laboratory floor, through a turbine flowmeter, and finally to an exhaust stack mounted on the outside of the building.

The long vertical nitrogen channel is designed for a uniform, stable flow of gas through the experiment. With experience, we may be able to reduce the bulk and complexity of this design. The open-loop operation of the nitrogen system does not represent a significant expense to the HAP program, given the short duration of experimental runs. Provision has been made for later conversion to closed-loop gas cooling, primarily to allow the use of helium, which has cooling properties four times better than those of nitrogen.

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Thermal and Flow Diagnostics. A HAP laser requires the continuous cooling of internal components. A major design goal is to devise a system that does not induce thermal distortions in the internal optical components. This demands a thorough understanding of heat transfer within the experiment. Accurate measurement of the distribution of the thermal energy among the operating test-bed components is a critical input for the computer models used to guide the design of a successful HAP laser.

Point temperatures in the HAP experiments are measured with thermistors, precision semiconductor devices whose resistance varies inversely with temperature. A bare thermistor is small enough to fit within a 2-mm-diam spherical volume, which allows implantation in test-bed components and fittings with a minimum of physical disturbance. A precisely regulated 0.1-mA current is continuously applied to each thermistor. The resulting voltage drops across each thermistor are measured sequentially by a digital voltmeter and relayed to an IBM AT computer over a fiber-optic transmission line. In the HAP control room, the thermistor

temperatures are calculated using individual calibration data for each thermistor. The system is currently configured to scan up to 40 thermistors.

Because energy balance measurements on coolant flows typically require an accurate reading of relatively small temperature differences ($T_0 - T_i$), precise and repeatable temperature readings are needed. We developed equipment and techniques (including the partially automated calibration of individual thermistors) in the HAP lab that measure thermistor temperatures with a reliability of $\pm 0.04^\circ\text{C}$, relative to a highly accurate quartz thermometer. The calibration data are taken in a small isothermal bath with a measured temperature uniformity better than 0.01°C .

Coolant flow rates are measured with axial flow (turbine) and tangential flow (paddle wheel) flowmeters. Within each meter, the rotation of the freely spinning bladed rotor is sensed by an rf pickup. The resulting electrical pulses are amplified and sent to a frequency counter, whose output is sent in turn to the IBM AT computer over the fiber-optic transmission link. At this time, we can scan as many as 12 turbine and 4 paddle-wheel flowmeters. Factory calibration data, verified by "stopwatch and bucket" measurements, are incorporated in the computer software. Our routine flow rate measurements are accurate to $\pm 2\%$.

Individual thermistor measurements take ~ 0.75 s each, and flowmeter measurements, ~ 1 s each. For the current zig-zag test bed, a complete system measurement cycle is done in ~ 45 s. A computer program automatically displays all inlet and outlet temperatures, flow rates, circuit thermal power, and overall electrical input and thermal output powers after each cycle. The soft keys on the computer keyboard can be used to switch to an alternative display of up to 15 temperatures, flow rates, or powers as a function of time. The soft keys are also used to make any desired changes in experiment duration and to enter the flashlamp electrical power, which is obtained by integrating the voltage-current product on a second IBM computer. The raw data is stored on a floppy disk, and any displays can be regenerated using the data file after the experiment is terminated. A related family of programs allows us to test individual

thermistor flowmeter circuits and to acquire calibration data.

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Analysis

Introduction. The primary objective of thermal analysis in the HAP Program is to identify and quantify factors influencing thermal conditions of the slab. Slab cooling schemes to date have employed fluid flow, so much of this work was in the area of forced convective heat transfer. A secondary objective is to examine removal of waste heat from other components in the laser system.

Thermal-fluids studies conducted in 1985 were based on closed-form and numerical solutions for both developing and fully developed flow with heat transfer. Existing empirical correlations were extensively employed, especially for turbulent flow where analytical solutions are of limited usefulness. We applied these results to the design of the zig-zag and GCS test units and to the conceptual design of proposed HAP laser systems. The following thermal design criteria were the basis for analysis of slab cooling:

- Maintain an acceptably small linear temperature gradient along the slab faces, as specified by stress-optics analysis.
- Minimize the overall temperature level in the slab in order to sustain an acceptably low absorption coefficient in GSGG and to minimize thermal population of the lower laser level.
- Minimize nonlinear thermal behavior arising from (1) three-dimensional coolant flows (especially those due to nonuniform introduction of coolant), secondary flows in corners, and local disturbances around gaps and seams; (2) thermal and hydrodynamic boundary layer development; and (3) thermal energy conducted between the slab and its supporting structure, especially the ASE absorber.

In 1985, we calculated thermal operating parameters optimized for various flow geometries; designed coolant-channel entrances and exits; and studied the relation of nonlinear effects to fabrication tolerance, tapered flow channels, and thermal boundary layer development. Conceptual system designs for thermal and fluid flows were developed for a 10-slab GCS amplifier, the FEL driver, and a space-based laser radar system (LIDAR).

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Ten-Slab GCS Parameter Study. During 1985, the thermal operating parameters of a 10-slab GCS amplifier³⁸ that uses the existing HAP lab nitrogen system were optimized for various flow geometries. Computations were made with the IDFLOW code, which is under development at LLNL for thermal and flow analysis in both the developing and fully developed flow regions. The improvements we made to the code include: graphics capability, automation of nozzle and diffuser design, extended heat-transfer correlations, and variable preheating length to account for ASE heating of the coolant.

We investigated three flow geometries (Fig. 9-42):

- Parallel flow through 20 flow circuits (top to bottom through the channels on both sides of the 10 slabs).
- Series flow in two circuits, one on each side of the amplifier (horizontal flow across a slab, turning in a common slab holder, horizontal flow across the next slab, etc.)
- Alternating flow down one channel, up another, down a third, etc. In Fig. 9-42, two circuits use space only above and below the amplifier for nitrogen plumbing. The other two circuits also extend laterally, although not at the elevations occupied by flashlamps. Plumbing for alternating geometries with other than four circuits is even more complicated.

Inlet static pressures of 150 and 200 kPa (1.5 and 2.0 atm) and channel heights between 0.5 and 30 mm were considered. The 10 slabs were each 100 mm high, 230 mm wide, and 15 mm thick. Each flow

geometry and pressure combination was investigated for a variety of flow rates and channel heights. Two figures of merit were used to evaluate the various configurations: optimal use of the existing nitrogen system and slab optical performance. The nitrogen system can supply about 2 kg/s at 500 kPa (5 atm) inlet pressure. Initially, however, we set the flow limit at 1 kg/s total for the 10 slabs. This is the maximum flow available at 200 kPa, and our results showed that it was not necessary to resort to higher pressure.

The key optical performance goal was to minimize the fractional slab area that has a nonlinear temperature gradient in the flow direction, since a nonlinear temperature gradient is believed to produce

stress-birefringence problems. A low slab temperature was also desired to reduce thermal population of the lower laser level.

Two other key variables were calculated: the temperature gradient in the linear region, which can produce an optical wedge in the beam, and the laminar sublayer thickness. Obstacles larger than this thickness can disrupt the flow.

Seven of the 1000-plus calculated operating points for the 10-slab GCS amplifier are summarized in Table 9-6. The design point conditions for the current single-slab GCS experiment is also shown, after scaling to the 10-kW/m² heat flux used for this study. The peak slab-surface temperature is based on a 20°C inlet-gas temperature. Each GCS-10 case shown uses

Fig. 9-42. Parallel, series, and alternating flow geometries at the slab centerline elevation (top view).

Table 9-6. Selected design points calculated for parallel, series, and alternating flow geometries for the 10-slab GCS amplifier and scaled values from the single-slab GCS-1 experiment. The GCS-1 column is based on parallel flow over 10 GCS-1 slabs.

Design parameters	Parallel		Series			Alternating		GCS-1
Pressure (MPa)	0.15	0.2	0.15	0.2	0.2	0.15	0.2	0.2
Channel height (mm)	1	0.5	10	10	5	1	0.5	5
Flow circuits (10 slabs)	20	20	2	2	2	10	10	20
Total flow (kg/s)	0.80	0.52	0.36	0.56	0.24	0.40	0.26	3.8
Peak surface temp. (°C)	42	39	62	48	58	48	48	35
Nonlinear area (%)	0	0	7.8	9.8	3.3	0	0	23
Maximum obstacle (μm)	11	8	10	7	9	11	9	6
Optical ΔT/slab (°C)	5	8	2.5	1.7	2.4	9	8	6
Peak Mach number	0.36	0.44	0.5	0.5	0.5	0.41	0.58	0.5

considerably less gas flow than the current single-slab experiment. The peak glass surface temperatures are reasonable, and, in many cases, a small preheat length (perhaps supported by ASE) is adequate to ensure a linear slab profile in the flow direction.

We consider the parallel flow geometry to be the best because it has the least complex plumbing and the lowest slab temperatures; the flow rates are reasonable even though they are higher than in the other geometries. We rejected the alternating geometry because it would be difficult to construct the gas flow connections between the slabs, which are alternated at Brewster's angle. We rejected the series geometry because it has an extended nonlinear temperature region due to the wide channel needed to limit pressure losses over the long flow distance and because it is difficult to construct the slab holders for continuous flow between slabs.

We made several geometric simplifications in this study. We ignored plenum and header losses, assumed a diffuser recovery factor of 1.0, and ignored abrupt area reductions due to flow straighteners. Although these effects are not expected to change our conclusions, they should be considered during the design of a 10-slab amplifier.

Similarly, we ignored end-wall effects and tolerance variations in the channel, and we used published (approximate) correlations in the 1DFLOW code. The recently acquired TEXTAN code could be used to produce improved correlations, and the three-dimensional FIDAP code could be used to model end-wall effects and tolerance variations.

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Nozzle and Diffuser Design. Flow in the GCS slab-coolant channel can be affected by perturbations at the entrance and exit of the channel. We designed the entrance nozzle and exit diffuser for the GCS to provide well behaved, predictable flow conditions in the slab region. The configuration of the slab-coolant channel shown in Fig. 9-43 includes the nozzle and diffuser.

Fig. 9-43. GCS slab-coolant flow channel with nozzle and diffuser.

Nozzle Design. Heat in the slab region is transferred through a developing thermal boundary layer, a fully developed boundary layer, or a transitional boundary layer between the laminar and turbulent regimes. Developing and transitional flows are likely to occur in the slab region of the test unit. Because the heat transfer in these regimes varies strongly in the direction of flow, we must identify the origin of the boundary layer to determine the heat-transfer distribution.

A nozzle design was developed using two-dimensional wind tunnel data.³⁹ This approach provides good exit velocity uniformity (i.e., little boundary layer development in the nozzle). The design employed in the GCS test unit has an overall contraction ratio of 6:1. The nozzle is fed by a straight rectangular duct in which indeterminate external flow occurs. We expect flow in this region to be reasonably uniform and free from substantial swirl. Because nozzle performance is only weakly affected by conditions in the upstream duct, it was unnecessary to consider the flow conditions in the upstream section. A boundary layer trip, created at the junction between the exit of the nozzle and the start of the area of constant flow, also helps establish the origin of the boundary layer.

Diffuser Design. In subsonic flow, which is employed in the GCS test unit, maldistributions in the flow downstream of the slab region can propagate upstream and adversely affect heat transfer. To avoid this, the diffuser section must minimize flow-induced instabilities. Good pressure recovery is also desirable.

Data for straight-walled, two-dimensional diffusers were used for the design.⁴⁰ Operation of the diffuser in a stall-free state requires an area expansion ratio of less than 1:4. The exhaust duct of the nitrogen supply system is sufficiently large that this cannot be accomplished by a one-stage diffuser. As a consequence, we elected to use two parallel four-stage diffusers. Each stage is designed to operate without separation. Short sections of honeycomb flow straightener between stages are used to control the inlet conditions at each stage. This design eliminated the flow maldistribution

that was apparent in the gas supply system before installation of the diffuser.

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Effect of Fabrication Tolerance on Flow in Zig-Zag Channels. Variations in channel size on the order of the fabrication tolerance can cause local maldistributions in heat transfer. These arise from the effect of channel size on both the local flow rate and the temperature difference across the boundary layer. We conducted the analysis described below for the cases of fully developed laminar and turbulent incompressible flow in high-aspect-ratio ducts.⁴¹

The analysis compared the flow and heat transfer characteristics of parallel stream tubes of different but constant height (t_1, t_2) subject to a common pressure drop. The resulting relations between conditions in the channels to the ratio of channel heights are given by the following set of equations:

$$\dot{m}_1/\dot{m}_2 = (t_1/t_2)^a, \quad (32)$$

$$\Delta T_{m,1}/\Delta T_{m,2} = (t_1/t_2)^b, \quad (33)$$

$$\Delta T_{b,1}/\Delta T_{b,2} = (t_1/t_2)^c, \quad (34)$$

where $\dot{m}$ is the mass flow rate, ΔT_{b1} is the difference between the slab surface temperature and the local mean coolant temperature, and ΔT_m is the change of the mean coolant temperature. Since channel height is constant, ΔT_m is equivalent to the change of slab surface temperature in the flow direction. Subscripts 1 and 2 denote the respective channels. For laminar flow, $a = 3$, $b = -3$, and $c = 1$. For turbulent flow, $a = 1.5$, $b = -1.5$, and $c = -0.2$.

First, it is surprising to note that mass flow and temperature ratios depend only on channel height and flow regime. A dependence on fluid properties is conspicuously absent. As a consequence, the results are valid for any coolant, liquid or gas. Second, variations in channel size affect heat transfer in laminar flow much more strongly than they do in turbulent flow.

The behavior described here is difficult to isolate in actual practice since a means of directly observing thermal maldistributions is necessary. Simple tests of uniformity such as flow visualization are helpful, but insufficient to assure uniform heat transfer, because thermal maldistributions can occur in the presence of straight streamlines (e.g., when channel height varies transversely with respect to the flow).

We evaluated Eqs. (32) through (34) at the nominal design operating point of the zig-zag test unit, and found that the channel height variations we expected initially (0.15 mm) could cause nonuniformity in the surface temperature of the zig-zag slab of up to several degrees Celsius. As a consequence of this analysis, we decided that channel height uniformity should be controlled to within 0.025 mm. This is presently the practical limit of accuracy that can be achieved with our elastomer casting technology.

Author: D. A. Bender

Effect of Tapered Channel on Developing Turbulent Flow. Developing thermal flow in the constant area section of the GCS coolant channel causes the slab surface temperature to vary nonlinearly in the direction of flow. We conducted an analysis to determine how well a tapered (i.e., varying area) flow channel modifies turbulent boundary layer development. Our objective was to achieve uniform temperature distribution on the surface of the slab.⁴²

The analysis was based on external incompressible flow of a coolant with constant properties. A uniform heat flux q'' was defined beginning at $x = 0$, where flow is in the x direction. The length of the channel is L ; the area of the channel A decreases linearly with x and has the value A_1 at $x = 0$ and A_2 at $x = L$. The expression derived for slab surface temperature T as a function of x is:

$$T = C_{op}(1 - A^*)^{4/5} X^{1/5}, \quad (35)$$

where

$$C_{op} = 33.30 \text{Pr}^{2/5} c^{-1} \mu^{-1/5} \dot{m}^{-4/5} A_1^{4/5} L^{1/5} q'', \quad (36)$$

$$A^* = 1 - A_1/A_2, \quad (37)$$

$$X = x/L. \quad (38)$$

Pr is the Prandtl number, c is the heat capacity of the coolant, μ is the viscosity of the coolant, $\dot{m}$ is the mass flow rate, A^* and X are nondimensional, and C_{op} is temperature rise.

A nondimensional plot of slab surface temperature vs normalized channel length is shown in Fig. 9-44 for several contraction ratios. A value of $A^* = 0.65$ gives a slab surface temperature that is uniform to within $0.04 C_{op}$ over the last 75% of the heated region. The correlations used in this analysis⁴³ are valid for the contraction ratios indicated in Fig. 9-44. The magnitude of the temperature rise C_{op} is dependent on the operating point but the shape of the profile is not. Nonlinearity remains appreciable in the entrance region of the channel.

