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Lawrence Livermore National Laboratory

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Preface

The Laser Program Annual Report is part of the continuing series of reports documenting the progress of the unclassified Laser Fusion Program at the Lawrence Livermore National Laboratory (LLNL). As in previous years, the report is organized programmatically. The first section is an overview of the basic goals and directions of the LLNL Inertial Confinement Fusion (ICF) Program, and highlights the year's important accomplishments. Sections 2 through 7 provide the detailed information on the various program elements: Laser Systems and Operations, Target Design, Target Fabrication, Laser Experiments and Advanced Diagnostics, Advanced Laser Development, and Applications of Inertial Confinement Fusion.

The purpose of this report is to present our work in a brief form, but with sufficient depth to provide an overview of the analytical and experimental aspects of the LLNL ICF Program. More detailed expositions of the research can be found either in the many publications and reports authored by the ICF Program staff members or by contacting the author(s) listed at the end of each article.

Erik Storm

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In inertial confinement fusion, powerful laser beams compress and heat a small amount of DT fuel to high densities and temperatures, initiating a fusion burn. To accomplish this, the LLNL Laser Program builds high-power high-energy neodymium glass lasers (currently, the recently-completed Nova laser), develops the scientific concepts and computer calculations pertaining to fusion burn, fabricates fuel-filled targets designed for efficient fuel burn, conducts complex experiments and precisely measures the results, researches all aspects of components for future high-power lasers, and develops designs for future fusion power plants. Success in these areas, already demonstrated, will have both near- and long-term military and civilian benefits for the country. Recently, lasing action in the the x-ray region was obtained for the first time by exploding a thin selenium foil with the Novette laser's green light. The inset photo at left shows an exploding-foil x-ray target. In the large photo at left, a target is irradiated inside the Nova target chamber. In the photo below, part of Nova's 137-meter-long laser chains—the 20-cm amplifiers and the spatial filter—is portrayed.

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 Laser Fusion Program are to understand and develop the science and technology of inertial confinement fusion (ICF), and to utilize ICF in short- and long-term military applications, and, in the long-term, as a candidate for central-station civilian power generation. Some military applications, in the form of weapons physics experiments, are already being realized with the laser facilities constructed for the Laser Fusion Program. However, full realization of the military applications and the civilian potential of ICF will only be reached with the attainment of well-diagnosed, moderate- to high-gain (>0.1 ton TNT equivalent yield) thermonuclear performance. This accomplishment will experimentally prove the scientific feasibility of the ICF concept; in addition, it will generate large fluxes of x rays and neutrons, permitting a broad range of new military applications studies.

To reach our goal of demonstrating high-gain performance in the laboratory, we are conducting a technical program focused on two major complementary thrusts: experiments and advanced laser studies.

In the first thrust, we design and conduct experiments using our high-power Nd:glass laser systems whereby we develop our understanding of laser-target interaction phenomena. Our goal in these experiments is to demonstrate that the conditions required to drive a high-gain capsule can be attained with a driver of sufficient energy. In 1984, using the Novette laser system, we completed experiments showing the very favorable scaling of laser-plasma interactions with short-wavelength laser light. Our Novette experiments have unequivocally shown that short laser wavelength, i.e., less than $1\ \mu\text{m}$, is required to provide the drive necessary for efficient compression, ignition, and burn of DT fusion fuel. In other experiments with Novette, we made the first unambiguous observation of amplified spontaneous emission in the soft x-ray regime. We also conducted military applications and weapons physics experiments, which we discuss in detail in the classified volume of our Laser Program Annual Report.

In the second thrust, advanced laser studies, we develop and test the concepts, components, and materials for present and future laser systems. Over the years, this has meant providing the technology base and scientific advances necessary to construct and operate a succession of six ever-more-powerful laser systems (see Fig. 1-1). The latest of these, Nova, a 100-TW/100-kJ-class laser system, was completed in 1984. The Nd:glass laser continues to be the most effective and versatile tool for ICF and weapons physics because of its scalability in energy, the ability to efficiently convert its $1\text{-}\mu\text{m}$ output to shorter wavelengths, its ability to provide flexible, controlled pulse shaping, and its capability to adapt to a variety of irradiation and focusing geometries. For these reasons, many of our advanced laser studies are in areas appropriate to solid state laser technologies.

We must begin to develop a plan for an ICF facility for the next decade. We are confident that Nova experiments will strongly support our present projections for achieving high gain with a 5- to 10-MJ short-wavelength laser. We also predict that a high (>100) target gain can ultimately be attained with

lower (~ 2 to 5 MJ) input energies, the result of progress in understanding the physics and technology of designing and fabricating targets. We have thus continued analyses to examine the issues associated with the design and construction of a facility for high-gain target investigation, and to identify the research directions required to define such a facility.

In support of military applications, we define three categories of experiments. One category is to investigate and enhance our understanding of the physics of nuclear weapons, and to develop nuclear test diagnostics in support of the Weapons Program. We have made significant contributions in these areas with our current ICF facilities. The densities and temperatures attainable on a small scale with the high-power lasers developed for ICF are comparable to those otherwise achieved only in underground nuclear tests. Consequently, experiments on topics of physics and diagnostics of interest and applicability

Fig. 1-1. Solid state Nd:glass laser design has progressed rapidly since its invention. Construction of multimegajoule lasers is, in principle, a straightforward extension of existing technologies.

to the Weapons Program have already been conducted and will continue on a somewhat larger scale on Nova. When we achieve modest target gain with an appropriate future driver, additional military applications experiments will begin. This will add the second category of military applications: to produce high enough x-ray, neutron, and gamma-ray fluxes to simulate the effects of nuclear explosions on military targets, i.e., ballistic missiles, reentry vehicles, and satellites and their components. The third category of experiments is to provide a laboratory capability to help develop an understanding of the physics of nuclear-driven directed-energy weapons.

The longer term objective of the Program, to develop and utilize ICF technology for civilian energy production, has an extremely high payoff but is also one that entails considerably higher technical challenge and economic risk.

Progress and Accomplishments

Experiments

Nova Completion. The construction phase of the \$176-million Nova laser project was completed on schedule and on budget in December, 1984. Nova, depicted in Figs. 1-2, 1-3, and 1-4, is a three-wavelength system, capable of irradiating targets with 3-ns pulses that deliver 100 to 120 kJ at a wavelength of 1.05 μm , or 50 to 70 kJ at 0.53 or 0.35 μm . Over the next 2 to 3 years, Nova will be used to extend laser-plasma studies with plasmas 10 times larger than those produced with Novette. We expect that

these experiments will confirm our theoretical and experimental understanding of laser-plasma coupling and implosion physics for the types of capsules that are required for high-gain performance. We also expect to extend x-ray laser performance to shorter wavelengths, higher power levels, and greater coherence.

Laser-Target Interaction. Major progress occurred in understanding and quantifying laser processes and laser-target interactions. This progress includes not only obtaining very positive experimental data, but also improving our ability to model and scale these phenomena to conditions appropriate for high-gain targets. In an ICF target, high-gain conditions occur when 2 to 10 mg of DT fuel are compressed to densities exceeding 200 g/cm³ at 100 million thermonuclear burn is initiated in a high-temperature

Fig. 1-2 The Nova 10-beam ICF facility.

Fig. 1-3. Nova spaceframe showing the 31-cm amplifiers on the near right, the larger 46-cm and 74-cm hardware in the distance on the right, and, on the left, the small 6-, 10-, and 20-cm amplifiers at the beginning of the other 5 chains.

Fig. 1-4. Nova target chamber, showing 5 of the 10 spools containing the frequency-conversion crystals and the final focusing optics.

Fig. 1-5. Novette target chamber in the foreground, with the two laser chairs and pre-amplifier sections in the background.

(>3 keV) hot spot. To assemble the fuel in such a state, a precisely time-programmed, 10-ns duration pressure pulse, exceeding ≈ 9 Mbar at its peak, must be applied. Very little preheat of the fuel can be tolerated if efficient implosions are to occur. In order to achieve the energy densities that correspond to these high pressures, the laser intensity must be $>10^{14}$ W/cm². Experiments on the Argus and Shiva lasers (designed to explore the laser-plasma interaction and high-compression conditions required for high-gain performance) showed that hot electrons generated with 1.06- μ m laser light caused deleterious preheat. In some experiments, as much as 50% of the laser light appeared as very energetic (10 to 100 keV) hot electrons, which significantly compromised the quality and usefulness of the experiments. Theoretical work, confirmed by some small-scale experiments on Argus and at other facilities around the world, indicated that converting the laser light from its fundamental (or first-harmonic) frequency (1.06 μ m) to its second (0.53 μ m), third (0.35 μ m), or fourth (0.26 μ m) harmonic would greatly reduce the production of hot electrons. We were able to confirm this theory with Novette experiments.

Novette (see Fig. 1-5) provided as much frequency-converted (0.53 μ m) laser energy to targets as did the entire 20-beam Shiva laser at the fundamental frequency of 1.06 μ m. Utilizing up to 8.5 kJ (in a 1-ns pulse) of 0.53- μ m light from Novette, we studied short-wavelength, laser-target interaction phenomena and verified that positive scaling occurred under these short-wavelength, high-energy, large-plasma conditions (see Fig. 1-6). With targets irradiated at 0.53 and 0.26 μ m, we saw reductions of two to three orders of magnitude in the hot-electron production as compared with similar targets irradiated with the same energy at 1.06 μ m.

These experiments, as well as experiments on disks coated with high- or low-Z materials, have identified that stimulated Raman scattering (SRS) is the primary mechanism responsible for generating the hot electrons. A subset of these experiments also demonstrated that by varying the target Z, and the laser intensity and wavelength, we could keep the plasma collisional and keep the laser-plasma interaction conditions below the SRS threshold.

These experiments and their implications for *Nova* and future high-gain target experiments are discussed in more detail in Secs. 3 and 5 of this book.

Using Novette's two 0.53- μm beams, we executed a series of high-compression experiments. Because of the significantly lower hot-electron-preheat environment and, hence, the much higher quality of the implosions driven with 0.53- μm -light, targets in these experiments achieved final compressed-fuel densities about a factor of 2 higher than similar targets driven with the same energy on *Shiva* with 1.06- μm light (see Fig. 1-7).

The Novette plasma-coupling data and these compression results lead us to expect that by using temporally shaped pulses of 30- to 50-kJ of 0.35- μm light from *Nova*, we can compress DT fuel to the 200 g/cm³ density ultimately required for high-gain target performance. One of the more significant implications of the Novette experiments is that 0.5 μm may be a short enough wavelength for high-gain target performance. We plan to resolve these issues with *Nova*.

The performance of Novette in 1983 and 1984 certified the *Nova* laser design and showed that *Nova* will easily meet its baseline specifications.

X-Ray Laser Demonstration. Eliminating the hot-electron problem also allowed us to successfully complete a major x-ray laser campaign. This series of Novette experiments was in direct support of and in close collaboration with the laboratory's X-Ray

Laser Program. These experiments produced the first conclusive demonstration of laser gain in the soft x-ray region (see Fig. 1-8) and also showed clearly the utility and effectiveness of our ICF facilities for

Fig. 1-6. Summary of laser-plasma interaction results showing that laser-plasma coupling is very favorable at shorter wavelengths.

Fig. 1-7. Novette, operating at 0.53 μm , produced significantly higher compression than did *Shiva* at 1.06 μm . These data imply that final compressed fuel densities are $>20 \text{ g/cm}^3$.

Fig. 1-8. Novette experiments demonstrated the first laboratory laser gain in the soft x-ray region. (a) The two Novette beams focused to lines on the target foil. (b) X-ray laser emission at 206 and 209 Å. (c) Exponential growth of x-ray laser intensity with increasing selenium foil target length.

Progress and Accomplishments

conducting high-quality, well-diagnosed experiments important to the Weapons Physics Program.

Target Fabrication. The Target Fabrication Program in 1984 supplied the variety of targets required for Novette experiments, and, in addition, developed the target technologies required for high-density experiments on Nova. As part of that effort, we advanced our target-coating technologies by substantially improving our capabilities to produce smooth, uniform, low-Z ablaters. We can now produce high-quality plastic shells and have begun to develop techniques for producing thick (~ 10 to $30\ \mu\text{m}$) cryogenic fuel layers for Nova targets. As part of our support of the x-ray laser experiments on Novette, we developed the precision technologies required for fabricating and aligning the x-ray laser targets. The target fabrication technologies are described in Sec. 4.

Advanced Laser Studies

Introduction. Our laser technology development activities play a very important role in meeting the ongoing needs of the Program and its future goals. Our developments in optical materials, optical manufacturing and finishing, lens and mirror coatings, laser and optical components, frequency conversion, and analytical and numerical modeling are all significant (see Fig. 1-9). These developments will not only play an important role in supporting the

operation, maintenance, and improvement of existing laser systems, but will also form the technological basis for the development of concepts and architectures for advanced laser systems. These developments will contribute to more effective laser systems by improved performance, improved reliability, and lower operating costs.

Multimegajoule Driver. Our ultimate goal in advanced laser research is to provide the technical baseline for a multi-megajoule driver with which to prove the scientific feasibility of ICF. Table 1-1 summarizes the requirements for this driver and for a short-wavelength laser reactor driver, specifications that have not changed for several years. To obtain gains of several hundred, we will need to deposit several megajoules of energy on a target in a few tens of nanoseconds. For optimum performance, the input pulse must have both a specific shape in time and high beam quality to assure effective energy delivery to a target. Our recent experiments have verified that a laser wavelength of $0.5\ \mu\text{m}$ or shorter will be required. Driver cost is of major importance; the attractiveness of a particular driver concept will be inversely related to its cost. We have chosen a goal of approximately \$250 million for laser hardware to guide our laser design and research programs. Success of the multimegajoule study depends critically on the development of new materials, cost-effective components, and an efficient power-conditioning system coupled with a highly efficient power amplifier. We must evaluate present laser

Fig. 1-9. The ICF Program is engaged in broad-based research and development for advanced lasers.

concepts, and develop new ones for application to this high-gain system.

Solid State Lasers. In 1984, we examined many of the issues associated with the design and construction of a high-energy solid state laser. The results of this study have helped us to define the research directions that will be required to meet our performance and cost goals. The main conclusions of this study were:

- Multimegajoule laser drivers for fusion, with improved performance and reduced unit costs, can be constructed with modest extensions of existing solid state laser technology.
- Significant additional improvements in performance and reductions in system costs can be obtained by developing regenerative or multipass laser architectures.
- Substantial reductions in the cost of individual components will be possible with the development of manufacturing technology for high-volume production of identical parts for a multimegajoule laser system; additional cost reductions will be generated by an innovative research program directed at critical areas.

Laser Research and Development. In 1984, the Laser Research and Development Program was primarily concerned with Nova-related issues. However, some advances with particular relevance to high-energy laser development were made and they indicate the progress that is possible with a more substantial program in this area. We summarize the advances here.

We built an advanced rectangular-disk amplifier whose flexible design allows us to test a variety of components and ideas, i.e., new amplifier materials, coatings, and shapes of flashlamp reflectors. By careful redesign of the flashlamp cavity, the use of a lower-cross-section glass, and the use of a square aperture, we doubled the storage efficiency of a 31-cm laser head over a similarly sized head from Nova, reaching storage efficiencies of 5% at modest energy densities.

We demonstrated a concept for an electrically driven, large-aperture optical switch. We built and tested a 27-cm clear-aperture Pockels cell with plasma electrodes, which uses a KDP crystal similar to those used in the Nova frequency-conversion arrays. This optical switch will allow the construction and testing of multipass and regenerative

Characteristics	High-gain ($\sim 100\times$) experimental facility	Reactor driver (1 GW _e)
Pulse energy	5 to 10 MJ	2 to 6 MJ
Peak power	500 to 1000 TW	200 to 600 TW
Pulse duration	0.030 to 100 ns	~ 10 ns
Pulse shape	Diverse, flexible	Single, fixed
Pulse repetition rate	Single shot	Few Hz
Efficiency	Unimportant except for cost impact	$\geq 10\%$
Wavelength	0.25 to 1 μm	Fixed λ , $\approx 0.5 \mu\text{m}$
Cost goal	25 to 50 \$/J	30 to 60 \$/W _{avg} 100 to 150 \$/J

amplifiers for the first time at an aperture relevant to a high-energy laser. Other possible uses for this technology include the active control of the frequency-conversion process.

We developed damage resistant sol-gel antireflection coatings for Nova fused-silica optics and also applied these coatings to KDP whereby the frequency-conversion crystals are used without in-text-matching liquids and without Fresnel surface losses.

Finally, analytic and modeling work over the past year showed that solid state materials with proper thermal and mechanical properties can provide high-repetition-rate, high-power laser performance. These developments have opened a technological development path from a single-shot multimegajoule driver concept to a solid state ICF reactor driver.

Table 1-1. Requirements for future ICF laser irradiation facilities.

Systems Studies

Our initial and systems applications studies are directed toward evaluating and analyzing ICF applications, their potential, and their payoff, and toward producing thorough engineering studies of promising systems. Working with representatives of other Laboratory programs, we consider specific military and civilian applications of ICF. The results of this kind of system analysis leads to defining technically important, "high-leverage" areas in which ICF can contribute most effectively.

The Cascade fusion reactor concept is being developed as an alternative concept to correct some of the deficiencies in our earlier HYLIFE design and potentially offers impressive improvements. PulseStar is another concept being studied as a possible approach to a more efficient and economical ICF reactor design. These concepts, together with preliminary studies of

Program Resources

multimegajoule target chamber options and other studies are presented in Section 7.

Authors: E. Storm and L. W. Coleman

Program Resources

Financial resources for the Laser Fusion Program in FY 84 totaled \$63.7 million in new appropriated funding and \$98.5 million in expended funding. Program costs included \$60.1 million in operating funds and \$33.6 million in previously appropriated Nova construction funds. The operating budget represented almost a 10% increase over FY 83, approximately even with inflation. This increase was proportionately distributed over the laser development, and fusion theory and experiments categories.

Work force levels were essentially constant at 291 direct employees versus 296 for FY 83. The Nova construction project had an additional work force of 132 direct employees. A 10-year history of funding and work force for the operating program is shown in Fig. 1-10.

Costs for the direct work force were maintained at less than 50% of total costs, continuing our practice of maintaining maximum program flexibility by the use of a high percentage of outside procurements.

Capital equipment funding of \$4.7 million was available to support the operating program in its research endeavors. This level of funding, representing a 5% increase over FY 83, supported our need for increased capabilities in diagnostics, target fabrication, and laser technology development.

Fig. 1-10. Ten-year history of ICF Program resources.

Nova construction continued during the year and was completed in late December, 1984. We were authorized by DOE to delay the date from September to allow the completion of the x-ray laser experiments that were being performed on Novette. Nova had been fully funded in FY 83, and therefore no new funding was required this year.

Author: L. P. Altbaum

Program Facilities

Introduction

Completion of the final phases of supplemental construction at Building 391, Increment II, the Inertial Confinement Fusion Nova Laboratory, was the most significant accomplishment of several major ICF facilities projects concluded in 1984.

The consistent growth and development of the ICF Program, which this year utilized a total work force of more than 650 people, has been guided for the past five years by a highly successful and innovative long-range development plan. The success of this plan has resulted in a current total of 458 605 ft² of office, laboratory, and multipurpose space in 15 permanent buildings, and 26 699 ft² of comparable space in 13 temporary facilities.

Continuing construction, routine and emergency maintenance, and the ongoing renovation of buildings and temporary structures has been managed by the Facilities Group, which provides services in the areas of project design, planning and estimating, material procurement, and construction supervision. During the previous year these operations have required the generation, development, and subsequent completion of 390 separate job orders at a cost of \$2.4 million.

Major Projects

New Construction. Among several jobs concluded, some of the most important were: (1) completing all seismic and GSHA up-grade requirements for our crane systems; (2) installing the main nitrogen

cooling supply system for Nova; (3) installing the emergency air compressor system; (4) constructing and installing the viewing windows in the Nova master oscillator room and the laser bay; (5) upgrading power systems, necessitated by increased demands, for some of our laboratories; and (6) completing the access road and loading area for the central transportainer complex north of the Nova building.

Renovation. The continuing need to modify and improve existing ICF Program facilities demanded the undertaking of several renovation projects of varying size and complexity. These included: (1) remodeling several laboratories; (2) major remodeling and renovation of laboratories for the Advanced Laser Group; and (3) redesigning and revising the interior of the Facilities

Group offices and the Laser Program emergency command center.

Moving. With the completion of stepped phases of work and the redirection of programmatic goals, the evolution of projects often requires the removal or relocation of personnel and facilities and the refinement of project development. Major moving included: (1) relocating the Power Conditioning Group activities to the basement of the Nova building; (2) modifying, then shutting down and moving the Novette laser to the Nova building; (3) removing the Shiva spaceframe and flooring; (4) beginning demolition work in preparation for the new LLNL motor pool facilities; and (5) removing the Nova storage tent facilities.

Author: R. A. McCardle

Our current Nd-glass laser, Nova, consists of ten identical chains, each built on the master-oscillator/power-amplifier model, and each capable of producing 10 kJ of energy in one billionth of a second. Each beam starts as a finely tuned, but weak, oscillator pulse increased in energy by successively larger amplifiers. Lenses and mirrors direct and focus the laser beams' energy, timed for simultaneous focusing onto a tiny target filled with a deuterium and tritium mixture. Nova's light is $1.06\text{ }\mu\text{m}$ (infrared), but KDP crystals can convert it to either $0.53\text{ }\mu\text{m}$ (green) or $0.35\text{ }\mu\text{m}$ (violet) light to couple more efficiently to the fusion targets. In the large photo, a portion of the frequency-conversion array reveals the aluminum egg-crate design holding the nine KDP crystals on each side. Anti-reflective coatings on each crystal increase beam transmission capability. The inset photos at left and top zoom in on details of the egg-crate apparatus. The photo below shows that part of Nova's 4.6-meter-diameter aluminum target chamber where each spool contains a complete frequency-conversion array for final conversion of the beam from infrared to green or violet.

Laser Systems and Operations

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Section Editor: R. J. Foley

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Laser Systems and Operations

Introduction

J. T. Hunt and R. J. Foley

The two major accomplishments of the Laser Systems Group in 1984 were the completion of several diverse target campaigns with the 2-beam Novette laser, and the completion of the 10-beam Nova laser facility.

Novette. The Novette laser filled the need for a short-wavelength driver for multikilojoule target experiments during the time between the deactivation of Shiva and the completion of the 10-beam Nova laser system. Novette also served as a test bed for various Nova system components. A successful series of target experiments with Novette led to the world's first confirmed, laboratory demonstration of x-ray lasing. Fourth harmonic (4ω , $0.26\ \mu\text{m}$) wavelength scaling experiments were also performed on Novette using a 1-kl, 1-ns pulse. On August 30, 1984, the Novette laser was shut down and its two beams were used to complete the Nova laser project.

Nova. On December 19, 1984, the Nova laser construction milestone was completed by firing all 10 frequency-converted beams to the target chamber center. The Nova laser system was completed on the schedule authorized by the Department of Energy (DOE) and within the budgeted \$176 million. Following the Nova activation period that extends into early 1985, we intend to perform target experiments directed at obtaining a better understanding of weapons physics, inertial confinement fusion (ICF) physics, and the physics of x-ray lasers.

Nova Two-Beam Target Chamber. During 1984, we began planning the addition of a second target chamber to the Nova facility. This auxiliary two-beam chamber will essentially double the rate at which target experiments can be performed. Our experience is that the number of target experiments conducted in a given time is limited by target and diagnostic changeover, not by the laser shot rate. The plan allows selective diversion of two Nova beams into the new chamber to do first (1ω , $1.06\ \mu\text{m}$), second (2ω , $0.53\ \mu\text{m}$), and third harmonic (3ω , $0.35\ \mu\text{m}$) target experiments. The initial experiments will be a continuation of the x-ray laser series that was begun on Novette.

Nova

Nova Activation and Performance

At the beginning of 1984, the master oscillator room (MOR) was able to produce alignment beams and 1-ns pulsed Gaussian beams. The single-beam output of the MOR had been split and propagated to the input sensors of the first eight beamlines, and the operation of the automatic alignment control system had begun.

During 1984, we assembled, tested, installed, and activated over 400 major

subassemblies, comprising thousands of unique components, to complete the Nova laser and target systems with their associated pulsed-power, alignment, and diagnostic subsystems. Our activation plan allowed us to install assemblies in the laser bay, optical switchyard, and target areas while continuing the activation of successive stages of the laser beamlines. We first tested, installed, and activated 8 of Nova's 10 laser beamlines. On August 31, we began to dismantle Novette, and subsequently installed the remaining two beamlines from the Novette test bed. On December 19, 1984 we successfully demonstrated the completion of Nova construction by delivering

10 beams of frequency-converted, 3ω light into calorimeters inside the 15-ft-diam target chamber.

The Nova laser fusion facility is capable of irradiating targets at previously unavailable power levels with 1ω , 2ω , and 3ω light. With the combined efforts of our industrial suppliers and an outstanding Nova project team, we completed this effort both on schedule and within the authorized budget. We are pleased with the reliability and demonstrated performance of this system during its very first days of operation and look forward to complete activation for target experiments in early 1985.

Activation History. The sequence of events leading to the successful activation of Nova is shown in Table 2-1. In January, assembly, installation, and checkout activities were proceeding in the laser bay, optical switchyard, and target room. Using a pulsed beam, we began first to align and then to activate successive laser stages beginning with the input of each laser chain.

We activated the 220-ft-long driver sections through the mid-chain sensor of the first eight beamlines by February. Activation of the high-energy power amplifier sections of the chains commenced in March with the alignment of the first 31.5-cm amplifier heads, followed closely by the final 46-cm amplifier stages. We activated beamlines in parallel within each of three groups (consisting of 3, 3, and 2 chains, respectively). This grouping was based on the scheduled availability of assemblies and the efficiencies of activating beamlines in parallel. Installation of full-aperture mirrors in the optical switchyard and target area proceeded as these massive optical components became available.

The three months added to the schedule for completion of the x-ray laser target campaign on Novette were well spent. First, we characterized single-beam laser and frequency-conversion performance at various energy levels and pulse widths. Second, by using slower feed rates on our

Table 2-1. Chronology of Nova activation during 1984.

January	<ul style="list-style-type: none"> First output sensor successfully activated in switchyard. Favorable financial status allowed DOE to authorize target diagnostic instrumentation as part of project.
February	<ul style="list-style-type: none"> Driver sections for eight beamlines activated: 3.9 kJ delivered to mid-chain sensors in 1-ns pulse. First full beamline installed.
March	<ul style="list-style-type: none"> Target chamber successfully tested under vacuum. Pulsed activation of first three beamline power sections initiated.
April	<ul style="list-style-type: none"> Successfully activated first two full beamlines: >2 kJ each in a 700-ps pulse (1ω). Capacitor bank installation completed for eight beamlines. Automatic alignment of first full beamlines demonstrated. First fused-silica final focus lenses coated.
May	<ul style="list-style-type: none"> A total of four beamlines fully activated. Laser bay, switchyard, and target room converted to clean rooms. First two beamlines completed through switchyard to target room. Novette shutdown postponed by one month, delaying Nova scheduled completion until October 31.
June	<ul style="list-style-type: none"> Eight beamlines installed up to the target area. Eight output sensors activated. First KDP frequency-conversion crystals antireflection-coated successfully.
July	<ul style="list-style-type: none"> Eight beams produce 57 TW of power in a 1-ns pulse. Nova is now the world's most powerful laser. Further experiments on Novette require second postponement of shutdown. First frequency-converted 2ω light delivered to target chamber from single beamline (2.5 kJ, 1 ns). Installation of additional 31.5- and 46-cm power amplifiers initiated.
August	<ul style="list-style-type: none"> First 3ω light delivered to target chamber (3.1 kJ, 1 ns). Novette shutdown formally delayed by DOE until August 31. Nova construction completion scheduled for December 20.
September	<ul style="list-style-type: none"> Transfer of final two beamlines from Novette initiated. Installation of the first target diagnostic instrument started on target chamber (Dante soft x-ray spectrometer).
October	<ul style="list-style-type: none"> 1ω-diagnostic system installation completed for all 10 beamlines. Capacitor bank installation completed for all 10 beamlines.
November	<ul style="list-style-type: none"> All laser components installed, tested, and activated. Optical train to target chamber complete for 10 beamlines. Pulse-shaping capability demonstrated in MOR.
December	<ul style="list-style-type: none"> Ten beams fired simultaneously into target chamber center produce 18.2 kJ of 3ω light in a 1-ns pulse on December 19. The Nova construction milestone is complete.

diamond-turning machines, we achieved higher-quality surface finishes on our potassium dihydrogen phosphate (KDP) frequency-conversion crystals.

By December, all components necessary to demonstrate system operability were in place on the laser, in the optical switchyard, and on the target chamber. These components included the full complement of fused-silica focusing optics, diagnostic equipment, and the KDP frequency-conversion arrays that provide 2ω and 3ω light from the 1ω infrared output of the Nova laser chains. A series of activation test shots with the laser culminated when we fired 10 beams simultaneously into 44-cm calorimeters mounted inside the evacuated target chamber. This shot produced over 18 kJ of 3ω light in a 1-ns pulse; it established Nova as the world's most powerful 3ω laser.

We plan to operate Nova at higher energy and power levels during 1985, as we gain experience with the system in preparation for target experiments.

Initial Performance Results. For the construction milestone shot of December 19, 59.7 kJ of 1ω light in a 1-ns pulse were generated to produce 18.2 kJ of 3ω light at the target chamber center. Because the frequency-conversion arrays perform better at high intensity levels, it is advantageous to drive them with pulses that are temporally square rather than Gaussian. Accordingly, in September, we accelerated the development of non-Gaussian pulse-generation capability for the Nova MOR. This capability was operational for the milestone shot.

From the MOR, a high-contrast-ratio, linearly rising, ramped pulse was propagated to the input of the amplifier chains (Fig. 2-1). By ramping this pulse, we compensated for high-energy saturation effects, which cause a decrease in the gain in each amplifier chain over the duration of the pulse as stored energy in the neodymium-doped phosphate laser disks is depleted. As shown in Fig. 2-2, we were successful in producing nominally square pulses at the output of the chains and thus also at the conversion arrays.

The average intrinsic conversion efficiency for the arrays was approximately 36.3% for this shot. This figure is adjusted for the measured 16% transmission loss between the measured 1ω output energy and

▲ Fig. 2-1. High-contrast, linearly rising, ramped pulse generated in the MOR to compensate for gain saturation in the amplifier chain to produce a nominally square output pulse

Fig. 2-2. High-intensity pulse exiting the amplifier chain is nominally square prior to propagation into the frequency conversion array.

the target chamber center. This loss is due to beam obscuration, absorption, and coating losses for mirrors, lenses, and crystals in the path of the beams. The 36.3% conversion efficiency is lower than that theoretically predicted by the frequency-conversion models. We plan to investigate the sources of this discrepancy as we prepare Nova for target experiments next year.

We completed a series of initial experiments designed to characterize system performance based upon single-beam measurements. On these experiments, we produced over 15 kJ of 1ω light on a single chain, 7.2 kJ for 2ω light, and 7.4 kJ of 3ω light in 2.5-ns pulses. The comparable long-pulse (3-ns) criteria in the minimum Nova baseline are 8 kJ, 5 kJ, and 4 kJ respectively. We have, in addition, demonstrated single-array intrinsic conversion efficiencies of 62% for 2ω and 59% for 3ω . These initial experiments give us a high degree of confidence that Nova will meet its formal performance criteria, and we expect to gradually

increase Nova's performance as we prepare for target campaigns in 1985.

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Nova Frequency Conversion and Focusing System

New developments in crystal array technology provided significant improvements in the mechanical design and optical performance of the Nova $2\omega/3\omega$ array hardware. The final Nova array configuration was tested on the Novette laser and on the first arm of Nova.

Ten Nova $2\omega/3\omega$ crystal arrays were assembled and tested for crystal alignment and wave front distortion before installation on the Nova target chamber.

Ten Nova focus lens positioners were assembled and tested last year. The positioning accuracy and repeatability of each assembly were evaluated before installation on the target chamber. A cylindrical focusing system was also developed for installation in the Nova lens positioner assembly. Finally, 10 completed frequency conversion and focusing systems were activated.

Crystal Arrays. The demonstration in 1984 of a high-damage-threshold, room-temperature antireflective (AR) coating

for KDP (see "Sol-Gel Antireflective Coatings," in Sec. 6) was a significant technical breakthrough in the development of larger-aperture crystal arrays. In conjunction with the improved optical quality of diamond-turned KDP crystal surfaces, the development of this coating permitted a significant change in direction for harmonic generation array technology. The early Nova array designs, which used index-matching fluid between crystal and windows, were limited by wave front distortion induced by hydrostatic pressure. The early designs were also prone to fluid decomposition at high drive fluences and to incipient transverse stimulated Raman scattering losses. It was for these reasons that we developed a fluidless design, which was first tested in the bare-crystal, Type-II-Type-I, 4ω array installed on the Novette laser.