Optimization of Channel Size. In conjunction with the analysis of developing turbulent boundary layers in a duct with varying area, channel size was optimized for fully developed, compressible flow in a duct of constant area. The optimization minimized the combination of boundary layer temperature rise and the rise in the stream temperature of the coolant due to heating.⁴⁴ In compressible flow, mass flow per unit area is limited by sonic conditions. Thus, as channel size decreases, flow decreases and the rise in mean fluid temperature in the direction of flow increases. As channel size increases, the temperature rise across the boundary layer increases. The channel height t at which the sum of the temperature rises is at a minimum is given by

$$t = 3.406(bkL/c)^{5/6} \mu^{-2/3} (\dot{m}'')^{-1/6}, \quad (39)$$

where $b = 0.023 \text{Pr}^{0.4}$, and in any consistent set of units, k is the thermal conductivity of the coolant, L is the length of the constant area duct, c is the heat capacity of the coolant, μ is the viscosity of the coolant, and $\dot{m}''$ is the mass flux (mass flow rate per unit flow area).

The optimum channel size is determined almost entirely by geometry and fluid

Fig. 9-44. Slab surface temperature vs normalized channel length for flow in a converging channel.

properties. The only term that is operating-point dependent in Eq. (39) is m'' , which has little significance because it is raised to a power of $-1/6$. The optimum channel size calculated for the present GCS test bed is ~ 1.5 mm.

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Laser Materials

Introduction

An ongoing goal in the LLNL laser program is the development of solid state lasers with greater power, enhanced efficiency, and improved beam quality. In pursuit of this goal we are conducting a broadly based program to identify, fabricate, and characterize materials suitable for improved lasers, including advanced crystals and glasses.

For our initial technology demonstration lasers we are using GSGG:Nd,Cr crystals as well as silicate and phosphate laser glasses. As an integral part of this effort we are supporting investigations into advanced growth techniques, which are needed for obtaining large-sized plates of the materials required for high-average-power lasers. In addition to the work

described here, we are investigating the possibility of using fluoride and oxide crystals as advanced laser materials for a variety of applications. The results of this effort are described in Sec. 7.

Program accomplishments in 1985 include the following:

- Establishing facilities in LLNL, universities, and industry for growth and characterization of new materials.
- Developing improved phosphate glass compositions with long neodymium fluorescence lifetimes and high thermal shock figures of merit.
- Measuring GSGG:Cr,Nd properties and laser performance.
- Initiating a program with the Allied Corporation to produce $10 \times 20 \times 1$ -cm GSGG:Cr,Nd plates.
- Establishing a LLNL program to model and characterize experimentally the Czochralski method of growth, which is used to grow all commercial laser crystals.

Author: S. E. Stokowski

GSGG:Nd,Cr Crystals

Introduction. In 1985 we decided to speed the development of GSGG:Nd,Cr as a high efficiency laser material for the zig-zag and gas-cooled slab lasers. We

accordingly began work in three areas necessary for bringing co-doped GSGG to fruition: increasing GSGG boule size, determining a polishing procedure, and developing an edge-cladding. We initiated the growth of 5-in.-diam GSGG:Nd,Cr boules; determined a polishing procedure that gives smooth, high-strength surfaces; and found a glass edge-cladding that matches GSGG.

In addition we measured the spectroscopic, optical, and thermo-mechanical properties of GSGG⁴⁵ (see also Sec. 7) and characterized the quality of the GSGG material that we had received as a result of previous growth programs.⁴⁶ We found that the quality of our GSGG material is comparable to that of YAG, but with almost twice the lasing efficiency.

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Material Characterization. We set up measuring instrumentation and investigated the transmission, wave front distortion, and birefringence of material from a number of GSGG boules.

The material was obtained under two DOE contracts in GSGG:Cr,Nd crystal growth, which were completed early in 1985. Allied Corporation, Synthetic Crystal Products Division, and Airtron, a division of Litton Industries, each grew ten single-crystal boules of the material and delivered a number of laser rods cut from these boules for our evaluation. Final reports^{47,48} on work completed under these contracts were distributed widely within the DOE, DOD, and industrial communities as part of our technology transfer process. GSGG rods were also procured from Material Progress Corporation and from Union Carbide Corporation under separate LLNL purchase orders. In addition, a number of high quality YAG:Nd laser rods were purchased and tested for comparison to the GSGG:Cr,Nd laser rods. Our evaluation of the laser materials obtained through these contracts was reported at the 1985 Conference on Lasers and Electro-Optics⁴⁹ and in a detailed laboratory report.⁵⁰ The results are summarized below.

There was a notable 1- μ m absorption loss in the GSGG:Cr,Nd material. A number of the GSGG:Cr,Nd crystals obtained from industrial sources were found to exhibit an absorption band in the near-IR stretching

from about 0.85 to 1.20 μ m. The absorption varied significantly from sample to sample, and was often strong enough to affect lasing efficiency or even totally prevent laser action at 1.06 μ m. The absorption spectrum of a crystal exhibiting this effect is shown in Fig. 9-45.

After a long series of analytical studies we were able to establish a correlation between the strength of this absorption and the level of calcium impurity in the crystal⁵¹ as shown in Fig. 9-46. We now believe that the divalent calcium causes chromium to switch from the +3 to +4 valence state if it is not otherwise compensated by +4 impurities.⁵² Cr^{4+} on a tetrahedral site in the crystal is thought to produce the near-IR absorption band.

Fig. 9-45. Absorption spectrum of a GSGG:Cr,Nd crystal grown from a melt containing 100 wppm CaO.

Fig. 9-46. Absorption coefficient at 1.12 μ m vs calcium content of GSGG:Cr,Nd crystals. Two of the crystals were grown from the same melt before and after adding 50 wppm CaO.

Fortunately, most of the crystals we received had acceptably low absorption coefficients at the 1.06- μm laser wavelength, indicating that an appropriate impurity charge balance can be readily achieved. However, crystal growers will occasionally add 50 to 100 wppm of calcium oxide to the melt to prevent a spiral growth morphology, which is otherwise prevalent for this type of garnet crystal, although the mechanism by which the added calcium prevents spiral growth is not understood. Two of the data points shown in Fig. 9-46 were obtained from crystals grown from the same melt before and after 50 wppm of calcium oxide was added.⁵³

Chemical analysis of the starting materials used in the production of GSGG revealed the presence of calcium as a major impurity, typically at levels exceeding 100 ppm, in most batches of scandium oxide tested. The other starting materials had very little calcium by comparison. To obtain high purity Sc_2O_3 we contracted with the Ames Laboratory Materials Preparation Center to refine 98 to 99% pure Sc_2O_3 to 99.995% or better in kilogram quantities. They have delivered to us about 25 kg of high purity Sc_2O_3 .

We also undertook measurements of losses and lasing efficiency in GSGG:Cr,Nd laser rods. As part of our effort to characterize laser crystals grown by industry, the free-running laser efficiency was measured for each laser rod delivered to us. The early measurements yielded seemingly inexplicable variations in lasing efficiency. Some of this variation was due to differences in the strength of the 1- μm absorption band discussed above. However, significant rod-to-rod variations in lasing efficiency were observed even in materials with no apparent absorption loss and essentially equal concentrations of the dopants. Further investigation revealed that these variations were caused primarily by differences in crystal quality and the resulting losses in the laser resonator.

A series of measurements was made to investigate the absorption losses in the GSGG:Cr,Nd laser material. Our testing program revealed several sources of loss. The major losses were due to:

1. Absorption by impurities and defects, including Cr^{4+} .
2. Wavefront distortion due to refractive-index inhomogeneity.

Table 9-7. Summary of losses and free-running laser efficiency measurements on GSGG:Cr,Nd, YAG:Nd, and GGG:Cr,Nd laser rods.

Crystal number	Host crystal	Dopant conc. [ND]/[Cr] (10^{20} cm^{-3})	Spectro- photometer attenuation coefficient ($\% \text{ cm}^{-1}$)	Distortion waves at (rms) 1.06 μm	Depolarization loss ($\% \text{ cm}^{-1}$)	Diffraction limited transmission loss ($\theta/2 = 250 \mu\text{rad}$) (L)	Findlay-Clay round-trip insertion loss η_{abs} (%)	Free running efficiency measurements		
								Absolute η_s (%)	Slope η_0 (%)	Intrinsic
1038	GSGG	3.0/0.29	0.5	0.043	0.038	1.8	0.13	3.8	5.5	6.3
1035	GSGG	2.2/2.0	4.3	0.027	0.018	3.7	0.81	2.8	4.2	8.0
3002	GSGG	1.3/0.9	0.8	0.082	0.128	3.1	0.14	3.8	5.3	6.1
3003	GSGG	2.0/0.9	0.8	0.170	0.170	4.5	0.20	4.6	6.6	8.1
3004	GSGG	2.0/0.9	0.8	0.162	0.177	8.5	0.21	4.4	6.2	7.6
3006	GSGG	2.5/0.9	0.4	0.047	0.034	0.9	0.12	4.6	6.5	7.4
3007	GSGG	2.0/0.9	0.9	0.080	0.076	3.6	0.51	3.1	4.8	7.5
3008	GSGG	1.8/2.0	0.5	0.144	0.372	3.9	0.12	4.9	7.0	7.9
3010	GSGG	2.0/2.0	0.5	0.099	0.059	2.4	0.18	4.7	6.7	8.0
3011	GSGG	2.0/2.0	0.8	0.049	0.099	1.5	0.19	4.6	6.5	7.9
3012	GSGG	2.0/2.0	0.5	0.066	0.199	5.1	0.14	4.8	7.1	8.2
4011	GSGG	1.3/1.0	0.9	0.072	0.452	4.3	0.47	3.3	5.0	7.7
5004-3	GSGG	2.2/2.0	0.8	0.031	0.033	1.0	0.36	4.4	6.2	8.7
5004-4	GSGG	2.2/2.0	1.2	0.019	0.007	0.9	0.33	4.2	6.0	8.2
5005	GSGG	1.9/2.0	0.4	0.023	0.026	0.2	0.13	4.8	6.8	7.7
5007-2	GSGG	1.9/2.0	1.0	0.041	0.012	1.3	0.23	3.7	5.4	6.7
5008	GSGG	1.9/1.0	0.4	0.029	0.007	0.3	0.11	4.6	6.6	7.4
5009	GSGG	1.9/2.0	0.5	0.037	0.001	0.5	0.16	4.2	6.1	7.2
5010	GSGG	1.9/2.0	0.2	0.031	0.391	0.5	0.12	5.0	7.0	7.9
5014	GGG	$\sim 2.0/1.0$	2.1	0.150	0.816	11.6	0.64	2.7	4.3	7.3
0000	YAG	$\sim 1.5/0$	—	0.039	0.005	0.5	0.08	2.2	3.1	3.4
3013	YAG	$\sim 1.5/0$	0.3	0.029	0.005	1.1	0.05	2.8	3.6	3.8
4100	YAG	$\sim 1.5/0$	—	0.029	0.003	—	0.08	2.2	3.0	3.2

3. Scattering from defects such as inclusions, voids, and striations.

4. Strain birefringence due to core, dislocations, striations, and other causes.

5. Intrinsic absorption from the lower laser level.

Several measurement techniques were used to characterize the laser rods under study. These included:

1. Twyman-Green interferometry at $1.06\ \mu\text{m}$.
2. Circular polarimetry.
3. Spectrophotometry.
4. Diffraction-limited transmission (or pinhole transmission) measurements.
5. Laser resonator insertion (the Findlay-Clay technique).⁵⁴

Detailed descriptions of the measurement techniques and sources of loss have been reported elsewhere.⁵⁰ Major results of the loss measurements are given in Table 9-7. The measured transmission losses, wave front distortion, and birefringence of GSGG and YAG boules are summarized in Fig. 9-47.

The laser resonator used for the measurement of resonator insertion loss was also used as a free-running oscillator to measure the lasing efficiency of each laser rod. The best results were usually obtained with an output coupler reflectivity of about 40%. Values are given in Table 9-7 for all rods tested.

It can be shown within a reasonable approximation that the measured slope efficiency η_s is related to the resonator insertion loss parameter L and the output coupler reflectivity R_{out} by the expression:

$$\eta_s = \eta_0 \frac{\ln R_{\text{out}}}{\ln R_{\text{out}} - L}, \quad (40)$$

where η_0 is the intrinsic slope efficiency, which could be obtained from each rod if all sources of loss were eliminated. In this way we can determine the intrinsic slope efficiency for the laser material in each rod by using the measured slope efficiencies reported here and the Findlay-Clay resonator insertion loss parameters. This measurement is independent of the optical quality of the rods. The calculated intrinsic slope efficiencies are also given in Table 9-7.

We experimentally verified the efficiency expression [Eq. (40)] for a large number of GSGG:Cr,Nd, and YAG:Nd rods having a wide range of Findlay-Clay round-trip loss. Equation (40) gives good agreement with measured slope efficiencies, using $\eta_0 = 0.0775$ and $R_{\text{out}} = 0.408$. For GSGG rods with chromium and neodymium concentrations of about $2 \times 10^{20}/\text{cm}^3$ the average intrinsic slope efficiency was 7.75%. For the YAG rods measured, the average intrinsic slope efficiency was 3.63%. This agreement provides substantial confirmation

Fig. 9-47. Transmission loss, wave front distortion, and birefringence of GSGG:Nd,Cr and YAG:Nd rods. All specimens were 3 in. long and 0.25 in. across.

of the model. The intrinsic lasing efficiency of GSGG:Cr,Nd rods with identical neodymium concentrations was not significantly different for chromium ion densities of 1×10^{20} and $2 \times 10^{20}/\text{cm}^3$.

The measurements we have performed indicate that the losses characteristic of the GSGG laser material now available are generally higher than corresponding values for commercially grown YAG. Perhaps this is not too surprising, considering the fact that the development of GSGG crystal growth technology is in a relatively early stage. In spite of the higher losses, however, all of the GSGG laser rods tested have yielded higher lasing efficiency than the best YAG rod. When the losses are accounted for, the intrinsic lasing efficiency of some of the co-doped GSGG rods is found to be more than twice that of the YAG:Nd rods.

Author: J. A. Caird

Crystal Growth. In August 1985 we contracted with the Allied Corporation to grow 5-in.-diam boules of GSGG:Nd,Cr. We will cut $10 \times 20 \times 1.5$ -cm slabs from these boules for the medium average power laser (MAP) program. The Allied Corporation will grow eight boules, which should supply us with at least four good plates for

our laser tests. If the growth program is highly successful we may obtain a significantly greater number of plates.

In October 1985 we delivered a 9-in.-diam by 9-in.-deep iridium crucible to Allied, and in November they tested the large-growth system by successfully growing a 5-in. undoped GGG boule. (See Fig. 9-48.) Allied has attempted to grow the codoped GSGG material in this system, but there have been problems with the growth morphology. Spiraling of the boule appears when crystal growth exceeds 1 to 2 in. Allied is attempting to control this unwanted spiraling by increasing thermal gradients in the system in combination with chemical additions such as calcium. Both techniques are known to suppress spiraling. Adding calcium to GSGG:Nd,Cr has the unwanted side effect of inducing the formation of Cr^{4+} , which is a $1\text{-}\mu\text{m}$ absorber as discussed in the preceding article. Therefore, we believe that the addition of another ion may be necessary to compensate for Ca^{2+} , perhaps cerium or titanium, both of which can be found in trivalent or tetravalent states. We believe that the $\text{Ce}^{3+} \rightarrow \text{Ce}^{4+}$ and $\text{Ti}^{3+} \rightarrow \text{Ti}^{4+}$ reactions are energetically more favorable than $\text{Cr}^{3+} \rightarrow \text{Cr}^{4+}$.