The final Nova array design (Fig. 2-3), was initially tested on one arm of Nova with bare crystals and then with AR-coated crystals. A transmission increase of approximately 12% over a bare KDP array was observed in the harmonic output. The full 16% Fresnel loss was not recovered because these coatings are simple quarter-wave-thickness interference coatings designed to have a high transmission (approximately 99% per surface) at two wavelengths.

The sol-gel coating was applied by a spin-coating technique similar to that used in the semiconductor industry for photoresist deposition. As shown in Fig. 2-4, the coating solution is applied to the center of a spinning crystal located in the large stainless steel tank. The crystals used in the array, both doublers and mixers, are identical; up to the coating step, they can be placed in any location in the egg crate. The AR coating applied to doublers is a $1\omega/2\omega$ compromise and is the same on both surfaces of the crystal. However, the crystals used for mixing have a $1\omega/2\omega$ coating on the input side and a $2\omega/3\omega$ coating on the output side. An array load map is prepared before AR coating so that any defects, such as edge chips, are placed outside the usable aperture. Once the crystal location is determined, the correct AR coating can be applied.

The crystals are mounted in a deep, 10-cm web aluminum egg crate that has been single-point diamond machined to a

Fig. 2-3. Nova 80-cm KDP crystal array employs a deep-web aluminum egg crate with nine AR-coated crystals mounted on each side.

Fig. 2-4. KDP crystal spin-coater: square or triangular crystal plates are mounted on the rotating fixture in the center of the coating tank. Coating solution is applied by a syringe above the tank.

flatness better than $10\text{ }\mu\text{m}$, determined interferometrically before assembly. Each crystal is held in the egg crate so that it is parallel to the flat egg crate surface. A 2-mrad wedge is machined into the egg crate to prevent multiple-surface interference effects. Data taken on the 10 Nova array assemblies show that the crystals are held in the egg crate to a precision of approximately $\pm 30\text{ }\mu\text{rad}$ in the internal phase-match angle.

Each Nova crystal array assembly was tested for wave front distortion in the large-aperture interferometer. The typical single-pass transmitted wave front distortion per 27-cm cell was less than $\lambda/4$ at $0.6328\text{ }\mu\text{m}$.

We have characterized the conversion efficiency of large-aperture KDP arrays for the following four configurations:

- Type-II, 2ω , 5×5 , 15 cm^2 ; crystal thickness, $L = 1.8\text{ cm}$.
- Type-II-Type-II, 2ω , 3×3 , 27 cm^2 , $L_1 = L_2 = 1.0\text{ cm}$.
- Type-II-Type-II, 3ω , 3×3 , 27 cm^2 , $L_1 = L_2 = 1.0\text{ cm}$.

- Type-II-Type-I, 4ω , 5×5 , 15 cm^2 , $L_1 = 1.8\text{ cm}$, $L_2 = 0.8\text{ cm}$.

These arrays were tested on Novette or Nova and, with the exception of the 4ω array, the conversion efficiency followed theoretical predictions up to a drive intensity $I_w \approx 2\text{ GW/cm}^2$ with pulse widths τ_w , ranging from 0.1 to 2.5 ns. Rapid departure from theoretical values was observed above this intensity. Our efforts to understand the reasons for this behavior are focused on the following areas:

- Transverse stimulated Raman scattering in KDP.
- Longitudinal stimulated Raman scattering in air.
- Intensity-dependent polarization mix-ratio (ellipse rotation).

The 4ω -array performance was well below theoretical values over the operating intensity range.

Lens Positioner. Ten Nova lens positioners were assembled and tested, as shown in Fig. 2-5, for positional accuracy and repeatability under vacuum load and at the

Fig. 2-5. One of the 10 Nova lens positioners mounted on test stand for testing at operating attitude before installation.

Fig. 2-6. Top view of Nova focus configuration with cylinder lenses tilted with respect to optical axis. The individual line foci will form one long continuous line focus.

operating attitude. The x-y positional accuracy and repeatability were $<10\text{ }\mu\text{m}$ on all 10 units. Focusing (z) accuracy and repeatability to within $5\text{ }\mu\text{m}$ were achieved. Cross-coupling between the z and transverse (x-y) directions was measured; over a typical 2-cm focus range, it was $<25\text{ }\mu\text{m}$. However, over large z excursions, transverse-z cross-coupling of $\sim 50\text{ }\mu\text{m}$ was observed.

Nova Cylinder Lenses. The Nova target illumination geometry, a 100° cone, rules out the Novette scheme for illumination of x-ray targets shown in an article later in this section, "Target Focusing Configuration for X-Ray Laser Experiments." Because the axes of the focus lenses and the x-ray target are collinear, the line foci of a cylinder lens are parallel to the target in Novette. A side view of the Nova focus configuration

(Fig. 2-6) shows that if the Novette scheme is used, the tilted line foci would create unacceptable variations in illumination intensity across the target. A simple modification of the Novette configuration removes these undesirable line foci tilts when the cylinder lens (located in converging light) is tilted by a properly chosen angle. As the cylinder lens is inclined, one of the line foci rotates and the other is unaffected.

The x-ray laser target is mounted perpendicular to the target chamber axis and is tipped 18° away from the horizon. This is one of the target orientations (the other is parallel to the horizon) that preserves the maximum degree of symmetry. In this orientation, x-ray diagnostics can be positioned above and below the horizontal plane to avoid other diagnostics.

The ACCOS V lens design code was used with calculated values of line foci rotation to determine the actual focal lengths and tilt angles for the Nova cylinder lenses. Three different cylinder lens designs are required to achieve uniform illumination when the line foci are stacked end-to-end. To provide a 1.0-cm line length from each beam, the focal lengths of the negative cylinder lenses vary from -200 to -310 m , and the tilt angles range from 5 to 21° . Other stacking arrangements allow targets 1 or 2 cm long to be illuminated. The required focus lens motions are within the range of the lens positioner.

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Optical Components for Nova

Highlights. The highlights of 1984 in optical components are:

- The implementation of a new LLNL-developed sol-gel coating for fused silica and KDP.
- The achievement of low-scatter KDP surfaces by diamond turning.
- The completion of all optical components required for Nova construction.

Sol-Gel Coatings. For several years we have been developing alternatives to vacuum-deposited metallic oxide AR coatings (such as $\text{SiO}_2\text{-ZrO}_2$) because of their low threshold for laser damage. For

BK-7 and related borosilicate glasses, Schott developed the neutral solution process,¹ a chemical surface treatment that produces a broadband AR coating with a damage threshold at least as high as that of bare surfaces. But the neutral-solution process cannot be applied to the fused silica focusing optics required by the Nova third-harmonic light. For this purpose, alternatives were sought in sol-gel technology,² because early tests had shown that sol-gel coatings have high damage thresholds. Contributions to this technology were made by Schott, Corning, Battelle Northwest, and especially by Westinghouse.

Early in 1984 the LLNL ICF research group achieved a significant breakthrough, obtaining a colloidal sol that could be deposited as a quarter-wave layer on optics of any size or type at room temperature, and that required no further processing after deposition. (Previous approaches had required heating up to 500°C and acid etching, which limited the type of substrate to which the coating could be applied.) Three methods of deposition were developed:

1. Draining the fluid away from the optic submerged in a tank.
2. Withdrawing the optic from a tank of the fluid ("dip-coating").
3. Spin-coating.

The fact that an AR coating could be applied to KDP crystals in the Nova frequency conversion system (which has four KDP surfaces) made it possible to dispense with the index-matching fluid and surrounding windows, resulting in better optical quality and higher damage resistance. These changes are discussed in the preceding article, "Nova Frequency Conversion and Focusing System." To make this change possible, the diamond-turned crystal surfaces had to be made significantly smoother, as described below.

The coating solution, an alcohol-based colloidal sol, deposits a uniform layer, consisting of stacks of small silica spheres, on the surface. The sphere size is uniform and can be adjusted, within limits. The size criterion is that the spheres must be small enough not to scatter the shortest wavelength light of interest. For the Nova optics, the required sphere diameter is about 200 Å. The average index of refraction of the layer, which is part air and part silica,

is about 1.2 (the square root of the index of the substrate); when the optical thickness of the layer is $nt = \lambda/4$ (that is, when its geometrical thickness is $t = \lambda/4n$) the coating is perfect for that wavelength and is useful over a fairly broad range around that wavelength.

The thickness achieved is determined by the initial concentration of silica in the solution and (in the case of dip-coating) by the withdrawal rate. The maximum thickness obtainable for 200-Å particles is about 100 nm, which is roughly $\lambda/4$ for 0.50- μm light. To obtain thicker coatings, multiple layers are applied. For use at two wavelengths, the optical thickness is made $\lambda/4$ for an intermediate wavelength. Thus, a coating optimized for the second and third harmonics is a single coating 90 nm thick, while a coating optimized for the first and second harmonics is a double layer about 160 nm thick.

The colloidal sol coatings for Nova were applied to the 80-cm-diam fused silica focusing lenses (see Fig. 2-7) and debris shields, and to the KDP crystals; such coatings will also be applied to the 31.5- and 46.0-cm aperture fused silica input spatial filter lenses that will replace the current

Fig. 2-7. Nova focus lens being lowered into coating tank.

BK-7 lenses. This coating may be applied to any other laser optics for which high laser damage resistance is desired.

Like the neutral solution coating, these sol-gel coatings have high damage resistance; they have the drawback that, at the current state of the art, they are quarter-wave coatings rather than broadband. As

Fig. 2-8. Transmission of a quarter-wave coating peaked at 440 nm.

Table 2-2. Cumulative lasing (through 1984) of large-aperture coated optics in Nova. ▼

Component	Diameter or dimensions (mm)	Aperture (mm)	Nova quantity	Quantity delivered
5-cm rods	50	42	16	20
Tab 01 rotators	100	95	10	11
Tab 02 rotators	166	156	10	11
Tab 03 rotators	220	210	10	11
Tab 04 rotators	331	321	10	14
Tab 04A mirrors	604	576	10	13
Tab 04B mirrors	604	509	10	13
1.5% beam splitters	420 × 240	218	10	11
Tab 10 polarizers	103 × 60	50	60	65
Tab 01 polarizers	196 × 115	160	40	47
Tab 02 polarizers	330 × 190	160	10	13
Tab 03 polarizers	420 × 240	218	20	30
Tab 04 polarizers	700 × 400	335	20	40
RBD ^a mirrors	820	790	10	12
Diagnostic mirrors	940	908	10	11
Turning mirrors	940	900	24	28
Turning mirrors	1090	980	4	5
Objective lens	800	770	10	10

^aRBD = return-beam diagnostic.

Fig. 2-9. Spectral reflectance and uniformity of Nova target chamber turning mirrors (94 and 109 cm diam).

shown in Fig. 2-8, the result is that reflection losses reduce the two-surface transmission for the second and third harmonics to about 99%. The results we have achieved on fused silica agree with theory within the limits of photometric measuring accuracy.

The basic equipment used for coating the large fused silica optics was the drain tank provided by Westinghouse. For coating the KDP, two approaches were taken, dip-coating and spin-coating. (Most of the spin-coating production development was done by the frequency converter group in conjunction with their spin-cleaning equipment.)

KDP was much more difficult to coat than fused silica, for several reasons:

- KDP is water.
- It is hygroscopic.
- It is chemically active.
- It is piezoelectric, resulting in charges on the surface.
- The diamond-turned optical surfaces have some residual roughness.

The KDP coating process is size sensitive; that is, processes and handling that result in good coatings on 5-cm-diam samples may not consistently work on the 27- × 27-cm final parts. For large parts, dip-coating gave some good results, but reliable production consistency was never achieved; all the Nova array crystals (360 on the laser, plus about 50 spares) were successfully spin-coated. The crystal coating process was integrated with the array assembly process in a clean room. Further details on spin-coating are reported in the preceding article.

Future work on sol-gel AR coatings will include development of broadband coatings for fused silica and KDP and further work

on KDP undercoating (there seems to be some promise in the use of nonporous methyl silicone to prepare KDP surfaces for coating, but work done so far has been inconclusive).

Vacuum-Deposited Coatings. In all, 48 Nova mirrors, with apertures from 79 to 98 cm and diameters up to 109 cm, were coated in 1984. Table 2-2 gives the cumulative list of large-aperture Nova optics coated through 1984.

A coating design modification during the last half of 1984 reduced the variation in reflectance near 3ω , which results from slight variations in thickness. This improvement allowed us to approach our goal of $\pm 3\%$ reflectance over a 1-m diameter. The measured thickness variation across the mirrors is about $\pm 0.5\%$. Figure 2-9 shows a calculated plot of the modified 3ω coating design and the measured 3ω reflectance across the diameter of target-chamber turning mirrors (94 or 109 cm) to 88 cm.

Table 2-3 gives the specifications for the 109-cm diam, 33° -incidence mirror. Figure 2-10 shows a 109-cm-diam mirror on the large automatic scanning photometer at LLNL.

Neutral Solution Process. All 110 Nova BK 7 spatial filter lenses were AR coated using the neutral solution process (NSP). Table 2-4 lists the diameter and aperture of each type of lens. Ten of each of these lenses have been neutral-solution processed. As noted in the table, some lenses have

Table 2-3. Specifications and goals (desired values) of damage threshold and reflectances for p and s polarizations for the 109-cm-diam, 33° -incidence Nova mirror.

	1.053 μm		0.527 μm		0.351 μm	
	R_p (%)	Damage ^a (J/cm ²)	R_p (%)	R_s (%)	R_s (%)	$R_{p, \text{unit}}^b$ (%)
Specification	$\geq 98.5^b$	≥ 8	≥ 95.0	$\geq 37^c$	$\leq 54^c$	± 5
G.al	≥ 98.8	≥ 10	≥ 98.8	$\geq 34^d$	$\leq 50^e$	± 3

^aDamage threshold for a 1-ns pulse.

^bAverage $R_p + R_s \geq 99\%$.

^cAverage of ~ 20 data points taken on the major axis.

^dMinimum value on mirror surface.

^eMaximum value on mirror surface.

Fig. 2-10. 109-cm-diam turning mirror on large automatic scanning photometer at LLNL.

Table 2-4. Diameter and aperture of Nova BK-7 spatial filter lenses AR coated by the neutral-solution process. Ten lenses of each type were processed.

Nova

Tab number	Diameter (mm)	Aperture (mm)
06, output	800	770
05, input ^a	520	460
05, output	520	460
04, input ^a	363	315
04, output	363	315
03, input	240	208
03, output	240	208
02, input	180	150
02, output	180	150
01, input	130	91.7
01, output	130	91.7

^aReplaced with fused silica lenses AR coated by the sol-gel process.

been replaced with fused silica lenses AR coated using the sol-gel process.

Plastic Membranes for Debris Shields. During 1984 10 samples of 50-cm-clear-aperture membranes were made by replication for evaluation as alternatives to the standard fused silica debris shields. The membranes, made of a Fenwal UV-transmitting epoxy that was purified by settling, had a thickness of approximately 0.002 in. The parts were easily replicated, but they were very difficult to separate from the flat masters. This difficulty prompted a small study, whose results suggest that Lexan might be substituted for the epoxy. Lexan is a castable material, less expensive than epoxy. It is transparent to wavelengths as short as 0.25 μm ; it has been produced in 5-cm-diam test pieces, 3500 Å thick, with a transmitted wave front flatness of $\lambda/10$.

KDP Production and Finishing. Overall, KDP growth and production of 27-cm-square KDP blanks in 1984 were very successful. Cleveland Crystals Inc., the major supplier of blanks for Nova, completed all their production contracts on schedule. A total of 146 Nova blanks were completed. Cleveland's performance was the result of the following factors:

1. Sufficient quantity of high-quality boule material.
2. High yield of acceptable blanks from boule material. (Rejection rate of finished blanks due to defects and fabrication problems was less than 5%.)
3. Adequate production machinery to allow parallel effort on shaping and roughing operations.
4. Effective and efficient use of personnel on double shift operations.

The net result was a delivery rate of 20 to 24 blanks per month, twice the maximum delivery rate in 1983.

Our second vendor, Inrad, has delivered 60% of their contract requirements. Present delivery rates should complete the contract by May 1985.

KDP Final Finishing. For final finishing we had two outside vendors in 1984: Cleveland Crystals, with two diamond turning machines, and Optics West Division of Pneumo Precision, with one machine. One LLNL machine was also available. The production requirement was 140 whole equivalent crystals, plus spares. [Each nine-element frequency-conversion array uses five whole crystals and four diagonally cut half ("corner") crystals, or the equivalent of seven whole crystals. The entire laser uses 100 whole crystals and 80 half crystals.]

Considerable effort was put into bringing up and qualifying the Cleveland and Optics West machines and facilities. LLNL's machine started machining blanks in early 1984 at a rate of four to six per month. Optics West started delivery in late January and was operating at a rate of one blank per week. LLNL's Metrology Group assisted Cleveland Crystals in modifying the spindle slides, adding protective covers for the machine ways, and stiffening up the entire spindle housing and ways. This reduced the machine compliance from approximately 120 $\mu\text{in./lb}$ to 50 to 60 $\mu\text{in./lb}$. With these improvements, the machines produced crystal surfaces that met our original specifications.

However, as mentioned earlier, the development of a successful AR coating for KDP crystals led to the decision, in April, to eliminate the index-matching oil and surrounding windows from the array design. The use of bare crystals required much more stringent specifications on rms surface roughness and peak-to-valley (P-V) finish. The requirements on crystal and wave-front flatness, wedge, and phase matching were not significantly affected; the absolute thickness tolerance was relaxed somewhat.

Studies were conducted at LLNL to establish the minimum surface finish that would achieve the following conditions:

1. Minimum scattering from an uncoated surface, such that transmission was maximized with Fresnel losses ($\sim 4\%$ /surface)

Table 2-5. KDP machining specifications.

Variable	Original specification	New specification
Thickness (mm)	10.0 ± 0.0015	10.0 ± 0.1
Wedge (mm)	± 0.0015	Same
Phase angle (μ rad)	± 30	Same
Transmitted wave front flatness at $0.6328 \mu\text{m}$	1%	$\lambda/4$
Surface finish (Å)	250 rms 500-750 P-V	30-50 rms 150-250 P-V

being the only significant contribution. (Crystals cut to original specification gave finishes of 200 to 300 Å rms and 800 to 1000 Å P-V, resulting in additional scattering losses of 10 to 15%.)

2. A surface finish that would accept consistent and uniform AR coatings with less than 1% loss per surface.

In connection with this study, experiments were carried out on 5-cm-square KDP blanks to determine the new finishing parameters. Optics West and Cleveland Crystals continued producing crystals to the old specification, but slightly over thickness, to allow for remachining once the new specification was established. Approximately 90% of the machining operations could be completed this way and only finishing cuts would be needed. This maintained the production schedule while development was in process.

Surface finish is affected by a variety of parameters:

1. Random motion inherent in the machine, arising from stiffness (compliance), vibration, spindle accuracy, and smoothness of crystal translation past the tool.
2. Variable and controllable parameters, such as diamond tool quality, tool radius R , feed rate f (the rate of crystal translation past the tool, in in./min), and spindle rotation speed W .

The theoretical P-V roughness h of the surface produced by a single-point diamond machine is given by

$$h = \frac{F^2}{8R} \text{ inches} \quad (1)$$

where $F = f/W$, with W in rpm, and where R is in inches. The theoretical surface finish is independent of the depth of cut of the tool, assuming that the cut does not fracture the surface (a critical factor in machining KDP).

Typical parameters used in the initial KDP finishing requirements were $W = 1000$ rpm, $R = 0.1$ in., and $f = 0.5$ in./min, giving $h = 0.3 \mu\text{in.}$ (80 Å) P-V.

Random machine factors lead to a surface roughness considerably greater than indicated by h , however: a typical surface machined with the parameters just stated had a finish of ~ 150 to 250 Å rms or 750 to 1000 Å P-V. We machined a series of KDP

samples, using varied machining parameters, until we obtained a surface finish of about 30 to 50 Å rms and 150 to 250 Å P-V.

Sample performance was inferred from transmission measurements on the crystal. A 1.5-in.-aperture, 0.442-nm beam from a He-Cd laser was passed through the crystal and focused on a $100\text{-}\mu\text{m}$ pinhole. A detector measured the energy passing through the pinhole. Crystals exhibiting high P-V readings produced significant scattering and reduced the measured transmission to 60 to 35%. Machining parameters were varied until the transmission was increased to 90 to 92%.

The machining parameters that produced the desired transmission were $W = 1000$ rpm, $R = 0.2$ in., and $f = 0.1$ in./min, corresponding to $h = 0.01 \mu\text{in.}$ (3 Å) P-V. The actual surface finish, as determined by the profilometer, was 30 to 50 Å rms and 150 to 250 Å P-V, as desired. Table 2-5 summarizes the resulting changes in specifications established with these tests.

LLNL's machine was capable of producing this finish, and the Optics West machine demonstrated such performance under ideal conditions; Cleveland Crystals' machines could not meet the new specification, however. After the Optics West facility was upgraded with new air conditioning to maintain $\pm 1.5^\circ\text{F}$, and machining fixtures were stiffened, the machine consistently met the new specifications (see Fig. 2-11).

To meet our crystal production schedule, Cleveland Crystals continued to finish blanks to the original specification, leaving extra stock (0.001 in. per side) for final finishing at LLNL or Optics West. All facilities went on three-shift operation. In addition, to evaluate machining performance on a surface-by-surface basis, LLNL and Optics West each installed a WYCO noncontacting profilometer. This equipment gives rapid verification of surface finish; Fig. 2-12 shows typical crystal-surface profile data.

Fig. 2-11. Single-point diamond turning machine at Optics West. Machine has 18-in.-diam flywheel and 26-in. precision travel in x direction.

Sufficient crystals (100 squares and 80 corners) were finished to complete 10 assemblies by the scheduled date of December 7. Figure 2-13 shows typical transmitted wave front interferograms.

One additional factor was important in the timely production of the required crystals. Early in the machining process it was decided that all crystals could be machined as squares, and the necessary corner pieces cut after finishing. This reduced corner-blank machining costs and time by half. This operation was accomplished by using a modified Bridgeport mill with a 5-in. diamond circular saw blade and by supporting the crystal on a split vacuum chuck (see Fig. 2-14). The diagonal saw cut was made

Fig. 2-12. Typical W/CO profilometer plots of KDP surface finish. Graphs represent top, middle, and bottom of each crystal surface. Measurement covers $665 \mu\text{m}$ across surface.

Fig. 2-13. Typical interferograms of transmitted wave front flatness of crystal #228-09 at 0.6328 μm (He-Ne laser). Corresponding flatnesses are $\lambda/5$ horizontal and $\lambda/3.6$ vertical.

at 0.25 in./min. Crystals were flooded with mineral oil lubricant held at constant temperature throughout the cutting cycle. In this way 31 crystals were cut, producing 62 corners, with no breakage or loss. (The remaining corners were machined as corners, because the crystal growers were allowed to deliver some in this form.)

The rapid development and implementation of techniques for high-tolerance optical finishing of critical UV components, and the production of these components within a tight schedule, represents a significant achievement by LLNL and its vendors.

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Pulse Generation and Shaping

The general layout of the MOR was re-worked significantly in 1984 to permit the installation of the pulse shaping hardware, which is now in operation. Fig. 2-15 shows the new pulse shaping system that produces 1- to 5-ns shaped pulses, the picket fence generator that produces 1 to 4 Gaussian pulses, and the oscillator for the fiducial system. The pulse synchronization system was also added this year.

The principles of the pulse shaping system have been described previously.³ The pulse shape is defined by changes in

Fig. 2-14. Bridgeport mill modified to cut square KDP crystals to corners. Slotted vacuum chuck allows saw blade to cut crystal in single pass.

impedance along the length of a variable impedance transmission line. A cross section of the conductor geometry is shown in Fig. 2-16. The center conductor-to-ground spacing S determines the impedance at any point. The center conductor is a thin metal strip a little over 1 m long supported from the top by a series of 23 micrometers. The box surrounding the strip is the ground conductor. This assembly is connected at one end to a box containing the Auston switches, and impedance matched to a 25- Ω microstrip line. A pulse shape is selected by adjusting the micrometers to a previously determined position or by making adjustments while observing streak camera data. These two assemblies constitute the pulse shaper driver, shown in Fig. 2-17.

Fig. 7-15. A portion of the revised MOR layout showing the pulse shaping system, picket fence generator, and fiducial system oscillator.

Fig. 7-16. Conductor geometry of the variable impedance transmission line in cross section.

The top of the shaped pulse is defined by a continuously rising waveform from the pulse shaper driver that drives two Pockels cells. The rising and falling edges are defined by a square gate from the pulse slicer driver, which drives the two Pockels cells in the slicer. The slicer uses two Auston switches in a "tail biting" arrangement, which allows the use of a fixed length of charge line between the switches. Output pulse length is determined by pulse illumination timing on the switches.

The silicon switches are 17 mm wide with a 2.5-mm interelectrode gap. The dark current at 10 kV is kept low by reducing carrier lifetime near the contacts. In contrast to the method reported previously,⁴ we now sandblast the silicon surface to produce a high density of traps before applying contacts. The switches are attached contact side down on a 25- Ω microstrip line and illuminated from the back. This 25- Ω line is split into two 50- Ω lines at the output, and standard connectors and cable are used to couple the output to two 50- Ω Pockels cells.

To illuminate the switches, we use a second switchout to select one pulse from the rejected pulse train, which is diverted by the short-pulse oscillator switchout. This pulse is amplified by a factor of 100 in an yttrium-lithium-fluoride (YLF) amplifier consisting of an oscillator head driven by a 10-mm rod amplifier power supply. The output energy is about 9 mJ. This energy is distributed to the silicon switches by a series of pickoffs along the beamline. Moving

these pickoffs (a waveplate-polarizer assembly) sets the timing for the system.

The picket fence generator is a set of passive optics that splits a single pulse from the short-pulse oscillator into four beams and delays each beam before recombining them into one line. The system reflects the beams so that they emerge parallel to one another; the only alignment adjustment required is centering. Each pulse is independently delayed over a 10-ns range. The short-pulse beam normally goes through one of the delay arms with all the energy directed down the arm. Since the pulse timing is set by the position of a retroreflector, timing can be precise and repeatable to within a few picoseconds.

Operation of the pulse shaping system has been very successful. We have been able to use theoretical predictions of the required pulse shapes and streak camera data taken in the MOR to produce pulses that compensate for saturation. The rise and fall times are close to 150 ps, which makes a well-defined square pulse even at a 1-ns pulse width. Figure 2-18 shows streak data from the MOR and the output of the laser at 9 kJ. The contrast ratio between the initial and final intensities of the shaped pulse is about 3.5.

The versatility of the MOR is now greatly enhanced, and any pulse shape we now think will be required for experiments can be produced.

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Power Systems Energy Storage

Capacitor Bank. The capacitor bank for Nova was completed during the past year. At the beginning of 1984, the Nova capacitor bank was capable of supporting 8 of Nova's 10 arms. During the year, the Novette capacitor bank was moved to the Nova laser facility. In addition, the capacitor bank was expanded to accommodate the additional 31.5- and 46.0-cm amplifiers.

At the end of 1983, the capacitor bank was sized at 42.35 MJ at 22 kV. It consisted of 1258 circuits, driving 3476 flashlamps and 34 Faraday rotators. The energy from the 8376 capacitors was transferred by 84 switches. The bank was charged by six

Fig. 2-17. Variable impedance transmission line and box containing Auston switches and charge line; both are pressurized with SF₆ at 1 psi.

Fig. 2-18. MOR and output streaks for a 9-kJ shot.

MVA-type and eight KVA-type power supplies. By the end of 1984, the size of the bank was increased to 62.30 MJ at 22 kV. It now contains 1676 circuits driving 5336 flashlamps and 42 Faraday rotators. The number of capacitors was increased to 10 401 and the number of switches to 117. The number of power supplies remained the same.

From January to September, the bank was operated daily to support operations and activation of the laser itself. The bulk of the increase in the capacitor bank was installed during September and October. After installation, the procedures described in the *Laser Program Annual Report 83*⁵ were used to activate the added circuits.

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Power Conditioning Controls. The Nova power conditioning control system is divided

into three major areas: pulsed-power, system-timing, and facility controls. The pulsed-power controls command and monitor the various pulsed-power hardware including the KVA- and MVA-type power supplies, the high-voltage ignitron switches, and the lamp-check diagnostics (LCD) packages. The system-timing controls provide timing triggers for synchronizing laser components and diagnostic subsystems. The facility controls are used to monitor the safety interlocks, building security, and nitrogen supply subsystem. The majority of all pulsed-power and system-timing controls were completed during the year and installed on Nova. Much of the facility monitoring hardware has been installed,

Fig. 2-19. Power conditioning control system, installed configuration.

Fig. 2-20. Pulsed-power menu displayed on central operator console, power conditioning system.

and the software development has begun. Figure 2-19 illustrates the installed system configuration.

Operators now control the power conditioning system from the central console. Three main color menus are provided, one for each of the major control areas. The pulsed-power menu, shown in Fig. 2-20, allows operators to select a variety of laser components, groups of components, or laser arms by touching pictures of the desired components on the screen. The status of each component is color coded. Once operators choose a component configuration, they may select the desired pulsed-power sequence [Test, Pulse-Ionization Lamp Check (PILC), Dry-Run, or Shot]. The underlying control software automatically selects the appropriate hardware and performs the necessary sequencing and verifications. The system-timing menu allows the operator to directly control the timing sequences and to select various timing triggers and rates.

An auxiliary-control program, which runs on a DEC VT125 graphics terminal, allows direct access to each hardware device in the system and performs low-level control and maintenance operations. Unlike

programs displayed on the central power conditioning console, this program may be easily accessed at various locations throughout the laser facility by individuals with access privileges, and is particularly useful during activation and maintenance. Each menu displays device status in real time and is keypad driven. A color monitor attached to the terminal displays a power-supply histogram or flashlamp-current waveform plot.

Several utility programs were developed this year that allow various parameter and configuration files to be loaded and saved. Basic communication between the power conditioning system and the shot scheduler was also established.

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Oscillator Controls. The Nova oscillator controls front-end processor (FEP) was moved from the Nova MOR to the system interconnect room with other power conditioning FEPs. This allowed us to easily implement communications to the Nova VAXs through the multipoint memory. This is a change from the planned Novanet

connection. We decided to connect to the multiport to minimize software development and to meet time constraints.

The connection to the multiport memory did require additional software modules that transfer status and commands to and from the VAX. This lets us view and issue commands to the fast timing channels from a video terminal in the control room using the power conditioning AUXPC program.

A touch-sensitive terminal is used to operate Nova oscillator controls. The available menus and their functions are:

1. MAIN MENU—Displays available menus.
2. FAST MENU—Controls the 128 fast timing channels.
3. SLOW MENU—Controls the 12 slow timing channels.
4. CONF MENU—Controls six local fast timing configurations.
5. DISP MENU—Displays status of the 128 fast timing channels.

The Fast Menu (Fig. 2-21) allows the operator to select a fast timing channel and alter its delay time, repetition rate, and enable status. This menu also displays the system synchronization source. When Local Sync is displayed, all timing is generated internally. In Remote Sync, timing is received from a MOR power conditioning control system interface. Synchronization selection is accomplished with a toggle switch mounted on a printed circuit card inside the R-Bus interface chassis. Another feature on the Fast Menu is a local shot button, which will fire the fast timing channels that are enabled at the shot rate. This only has effect in local sync mode.

The Slow Menu allows modification of delay times and enable status of the 12 slow timing channels. This includes the

lamp simmer and peak current pulses to the oscillator lamp driver, and if gating pulses to the acoustic-optic (A-O) modulator and Q-switch driver modules. One of the slow timing channels initiates the fast timing, and another channel provides scope triggers at shot time.

The Configuration Menu allows us to store up to six fast timing configurations on a single magnetic tape. The command to save a configuration copies the current status of all fast timing channels to the tape. When loading a configuration, the operator may choose whether to load the entire configuration (times and rates), or to load just rates. All save and load operations query the operator for confirmation.

The Display Menu is useful in verifying the status on a number of fast timing channels. This menu is divided into 4 pages of 32 channels each. Channel selection is made with buttons at the bottom of the screen. The return button returns the user to the previous menu.

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Beam-Alignment and Diagnostic Systems

Introduction. Three of the four alignment subsystems identified in Fig. 2-22 were completed in 1984; only the harmonic-conversion and target alignment subsystem remains to be activated. The major alignment components were described in earlier Laser Annals⁶⁻⁸ but the steps required to align the Nova system from oscillator to target have never been detailed. We do that here.

Construction of the baseline alignment control system was completed, on schedule, in 1984. In particular, we completed fabrication of the 32 stepping motor controllers (SMCs), which were then used in the remaining activation of the beamline and harmonic-conversion and target sections of the laser alignment system. We have completed two of the three levels of Nova operator controls. Details of the basic system design have been given before,⁹⁻¹¹ so only new capabilities and comments on activation are covered here.

Most of the Nova laser diagnostic sensor packages and data acquisition hardware¹²⁻¹⁴ are installed and operating. Two exceptions

Fig. 2-21. Power conditioning Fast Menu allows operator to select and modify a fast timing channel and its attributes.

Fig. 2-22. Nova alignment subsystems (details in Figs. 2-23 and 2-25).

are the streak-camera interface module for the harmonic-output diagnostics spool and the opposed-port target-plane diagnostic. The final designs for these packages are described below.

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Alignment Systems

Front End Alignment. Figure 2-23 shows the location of alignment components for the front-end subsystem. Sensor package S1 provides both near-field (NF) and far-field (FF) images of the continuous wave (cw) alignment oscillator or the pulsed oscillators in the MOR. Figure 2-24(a) shows the FF image at S1 including the sensor reticle. Comparisons of the NF image with the insertable sensor cross hair, and of the FF image with the reticle, determine the extent of any centering and pointing errors, which are removed by the mirror pair ahead of the sensor. The beam continues to the pre-amplifier section in the laser bay and is sampled by S2. Again NF and FF images are used to detect centering and pointing errors for correction by the preceding mirror pair.

The beam significantly overfills an aperture located after S2, so its perimeter cannot be seen past that point; the shadow of the S2 cross hair is used to designate the center of the beam for subsequent steps. Sensor package S3 and the associated mirror pair

serve to position and aim the beam through the two-way ($2 \times$) splitter. Figure 2-24(b), a NF image at S3, shows the diffracted silhouette of the S2 cross hair well centered on the S3 cross hair. The aperture between S2 and S3 is placed on beam center. This aperture carries its own insertable cross hair for comparison with the S2 cross hair as viewed from S3.

Between S3 and S4, the beam is centered on two insertable cross hairs while being viewed from S4 (S5 is in the other half of the system). Front-end alignment is complete when the beam is positioned and pointed by S4 for propagation into the five-way ($5 \times$) splitter assembly.

Amplifier Chain Alignment. Amplifier chain alignment components are also shown in Fig. 2-23. A portion of the alignment beam enters each of the 10 chains and is centered and pointed with respect to the input sensor (IS) as already described for the MOR and preamplifier. The overfilled aperture at the input to the chain is also positioned as described above. Since the beam to be propagated through the chain is defined by this aperture, the aperture cross hair is used to identify the beam center for subsequent steps. The mid-chain sensor (MCS) does not provide data for automatic alignment functions, but is used by operators for occasional monitoring of the beam position relative to insertable cross hairs 2 through 5.

Fig. 2-23. Sensors and other alignment components for front-end and amplifier chain alignment.

Fig. 2-24. (a) Pointing image typical of S1, S2, S3, S4, S5, and IS of Fig. 2-23. (b) Centering image typical of sensors listed in (a). (c) Output sensor stage of cross hairs 7 and 8. (d) Pinhole image with focused beam superimposed.

In the second half of the amplifier chain, cross hairs 7 and 8 are the alignment references, and the 31.5-cm turning mirrors in the middle of the chain implement corrections. The silhouette of the chain aperture cross hair is first centered on cross hair 8, and then the direction of the beam is adjusted so that the silhouette of cross hair 7 matches cross hair 8. Figure 2-24(c) shows a beam with these two cross hairs in good alignment. A square is added to cross hair 8 because the vertical part of the cross hair is hidden in the shadow of the split 46-cm amplifier disks. This view is obtained by the output sensor (OS) package shown in Fig. 2-25.

The only remaining task for complete alignment of the front end and amplifier chains is the positioning of the spatial filter pinholes to coincide with the focused beam in each spatial filter. The far-field viewing mode of each alignment sensor can provide an image of the beam in the pinhole plane of all preceding spatial filters. Viewed from the output sensor, however, part of the residual error from the alignment of the beam to cross hairs 7 and 8 appears as an offset from the center of the field of view. The output lens of the last spatial filter is translated to repoint the beam slightly and

remove this offset. Then light from the alignment beam is defocused enough to overfill (i.e., backlight) the pinholes. Defocusing is accomplished by inserting a weak lens in the beam ahead of the spatial filter. In each amplifier chain, for example, a weak positive lens introduced into the beam at the input sensor provides back-lighting for the last six pinholes in the chain. Each pinhole is inserted in turn for viewing. Figure 2-24(d) is an output sensor image of pinhole 4 in beam 4 after alignment. The photograph is a double exposure that also shows the focused beam as it appears before the backlighting lens is inserted.

Output Alignment. Components for aligning the laser output are shown in Fig. 2-25. The beam that exits the final spatial filter travels to the target chamber, where it must arrive centered on the crystal array and propagating in the same direction each time. The first and last 74-cm turning mirrors in each beam accomplish this alignment. The position of the beam is observed by inserting both halves of a full-aperture centering mirror (CM) immediately behind the crystal array to send light back to the output sensor. This mirror is registered with the array and carries a square mask a little

Fig. 2-25. (a) Sensors and other components for output alignment and target alignment.

Fig. 2-26. (a) Image of a beam on its centering mirror looking through the KDP array. (b) Comparison of reflections from the KDP and a corner cube. (c) 2ω alignment beam on the centering mirror viewed from the harmonic module. (d) 2ω beam on reticle viewed with the target plane imager.

larger than the square on cross hair 8. Output centering is accomplished by translating the beam until the two squares are concentric, as shown in Fig. 2-26(a).

The direction of propagation as the beam approaches the target chamber is critical, because the beam must pass through the KDP array at the correct angle for maximum conversion efficiency. Therefore, the exit surface of the array itself is used as the output pointing reference. The stepper motors that tip and tilt the array are driven to their "beam normal" setpoints, motor positions that were declared during system activation with the array surface perpendicular to the nominal beam path. The weak reflection from the array and a simultaneous reflection from an insertable corner cube are observed in the far-field mode of

the output sensor. Beam pointing is adjusted until the two reflections coincide, indicating that the beam as sampled by the corner cube is perpendicular to the array. Figure 2-26(b) shows a typical far-field view of these two reflections. In fact, the array produces a group of reflections, because the nine pieces of KDP are not precisely parallel to each other. The reflection pattern is the same each time, however, so the beam-pointing can be repeated with an accuracy much better than the spread of the array reflections.

Harmonic and Target Alignment. The angles through which the KDP array must be tilted to obtain maximum conversion of the $1.05\text{-}\mu\text{m}$ laser output to $0.53\text{ }\mu\text{m}$ (2ω) or $0.35\text{ }\mu\text{m}$ (3ω) are known to a good approximation from crystal fabrication

specifications verified by off-line measurements. However, to optimize the conversion for the on-line conditions of temperature and beam alignment, each array can be angle-tuned using a 2-cm-diam 1ω probe beam.

The beam used for 2ω tuning originates in the local illuminator (LI), shown in Fig. 2-25. The LI beam is introduced into the center of the main beamline by an insertable mirror at the input of the last spatial filter, and is aligned to match the propagation direction of the main beam as viewed from the output sensor. The KDP array is angle-tuned to maximize the 2ω power generated by this probe beam in the center crystal of the array. The signal, which is measured by an insertable diode in the diagnostics spool immediately ahead of the target chamber focusing lens, is normalized to the 2ω power generated by a separate crystal in the local illuminator.

Tuning for generation of 3ω light requires a more intense 1ω probe. This is obtained by diverting the short-pulse output from the MOR along a path that bypasses the preamplifier and amplifier sections and goes instead to the local illuminator. From there the propagation path and the procedure are the same as for 2ω tuning, except that a different detector is used (see the end of the article "Diagnostic Systems," later in this section).

After the preceding steps are completed, the laser system is ready to produce 1ω , 2ω , or 3ω light, but to be useful, the beams must be properly directed onto a target. For beam alignment we use a reticle (a thin piece of glass with a reticle pattern on one diffusing surface) as a surrogate target. The reticle is positioned in the plane at which focused beam positions have been specified for the next shot. The references for determining target position near chamber center are the target alignment viewers (TAV). They are used in conjunction with target alignment illuminators (not shown in Fig. 2-25) and the target positioner (TP) to locate surrogate or shot targets.

If the target is to be irradiated with 1ω light, the system 1ω alignment beams need only be positioned on the reticle before the surrogate target is replaced by the shot target and the system is fired.

For a harmonic shot, a full-aperture harmonic alignment beam is needed. A

wavelength-independent pointing reference (PR), consisting of a Cassegrainian telescope and a charge-injection device (CID) camera, samples the beam, and the reference is adjusted to bring the 1ω far-field spot to the image center. Then the 1ω beam is turned off, and a harmonic alignment beam from the harmonic module (HM) is introduced near the center of the final spatial filter in the chain. Motor-driven mirrors external to the chain center the harmonic beam, identified by the wide cross hair in Fig. 2-26(c), on the centering mirror square. The harmonic module collects light reflected by the centering mirror to form this image. The same motorized mirrors also point the beam to the center of the pointing reference image.

Now the full-aperture harmonic alignment beam is concentric and collinear with the 1ω beam path and will accurately represent the pulsed harmonic light for positioning on the target. The instrument used for viewing alignment beams of any wavelength near the center of the chamber is the target plane imager (TPI). A reticle illuminator (RI) provides background light for imaging the surrogate target reticle with the TPI. Some of the alignment beam light focused onto the reticle is scattered into the $f/3$ collection optics of the TPI, forming an image of the beam on the reticle as in Fig. 2-26(d). The target focusing lens is translated to adjust the position and size of each beam in the reticle plane to match the specifications for the target shot.

Finally, the shot target is put in place with the TAVs as references, and the laser is fired.

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Alignment Controls

Device Control. The capabilities of the SMC were extended to accommodate higher performance requirements for some target-room components. Figure 2-27 shows the interface options available in our final design. While small motors can control 95% of the devices in the system, heavy target-room components such as the target positioner and target-plane imager require

Fig. 2-27. Stepper motor controller interfaces.

low-speed, high-torque stepping motors. The final focusing lens not only requires high torques but must also be adjusted quickly over long distances to permit multi-wavelength operation. The required stepping motors are attached using boosting amplifiers in series with ordinary drive channels. The resulting high-performance channels are interfaced to sophisticated out-board equipment. The channel capabilities are assigned in the software to achieve minimum overall cost while satisfying system requirements. Installation of the full complement of SMCs represents approximately two-thirds of the overall hardware installation and completes the first of three levels of operator controls provided by this system.

The laser diagnostic section of the lens spool contains several sensors, which are alternately placed at coincident locations in the center of the beamline. Similarly, alignment devices inserted into the target chamber center must be protected from collision. We designed an auxiliary interlock package, which is attached in series with the control channels, to arbitrate potential interference and properly cycle device motions to prevent damage. We used a hardware approach to ensure fail-safe operation. Truth tables describing valid device states and

motion sequences are programmed into read-only memory (ROM) chips, which are then interpreted by the control logic. This approach permits the hardware to be used in several different configurations by re-programming the ROM.

VAX operator console software for remote manual alignment was completed to include the remaining alignment system elements: mid-chain and output sensors, local illuminators, harmonic modules, KDP arrays, focusing lenses, pointing references, and diagnostic spools. These graphic menu controls, along with the output alignment selection and status displays, complete the second level of operator controls.

Automatic Alignment. Configuration of the laser for routine operations and closed-loop alignment constitute the third and topmost level of operator controls. Whereas the first two levels of control are employed for manual alignment methods, the third-level software aligns all the mirrors and spatial filter pinholes and other motorized components automatically. An integrated status display for the automatic alignment system shows simultaneously the state of all 294 pointing and centering loops in the system. This information is particularly important because all 10 laser chains are aligned in parallel, and particular tasks are not necessarily

completed at the same time on each beam. We have activated 174 of these loops (including mirror gimbals) up to the output sensor, and all 75 spatial filter pinholes.

Development work during 1984 resulted in very reliable second-generation image analysis methods. Cross hair analysis was extended to permit finding the centers of large-aperture cross hairs and of the centering mirror. Analysis of round pinholes was changed from the less reliable centroid method used on Novette to a more robust method based on correlating the pinhole perimeter to a reference ring. The diameter of the ring is scanned through the expected value of the pinhole diameter to ensure optimum correlation sharpness and accuracy. This method is insensitive to intensity variations and to imperfections in the pinhole.

The final spatial filter is a large rectangular pinhole. The sensor magnification used to view this pinhole does not allow the entire rectangle to be visible at once. In addition, the split in the 46-cm amplifier disks causes diffraction lines on the backlighting beam. We developed a method using mathematical "compass filters," which recognize positive or negative intensity gradients perpendicular to a chosen axis direction, to detect two diagonally opposite corners of the pinhole while discriminating against the diffraction lines. Each corner is detected individually by finding the two edges that form it. The center of the pinhole is then defined as the point halfway between the corners.

During 1984 we began to use the automatic alignment system to measure the response of the laser to various perturbations and to identify marginal hardware items before they could affect laser operations. In June, for example, we ran 39 closed loops continuously for 64 h over a weekend. The variation in outside temperature from day to night exceeded 35°F, and a weak 24-h oscillation was seen in some alignment loops even though the facility's temperature control system functioned normally. One alignment loop attempted to make a large monotonic correction during this test. This was caused by an unreliable motor drive circuit that might otherwise have gone undetected for some time.

Work on the automatic alignment system was directed toward functions that would help maintain the laser system construction

schedule. Continued development in 1985 will concentrate on the remaining harmonic and target alignment tasks that will do the most to reduce shot-preparation time.

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Diagnostic Systems

Opposed-Port Target-Plane Diagnostic.

Knowledge of the intensity distribution of the Nova beams in the target plane is important in interpreting target performance. The opposed-port diagnostic (OPD) will provide this data. The design criteria for this diagnostic require it to do the following:

- Operate with full-power 1 ω , 2 ω , or 3 ω pulses.
- Record an image of the target plane with a resolution of 20 μ m and a field of view of 2 mm.
- Add no intensity-dependent distortion to the image.
- Simultaneously record four target equivalent planes with multiple exposures of each plane.
- Record the energy of the pulse.
- Record the temporal pulse shape.
- Cause minimal interference with other target chamber diagnostics.
- Be insertable by the polar handler (described below).

To meet these criteria we use the inverse Cassegrainian telescope shown in Fig. 2-28. The telescope, selected as the baseline design in 1983,¹² uses the Fresnel reflections from the uncoated primary and secondary mirrors while dumping most of the energy into absorbing glass behind the mirrors. At the exit of the telescope section, the beam is collimated at a diameter of 76 mm by the fused silica tertiary, whose position is adjusted to suit the wavelength of the shot. The intensity at the lens is approximately 0.5 GW/cm² for a 15-kJ, 1-ns laser pulse, so intensity-dependent index changes in this lens and subsequent optics are negligible.

The collimated light passes from the telescope section through a vacuum window into the camera module. For each laser pulse, a multiple image camera (MIC) records 16 images—4 exposure levels in

each of 4 planes, which range from target-chamber focus to 1.5 mm past best focus in 0.5-mm steps. The module also collects pulse-energy and temporal pulse-shape data. The major components of the camera module are identified in Fig. 2-29.

Earlier MIC designs used two orthogonal, partially reflecting wedges to supply the array of photographic images. With this design, one can photograph only multiple planes ahead of best focus, since the light

energy drops rapidly with each additional reflection. The drop in energy is compensated by the decreasing area of the imaged focal spot as the light approaches best focus. (Operation of the earlier MIC is explained more fully in Ref. 15.)

As in previous MICs, the MIC for the opposed pcr diagnostic uses a wedged mirror with a 50% front surface reflection to angularly separate the exposure levels. However, the new MIC overcomes earlier image-plane

Fig. 2-28. Opposed-port target-plane diagnostic package.
(a) Schematic;
(b) photograph.

Fig. 2-29. Opposed-port target-plane diagnostic camera module.

restrictions by using an array of beam splitters to supply the multiple plane images. Since the reflectivities of the beam splitters can be controlled independently, it is possible to have increasing values of reflected energy to photograph planes past best focus (where the focal spot diameter is increasing). The first three beam splitters are wedges with uncoated first surfaces and antireflection coated second surfaces. The materials are fused silica and Schott BK7 and SF8 glasses, giving reflectivities of 3.5, 4.4, and 7.4% respectively at 2ω . The final reflector has a multilayer dielectric coating with about 12% reflectivity. The unwanted transmitted beam is absorbed by an appropriate colored-glass beam dump. Neutral-density filters are inserted just ahead of the film plane to control exposure levels over a wide range of laser pulse energies.

Diverse requirements constrained the mechanical design of the OPD. The space available inside the target chamber is limited by other reentrant target diagnostics and their lines of sight. For this reason, the OPD cone had to be as small as possible. Outside the target chamber, the space available for installation is restricted by the beam tubes, so the optics for the camera module were arranged in a plane normal to the cone axis to minimize their length.

Access to the lower mounting positions is restricted by the overhanging final-focus

optics assemblies, which are cantilevered off the target chamber. For this reason, the OPD was designed to be moved from one opposed port to another by the polar handler, which is also used for moving the final focus optics assemblies (see "Nova Target Systems," later in this section). The weight-handling limit of the polar handler requires that the 1800-lb OPD be divided into two sections that can be transported separately from port to port. The separation point is identified in Fig. 2-28(a). Finally, the envelope of space within which the polar handler must operate restricts its stroke to less than the length of the telescope section of the OPD. Therefore, the forward cone of the OPD can be retracted 50 cm from its deployed position for insertion. The conical section is supported on spring mounts when retracted and is located by kinematic mounts to restore its aligned orientation when extended.

To conserve space and weight, we designed a primary lens bezel that weighs only half as much as the weight it supports and that extends less than 12 cm radially beyond the 74-cm-diam opti. Pyramidal surface bumps on the cast-elastomer mounting pads provide compliance for holding the 450-lb optic with less than one-quarter wave of distortion.

The OPD package must be carefully positioned and focused to image the desired region in the target chamber. For this purpose, there are two alignment aids at the front end. An insertable 1-mm ball, which can be viewed using the target alignment viewers,¹⁴ permits location of the OPD with respect to the center of the target chamber. An insertable reticle at the focal plane of the OPD, 5 in. from the locating ball, serves as a focusing reference. Both these alignment aids are deployed using ac synchronous motors located outside the vacuum, between the telescope section and the camera module.

Two additional reticles are used to monitor the alignment of the primary and secondary mirrors. One reticle is inserted by an ac motor, as just described, and registers to the center of the secondary. The other resides permanently in the center of the primary. An autocollimator in the camera module that looks back along the telescope axis (see Fig. 2-29) is used to monitor reflections from the two reticles. Both the primary and secondary mirror angles can be

adjusted by driving appropriate motors. Active realignment capability is provided, instead of relying on preservation of a fixed initial alignment, for three reasons. First, the static deflection caused by gravity varies significantly from port to port. Second, the separation of the telescope section into halves for transfer between ports results in some loss of alignment. Third, minor changes in alignment occur when the cone is retracted and reextended. The active system assures that proper performance of the unit can be restored after handling.

$2\omega/3\omega$ Incident Diagnostic Spool. The diagnostic spools described last year¹³ have been built and installed on Nova. These spools measure the following quantities:

- $2\omega/3\omega$ incident energy.
- 2ω signals for KDP array alignment.
- Nova focus-lens and debris-shield transmission data.
- Lens reflectivity.

Three additional kinds of data are to be measured with a module mounted externally to the diagnostic spool:

- Unconverted $1\omega/2\omega$ energy.
- $2\omega/3\omega$ temporal pulse shape.
- 3ω signals for KDP array alignment.

The module receives its signals from a 15-cm-diam pickoff mirror in the diagnostic spool. The mirror directs the central portion of the beam into the module, as shown in Fig. 2-30. When the module is configured for a shot, the sampled part of the beam

strikes a diffuser that scatters light onto photodiodes filtered to provide a measurement of 1ω , 2ω , or 3ω energy.

Part of the scattered light also illuminates the end of a fiber optic cable that transports the signal to a streak camera for temporal measurement of the incident pulse. The fibers from five modules are multiplexed onto a single streak camera; to these signals are added an optical time-reference fiducial and a comb of pulses, equally spaced in time, to provide a camera sweep rate calibration. Figure 2-31 shows the output from a multiplexed streak camera. The signal level on each fiber is adjusted by rotating a filter wheel in front of it. For very-low-level shots, such as those used for timing the streak camera, the diffuser can be removed so that the fiber is illuminated directly.

To tune the KDP array for 3ω output, a small-diameter (~ 2 cm) pulsed probe beam is routed from the oscillator along a path that bypasses the laser amplifier chain; the beam is reinjected at the input of the final spatial filter. The beam passes through the frequency conversion crystals, which convert some of it to 3ω . The diagnostic spool 3ω pickoff mirror directs this partially converted probe beam into the module. A wedge, inserted in the optical path in place of the diffuser, deflects the beam onto a 3ω photodiode. As the crystals are tipped through their tuning range, the photodiode signal is recorded by the laser diagnostic

Fig. 2-30. A module attached to each incident diagnostic spool provides multiple wavelength energy measurements and a streak-camera interface.

Fig. 7-11. Signals recorded on a streak camera multiplexed to five beams, including the comb of pulses used to calibrate the streak speed.

data acquisition system. The peak conversion position is determined by fitting a tuning curve to the data, and the crystal array is set to this angle.

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Nova Target Systems

Nova target systems include the target chamber with its support and vacuum systems, the target positioner, target alignment optics, final calorimeter supports, and handling equipment for lens positioners, KDP arrays, beam diagnostics, calorimeters, diagnostic controls, and data acquisition systems associated with the target chamber. During 1984, we completed the installation of all system components. Initial installation was made with temporary, portable fixtures.

Target Chamber Installation. The target chamber was installed as described in the *Laser Program Annual Report* £3.¹⁶ Following installation, some modifications were made to the spaceframe; we also experienced a mild earthquake. Either could have caused misalignment of the chamber, but a re-survey showed no detectable change.

Vacuum System. The installation of both vacuum systems, main and auxiliary, was

completed. The pumping speeds achieved are better than anticipated, reaching 10^{-5} Torr in 25 min and a base pressure of 2×10^{-7} Torr after 48 h. These results assure our ability to achieve 10^{-6} Torr on an operational basis for cryogenic target experiments. Some condensation occurs during the evacuation of this large vacuum chamber when moist air in the chamber is cooled to its dew point. The cooling is a result of the adiabatic expansion of the chamber air while the vessel is being pumped. The condensation problem will be solved by the addition of a desiccant dehumidifier in the duct that supplies fresh air to the chamber.

Target Positioner. The target positioner *x-y* slides were calibrated prior to assembly. They showed excellent linearity and orthogonality with imperceptible backlash, and the assembly was completed with very few modifications. The assembled unit was installed on the target chamber and in situ calibration was started. We will check linearity, orthogonality, backlash, and resonance.

Handling Equipment. The installation and servicing of the massive components used in the frequency conversion, alignment, beam diagnostics, and focusing systems require highly specialized, massive handling equipment. A unit was built to accept any component from the overhead bridge crane, rotate it to align with a specific beamline, and advance it radially for permanent attachment. The maximum rated load of this polar handler is 680 kg; it will be used to install and service the KDP spools, main focusing lenses, debris shields, 44-cm calorimeters, and the opposed port diagnostic package.

Debris shields are serviced with a special attachment that replaces the swing arm on the handling fixture. Use of the swing arm requires that beam pipes near the target chamber be removed; this attachment allows us to service debris shields with the beam pipes in place.

Lens positioners weigh 2000 kg each and must be handled separately. The fixture we designed to perform the original installation of the lens positioners imposed some constraints on the sequence of any future installation. We do not envision a need to remove any lens positioner during the operational life of the system, but the removal of any lower unit would require us to remove an upper one first.

Fig. 2-32. Nova target chamber alignment pylon: (a) extended and (b) retracted.

Fig. 2-33. Nova target chamber calorimeter configured to filter out unwanted 2ω light.

Alignment Components. An alignment pylon [Fig. 2-32(a) and 2-32(b)] was designed and fabricated to verify the exact location of the center of the target chamber. This pylon precisely and repeatably places a spherical target at the chamber center. From this reference, the target alignment and viewing instruments (TAVIs) and all diagnostics are aligned. The alignment aid positioner, a unit identical to the target positioner, will be built to support the various components that are necessary to align the laser beams to the center of the target chamber. It will be installed in the bottom port of the chamber.

Calorimeter Mounting. To calibrate the sensors in the diagnostic spool (described in the preceding article "Beam-Alignment and Diagnostic Systems"), we need an absolute measurement of the 2ω and 3ω light reaching the target. This measurement must be taken inside the target chamber. To accomplish this, we mounted a 44-cm-aperture calorimeter (Fig. 2-33) inside the chamber for each beam. The calorimeters have

insertable filters to separate out unwanted wavelengths (1ω or 2ω), an uncoated, concave reflector that returns 3% of the beam to the output sensors during calibration for reflected diagnostic calibrations, and a blast shield to protect the unit from target debris, x rays, and scattered light during a target shot. Each filter and the blast shield can be manually inserted from outside the chamber while it is under vacuum. Microswitches provide status and interlock capability.

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Nova Target Diagnostics Control System

During the past year the Nova target diagnostics control system was finished and put in service. The diagnostics loft constructed to the north of the target room provides the environmental conditions required to collect reliable target diagnostic data. These improvements include equipment cooling and isolation of the power source with strict control of instrumentation grounds to eliminate data corruption due to electromagnetic pulses from the laser power-conditioning system or from target implosion effects.

After the loft area was completed and Novette was shut down, the instrumentation from Novette was moved to Nova. The electronic instruments for the filter-fluorescer x-ray spectrometer (FFLEX) and for half of the Dante soft x-ray spectrometer diagnostic were installed in the Nova switchyard area to allow short coaxial cable connections from these diagnostics to the digitizing equipment on the south side of the target chamber. All other instrumentation that is not chamber-mounted was installed in the loft; this allows short connections to all diagnostics on the north side of the chamber. The short connections maximize signal bandwidth and fidelity and keep noise to a minimum.

Data Acquisition. We installed six front-end processors (FEP), a number of computer-assisted monitoring and control (CAMAC) crates, and the Novanet fiber-optic, digital data network to the loft area during the past year. We also installed an operator console consisting

of a single Ramtek screen and touch panel, which are slaved to one of the consoles in the Nova control room (Fig. 2-34). All equipment in the loft is fiber-optically coupled to the rest of the laser system to ensure that signals recorded in the loft are not contaminated by electrical noise from other sources. Approximately 20 fiber-optic cables carry digital and pulse information to and from the loft area. These signals range from optical trigger signals, and Novanet 10 million bits/s digital data to the lower-rate, EIA RS-232 standard signals that support computer terminals and the operator touch panels.

Major software tasks completed in the past year to support the data acquisition system include connection of the target diagnostic system to the power conditioning system and the integration software system. The integration software supplies all other Nova control systems with information about the state of the laser, and it sequences major tasks to allow a logical progression through a laser system shot. At various times prior to the shot, checks are made by all systems to ensure that correct operations have been performed in the proper sequence. The target diagnostic system is fully software-interlocked to halt the shot sequence if critical diagnostics fail to pass checks or if appropriate trigger signals are not received during the countdown to a target shot.

Software was completed to support *diagnostic-to-hardware mapping*. This allows the system operator to deal only with the set of hardware that supports a selected diagnostic. Additionally, the data analysis codes can easily extract just the data pertaining to a single diagnostic. The hardware mapping software therefore increases operator efficiency and accelerates the computer analysis and reduction of shot results.

Trigger System. We completed the design of the target diagnostics timing system and installed all the associated hardware and software. This system provides isolated electrical triggers to all target diagnostic instruments. It also provides timing fiducials that are summed with many diagnostic signals to allow cross-timing between channels to within 50 ps. This system receives a fraction of the laser pulse from the master oscillator room. This fraction of the MOR pulse is transmitted to the diagnostics loft via fiber optics, where the trigger system

Nova Target Diagnostics Control System

converts the optical pulse to an electrical equivalent, compensates for amplitude variations, monitors and records timing drifts, and distributes the triggers to four subsystems (Fig. 2-34). First the instrument trigger modules fire the time bases of digitizers, oscilloscopes, and standard recording equipment with a maximum jitter of 200 ps. Then the second subsystem provides a separate streak camera trigger source that is mandated by the 50-ps jitter requirements in ultrafast streak cameras. Third, the instrument fiducials are generated for inclusion with the digitizer and oscilloscope data to allow precise cross-timing of diagnostic data. Finally, a signal simulator is used to allow accurate timing of diagnostics prior to the target shot. This is done, without using the main laser, by precisely generating and timing a simulated diagnostic signal, and applying it to the signal cables that will be used on the actual target shot.

Software was also written to support the simulator electronics and a portable handheld terminal that allows the diagnostic

technician to adjust and control the instrumentation while moving anywhere in the diagnostics room. During the pre-shot setup sequence and during the target shot, the triggering and timing system monitors critical pulse amplitudes and delays, and alerts the operator to improper conditions. At the same time, this system captures the actual trigger parameters necessary for later data analysis.

Vacuum Controls. We have installed the basic set of hardware required to support the Nova vacuum control system: 7 LSI-11/23 processors [1 FEP, 6 input/output processors (IOPs)] and approximately 30 local control panels. The vacuum control system is similar to our other FEP-based systems except that it emphasizes real-time control and fast response to protect high-cost diagnostic equipment.

The typical diagnostic instrument and the main target chamber are connected to the vacuum control system through the IOP units as shown in Fig. 2-34. The IOP functions are to supply 24-V dc to operate the

Fig. 2-34. Block diagram with major components of the target diagnostics vacuum control, timing, and data acquisition subsystems.

valve solenoids, to sense the valve state switches used to obtain valve position, and to read, linearize, and scale the various vacuum gauges. The IOP also controls an array of compressors that are used to pump the chamber and the diagnostic instruments. *High vacuum is provided by cryogenic cold head pumps; the vacuum system controls and monitors both the cold heads and the gaseous helium compressors used to maintain a sufficiently low temperature for proper cryogenic pumping.*

Configuration and interlock information is maintained in a software data base in the VAX and downloaded to the main FEP and all IOPs prior to startup. The control panels are attached to the main FEP and are positioned where convenient in the diagnostics loft or in the target room. These local control panels are identical units except for a plastic overlay that customizes the panel to the vacuum map of the chamber or to a particular diagnostic. Communication between the panels and the main FEP is via hard wire or fiber-optic connections, using an RS-232 serial protocol. Communication between the VAX, FEPs, and IOPs is via Novanet at 10 million bits/s. When the main target chamber and the individual diagnostics are operating, status is displayed and operator commands can be entered through the local control panels.

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Mechanical Systems

Introduction. During 1984, we completed the Nova construction project. Our milestone commitment, a full energy shot at 3 ω , was accomplished on December 19, 1984. To make this shot, we had to assemble, install, and activate approximately 500 major components, and integrate them into a whole facility of 10 beams. Eight of the Nova beams were completely installed prior to the Novette shutdown on August 31, 1984, except for the addition of one 46-cm amplifier and one 31.5-cm amplifier on each beam. The last two beams were installed after Novette was dismantled. The components

from Novette were inspected, rebuilt if necessary, and installed on Nova.

All of this equipment was processed through clean rooms: amplifiers and the more sensitive components in a class-100 clean room,¹⁷ and the remaining components such as rotators and spatial filters in a class-10 000 clean room. Cradles and support brackets were mounted on the spaceframes and accurately surveyed to accept the components. Beam pipes were installed last, during beam alignment. A radial clearance of 2.5 cm in the tubes was allowed in the design to clear the 74-cm beam; this proved adequate to allow full aperture clearance.

The recirculating nitrogen gas system used for cooling the laser was operated in an open loop mode while completed laser chains were being added. *When all 10 laser chains were complete, the gas system was converted from open to closed mode and is now successfully operating in the recirculating mode described in the Laser Program Annual Report 83.*

The control system used to assure quality parts and timely supply was sharply tested throughout the year. Even though we had extensive experience with the Novette project, handling large numbers of critical parts resulted in unexpected problems. Examples of problems that were not apparent until parts were needed included deterioration damage, loss in storage, and poor quality of silver-plated reflective surfaces. Responding to these problems required extraordinarily quick action to meet our strict schedule requirements.

During the year, mechanical technician manpower grew, but the design team manpower decreased slightly, since design activity consisted primarily of follow-through on the construction effort. We also supported late designs in other subsystems, such as the opposed port diagnostic described earlier in this section.

Assembly of Components. The Shiva laser bay was used as the class-10 000 clean room for the assembly and interferometry testing of the Faraday rotators, Pockels cells, alignment cross hairs, turning mirrors, beam shutters, cradles, and spatial filters. All other components, amplifiers, and interstage hardware were assembled in the Nova class-100 clean room. A total of 165 amplifiers were assembled. The critical

components were the disk amplifiers. The largest of these, the 46-cm amplifier, was the most difficult to assemble because the split disk is very sensitive to stress birefringence. Some modest changes to the disk mounting hardware, combined with major changes to assembly procedures, had to be implemented.