Author: S. E. Stokowski

Fig. 9-48. A 5-in. undoped GGG boule.

Optical Finishing of Large Garnet

Slabs. Our development of an optical fabrication method for garnet laser components began with a systematic evaluation of existing methodologies. As a basis for a preliminary assessment we employed polishing methods previously used for lasers⁵⁵ and bubble memory substrates.^{56,57} Invariably, these failed to meet one or more of our objectives for flatness ($< \lambda/20$ over 400 cm^2), surface finish ($< 3 \text{ \AA}$ rms), reasonable material removal rate, and low subsurface damage. Keeping subsurface damage to a minimum is important to maintain high slab strength.

We then began a broad survey of pitch types, polishing compounds, and slurry pH. Our initial studies were performed on 3-cm-diam $\times$ 1-cm-thick GGG samples. When we identified promising techniques we scaled the process up in a series of steps, first to a moderate-sized slab measuring $0.5 \times 4 \times 8 \text{ cm}$, and then to a large slab with dimensions of $1.5 \times 10 \times 16 \text{ cm}$. Compatibility of the technique with GSGG:Nd,Cr was then established with intermediate-sized pieces measuring $1 \times 4 \times 10 \text{ cm}$.

We then investigated some less common oxides. We used standard grinding steps.⁵⁵ After cutting, we began with Blanchard surface grinding. A 100-grit, bound-diamond abrasive in water was used to remove substantial amounts of material (200 to $500 \mu\text{m}$) and generate a flat surface. After grinding came polishing, using 30-, 15-, 9-, and $3\text{-}\mu\text{m}$ alumina in water on cast iron (30 and $15 \mu\text{m}$) or Pyrex (9 and $3 \mu\text{m}$). In these polishing steps we removed 130, 70, 50, and $20 \mu\text{m}$ of the garnet material, respectively. Finally, further polishing produced a smooth surface with an acceptably low value of optical scattering.

We considered several approaches for the final polishing step. Double-sided planetary polishing and the use of fabric polishing pads were not attempted because we felt that figure control would be inadequate. We concentrated on methods that could be applied using a standard pitch lap. A method, detailed elsewhere,⁵⁸ using a combination of titania and silica compounds was found to be most effective.

To estimate the strength of slabs prepared using this technique, we assessed the

subsurface damage using a strength method.⁵⁹ Disks ground and polished in exactly the same way as the slabs were loaded to failure in biaxial tension. The size of the maximum flaw was then calculated using fracture mechanics, and a safe operating stress for the slab was estimated using Weibull statistics. For the grinding and polishing procedure outlined here, the mean strength of ten specimens was 516 MPa, giving an estimated flaw size of $5.4 \mu\text{m}$. The Weibull statistical parameters were $m = 6.59$ and $S_0 = 556 \text{ MPa}$. Scaling to a $1.5\text{-} \times 10\text{-} \times 20\text{-cm}$ slab size and assuming failures from surface defects rather than edge flaws, we obtained a design stress estimate of 186 MPa for a failure probability of one component in ten.

Author: J. E. Marlon

Major Contributor: E. R. Prochnow

Edge-Cladding. We contracted with the Hoya Corporation to develop a glass edge-cladding to suppress ASE in GSGG slabs. Hoya has identified a glass that matches the index of refraction and thermal expansion of GSGG. It is a high-lead glass of the composition (in mole %) 65 PbO, 17.5 Al_2O_3 , and 17.5 SiO_2 , with 10 wt% CuO added. Hoya has applied this glass to test plates of GSGG and found good adherence. At first they had some problem with interface bubbles, but they have been able to lower the bubble count to about $100/\text{cm}^3$ and bubble sizes to less than $20 \mu\text{m}$. This bubble density is marginal, since it causes a scattering of 10^{-3} of the incident energy, but a further decrease in bubble density is anticipated. Hoya has also developed the molding equipment necessary for cladding the large plates.

Author: S. E. Stokowski

Glasses for High Power Lasers

Neodymium-doped laser glasses are useful in HAP lasers where high pulse energy is needed and a low pulse repetition rate in the range of 1 Hz is tolerable. When compared with crystalline laser materials, laser glasses are easier to obtain in large

Table 9-8. Measured properties of new phosphate glasses.

Laser Materials

Glass	Cross section (pm ²)	r at 10 ²¹ Nd ions/cm ³	Thermal conductivity (W/m·K)	Thermal expansion (10 ⁻⁷ /K)	Young's modulus E (GPa)	Fracture toughness K _{IC} (MPa·m ^{1/2})	Thermal shock FOM (W/m) ^a
Ultraphosphates (high P ₂ O ₅ content)							
UP-16	4.0	210	0.49	131	45	0.33	31
UP-91	3.3	188	0.65	90	61	0.55	85
Silico-phosphates (containing SiO ₂)							
UP-160	3.8	186	0.69	80	56	—	95
H-310	3.9	198	0.80	80	67	—	95
Z-85	3.2	184	0.84	93	69	—	85
Commercial phosphates							
LHG-5	4.1	135	0.77	98	64	0.64	102
LHG-8	4.0	150	0.58	112	52	0.49	63
LG-750	4.0	145	0.62	114	52	0.50	62
Q-100	3.8	180	0.78	96	73	—	71

^a Assuming 40-μm defect size.

sizes with high optical quality. The broad neodymium absorption bands give good pump efficiency. In addition, some phosphate glasses can be made athermal, which means that they exhibit no change of optical path with temperature. Their main disadvantage is low thermal conductivity, which limits their application to low repetition rate systems.

For the HAP laser application we desire laser glasses with long neodymium fluorescence lifetimes and a high capability of withstanding thermal shock. Schott Glass Technologies, Inc., and the Hoya Corporation have been studying new phosphate laser glasses in support of the ICF program.⁶⁰⁻⁶² They have concentrated on identifying glasses with long neodymium fluorescence lifetime at high neodymium concentration, which is also a requirement for HAP lasers. An additional requirement for HAP lasers is a high resistance to thermal shock. Therefore, we have commissioned measurements of the thermomechanical properties of many laser glasses.

Industrial producers of laser glass have emphasized phosphate glasses over silicate glasses because silicates have a more rapid decrease of neodymium fluorescence lifetime with increasing neodymium concentration than phosphates. Schott and Hoya have investigated some promising phosphate glasses that are different from the usual commercial products. Schott^{60,61} undertook an extensive study of ultraphosphate compositions, which are glasses containing a high percentage of P₂O₅ (> 70 mole %). Hoya has concentrated on

silico-phosphate glasses, which contain about 10 mole % of SiO₂. Both these glass types show long neodymium lifetimes at high neodymium concentrations (Table 9-8; see also Table 6-8 in Ref. 62). Furthermore, their thermal shock figure of merit is comparable with that of the best commercial glass available, LHG-5. Schott and Hoya are planning large melts of these new glasses to see if they are amenable to commercial production.

For our experiments in the zig-zag and flow-cooled disk test systems we ordered a number of 20- × 70- × 2-cm plates of LHG-5 from Hoya. Two plates of LHG-5 doped with 5 × 10²⁰ neodymium ions/cm³ were received from Hoya in 1985; four additional plates are on order.

Author: S. E. Stokowski

Growth Techniques for Laser Crystals

Introduction. Laser crystals can be grown by a variety of methods, including Czochralski, vertical Bridgeman, heat-exchanger, horizontal-zone Bridgeman, and float-zone. The Czochralski method is used in the commercial production of crystals because it gives the largest crystals and the highest quality. The method consists of dipping a seed crystal into a melt that is maintained at the fusion temperature and then pulling a crystal directly from the melt. Figure 9-49 shows the basic components of a Czochralski growth

system. Heat is generated at the iridium crucible wall by means of a magnetic field produced by an external induction coil. The lowest temperature in the melt is at the crystal-melt interface because of conduction through the crystal and a higher radiative cooling rate near the melt center. Melt convection, both natural and forced, is critically important for controlling the growth morphology. Forced convection is provided by rotating the crystal, typically at 5 to 40 rpm.

As we scale up laser crystal growth, we need to determine what limitations there are to increasing crystal size, and whether any corresponding modifications of the Czochralski system are necessary. Unfortunately, the crystal growth industry has been slow in making progress because of a heavy reliance on empiricism. A rigorous quantitative understanding of heat and fluid flow in the Czochralski system has not yet been developed. We have therefore begun to develop an in-house capability to conduct research and process development in large-dimension crystal growth by the Czochralski method.

We are also investigating new techniques for growing large laser crystals. In 1985 we supported work at the Allied Corporation to develop a method of adding dopant ions to a Czochralski melt while the crystal is growing. In another program, we had Crystal Systems test the feasibility of growing a garnet laser material, GSAG:Cr,Nd , by the heat-exchanger method, which they use commercially for growing 8-in.-diam sapphire boules. These programs were partially successful, but some new problems were uncovered that may limit the usefulness of these new processes. A detailed evaluation of these programs is given in the rest of this article.

Czochralski Research and Development

During 1985 we became aware of a lack of quantitative understanding of the Czochralski growth method for oxide crystal growth. This deficiency is particularly troublesome in light of our requirements for increasingly large crystals for high average power lasers. The problem exists because the laser crystal industry does not yet have a large enough commercial market to support the critical mass of people necessary to generate substantial improvements in the Czochralski method.

Fig. 9-49. Functional components of a Czochralski furnace.

In contrast, the understanding of heat and fluid transport during the growth of silicon and GaAs is comparatively advanced due to developments in the last ten years, which include advances in finite element techniques as well as increased availability of advanced computing hardware. However, because of insufficient commercial need, these methodologies have not been applied to oxide crystal growth. As a consequence, crystal growers are still vexed by problems such as garnet helix formation and the difficulty of continuous melt replenishment. These problems might be solved by using existing analytical models to develop a better understanding of the interaction of heat and fluid transport with crystal growth kinetics.

For this reason we have recently begun to develop an in-house capability for growing garnet crystals by the Czochralski method. Planned activities include the development of industrial processes as well as basic research, which will include the investigation of advanced physical/computational models for crystal growth. The primary objectives of this work are to

achieve an understanding of heat and fluid flow in an experimentally verifiable configuration and to apply this understanding to problems of immediate interest: (1) controlling heat rejection from a growing boule to minimize a tendency toward helical growth; and (2) developing a means for continuously replenishing a melt to maintain a constant melt level.

The technical approach is as follows:

1. Acquire a state-of-the-art Czochralski machine capable of growing 5-in.-diam gallium or aluminum garnets. The Material Progress Corporation, following our specifications, constructed such a unit, which was delivered to LLNL in 1985.
2. Acquire the FIDAP code, an advanced fluid dynamic computer program designed for analysis of fluid and heat flow in Czochralski systems.
3. Through a research subcontract to MIT, adapt an existing thermocapillary model for Czochralski growth for purposes of developing process control strategies for cylindrical GSGG:Nd,Cr and flat-interface Nd:YAG materials.

FIDAP is a versatile, finite-element, free-boundary code for solving 2-D or axisymmetric fluid-flow and heat transfer problems. The MIT model neglects fluid-dynamic effects, but is capable of analyzing the radiative exchange in the gray-body furnace enclosure and determining the steady-state shape of the meniscus, which are the primary factors controlling high-temperature oxide growth. These models are complementary and will be used together in the design of processes to

control the morphology and composition of laser garnets. The Czochralski unit now in place at LLNL will provide experimental verification of the models.

Czochralski with Dopant Addition During Growth. Reducing the zonal variation of chromium content in GGG:Cr,Nd laser glass is an important goal. While chromium has a nearly unity distribution coefficient in GSGG,⁶³ its distribution coefficient in GGG is much larger, approximately 3.5. This results in a pronounced gradient of chromium concentration in GGG:Cr,Nd crystals grown by the standard Czochralski technique. This concentration gradient must be minimized to optimize laser performance.

A modified Czochralski technique for the growth of this crystal has been under development at the Allied Corporation under contract to LLNL. The method circumvents the greater-than-unity distribution coefficient by controlled addition of chromium to the melt during growth.⁶⁴ The method may be of interest in the growth of other chromium-doped crystals, since a greater-than-unity distribution coefficient is not unique to GGG.

The technique has been successful in maintaining the chromium concentration constant to within 1% in several crystal growth runs. However, the process as currently configured creates a large number of iridium inclusions in the material. Methods of circumventing this side effect are under study.

Heat Exchanger Method. The heat-exchanger method (HEM) or vertical gradient freeze method shows some promise for growing large plates of laser material. Crystal Systems has commercialized this method in the growing of large (8-in.-diam) sapphire (Al_2O_3) boules. In this method the crystal is grown by freezing the melt slowly up from the crucible bottom (Fig. 9-50). The temperature is controlled by cooling the crucible bottom with a stream of helium gas and by heating the crucible radiatively in a graphite furnace.

One substantial advantage of HEM is that growth striations are in the plane of the subsequently cut plate. These striations result from slight variations in the crystal composition and cause refractive-index variations and stress birefringence. Ideally, we would like to propagate our laser beams perpendicular to these

Fig. 9-50. Heat exchanger method for growing large crystals.

striations. With the striations in the plane of the plate this is easy to do with both the zig-zag laser and the flow-cooled disk geometry. On the other hand, crystals grown by the Czochralski technique have striations perpendicular to the disk plane. In the disk laser the beam will propagate nearly parallel to these striations, which may result in unacceptable scattering loss.

Realizing the high potential benefits of producing large plates by HEM, we contracted with Crystal Systems to develop a method for growing large garnet crystals. Because their system operates in a reducing atmosphere, they could not grow gallium garnets because of the volatility of gallium suboxide under those conditions. Therefore, we chose GSAG:Nd,Cr as a test crystal for this system.

Crystal Systems made 12 growth runs in 1-in. tungsten crucibles charged with GSAG. The results were not encouraging. Although Crystal Systems was able to produce crystals with sizes approaching 1 in., the resulting absorption spectra were unusual. Figure 9-51 shows a typical spectrum of an HEM-grown GSAG:Nd,Cr crystal. The neodymium lines match those observed in a Czochralski grown boule. However, there is an appreciable blue-UV absorption, which makes the crystal appear yellow. Furthermore, the chromium bands are very weak, even though the melt contained sufficient chromium to give about $1\text{--}2 \times 10^{20}$ chromium ions/cm³. This low chromium concentration was confirmed by spark-source mass spectrometer analysis. The mechanism for the chromium loss is unknown at this time.

Author: J. F. Cooper

Major Contributor: S. E. Stokowski

Slab Fracture

Introduction

Catastrophic component fracture from thermal stresses constitutes a fundamental limitation on operation of solid state lasers at medium- and high-average power. Solid state slab lasers using glasses and crystalline materials as the active elements have been constructed at 0.1 μ and at several

Fig. 9-51. Absorption spectrum of a GSAG:Nd,Cr sample grown by the heat exchanger method.

other laboratories. Although we have yet to fracture a slab in service, a number of these devices have been pumped to failure elsewhere. Informal reports indicate that the slabs have fractured at stresses lower than those expected by the investigators. This tendency results from inappropriate estimates of design stresses for large components, based on fracture data from small specimens. Scaling of these small-specimen fracture data to determine operating parameters for large slabs of brittle materials requires thorough component stress modeling and a careful probabilistic treatment of fracture data.

Fracture Behavior of Slabs

Introduction. In discussing slab fracture behavior, it is important to review the criteria for slab fracture. We use fracture mechanics to relate component strength to the depth of the most severe flaw. We then analyze the nature of these strength-controlling flaws and the manner in which they are formed during optical fabrication. Subsequently, we develop an outline of the sequential process that produces the commonly observed fracture morphology, in which the slab is nearly broken in half along its center plane. We then develop slab failure probability expressions and give specific examples for the scaling of fracture data for several interesting laser materials: GSAG and Hoya Glass Corp. LHG-5 phosphate laser glass. These fracture data are assessed to assign appropriate operating stress as for large slabs.