The 46-cm amplifier has an inherent stress level associated with the split disks. Stress concentrations induce local changes in the index of refraction. To reduce wavefront distortion, we altered the method of assembling the disk halves. We selected disk halves by matching thickness and wedge angle, and modified the disk holder and the assembly fixture. These changes resulted in a disk with a surface figure of $\lambda/2$. The assembly procedure was to first place the disk holder horizontally in the fixture, then align the disk halves above the holder, and finally, to raise the holder to fit around the disks. The disks have radial support points, which are tuned to optimum stress levels. The tips of these supports, which originally were metal, have been converted to silicone for better friction and load distribution. The operational effectiveness of these corrective actions is not yet known.

It takes about three days to properly assemble a 46-cm amplifier. This is the longest assembly time required for any of the six amplifier sizes used in a Nova chain. Since amplifier assembly was planned to coincide with the phased installation of laser beams, it was necessary to assemble several different sizes of amplifiers simultaneously. Each component had to be validated for cleanliness and optical integrity before integration into the system. The last amplifier, a 46-cm unit, was photographed (Fig. 2-35) on November 30, 1984, as it was being lifted into place for installation on the spaceframe.

Amplifiers were mechanically aligned using standard optical techniques with respect to established bench marks and beam paths. The final alignment was made with the alignment laser. Power cables and nitrogen cooling lines were connected, and then the interstage hardware was inserted.

Twenty-eight amplifiers from the Novette laser were the last to be processed. All these amplifiers were disassembled, cleaned, reassembled, and validated. The large reflective

silver-plated areas are at risk in the amplifiers. However, examination of the disassembled Novette amplifiers showed that the silver-plated surfaces had withstood two years of operation and approximately 450 shots very successfully. Very little tarnish, and the consequent degradation of reflectivity, had occurred during this time.

This assembly was a task of long duration—it was a constant, dedicated effort to quality and is the first quality assurance effort of such major proportions in the history of the laser program at LLNL and probably in the history of large optical systems throughout the world.

Spaceframe Stability. After four years of design, analysis, construction, and measurement, the system spaceframe was put to use. Only calculations and subsequent measurements, not yet known, led us to believe that the system would be stable enough over the required time span to perform meaningful experiments. We have

Fig. 2-35. Technicians removing the protective clean covering immediately before installation of the last 46-cm amplifier on the top beam of the south laser frame—23 ft above the floor.

Fig. 2-36. 31.5-cm Faraday rotators and 31.5-cm box amplifiers (right) and 31.5-cm spatial filters on the south laser frame (left foreground).

observed stable conditions for alignment purposes of up to 12 hours, and in one case, over an entire weekend. In April 1984, we experienced an earthquake and the system performed as expected. Measurements indicated a 0.02-g equivalent acceleration at LLNL. There was one fully assembled and aligned beam on the frame at the time. It moved with the earthquake and returned perfectly to its original alignment.

The total path length from the master oscillator room to the target is 280 m. Components are supported on six separate spaceframes in four rooms with independent temperature control systems. We have observed no thermal drift or vibrational instabilities in any part of this system. We will continue to investigate spaceframe stability to finer tolerances.

Installation of Components. We developed a procedure for installing the beamline components and mirrors that ensured

the proper mechanical alignment. The criteria for mechanical alignment were specified in terms of distance from the beam centerline: ± 1.5 mm for installation of component cradles, ± 0.5 mm for beamline components, and ± 1 cm for turning mirrors. Before installing the beamlines, we first surveyed the relationship of the spaceframes to each other and to the building. This information was used to position beamlines on the spaceframe. A fixture was constructed that allowed an entire vertical row of cradle-mounting brackets to be installed and aligned at one time. Once the mounting plates were in place, the cradles were installed and a cradle cross hair was aligned using a transit set at beamline center.

The components on the driver side of the laser were installed in sequence. First, we installed the permanent cross hairs and the assemblies that off-set the beam, second the spatial filters, and then the remaining components. This sequential installation allowed us to perform the optical alignment in parallel with the mechanical installation. A similar procedure was followed during the power-side component installation. Figure 2-36 shows the center aisle of the laser room looking toward the switchyard, with driver-side components on the left and power-side components on the right.

We verified the proper position of the center of the target chamber by referencing it to the master survey monument in the switchyard. This permitted proper placement of the chamber with respect to the 10 beamlines. Figure 2-37 shows the target chamber from above; Fig. 2-38 shows the switchyard, where beams are turned 90° from the laser into the target room. A helium-neon alignment laser was positioned on the far side of the target chamber, and the beam directed back through the target chamber and laser path. Surrogate alignment mirrors (5-cm diam) were used to verify the turning mirror locations. Turning mirrors were positioned within ± 1 cm in the target and switchyard rooms. After installation and mechanical alignment within this tolerance, the mirrors were finally positioned using the system cw alignment laser.

By using a systems approach for installation and alignment, we ensured that the laser

Fig. 2-37. Target chamber with upper beamlines installed. Turning-mirror box locations, designed into frame construction, were verified by shining an alignment laser through the center of the target chamber.

Fig. 2-38. Turning-mirror towers with diagnostics mounted outside (foreground) contain two main beam mirrors and one diagnostics mirror per beam; note remotely insertable calorimeters (left, target-room wall).

beam will transit the entire 280-m beam path to the target chamber without interference.

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Integrated Controls

Integrated controls consist of the overall design, hardware, and software that the four major Nova control subsystems—power conditioning, alignment, laser diagnostics, and target diagnostics—utilize to perform and coordinate their tasks. The accomplishments within each of these control subsystems are described earlier in this section in the articles “Power Systems Energy Storage,” “Beam-Alignment and Diagnostic Systems,” “Nova Target Systems,” and “Nova Target Diagnostics Control System,” which includes a report on the increasingly important vacuum control components.

During 1984, all portions of the Nova control system that were necessary for the support of laser activation and completion of the Nova project were finished and placed in service on time. The Nova control system has been unique in providing, on schedule, the capabilities required in the central control room and in various local control areas throughout the facility. The ambitious goal of deploying this system early enough to use it as an aid in the activation of the laser¹⁸ was accomplished; thus the control system made a major contribution to the completion of Nova activation on schedule.

Support and enhancement activities continued during the year on the VAX computer systems, central control room, operator consoles and displays, Novanet data communications network, system-level software for both the VAX and LSI-11 computers, Praxis control system computer language, software management tools, and the development system, which includes office terminals. Computational support was also supplied for a wide variety of test fixtures required by the optical and mechanical subsystems.

Significant new advancements were made in four areas in integrated controls this year: the integration software (which

includes the shot scheduler), the Praxis language, software quality assurance audit, and software development and data handling. A description of the accomplishments in each of these areas follows.

Integration Software. The first phase of the integration software was completed and became operational. The integration software allows the four control subsystems to operate in a coordinated fashion. It includes the shot scheduler and the data shared by the subsystems. In previous years, operators were required to coordinate the subsystems to operate the laser facility. Now, some of that coordination is performed by the integration software.

The shot scheduler, which is part of the integration software, facilitates the scheduling of a complex array of laser control subsystem tasks and ensures that they are done at the correct times and in the correct sequence. The scheduler software manages the exchange of interdependent information among subsystems and verifies that there are no task conflicts before a subsystem operation is allowed; it replaces a number of manual coordination activities during laser facility operations. For example, if laser component alignment is sequenced through the scheduler, the scheduler first verifies that there are no conflicts with operations running in any of the three other subsystems before passing the alignment tasks to the alignment subsystem for performance.

In a complex, ever-changing, experimental environment such as Nova, new experiments impose new control system requirements. As operating experience increases, new task dependencies are understood. The scheduler is flexible enough to easily accommodate new control system requirements and to update task dependency information for smoother operations.

The scheduler provides a means of executing control system tasks either in a desired order or concurrently. Executing control system tasks in parallel reduces total task execution time, thereby increasing control system throughput. Coordinating control system tasks by the scheduler, instead of by operators, provides uniform, reproducible conditions, and reduces operator requirements and errors.

Some control system tasks implemented by the subsystems are now being

coordinated by the integration software rather than directly by the operators. For example, the power conditioning subsystem will not fire the laser until the target diagnostics subsystem is ready for the target shot. The countdown can begin only after the target diagnostics equipment is ready. If the target diagnostics subsystem detects a problem during the countdown, the countdown will be suspended by the power conditioning subsystem and the nature of the problem will be indicated to the operator. The operators still have the capability of overriding the operational interlocks in the control system to fire the laser.

The integration software is ready for additional coordination of control system tasks as we gain experience with operation of the laser facility.

Praxis Enhancements. Substantial additions and enhancements were made to the Praxis computer language to increase software capabilities on the Nova control system. Praxis is a modern, block-structured language,¹⁹ originally designed and written for Nova by Bolt, Beranek and Newman. The Praxis language is used almost exclusively within all the computers of the Nova control system, and its use has been a major contributing factor in the success of the system.

Many of the improvements²⁰ in Praxis add support for previously unimplemented features of the language; now statements that may be more effective in a particular segment of code can be utilized. Additional checks were added in the language compiler to verify that the code being compiled is logical and that its execution is clearly stated in the source program. Some of these checks are performed when the code is compiled, and other checks are performed during execution. Some of the improvements allow more work to be done in less computer memory.

The improvements made to Praxis were essential for the features required of the vacuum control system (described in "Nova Target Systems," earlier in this section), and they are important for optimization, reliability, and additions to all other parts of the Nova control system.

Software Quality Assurance Audit. In previous years, a useful software quality assurance plan had been developed.²¹ The

technical content of this plan is the primary guide for all software development on the Nova control system. Through a continuous, interactive review of the plan, and of the extent to which the software meets the plan, we have been able to successfully manage the entire scope of the software effort for the Nova control system.

An internal management review of the software for the Nova control system was conducted by the LLNL Quality Assurance Office during May 1984. Through interviews with both the operational staff and the control system engineers, the Nova software was found to be useful and timely, and to meet system requirements. Particular note was made of the especially good conformance to the technical aspects of software design control.

New Development and Data Handling Capabilities. New equipment, added to the Nova central computers, has provided a means for handling large volumes of data flow within the control system and has improved the reliability and rate of development of control system software.

Larger and faster disk storage devices were added to the development computer. Data storage capacity was increased from 1.2 to 4.0 gigabytes with the addition of six disk drives. These drives operate at approximately twice the data rates of the disks they replaced. Improved software maintenance was made possible by the increased storage capacity. As described in the Nova software quality assurance plan, the design for software development was to maintain all software at three levels of testing, allowing a program to be modified at one level while older versions of the program are being tested and used from the other levels. This new data storage equipment makes the long-term software development plan practical.

A fourfold increase in main memory on the development VAX to 16 megabytes has enabled faster development and testing and more parallel activity. The computer now supports approximately 60 concurrently executing programs for many users in the Nova office building and surrounding development laboratories.

DECNET, an intercomputer communication system supplied by Digital Equipment Corporation, was added to all of the VAX

computers in the Nova facility. This network package has provided the means for transferring large amounts of data gathered from the laser facility by any of the three on-line control VAX computers. The data flows back to the development VAX for processing without disturbing the on-line control system that utilizes Novanet. DECNET also provides the principal means by which software development is monitored and managed throughout all of the VAX computers within the facility, and by which partially tested control system programs are transferred to the computers that run the control system. The rate at which new features can be reliably incorporated into the control system is proportional to the number of times per day that programs can be edited, regenerated, and transferred to the appropriate machine for testing. In tests of a large software package, the turnaround for these operations is now three times faster.

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Nova Project Management

Introduction. The major Nova project milestone, "Nova Operable/All Project Construction Complete," was reached on December 19, 1984, with a 17 kJ, 3 ω test shot that delivered 10 beams simultaneously into the Nova target chamber.

The construction completion milestone was accomplished on schedule and about \$1.5 million under the March 1, 1983, baseline construction budget. During 1984, the Nova project continued to operate under this budget except that three major modifications to the plan, involving release of contingency funds, were authorized by DOE. These modifications were:

1. Expenditure of \$3 million for additional target diagnostics.
2. Expenditure of \$2.4 million to extend the construction period from September to December 1984 to allow for an additional three months of x-ray laser experiments utilizing the two arms of Nova installed in the Novette test bed.

3. Expenditure of \$3 million for Nova activation during the period from October 1984 through February 1985.

The baseline financial plan for the total project updated to December 1984 is illustrated in Fig. 2-39. The comparison between project financial status on December 31, 1984, and the baseline budget adopted 22 months earlier is shown in Table 2-6. The Nova laser/target system figure of \$133 878 includes the additional expense of extending the construction period three months to December 1984.

Project Management Controls and Tracking. Nova project management controls continued to provide effective cost and schedule tracking during 1984. Additional cost and schedule aids, described below, were used to ensure project completion on schedule and within budget. The need to draw against contingency funds was reassessed each month so that additional work within project scope could make use of these funds as early as possible.

Contingency funds of \$4.7 million were held in reserve on March 1, 1983. This was 10.1% of the Nova project uncommitted balance at that time, and was considered a prudent level. Later, when projections indicated a project underrun and excess contingency, a detailed procurement and labor plan for the expenditure of an additional \$3 million for target diagnostics was submitted by the project office to DOE and was approved. Likewise, the use of excess contingency funds to extend construction three months to accommodate the Novette x-ray laser experiments was approved. In mid-1984, after projecting additional excess contingency, a detailed Nova activation plan for \$3 million was proposed and submitted to DOE. This provided for training, system integration, target diagnostic checkout, laser alignment, diagnostic checkout, administration, and a reserve for punch list items. In October 1984, this activation plan was approved.

During 1984, the Nova systems engineers utilized detailed schedules for installing each part of each Nova chain from the splitter arrays to the focusing lenses. These installation schedules were used by the Nova engineers in conjunction with task schedules and remaining procurement schedules to install chain items on time.

Novette Laser System

Total project	03/01/83 Baseline	171.300
December 1984	Estimate at completion (EAC)	175.709
Planned % complete 100%	Variance over (-) under (+)	-4.409 ^a
Actual % complete 100%	Remaining contingency	0.291
	Total estimate at completion (TEC)	176.000

Fig. 2-39. Nova project financial and completion schedule.

^aIncludes 5.902 additional; 2.992 target diagnostics and 2.910 activation.

Table 2-6. December 31, 1984, Nova project financial status and March 1, 1983, project baseline budget (in thousands).

The critical path continued to be the delivery of finished KDP crystals for the frequency conversion arrays. The optics group resolved finishing problems by using alternate scheduling methods. Arrangements were made to perform the prefinishing and the final diamond turning of KDP crystals at separate facilities, according to their level of expertise in this advanced technology.

Cost/Schedule Summary. The Nova project construction was completed on schedule and within budget. About \$3 million of Nova project funds is being used for system activation leading to the "Nova Activation Complete/Start of Experimental Program" milestone scheduled for February 28, 1985. This milestone is expected to be completed on schedule within the activation budget.

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	Baseline 3/01/83	Financial status 12/31/84
Nova laser/target system	\$135,310	\$133,878
Nova laboratory/office building	35,990	35,929
Additional target diagnostics	—	2,992
Nova activation	—	2,910
Contingency	4,700	291
	<u>\$176,000</u>	<u>\$176,000</u>

Novette Laser System

Introduction

The last Novette target shot was fired and Novette facility operations were terminated at 6 p.m., August 31, 1984, exactly three years and one day after Argus, the laser that preceded Novette in the same building, was shut down. Unlike Argus, which was dismantled rapidly to make way for

Novette Laser System

Fig. 2-40. History of Novette high-energy laser shots.

Fig. 2-41. Distribution of high-energy laser shots by type of activity over the Novette project lifetime.

Table 2-7. Novette project milestones

Jan. 1 1982 to Aug. 31, 1984	
Beginning of Novette laser operation	Jan. 1, 1982
Highest Novette output power	Jan. 24, 1983
1 chain 12.6 TW 1w, 5.8 TW 2w	
2 chains 24.8 TW 1w, 11.0 TW 2w	
First x-ray laser target	Feb. 1, 1983
First 1-ns target	May 19, 1983
Highest Novette output energy	Oct. 14, 1983
2 chains 17.7 kJ 1w, 8.45 kJ 2w	
4w shot series	Jan.-Apr., 1984
Nova array test series	Feb.-May, 1984
Highest 1 chain 2w output energy 4.6 kJ	Apr. 11, 1984
Highest 1 chain 1w output energy 10.0 kJ	Apr. 24, 1984
Highest 1 chain 3w output energy 2.8 kJ	Apr. 24, 1984
Highest 1 chain 3w output power 2.7 TW	Apr. 24, 1984
Highest 1 chain 4w output 1.6 kJ, 1.6 TW	Apr. 26, 1984
Second x-ray laser target series	May-Aug., 1984
Final shot	Aug. 31, 1984

Novette the two Novette amplifier chains were carefully removed and reinstalled as the final two arms of Nova. Preceding shut-down and disassembly, the year was almost entirely devoted to the target studies, reported in Sec. 5, "Laser Experiments and Advanced Diagnostics." During the 20 months that Novette was used as a target irradiation facility, 674 high-power shots were made: this is an average of 34 shots per month, of which 66% were target shots (Fig. 2-40). The occasional large decreases in target shots per month, marked the beginning of each target campaign when an entirely new target was introduced. Daily shot rates climbed as high as six shots into the target chamber in one day, close to the maximum of eight per day, set by amplifier cooling time. The distribution of Novette high energy shots over the 20-month project lifetime is shown in Fig. 2-41, and a short history of Novette is summarized in Table 2-7.

Novette operation and performance during 1984 are described in the following two articles. The first article emphasizes the 4w experiments in which a unique, tandem KDP array converted the frequency of one

of Novette's beams from 1ω to 4ω and over 1.6 kJ of 4ω ultraviolet light was delivered to targets in <1 ns. The focusing arrangements for the highly successful x-ray laser study done at the Novette facility are described in the second article. Novette was generally acclaimed as the most successful target irradiation laser facility yet constructed at LLNL, an accolade that bodes well for the Nova facility just coming into operation now.

Author: K. R. Manes

Novette Operations and Performance

From January to the end of August 1984, the major Novette activities were frequency conversion tests and 4ω target irradiation experiments. Our efforts to produce 3-ns pulses correctly shaped for high gain and minimal hot electron production in these targets established the principles for the later use of this shaping technique on Nova, and gave us sufficient information to complete the 4ω series with 1-ns pulses.

When the conversion of the Novette facility for x-ray laser experiments was completed at the end of May, staffing levels in Novette operations dropped as the needs of the Nova installation work increased. Even with fewer personnel, the last three months of operation saw us reach the highest rate of successful target shots in the whole Novette project period. In addition to the frequency-conversion testing and target experiments, several subsidiary activities were undertaken or continued through the end of August. These included demonstrating that the copropagating back-lighter configuration could be aligned through the laser system with existing

alignment sensors and that 1-ns pulse propagation at higher output fluences (0.7 J/cm^2) did not result in either pinhole closure or mingling of energy from the copropagating pulses. In August, we performed several polarization tests on Novette for comparison with similar Nova shots and found no significant performance differences between the two systems. The Novette target area was also used as a training and prototyping test bed for the target diagnostic used with the $22\times$ microscope.

Frequency Conversion. Frequency conversion tests were conducted on both chains of Novette but with quite different goals. On the north chain, we tested the first Nova crystal array for conversion efficiency, output beam quality, and durability at both 2ω and 3ω under actual operating conditions. On the south chain, we did wavelength-scaling target experiments at 4ω . We measured the 3ω or 4ω UV energy by propagating an annular beam through an aperture at target chamber center, using the physical separation of the 1ω , 2ω , and UV focal spots arising from the dispersion of the target illumination lenses to block 1ω and 2ω energy while transmitting UV. This technique was designed for Nova,²² but was tested first on Novette (Fig. 2-42).

Both the north and south chain frequency conversion experiments required the installation of fused silica optics in the target chamber because BK-7 solarizes rapidly when exposed to high UV fluences. The frequency conversion tests of the Nova crystal array were the first to incorporate all the modifications made to the original Novette 74-cm, 3×3 array configuration: the quadrature- 2ω -cascade- 3ω configuration, the sol-gel-coated KDP crystals, and the deep-web crystal support structure. In general, conversion efficiency peaked near

Fig. 2-42. Calibration configuration for measuring UV energy in Nova frequency conversion tests conducted on Novette.

2 GW/cm², similar to the single-crystal Novette results reported earlier.²³ The sol-gel coatings on the fused silica optics in the west target chamber had visible discoloration after the coatings had been exposed to 1 kJ of 3 ω light.

Target Irradiation Experiments. To perform the short series of 4 ω wavelength-scaling target shots, we added a second layer of appropriately cut KDP crystals to the south chain 5 $\times$ 5 array in January. It was also necessary to add a completely new set of crystal-tuning and target-energy diagnostics.

Pulses of 1 ω light of sufficient power to cause measurable conversion must be delivered to the array input to tune crystal assemblies. Tuning the 4 ω array needed

higher power than the target local illuminator could deliver, so we installed a new beamline from the MOR along the top of the south laser spaceframe to the output spatial filter.

For shot-time energy measurement, we adopted the baseline Nova technique of placing a reflector in the center of the beam; we used a fused silica prism to separate the wavelengths of the reflected beam, measuring not only the 4 ω energy but also the residual 2 ω energy. As on Nova, the shadow of the reflector in the center of the beam, dispersed by the fused silica target optics, also protected the targets from exposure to 1 ω and 2 ω light. The 4 ω target-energy diagnostic was calibrated by using a calorimeter with a Pyrex absorber to

Fig. 2-43. (a) Typical preamplifier input signal and (b) temporal shape of Novette output pulse; and MALA-PROP simulations of (c) 1 ω preamplifier input and (d) 2 ω output temporal profile for laser output of 7.8 kJ, 1 ω .

measure the transmission through the chamber-center aperture.

The initial 4ω experiments used a 3-ns rectangular pulse of 8 to 10 kJ of 1ω energy. At this fluence level, significant gain saturation in the laser chain (Fig. 2-43) required that a ramped pulse shape be delivered from the MOR. We generated the desired pulse shape with a variable impedance transmission line (VITL) and a silicon switch to drive a Pockels cell (see "Pulse Generation and Shaping" earlier in this section). These experiments were used to prove the principle of this pulse shaping method, now being perfected for use on Nova.

In the prototype form, it was too difficult to achieve the correct output pulse shape within the time constraints of the 4ω

experimental series. Calculated input intensity for 8- and 10-kJ square output pulses are displayed in Fig. 2-44. One of our best efforts, shown in Fig. 2-45, resulted in a 3.4-ns 2ω output pulse. The 4ω peak intensities corresponding to these square output pulses were 200 to 500% of average. Even

Fig. 2-44. Calculated intensity inputs for square output pulses of 8 and 10 kJ.

Fig. 2-45. (a) Shaped preamplifier input signal and (b) resulting 7 kJ 2ω output pulse; and MALAPROP simulations of (c) input and (d) output at 1ω and 2ω .

with these intense peaks, the number of hot electrons generated in these target experiments was extremely low, so we were able to use 1-ns Gaussian pulses from the MOR for the remaining shots.

4 ω Experimental Results. The operationally significant results of the 4 ω experiments include:

- Quasi-Gaussian pulses of 1-ns width at half power with energies exceeding 1600 J were delivered to the target chamber center.
- Ten to twenty percent of the total beam energy was missing, confirmed by energy balance measurements that were continued for several weeks after the completion of the target experiments. We ascribe this loss to the crystal surface roughness that resulted in wide-angle scattering of 4 ω light.
- Tuning-angle corrections, generated by scanning the crystal array at different ambient temperatures, allowed us to correct for laser bay temperature variations prior to each shot.
- The fused silica target chamber optics were visibly damaged after two 3-ns shots at 800 J and two 1-ns shots at 1 kJ. At the end of the 4 ω experiments, damage to the target optics included discoloration of the sol-gel coating and tracking ("angel hair") of the bulk material.
- The Pyrex calorimeter was visibly solarized, but the calibration was not affected.

X-Ray Laser Series. The last four months of the operation of Novette were dedicated to the x-ray laser experimental series.²⁴ Even prior to the end of frequency conversion experiments in early April, work was in progress to prepare for this final series.

Much of the equipment related to the delivery of properly mixed gases to some of the x-ray laser targets was installed and vacuum-tested. At the same time, a new optical beamline that triggers fast silicon switches used in diagnostics was installed between the preamplifier and the target chamber. We installed three spectrographs gated by the new optical trigger line and one streak spectrograph on the target chamber, and we removed most of the plasma diagnostic equipment used for wavelength scaling and compression experiments. In addition, because the x-ray series was scheduled to require only short laser pulses (300 ps or less), we removed one 46-cm disk amplifier from each of the two chains to

reduce ΔB (the phase shift resulting from intensity-dependent changes in index of refraction) in the 46-cm stage. For reproducibility of laser alignment, we installed a ground glass plate on a separate manipulator. This eliminated the need for an alignment aid on each target and simplified target construction.

Test Bed for Nova. The use of two Nova chains on Novette for a year made the installation of Nova significantly easier. We corrected a number of design flaws, and addressed several of the more serious problems in time to correct them for the Nova installation.

During the preparation of components for Novette installation, it became evident that the infant mortality of flashlamps in the 46-cm amplifiers was excessive. Initially, we were not able to deliver units in which all 16 lamp circuits were actually functional. Lifetime on line was relatively short compared to lamps in other amplifiers. We found that new designs for the flashlamps, the box amplifiers, and the power conditioning circuits had resulted in a convergence of the natural resonances of each, which tended to increase the mechanical stresses on the lamps, leading to premature failure. Early detection of this problem allowed us to prevent the same difficulty on Nova.

Another major problem inherent in the Nova baseline was damage to the crystal array index-matching fluid used with our first arrays. Incident 1 ω energies of more than 1 kJ caused the explosion of tiny aluminum particles suspended in the fluid, breaking the fluid down and leaving free carbon that continued the damage process on later shots. Elimination of these aluminum particles proved very difficult, so we searched for other means of reducing the reflections at the crystal surfaces. We developed a sol-gel coating to replace the index-matching fluid; this coating was applied to the surfaces of the KDP crystals used on all Nova chains. As a temporary expedient on Novette, however, we used bare KDP crystals enclosed in the original window assemblies but without the index-matching fluid. This allowed us to progress toward high-energy target shots.

As output energies were raised above 7 kJ, we began to notice bulk damage in the input lenses to the final two spatial filters in each chain. Since all our optics were

high-purity borosilicate glass from several manufacturers, damage-tested here many times, this was unexpected. We were concerned not only about degradation of beam quality, but also about the safety of personnel; these optics were the vacuum windows for the spatial filters, and the bulk damage could grow to the point where the lenses were dangerously weakened. Fused silica was tested as a replacement material and is now part of the Nova baseline configuration.

A calibration problem at the mid-chain (30-cm) sensor and at the output sensor was finally tracked to the response of several partially transmissive mirror coatings to the relative humidity of the ambient air. The 2% transmissive, 30-cm mirror in front of the mid-chain sensor was normally bathed in the dry nitrogen used to cool the disk amplifiers. The calibration measurements were taken at the mid-chain sensor by inserting a calorimeter into the main beam, thus exposing the mirror to laboratory air. The change in humidity gave repeatable calibrations, and the amplifier gains calculated from these measurements were consistently 25% low. More subtle variations of mirror reflectivity occurred at the output sensor, which was exposed to normal room humidity that varied between 30 and 60%. We replaced our original 2ω sensor with a calorimeter located where it was less subject to these fluctuations.

The Novette facility was also a development test bed for the Nova system software. Communication problems within the computer system were eliminated, process priorities were optimized, and basic

techniques for facility control were developed. For instance, an early code for automatic alignment of the laser chains was tried on Novette with only minimal success because it did not control enough of the hardware. Software control over the intensity of the alignment laser, together with improvements in image processing, made realignment a routine operating feature of the Nova facility.

Laser Performance. Novette output sensors provided us with $3 \times$ images of the 1ω and 2ω beams at planes equivalent to the volume surrounding the target chamber center. The 1ω beams were sampled at the first turning mirrors, while the converted 2ω light was recorded after reflection from the uncoated, inner surface of the second element of each of Novette's doublet lenses. We were thus provided with excellent images of the 1ω beams. Because each 2ω beam was reflected from a large, thin lens whose figure was questionable and then traversed the crystal array, whose elements introduced random phase changes, representations of the 2ω beams were more suspect.

Typical 1ω and 2ω focal spot images (from a 1-ns target shot) display a 2ω focus that is visibly larger than its 1ω counterpart (Fig. 2-46). The 1ω focus has a weak horizontal bar produced by the vertical split in the near-field beam. The 2ω far-field photograph shows the X-like diffraction pattern of the crystal arrays.

Reduction of the film images to intensity vs position (Fig. 2-47) enabled us to determine the effective spot sizes. These images

Fig. 2-46. Typical 1ω and 2ω focal spot images made during a 1-ns high-energy laser shot on Novette.

Novette Laser System

Fig. 2-47. Shot intensity profiles (south beam) reduced from images in the preceding figure and used to determine spot sizes.

Fig. 2-48. Power-induced broadening of Novette focal spots.

are strongly modulated at the diffraction limit (see structure 5 to 10 μm in size in this figure) with roughly Gaussian envelopes resting on weak pedestals several hundred microns in diameter.

Radii measured from beam centroid, within which 25%, 50%, and 80% of the energy was delivered, have been calculated for more than 70 images. Figure 2-48 contains plots of these radii measured on 1-ns south beam shots taken during October, November, and December, 1983. An in-house statistical optics code,²⁵ NORA, predicts that the 1ω focal spot grows rapidly as

we approach the onset of nonlinear index effects at 9 TW. The 2ω focal spot size is larger because of the additional ΔB between the output of the spatial filters and the target, and because of the additional passive aberration of the sensor optical path noted above.

Targets were often located about 1 mm from focus on Novette. Figure 2-49 shows the appearance of the north beam 1 mm from focus during a relatively low-power shot. The 3×3 array pattern is well resolved, suggesting that there is little distortion in the diagnostic system. A

Fig. 2-49. Shot intensity profile (north beam) taken 1 mm from focus.

MALAPROP computer code calculation of a Novette beam at a higher power, but near the same plane (Fig. 2-50), reminds us that the highly structured beam we measured is to be expected.

Several other measurements confirmed the size of the beam foci. X-ray micrographs of disk targets showed emitting regions that were used to determine best focus; these also indicated the diameter of the 2ω focal spot. Low power rod shots were used to burn metallic foils, and accurate measurements of these burn radii (Fig. 2-51) agreed with the x-ray micrographic determination. Finally, tests of a prototype target-plane imager, which consisted of a coherent fiber-optic bundle mounted at chamber center, gave similar radii.

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Target Focusing Configuration for X-Ray Laser Experiments

X-ray laser experiments imposed a new demand on the Novette focusing optics. These optics had to provide highly uniform, double-sided illumination on a target region 1.0 cm long by 100 to 200 μm wide. This line focus requirement had to be achieved without degrading the diagnostic reflection from the last surface of the focus lens and without potential ghost focus problems.

Fig. 2-50. MALAPROP calculation of intensity profile of Novette beam 0.9 mm from focus.

Fig. 2-51. Novette 2ω low-power burn radii.

Nova Two-Beam Target Chamber

Auxiliary optics outside the target chamber or tilts of the focus lens could not be used because of their effects on the diagnostic reflection.

The only optical configuration that preserves the diagnostic reflection is shown in Fig. 2-52. A negative focal length cylinder lens is placed between the focus lens and the debris shield, with the concave surface facing toward the focus lens. Any ghost reflections from the cylinder lens or debris shield are degraded by astigmatism, making them less hazardous. In practice, the uniformity of illumination is probably about the same for a positive or a negative cylinder lens. The minimum Novette focused spot was approximately 50 to 75 μm in diameter, and the fabrication errors in the 80-cm-diameter precision cylinder lens produced a line focus 25 μm wide. A negative cylinder lens design was chosen, however, to optimize the illumination uniformity in the case of line widths of several hundred microns.

For an optical beam diameter D and a focusing lens of focal length F_1 , a cylinder lens of focal length F_2 produces a line length L given by

$$L \approx \frac{DF_1}{F_2} \quad (2)$$

In order to get uniform illumination over a 1.0-cm line, L was set to 1.4 cm. Therefore, for $D = 74$ cm and $F_1 = 300$ cm, F_2 is about 160 m. The departure of the cylinder from a plane surface is only about 0.1 cm. As mentioned earlier, the passive optical aberrations of this cylinder lens are negligible

in comparison to the effects of beam divergence and optical fabrication errors.

Author: L. G. Seppala

Nova Two-Beam Target Chamber

Introduction

We plan to add a two-beam target chamber to the Nova laser fusion facility to increase our experimental flexibility and productivity. With construction completed in December 1984, the 10-beam Nova laser system and its target chamber are in great demand because of the increasing number of experiments that are important to the ICF, weapons, and x-ray laser programs.

We investigated the limitations on the effective target experiment rate for Nova and determined that target-related factors are more important than laser-related factors. We noted that target experiment campaigns frequently require different target diagnostic instrumentation and beam-focusing arrangements. Reconfiguring the diagnostics on the Nova target chamber can be a lengthy process, during which no useful target experiments can be performed. We have demonstrated that we can fire the full Nova laser at a much greater rate than targets can be installed in the chamber, aligned, fired, and diagnosed.

Adding a second target chamber to Nova will enable us to perform experiments in one chamber while another is being reconfigured. This will allow us to use Nova's excess capacity to perform experiments simultaneously when both chambers are in operation, effectively doubling our target experiment capacity and flexibility.

The addition of a full 10-beam chamber to Nova has been proposed. However, the success of the now-dismantled Novette system showed that a two-beam chamber can be beneficial when used to perform certain important experiments in the ICF and x-ray laser programs. This chamber will also have the capability to test diagnostics before installation on the 10-beam chamber.

We plan to move the two-beam chamber (Shiva target chamber) from Novette to the Shiva laser bay and install beam-propagation

Fig. 2-52. Focusing configuration for Novette x-ray target experiments.

elements, controls, laser, and target diagnostics support. Two of Nova's 10 beams will be selectively diverted into this chamber, while near-simultaneous use of the larger chamber for 10-beam or 8-beam experiments continues.

To accommodate x-ray laser experiments, the two-beam chamber will initially emphasize delivery of 2 to 5 TW of 0.53- μm light in 200- to 400-ps pulses. Target focusing optics will provide a line focus with a variable length line up to 3 cm long and approximately 200 μm wide.

Author: G. J. Suski

System Performance

Overview. We will divert 1.05- μm (1ω) light from two Nova beams in the optical switchyard and propagate it into the two-beam target chamber area (shown in the following article, "Mechanical Plan Layout"). There the beams will be frequency converted and focused for target experiments. Both 1.05- μm energy and beam quality at the frequency conversion arrays are expected to approach that of the 10-beam chamber arrays. However, some potentially degrading effects caused by the additional propagation distance to the two-beam chamber are under study. The longest beam path of the 2-beam chamber exceeds the longest of the 10-beam chamber by 120 ft.

The additional propagation distance may affect beam quality through the nonlinear processes occurring in the long air path. Processes currently under study are (1) nonlinear growth of intensity modulation caused by the n_2 of air, (2) stimulated Raman scattering, which generates small shifts in the fundamental wavelength of the beam that subsequently affect frequency conversion efficiency, and (3) an intensity-dependent rotation of the polarization ellipse of the partially depolarized Nova output beams. Alternatives are available to alleviate these effects if any should seriously limit the beam energy delivered to the two-beam chamber.

The 1.05- μm light will be converted by using two 3×3 arrays of KDP crystals mechanically similar to the Nova arrays, but with crystal thicknesses chosen to optimize 2ω conversion efficiency. (The Nova

10-beam chamber arrays were optimized for 3ω performance.)

Performance Specifications. Target irradiation experiments at 1.05 μm (1ω), 0.53 μm (2ω), and 0.35 μm (3ω) will be possible by using one of the two beams for x-ray backlighting. The 1ω experiments will require the addition of final-stage 1ω retropulse isolators in the beamlines. The final beam aperture will be the 74-cm output aperture of the Nova chain. Spherical focusing optics will be identical to those of the 10-beam target chamber. Counter-rotating fused-silica cylinder lenses will be used to provide a variable-length line focus (0 to 3 cm long) with a nominal 200- μm line width. The cylinder lenses will be mounted between the frequency conversion arrays and the $f/4$ spherical focus lenses.

The two frequency conversion arrays will each use two 3×3 arrays of 27-cm-wide KDP crystals. The first will be 10 mm thick and the second 18 mm. These will be Type-II quadrature arrays optimized for 2ω performance. The arrays can be reversed for 3ω operation. In this mode, the 10-mm-thick crystals will serve as mixers to generate 3ω light with approximately 10% lower performance than the Nova arrays.

The pulse length will be the same as in the 10-beam Nova system. Gaussian pulses will range from 0.1 to 5 ns and shaped pulses from 1 to 5 ns. A 30- to 100-ps short pulse capability is also expected.

We predict that the maximum power on target per beam will be approximately 5 TW at 0.53 μm , 2 to 3 TW at 0.32 μm (slightly less than Nova 10-beam chamber array performance), and 8 to 9 TW at 1ω with chain isolators. Estimates of the maximum expected energy on target per beam at pulse lengths of 0.2, 0.5, and 1 ns are given in Table 2-8. For short pulses (200 ps), 2ω and 3ω performance per beam will be 1200 and

Table 2-8. Expected maximum energy on target per beam.

		Energy at 200 ps (J)	Energy at 500 ps (J)	Energy at 1 ns (J)
Nova				
10- $\times$ 10-mm	2ω	950	2200	3000
	3ω	600	1500	3000
Two-beam				
10- $\times$ 18-mm	2ω	1200	3000	5200
	3ω	550	1300	2600
Novette				
18 mm	2ω	850	2100	4200
	3ω	—	—	—

Table 2-9. Technical baseline: Nova two-beam target chamber.

Nova Two-Beam Target Chamber

Irradiating wavelengths	0.35, 0.53, 1.05 μm with addition of isolators
Final beam aperture	74 cm
Focusing optics	$f/4$ fused-silica lenses, counter-rotating cylinder lenses to provide line focus 0 to 3 cm long
Frequency conversion crystals	2ω : 3×3 Type-II quadrature array, 10 and 18 mm thick 3ω : 3×3 array, 18-mm-thick doubler, 10-mm-thick mixer
Pulse length	0.1- to 5-ns Gaussian; 1- to 5-ns shaped pulses; 30- to 100-ps
Maximum power on target per beam (short pulse)	1ω : 8 to 9 TW with isolators 2ω : 2 to 5 TW 3ω : 2 to 3 TW
X-ray backlighting	Easily implemented using beam 6 as the backlighter
Target area switching time	6 h
Target shots begin	September, 1985

Fig. 2-53. Calculated 2ω conversion efficiency for 10- $\times$ 18-mm quadrature frequency conversion array.

Fig. 2-54. Projected 3ω performance of the Nova 10- $\times$ 10-mm arrays.

550 J, respectively. Corresponding 1-ns performance levels are projected at 5.2 and 2.6 kJ. This compares to the projected long-pulse performance for the Nova 10- $\times$ 10-mm arrays of 4.2 and 3 kJ.

Table 2-9 summarizes the major criteria for the two-beam chamber extension. When this beam chamber becomes operational, we can expect an average of one experiment on each chamber per two-shift day of operation. We anticipate that this rate will improve to as much as one shot per chamber per shift, if we have the scientific staff to support such a schedule. Actual switch-over time between chambers should be less than one-half hour. Delays due to thermal settling of the amplifier glass and final target alignment account for the difference between switchover time and shot rate.

Conversion Efficiency. The first target experiments in the two-beam chamber will emphasize x-ray laser studies. The design of the KDP frequency conversion arrays was chosen to optimize 2ω conversion efficiency while maintaining compatibility with the crystals used on Novette (18 mm thick) and Nova (10 mm thick).

The results from Novette and Nova indicate lower frequency-conversion efficiency than predicted at high intensities. We are beginning to understand the causes of the degraded performances, and further experiments are planned to obtain confirmation. However, in estimating the performance of the array design planned for the two-beam chamber, we used data scaled according to the measured Nova and Novette results. Therefore, we believe that the performance projections for the current baseline are realistic.

Conversion efficiency as shown in Fig. 2-53 is approximately 65% over the range of 1.5 to 3.0 GW/cm^2 (5 to 10 TW) of 1ω light. Eventual improvements toward the theoretical limits of 70 to 80% may be possible as the quality of the Nova 1ω beam is improved.

For 3ω operation, the arrays will be reversed to allow use of the 10-mm crystals as the mixer or second array element. Results from Nova 10- $\times$ 10-mm arrays indicate significant roll-off from the theoretical prediction at about 1.5 GW/cm^2 . Consistent with these results, we project the 3ω performance of the two 10- $\times$ 10-mm arrays to be greater than 50% at 1.5 GW/cm^2 . This implies that the energy per beam on target will be slightly less than projected for the Nova 10- $\times$ 10-mm arrays shown in Fig. 2-54.

A universal deep-web "egg-crate" design is used to support all our crystal arrays. Since the support assemblies are mechanically identical with respect to mounting and positioning, several options in addition to the baseline array can be accommodated. Specifically, the baseline 10×18 -mm arrays may be interchanged with the 10×10 -mm arrays for the 10-beam chamber, which are optimized for 3ω . The 5×5 array of 10-mm thick, 15-cm-wide crystals, which was operating in cascade on Novette to generate 4ω light, will be retained and used to generate ~ 1 kJ of 4ω light in a 1-ns pulse.

Final Focusing Configuration. Figure 2-55 shows the nine major elements in the optical train, as the beams approach the target chamber. The elements are:

- An optional full-aperture 1ω calorimeter, which is removed for target experiments.
- The KDP frequency conversion array.
- An optional beam dump to absorb unwanted 1ω light (during 2ω experiments).
- An insertable centering mirror used during beam alignment.
- The cylindrical lens doublet described in the next section.
- A 74-cm calorimeter used to measure 2ω or 3ω energy after the arrays.
- A beam-diagnostic assembly that measures 2ω and 3ω incident energy and temporal profile.
- A spherical $f/4$ fused-silica focus lens that forms the vacuum barrier for this target chamber.
- A 44-cm calorimeter that can be placed inside the target chamber for calibrating the $2\omega/3\omega$ diagnostic sensor package.

Final Focusing Optics. The variable line focus for x-ray laser experiments is provided by two counter-rotating negative cylindrical lenses in each beam path. The

Fig. 2-53. Optical and diagnostic assemblies for beam propagation into the target chamber.

Fig. 2-56. Effect of cylinder lens choice on ∇ minimum line width.

Fig. 2-57. Counter-rotating negative cylinder lenses will provide variable-length line focus.

Nova Two-Beam Target Chamber

negative cylindrical lens pair was designed to provide minimum spherical aberration while keeping a common design for both cylindrical elements as shown in Fig. 2-56. The lens material is fused silica to prevent 3ω solarization. Maximum line length at the target occurs when the lenses are colligned, and minimum line length is achieved when each is counter-rotated about 45° , as shown in Fig. 2-57. Accuracy of rotation is sufficient to maintain a line focus tilt of less than $30\text{ }\mu\text{m}$. The width of the line focus is expected to be approximately $200\text{ }\mu\text{m}$ with the aberrations in the laser system dominating.

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Mechanical Plan

Introduction. Conceptual planning for a Nova target chamber with two beams included consideration of the available locations. Three sites were studied: the Shiva target room, adjacent to the west wall of

the Nova switchyard; the Shiva laser bay, west of and adjacent to the Shiva target room; and the Novette laser bay in Building 381, southwest of Nova. The value of the third location was that the target chamber, support, and auxiliary equipment would not have to be moved. The long, awkward beam transport path was judged to be a major disadvantage.

We decided to locate the new, two-beam target chamber in the Shiva laser bay, primarily to keep the adjoining space, the Shiva target room, free for existing and future uses. (See Fig. 2-58.) Currently, the target room is in use as an optical coating laboratory. Projections indicate a need to keep this room open for possible installation of a second 10-beam target chamber.

Once the decision to locate the target chamber in the Shiva laser bay was made, it was necessary to find the correct arrangement of beams and mirrors to minimize the path length difference between beams six and seven. The distance from the west Nova laser bay wall to the center of the two-beam target chamber is 105.6 m on beam seven and 91 m on beam six. (See Fig. 2-58.) We also needed to minimize

Fig. 2-58. New Nova 2-beam target chamber room, switchyard with mirror switchout towers, the Nova 10-beam target room, and the laser bay.

the path length difference between the new 2-beam target chamber and the existing Nova 10-beam target chamber to avoid operational and diagnostic differences between target chambers.

It was necessary to use existing Novette hardware for the new installation to meet our rigid schedule. A spaceframe to support all the target system equipment was needed and could possibly have been a long-lead item. This delay was avoided by moving the Novette frame to the new location. We oriented the frame with respect to the two incoming beams so that the beam path differences were small and easy to equalize and so that the target elevation was compatible with that of the incoming beams. (See Elevation, Fig. 2-59.) These conditions were met with a few modifications of the Novette target frame.

Spaceframe. The target chamber spaceframe was mechanically sawed into seven segments. These were sized to be compatible with crane handling and doorway openings for the move from the

Novette building to the Shiva laser bay. Sawing was used to make clean, square joints for rewelding and to minimize particulate generation. Both the Novette room and the Shiva room are clean rooms and had to be kept clean for future use. We removed the Shiva spaceframe with conventional burning, using exhaust hood procedures to protect the room. Our initial plan for removal of the Shiva spaceframe was to saw it also, but burning was much faster, and it was appropriate because this frame was not to be reused.

The Novette target spaceframe, modified slightly by reducing platform space, was installed and welded in the former Shiva laser bay as shown in Fig. 2-59. Sixty percent of this room is taken up by the Nova two-beam target chamber facility. The remaining space will be used as a class-10 000 clean room for the maintenance of Nova hardware and to house the Zygo interferometer.

Mirror Supports. Three additional mirror support towers hold the mirrors that transport the beam from the Nova switchyard to

Fig. 2-59. Plan view and elevation of the Nova two-beam target system, with class-10 000 clean room for Nova component maintenance (far left).

Fig. 2-60. Nova switchyard showing existing support towers and mirror support towers added for the two-beam system.

Fig. 2-61. Nova two-beam target frame orientation and final turning mirrors.

the new target chamber. Two of these are located in the Nova switchyard (Fig. 2-60) and the third is in the two-beam target chamber room. The tower for beam seven is located in the north bank of five Nova beams at an elevation that allows the beam to travel straight through the room previously used to house the Shiva target chamber. It goes to a mirror on the two-beam target chamber frame and from there directly into the new target chamber with no elevation change. (See Fig. 2-61.) The mirror support tower for beam six is located near the south bank of five Nova beams at an elevation 5.0 ft higher than that of beam seven. We equalize the path lengths of beams six and seven by making adjustments to beam six. An intermediate mirror, on the third tower in this system, is located in the Nova two-beam target chamber room. (See Fig. 2-61.) This mirror has an adjustable base to set fine adjustments for path equalization. It is mounted 5.0 ft higher than the target chamber and aims the mirror downward to the final turning mirror.

Two mirrors in the switchyard turn the beam east 90° from its northward path. One mirror on each beam moves in and out of the Nova beam path to switch the beam between the two target chambers. These mirrors are mounted on accurate

rails and are driven with a motor and ball screw. (See Fig. 2-62.) They are designed with a positive stop location for repeatable alignment and they travel approximately 60 in. in 3 min. This allows easy switching between the two target chambers. The mirror adjustment mechanism for path length equalization is identical to this except that the travel path is only 12 in. long.

The two mirror tower supports in the switchyard are designed with bolted connections; thus, they can be easily transported through the laser bay without disturbing operations and they can be erected cleanly without welding. Since bolted construction is also adaptable to partial erection, these two tower supports can be installed with minimal interruption of operations. These structures are designed with the same stability criteria as other supports in the system. An attempt was made in the design to allow for clear personnel access and traffic flow.

Since the target room has no construction restrictions like those in the switchyard, the single mirror support tower there is designed with less expensive, welded construction.

Mirrors. Seven additional mirrors are needed for the two-beam target system. Six are 94 cm in diameter and one is 106 cm in diameter (Nova standard sizes) and they

Fig. 2-62. Insertable mirror that switches the beam from the Nova 10-beam target chamber to the 2-beam target chamber.

Nova Two-Beam Target Chamber

will be mounted in Nova bezels. Their incident angles are within the necessary tolerances, so the optical coatings can be the same as they are on Nova.

Beam Tubes. Approximately 400 ft of 36-in.-diam beam tube must be added to shroud the beam from the switchout mirrors to the target chamber. (See Fig. 2-59.) The beam tubes are designed to be vacuum tight in case it becomes necessary to evacuate them, to introduce other gases, or to use the tube as a relay spatial filter. The largest span is approximately 60 ft, measured across the Shiva target chamber room. At the lowest point, the beam tube is 7.0 ft above the floor in the Nova optical coating

laboratory. Intermediate supports from the beam tube to this floor or to the basement floor are not acceptable: they would preclude starting construction of a 10-beam target system in this room while operating the 2-beam system. This 60-ft span will be effectively shortened by 10 ft with 5-ft support extensions from the east and west walls to minimize deflection. This will raise the fundamental frequency of the beam tube and make it stable enough to use it as a spatial filter, collimating the beam through a pinhole, if we should choose to do so.

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Target designers develop scientific theories pertaining to thermonuclear microexplosions and develop computer codes that will simulate laser-target interaction. A variety of targets are designed for experiments involving radiation by large lasers such as Novette and Nova; developing these intricate targets usually requires sophisticated numerical simulations. While the primary focus is on designing targets for laser fusion, much attention has also been given to the x-ray laser target. The large photo at left is an XRASER code simulation showing contours of x-ray laser gain within the butterfly-shaped lasing cavity at the center of the target. Each letter refers to a value and two parameters are important in optimizing the design: the value in the center of the cavity and the variation in the value as it moves away from the center. The top inset photo shows a portion of a ray direction plot for spectral-line transfer simulations. The photo below shows a mesh plot used for two-dimensional line-transfer calculations, which are done to predict x-ray laser gain.

Target Design

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Section Editor: A. Friedman

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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 these theoretical and computational tools to design targets and simulate experiments.

Effort in laser-target theory and analysis during 1984 was split about equally between the very successful experiments on Novette and the upcoming experiments for Nova. We made major progress on Novette toward obtaining a quantitative understanding of laser-plasma coupling for submicron laser light. A particularly successful series of experiments, designed in collaboration with personnel of the Laser Experiments and Advanced Diagnostics Program, demonstrated the suppression of Raman scattering by collisional damping. For these experiments we correctly predicted target density profiles, average plasma temperatures, and collisional thresholds for Raman scattering. Our conclusion from all the submicron-wavelength coupling experiments is that laser light with a wavelength between 0.25 and 0.5 μm will provide the plasma conditions required for high-gain inertial confinement fusion (ICF) targets. More precise definition of the required wavelength must await experiments on Nova with high-gain-target plasmas.

We played a major role in the invention and design of the exploding-foil target that led to the world's first laboratory demonstration (on Novette) of soft x-ray lasing. Besides this neon-like selenium collisional excitation scheme for demonstrating x-ray lasing action, we also made significant progress in modeling photoresonant and recombination schemes.

Design work for Nova concentrated on targets for the study of irradiation symmetry, implosion stability, and the physics of long-scalelength plasmas, all important issues for high-gain ICF targets.

In 1984 we implemented several numerical tools important for our analysis and design of these increasingly complex and sophisticated targets. A 3-D ray-tracing package and a fully Eulerian version of LASNEX are now being widely used in our work. The Eulerian LASNEX permits us to calculate targets in which large shear flows occur. An analytic Thomas-Fermi-based equation of state model has been implemented and is being checked out. We have begun work on a Monte Carlo atomic physics model that could lead to an inexpensive way of greatly increasing the precision of line-emission calculations for high-Z plasmas.

We have continued our low-level effort on fusion driven by charged-particle beams, and have demonstrated that it is possible, in principle, to increase the power of a heavy-ion beam, by one to two orders of magnitude by longitudinal compression without developing excessive longitudinal beam temperatures. Such beam-power increases are crucial if heavy-ion beams are to achieve the required intensity on target.

The following articles report details of 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 Sec. 5, "Laser Experiments and Advanced Diagnostics."

Laser Targets

Introduction

The laser target design group was engaged in three main tasks in 1984: (1) analyzing Novette implosion and hohlraum-scaling data, (2) planning for the first experiments on Nova, and (3) designing laboratory x-ray laser targets and experiments.

The Novette implosion and hohlraum-scaling data are mostly classified and are therefore not discussed in detail here. We achieved average final/initial pusher ρr ratios of about 50, some 3 times higher than the value achieved in the best Shiva shots. These ρr values imply a fuel compression to 100 times liquid density, although this figure and other aspects of the experiments are subject to further interpretation because of detailed questions of target symmetry and stability.

Our main long-term goal for Nova is to produce a so-called hydrodynamically equivalent target (HET)—that is, a target whose hydrodynamic behavior (implosion velocity, convergence ratio, symmetry and stability requirements, etc.) is very much like that of a high-gain target, but one that is scaled down in size to match the energy available from Nova and is too small to achieve enough hot-spot ρr to ignite the cold, near-Fermi-degenerate fuel around it. Our goal for Nova's first year is to do experiments that will teach us how to achieve the symmetry and stability conditions required by an HET.

We have been developing theoretical tools to improve our modeling of the stability problem: these are described in the articles "RIPPLE—A Spherical-Harmonic Postprocessor for 1-D LASNEX Simulations," "2-D LASNEX Performance on an Analytically Solvable Rippled-Shock Problem," and "Richmyer-Meshkov Instabilities in ICF Capsules."

The laser target design group also began a new endeavor in 1984, the design of laboratory x-ray lasers (XRLs). Our expertise

in modeling targets driven by high-power lasers combined with efforts in inventing new target designs and improving others culminated in the spectacular success achieved at Novette in the summer of 1984 with the world's first demonstration of soft x-ray lasing in the laboratory.

In "Design of an Exploding-Foil X-Ray Laser Target," we describe the design of the exploding-foil target that led to the successful selenium collisional excitation scheme used in the Novette XRL experiments. In "Design of a Photoresonantly Pumped Fluorine X-Ray Laser," we describe modifications made to the gas-cavity target for another scheme, in which hydrogen-like fluorine is photoresonantly pumped by a laser-driven, beryllium-like manganese flashlamp. While evidence for lasing in those experiments is somewhat ambiguous, we learned much from the effort. A third scheme, a recombination XRL ($n = 4$ to $n = 3$ in hydrogen-like magnesium) driven by rapid cooling due to target expansion, also showed signs of success. The data from that experiment are still being analyzed, so we postpone discussion of it until next year.

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RIPPLE—A Spherical-Harmonic Postprocessor for 1-D LASNEX Simulations

The first-order hydrodynamics code RIPPLE provides a lower-cost alternative to 2-D LASNEX for problems involving small deviations from spherical symmetry. RIPPLE, developed as a postprocessor for 1-D LASNEX, solves the linearized hydrodynamic equations using the LASNEX results as zero-order solutions. Linearity is not a severe limitation, since a target designer is frequently only interested in estimating the severity of instabilities and in determining the spatial wavelengths that exhibit greatest instability.

Rayleigh-Taylor instability of shell structures in ICF targets is a *primary limiter* of target performance. In certain ranges of target

parameters, notably for thin shells that have large aspect ratios $A = r/\Delta r$, instable growth amplifies asymmetric perturbations (such as shell fabrication imperfections) enough that the shell becomes highly distorted. When this happens performance can be significantly degraded from that predicted by symmetric, one-dimensional design calculations. Performance is therefore limited to what can be obtained with target parameters that do not lead to large instability growth. Since one-dimensional calculations often predict that target performance increases strongly with decreasing shell thickness (i.e., with increasing A), it is important to be able to predict accurately under what conditions instability will cause excessive shell distortion.

Useful order-of-magnitude estimates of instability growth have been obtained¹ by assuming that unstable modes have radially localized structures and that the equation of state (EOS) is incompressible. More accurate formulations, such as that in RIPPLE, must account for a number of additional considerations. For one, instability growth is sufficiently sensitive to target parameters, and calculations of growth are sufficiently sensitive to physical and numerical approximations, that instability calculations should be carried out with the same accuracy and using the same approximations as the one-dimensional design calculations. In particular, in perturbation calculations it is necessary to have nonlocal interaction of perturbation mode structure when communication between adjacent shell surfaces or between multiple shells is important. Further, it is necessary to treat the EOS and the artificial viscosity in a way that is consistent with the one-dimensional calculations.

RIPPLE solves a linearized system of hydrodynamic equations, using 1-D LASNEX results as zero-order solutions. The first-order solutions given by RIPPLE are three-dimensional, angle-dependent departures from the zero-order solutions. (The alternative to linearization would be full, nonlinear, two- or three-dimensional calculations. While such calculations should be capable of predicting the degrading effects of large-amplitude distortions, efforts to treat even small departures from sphericity with large, multidimensional design codes have often been frustrated by the finite size of the zones in the azimuthal direction.) RIPPLE's

linearized calculations predict the small-amplitude behavior, from which one can determine the time at which nonlinearities become important. Linearization offers relative computational economy, the ability to linearly superpose the results of calculations of the effects of different initial and/or input perturbations, and the ability to calculate in a postprocessing mode. The remainder of this article describes the structure of RIPPLE and presents an example of its use.

Linearized Hydrodynamic Equations.

The procedure followed to obtain consistent first-order equations of motion for first-order perturbation quantities is to expand all angle-dependent quantities in spherical harmonics $Y_l^m(\Omega)$. First-order quantities are then defined at appropriate radial positions on the one-dimensional, Lagrangian LASNEX grid. Spatial differencing is done so that the first-order equations are exactly time derivatives of departures of variables from their zero-order values on a fictitious multidimensional Lagrangian grid with vanishing angular (nonradial) zone intervals. The resulting scheme is fully Lagrangian. We find this approach more robust and more accurate than mixed Eulerian-Lagrangian first-order treatments of ICF implosions.

If $\xi_l(t)Y_l^m(\Omega)$ is the l th component of the perturbed displacement or the grid surface at unperturbed position r , angular position Ω , and time t , then the equations of motion for first-order quantities for any l are as follows (all first-order quantities except ξ have subscript 1; zero subscripts indicate zero-order quantities):

$$\frac{d\xi_l}{dt} = \frac{4\pi r^2}{M_0} \left\| \frac{2\xi_l}{r} + (\nabla \cdot \xi)_l \right\| \times \left(P_{l+1} - P_l + P_{l-1} - P_l \right) \quad (1)$$

$$\frac{d\xi_l}{dt} = \frac{\xi_l}{\tau_l} \quad (2)$$

$$\frac{d}{dt} (\nabla \cdot \mathbf{P}_l)_{12} = -l(l+1)r \left\| \frac{\xi_l}{r} + r(P_{l+1} - P_l) \right. \\ \left. - \frac{1}{6} [P_{l+1}(r_{l+1}^2 + r_{l-1}^2 - 2r^2) \right. \\ \left. + P_{l+1}2r^2 - r_{l-1}^2 - r_{l+1}^2] \right\| \quad (3)$$

$$\frac{d}{dt} (\nabla \cdot \xi)_{\Omega_i} = \frac{(\nabla \cdot P_1)_{\Omega_i}}{M_{0,i} r_i^2}, \quad (4)$$

$$\rho_{i,j} = \rho_{0,i} (\nabla \cdot \xi)_i, \quad (5)$$

$$\frac{dS_{1,i}}{dt} = \frac{1}{M_{0,i} T_{0,i}} \times \left[-\frac{dQ_0}{dt} \left(\frac{T_1}{T_0} + \frac{M_1}{M_0} \right) + \frac{dQ_1}{dt} \right], \quad (6)$$

$$P_1 = \left(\frac{\partial P}{\partial \rho} \right)_\rho r_1 + \left(\frac{\partial P}{\partial S} \right)_S S_1, \quad (7)$$

$$T_1 = \left(\frac{\partial T}{\partial \rho} \right)_\rho \rho_1 + \left(\frac{\partial T}{\partial S} \right)_S S_1. \quad (8)$$

Here ρ , p , T , M , Q , and S are density, momentum, pressure, temperature, zone mass, heat added, and specific entropy, respectively. These are equations for the perturbed values of (1) radial momentum, (2) radial position, (3) angular components of momentum, (4) angular position, (5) density, (6) specific entropy (in effect, energy), and (7), (8) EOS.

Particular care was taken to make the difference equations that determine the first-order artificial viscosity contributions to P and dQ_1/dt consistent with the zero-order equations in LASNEX. The terms in Eq. (3) containing various sums and products of r , p , T , and r_1 have the particular form shown to be consistent with the geometry of zero-order, multidimensional quadrilateral Lagrangian zoning in the limit of vanishing angular zone dimensions. Consistency also requires that values of the thermodynamic derivatives that occur in Eqs. (7) and (8) be passed from LASNEX along with all other zero-order quantities required by RIPPLE.

At each LASNEX time step, all quantities needed by RIPPLE are stored in files that are read later. The first-order hydrodynamic equations are sometimes stiffer than the zero-order equations because of short-wavelength angular motion. This increased stiffness enters through the factor $l(l+1)$ in Eq. (3). Therefore, to maintain numerical stability it is sometimes necessary to use smaller time steps in RIPPLE than in LASNEX and to interpolate zero-order

quantities between the values calculated at each LASNEX time step.

Example Use of RIPPLE. The sample problem discussed below considers a single pusher shell of unit initial density and aspect ratio $A \approx 4$, containing low-density gas. Figure 3-1 shows the time development of perturbations of the inside surface of the shell. The successive rows of Fig. 3-1 represent a succession of times during the run. The first frame in each row shows the zero-order density profile at that time, and successive frames show the radial component of the perturbed displacement, ξ_r , for spherical-harmonic l values of 0, 10, 20, and 40. Many first-order quantities are calculated in RIPPLE, but we show ξ_r because it is a readily interpreted measure of distortion.

The initial perturbation used here for all l is specified by giving ξ_r an approximately Gaussian form [as seen in row (a) of Fig. 3-1], with maximum displacement at the inside surface of the pusher shell and with width small enough that the outside surface is undisturbed, and by taking $(\nabla \cdot \xi)_{\Omega_2} = 0$. The perturbed masses M_1 are then calculated so that the initial perturbed density is zero. When this initialization procedure is followed and the different material regions are homogeneous, only the initial values of ξ_r at material interfaces have physical significance.

Note that the maxima of all the initial ξ_r , which occur at the inside pusher surface, have the value unity. The linearity of the calculational method makes this value arbitrary, and with this choice ξ_r represents (at any time) the factor by which any surface in the target has been distorted relative to the amplitude of the initial perturbations of the inside surface of the pusher.

Row (a) represents a time shortly after the beginning of illumination. The outside surface has just begun to ablate, but the perturbed part of the shell has not begun to move.

In row (b), the first shock has progressed almost to the inside surface of the pusher (note that the vertical scale on which p_r is plotted has changed). The initial perturbation has not yet had a physical effect.

A short time later, in row (c), the shock has broken out of the inside surface, and a negative (i.e., phase-reversed) spike

in ξ_r has appeared just inside the inside surface of the pusher, for all values of l . This phase reversal occurs because the jump-off speed of the pusher into the lower-density fuel exceeds the shock speed in the pusher. At the lower l values, the perturbation is little different from its symmetrical ($l = 0$) form; but at $l = 40$ and above, there is a small difference in the magnitude of the phase-reversed spike, which is evidence of the Richtmyer-Meshkov effect.

In row (d), the imploding pusher shell has moved about halfway to the origin. At lower l the perturbation is approximately just the phase-reversed negative spike generated when the first shock broke out of the inside surface, which has been convected inward with the imploding shell. However,

at larger l there is a broader, negative feature outside this spike, which is the outside-surface Taylor-unstable mode that is driven by ablation pressure. As this mode grows during inward acceleration, and because it couples to the inside surface of the shell, it adds to the distortion of the inside surface. This contributes to the amplitude from which instability of the inside surface begins to develop when deceleration of the shell begins.

In row (e) the shell is still moving inward, but significant deceleration has already occurred. One can distinguish three regions or relatively large spatial density gradients, corresponding to the different regions of hot fuel, cold fuel, and pusher. These regions drive Taylor growth. At $l = 10$ and $l = 20$, the growth of ξ_r , which

Fig. 3-1. Zero-order density profiles, and first-order radial displacement profiles for $l = 0, 10, 20$, and 40 . Each row represents a single time in an implosion. Note vertical-scale changes.

is driven by these density steps, takes the form of separate ξ_r peaks. At $l = 40$ and above, there is relatively little inside-surface instability, because of relatively weak coupling between the inside surface and the unstable behavior of the outside surface. At this time the outside-surface mode is seen at all l , but the inside-surface modes are most important because they contribute directly to fuel-pusher mixing.

Row (f) is at approximately peak compression; the pusher has been compressed by a factor of about 10. The outer density peak is the remnant of the cold pusher shell; the inner peak is an outward-propagating shock. At $l = 10$ and $l = 20$, the multiple-peak structure of ξ_r inside the pusher has grown a bit from row (e). At $l = 40$, there has been more growth of ξ_r inside the pusher from row (e), but the amplitude still is not as large as at $l = 10$ or $l = 20$. At $l = 80$ (not shown), ξ_r is still small everywhere.