Crack Propagation Criteria. Fracture occurs when the surface stress exceeds the fracture strength σ_f of the component. For a given thickness, the highest thermal power can be withstood by components

for which the thermal stress-resistance parameter R_T is maximized⁶⁵:

$$R_T = \frac{\sigma_1 \kappa (1 - \nu)}{\alpha E} \quad (41)$$

where κ is the thermal conductivity, ν is Poisson's ratio, α is the thermal expansion coefficient, and E is Young's modulus. Of the material properties that make up R_T , the component strength σ_1 is the only parameter that varies widely within a given class of materials.^{66,67} Current solid state laser materials can all be classified as brittle; these materials have strong chemical bonds, giving high theoretical fracture strengths, yet actual components fracture at stresses on the order of one thousandth of their theoretical strength. This discrepancy is due to the low inherent resistance to crack propagation in brittle materials, a property characterized by the fracture toughness K_{Ic} .

Unlike designing with ductile materials, brittle laser material designs require a detailed understanding of the statistical nature of fracture in brittle materials. This understanding is required to comprehend

the variability in strength and to predict the tendency to lower strengths for larger components. This variability in the strength of components that have been fabricated by identical processes is due to the distribution in the sizes of the many small flaws created during fabrication. In metals, the voids and machining flaws that inevitably exist in components do not significantly decrease the observed strengths; this is because crack-tip plasticity reduces the stress concentration at crack tips. The situation is entirely different in brittle materials, which have limited or no crack-tip plasticity. In the brittle laser materials of interest, small machining flaws can concentrate the thermal stress developed during laser operation to a critical level that is capable of initiating unstable crack propagation and component fracture. The component strength σ_1 is related to the depth of the largest surface flaw a by the fracture toughness, using the standard relation^{68,69}

$$\sigma_1 = YK_{Ic}(a)^{-1/2} \quad (42)$$

where Y is a crack geometry factor, near unity. Equation (42) defines the stress at which the largest flaw will become energetically unstable and propagate catastrophically. Clearly, the strength of a component is maximized by a reduction in the size of the largest flaws in the area of high tension. These flaws originate during preparation of the optical surface; a discussion of the manner in which the flaws develop and their characteristics follows.

Nature of Surface Defects. Surface finishing of brittle optical components commences with a multipoint grinding operation in which material is disrupted and removed by the interaction of sharp, rigidly held particles with the component surface. The process has been studied extensively by Marshall and others.^{70,71} Above a threshold load (~ 1 N), penetration by the indenting particles creates a characteristic crack system consisting of median and lateral components. The cracks originating from machining grooves can be conceptually modeled as a linear extension of the axisymmetric crack system introduced by a sharp static indenter (see Fig. 9-52).

The cracks evolve during recovery of the underlying material as the penetrating

Fig. 9-52. (a) Schematic of axisymmetrical cracks induced by Vickers indentation. (b) Top view of indentation cracks showing (horizontal) lateral crack system. (c) A linear extension of the axisymmetrical indentation pattern. Median radial cracks extend from the bottom of the machining groove; lateral cracks are also shown.

particle is removed. During penetration an impression is created whose volume is accommodated by an elastic field surrounding a localized plastic zone (Fig. 9-53). The residual stress in the irreversibly (plastically) deformed zone is tensile, whereas the surrounding elastic volume is in residual compression. The cracks extend as the penetrating particle is removed, and the constraining elastic stress drops. In the severe contact loads typical during grinding, chipping and material removal occurs by the intersection of the lateral components of the crack system with the surface or with cracks from adjacent penetration sites. Although the lateral cracks generally control material removal during grinding, the *median* portion is the most important element of the crack system from the strength perspective because the median cracks are oriented normal to the surface, and thus constitute the most severe stress concentrations that control the strength.

An important element of this model is the lack of correlation between the depth of the shallow machining grooves, which are subsequently removed during optical polishing, and the depth of the underlying median cracks, which can be quite deep and may not be removed during polishing.

In contrast to the relatively brutal mechanical operation of fixed-abrasive grinding, fine optical polishing is a combination of a mechanical process and a chemical reaction. The polishing particles interact with the surface and produce a stress-induced chemical reaction between the surface and the polishing medium in the presence of water. The hydrolyzed surface then repolymerizes as the contact pressure is relieved, resulting in a flat, amorphous, hydrated layer. This gel layer can smooth over the grinding damage without necessarily removing it. The resulting high-quality optical surface may have inferior strength. Furthermore, the crack-opening displacement of these defects is typically less than the wavelength of light. Hence, the defects may not be detectable by optical means. The strengths of seemingly identical components may thus vary substantially from one another.

Fracture Morphology. When the stress at the slab surface reaches the value required to propagate the largest surface flaw, the slab will fracture. A characteristic

morphology is evident in slabs that have been pumped to failure during high-power, high-repetition rate operation (see Fig. 9-54). These demarcations consist of several crack systems running perpendicular to the large faces in a random pattern that join a striking broad crack plane, which extends parallel to the large faces along the slab center plane. The description of the slab being "nearly broken in two" is common because these crack planes typically occupy such a large area along the slab center plane. These features on the fractured amplifier slabs can be interpreted to illustrate the sequential failure process.

The stress state on the large faces of the slab is equi-biaxial tension degenerating to uniaxial tension at the edges. This surface tensile stress will propagate a flaw located on the surface or along the edge when the stress intensity for the worst flaw reaches the critical value K_c given by Eq. (42). This relation energetically defines the regime in which the crack will propagate unstably and at speeds approaching the acoustic velocity in the material. We assume here that a single flaw at one of the large surfaces is the critical defect [see Fig. 9-55(a)]. For well-prepared components we might expect multiple crack propagation induced by the superposition of the applied stress and a shock wave created by the first propagation event. When the crack begins propagating, the slab is no longer constrained (as it was before the crack

Fig. 9-53. Scanning electron micrograph of a groove inscribed by a Vickers indenter in soda-lime glass at 5 N load illustrating a shallow machining groove, lateral cracks, and a large median crack. If only the machining groove and the lateral cracks are removed by lapping and polishing, the deep median crack will remain, leading to degraded component strength.

Fig. 9-54. Fractured LHG-5 phosphate glass slab, Spectra Technology, Bellevue, Wash. Failure along the center plane is apparent.

Fig. 9-55. (a) Schematic of flaw in zig-zag slab, side view. (b) and (c) Crack extending to midplane where it changes orientation to parallel to the slab faces, to accommodate unconstrained thermal stresses. (d) Schematic of fractured slab illustrating large center plane crack.

formed) and can now deform to accommodate the thermal stresses. Recall that the crack propagates at near the acoustic velocity, about 10^6 times faster than the heat transfer rate: the temperature distribution is essentially fixed within the time frame of these fracture events. The critical defect propagates normal to the slab surface toward the center plane in a macroscopic stress field, whose sign changes from

tension at the surface to compression near the center plane. Crack propagation in brittle solids occurs almost exclusively by Mode I (opening mode) crack surface displacement,⁷² where the crack plane separation is normal to the maximum principal tension. The propagating crack follows a direction that minimizes shear loading.⁷² As the crack propagates from the surface towards the compressive field, the orientation of the principal tension becomes normal to the surfaces of the large faces (to allow the now unconstrained thermal stresses to be relieved), and the crack front changes orientation; the center plane crack forms [Figs. 9-55(b) to 9-55(d)].

Despite the interesting slab separation along the centerline, the only critical issues in regard to slab design are the state of stress at the surfaces, the size and nature of the surface flaws, and the presence or absence of water. As soon as the stress reaches the value required to propagate a crack, the useful life of the component is over. The subsequent fracture events, in this context, are irrelevant.

Fracture Testing and Scaling

Introduction. In the absence of a dependable nondestructive technique for assessing the extent of subsurface damage, we have adopted a method for estimating a safe operating stress, based on a Weibull extreme-value statistical treatment of fracture data. Following is a description of the methods used for the fracture tests and the results. We also describe the development of the Weibull method and interpret its results.

Test Methods and Results. We performed fracture testing on GSGG and on LHG-5 phosphate glass. We stressed small samples of GSGG ($3.0 \times 0.6 \times 0.3$ cm) in four-point flexure to failure. The edges of the specimens were carefully polished to give rounded edges, thereby reducing the number of edge failures. The data presented in Table 9-9 are restricted to surface-initiated fractures only. These surfaces were prepared by standard optical finishing wherein the surface was subjected to a bound-abrasive grinding, 30-, 15-, and 9- μ m loose-abrasive grinding steps, and a polishing, step of about 30 min using Linde B alumina. This

technique can give optically flat surfaces, but we made no particular effort to minimize subsurface damage.

We prepared a second set of GSGG specimens, using a variety of techniques to minimize subsurface damage. We removed large amounts of material at each grinding step to reduce the damage from the previous step. An acid-etching treatment was incorporated before the 9- μ m loose-abrasive grinding to remove the damaged layer.⁷³ The samples for these tests were 2.5-cm-diam disks used in a ball-on-disk-on-3-ball configuration, similar to the concentric ring method,⁷⁴ which gives the equi-biaxial tension correctly simulating the stress state in the laser-slab faces. A further advantage of this test method is that the edges are not under stress. For characterization of the surface flaw population and for the development and evaluation of techniques to minimize subsurface damage, this test configuration is preferred; it is used whenever sufficient material of the appropriate geometry is available.

The data for these specially polished disks cannot be directly compared with the uniaxial data. We expect the strength of the conventionally polished specimens fractured in uniaxial tension to be approximately 10% higher than those from equivalent biaxial tension tests.⁷⁵ Consequently, we have reduced our estimate of the strengths for samples fractured in uniaxial tension by 10%; this will reconcile these data with the actual biaxial stress state at the slab surfaces. Results for these tests are given in Table 9-10.

When we compared two sets of fracture data for LHG-5, one sample set had a standard-high-quality optical finish. The second sample had an experimental polish in which the samples were etched after each polishing step in an attempt to remove the subsurface damage (Table 9-11). The surface finishing for these specimens was performed by Zygo Corporation using proprietary techniques. These strength tests were performed on 5-cm-diam disks in the ball-on-disk-on-3-ball configuration.

We also examined the effect of the size of the edge radius on fracture strength for LHG-5 phosphate glass. The result, shown in Fig. 9-56, illustrates two trends: rounded edges are preferred to distinct bevels, and large bevels are in general stronger than

Rank	Φ	σ (MPa)
1	0.063	91
2	0.188	107
3	0.313	114
4	0.438	128
5	0.563	154
6	0.688	161
7	0.813	194
8	0.938	232

Rank	Φ	σ (MPa)
1	0.033	202
2	0.100	241
3	0.167	322
4	0.233	341
5	0.300	348
6	0.367	364
7	0.433	406
8	0.500	425
9	0.567	437
10	0.633	461
11	0.700	472
12	0.767	498
13	0.833	610
14	0.900	612
15	0.967	639

Rank	Φ	Standard polish σ (MPa)	Experimental polish σ (MPa)
1	0.056	200	242
2	0.167	245	282
3	0.278	248	348
4	0.389	254	392
5	0.500	264	406
6	0.611	295	419
7	0.722	307	444
8	0.833	316	496
9	0.944	363	511

Table 9-9. Uniaxial fracture data for GSGG, standard optical polish $A = 1.2 \text{ cm}^2$.

Table 9-10. Biaxial fracture data for GSGG, special polish, $A = 1.1 \text{ cm}^2$.

Table 9-11. Biaxial fracture data for LHG-5 phosphate glass, $A = 0.9 \text{ cm}^2$.

small, sharp bevels. The fracture data are now interpreted in the context of scaling to slab-laser component sizes.

Weibull Statistical Treatment. The severity of surface flaws introduced by optical fabrication follows a statistical distribution whose characteristics critically depend on the specifics of the mechanical surface preparation. For example, the strength of eight optically polished small GSGG samples which had been prepared identically and simultaneously, varied by more than a factor of 2, from 91 to 232 MPa.⁷³ As a further example, the mean strength of identical optically polished GSGG samples, which had been damaged, acid-etched to remove most of the subsurface damage had a mean strength (2320 MPa) over 15 times greater than the mean

Fig. 9-56. Fracture stress in four point bend configuration vs bevel (or radius) size for LHG-5 phosphate glass. Trend indicates higher strengths for larger bevels, and higher strengths for rounded edges compared to sharp bevels.

strength of these optically polished specimens (148 MPa). Our methods for reducing this subsurface damage are the subject of a separate publication.⁷³ These examples illustrate the variability in component strength due to the statistical variation in the sizes of the flaws introduced by the mechanical abrasion associated with optical surface preparation.

An effective dependence of strength on component size occurs because large components have more flaws than small components. Within a given distribution of flaw sizes, which result from a given preparation technique, the likelihood of having a particularly severe defect is clearly higher for large components than for small ones. The strength scales with the volume, area, or edge length of a specimen, depending on which flaw distribution(s) govern the strength.

Since the opening displacement of the crack planes below the surface is often too small to permit optical detection by light scattering, no reliable method for characterizing the extent and depth of the damaged layer exists. Thus, the precise

strength of a given component cannot be predicted before the component is fractured. We therefore adopted a statistical method for assigning a failure probability (at a given stress) to identify reasonable operating parameters for large-slab lasers. This approach is based on the assumption that the component will fail when the weakest element, that is, the element containing the largest crack, fails. The failure probability, Φ , is expressed as⁷⁶

$$\Phi = 1 - \exp \left[- \int_V dv \int_0^{\sigma} g(s) ds \right] \quad (43)$$

where V is the component volume, dv is the element volume, σ is the stress, and $g(s)$ is usually fitted by a two-parameter Weibull distribution such that Eq. (43) becomes

$$\Phi = 1 - \exp \left[- V^*(\sigma/S_0)^m \right] \quad (44)$$

where S_0 and m are scale and shape parameters and V^* is the ratio of component volume to fracture sample volume. If we instead assume that the failures all

occur from surface flaws, then the general equation becomes

$$\Phi = 1 - \exp \left[-A^* (\sigma/S_0)^m \right] \quad (45)$$

where A^* is the ratio of component surface area to fracture sample surface area. Equation (45) is typically solved graphically; the measured failure stresses are ranked in increasing order and assigned a failure probability using

$$\Phi = \frac{n - 0.5}{N} \quad (46)$$

where n is the rank and N is the total number of strength data. Equation (45) is then plotted as $\ln[1 - \Phi]$ vs $\ln \sigma$; these units are chosen to linearize the exponential in Eqs. (43) to (45). The slope is m and for a zero abscissa, the ordinate is $\ln S_0$. The graphs are called Weibull plots and illustrate the trend in likelihood-of-failure as a function of stress. Once these parameters have been determined, an estimate of the failure probability of large components can be made by substituting the appropriate area ratio into Eq. (45). With this approach, we can assign appropriate operating parameters for large, brittle components on the basis of fracture data from small samples.

At the intersections of the large slab faces with the shorter faces (the edges), the biaxial tension degenerates into uniaxial tension. The stress at the edge is a fraction of the maximum stress in the slab faces, but depends on cavity and slab mounting particulars; for many designs the tension at the edges is significant. This fact, coupled with the generally poorer surface finish at the edges, requires attention to avoid edge-initiated rupture. If the failures are assumed to all initiate from edge defects, the failure probability of a component is

$$\Phi = 1 - \exp \left[-L^* (\sigma/S_0)^m \right] \quad (47)$$

This expression is similar to Eq. (44), except that the volume ratio is replaced by the ratio L^* of the length of edge exposed to high tension in the component to that in the test samples.

To ascertain the probability of component failure from both edge and surface

defects, the flaw populations are treated separately by calculating individual failure probabilities for each mechanism at each appropriate stress level, and then combining the two probabilities. For the slab laser, the combined probability Φ of fracture from edge defects ϕ_e and surface defects ϕ_s is

$$\Phi = 1 - (1 - \phi_e)(1 - \phi_s) \quad (48)$$

These expressions are used to scale the fracture data to assign safe operating stresses.

Interpretation of Results. We illustrated these scaling procedures using published data for optical surfaces on GSGG (Ref. 73) and unpublished work by the author on LHG-5 glass. Data on GSGG are given in Tables 9-9 and 9-10. The Weibull parameters for the standard optical finish are $m = 3.68$ and $S_0 = 164$ MPa; the special finish yields values of $m = 3.72$ and $S_0 = 472$ MPa. Note that the slope of the two Weibull plots for GSGG are nearly identical; the data for the special polish are simply displaced to higher strengths. These results can be scaled to estimate the probability of failure for larger slabs at a chosen stress.

For a slab in which $0.5 \times 1.2 \times 10$ cm is being pumped, the area of highest tension (the large faces) is 24 cm^2 . Assuming that the slab can be fabricated by processes identical to the one described above, we can estimate the stress for which one-tenth of the components will fail ($\Phi = 0.1$) for each surface treatment. These stresses are 34 MPa for the conventional polish and 110 MPa for the special polish.

Much larger GSGG slabs have been proposed for high-average-power systems. A $1 \times 10 \times 20$ -cm slab that is being pumped along 15 cm of its length has an area of 300 cm^2 in high tension. For $\Phi = 0.1$, the design stress is 19 MPa for conventional, and 55 MPa for special polishing. The stresses for several failure probabilities and slab sizes in GSGG are summarized in Table 9-12.

Our results for the optically polished LHG-5 phosphate glass give $m = 6.77$ and $S_0 = 296$ MPa. For the special polish, $m = 4.79$ and $S_0 = 423$ MPa. The area under high tension depends on the load, and averages 0.9 cm^2 . For small slab-lasers

Slab Fracture

Table 9-12. Calculated failure probabilities, GSGG.

Pumped slab dimensions	Φ	Design stress (MPa)	
		Standard polish	Low subsurface damage polish
0.5 × 1.2 × 10 cm	0.5	56	180
	0.1	34	110
	0.01	20	58
1.0 × 10 × 20 cm	0.5	32	92
	0.1	19	55
	0.01	10	29

Table 9-13. Calculated failure probabilities, LHG-5 phosphate glass.

Pumped slab dimensions	Edge failure probability (ϕ_e)	Design stress (MPa)	Surface failure probability (ϕ_s)	Design stress (MPa)	
				Standard polish	Experimental polish
0.5 × 1.2 × 10 cm	0.5	113	0.5	175	201
	0.1	75	0.1	133	136
	0.01	45	0.01	94	83
1.5 × 10 × 50 cm	0.5	79	0.5	101	93
	0.1	53	0.1	76	63
	0.01	31	0.01	54	38

Table 9-14. Uniaxial fracture data for LHG-5 phosphate glass-edge failures. Bevel sizes 0.9 to 1.25 mm; $L = 1.9$ cm.

Rank	Φ	σ (MPa)	Bevel size (mm)
1	0.071	135	1.0
2	0.214	185	0.9
3	0.357	195	0.9
4	0.500	210	0.9
5	0.642	245	1.25
6	0.786	250	1.0
7	0.929	305	1.0

(again we used 0.5 × 1.2 × 10 cm), $\Phi = 0.1$ gives a design stress of 133 MPa for the standard polish and 136 MPa for the experimental polish. For a large laser, 1.5 × 10 × 70 cm, in which 50 cm is pumped, $\Phi = 0.1$ gives 76 and 63 MPa for standard and experimental polishing, respectively. Table 9-13 summarizes these findings. These operating stresses are somewhat higher than those found for optically polished GSGG. Although the fracture toughness of GSGG is higher than that of phosphate glass, the degree of subsurface damage is also higher in this group of GSGG samples, particularly in those in which no effort was made to minimize the damage. Rather than being a general result, this illustrates the importance of the specific grinding and polishing procedure in determining component strength.

Extrapolating the scaling over several orders of magnitude in area becomes suspect, particularly when the extrapolation is based on only a few data. Notice that for scaling to low failure probabilities and large slabs, the design stress for the

experimental polish on LHG-5 phosphate glass is lower than that for the standard polish, due to the higher shape factor ($m = 6.67$) of the least-squares-fit to the nine data points. As a guideline, we should test about the same total area of sample as the area of the component; this condition is clearly not satisfied by our data set. Caution should be exercised in assigning design stresses when the number of data is limited. For example, using the lower shape factor ($m = 4.79$) from the experimental-polish data set reduces the design stress estimates (at $\Phi = 0.1$) for the standard polish from 133 to 95 MPa for the small slab and from 76 to 44 MPa for the large slab. Using these reduced values is prudent. The results of these Weibull scaling procedures illustrate the inaccuracy in the common assumption that large components will have strengths similar to those of small fracture specimens.

For a particular bevel size, the fracture data can be plotted to determine the statistical factors. For carefully polished edge bevels in the size range 0.9 to 1.25 mm, the shape and scale factors are $m = 4.57$ and $S_0 = 238$ MPa for seven specimens with an edge length, 1.9 cm, exposed to uniformly high tension. For edge-initiated fracture, the design stresses at the edges of a 0.5- × 1.2- × 10-cm slab are 75 MPa, for $\Phi = 0.1$. For the large slab (1.5 × 10 × 50 cm) the design stress is 50 MPa for $\Phi = 0.1$. Results for other failure probabilities are displayed in Table 9-14.

For a 1.5- × 10- × 50-cm LHG-5 slab, fabricated with a high quality optical polish, the nominal stress for 80-kW lamp input power is about 96 MPa at the slab faces; the stresses at the edges are estimated to be one half this level, 48 MPa. The individual failure probabilities are then $\phi_e = 0.07$ for 1-mm polished edge bevels and $\phi_s = 0.42$. The combined failure probability given by Eq. (48) is $\Phi = 0.46$.

Brittle component design involves careful analysis of the stress state in the component and adequate fracture testing to characterize the flaw populations of concern. Surface flaws induced by mechanical abrasion of the component during preparation of optical surfaces are the strength-controlling flaws. We analyzed fracture data using Weibull statistics; this enabled us to assess the effect of slab size and

surface finish on the probability of component failure. Specific examples for LHG-5 phosphate glass and GSGG illustrate the trend toward lower strength for large slabs. We conclude that the surface stress, the component size, and the characteristics of the surface and edge flaws are the important parameters in designing slab lasers to avoid catastrophic failure.

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Evanescent-Wave Control

Introduction

The laser beam passes through a zig-zag amplifier by TIRs to cancel thermal optic distortions within the laser material. However, the electromagnetic field associated with the laser pulse is not discontinuous at the glass-coolant interface. The field penetrates the coolant for a distance of a few wavelengths, and this portion of the field is called the evanescent wave. Imperfections at the interface or bubbles in the coolant can lead to coupling losses and phase distortions of the laser beam. A clear coating that contains the evanescent-wave can protect the zig-zag laser from such imperfections.

Several low-index materials that meet most of the requirements for evanescent wave coatings (EWC) have been identified. However, because the long-term effects on these materials under operating conditions are presently unknown, we also considered the use of strengthened laser glass mounted on low-index seals with the coolant liquid serving as the EWC.

General Considerations

Evanescent-wave coatings serve several purposes. They provide:

- High reflectivity for the laser beam without phase distortions.
- Low effective reflectivity for ASE.
- A barrier between the laser slab and seal.
- Protection of slab surfaces from contamination and corrosion by the coolant.

- Improved fracture strength by tailoring the stress distribution in the surface layer and by preventing chemically enhanced crack propagation.

Reflective coatings, liquids, and low-index coatings all have the properties needed for EWC. Highly reflective multi-layer dielectric coatings do not appear to be suitable because they exhibit enhanced ASE losses and loss in pump light due to transmission losses at various wavelengths and angles of incidence. Liquids could be used either as thin layers or as coolants. The simplest solution is to use a clean coolant with strengthened, low-solubility laser glass mounted on small seals. The EWC requirements for low-index seals are much less demanding than those for low-index coatings. The seals must be resistant to damage from laser light, be compatible with the coolant, and have an index of refraction close to that of the coolant. It should be possible to obtain small seals in the form of wetted knife edges; the mechanical strength would be provided by the knife edge and the seal could be quite resilient. Clean coolant systems will be required even with EWCs to avoid contamination of surfaces and loss in pump efficiency. Concerns about the solubility of the laser materials and crack propagation can be addressed by selecting different coolants, such as fluorocarbons, alcohols, or glycol mixtures.

We investigated low-index dielectric and plastic coatings, primarily as protection for the slab surfaces. The required index of refraction of these coatings depends on the refractive index of the laser material, the thickness of the coating, and the geometry of the laser slab. The index must be low enough for TIR, but not so low as to enhance ASE losses. The coatings must adhere to the laser material and be implementable at reasonable cost. In addition, these coatings must be transparent, damage resistant to flashlamp and laser light, and compatible with the coolant and seal. They must have low, uniform thermal resistance, low thermal expansion, and a high modulus of elasticity for stress control.

Index Requirements of EWC

For the zig-zag laser geometry with slab faces at Brewster's angle, the index of

refraction for the coolant or EWC must be below a critical value that is determined by the beam-fill factor and allowable reflectivity at the beam entrance and exit faces. The definitions of angles, dimensions, and pertinent relations are illustrated in Fig. 9-57. This geometry provides complete beam overlap for face angles $\gamma < 90^\circ - \theta$, and is, in essence, a double-pass system with respect to energy extraction for very short laser pulses. For laser pulses that are long relative to the traversal time between reflections, the beam overlaps and the effective beam intensity is essentially twice that determined from the unfolded beam dimension w_i . For face angles $\gamma > 90^\circ - \theta$, the beam overlaps only the fraction $f < 1$ of the pumped laser volume, while the volume fraction $(1 - f)^2$ is not intercepted at all and the volume fraction $f - f^2$ experiences single-pass extraction.

The angles of the zig-zag laser with Brewster faces are confined to a narrow regime as shown in Fig. 9-58 for LHG-5 laser glass with an index of 1.53. Figure 9-58 illustrates the dependence of the reflection angle θ within the zig-zag laser on the incident angle α for various face angles γ ranging from 25° to 35° . The dashed lines indicate the reduction of the beam

overlap fraction from $f = 1$ to 0.8, and the single dotted line represents the reflectivity of each slab face for parallel polarized light. For full-beam overlap ($f = 1$) and incidence at Brewster's angle ($\alpha = \alpha_B$), the reflection angle θ is restricted to less than 61.5° ; it depends on the face angle γ in the form $\theta = \beta_B + \gamma < 90^\circ - \alpha_B/2 = 61.5^\circ$. The refractive index of the EWC must be below a critical index n_c , which depends on the refractive index n of the laser material, the reflection angle θ , the thickness x of the EWC, the wavelength λ , and the required intensity reduction I/I_0 in the form⁷⁷

$$n_c \leq \sqrt{n^2 \sin^2 \theta - \left[\frac{\lambda}{2\pi x} \ln \left(\frac{I_0}{I} \right) \right]^2}. \quad (49)$$

The critical indices n_c for evanescent wave coatings with thickness $x = 1 \mu\text{m}$ are shown in Fig. 9-58 for various intensity reductions I/I_0 of $\lambda = 1.06\text{-}\mu\text{m}$ evanescent waves. Doubling the coating thickness is equivalent to squaring the intensity ratio in this plot, e.g., a $1\text{-}\mu\text{m}$ $I/I_0 = 0.1$ coating is equivalent to a $2\text{-}\mu\text{m}$ $I/I_0 = 0.01$ coating or a $3\text{-}\mu\text{m}$ $I/I_0 = 0.001$ coating. With the restrictions $f = 1$ and $\alpha = \alpha_B$, the index of refraction for the EWC must be less than 1.34 for a very thick ($x \gg \lambda$) coating and

Fig. 9-57. Geometrical relations of zig-zag laser with Brewster faces ($N = 4$ bounces).

Fig. 9-58. (a) Various face and reflection angles using LHG-5 laser glass ($n = 1.53$, $\lambda = 1.06 \mu\text{m}$) and (b) the corresponding index requirements of 1- μm -thick EWCs.

less than 1.21 for a 1- μm coating that provides a 1:1000 intensity reduction.

These indices are very low and drastically reduce the choice of materials. However, it is possible to compromise both with respect to beam overlap fraction f and reflectivity R at the entrance and exit faces. For Brewster's angle and full beam overlap, the critical coating index is determined by

$$n_c = \sqrt{0.5n^2 \left(1 + \frac{1}{\sqrt{1+n^2}} \right) - \left[\frac{\lambda}{2\pi x} \ln \left(\frac{I_0}{I} \right) \right]^2} \quad (50)$$

The change in the coating index with beam overlap fraction f and reflectivity R can be approximated by

$$\Delta n_c \approx 0.15n^2(1-f) + 0.2 \frac{R^{1/3}}{n^2} \quad (51)$$

A 10% reduction in f , corresponding to a 1% reduction in maximum extraction efficiency, would raise the critical indices by about 0.03. In addition, a 1% reflection loss per surface would raise the indices by about 0.02 to a total of 0.05.

The critical indices of coatings for various laser materials and the corresponding

critical angles for the glass slabs are shown in Fig. 9-59. They increase almost linearly ($\Delta n_c \approx 0.8 \Delta n$) with the index of the laser material, while the critical angles are confined to narrow ranges.

The angle δ relative to the slab axis is typically in the range of $\pm 10^\circ$ and, if necessary, can be shifted to $\delta = 0$ with relatively minor impact on laser performance. The arrows in Fig. 9-59(b) indicate the direction in which the angles can vary without violating the conditions for $f = 1$ and $R = 0$. However, this flexibility exists only for designing a slab. For an existing slab, the angles vary significantly with the number of traversals N , as shown in Fig. 9-60 for the 320- $\times$ 15-mm LSG-91H slabs with a face angle of 28° . This slab is primarily a 10-bounce slab that requires a coolant or thick coating with $n_c < 1.38$. Increasing the traversals to 12 would require a thick coating with $n_c < 1.31$.

Low-Index Coatings. A number of organic and inorganic materials have been tested so far. The index of refraction of ideal thick coatings must be less than 1.37 for silicate and 1.34 for phosphate laser glasses [Fig. 9-59(a)]. For 1- μm coatings, these two glasses would need indices of

less than 1.31 and 1.28 to reduce the intensity of the evanescent wave to 1%.

Sol-gel and other porous inorganic coatings can have effective indices of less than 1.3. However, they must be

overcoated to prevent the pores from filling with coolant. Also, some materials are soluble in the coolant and need a physical barrier. Overcoats in the form of a thin film of FEP (polytetrafluoroethylene)

Fig. 9-59. (a) Maximum index for 1- μ m-thick coatings and (b) critical zig-zag angles for various indices n of the laser material with full beam overlap and no reflection losses.

Fig. 9-60. Angles, beam width, and EWC index requirements vs N passes of a 28°, 320- $\times$ 15-mm LSG-91H slab ($n = 1.55$).

provide such a barrier, but the glue that is used to affix the tape penetrates the pores and raises the index of refraction of the coating close to that of the glue (1.41). One can actually observe a time varying increase in the index of refraction as the glue creeps into the pores. This effect was observed both on 1- μ m-thick sol-gel coatings ($n_c \approx 1.23$) and 2- μ m-thick CaF_2 coatings ($n_c \approx 1.22$) that have packing fractions of about 50%.

We evaluated two other vapor-deposited coatings. Scanning electron microscopy showed that a 2.3- μ m-thick cryolite (Na_3AlF_6) coating was severely crazed and that a 2.3- μ m-thick AlF_3 coating had a much smoother surface than those of the sol-gel and CaF_2 coatings. Both materials exhibit an effective index of 1.32, which corresponds to packing fractions of 94% and 84% for cryolite and AlF_3 , respectively. AlF_3 is water soluble and needs a protective overcoat. In this case, however, FEP tape can be used for the protective layer because the glue does not penetrate the relatively small pores of the AlF_3 coating. At least, we saw no deleterious effects either during a period of 2 weeks after application or when increasing the pressure on the FEP tape. However, the tape does separate from the substrate if its edges are exposed to water, so the seal at the edge of the FEP tape must be watertight. Although this seal makes mechanical assembly somewhat more complex, use of an FEP tape overcoat is a definitely viable backup option for all porous coatings.