Application of such calculations to a specific target would involve analysis of initial perturbations of the inside and outside surfaces into a sum of spherical-harmonic l and m components. These components would then be multiplied by their amplifications during the calculation (the amplitudes of the ξ_r , as shown here), and re-summed at about the time of peak compression to predict the form and magnitude of target distortion.

Quantitative Interpretation. Suppose that for the target in the sample calculation the dominant initial perturbations were near $l = 20$, on the inside surface, and had the effect of introducing an $l = 20$ component of ξ_r , whose magnitude was about 10^{-3} times the inside-surface radius. Row (f) indicates that the maximum of the highest peak of ξ_r is almost 30 for $l = 20$. Thus the physical value of the radial distortion would be approximately 0.03 times the initial radius, or 0.3 times the radius at peak compression. This largest peak occurs not at the inside surface of the pusher but near the transition from the central, hot-fuel region to the outer, cooler-fuel region. The relative radial distortion of 0.3 could be disruptive; multidimensional, nonlinear calculations would be needed to assess the effects in detail for $l = 20$. The major negative effects in

this case would be mixing of hot and cold fuel. By contrast, if the magnitude of the initial perturbation were 10^{-4} times the inside-surface radius, these calculations would predict no significant degradation by hydrodynamic instability.

In many situations, the knowledge that the worst spatial wavelength for instability growth lies near (say) $l = 20$, and that this mode exhibits growth by a factor of (say) 30 during the course of the implosion, will be enough for a target designer, and a full, nonlinear 2-D LASNEX calculation may not be required. Since a single RIPPLE calculation (with a half-dozen l values) costs far less than a single 2-D LASNEX calculation (with only a single l), it is clear that RIPPLE may become an important design tool.

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2-D LASNEX Performance on an Analytically Solvable Rippled-Shock Problem

How do small deviations from perfect symmetry evolve under the influence of the violent hydrodynamic processes associated with an ICF implosion? Early Rayleigh-Taylor experiments or Nova will attempt to monitor the evolution of such perturbations. In the absence of such experiments, ICF target designers are forced to trust the results of 2-D LASNEX calculations to assess the effects of deviations from symmetry. There are very few analytically solvable hydrodynamics problems that can be used to check LASNEX simulations of fluid motion in slightly perturbed cases; one such test problem is reported here.

A first-order perturbation solution of a problem of a plane shock wave propagating through an infinite, homogeneous medium can be obtained analytically. This solution seems to be absent from the literature, although it is a simple extension of the solution used in analyzing the stability of a plane shock front.^{1,2} Suppose the front propagates into a medium that has a very slight density fluctuation, in the form of a plane wave with a wave vector parallel to the shock front (no density gradients in

direction of shock propagation). After the shock has propagated some distance, a steady state is reached in which the shock front is rippled like a piece of corrugated sheet metal, protruding slightly into the lower-density regions and lagging slightly in the higher-density regions. In the steady state, the flow behind the shock consists of an evanescent sound wave, a shear wave, and an entropy fluctuation (a density fluctuation with no pressure fluctuation). The sound wave relaxes within a few transverse wavelengths, while the shear and entropy waves persist indefinitely behind the shock front and play a crucial role in the so-called Richtmyer-Meshkov instability.^{5,6}

The perturbation analysis is carried out in the rest frame of the unperturbed shock front. For a given transverse wave number and perturbation frequency, there are four first-order modes in the homogeneous flow on each side of the shock front. Four modes arise on each side because four conservation equations must be satisfied—for mass, two components of momentum, and energy. The four modes are an upstream- and a downstream-propagating sound wave, a shear wave, and an entropy wave. (In three dimensions, with an additional momentum-conservation equation, there is a second shear wave, which will never couple to the other four modes.) At the shock front itself there are four first-order jump conditions (also corresponding to the four conservation laws), which depend on the first-order amplitude of the ripple on the shock front.

The jump conditions place linear constraints on the nine first-order quantities—the amplitudes of the four modes on each side of the shock front and the amplitude of the ripple on the front itself. The resulting five independent first-order quantities are determined by the five types of first-order perturbation that can be input to the shock front—an arbitrary combination of the four first-order modes upstream of the front, and an arbitrary upstream-propagating sound wave from the downstream side of the front. Once these five perturbation amplitudes are given, the jump conditions determine the other four first-order amplitudes—those of the downstream-propagating sound wave, the

shear wave, the entropy wave downstream of the front, and the shock-front ripple. The test case described here is a special case or this general perturbation situation, corresponding to a zero-frequency entropy wave input upstream of the shock front, with none of the other four possible first-order input modes.

A variety of LASNEX runs were made to understand the characteristics of the numerical methods employed in the HYDRO subroutine⁷ when solving this type of problem. The basic LASNEX problem consisted of a 10^7 -cm/s piston pushed into a cold, 200- μ m-thick slab of $\gamma = 5/3$ perfect gas. When the resulting strong shock had propagated 110 μ m through the slab, it moved from a perfectly homogeneous region into a region with an extremely small, sinusoidal, transverse density fluctuation. The sudden encounter with the perturbed section produced a transient response that could not be calculated with the simple perturbation theory outlined above. This unwanted transient response was tolerated in order to avoid an even less desirable interaction of the perturbation with the piston, which is in sonic communication with the shock front. The test problems ended before first order acoustic waves generated by the sudden onset of the perturbation could be reflected from the piston into the region of interest. The transverse zoning covered one-half of a transverse wavelength, with symmetric boundary conditions to assure a periodic solution. Transverse wavelengths of 10 and 25 μ m, 13 and 26 transverse zones, 200 and 400 longitudinal zones, and normal and 10-times-normal artificial viscosity were run separately and in various combinations.

The first-order part of the LASNEX solution was separated from the zero-order part (a one-dimensional shock front) by performing a discrete Fourier transform on any desired field quantity (density, velocity, etc.) in the transverse direction. The first-order transverse velocity amplitude was the first Fourier sine component, while all other first-order amplitudes of interest were first Fourier cosine components. The initial perturbation was entirely in the first transverse Fourier component; one measure of the linearity of the calculation was that the coupling to higher Fourier components was

negligible. These first fourier components of the LASNEX solution were compared directly with the results of the analytic solution to assess LASNEX performance.

LASNEX reproduced the analytic results qualitatively, exhibiting a branching of the perturbation into an evanescent sound mode and persistent shear and entropy modes after shock passage. The quantitative differences between LASNEX and theory provide considerable insight into the limitations of the numerical methods used in LASNEX for solving this type of problem.

The most noticeable discrepancy between LASNEX and theory is the presence of a high-amplitude oscillation (evident in Fig. 3-2) at the spatial Nyquist frequency for the LASNEX mesh. This high-frequency noise in the first-order component of the

solution is related to the ringing behind any shock front caused by the use of a quadratic artificial viscosity. Because a quadratic viscosity cannot damp small disturbances, the density behind a shock front will oscillate with a spatial period near twice the zone thickness (i.e., in the Nyquist mode) and with a temporal period corresponding to the sound transit time across this distance. With the LASNEX default artificial viscosity, these density oscillations have an amplitude between 0.5% and 1% of the post-shock density. The planes of constant phase of the ringing are ordinarily parallel to the shock front.

However, when there are first-order differences in the sound speed in the transverse direction behind the shock, a phase difference will accumulate between different

Fig. 3-2 First-order density at three times after the shock (moving to left) strikes the perturbation, which begins at zone 90. Smooth curves, theory; noisy curves, LASNEX.

transverse locations. The result is a slow leakage of the ringing amplitude from the zeroth to the first transverse Fourier component. Since the fractional amplitude of the oscillations is about 1%, roughly 20 $(100/2\pi)$ cycles of the oscillation are required for the leakage into first order to become comparable with the other first-order amplitudes in the problem. The leakage model is in reasonable agreement with Fig. 3-2, which shows a roughly constant rate of growth in the first-order ringing at each point behind the shock front. (The three snapshots in the figure correspond to about 40, 60, and 80 cycles of the ringing since the shock encountered the perturbed section.) When twice as many longitudinal zones were used, the amplitude of the first-order ringing doubled since the ringing frequency doubled while the amplitude in

zero order remained the same. Running with 10 times the normal artificial viscosity dramatically decreased the zero-order ringing amplitude, with a corresponding reduction in the first-order amplitude.

The second discrepancy between LASNEX and theory is quantitative; even if the high-frequency noise were filtered out of Fig. 3-2, the first-order density computed by LASNEX would differ by at least 50% from the analytical result. The discrepancy is easier to see in Fig. 3-3, where the first-order entropy is plotted (the high-frequency ringing is an acoustic mode, so it is absent from a plot of entropy). The sudden onset of the perturbation in the LASNEX problem gives rise to the damped-exponential transients in the figure, which are not calculated in the theory. Even after the transient has damped, the 200-zone LASNEX calculation

Fig. 3-3. First-order entropy (the shock has reached zone 10).

shows a 50% error in first-order entropy, while the 400-zone calculation shows a 25% error. The fact that the error is halved when the zoning is doubled strongly suggests that the culprit is the coarseness of the zoning. However, the large magnitude of the error is somewhat surprising, considering that both problems had 13 transverse zones, or 26 zones per transverse wavelength.

The transverse wavelength of the perturbation for the runs of Fig. 3-3 was 10 μm . In the 200-zone problem, each zone had a longitudinal dimension of 1 μm . For the default LASNEX artificial viscosity, this implies a shock front thickness of 1 μm , which corresponds exactly to one radian of the transverse periodicity. Since it is clear that the analytic theory would require modification by terms of order the ratio of shock thickness to transverse wavelength to account for shock fronts of finite thickness, the rather large discrepancy in the LASNEX calculations is easy to understand. The run made with 10 times the normal artificial viscosity had a shock front 3 times thicker, and the discrepancy in first-order entropy between LASNEX and theory was correspondingly about 3 times as great. Furthermore, the run with 25- μm transverse wavelength and 200 longitudinal zones showed about a 20% discrepancy with theory, as opposed to the 50% discrepancy shown for the 10- μm run whose results are shown in Fig. 3-3.

However, when the longitudinal zone dimension fell below the transverse zone dimension, the discrepancy between LASNEX and theory was not reduced by further reduction in the longitudinal zone dimension. Similarly, when the transverse zone dimension was less than the longitudinal, LASNEX accuracy was not increased by finer transverse zoning. Hence, for this test problem, initially square zones maximized LASNEX accuracy at minimal computational cost. The fractional error in the first-order entropy (the error in the persistent entropy mode) was found to be $\sim 0.8 \cdot \max(k\Delta z, \Delta r)$, where k is the transverse wave number, Δz is the longitudinal zone dimension, and Δr is the transverse zone dimension. The fractional error in the amplitude of the first-order shear wave behind the shock front was about 4/5 as large as that in the first-order entropy wave.

These results are directly applicable to realistic calculations of the Richtmyer-Meshkov instability. In some realistic calculations, the heat transport behind the shock is large enough to damp the spurious high-frequency ringing. It is also possible to damp the ringing with a linear artificial viscosity, although extreme care must be taken that this does not also damp first-order perturbations of interest. For more complex situations, the results of the tests reported here can serve only as a guide for estimating the accuracy of 2-D LASNEX calculations of small-amplitude perturbation growth.

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Design of an Exploding-Foil X-Ray Laser Target

In 1984 the Target Design Program's expertise in the modeling of the interaction of high-power lasers with plasmas contributed to an ongoing effort to produce the world's first laboratory x-ray laser. Our theoretical and experimental efforts led to success this past summer. This article supplements the preshot theoretical approach and postshot analysis of Ref. 8.

The lasing scheme we attempted is the collisional excitation of neon-like selenium. The $j = 0\ 2p^33p$ levels (see Fig. 3-4) are populated directly by collisions from the ground state, and the $j = 2\ 2p^33p$ levels are populated by cascades from higher levels. Dipole radiative decay from the $3p$ levels to the ground state is forbidden, while the $2p^33s$ levels radiatively decay very rapidly, leading to a $3p\text{--}3s$ population inversion and possible gain.

The target we proposed is a thin foil ($\sim 750\ \text{\AA}$) of selenium ($Z = 34$) deposited on a minimal amount (1500 $\text{\AA}$) of Formvar ($\text{C}_{11}\text{H}_7\text{O}_2$) for structural integrity. The Novette 2w beams irradiate the film from both sides in a $1 \times 0.02\ \text{cm}$ line focus, with an intensity of $5 \times 10^{15}\ \text{W/cm}^2$ in a 450-ps (FWHM) Gaussian pulse. The foil explodes as the laser burns through it, creating a roughly cylindrical plasma (200- μm diam) of nearly uniform electron density. The smooth density profile allows the x-ray laser beam to proceed straight down the long line focus direction and stay in the

Fig. 3-4. Simplified energy level diagram [$\{l_1, l_2\}$ notation] for neon-like selenium. Only 8 of the hundreds of levels included in the calculations are shown.

high-gain region. The electron temperature of 1 keV (due to laser heating) and the electron density of mid- 10^{20} cm^{-3} ionizes a substantial fraction of the selenium to the neon-like state. The fraction is high enough to lead to substantial gain (as discussed below), but low enough that the optical depth along the 200- μm transverse dimension is low enough to avoid the trapping of 3s-2p radiation, which would repopulate the 3s lower level and destroy the gain.

The two main design tools were the 2-D hydrodynamics code LASNEX,¹⁰ which simulates the laser-foil interaction, and the 2-D x-ray laser code XRASER,¹¹ which uses the hydrodynamics output from LASNEX (densities, temperatures, velocities, etc.) along with detailed atomic physics data to

calculate atomic level populations and gain.

The design goal of producing a plasma with an electron density gradient scale-length L of at least 100 μm was motivated by simple considerations. Consider a plasma with the linear density profile

$$n_e = n_0 [1 - (x/L)] \quad (9)$$

and obeying the usual dispersion relation

$$\omega^2 = \omega_p^2 + k^2 c^2 \quad (10)$$

An x ray initially propagating in the x direction in plasma of density n_0 has, by Snell's law, a constant k_x given via Eqs. (9) and (10) by

$$k_1 = (\omega/c) (1 - d)^{1/2} \quad (11)$$

where

$$d \equiv \omega_{pe}^2 / \omega^2 \quad (12)$$

For $n_0 = 10^{21} \text{ cm}^{-3}$ and a 50-eV x ray, we find $d = 1/2500$. The group-velocity equation

$$V_g = c^2 k / \omega \quad (13)$$

leads us via Eq. (11) to the following expression for the x-ray trajectory:

$$x = ct(1 - d)^{1/2} \quad (14)$$

Since $k_y = (k^2 - k_1^2)^{1/2}$, Eq. (13), which leads to

$$dy/dt = c^2 k_y / \omega \quad (15)$$

along with Eqs. (9), (10), and (11), yields

$$y = c^2 t^2 d / 4L \quad (16)$$

or, using Eq. (14), with $d \ll 1$, to

$$y = x^2 d / 4L \quad (17)$$

Thus, the trajectory of an x ray starting in the x-direction at $x = 0$, $y = 0$ and refracted down the density gradient is a parabola. For a typical 10-mrad spectrograph acceptance angle (a value comparable to the expected geometric divergence of the x-ray laser tube), an $x = 1 \text{ cm}$ line focus implies $y_{\text{max}} = 0.01 \text{ cm}$. Equation (17) then implies $L_{\text{min}} = 0.01 \text{ cm}$. In general,

$$L = \frac{(x/1 \text{ cm})(v_0/10^{21} \text{ cm}^{-3})(100 \mu\text{m})}{(\text{divergence}/10 \text{ mrad})(\text{XRL photon energy}/50 \text{ eV})^2} \quad (18)$$

For an exponential profile, $n = n_0 \exp(-y/L)$, Eq. (13) leads to

$$dy/dt = cd^{1/2} [1 - \exp(-y/L)]^{1/2} \quad (19)$$

which yields

$$\frac{y}{L} = -\ln \left[1 - \tanh^2 \left(\frac{xd^{1/2}}{2L} \right) \right] \quad (20)$$

for the ray trajectory. For the same conditions as before, Eq. (20) implies $L_{\text{min}} = 0.009 \text{ cm}$, essentially the same as for a linear density profile.

A 1-D LASNEX simulation of an exploding foil, coupled with the XRASER postprocessor SPECTRE, which numerically solves for photon ray trajectories, graphically shows the refraction problem. Figure 3-5(a) shows the foil density profile and the resulting ray refraction before the driving laser pulse has burned through the foil. The steep density gradient causes severe refraction. After burn-through [Fig. 3-5(b)], the density profile, flat over about 0.01 cm, cures the refraction problem.

Experiments were performed at KMS Fusion to test the LASNEX modeling of the exploding-foil target and to demonstrate that a flat density profile is indeed produced. In Ref. 8 we discussed the good agreement between the density profile from a 1-D LASNEX simulation and the axial density profile inferred from a holographic interferogram made at KMS. In Fig. 3-6 we show similar agreement; from two other shots, taken with short probing laser pulses relatively early in the main driving pulse. This figure shows the ability of LASNEX to track the evolution of the exploding foil.

Fig. 3-5. (a) Density profile and x-ray refraction before the driving laser pulse has burned through the foil. (b) Density profile and x-ray refraction at the peak of the laser pulse. Foil has burned through and refraction problem has been cured.

Fig. 3-6. Density profiles in selenium/Formvar targets at various probing times. Agreement of LASNEX prediction with data is good.

Fig. 3-7. 2-D LASNEX density profiles (solid lines) agree well with observations at $t_{\text{probe}} = 450$ ps.

Figure 3-7 shows a 2-D LASNEX profile that is in excellent agreement with the 2-D data.

In Ref. 8 we also demonstrated good agreement between the modeling and various data with regard to the electron temperature ($T_e \sim 1$ keV) and the ionization balance (20 to 40% of the selenium is in the

near-n-like state). Given the correct densities, temperatures, and ionization balance, we can calculate and verify the most important quantities: the gains of the various transitions. Figure 3-8 shows gain isocontours near the peak of the driving laser pulse. Only the selenium part of the selenium/Formvar foil is displayed—hence the

Fig. 3-8. Gain isocounts at (a) 183 Å ($j = 0 \rightarrow 1$), (b) 206 Å ($j = 2 \rightarrow 1$), and (c) 209 Å ($j = 2 \rightarrow 1$). Gain along successive contours decreases by a factor of 0.8 from values on central contours.

Fig. 3-9. Predicted and observed time history of X-ray laser line emission at 209 Å.

apparently asymmetric plasma. Figure 3-8(a) shows the gain at 183 Å, which was predicted to have a maximum value of about 10 cm^{-1} . Figures 3-8(b) and 3-8(c) show gains for the transitions at 209 and 206 Å of about 3 to 4 cm^{-1} .

The observation⁴ of gain of $5 \pm 1 \text{ cm}^{-1}$ for both the $j = 2 \rightarrow 1$ lines lends considerable credence to the entire complex modeling chain, since the observed value is within 50% of predictions.

The modeling also successfully predicted the time history of the lasing signal. Figure 3-9 shows the theoretical signal (normalized to match the measured gain), which agrees well with the observed signal. The signal rises with the fraction of neon-like selenium, which increases as the temperature rises as does the electron excitation into higher levels, eventually cascading into the $j = 2$ upper laser state. The signal falls

as the foil decompresses, because of lower gain as the foil density decreases and because of decreased excitation as the foil cools.

Despite many speculations,⁸ the $j = 0$ to $j = 1$ line, whose gain was at least an order of magnitude lower than the predicted 10 cm^{-1} , is still a mystery. We are similarly puzzled by the absence of gain for fluorine-like 3p-3s transitions and for other neon-like (3-3, 2s-hole) transitions for which gain was predicted.

Thus, while our modeling achieved considerable success in predicting observed densities, temperatures, ionization states, and even gain on the observed lasing transitions, we need much deeper understanding of the experimental results to give us full confidence in designing laboratory X-ray lasers with shorter wavelength.

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Design of a Photoresonantly Pumped Fluorine X-Ray Laser

Introduction. In this article we describe recent design work on photoresonantly pumped X-ray laser (XRL) targets, done primarily in connection with the 1984 Novate XRL experiments. We also summarize preliminary results from those experiments, and discuss the analysis of experiments with scaled-down targets, intended to test the basis for the photoresonant pumping scheme.

Initial design studies for XRLs powered by large optical lasers, such as those used by the LLNL ICF program, were carried out by Hagelstein.^{12,13} Recent experimental work centered on the Shiva and Novette lasers, as described by Matthews and Kauffman.¹⁴ A recent general review of the field is given by Hagelstein.¹⁵

Overview of the Scheme. Photoresonance schemes are based on wavelength coincidence between two atomic transitions. Radiation emitted by a thermal process in one line can excite a lasing transition arising from the upper state of the second line. In the present case an L-shell line of beryllium-like manganese ($23+$) excites the Ly β line ($n = 1$ to $n = 3$) of hydrogen-like fluorine ($8+$). Of several previously suggested coincidences,^{13,16} this scheme is the only survivor of high-resolution spectroscopy¹⁷ done at KMS Fusion in early 1984. The wavelength mismatch between the manganese and fluorine transitions was found to be less than 1.5 mÅ, within the expected Doppler width of the pumping line and close enough for resonance excitation.

To use the line coincidence efficiently, it is desirable to create the manganese radiation in a hot ($T \sim 300$ eV) region in which collisional excitation is rapid, but to locate the fluorine in a separate, cold ($T \sim 20$ eV) region so that collisions do not destroy the population inversion. For this reason, we use a sandwich design for the target, with two laser-irradiated manganese foils ("flashlamps") on the outside and a gas cavity containing the fluorine on the inside.

Lasing is sought in the two 3d-2p transitions of hydrogen-like fluorine at 81 Å (153 eV). This is the shortest wavelength of the three schemes recently pursued at LLNL^{18,19} (see also the articles "Collisional Excitation Laser Scheme" and "Exploding Foil Recombination Laser," in Sec. 5). The transitions and energy levels involved are shown in Fig. 3-10. Kinetics data, from the baseline simulation at 400 ps (see discussion below and Table 3-1), are also shown. Following the resonance pumping, population is transferred to the 3d state by electron collisions, which are quite rapid owing to the small inelasticity (<1 eV). The population

Fig. 3-10. Hydrogen-like fluorine energy levels and laser scheme. Major transitions are indicated by wide arrows. Level populations and net flow rates are also given (see text for units and further details).

Note: Level populations (in parentheses) are in 10^{14} cm^{-3} . Broad arrows indicate primary processes; narrower, heavy arrows indicate minor processes. Net flow rates are indicated next to these arrows. Some one-way flow rates are also shown, with short arrows indicating flow direction. All flow rates are expressed as percentages of the 1s-3p pump rate, $1.1 \times 10^{26} \text{ cm}^{-3}\text{s}^{-1}$.

inversion hinges on the fact that the 3d-2p radiative decay rate is significantly less than 2p-1s rate, thereby allowing the lower laser level to empty quickly.

The physical design for the target, shown in cross section in Fig. 3-11, evolved from earlier XRL experiments on Novette. The fabrication technique is described in the article "X-Ray Laser Design and Fabrication," in Sec. 4. The operation is as follows: the manganese foils are irradiated by short (200 ps), intense ($\sim 2 \times 10^{14} \text{ W cm}^{-2}$) pulses from the two Novette beams. This creates, on each foil, a hot corona ($T_e \sim 2 \text{ keV}$) and an advancing conduction front, in which x rays are produced. This x-radiation,

includes lines and continuum from a range of manganese ionization stages; in particular, it includes the desired pumping line. The fluorine is initially gaseous (in inert form, as CF_4 or SF_6) to yield the proper electron density at lasing time, and is held in a cavity that is hydrodynamically isolated from the flashlamps. The x rays from the flashlamps flood the target, since the membranes and the gas are fairly optically thin ($\tau < 0.1$). They heat the fluorine, strip some of it to the hydrogen-like stage (primarily by K- and L-shell photoionization), and pump the $n = 3$ level of the resulting ions. The silicon knife edges extending into the gas cavity are intended to absorb some of the fluorine Ly α photons, thereby decreasing the trapping caused by repopulation of the lower laser level by this radiation. During laser irradiation, the initially solid parts of the target expand hydrodynamically as a result of heating by laser-energy deposition, conduction, and x-ray absorption. The expansion of the cavity walls could be detrimental if the lasing region is disturbed before the fluorine is sufficiently stripped or before the pump reaches its maximum strength. To delay expansion, we inserted a gap between the cavity and the flashlamp. The motions of the gas-cavity walls and membranes are slower than that of the flashlamp and are therefore of less concern. A trade-off in this design was a decrease in x-ray intensity in the fluorine cavity, resulting from the greater distance to the flashlamp.

Table 3-1. Optimal target parameters and simulation results.

Parameter	Value
Laser pulse width (FWHM)	200 ps (Gaussian profile)
Time of pulse peak in simulation	300 ps
Laser energy	600 J (for each of 2 beams)
Laser focus dimensions	100 $\mu\text{m} \times 1 \text{ cm}$
Laser intensity on target	$2 \times 10^{14} \text{ W cm}^{-2}$
Manganese flashlamp thickness	1300 Å
Si_3N_4 membrane thickness	1500 Å
Flashlamp-cavity gap	50 μm
Cavity width	100 μm
Knife-edge gap	65 μm
Knife-edge angle	55°
Fluorine density	$5.5 \times 10^{18} \text{ cm}^{-3}$
Pump line strength	0.003 photons/mode
Conditions in fluorine cavity at 400 ps:	
Electron density	$4 \times 10^{19} \text{ cm}^{-3}$
Electron temperature	20 eV
Ion temperature	25 eV
f (hydrogen-like fluorine)	0.15
f (hydrogen-like fluorine in 3d)	0.006
Velocity of imploding walls	$1.5 \times 10^7 \text{ cm/s}$
Gain in 3d-2p transition	1-6 cm^{-1}

Fig. 3-11. Major components of the photoresonance XRL target. The lasing path (perpendicular to the figure) is 1 cm long.

Design Issues and Numerical Simulations of Target Performance. Our goal was to maximize the predicted XRL gain coefficient, given the capabilities of the Novette laser and the materials and techniques used in target fabrication. We also attempted to ensure that the design was not overly sensitive either to known sources of variation, such as fabrication tolerances and beam spatial uniformity, or to uncertainties in the basic physics, such as the details of the fluid motions.

We relied on time-dependent simulations using the LASNEX and XRASER computer codes to model the target performance. Our interest centered on predictions for the gain in the lasing transitions, the hydrodynamic behavior of the target, and the output spectrum between 65 and 120 Å, which was seen by the primary diagnostics looking down the lasing channel (see the article "Resonant Photoexcitation Laser Experiments," in Sec. 5, for a description of the experimental setup). The simulations were mainly plane-parallel and one-dimensional. In addition, we found it useful to perform a large number of one-zone ("zero-dimensional," or 0-D) level-kinetics calculations to model the central region of the target, since such calculations require much less computer time. Some work on 2-D simulations was also done.

The first step in the modeling procedure is a complete simulation of the target behavior with LASNEX,³¹ including laser absorption, radiation emission and absorption, and hydrodynamics. Since LASNEX has only approximate (screened-hydrogenic, average-atom) atomic models, it cannot treat the level kinetics and energy balance of the lasing medium accurately. Therefore, we go through a second simulation step with XRASER,¹² which uses detailed models for individual ionization stages and calculates the line transfer required to account for trapping effects. XRASER uses a discrete-ordinate (S_N) method, including partial redistribution and bulk Doppler shifts for the line transfer. In the XRASER pass we treat only the fluorine region, and use the continuum radiation field and hydrodynamic quantities (mass density, velocity, and ion temperature) calculated by LASNEX and stored for each spatial zone at pre-chosen time intervals.

We used an analytic formula to account for the 2-D geometrical dilution in the

ionizing radiation as it propagates from the flashlamp into the gas cavity. This correction is applied to the radiation field used by XRASER, but is not used in the LASNEX simulation. We treat the pumping line strength as a parameter, whose value is taken from the 1983 (Ref. 21) and 1984 Novette experiments. Measured values of this line strength range between 1.7×10^{-3} and 1.1×10^{-2} photons/mode for a single flashlamp, at the irradiated side. We expect about half of this radiation to penetrate through the flashlamp to the side facing the gas cavity. We include the following processes and their inverses in the kinetics: photoionization (assuming Auger ejection when allowed), bound-bound radiative excitation, and collisional ionization and excitation. Details of the rate coefficients and numerical techniques used in XRASER are being prepared for publication by Hagelstein.

Table 3-1 summarizes the optimal target parameters and the key predicted quantities. The main physical considerations in the optimization are briefly discussed here. To maximize the fraction of hydrogen-like fluorine, we used long (200 ps) laser pulses and high intensities. The flashlamp thickness was chosen to maximize the amount of ionizing radiation reaching the lasing cavity by balancing early burn-through, with thin foils, against x-ray reabsorption, for thicker foils. Use of a long pulse necessitated a large gap between the flashlamp and the gas seals, and a wide laser cavity; both measures help ensure that the lasing medium is not catastrophically crushed before the time of peak gain. Figure 3-12 shows the hydrodynamic behavior as calculated in a 1-D LASNEX simulation along a horizontal cut through the center of the target, oriented as in Fig. 3-11. Approximately half of the initial cavity volume remains undisturbed at 400 ps, when we expect the peak product of pump intensity and hydrogen-like fraction. Further simulations showed that the design is hydrodynamically safe (in the sense that there is still a satisfactory undisturbed cavity volume) with the 50% variation in laser intensity possible along the line focus. We also found, from 1-D simulations on a diagonal cut, that expansion of the silicon knife edge during target irradiation should not significantly disturb the central region of the

Fig. 3-12. Hydrodynamics of XRL target as simulated by LASNEX. Selected mesh lines are plotted vs time. Above is the laser pulse shape.

cavity. Preliminary results of 2-D simulations confirm this expectation.

Finally, we optimized the fluorine density in the cavity to balance the increase in gain resulting from increased hydrogen-like fluorine density, and better collisional transfer, against the decrease in gain as trapping becomes critical in destroying the inversion ratio. Figure 3-13 shows the calculated gain vs fluorine density, with and without trapping. Without trapping, the nonlinearity is due mainly to collisional effects. At low densities, collisions are insufficient to transfer the population from $3p$ to $3d$, while too high a density leads to collisional depopulation of the $n = 3$ levels. The effect of trapping

in the 1-D simulations is dramatic. The effect increases with density as the optical depth in $Ly\alpha$ increases, and ultimately sets the upper limit on the density. The 1-D simulations are expected to give an upper limit on the trapping effect. Initial 2-D calculations indicate that the gain lies about halfway between the 0-D and 1-D results.

Pillbox Experiments. To test some of the basic processes of the XRL scheme, we designed scaled-down, 0.1-cm-long targets having a single flashlamp. These targets have rectangular cross sections, and we refer to them as pillboxes. The transparent side window of the gas cavity was viewed by a streak crystal spectrometer that

covered the main ground-state lines of helium-like and hydrogen-like fluorine. The experiments were performed in May 1984 at KMS fusion (again see the articles "Collisional Excitation Laser Scheme" and "Exploding Foil Recombination Laser," in Sec. 5). Targets were shot using both manganese and nickel flashlamps to look for a differential effect that could be ascribed to resonance pumping.

For short (160-ps) laser pulses, the emission from both hydrogen-like and helium-like fluorine was weak; for the 200-ps shots, strong emission lines from helium-like fluorine and weaker emission from the hydrogen-like stage were seen. Both observations suggest that the ionization state of the gas was lower than predicted by the 1-D modeling. We interpret this as arising from the dilution of the flashlamp radiation as it propagated through the gas cavity. This is essentially a 2-D effect, which we included in subsequent modeling. The 200-ps manganese-flashlamp shots show fairly strong emission in fluorine Ly α and Ly β , and emission in several helium-like lines. The nickel-flashlamp shots, on the other hand, show no emission in Ly β , but moderate emission in Ly α and strong emission in the helium-like lines. The difference in the Ly β flux between the manganese and nickel shots suggests that (as expected) resonance pumping occurs with the manganese flashlamp but not with the nickel flashlamp.

We modeled in detail two shots in which particularly good data were obtained. For the manganese shot, the observed Ly β /He α ratio could be matched by a simulation with a pumping line strength of 2×10^{-3} photons/mode at the flashlamp. The nickel data are consistent with a simulation with no pump. In both cases the other calculated line ratios are within a factor of 3 of the observed values. Although some alternative excitation mechanism for the Ly β line may exist, resonance pumping by manganese seems the most likely.

Several interesting points can be made concerning the other observed lines. In the simulations, the main excitation mechanism for the He α line is by K-shell photoionization of lithium-like fluorine. There also seems to be significant excitation of Ly α in the nickel target by K-shell ionization of the triplet $n = 2$ levels of helium-like fluorine. A feature appears in the data on the

Fig. 3-13. Predicted gain of the fluorine-based XRI target vs initial fluorine density.

long-wavelength side of the He α line, which we attribute to K α fluorescence of lithium-like fluorine following photoionization of the beryllium-like stage. Thus, besides seeing evidence for the resonance pumping of hydrogen-like fluorine by manganese, which is the basis for our laser scheme, we also saw a series of photoionization-excited lines that are characteristic of a low-density radiation-dominated plasma.