A solution-deposited AlF_3 film $\sim 0.2 \mu\text{m}$ thick protected by an FEP tape exhibited poor reflection characteristics. Either the glue of the FEP tape penetrated the coating or the coating was too thin.

Vapor-deposited dielectric films allow some tailoring of the surface stresses, but plastic films appear more practical and cost less. A sublayer to control stress on plastic films could be added, if necessary.

Organic Coatings. An organic polymer appears to be the most likely candidate material for EWCs. These coatings are water resistant and can be applied uniformly in virtually any thickness desired. We can also tailor a number of the required physical and mechanical properties discussed above. While the high particulate content of commercial polymers is of

great concern, the most difficult requirement to meet is that of refractive index. Very few organic materials possess a low enough index; possible candidates are virtually limited to the fluorocarbon or fluorosiloxane families.

Many of the fluorocarbons are not optically clear, and the difficulties of fabricating materials such as FEP or FEP-related polymers are considerable. The high temperatures required to sinter FEP powder and related materials present too many difficulties to allow them to be seriously considered for application to laser glass at this time. In addition, the limited solubilities of many fluorocarbons represent a considerable barrier to coating techniques relying upon that property.

An experimental polyfluorourethane (FPU) sample was obtained from the 3M Company. One component, Dymar FC-2202 perfluoropolyether diol, has terminal groups of $-\text{OCF}_2\text{CF}_2\text{OH}$, randomly polymerized backbone units of $-\text{OCF}_2-$ and $-\text{OCF}_2\text{CF}_2-$, and a refractive index of 1.30. A second component, L-9183, is a diisocyanate with the same backbone as FC-2202, but it has terminal groups of $-\text{OCF}_2\text{CH}_2\text{NCO}$ and a refractive index of 1.3075.

L-9019 is another type of fluorinated alcohol with backbone units primarily of $-\text{OCF}_2(\text{CF}_3)\text{CF}_2-$, end groups of $-\text{OCF}(\text{CF}_3)\text{CH}_2\text{OH}$, and a refractive index of 1.31. This was the alcohol of choice because its functionality of slightly greater than 2 allowed the formation of a cross-linked film when reacted with L-9183 diisocyanate. We obtained only an oil from the FC-2202 diol/L-9183 reaction because a linear polymer was formed.

In-depth optimization of the system was not possible because only very limited quantities of the components were available. Initially, a 50:50 mixture of L-9183:L-9019 resulted in tacky, weak films containing unreacted isocyanate functionality (as indicated by IR analysis). Increasing the ratio to 55:45, 57.5:42.5, and finally to 60:40, produced a successively drier, stronger film, although further study is certainly justified in this area once larger quantities of materials are available. Because an overcoat is necessary to protect the FPU from water, the slight tack of the film is actually an advantage in

promoting adhesion between the FPU and the overcoat film.

Water immersion tests showed that the cured FPU films developed cloudiness over several weeks. Drying the film under ambient conditions for several hours reversed this process. The cloudiness is attributable to the presence of a second phase (water) within the FPU film. FEP, a DuPont product with a structure related to that of Teflon, was chosen as the overcoat film. A 1-mil film was applied to the cured FPU coating as a water barrier. We showed that, after etching, a silicone sealant (Dow Corning 730) would satisfactorily bond to both the FEP overcoat and the aluminum

structural members of the assembly. This watertight seal prevents coolant leakage.

We also found that McGhan NuSil's 2655 silicone could be sprayed over the FPU as a protective overcoat, although it is permeable to water and provides limited protection in a aqueous environment.

Very good quality FPU films were obtained by spraying a 50% solid solution (in Freon 113 filtered through a 0.5- μm syringe filter) from an artist's airbrush. Multiple overlapping passes produced the most even coating. These coatings were cured in a vacuum oven at 70°C for about 10 h.

Laser damage and flashlamp testing show that FPU is a viable candidate for EWC, although maximum performance can be achieved only if all cleaning, spraying, and curing are done in a clean room because dust is a prime catalyst for the degradation of polymers in laser applications.

Phase Distortions

A major effect of the evanescent wave on zig-zag amplifier performance is phase distortion caused by variations of the index of refraction at the slab-coolant interface. The phase shifts ϕ for reflection angles θ greater than the angle of TIR are given for parallel and perpendicular polarization by⁷⁸

$$\phi^{\parallel} = 2 \arctan \left(\frac{n \sqrt{n^2 \sin^2 \theta - n_c^2}}{n_c^2 \cos \theta} \right) \quad (52)$$

and

$$\phi^{\perp} = 2 \arctan \left(\frac{\sqrt{n^2 \sin^2 \theta - n_c^2}}{n \cos \theta} \right) \quad (53)$$

The phase shifts for various interfaces of a glass with index $n = 1.55$ and air ($n_c = 1$) or water ($n_c = 1.33$) are shown in Fig. 9-61(a) for parallel polarized light. The difference in phase shift for air and water indicates a differential phase shift for an air bubble in water (or, equivalently, a water droplet on an air interface) of about 80° at typical reflection angles of $\theta \approx 62^\circ$. This relatively large phase shift could lead to

Fig. 9-61. (a) Phase shifts for parallel polarized light at various internal reflection angles; (b) and (c) associated intensity modulations for an 80° phase shift of a 1-mm bubble.

large beam modulations for millimeter-size bubbles, as shown in the MALAPROP simulations [Fig. 9-61(b) and (c)]. However, visual observations indicate less severe modulations; the differences could be caused by MALAPROP's perfect symmetry and perfect beam, which may not exist in the experiment. Small bubbles are difficult to observe because diffraction reduces the intensity modulations after relatively short propagation distances. The primary effects of small bubbles will be a loss in

effective beam area equal to the size of the bubbles.

A somewhat different condition exists for an index mismatch of a straight line seal. Figure 9-62(a) shows the differential phase shifts between water and sealants of various indices n_s . At typical reflection angles of $\theta \approx 62^\circ$ or greater, the phase shifts can be kept below 18° for seals with indices ranging from 1.30 to 1.35. The modulations of a 5-cm beam for a phase shift of 10° are shown in Fig. 9-62(b)

Fig. 9-62. (a) Phase shifts between water ($n_c = 1.33$) and sealants of index n_s for $n = 1.55$ laser glass; (b) and (c) associated intensity modulations for a straight-edge seal, $\Delta\phi = 10^\circ$.

and (c) for propagation distances of 1 and 5 m. In this case, the intensity modulations persist for long distances, but can be kept at tolerable levels relative to other beam imperfections.

Flashlamp Exposure Tests

Flashlamp exposure tests of coating materials were conducted by placing the samples in quartz tubing 10 cm from a 50-cm arc flashlamp. Water flowing through this tubing simulated the conditions in the zig-zag laser. The flashlamp was operated at 31 kW (100 Hz), which provides sufficient average power per unit sample area, but relatively low peak intensity. Under nominal conditions of the zig-zag laser [2.1 kV, 1 pulse per second, $(LC)^{1/2} \approx 0.14$ ms], the EWC is exposed to $\sim 1 \text{ J/cm}^2$ and 1 W/cm^2 for single-sided illumination and about twice that within the pump cavity. In contrast, the flashlamp tester under standard conditions provides an average loading of 10 W/cm^2 , but only 0.1 J/cm^2 per pulse with $(LC)^{1/2} \approx 0.2$ ms. Because of the low energy deposition, an additional 800 000 test shots were made at 4 times higher energy, but at the same power. The results are summarized in Table 9-15. PSW 2062 failed, and part of the edge of the urethane coating had a white, cloudy appearance. The FPU and silicone-overcoated FPU samples did not show any deterioration.

However, localized brownish discolorations were observed after exposure of both a FEP and a R2655 silicone-overcoated FPU layer in the zig-zag test bed. These discolorations are most likely caused by contamination, which could be prevented by using cleaner assembly areas.

Laser Damage Tests

A laser beam passing through the FPU coating on a flat LSG-91H sample in 50-ns, $1.06\text{-}\mu\text{m}$ pulses damaged local imperfections at 29 J/cm^2 . The damage sites were small, round, brownish discolorations at dispersed locations that perhaps contained particulates or uncured material. The bulk of the material was undamaged at 29 J/cm^2 . Additional damage tests were made with the coating exposed only to the evanescent wave (Fig. 9-63). At an entrance angle of $\alpha = 26^\circ$, the internal reflection angle is $\theta = 62^\circ$, just 3° above the TIR angle. Damage to imperfections at the glass-FPU interface for this condition occurred at an entrance beam fluence of $\sim 12 \text{ J/cm}^2$. It was somewhat surprising that this damage occurred at a lower fluence, but, because the beam overlaps near the surface, standing wave patterns could lead to a fourfold increase in beam intensity for perpendicular polarized light. For the parallel polarized light used in this experiment, local increases in intensity of

Table 9-15. Flashlamp tests of coating materials for the zig-zag laser.

Material	Cure or deposition	Refractive index	Thickness (μm)	Flashlamp test		Remarks
				Low E	High E	
CaF ₂	Vapor	1.22	2	— ^a	—	FEP glue penetrates
Sol-gel	Solution	1.23	1	—	—	FEP glue penetrates
Cryolite	Vapor	1.32	2.3	—	—	Crazed
AlF ₃	Vapor	1.32	2.3	Passed	—	Water soluble
AlF ₃	Solution	—	~ 0.2	—	—	Poor TIR
FEP	Film	1.34	~ 50	Passed	—	—
FEP glue	—	1.41	—	Passed	—	Separates in water
FPU	70°C 3 d	1.30	>10	Passed	Passed	Soft
FPU + R-2655	70°C 3 d	1.30	>10	—	Passed	Silicone overcoat
FC-721	Dip coat	1.37	>2	—	Passed	Brittle, softens in water
PSW-2121	Ambient	1.38	>10	Passed	Passed	Durable
Fluorosiloxane	Ambient	—	<1	Passed	—	Too thin
Fluoro-epoxy	Ambient	1.41	<10	—	—	—
Sylgard 184	Ambient	1.41	<10	—	—	—
Silicone 88	Ambient	1.41	<10	—	—	—
GE-655	Ambient	1.43	<10	—	—	—
Urethane 35-29x	Ambient	1.47	<10	Passed	— ^b	—
PSW-2019 +	Ambient	1.49	<10	Passed	—	—
PSW-2062	Ambient	1.49	<10	Failed	—	—
Cyano-acrylane	Ambient	1.49	<10	Passed	—	Separates in water

^a Not tested.

^b Scatter at edge.

$2(1 + |\cos 2\theta|) \approx 3$ could be expected near the surface within the glass prism. In addition, the effective beam area is increased $\sim 6\%$ by refraction. This leads to a local increase in beam fluence of $(3/1.06) \times 12 \text{ J/cm}^2 \approx 34 \text{ J/cm}^2$, which is comparable to the result at normal incidence. However, one would have expected still higher damage values because the beam intensity falls off rapidly within the EWC unless the imperfections are right at the interface.

Even though cleaner applications and better curing could probably increase the damage threshold, the present 12 J/cm^2 for TIR conditions are sufficient for near-term applications.

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Flashlamps for High-Average-Power Lasers

Introduction

Rare-gas flashlamps satisfy most of the requirements of pump sources for HAP lasers. They convert electrical input to optical output with high efficiencies, typically 70 to 80% for xenon-filled flashlamps.⁷⁹ Flashlamps can deliver very large fluences reliably, so large amounts of energy can be stored in the laser medium. When operated in a repetitively pulsed mode using liquid cooling, quartz flashlamps produce high average optical output power per unit area for many shots.

The major difficulty with flashlamp pulse sources is the low efficiency⁸⁰ of energy transfer to the active laser ion. Flashlamp output is not very directional, so it is difficult to design and construct lamp reflectors and pump cavities for reasonably efficient and uniform pumping. Flashlamp emission occurs in a band that is much broader than the absorption band of the laser medium; because of the resultant energy differences between pump and laser photons, large amounts of heat energy are deposited in the laser medium.

Fig. 9-63. Test prism for damage tests.

Flashlamp output can be spectrally narrowed by using reflective coatings on the envelopes. The coatings recycle energy in the unused portions of the output spectrum back to the lamp plasma, while passing radiation in the neodymium-pump bands. The use of coated lamps has great potential to improve solid-state laser efficiency and to reduce heating of laser media.

Operating Limites and General Characterizations

The maximum electrical energy that can be dissipated by a pulsed flashlamp without exploding has been determined empirically. The explosion energy E_x varies with the lamp dimensions and with pulse length⁸¹:

$$E_x = C l' d \tau^{1/2}, \quad (54)$$

where $C = 20\,000 \text{ J/cm}^2$, l' is the arc length, d is the lamp bore diameter, and $\tau = (LC)^{1/2}$. L and C are the single-mesh discharge circuit inductance and capacitance, respectively. We refer to the pulse length as $3\tau = 3(LC)^{1/2}$, which is approximately equal to the 10–10% width of the input power for a critically damped pulse. At fixed pulse width, the explosion fraction

$$f_x = \frac{E}{E_x} \quad (55)$$

is simply a measure of the input energy per unit area of flashlamp bore, where E is the pulse energy.

Lamp lifetime depends strongly on input energy. For lamps with fused-quartz envelopes, the number of shots N obtained during service life is estimated by⁸⁰

$$N = f_x^{-\beta}, \quad (56)$$

where β ranges from 8 to 10. For $f_x > 0.25$, lamp lifetime is usually limited by envelope failure. At these high lamp loadings, the inner surface of the envelope becomes crazed, and envelope material recondenses to a white powdery material in the tube.⁸² With continued operation, cracking becomes more pronounced until the envelope fails.

At lower pulse energies, lamp life is limited by electrode sputtering and the resultant darkening of the flashlamp envelope. In modern lamps, porous, sintered tungsten impregnated with low work-function materials reduces cathode sputtering, which is more severe than anode sputtering. Increased gas-fill pressures⁸³ and operation of the lamp in the simmer mode^{83,84} also reduce sputtering. With these techniques, lamp lifetimes of 10^8 shots have been obtained.^{84,85} Because improvements should be additive, the combined effects of changes in cathode materials, temperatures, gas-fill pressures, simmer operation, pulse shape, and pulse length should in the future extend lamp lifetimes to 10^9 shots.

Thermal stress induced in the envelope limits the average power that can be dissipated by a flashlamp.⁸⁰ Envelope heating occurs at the inside surface through electron-ion recombination and through absorption of UV radiation.⁸⁶ The ratio b' of thermally induced stress to tensile yield stress has been calculated for a cooled outer envelope surface as a function of thermal power per unit area deposited at the inside surface.⁸⁰ By limiting b' to an acceptable value (<1), and assuming that all input energy that does not appear as radiant output appears as heat energy in the flashlamp envelope, the maximum average electrical power per unit area of flashlamp bore P_A may be calculated from

$$P_A < P_{A_{\max}} = \frac{2R'_T b'}{\zeta(1 - \eta)}, \quad (57)$$

and

$$R'_T = \frac{(1 - \nu)KS_T}{\alpha e}, \quad (58)$$

where

$$\begin{aligned} R'_T &= \text{thermal shock parameter,} \\ \zeta &= \text{envelope thickness,} \end{aligned}$$

$$\begin{aligned} \eta &= \text{flashlamp radiant efficiency} \\ \nu &= \text{Poisson ratio,} \\ K &= \text{thermal conductivity,} \\ S_T &= \text{tensile yield stress,} \\ \alpha &= \text{coefficient of thermal expansion,} \\ e &= \text{Young's modulus.} \end{aligned}$$

The thermal shock parameters for fused quartz and for sapphire are about 14.5 and 100 W/cm, respectively.⁸⁰ For flashlamp envelopes 1 mm thick, and using $b' = 0.35$ and $\eta = 75\%$, we estimate the maximum average electrical input power per unit area to be 400 W/cm² for fused quartz and 2750 W/cm² for sapphire envelopes. Clearly, sapphire is an interesting material for high average powers.