Preliminary Results from 1984 Novette Experiments. During the 1984 Novette XRI experiments, 10 photoresonance targets were shot (again see "Collisional Excitation Laser Scheme" and "Exploding Foil Recombination Laser," in Sec. 5). Only four of these shots produced good data from both of the primary diagnostics looking down the lasing channel. Two shots show evidence (in the streak spectrometer) for an early line at 81 Å. In one case the line is particularly strong, perhaps indicating

amplification. An absolute sensitivity calibration, in progress in early 1985, will shed further light on this question. Several complications in the interpretation of these data arise from the unreliability of occurrence of this line, and from possible contamination by a second-order helium-like carbon line at 80.6 Å, unresolvably close to 81 Å. Several targets were shot using SF_6 in the lasing cavity, rather than CF_4 , in an attempt to eliminate carbon contamination. In the shot with good data, no early line was seen; but a null shot, using CF_4 with a chromium flashlamp, did not show an early line either. There is no clear conclusion from this study as to whether the early line seen on the two of the Mn- CF_4 shots was from the fluorine laser transition or from carbon. The gated diagnostic generally saw lines from carbon in second order even on the SF_6 shots. Thus there appears to have been some other, as-yet unidentified source of carbon contamination in the target. It is not clear why the two spectrometers have seemingly discrepant spectra from these targets.

Conclusions. We have designed a photoresonantly pumped XRL for which we predict gains of approximately 3 cm^{-1} at 81 Å. Such a scheme should work with a laser such as Novette. Improvements in the design are anticipated, particularly in increasing the flashlamp brightness and minimizing trapping. We observed evidence for resonance pumping of fluorine by manganese in scaled-down, "pillbox" experiments at KMS. There is evidence from the 1984 Novette experiments for detection of the lasing line, with positive gain. However, due to difficulty of fabricating and fielding these targets, the preshot predictions of lower gain compared to the other schemes, and the more encouraging results then being obtained with the exploding-foil, collisional-excitation scheme, only a few shots were actually made, and the data were not reproducible from shot to shot. Therefore, we cannot draw a definitive conclusion as to the success of this scheme in making an XRL.

Although this scheme is more complicated than the successful collisional-excitation scheme, it deserves further work: it would lase at much shorter wavelength,

and it is based on different physical processes that may scale better to still shorter wavelengths.

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Richtmyer-Meshkov Instabilities in ICF Capsules

Introduction. The Richtmyer-Meshkov (RM) instability occurs in systems subjected to shocks. As a shock wave crosses the interface between two fluids, small perturbations at the interface begin to grow and, eventually, lead to mixing of the two fluids; or, if one of the fluids is a shell of finite thickness, perturbation growth results in the breakup of the shell. These effects are similar to those caused by the Rayleigh-Taylor (RT) instability, which occurs in systems undergoing constant acceleration. Since ICF capsules undergo both shocks and constant accelerations, designers must take both RT and RM instabilities into account.

Earlier, we developed²²⁻²⁵ an analytic theory of RT instabilities in an arbitrary number N of stratified fluids. This article summarizes our work in 1984, during which we extended our theory to include RM instabilities and to provide a uniform analytic treatment of both types of instabilities.²⁶

General Evolution Equations. For the RT instability we derived^{24,25} the equation

$$\eta_i(\tau) = \sum_{j=1}^{N-1} \sum_{l=1}^{N-1} W_{(l,i)} W^{-1}_{(l,j)}(L_i) \cdot \left[\eta_j(0) \cosh(\gamma_j \tau) + \frac{1}{\gamma_j} \dot{\eta}_j(0) \sinh(\gamma_j \tau) \right], \quad (21)$$

which describes the evolution with time τ of a perturbation of amplitude η_i at interface i . The summation over l indicates that all eigenmodes (of which there are $N-1$) contribute to the evolution, while the summation over j indicates that initial conditions at all $N-1$ interfaces contribute to $\eta_i(\tau)$. This results in what we have called

feed-through: surfaces with zero initial perturbation will eventually become perturbed by virtue of non-zero initial perturbations on other surfaces.

The corresponding equations in the case of the RM instability are obtained by treating a shock as an instantaneous acceleration, i.e., $g = \Delta v \delta(\tau - \tau_s)$, where Δv is the jump in velocity induced by a shock arriving at $\tau = \tau_s$. In this case the perturbations evolve according to²⁰

$$\eta_i(\tau) = \eta_i(0) + \eta_i(0)\tau + \Delta v \sum_{l=1}^{\infty} \sum_{j=1}^{\infty} \Gamma_{ij}^2 W(l,j) W^{-1}(l,j) \times [\eta_i(0) + \eta_i(0)\tau_s + \tau_s(\partial_t \tau - \tau_s)] \quad (22)$$

The notation is the same as in Eq. (21), with the addition of $W^{-1} = \gamma_i^2/g$. As in Eq. (21),

all modes and all interfaces contribute to the evolution of perturbations at any given interface.

The Classical Case ($N = 2$). In this case of two infinitely thick fluids there is only one interface, at which perturbations evolve according to

$$\eta(\tau) = \eta(0) + \eta(0)\tau + \Delta v \left(\frac{\rho_2 - \rho_1}{\rho_2 + \rho_1} \right) \times [\eta(0) + \eta(0)\tau_s + \tau_s(\partial_t \tau - \tau_s)] \quad (23)$$

Here $k = 2\pi/\lambda$, where λ is the wavelength of perturbations, and ρ_1 and ρ_2 refer to the densities of the two fluids; the shock moves from ρ_1 to ρ_2 at time τ_s .

From Eq. (23), it is clear that if the shock proceeds from a low-density fluid to a high-density fluid ($\rho_1 < \rho_2$), perturbations will grow linearly with time. In the opposite

Fig. 3-14. Possible histories of the RM instability amplitude for $N = 2$. Only cases (a)-(ii) are allowed if the shock proceeds from high to low density. Only cases (ii)-(o) are allowed in the reverse case.

case, where $\rho_1 > \rho_2$, the last term in Eq. (23) becomes negative, indicating that perturbations at the interface undergo a phase reversal. If the system continues to coast, the perturbations continue to grow, but with a phase opposite to their preshock phase.

Equation (23) is a generalization of Richtmyer's result,²⁷ in that it accommodates arbitrary initial conditions and can therefore treat two-shock systems. We have analyzed all possible evolutions of perturbations in two-shock systems, with the results shown in Fig. 3-14. There are 15 possible cases: the first 9, Figs. 3-14(a) to 3-14(i), are allowed if $\rho_1 > \rho_2$; the last 7, Figs. 3-14(i) to 3-14(o), are allowed if $\rho_1 < \rho_2$. The most interesting cases are those in Figs. 3-14(c), 3-14(g), and 3-14(n): in these cases the perturbation amplitude η is frozen out (does not grow) after the passage of the second shock. To achieve this freeze-out, the second shock must arrive at a specific time (see Ref. 26).

Single Shell with Two Free Surfaces. The simplest example of feed-through appears in the system shown in Fig. 3-15: a shell of thickness t and density ρ has perturbations of some initial amplitude η_0 on one free surface, while the other free surface has $\eta_0 = 0$. If a shock proceeds from the perturbed surface toward the smooth surface, the perturbations at each interface evolve according to

$$\frac{\eta_{\text{pert}}}{\eta_0} = 1 + \left(\frac{\Delta v k}{1 - e^{-2kt}} \right) (1 + e^{-2kt}) \tau \quad (24)$$

Fig. 3-15. The simplest case for $N = 3$: a fluid layer of density ρ and thickness t with two free surfaces. Initially one surface is perturbed and the other is smooth.

Fig. 3-16. Density profile for $N = 5$, with densities $(0, 1, 2, 4, 0)$ and thicknesses $(\infty, t/3, t/3, t/3, \infty)$.

$$\frac{\eta_{\text{smooth}}}{\eta_0} = 2 \left(\frac{\Delta v k e^{-kt}}{1 - e^{-2kt}} \right) \tau \quad (25)$$

Equation (25) shows how perturbations develop at a surface that initially had no perturbations. The feed-through is more effective if $kt \ll 1$, i.e., at longer wavelengths or, equivalently, in thinner shells. Feed-through occurs in both the RT and RM instabilities.

An Example Involving Shocks and a Constant Acceleration. ICF capsules have shells with continuous density gradients, which we can approximate by using large N . An example with $N = 5$ will illustrate our technique for handling both shocks and accelerations.

The density profile, shown in Fig. 3-16, is $\rho = (0, 1, 2, 4, 0)$; the units are arbitrary. We set $\lambda = 3t$. If the shell moves according to the velocity history shown in Fig. 3-17(a), the perturbations on the outer and inner surfaces evolve as shown in Fig. 3-17(b).

Figure 3-17(a) shows the shell at rest until $\tau = 1$, when a shock imparts a velocity $\Delta v = 1$. The units are immaterial; negative velocities indicate motion from outside to inside. From $\tau = 1$ to $\tau = 2$, the shell coasts. From $\tau = 2$ to $\tau = 4$, it accelerates inwards; then it coasts again until $\tau = 5$, when a strong shock, directed from inside to outside, brings it to rest. During this time the shell travels a distance equal to 12 times its thickness.

Figure 3-17(b) shows perturbations on the outside growing steadily until the final shock, when they reverse their phase almost instantaneously but continue to grow. Perturbations on the inner surface undergo phase reversal after the first shock. This is as expected, since at that surface the shock proceeds from $\rho = 4$ to $\rho = 0$. As the outer perturbations grow large they take control of the inner perturbations, and both surfaces ultimately develop perturbations 100 to 1000 times the initial amplitudes.

A number of other examples, treating different velocity histories and different assumptions about surface finish, are given in Ref. 26.

Summary. We have developed a theory to calculate how perturbations evolve in systems undergoing shocks or accelerations. The theory can handle arbitrary density profiles and arbitrary velocity histories. A

number of effects, particularly ablation and compression, are missing from this theory. By adding the ability to handle shocks, and the resulting RM instability, we have come closer to simulating actual ICF targets.

Author: K. O. Mikaelian

Charged-Particle Fusion

Introduction

In 1984 we looked at certain crucial elements of charged-particle-driven fusion. Some of our results could also be useful for laser drivers. In particular, we have

- Outlined an alternative strategy for achieving illumination symmetry in directly driven targets.
- Begun studies on target concepts that could give useful target gain at input energies as low as ~ 1 MJ.
- Made progress towards an in-principle solution of the problem of beam compression in heavy ion fusion.

Many charged-particle fusion schemes permit near-spherical target illumination by incident particle beams, but exhibit residual asymmetries centered around an axis determined by the particular driver or reactor scenario. In such illumination geometries, it should be possible to vary the angles of beamlet incidence on spherical targets and, under some circumstances, the beam currents (or energy fluxes) themselves, to eliminate all long-wavelength (low-index spherical harmonic) asymmetries down to specific values. Some, but not all, shorter-wavelength asymmetries can also be eliminated by this method, which is described in the following article, "Reduction of Illumination Asymmetries in Directly Driven Ion Beam and Laser Targets"; this is also the case for spherically symmetric ("Platonic solid") beamlet placement. Elimination of all short wavelength asymmetries to a level low enough to permit spherical target implosion would now depend on effects such as transverse beam profile, dispersions in particle (photon) incident angles, particle energy (laser frequency), and thermal effects of target plasmas. One advantage of our strategy is that by concentrating on the variation of beamlet placement and beam

Fig. 3-17. (a) Velocity history of a shell whose density profile is shown in Fig. 3-16. Negative velocities are directed from outside to inside. (b) Evolution of perturbations at the shell surfaces. At both surfaces, perturbation has initial amplitude $\eta(0)$ and wavelength $\lambda = 3t$.

currents (fluxes) for suppression of long wavelength asymmetries, we could eliminate more such asymmetries than with Platonic-solid placement. And this could be done without using up some of the free parameters in the other smoothing physical effects mentioned above, thus retaining freedom to use those effects exclusively for the reduction of shorter-wavelength asymmetries.

Even where the driver scenario does not lead to a preferred axis in the illumination scheme, our results might still be useful if we adopt a similar strategy for using only beam placement and beam-current variation to eliminate the long-wavelength asymmetries.

For heavy ion fusion, final longitudinal compression of the particle beam bunch before the beam is focused onto the target is envisaged as one means to increase peak power by an order of magnitude or more.

Initial studies of this process raised concern regarding the distortion of longitudinal phase space by the radial dependence of the induced longitudinal electric field. Our more detailed studies, described in the article "Studies of Beam Compression Including Transverse Effects for a Proposed Heavy-Ion Fusion Test Facility," include comparisons of one-dimensional and two-dimensional models and show that the effect is not as great as once predicted.

Progress towards an in-principle solution to the beam compression problem should be continued if we are to obtain the lowest-cost final system.

Author: J. W-K. Mark

Reduction of Deposition Asymmetries in Directly Driven Ion-Beam and Laser Targets

We have developed a procedure for reducing energy-deposition asymmetry in spherical targets driven directly by ion or laser beams. This work is part of a strategy for achieving illumination symmetry in such targets, which we propose as an alternative to those in the literature.²⁸⁻³³ This strategy allows an axially symmetric placement of beamlets, which would be convenient for some driver or reactor scenarios. It also allows the use of beam currents or energy fluxes to help reduce deposition asymmetry.

A motivation for our work is that many charged-particle fusion schemes allow a near-spherical target illumination, but leave residual asymmetries centered around an axis that is determined by the driver or reactor scenario. For example, this occurs in the PBFA pulsed-power facilities at Sandia, Albuquerque (and in related facilities), and in heavy ion reactor concepts such as the German HIBALL. In both these cases, driver and reactor considerations discourage the use of beam ports directly above or below the fusion chamber. In some reactor chambers, such as those with vertically flowing liquid metals, it might simplify the design to have fewer vertical planes occupied by beam ports. Our strategy provides just such flexibility.

Even if no preferred axis arises in a particular illumination scheme, beam

placement and beam current can be used to eliminate long-wavelength (low spherical-harmonic) asymmetries in the deposition profile. Typically, our strategy eliminates more long-wavelength asymmetry per beamlet (in the form of coefficients of spherical harmonics) than does symmetric (Platonic-solids) beamlet placement.

Low-spherical-harmonic asymmetries are eliminated, up to values of l and $|m|$ that depend on the number of beamlets and the number of different beamlet strengths used. Many harmonics of higher l and $|m|$ are also eliminated, depending on what particular asymmetries arise from the choices of beamlet orientations and strengths. The convergence ratio of the target determines how much smoothing of higher-order terms is required. Our procedure shows what penalties must be paid, in terms of higher harmonics not removed, if the driver or reactor system makes it desirable to eliminate certain symmetries in placement of beamlets.

In general (even with symmetric beam placement), higher harmonics can be removed using physical effects such as transverse beam profiles,^{29,31} dispersion in beamlet angles, particle (or photon) energy, thermal effects of target plasmas, or focal shift relative to the target center.^{24,30,32} For laser drivers, additional techniques include refraction smoothing³² and smoothing due to decoupling of absorption and ablation layers. Because our strategy can eliminate more long-wavelength (low-harmonic) asymmetry than can symmetric beam placement, we can use the free parameters in the other physical effects exclusively against shorter-wavelength asymmetries.

Here are two examples of the smoothing of lower harmonics. We could eliminate all harmonics with $1 \leq l \leq 5$ and $|m| \leq 5$ by appropriate choice of beamlet angles with 18 beamlets of identical strength. With 32 beamlets, of two different strengths (in the ratio $\approx 15/8$), we could eliminate all harmonics with $1 \leq l \leq 7$ and $|m| \leq 7$. Convenient arrangements for several ion-beam and laser driven illumination schemes are described later in this article.

Our procedure consists of expanding the deposition profile of each beamlet in spherical harmonics and choosing the beamlet orientations (θ_{lm}, ϕ_{lm}) and amplitudes ϵ_{lm} so that harmonics of specific (l, m) vanish in

Fig. 3-18. Geometry used in deriving the results given in this article. The angle γ_n is measured between the axis of beamlet n and an arbitrary point on the target (r, θ, ϕ) .

the net deposition profile. We assume that all beamlets are directed at the center of the target. The same shape is assumed for each beamlet: only the beamlet strengths vary, so that we let (θ_n, ϕ_n) represent the orientation of the central axis of beamlet n . With the geometry of Fig. 3-18, (r, θ, ϕ) designates an arbitrary point on the target surface and γ_n represents the angle between the directions (θ_n, ϕ_n) and (θ, ϕ) . Assuming that the deposition profile $E_n(r, \theta, \phi, t)$ of beamlet n depends only on (r, γ_n, t) , we expand E_n as

$$E_n(r, \theta, \phi, t) = \epsilon_n \sum_{l=0}^{\infty} E_l(r, t) P_l(\cos \gamma_n) \quad (26)$$

where the $P_l(x)$ are Legendre polynomials. In accordance with our assumptions, the $E_l(r, t)$ are common to all the beamlets, while the γ_n and ϵ_n may vary. To sum over beamlets we use a particular definition (see,

e.g., Ref. 34) of spherical harmonics and the expansion of $P_l(\gamma_n)$ in the spherical harmonics $Y_{lm}(\theta_n, \phi_n)$ and $Y_{lm}(\theta, \phi)$: we thus obtain the net deposition profile

$$E(r, \theta, \phi, t) = \sum_{l=0}^{\infty} E_l(r, t) \sum_{m=-l}^l D_{lm} Y_{lm}(\theta, \phi) \quad (27a)$$

where

$$\begin{aligned} D_{lm} &= \frac{4\pi}{2l+1} \sum_n \epsilon_n Y_{lm}^*(\theta_n, \phi_n) \\ &= A_{lm} \left[\frac{4\pi}{2l+1} \frac{(l+m)!}{(l-m)!} \right]^{1/2} \quad (27b) \end{aligned}$$

Retaining some symmetry in the beamlet distribution would by itself cause many of the coefficients A_{lm} or D_{lm} to vanish.

Charged-Particle Fusion

Table 3-2. Symmetries for which particular A_{lm} coefficients vanish (we assume ϵ_n).

Symmetry	Condition: for every n , there exists an n' such that the relations below are valid	Result: $A_{lm} = 0$ unless quantity below is an integer
About origin	$\theta_{n'} = \pi - \theta_n$ $\phi_{n'} = \pi + \phi_n$	$l/2$
About $z = 0$ plane	$\theta_{n'} = \pi - \theta_n$ $\phi_{n'} = \phi_n$	$(l \pm m)/2$
Of rotation about z -axis ^a	$\theta_{n'} = \theta_n$ $\phi_{n'} = \phi_n + (2\pi/N_2)N_1$	$(m/N_2)N_1$
About $z = 0$ plane with rotation ^a	$\theta_{n'} = \pi - \theta_n$ $\phi_{n'} = \phi_n + (2\pi/N_2)N_1$	$(l \pm m)/2 + (m/N_2)N_1$

^aHere N_1 and N_2 are integers.

Table 3-2 summarizes the results of symmetry about the origin, about the $z = 0$ plane, rotational symmetry about the z -axis (over specific angles that are rational submultiples of 2π), and a combination of these symmetries. [In considering the results in Table 3-2, it should be kept in mind that the results for rotational symmetries—the last two entries in the table—arise from consideration of the interaction of *all* beamlets, not just two beamlets. This latter result represents a particular case of the result obtained in Eqs. (29) through (31), below.]

In what follows we illustrate some applications of these symmetries and additional helpful effects of choosing the θ_n such that the $\cos \theta_n$ are at the Legendre-Gauss numerical quadrature points.

To eliminate specific harmonics of low (l, m) , we assume that θ takes on the values θ_k , $k = 1, 2, \dots, K_0$, and that for each θ_k we have beamlets at angles $\phi = \phi_k + \phi_{jk}$, $j = 1, 2, \dots, J_k$. The total number of beamlets is then $N_b = \sum J_k$. Using properties of spherical harmonics,³¹ and doing some re-summing, we reduce Eq. (27) to

$$\begin{aligned}
 E(r, \theta, \phi, t) = & A_{00} E_0(r, t) \\
 & + \sum_{l=2}^{\infty} \sum_{m=-l}^l A_{lm} F_l(r, t) P_l(\cos \theta) \\
 & + 2 \sum_{m=1}^{\infty} \sum_{l=m}^{\infty} \sum_{n=-l}^l \operatorname{Re}(A_{l,m,n} e^{im\phi}) E_l(r, t) P_l(\cos \theta)
 \end{aligned} \quad (28)$$

Here $P_l^m(\cos \theta)$ is an associated Legendre function, and

$$A_{lm} = \frac{(l-m)!}{(l+m)!}$$

$$\times \sum_{k=1}^{K_0} \epsilon_k B_{mk} \exp(-im\phi_k) P_l^m(\cos \theta_k) \quad (29a)$$

$$B_{mk} = \sum_{j=1}^{J_k} \exp(-im\phi_{jk}) \quad (29b)$$

We assume that the ϵ_n depend only on k and not on j , so that ϵ_k replaces ϵ_n .

In Eq. (28), we have chosen the (θ_k, ϕ_{jk}) and the ϵ_k so that the A_{lm} vanish for $1 \leq l < 2K_0$ and $1 \leq m < J_0$. This is why the first sum (the $m = 0$ case) in Eq. (28) starts with $l = 2K_0$, rather than with $l = 0$, and why the second sum starts at $m = J_0$, rather than $m = 1$. That is, we have taken

$$\phi_{jk} = 2\pi(l_0 + j), \quad j = 1, 2, \dots, J_k \quad (30)$$

so that

$$\begin{aligned}
 B_{mk} = 0, \quad \text{if } m/J_{k1} \neq \text{integer} \\
 = J_k, \quad \text{if } m/J_{k1} = \text{integer}
 \end{aligned} \quad (31)$$

Then if $J_0 = \min(J_k)$, the A_{lm} vanish for $1 \leq m < J_0$. If all the J_k were equal, this same result would follow more directly from rotational symmetry, as can be seen from Table 3-2. Our derivation of Eq. (31) is similar to that of Schmitt.³¹

Our choice for the θ_k is obtained from the following considerations. For $m = 0$, we have $B_{0k} = J_k$, and

$$A_{00} = \sum_{k=1}^{K_0} w_k P_0(\cos \theta_k) \quad (32)$$

where $w_k = \epsilon_k J_k$. This sum vanishes for $1 \leq l \leq (2K_0 - 1)$ if we choose the $x_k = \cos \theta_k$, and the w_k to be the positions and

weights of a Legendre-Gauss numerical quadrature formula of order K_0 . [This quadrature is exact for polynomials of order $l \leq 2K_0 - 1$; this is implied by Eq. 25.4.29 of Ref. 35 and by the orthogonality relations of the Legendre functions (p. 774 of Ref. 35).]

We should consider both even and odd values of K_0 since both are convenient for driver concepts that allow beamlets to impinge from all lateral directions around the reactor chamber. In the simplest case of interest, for which $K_0 = 2$ and $J_1 = J_2 \geq 4$, deposition harmonics with $1 \leq l \leq 3$ and $|m| \leq 3$ will vanish provided that $\theta_1 \approx 54.74^\circ$ and $\theta_2 \approx 125.26^\circ$. Since $w_1 = w_2$ for $K_0 = 2$, the beamlets can all have the same ϵ_i ; this is useful for the PBFA or HIBALL (rf accelerator) geometries.

Figure 3-19(a) illustrates the beamlet placement for $K_0 = 2$, $J_1 = J_2 = 10$ (the HIBALL-like case). Note that having $J_1, J_2 > 2K_0 = 4$ does not help remove even-integer $l \geq 4$, but $m \neq 0$ deposition asymmetries of the form $P_l(\cos \theta)$. This is true despite the fact that many higher harmonics with $m = 0$ can be eliminated (we could take $J_1 - J_2$ to be an even integer and choose the ϕ_{0i} so that the first three symmetries of Table 3-2 hold: A_{lm} would then vanish unless l, m are even and m is an integer multiple of J_1). Thus this illumination generates more asymmetry in the θ direction than in the ϕ direction. Of course, if θ_1 and θ_2 were not at the Legendre-Gauss quadrature points, it could not even be guaranteed that the $l = 2$ harmonics would be minimized in this case.

On the basis of target-deposition symmetry alone (the point of view taken here), it

is preferable to have $J_0 \geq 2K_0$, but to have the two not too different from one another. Doing this as an alternative to the case of Fig. 3-19(a), we could obtain somewhat better deposition symmetry, i.e., the elimination of harmonics for $1 \leq l \leq 5$, all odd l , and $|m| \leq 5$, with only 18 beamlets. Such an illumination scheme could be achieved with $K_0 = 3$; $\theta_1 \approx 39.23^\circ$, $\theta_2 = 90.00^\circ$, and $\theta_3 \approx 140.77^\circ$; $J_1 = J_2 = 5$, $J_3 = 8$; symmetry about the origin; and $\epsilon_1 = \epsilon_2 = \epsilon_3$. The Legendre-Gauss weights in this case are in the exact ratio $w_2/w_1 = w_3/w_1 = 8/5 = 1.6$. With the same basic conditions (18 beamlets, $K_0 = 3$, and the same values of θ), we could eliminate all harmonics for $1 \leq l \leq 5$, $|m| \leq 5$, all odd l and m , and $m/5 = \text{integer}$ with $J_1 = J_2 = J_3 = 6$ and $\epsilon_2/\epsilon_1 = \epsilon_3/\epsilon_1 = 8/5 = 1.6$. This is illustrated in Fig. 3-19(b). Clearly other such combinations are possible.

For some driver scenarios, it might be convenient to arrange the beamlets into several wide bunches, directed at the target from two sides. Take the case $K_0 = 4$ and $J_0 = 8$, which eliminates deposition harmonics with $1 \leq l \leq 7$ and $|m| \leq 7$, as well as those shown for the first three symmetries in Table 3-2. The beamlets are arranged in four cone-like bunches with cone angles $\theta_1 \approx 30.56^\circ$, $\theta_2 \approx 70.12^\circ$, $\theta_3 \approx 109.88^\circ$, and $\theta_4 \approx 149.44^\circ$. The Legendre-Gauss weights satisfy $w_1 = w_4$, $w_2 = w_3$, and $w_2/w_1 = 1.8748$. With $J_1 = J_2 = 8$ and $J_3 = J_4 = 15$ (for a total of 46 beamlets), we note the apparent coincidence that $J_2/J_1 = 1.8750 \approx w_2/w_1$; thus for this case the $\epsilon_i J_i = w_i$ are the Legendre-Gauss weights without our even having to choose different ϵ_i .

Fig. 3-19. Three beamlet-placement schemes based on the strategy outlined in this article: (a) $K_0 = 2$, $J_1 = J_2 = 10 \cdot 2K_0$ (similar to the HIBALL reactor scenario); (b) fewer beamlets than in (a), but greater deposition symmetry; (c) two open bunches of beamlets from two sides; this gives still greater symmetry because of greater number of beamlets (32).

Alternatively, one could also accomplish the same result with only 32 beamlets by choosing $J_1 = J_2 = J_3 = J_4 = 8$ and $\epsilon_2/\epsilon_1 = w_2/w_1 \approx 15/8$; Fig. 3-19(c) illustrates this latter case. This comparison shows again the advantage (e.g., 32 versus 46 beamlets) for judicious use of different beamlet currents wherever convenient for the driver.

Further studies are needed to confirm that higher-order asymmetries can be kept to tolerable levels by the use of the various physical effects mentioned earlier in this article.

Author: J. W.-K. Mark

Studies of Beam Compression Including Transverse Effects for a Proposed Heavy-Ion Fusion Test Facility

For applications as an ICF reactor driver, a heavy-ion beam pulse must be longitudinally compressed by one to two orders of magnitude to achieve the peak power required to ignite a target.³⁰ This process, among others, will be tested³⁷⁻⁴¹ in a facility known as the "High-Temperature Experiment" in heavy-ion fusion. Beam compression is a critical element of an accelerator for heavy-ion fusion; it occurs primarily after the main phase of acceleration and before final focus onto target.

Here we examine the compression of a drifting heavy-ion pulse with the aid of

particle simulations. Space-charge forces play a vital role in halting the compression before the final-focus lens system is reached. Compression must take place with minimal growth of both transverse emittance and longitudinal momentum spread. Of particular concern here is the possibility of distortion of longitudinal phase space because of the radial dependence of the longitudinal self electric field.

We describe initial theoretical results for an in-principle solution to this problem. Further refinements are necessary before the least costly solution can be chosen.

Compression. An accelerated beam pulse has finite extent in space and time (in what follows, we take the z -coordinate as lying along the direction of beam propagation). Beam compression is initiated by firing the accelerating modules so that the tail end of the beam moves faster than the head. Especially when the variation of velocity is a linear function of distance (measured from the center of the pulse), this differential velocity can be viewed as a velocity "tilt" in longitudinal (z, v_z) phase space [see Fig. 3-20(a)]. At the end of beam compression, the tilt is removed and the compression is halted (and even reversed) by the buildup of longitudinal space-charge repulsion in the beam bunch (in the absence of space charge, the free-streaming velocity differentials would instead be the determining factor for minimum beam length). Of course, we would have to impose external fields on the compressed beam to keep it compressed for any

Fig. 3-20. The (z, v_z) phase space of the 1-D particle-code simulation (see text) at (a) $t = 0$ (initial velocity tilt), (b) $t = 0.7 \mu\text{s}$, (c) $t = 1.5 \mu\text{s}$ (final compressed state).

length of time. In practice, the beam bunch will be sent through the final focus lens system onto the target before expansion occurs. This entire process must be carefully controlled to produce the desired small spot on the fusion pellet. During compression, the transverse dimensions of the pulse must be controlled by alternating-gradient quadrupole focusing.

Apart from the decrease in pulse length, it is also necessary to observe the requirements on the final beam pulse necessary to focus it onto a target at a particular focal distance. For example, chromatic aberrations (which arise because different momenta have different focal lengths) caused by the final lens system impose a limit, which for rough illustrative purposes can be summarized as

$$r_{\text{spot}} > 2 \left(\frac{\Delta p_z}{p_z} \right)_{\text{lens}} (r_{\text{max,lens}}) \quad (33)$$

Here $r_{\text{spot}} \approx 0.2$ cm is the envelope radius of the desired focal spot, $(r_{\text{max,lens}}) \approx 10$ cm is the envelope radius at the final focus lens, and $\Delta p_z/p_z$ is the half width in the dispersion of z momenta (accumulated from all sources) at final focus. These typical numbers imply that $\Delta p_z/p_z < 0.01$.

A second condition is that the transverse emittance ϵ must be sufficiently small. To limit the effect of spherical aberrations in the final lens, the half-angle of the focal cone must satisfy

$$\theta_{\text{qph}} \lesssim 0.02 \text{ rad} \quad (34)$$

This translates into the condition $\epsilon \lesssim 4 \times 10^{-3}$ cm·rad, which provides another constraint on the resultant beam after final compression. [The emittance ϵ measures the phase area (divided by π) occupied by the beam in a typical transverse phase space ($x, dx/dz$).]

Simplified Analytic Treatment of Compression. A simplified analytic solution obtains if we adopt the model of a one-dimensional beam that has no longitudinal momentum spread Δp_z at any z before beam compression. The space-charge self field of the beam bunch is approximated by the electric field

$$E_z = -g \frac{\partial \lambda}{\partial z} \quad (35)$$

where $\lambda(z, t)$ is line charge density; g is a constant in the limit in which the longitudinal length scales of beam shape and disturbances are long compared with actual beam radius a or beam pipe radius b . In this regime g has the value

$$g \approx \frac{1}{4\pi\epsilon_0} \left(\log \frac{b^2}{a^2} + \frac{1}{2} \right) \quad (36)$$

In this idealized limit, the beam pulse obeys the equations

$$\frac{\partial \lambda}{\partial t} + \frac{\partial \lambda v_z}{\partial z} = 0 \quad (37a)$$

$$\frac{\partial v_z}{\partial t} + v_z \frac{\partial v_z}{\partial z} = -\frac{Ze}{m} g \frac{\partial \lambda}{\partial z} \quad (37b)$$

where $v_z(z, t)$ is beam internal velocity and Ze and m are beam-ion charge and mass.

An exact solution with a beam parabolic in λ is given by

$$\lambda(z, t) = \frac{\lambda_0 L_0}{L(t)} \left[1 - 4 \frac{(z - \bar{z})^2}{L^2(t)} \right] \quad (38a)$$

$$v_z(z, t) = \bar{v} + (z - \bar{z}) \frac{d}{dt} \ln L(t) \quad (38b)$$

$$K^{1/2} (t - t_0) = -L_0^{1/2} \left\{ \frac{1}{\xi} (1 - \xi)^2 + \tanh^{-1} (1 - \xi)^{1/2} \right\}^{\frac{1}{2}} \frac{L_0^{1/2} L(t)}{L_0^{1/2} L_0} \quad (38c)$$

where z and $\bar{v}$ are constants, and where $L(t)$ is determined by Eq. (38c) in terms of the initial ($t = t_0$) length L_0 and final length L_f . The constant K is given by

$$K = \frac{L_0 L_f}{L_0 - L_f} \left(\frac{dL}{dt} \right)_{t_0}^2 = 16 \frac{Ze g}{m} \lambda_0 L_0 \quad (38d)$$

and $Q = 2L_0 \lambda_0/3$ is the total beam charge.