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Lamp Parameters in HAP Lasers

The pump requirements of solid state HAP lasers are set by the laser operating characteristics, which are strongly dependent on the optical, thermal, and mechanical properties of the gain medium.⁸⁰ The demands made upon the flashlamp pumps are different for different gain media, as shown below. The comparisons are made for lasers operating at the maximum pulse energies allowed by parasitic losses, and at the maximum repetition rates allowed by the thermal power that can be dissipated by the laser slab without fracture.

The inversion loss rate through ASE is related to the product $\bar{\alpha}nl$, where $\bar{\alpha}$ is the plate-averaged gain coefficient, n is the index of refraction, and l is the slab length. When $\bar{\alpha}nl = 3$, ASE reduces amplifier stored energy by approximately 10%.⁸⁰ The constraint $\bar{\alpha}nl = 3$ can be rewritten in terms of the average stored-energy density

$$E_d = \frac{3S}{nl} \quad (59)$$

and the saturation parameter $S = hv/\sigma$, where hv is the laser photon energy and σ is the stimulated emission cross section.

We now calculate the explosion fraction, assuming that the slab is pumped to the ASE-limited storage energy. We use the

ratio a of slab single sided surface area to flashlamp bore area,

$$a = \frac{wl}{\pi dl'}, \quad (60)$$

where w is the width and t is the thickness of the slab. The storage efficiency of the laser η_s is defined as the ratio of the energy stored in the slab to the electrical energy of the flashlamps:

$$\eta_s = \frac{E_d w l t}{E}. \quad (61)$$

By combining Eqs. (54), (55), (59), (60) and (61), we find the explosion fraction,

$$f_x = \frac{3\pi S t}{n_s C \tau^{1/2} n l} a. \quad (62)$$

When designing lasers, the ratio a may be varied by adjusting the number or size of the flashlamps pumping a given slab to obtain operation at the desired f_x . One may choose low f_x to optimize lamp lifetime or a somewhat higher f_x to optimize neodymium-pumping efficiency. Alternatively, lamp size may be chosen to optimize uniformity of slab illumination. However, the constraint on a that the flashlamps be large enough to dissipate the required average powers must be satisfied.

Several processes contribute to heating flashlamp-pumped gain media, including relaxation from the pumped states to the upper laser level and relaxation from the lower laser level. There can be additional direct heating of the ion host material from absorption in the UV and IR unless care is taken to filter the flashlamp radiation. The specific thermal load parameter χ is defined as the ratio of the heat energy density released in the medium (both during pumping and after pulse extinction) to the stored energy density at the time of extraction.⁸⁷ Values of χ range from 0.5 to 2.^{80,87} The definition of χ is applied to relate the average thermal power density P_v to the energy storage density in the slab:

$$P_v = \chi E_d f, \quad (63)$$

where f is the laser repetition rate.

Thermal gradients develop in flow-cooled laser slabs as heat is deposited by

absorption of the flashlamp radiation. The thermal gradients cause surface tensile stresses, which must remain below the tensile yield stress to avoid rupture. The ratio b of the surface stress to the tensile yield stress has been calculated for face-cooled slabs that are uniformly heated⁹⁰:

$$b = \frac{P_v t^2}{12 R_T}. \quad (64)$$

The maximum P_v can be calculated by choosing b , preferably <0.3 for long slab lifetime. In practice, laser slabs are not pumped uniformly. More energy is stored and more heating occurs near the surfaces, and the surface stresses are lower than if the heating were uniform. We ignore this detail here, but note that our calculations for allowed average powers are conservative.

Using Eq. (63) to relate the thermal power density to the energy storage density, the repetition rate is

$$f = \frac{4n l R_T b}{S \chi t^2}. \quad (65)$$

The flashlamp average electrical power per unit bore area,

$$P_A = \frac{E f}{\pi d l'}, \quad (66)$$

can be expressed, using Eqs. (59), (60), (61), and (65), as

$$P_A = \frac{12 R_T b}{\chi \eta_s t} a. \quad (67)$$

High average power demands will be placed on the flashlamps when thin-cut slabs of material with a large thermal-shock parameter are pumped. The ratio a must be kept small enough that P_A does not exceed the limit imposed by induced surface stresses in Eq. (57). Therefore,

$$a < a_{\max} = \frac{1}{6} \frac{R'_T}{R_T} \frac{b'}{b} \frac{t}{\delta} \frac{\eta_s}{(1 - \eta)} \chi. \quad (68)$$

When a exceeds $a_{\max}$, the flashlamp envelopes are no longer large enough to handle the required average powers.

Given that the flashlamps must be of at least some minimum size to satisfy the average power requirements, and that the pulse energy is set by the ASE limit, the maximum explosion fraction f_x at which the lamps can be operated is reached when $a = a_{\max}$:

$$f_x < f_{x,\max} = \frac{3\pi St}{\eta_s C \tau^{1/2} n l} a_{\max} \quad (69)$$

We compare the flashlamp operating conditions for different slab materials by calculating f_x/a , P_A/a , $a_{\max}$, $f_{x,\max}$, and f for a particular laser slab geometry. We also find $a_{\min}$, the ratio $w/\pi d l'$ at which $f_x = 0.02$. Although this particular value of f_x was chosen somewhat arbitrarily, it reflects the need to run the flashlamps above some minimum f_x to obtain reasonable lamp radiant efficiencies. If the total flashlamp surface area exceeds the total slab surface area, it will be difficult to efficiently transfer flashlamp emission to the slab. Therefore, we also somewhat arbitrarily require that $a_{\min} \geq 0.5$; this gives the lower limit on the explosion fraction, $f_{x,\min}$, which may be greater than 0.02.

For the calculations, we assume the laser slab is 1 cm thick and 25 cm long, the lamps maximum sustainable electrical power per unit area is 400 W/cm², and the pulse width is $3\tau = 300 \mu\text{s}$. We also assume an amplifier storage efficiency of 10%, a specific thermal load parameter of 2, and that the slab is pumped to the ASE limit, with average power set by $b = 0.2$.

The results in Table 9-16 are not general, but they do give us a feel for the different

flashlamp operating conditions. The relatively low average powers for Nd:glass can be met over a broad range of flashlamp sizes and f_x at low repetition rates. The minimum flashlamp-bore area is not determined by average power requirements, as for the other materials, but rather by the pulse-energy requirements. We have arbitrarily chosen the maximum f_x to be 0.25, so that reasonable flashlamp lifetimes are obtained. To obtain longer lamp lifetimes, the maximum f_x should be made much smaller, less than 0.1. Also, there is considerable freedom to adjust lamp area and f_x for optimum lamp neodymium-pumping efficiency.

For GSGG and YAG, the pulse-energy requirements are low because of low saturation parameters. Because of high thermal-shock parameters, however, GSGG and YAG lasers are run at high average powers, and the flashlamp requirements are met over a narrow range of small f_x at high repetition rates. We expect that flashlamps pumping Nd:YAG and Nd:GSGG HAP lasers will be run close to the limit of input power per unit area, at least for lamps with fused-quartz envelopes. In this particular example for Nd:YAG, the calculated $a_{\min}$ exceeds $a_{\max}$. This tells us that the lamps cannot be run at $f_x > 0.02$ without some compromise in average power (e.g., reducing the repetition rate). For our Nd:GSGG example, pulse energy requirements can be met over the range $0.02 < f_x < 0.047$.

Alexandrite lasers are pumped at very high average powers at moderately high f_x and repetition rates. The average-power requirements are met by making the

Table 9-16. Material parameters and calculation results that demonstrate the different operating conditions of flashlamps pumping different gain media.

Material	n	R_T (W/cm)	S (J/cm ²)	f_x/a	P_A/a (W/cm ²)	$a_{\max}$	$a_{\min}$	$f_{x,\max}$	$f_{x,\min}$	f (Hz)
Nd:YAG	1.81 ^a	~7.90 ^a	0.21 ^b	0.0022	95	4.3	9.1	0.010	0.02	680
Nd:GSGG	1.92 ^a	~7.00 ^a	1.0 ^c	0.0098	84	4.8	2.0	0.047	0.02	130
Nd:glass (LHG-5)	1.54 ^d	1.0 ^e	4.6 ^d	0.056	12	4.5 ^f	0.5	0.25 ^f	0.028	3.3
Alexandrite	1.73 ^a	~35.7 ^a	20 ^g	0.218	428	0.95	0.5	0.21	0.11	30

^a Ref. 80.

^b Ref. 88.

^c Ref. 89.

^d Ref. 90.

^e Ref. 91.

^f The average power requirements imply $a_{\max} = 34$. For Nd:glass, $a_{\max}$ is limited by pulse energy considerations, and we have arbitrarily selected $f_{x,\max} = 0.25$.

^g Ref. 92.

flashlamps large. The working range of f_x extends from 0.11 to 0.22; it is more limited than the range for Nd:glass because of the higher average power requirements.

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Flashlamp Coatings and Laser Efficiency

The efficiency of flashlamp-pumped solid-state lasers is limited by the mismatch between the broad emission spectra of flashlamps and the more narrow pump bands of active ions. For example, flashlamp spectra lie within the transmission band of clear fused quartz from 0.2 to $4.5\text{ }\mu\text{m}$ while the major neodymium-pump bands lie between 0.4 and $1.0\text{ }\mu\text{m}$. The neodymium absorption band at $\sim 0.36\text{ }\mu\text{m}$ usually does not contribute to pumping because of cerium-doping of the laser glass or lamp envelopes to reduce solarization, and because pump-cavity reflectivity is low in the UV. The data presented in "Diameter Scaling" in Sec. 7 show that 35 to 55% of the xenon-filled flashlamp radiation falls outside the major neodymium-pump bands. A representative flashlamp spectrum is superimposed with neodymium-absorption features in Fig. 9-64.⁹³

Reflective coatings placed on the envelopes of flashlamps increase output in the laser pump-band range by reflecting unused portions of the lamp emission spectra back to the lamp plasma to be absorbed. Flashlamp plasmas are near local thermodynamic equilibrium, where energy is exchanged rapidly between particles. Energy added to the plasma through optical absorption should therefore have the same effect as an equivalent amount of added electrical energy. Flashlamps are very efficient converters of electrical to optical energy, and they should also convert optical input into optical output with the same high efficiency. Thus, the reflective coating technique should significantly increase flashlamp emission over a desired wavelength range at the expense of emission at other wavelengths.

To demonstrate the potential of coated lamps for improving laser efficiency, we calculated the fractional enhancement in the 0.4- to $1.0\text{-}\mu\text{m}$ neodymium-pump band, assuming that either the UV, the IR, or both have been recycled. The fractions of radiant output in the bands $\lambda < 0.4\text{ }\mu\text{m}$, $0.4 < \lambda < 1.0\text{ }\mu\text{m}$, and $\lambda > 1.0\text{ }\mu\text{m}$ are designated f_{UV} , f_{pump} , and f_{IR} , respectively. These fractions, and the instantaneous radiant efficiency η , have been characterized as functions of P_A using the data

Fig. 9-64. Representative emission spectrum of a xenon-filled flashlamp and the absorption spectrum of Nd:glass.

presented in "Diameter Scaling" in Sec. 7. Because η and the flashlamp output spectra of different diameter lamps have been shown to have a common functional dependence on P_A , the results apply to different bore diameters as well as to different pulse lengths. Optical recycling increases the total power per unit area P_T delivered to the lamp plasma. Assuming that electrical and optical inputs are equivalent,⁹⁴

$$P_T = P_A + \eta(P_T) \times [\delta_{UV} f_{UV}(P_T) + \delta_{IR} f_{IR}(P_T)] P_T, \quad (70)$$

where δ_{UV} and δ_{IR} are either 1 or 0, depending on whether UV or IR recycling is included or not. The efficiency η and f_{UV} and f_{IR} are now functions of the total input power to the plasma.

By selecting P_T , it is straightforward to calculate P_A . Then, the fractional enhancement factor F is simply the power in the pump band obtained with recycling divided by the pump band power obtained without recycling:

$$F = \frac{f_{\text{pump}}(P_T) \eta(P_T) P_T}{f_{\text{pump}}(P_A) \eta(P_A) P_A}, \quad (71)$$

The calculations of fractional enhancement (Fig. 9-65) show that neodymium-pump band emission may be increased by as much as 45%. Enhancement from combined UV and IR recycling is nearly

independent of P_A . Recycling the IR gives decreasing enhancement as P_A increases because the fraction of emission in the IR becomes smaller. The opposite is true for recycling UV, and only at very low P_A does IR recycling produce greater enhancement than UV recycling. The total enhancement from combined recycling is greater than the product of the individual UV and IR enhancements by 1 to 4% because leakage of the enhanced lamp output out of the missing reflected band does not occur with dual UV and IR reflectivity.

The results in Fig. 9-65 are idealized pump-band enhancement factors; measured enhancements will be lower for several reasons. Actual lamp coatings do not have 100% reflectivity in the UV and IR or 100% transmission in the neodymium-pump bands at all angles of incidence. For example, limits on the thickness of coatings restrict IR-reflectivity to ~ 1.0 to $1.6 \mu\text{m}$. Also, reflection spikes in the pump bands and transmission holes in the UV and IR appear with the detuning of coatings at large angles of incidence.

Lamp plasmas are not strong absorbers in the UV, especially at low lamp loadings. Therefore, UV emission undergoes several reflections on average before absorption. The transmission losses over several reflections reduce the amount of UV absorbed by the plasma. Another process that reduces enhancements of real coatings is the useless absorption of reflected UV and IR at the lamp electrodes.

The enhancement factor depends strongly on the lamp radiant efficiency because recycling is repetitive. To maximize the benefit of coated lamps, it is important to obtain the highest possible lamp radiant efficiency by using simmer currents or by applying preionization pulses. It is also important to use solarization-resistant envelope material to keep absorption losses low.

Optical recycling has been demonstrated⁹³ using flashlamps coated with multilayer dielectric films. Coatings reflective in the UV have enhanced neodymium-pumping efficiencies by as much as 12% at $f_x = 0.27$ at pulse lengths of $\sim 600 \mu\text{s}$. Less enhancement was obtained at reduced pulse energies. In other tests, the electrical input power of a coated lamp was reduced relative to that of an

Fig. 9-65. Calculated fractional enhancement in Nd-pumping efficiency vs flashlamp electrical power per unit area. Optical and electrical input to the lamp plasma are assumed to be equivalent.

uncoated lamp, so the emission spectra of the two lamps were virtually identical over the pass-band of the lamp coating. In experiments of this type, the flashlamp plasma was found to use electrical and optical input energies equivalently, within experimental error. Coatings survived the high radiant fluences present at the outside surfaces of flashlamp envelopes.

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Coating Fabrication and Durability

To effectively increase laser efficiencies by using coatings on flashlamp envelopes, the coatings must meet demanding reflectance and transmission requirements over broad spectral regions. Flashlamp emission within the neodymium-pump band is enhanced by recycling emissions that fall in other spectral regions back to the lamp plasma. Coatings must therefore be highly reflective in the UV and IR, but highly transmissive within the neodymium-pump bands. Multilayer dielectric films best meet the required spectral characteristics of flashlamp coatings. We consider here the properties and fabrication of dielectric films on lamp envelopes and the damage that can occur to lamp coatings within the HAP laser pump-cavity environment.

Dielectric Coatings. Optical coatings are constructed by stacking thin layers of dielectric materials with different refractive indices; the layer thickness is chosen to give constructive interference of the reflected waves. For example, high reflectivity coatings can be made of alternating quarter-wave layers of high- and low-index materials. Coatings that are reflective in the UV are thinner than coatings reflective in the IR.

Dielectric materials are heated by electron bombardment and deposited onto substrates in a vacuum chamber. Coating thickness depends on the source and coating geometries, the density of the evaporant stream, condensation coefficients, and exposure times. Coating uniformity depends on the angular distribution of evaporant from the source, the distance from the source to the substrate, the angles of incidence, and collisions of

evaporant molecules with each other or with a background gas. It is particularly desirable to have uniform coatings on lamps that pump HAP lasers because uniform illumination of laser slabs avoids optical distortions from thermal gradients.

Obviously, it is more difficult to coat a cylindrical surface, but uniform coatings have been obtained by using precision-ground lamp surfaces, rotating the substrate, and masking to control the angle of incidence. The entire surface of a cylindrical substrate must be exposed to the source during the deposition process. This is done at considerable mechanical complexity and cost by rotating the lamp behind a slit that runs the length of the lamp. For uniform deposition, the slit must be properly aligned with the lamp and the coated cylindrical surface must be straight and concentric with the axis of rotation. Dielectric coatings deposited at large angles of incidence are inhomogeneous and have inferior mechanical properties, so evaporant material reaching the lamp surface at angles much greater than 30° to the normal must be masked.

Dielectric coatings do not exactly match ideal specifications. For example, the region of IR reflectivity is limited to ~ 1.0 to $1.6\ \mu\text{m}$ because coating layers (whose thickness scales with the reflected wavelengths) that are too thick are subjected to internal stresses that cause crazing, cracking, clouding, or peeling.⁹⁵ Third-harmonic reflections also limit the IR-reflective range, and coatings with high reflectivity in the 1.0 - to $1.6\text{-}\mu\text{m}$ range have spikes of high reflectivity at wavelengths of $\sim 0.5\ \mu\text{m}$ and below. If the transmission range of the coating is to be extended downward to $0.4\ \mu\text{m}$ to use more of the neodymium absorption bands, the IR-reflective region must be abbreviated.

Spectral features shift to shorter wavelengths as incidence angle increases, so the IR-reflective region at normal incidence must be shifted out to longer wavelengths than would otherwise be optimum. Also, as mentioned above, the detuning of the coating at large angles of incidence reduces transmission in the pump bands and decreases reflectivity in the UV and IR.

Reflectivities may fail to meet design specifications, especially when thick IR reflectors are used. Thick coating layers tend

to be inhomogeneous and the dispersion reduces neodymium-pump band transmission.

Damage in the Pump-Cavity Environment. Flashlamp coatings in HAP solid-state lasers are subjected to many physical and chemical processes acting together or separately to cause damage. Coatings are exposed to broadband flashlamp radiation at high fluences for many shots. Absorption of radiation, which is strongest in the UV, can cause catastrophic coating failure. At high average powers, heating of the flashlamp envelopes induces tensile stress that can cause cracking or crazing in the coatings. Impurities in the water can attack the coatings through dark or photochemical reactions. The deposition of particles or minerals over long periods of time can also damage coatings. Some dielectric materials, including silica, are soluble in water, although the time required for coatings to dissolve in flowing, unsaturated water is unknown.

Absorption of flashlamp radiation by the dielectric lamp coatings can cause local heating of the material to the point of stress fracture, melting, or evaporation. Absorption occurs as a result of both intrinsic material properties and small particulate impurities. To avoid damage to flashlamp coatings from intrinsic absorption, dielectric materials should be chosen with UV-absorption edges at or below the 0.2- μm UV-absorption edge of natural clear-fused-quartz envelopes. As a counter-example, coatings of titania and silica were severely damaged on single flashlamp shots at low energies.⁹³ The UV-absorption edge of thin films of titania occurs at 0.330 μm , (Ref. 96) well above 0.2 μm . Lamp coatings of hafnia and silica with low UV-absorption edges of 0.225 and <0.200 μm , (Ref. 97) respectively, have survived at least several shots at substantial lamp loading (18 J/cm² of optical output) without noticeable damage.⁹³ Nevertheless, the effects of low UV absorption by the coatings over many shots is unknown.

The choice of lamp envelope materials can moderate coating exposure to UV. Titanium-doped quartz absorbs below 0.220 μm , (Ref. 83) so the absorption edge of titanium-doped envelopes is well matched to the absorption edge of hafnia (and to

those of scandia and yttria, two other high-index materials with low UV-absorption edges).⁹⁷ Thus, intrinsic UV absorption is minimized without loss of recyclable UV radiation. Cerium-doped envelopes, which absorb below 0.380 μm , (Ref. 83) would protect coatings from UV, but would leave little UV for recycling.

The envelope material should be solarization resistant. UV radiation causes solarization in fused quartz that contains ionic impurities (especially Al³⁺, Li⁺, and Ge⁴⁺),⁹⁸⁻¹⁰⁰ The color centers formed from solarization reduce laser efficiency because they absorb light that would otherwise be used to pump neodymium ions or be recycled in the lamp plasma. Solarization reduces average transmission by 5 to 10% in most natural quartz envelopes¹⁰¹ after a period of lamp operation; the extent of quartz-envelope solarization varies with the density of impurities and with the manufacturer. Doping quartz with cerium^{101,102} inhibits solarization, but reduces recycled UV (see above).

Another envelope material, suprasil, is made from gaseous sources of silicon and oxygen, and contains extremely low levels of ionic impurities; its high purity would nearly eliminate solarization. However, the UV-absorption edge of suprasil is at ~0.175 μm , (Ref. 83) which increases the UV exposure of the lamp coatings below the UV-absorption edge of the coating materials. Because we do not know the lamp loadings at which emission at short UV wavelengths causes either short- or long-term damage to coatings, we must experimentally assess the usefulness of suprasil envelopes for coated lamps.

Absorption at impurity sites can also cause coating damage. Avalanche ionization, which is responsible for much damage in high-power laser optics, will not damage flashlamp coatings because flashlamp radiation is incoherent and peak power densities are relatively low. Avalanche ionization occurs first at atomic or particulate impurity sites or, in pure materials, along cracks or pits.¹⁰³ Flashlamp coatings have a higher tolerance for impurities and substrate defects than do high-power laser optics, so techniques already used in the manufacture of high-power laser optics will keep impurities sufficiently low in flashlamp coatings.

Dielectric films sustain damage when tensile yield stress is exceeded. Fortunately, the tensile yield stress of thin films can be many times higher than that of the bulk material.⁹⁵

With low coating absorption, laser operation induces only low thermal stresses between flashlamp coatings and the outer envelope surfaces. However, the inside surfaces of flashlamp envelopes will be heated at high average powers, and external tensile stress will develop at the outer envelope surfaces and in the lamp coatings. To avoid coating damage, we must rely on experiments to tell us whether the average thermal powers dissipated in coated lamps must be kept lower than those for uncoated lamps. The calculations would require Young's modulus and Poisson's ratio for the coating materials, but these values are generally unknown.

HAP flashlamps are usually water-cooled, so coatings must work well submerged in water. Submersion could change coating performance or cause coating degradation. However, we do not expect large changes in the spectral characteristics of submerged lamp coatings because most become saturated with water from the atmosphere shortly after manufacture,¹⁰⁴ and it is likely they are designed using the refractive indices measured for water-saturated layers.

If cooling water is inadequately filtered, suspended particles can be deposited on coatings and heated by absorption of lamp radiation, causing damage. If cooling water is inadequately pure, species in solution may chemically attack the coatings (e.g., corrosion of fused silica by hydrofluoric acid, decomposition of impurities to reactive species by flashlamp radiation). Damage from deposition of soluble materials is unlikely when the lamps are running because solubilities usually increase as water is heated.

A potentially more serious water-related problem is that flashlamp coatings may eventually delaminate because their dielectric constituents are soluble. For example, at room temperature in neutral solution, measured solubilities of amorphous silica range from 0.01 to 0.015%.^{105,106} Solubility increases rapidly with temperature to ~0.04% at 100°C, and also increases with

pH. Soluble coating materials exposed to unsaturated water will eventually dissolve. The rate of dissolution is not known, although rates of solution are dependent on the specific surface area; the saturation of silica solutions requires years for bulk silica, but only days for powdered silica.¹⁰⁵ Even so, serious damage to coatings could take place long before a significant fraction of the coating mass had gone into solution. Solubility damage needs to be examined experimentally; if necessary, protective jackets could be developed to seal coatings from water.

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Pulsed Power

Introduction

Our electrical pulsed-power system is configured to meet known and forecast requirements for solid-state HAP lasers. We currently support the three major test-bed laboratories described in preceding articles: the zig-zag, gas-cooled slab, and flashlamp test units. Table 9-17 summarizes the general region of interest for pulsed power, and gives the operating values demonstrated thus far for these parameters. The electrical and electronic components of the power system are traditional with some innovative augmentations that are described in following articles. Electrical power is taken from the 13.8-kV utility grid and conditioned to provide independent pulsed power to each of the three test beds. The block diagram in Fig. 9-66 shows the basic sections of the electrical system.

During the design, implementation, and operation of the electrical systems, we

Table 9-17. HAP pulsed-power region of interest.

Parameter	Desired range	Demonstrated range
Voltage (peak) (kV)	1-7	1-7
Current per load (kA)	0.5-5	0.5-5
Pulse width (μ s)	10-400	200-300
Pulse rate (kHz)	0-10	0-0.125
Energy per pulse (kJ)	2-4	2-3
No. of flashlamp loads	1-30	1-4
Peak power (total) (MW)	0-1	0-0.1

Pulsed Power

paid special attention to the criteria for safety, reliability, and flexibility.

Safety. All high-voltage areas are either physically locked or electrically interlocked to prevent inadvertent access by laboratory personnel. Additionally, all operator controls and instrumentation are isolated from interfaced pulsed-power circuitry with fiber optics.

Reliability. The pulsed-power system is designed to support laser operation rates ranging from a single shot to over 1000 shots per second. Electrical components (e.g., resistors, capacitors, inductors, and diodes) are sized to meet worst-case electrical and environmental conditions. Single-fault conditions such as laser flashlamp failures must be detectable and must not result in failure of the electrical system.

Flexibility. The system is designed to accommodate a wide range of electrical

parameters such as voltage, current, energy per pulse, pulse shape, and repetition rate. Parameter changes are accomplished by motor-driven variable-voltage transformers, segmented capacitors and inductors, and timing control of silicon-controlled rectifiers (SCRs). Test data are converted to digital form and sent to computers for processing and display.

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Test-Bed Power Conditioning

The basic pulsed-power circuit used with the test-bed flashlamp circuit is shown in Fig. 9-67. When testing is initiated, the simmer supply ionizes the

Fig. 9-66. Block diagram of HAP laser pulsed-power system.

Fig. 9-67. Basic resonant charge circuit.

flashlamp gas with a single high-voltage low-energy pulse. Then, a continuous simmer current supplied to the lamp maintains an ionized condition, and a stream of pump pulses is delivered to the lamp by sequentially charging and discharging the energy storage capacitor.

We chose the resonant charge method to supply pulsed power: a dc power supply charges a pulse-forming capacitor, and a closing switch discharges the capacitor through a shaping network and the flashlamp load. This type of pulse generation is very efficient, especially when augmented with the energy recirculating techniques described below. Note that the resonant charge SCR isolates the load from the power supply. Parallel circuits thus can share a single power supply without the need for additional isolation.

The electrical load for each of the test beds is one or more flashlamps. Shaped pulses of electrical power are delivered to loads to meet specific test objectives. At single shot or very low repetition rates, the explosion energy of the flashlamp is the limiting factor. At higher pulse rates, power loading on the flashlamp wall becomes the limiting factor. In practice, we limit the pulse power to 25% of explosion energy or to a maximum of 200 W/cm² for conventional flashlamps. Novel flashlamp designs are expected to extend this wall-loading limit by an order of magnitude.

Flashlamp explosion energy E_{exp} and wall loading are given by

$$E_{\text{exp}} = 2 \times 10^4 LD\tau^{0.5}, \quad (72)$$

$$\text{Wall loading} = \frac{I/\text{pulse} \times \text{pulses/s}}{\pi DL}, \quad (73)$$

where

L = flashlamp length (cm),

D = flashlamp diameter (cm),

$\tau = (LC)^{1/2}$.

The HAP range of interest for these parameters is

- Flashlamp length between 25 and 50 cm.
 - Flashlamp diameter between 1 and 2 cm.
 - Pulse width 3τ between 50 and 500 μs .
- These parameters translate to the following electrical parameters:
- Power-supply dc voltage, 3 to 4 kV.
 - Resonant-charge peak current, 100 to 200 A.
 - Resonant-charge inductor, 80 to 640 mH.
 - PFN capacitor, 100 to 300 μF .
 - PFN inductor, 30 to 525 μH .
 - Lamp peak current, 1 to 10 kA.

For a typical flashlamp 1 cm in diameter and 25 cm long, Fig. 9-68 shows the acceptable operating realm for 200- μs pump pulses.

Fig. 9-68. Operating region for conventional 25-cm-long, 1-cm-diam flashlamp, $\tau = 200 \mu\text{s}$.

Fig. 9-69. Schematic of 600-kV-A HAP power supplies.

Figure 9-69 is a schematic of the HAP 300- and 600-kV·A dc power supplies. Each power supply has a motor-driven variable-voltage transformer that is controllable from the HAP control room.

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Pulse Shaping

Shaping the flashlamp current pulse offers several advantages. Shaped pulses can reduce the thermal loading of the laser medium and improve the wall plug efficiency of the laser system. We are investigating two types of pulse shaping. The first is a clamped LC discharge and the second is a square current pulse drive.

The electrical drive for our flashlamps is an LC discharge loop. The curve of stored energy in the laser medium peaks before the flashlamp turns off in a normal

LC discharge. Turning off the flashlamp at the peak of the stored energy curve shapes the pulse. This reduces thermal loading for a given stored energy (Fig. 9-70); thus, the total energy required to run the system is reduced.

At high repetition rates, our flashlamps are run at a low percentage of flashlamp explosion energy, which reduces the pumping efficiency. A shorter, higher peak-intensity LC discharge driver could be used to counteract this. However, we expect a greater improvement by pulse shaping with a square current pulse, where the lamp is operated at its highest pumping efficiency during the length of the pulse.

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Diagnostics

Instrumentation for the pulsed-power system provides test personnel with the real-time status of electrical parameters, and also records test-bed parameters

Fig. 9-70. Flashlamp current, voltage, and power with the LC discharge loop drive.

Fig. 9-71. Recorded flashlamp voltage and current and the calculated electrical power.

for diagnostic evaluation of system performance. The following displays are continually updated:

- Analog meter display of power-supply voltages and currents.
- LED display of interlock status.
- Oscilloscope display of flashlamp current pulses.
- Visual flowmeters of flashlamp coolant.
- Computer readouts of device temperatures.

Each test bed has a 40-channel digitizer and multiplexer connected with fiber optics to a small computer in the control room. Analog measurements are sampled at up to 10 times per second and recorded for post-test data reduction. These measurements are typically associated with coolant temperatures and flow rates, but can readily be connected to other sensor packages.

Flashlamps represent a nonlinear load that varies with pulse energy, pulse rate,

and simmer current. The actual energy per pulse must be within prescribed limits to validate predicted test-bed performance. Thus, each flashlamp circuit is instrumented with a magnetic current probe, and a selected circuit is connected to a digitizing oscilloscope (Tektronix 7020). High-voltage probes from the digitizing scope are also attached to the selected flashlamp. For a single pulse, or for the ensemble average of a sequence of pulses, the scope records current and voltage, and calculates the integrated electrical energy delivered to the flashlamp (Fig. 9-71). Selected digitized current and voltage waveforms are transferred to the control room computer for additional analysis.

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data management; IDFLOW thermal
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fluids code; LASNEX hydrodynamics
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GETBOOK, UPDATE, HERMES,
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