This simplified solution provides a useful guide to an in-principle solution to the beam compression process provided that it survives under the realistic, multidimensional environment of actual pulses. For example, Δp_z is not initially zero, and the constant factor g is replaced by a radius-dependent term in actual beams, so that particles at different radii r are acted on by different longitudinal self-forces and

experience different evolution, resulting in further growth of Δp_z .

1-D Particle-Code Simulation. An estimate that g varies over r by $\pm 20\%$ results in the evolution of phase space (z, v_z) and line density $\lambda(z, t)$ shown in Figs. 3-20 and 3-21. These results were obtained using our one-dimensional particle-simulation code BLISS, which uses the field of Eq. (35). Although actual transverse particle motions average over actual spread of g with radius, this one-dimensional model assigns each particle a g -factor appropriate only to its initial position.

As expected, the initial momentum tilt of Fig. 3-20(a) is stopped in Fig. 3-20(c), but at the expense of some momentum spread $\Delta p_z/p_z$. The simulated problem corresponds to the situation typical of the High-Temperature Experiment, in which a $1.875\text{-}\mu\text{C}$, 125-MeV beam of singly ionized sodium will be compressed from a length of

10 m to a final length (according to our analytic model) of 1 m. Since the model shows that space-charge effects are important only in the last 2 m of compression, our simulation concentrates on this phase.

2-D Particle-Code Simulation. We used a 2.5-dimensional LLNL version of the particle code MASK⁴² to repeat the one-dimensional simulation, with realistic variation of g with r automatically included in the force calculations. The final momentum spread in this simulation (Figs. 3-22 and 3-23) is roughly one-fourth as great as in the one-dimensional simulation, with its artificially spread g -factor.

The two-dimensional simulation again uses a $1.875\text{-}\mu\text{C}$, 125-MeV beam of Na^+ ions. The beam radius has the value $a = 1\text{ cm}$, and the perfectly conducting pipe has radius $b = 2\text{ cm}$. The boundary conditions at $z = \pm \infty$ are approximated accurately by doubling the periodicity length

Fig. 3-21. Spatial line densities $\lambda(z, t)$ of particles in the 1-D simulation of Fig. 3-20 at (a) $t = 0$, (b) $t = 1.5\text{ }\mu\text{s}$.

Fig. 3-22. The (z, v_z) phase space of the 2-D particle-code simulation (see text) at (a) $t = 0$ (initial velocity tilt), (b) $t = 0.7\text{ }\mu\text{s}$, (c) $t = 1.4\text{ }\mu\text{s}$ (final compressed state).

Fig. 3-23. The spatial positions in (r, z) coordinates of 10% of the particles in the 2-D simulation of Fig. 3-22, at corresponding times. Note the different horizontal and vertical length scales.

of the simulation region in z , because the Green's function falls off exponentially. The transverse quadrupole focusing forces are approximated by the axially symmetric radial electric field

$$E_r = -5.266 \times 10^6 \frac{r}{b} \text{ V/m} \quad (39)$$

A total of 46 751 particles are used on a 20×512 grid, representing a region $2 \text{ cm} \times 2 \text{ m}$ in the r and z directions, respectively. The initial line density $\lambda(z, 0)$ is made nearly parabolic by the superposition of 20 constant-charge-density and constant-radius ($r = a$) components. Each such component is loaded with a Maxwellian velocity distribution with

$$\langle v_{ri}^2 \rangle = \frac{2i-1}{20} \theta, \quad i = 1, \dots, 20 \quad (40)$$

where θ characterizes the required $\bar{v}_{ri}^2$ at $z = 0$ for the parabolic distribution. Note that $(2i-1)/20 = 1$, so that loading according to Eq. (40) gives the correct $\bar{v}_{ri}^2$.

In Figs. 3-22 and 3-23, 10% of the particles are plotted. As the beam is compressed, an efficient use of the simulation grid is made by allowing the beam to expand in r while the space-charge self-fields increase to a dominant level, so that the beam radius almost reaches the pipe radius. This exaggerates the deleterious effects of radial variation in g compared with the realistic accelerator situation, in which the transverse focusing forces must vary in time (actually the beam travels into regions of increased focusing) in a way that allows the beam to maintain specific clearances from the pipe. Another conservative factor is that

we initiated the simulation with a near equilibrium rather than a true equilibrium, so that extra dispersions arise. Even with all these extra conservative factors, the final $\Delta p_z/p_z$ is smaller than estimated in the one-dimensional model.

Summary. Transverse beam spread need not be an impediment to the longitudinal beam compression needed in heavy-ion fusion to increase peak power by one to two orders of magnitude. Additional work is necessary to develop the lowest-cost engineering solution. In particular, the parabolic beam line density $\lambda(z, 0)$ introduced above (because of the existence of an analytical model for this form of λ) uses focusing forces inefficiently; a nearly flat $\lambda(z, 0)$ is preferable. Also needed are additional studies of multidimensional effects introduced by beam misalignments and by the accumulation of realistic jitters in acceleration forces. Effects of initial beam temperature in the longitudinal direction must also be introduced.

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Plasma Physics

Introduction

The Plasma Physics Group had a very productive year. Extensive collaboration with the ICF experimental program resulted in significant advances in the understanding of laser-plasma coupling. These advances include the characterization of Raman scattering in long-scalelength plasmas and the

demonstration of its correlation with hot-electron generation and its suppression by collisional damping. Most importantly, a very favorable scaling of the laser-plasma coupling with wavelength has been shown to extend to longer-scalelength underdense plasmas.

We also made significant advances in the theory and simulation of the coupling physics. We improved our understanding of many important processes, including the Raman instability at a local density maximum, Raman sidescattering, ponderomotive and thermal filamentation, electron-ion collisions in partially ionized plasmas, and an unexpectedly strong heating by high-intensity light. A Monte Carlo atomic physics model was developed and applied to improve the calculation of line emission in high-Z plasmas. An implicit 2-D electromagnetic particle code was implemented to expand our ability to simulate low-frequency kinetic phenomena in plasmas.

The development of reliable scalings necessitates a strong interaction between calculations and well-diagnosed experiments. Hence a large fraction of our effort was again devoted to the design and interpretation of coupling experiments. Detailed results of this collaboration with the Laser Experiments and Advanced Diagnostics program are reported under "Novette Experiments" in Sec. 5. An example will suffice to illustrate the progress. By use of LASNEX calculations to determine the plasma conditions, and plasma simulations to compute the Raman scattering, a series of experiments were designed to demonstrate the suppression of Raman scattering by collisional damping. The experimental results confirmed the predictions. In addition, the plasma density profile computed with LASNEX agreed well with the profile inferred from spatially resolved measurements of the bremsstrahlung, and the electron temperature in the underdense plasma predicted by LASNEX agreed well with the temperature inferred from the long-wavelength cutoff of the Raman signal. In short, we are obtaining an improved knowledge of the plasma conditions in laser-irradiated targets, and we are learning how to control the important coupling mechanisms more reliably.

Although some very important scalings have been successfully predicted, improved models are needed to optimize target designs. Hence we continue to develop theoretical and computational models of coupling physics; our work in this area is about equally divided between theory, computer simulation, and code development. Details of this work are presented in the following articles.

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Ponderomotive and Thermal Filamentation of Laser Light

As targets are irradiated with longer, more energetic pulses of laser light, longer-scalelength plasmas are produced. Filamentation is a potentially important process in such plasmas. In this instability, perturbations in the intensity profile of an incident light beam grow in amplitude, causing the beam to break up into intense filaments. The instability arises when a local increase in the light intensity creates a depression in plasma density either directly, via the ponderomotive force, or indirectly, via enhanced collisional absorption and subsequent plasma expansion. The density depression refracts the light into the lower-density region, enhancing the intensity perturbations. The instability is termed either ponderomotive⁴³ or thermal⁴⁴⁻⁴⁶ filamentation, depending on which mechanism generates the density depression. The analogous process involving the entire beam is called self-focusing.

Filamentation can significantly affect laser-plasma coupling. Intensity enhancements can introduce or modify other instabilities, change the location of the energy deposition, and possibly aggravate deleterious collective effects such as hot-electron generation. Spatial structure in the irradiation pattern can enhance magnetic field generation,⁴⁷ modify energy transport, and even degrade the symmetry and stability of implosions if the scalelength of the structure is sufficiently long. Filamentation can also complicate the interpretation of coupling and transport experiments, since the intensity in the underdense plasma is no

longer a controlled variable and may be particularly sensitive to details of the beam profile.

Theoretical Description. A simple treatment suffices to illustrate the essential features of filamentation. Consider a plane light wave with intensity I propagating in the z direction in a plasma with uniform density n_0 . We first calculate, in the quasi-static limit, the density perturbation δn_e induced by a perturbation in the intensity profile $I = I_0(1 + \alpha \cos ky)$. In the case of ponderomotive filamentation,

$$n_e = n_0 \exp\left(-\frac{I}{2n_{cr}\theta_{ec}}\right) \quad (41)$$

where n_e is the electron density, n_{cr} is the critical density, θ_e is the electron temperature, and c is the speed of light. This equation is easily obtained by balancing the variation in electron pressure with the variation in the light wave pressure. Hence Eq. (41) gives

$$\frac{\delta n_e}{n_0} \approx -\frac{\alpha I_0 \cos ky}{2n_{cr}\theta_{ec}}, \quad (42)$$

where we have assumed that $I_0 \ll 2n_{cr}\theta_{ec}$.

In the case of thermal filamentation, we balance the electron heat flow with the absorbed intensity:

$$\nabla \cdot (\kappa \nabla \theta_e) = -\kappa I, \quad (43)$$

where κ is the classical electron thermal conductivity and κ is the spatial decay rate due to collisional absorption. The perturbation in intensity drives a perturbation in electron temperature, which leads to a variation in electron density. In particular, we postulate that

$$I = I_0(1 + \alpha \cos ky),$$

$$\theta_e = \theta_0(1 + \beta \cos ky), \quad (44)$$

$$n_e = n_0(1 - \beta \cos ky),$$

where we have invoked pressure balance transverse to the beam, i.e., the quasi-static limit. Only first-order corrections are

retained. With these variations in θ_e and n_e we also have

$$\kappa^T = \kappa_0^T \left(1 + \frac{5}{2} \beta \cos ky\right), \quad (45)$$

$$\kappa = \kappa_0 \left(1 - \frac{7}{2} \beta \cos ky\right), \quad (46)$$

where κ_0^T and κ_0 represent the thermal conductivity and absorption coefficient evaluated at θ_0 and n_0 , the electron temperature and density in the absence of an intensity perturbation.

Substituting Eqs. (45) and (46) into Eq. (43) gives, to lowest order,

$$\frac{\partial}{\partial z} \left(\kappa_0^T \frac{\partial \theta_0}{\partial z} \right) = -\kappa_0 I, \quad (47)$$

and to next order

$$\beta = \frac{\kappa_0 I_0 \alpha}{6\kappa_0 I_0 + k^2 \kappa_0^T \theta_0}. \quad (48)$$

Equation (47) determines θ_0 . Equations (44) and (48) yield

$$\frac{\delta n_e}{n_0} \approx -\frac{\kappa_0 I_0 \alpha \cos ky}{k^2 \kappa_0^T \theta_0}. \quad (49)$$

Here we have assumed that $k^2 \kappa_0^T \theta_0 \gg 6\kappa_0 I_0$, i.e., that the scalelength for temperature variation in the direction of propagation of the light wave is much longer than the wavelength of the intensity modulation in the transverse direction.

Comparison of Eqs. (42) and (49) shows that

$$\left(\frac{\delta n_e}{n_0}\right)_t = \left(\frac{\delta n_e}{n_0}\right)_p \frac{2n_{cr}\kappa_0^T}{k^2 \kappa_0^T}, \quad (50)$$

where the subscripts t and p denote thermal and ponderomotive, respectively. For a high- Z plasma,

$$\frac{2n_{cr}\kappa_0^T}{k^2 \kappa_0^T} \approx \frac{1}{7} \frac{v_{ci}^2}{v_e^2}, \quad (51)$$

where ν_{ei} is the collision frequency appropriate to the high-frequency resistivity and v_e is the electron thermal velocity. Hence thermal filamentation dominates for wavelengths $\lambda \gg 10\lambda_{ei}$, where $\lambda_{ei} = v_e/\nu_{ei}$. To include the implicit Z -dependence of the thermal conductivity, the right-hand side of Eq. (51) is multiplied by a factor of about $(1 + 3.3/Z)$.

The growth of these zero-frequency intensity modulations and their self-consistent density fluctuations can easily be obtained from the dispersion relation⁴⁸ familiar from discussions of Brillouin scattering:

$$\omega^2 - k^2 c_s^2 = \frac{k^2 v_{os}^2 \omega_{pi}^2}{4} \left[\frac{1}{D(\omega - \omega_{bi}, \mathbf{k} - \mathbf{k}_i)} + \frac{1}{D(\omega + \omega_{bi}, \mathbf{k} + \mathbf{k}_i)} \right]. \quad (52)$$

Here $D(\omega, \mathbf{k}) = v_{os}^2 - k^2 c_s^2 - \omega_{pi}^2/c^2$ is the sound speed, ω_{pi} (ω_{pe}) is the ion (electron) plasma frequency, ω_0 and $\mathbf{k}_0$ are the frequency and wave number of the incident light wave, and v_{os} is the oscillatory velocity of an electron in the incident light wave. The derivation assumes that there are no variations along the direction of the electric field vector of the light wave. To obtain ponderomotive filamentation, we look for zero-frequency fluctuations with a wave vector orthogonal to the direction of propagation $\mathbf{k}_0$. In particular, we set $\omega = 0$ and take $\mathbf{k} = \mathbf{k}_\perp + i\mathbf{k}_\parallel$, where $\mathbf{k}_\perp \cdot \mathbf{k}_0 = 0$ and where $\mathbf{k}_\parallel$ is parallel to $\mathbf{k}_0$. Equation (52) then gives the dispersion relation

$$4k_\perp^2 k_0^2 + k_\parallel^4 - k_\parallel^2 \frac{\omega_{pe}^2}{c^2} \frac{v_{os}^2}{2v_e^2} = 0. \quad (53)$$

Growth is maximized (in the sense that $\partial k_\parallel / \partial k_\perp = 0$) when

$$k_\perp = \frac{1}{2} \frac{v_{os}}{v_e} \frac{\omega_{pe}}{c}, \quad (54)$$

$$k_\parallel = \frac{1}{8} \left(\frac{v_{os}}{v_e} \right)^2 \frac{\omega_{pe}^2}{k_0 c^2},$$

where we have also assumed that $Z\theta_e \gg \theta_i$.

Since laser beams often have sizable intensity structure, we conservatively define an intensity threshold by requiring that E^2

exponentiate once: i.e., we set $2k_\parallel L \approx 1$, where L is the plasma length. Then

$$I_T \approx \frac{2 \times 10^{15}}{L_\mu \lambda_\mu} \frac{n_{cr}}{n_0} \theta_{kev}, \quad (55)$$

where θ_{kev} is the electron temperature in keV, L_μ (λ_μ) is plasma length (wavelength) in microns, and $n_0 \ll n_{cr}$. There is likewise a collisional threshold given by the condition $2k_\parallel l_{abs} \approx 1$, where l_{abs} is the collisional absorption length. Taking $\ln A \approx 7$, we estimate

$$I_T \approx \frac{3 \times 10^{12}}{\lambda_\mu^3} \frac{Z}{\theta_{kev}^2} \frac{n_0}{n_{cr}}. \quad (56)$$

As shown in Eq. (50), the dispersion relation for thermal filamentation can be obtained from Eq. (53) by multiplying the intensity term by approximately $(7l_i^2 \lambda_{pi}^2)^{-1}$. Then

$$4k_\perp^2 k_0^2 + k_\parallel^4 - \frac{v_{os}^2}{14v_e^2} \frac{\omega_{pe}^2}{c^2 \lambda_{pi}^2} = 0, \quad (57)$$

where we have again assumed that $Z \gg 1$. The growth is now maximized for long wavelength:

$$k_\perp \approx \left(\frac{\omega_{pe}}{3.6c} \frac{v_{os}}{v_e} \frac{1}{\lambda_{pi}} \right)^{1/2}, \quad (58)$$

$$k_\parallel \approx \frac{\omega_{pe}}{7.5k_0 c} \frac{v_{os}}{v_e} \frac{1}{\lambda_{pi}}.$$

We note that the maximum spatial gain coefficient for thermal filamentation exceeds that for ponderomotive filamentation when $v_{os}/\omega_{pe} > v_{os}/c$. Hence the thermal mechanism is especially competitive in the denser, more collisional plasmas produced by short-wavelength laser light.

We define a threshold intensity for thermal filamentation, again using the condition $2k_\parallel L = 1$. Taking $\ln A = 7$, we estimate

$$I_T \approx 2 \times 10^{14} \left(\frac{n_{cr}}{n_0} \right)^2 \frac{\theta_{kev}^2 \lambda_\mu^2}{Z^2 L_\mu^2}. \quad (59)$$

We estimate a collisional threshold using the condition $2k_\parallel l_{abs} = 1$, which yields

$$I_T \approx 6 \times 10^{13} \frac{n_0}{n_{cr}} \frac{\theta_{keV}^2}{\lambda_\mu^2} \quad (60)$$

Discussion. To illustrate the numbers, consider conditions typical of recent experiments⁴⁹ in which gold disks were irradiated with 1-ns pulses of 0.53- μm laser light. We take $I \approx 2 \times 10^{15} \text{ W/cm}^2$, $L_\mu \approx 150$ at $n_0/n_{cr} \approx 0.1$, $\theta_{keV} \approx 2$, and $Z \approx 50$. The minimum growth length ($l_g = 1/k_i$) for ponderomotive filamentation is then about 60 μm for a filament with a wavelength of about 10 μm . There are several growth lengths for ponderomotive filamentation; its threshold intensity (according to our definition of threshold) for these conditions is about $5 \times 10^{14} \text{ W/cm}^2$. Thermal filamentation is near its threshold. The minimum growth length for thermal filamentation at $n = 0.1n_{cr}$ is about 300 μm for a filament with a wavelength much longer than about 20 μm .

Filamentation can take place in longer-scalelength plasmas at lower intensities. For example, with $L_\mu \approx 10^3$ at $n_0 = 0.1n_{cr}$, 0.53- μm laser light, $\theta_{keV} = 2$, and $Z = 37$, the threshold intensity is about 10^{14} W/cm^2 for either ponderomotive or thermal filamentation.

In making these estimates, we have treated the filamentation instability as convective. Under some circumstances the instability can become absolute⁵⁰—for example, with feedback via reflection from the critical surface. Improved linear theories will also allow for plasma inhomogeneity, plasma flow,⁵¹ radiation cooling,⁵² and the stabilizing influence of angular spread in the beam.⁵³ The quasi-static assumption is sometimes a poor approximation. The time required for the density to adjust to changes in the intensity or temperature can be quite significant in a large filament. Finally, improved calculations of thermal filamentation require a better understanding of electron heat transport, which we have assumed to be classical.

Ponderomotive filamentation is closely related to Brillouin scattering, except that in Brillouin scattering the density fluctuation is associated with a resonant mode, an ion acoustic wave. Hence the Brillouin instability has a larger gain coefficient (or growth rate) than ponderomotive filamentation, provided that the ion waves are weakly damped. Some calculations⁵⁴ indicate that

Brillouin sidescattering can prevent filamentation. In practice, however, filaments may grow from a much higher noise level than the scattering does, since laser beams often have sizable structure to start with. Furthermore, filamentation is not particularly sensitive to plasma inhomogeneity, whereas Brillouin scattering is.

Although filamentation is difficult to quantify, evidence for filamentation is rapidly accumulating in laser-plasma experiments. Some of the evidence is indirect, taking the form of inferences from structure⁵⁵⁻⁵⁷ in x-ray pictures of the heated plasma or in images of the back-reflected light, or of inferences from the angular distribution⁵⁸ of half-harmonic light or from frequency shifts⁵⁹ in the reflected light. Recently, filamentary structures have been observed directly by optical shadowgraphy,⁶⁰ by imaging of second-harmonic emission,⁶¹ and by Thomson scattering⁶² from electron plasma waves generated in the filament walls. Some experiments⁶⁰ with short-wavelength light also indicate that implosions are degraded when signs of filamentation are present.

Summary. Filamentation of laser light is a potentially significant process in targets with large regions of underdense plasma. Some laser-plasma experiments show clear signs of this process. The competition and interplay of filamentation with other coupling processes is a rich topic for experiment and theory.

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Raman Scattering at Plasma Density Maxima

Intense laser light incident on plasma is susceptible to numerous parametric instabilities.⁶³ One such instability is stimulated Raman scattering (SRS), in which the incident light is converted into light of lower frequency while simultaneously exciting Langmuir waves. In this article we discuss SRS in plasmas that contain local density maxima, and show that particularly low thresholds apply for sidescattering in that case.

The matching conditions, which require that the sum of the scattered-light and Langmuir-wave frequencies equal the laser frequency, confine SRS to regions where the

plasma density is at or below one quarter of the critical density for the laser light. In an infinite, homogeneous plasma, the SRS growth rate is given by

$$\gamma_0 = (k_2 v_0 / 4) (\omega_p / \omega_1)^{1/2} e_0 \cdot e_1 \quad (61)$$

Here v_0 is the oscillatory velocity of an electron in the pumping electric field E_0 ($v_0 = eE_0/m\omega_0$), e_0 and e_1 are the polarization vectors of the pump and scattered light waves, k_2 is the wave vector of the Langmuir wave, and ω_1 is the frequency of the scattered light wave. In such a plasma the threshold for SRS is given by $\gamma_0 > (v_1 v_2)^{1/2}$, where v_1 and v_2 are, respectively, the damping rates of scattered light and plasma waves arising from collisional and/or Landau damping. This threshold is readily exceeded in situations of interest to laser fusion, except in the highly collisional plasmas that can arise when high-Z targets are irradiated with light of UV or shorter wavelengths (see the article "Summary of Absorption by Gold-Disk Targets," in Sec. 5).

With this exception, SRS is limited not by damping but by plasma inhomogeneity in typical disk-irradiation experiments, such as those recently performed on Novette. In an inhomogeneous plasma, SRS is generated in a resonant region where the matching conditions are satisfied locally, and its growth is then limited by convection of the unstable waves out of this region, whose size is proportional to the local density-gradient scalelength $L' = [d \ln(n)/dx]^{-1}$. Rosenbluth⁶⁴ has shown that light or plasma waves from thermal or other noise sources incident on the resonant region are amplified by a factor $\exp(2\pi G)$ in intensity; here $G = \gamma_0^2 / \kappa' v_1 v_2$, where v_1 (v_2) is the component of the group velocity of the scattered-light (plasma) wave in the gradient direction, and κ' is a mismatch factor whose value depends on the size of the resonant region. Furthermore, absolute instabilities are found in which reflection of the scattered light wave from its critical surface plays an essential role.⁶⁵

The value of the scalelength L' defined above diverges where the density profile is locally flat, and the above theory no longer applies. The resonant region remains finite, but its extent now depends on the curvature scalelength $L = \{[2(n)^{-1} d^2 n / dx^2]\}^{-1/2}$.

Such density maxima, or flat spots, necessarily occur in experiments in which thin foils become underdense to the laser light because of expansion. Density inhomogeneities superimposed on the expanding corona of a solid target can be expected to give rise to local density flat spots when the background scalelength is sufficiently long.

Using a simple fluid description of the plasma, we can describe SRS by the following pair of coupled wave equations for the scattered wave jitter velocity $\mathbf{v}_1 = e\mathbf{A}/mc$ and the electron density perturbation δn :

$$[c^2 \nabla^2 + \omega^2 - \omega_p^2(x)] \mathbf{v}_1 = \frac{1}{2} v_0 \omega_p^2 (\delta n / n) \quad (62)$$

$$[3v_1^2 \nabla^2 + (\omega_0 - \omega)^2 - \omega_p^2(x)] (\delta n / n) = \frac{1}{2} \nabla^2 (\mathbf{v}_0 \cdot \mathbf{v}_1) \quad (63)$$

Here ω is the frequency of the scattered light wave, $\mathbf{v}_0$ is the jitter velocity of the pump, and $v_1 = (T_e/m)^{1/2}$ is the electron thermal velocity. We write the plasma density in the vicinity of its maximum as

$$\omega_p^2(x) = \omega_{p0}^2 [1 - x^2/L^2] \quad (64)$$

and we assume that the pump is incident in the x direction. SRS will be generated in the vicinity of the matching point given by

$$\begin{aligned} \kappa(x) = \frac{1}{\sqrt{3}v_1} \{ (\omega_0 - \omega)^2 - \omega_p^2 - 3v_1^2 k_\perp^2 \}^{1/2} \\ - \frac{1}{c} (\omega_0^2 - \omega_p^2)^{1/2} \\ \pm \frac{1}{c} (\omega^2 - \omega_p^2 - c^2 k_\perp^2)^{1/2} = 0 \quad (65) \end{aligned}$$

where the scattered light has transverse wave vector $k_\perp$ and where the plus and minus signs correspond to forward and backward scattering, respectively. If $v_1/c \ll 1$, the matching density is given by $\omega_p = \omega_0 - \omega$, and the x component of the plasma-wave $\mathbf{k}$ -vector is

$$\begin{aligned} k_{z1} = \frac{1}{c} (2\omega\omega_0 - \omega^2)^{1/2} \\ \pm \frac{1}{c} (2\omega\omega_0 - \omega_0^2 - c^2 k_\perp^2)^{1/2} \quad (66) \end{aligned}$$

We can approximate ∇^2 on the right hand side of Eq. (63) by $-k_z^2$, and we can treat ω_p^2 on the right hand side of Eq. (62) as constant.

In this profile we expect the strongest instability to occur when the matching point is close to the density maximum—that is, for

$$\omega \approx \omega_0 - \omega_{pi} \quad (67)$$

Then the condition that the scattered light wave be propagating restricts $k_\perp$ to the range

$$k_\perp^2 \leq \frac{1}{c^2} (\omega_0^2 - 2\omega_0\omega_{pi}) \quad (68)$$

with the upper limit corresponding to sidescattering, which we expect to be the most unstable case.

In previous work on a parabolic profile,⁶⁴ Eqs. (61) and (62) were simplified by using the Wentzel-Kramers-Brillouin (WKB) approximation on each wave to obtain a pair of coupled equations for the mode amplitudes. Use of the WKB approximation rules out the consideration of sidescattering, or of any interaction between forward scattering and backscattering⁶⁵; and the approximation is inapplicable if $\omega_{pi}^2 - \omega_p^2 < 3v_i^2 k_\perp^2$, so that the turning point of the plasma wave is below the critical density. Within this (thus significantly restricted) WKB approximation, Rosenbluth found that forward scattering can only be convectively unstable, but that backscattering is absolutely unstable, with a threshold given by

$$\frac{\gamma_0^2}{v_i v_s \kappa''^{2/3}} \geq 1 \quad (69)$$

where κ'' is the second spatial derivative of the mismatch defined in Eq. (65) at the density maximum, where $\kappa = \kappa' = 0$. Equation (69) reduces to approximately⁶⁶

$$(v_0/c)^2 (v_i/c)^{-2/3} (\omega_0 L/c)^{4/3} \geq 1 \quad (70)$$

In the limit of sidescattering v_i vanishes and the WKB approximation fails, but we expect that the correct scaling is obtained by setting

$$v_i \sim (c^2/\omega) (\omega_{pi}/cL)^{1/2} \quad (71)$$

in which we insert the minimum wave number from an exact solution of the wave equation into the expression for the group velocity. This provides a threshold given by

$$\left(\frac{v_0}{c}\right)^2 \left(\frac{\omega_0 L}{c}\right)^{2/3} \geq \max \left[\left(\frac{v_i}{c}\right)^{2/3} \left(\frac{\omega_0 L}{c}\right)^{1/6}, \left(\frac{v_i}{c}\right)^{4/3} \left(\frac{\omega_0 L}{c}\right)^{1/3} \right] \quad (72)$$

for sidescattering at the density maximum.

The threshold in Eq. (72) unsatisfactorily vanishes with electron temperature: in contrast, the threshold for SRS on a nonvanishing density gradient is independent of temperature. For instance, the sidescattering threshold is approximately⁶⁷

$$\left(\frac{v_0}{c}\right)^2 \left(\frac{\omega_0 L}{c}\right)^{4/3} \geq 1 \quad (73)$$

We thus cannot believe Eq. (72) to be valid for sufficiently cold plasma, which suggests that an examination of the cold-plasma limit would be valuable to determine the range of validity of Eq. (72).

These considerations motivate an alternative approximation, in which we neglect the $3v_i^2 \nabla^2$ (dispersion) term in Eq. (63); doing this, we obtain what can be thought of as either a cold-plasma or a strong-damping limit. In the case of a linear gradient,^{65,68} this approximation readily yields the same growth rate and threshold as the warm-plasma analysis. This is not so at a density maximum, where we obtain a nonvanishing threshold, in sharp contrast with the predictions of Eqs. (71) and (72). In the rest of this article we investigate the cold-plasma approximation in some detail. The transition from cold to thermal plasma is not yet fully analyzed.

On eliminating δn between Eqs. (62) and (63), and choosing suitable dimensionless variables, we obtain

$$\left\{ \frac{d^2}{d\xi^2} - \mu + (1 + \rho)(\Omega - iN_1) + \xi^2 - \frac{\Gamma}{[\xi^2 - \mu - (1 - \rho)(\Omega - iN_2)]} \right\} v_i = 0 \quad (74)$$

where

$$\Omega = \left(\frac{\omega_0 L}{c \omega_{pi0}} \right) \left(\omega - \frac{\omega_0}{2} - \frac{c^2 k_{\perp}^2}{2 \omega_0} \right), \quad (75)$$

$$\rho = c^2 k_{\perp}^2 / \omega_0^2, \quad (76)$$

$$N_i = (\omega_0 L / c) (v_i / v_{pi0}), \quad (77)$$

$$\xi = (\omega_{pi0} / c L)^{1/2}, \quad (78)$$

$$\mu = \frac{L}{\omega_{pi0} c} \left[\omega_{pi0}^2 - \left(\frac{\omega_0}{2} - \frac{c^2 k_{\perp}^2}{2 \omega_0} \right)^2 \right], \quad (79)$$

$$\Gamma = \frac{1}{\xi} \left(\frac{v_0}{c} \right)^2 (k_{\perp} L)^2. \quad (80)$$

If $\mu > 0$ (< 0), the plasma is overdense (underdense) to the sidescattered light wave. Equation (74), and the condition that the solutions are outgoing as $\xi \rightarrow \pm \infty$ (corresponding to an absolute instability), together form a Sturm-Liouville system that determines the eigenvalues $\Omega = \Omega(\Gamma, \mu, \rho, N_i)$ for the frequency and growth rate of the instability. We have written a numerical code to determine Ω from this system and have used it to check various analytic approximation schemes described below.

Overdense Case. If $|\mu| \gg 1$ and $\mu > 0$, then the matching point for SRS lies below the density maximum and the scattered wave will be evanescent above this point. We can solve in the region $\xi < 0$ by linearizing about the matching point, which for sidescattering (the most unstable case) lies near $\xi = \mu^{1/2}$. Defining

$$\xi = (4\mu^{1/2})^{1/2} + \frac{1}{2} \xi', \quad \left[\mu + (1 - \rho) \Omega - i N_2 \right]^{1/2},$$

$$\Omega' = (4\mu^{1/2})^{1/2} + \frac{1}{2} \Omega',$$

$$> [2\Omega - i \{ (1 + \rho) N_1 + (1 - \rho) N_2 \}]. \quad (81)$$

$$\Gamma' = 4^{-1/2} \mu^{-1/2} \Gamma,$$

we obtain

$$\left(\frac{d^2}{d\xi'^2} + \Omega' + \frac{\xi'}{2} - \frac{\Gamma'}{2\xi'} \right) v_1 = 0, \quad (82)$$

which maps into the linear-profile problem, for which

$$\text{Im} \Omega' = -\Gamma'^{1/2} + \frac{1}{2} \Gamma'^{1/4}, \quad (83)$$

so that

$$\text{Im} \Omega = -\Gamma^{1/2} + (\mu/2)^{1/2} \Gamma^{-1/4} + \frac{1}{2} (1 + \rho) N_1 + \frac{1}{2} (1 - \rho) N_2. \quad (84)$$

The inhomogeneity threshold is thus

$$\Gamma > (\mu/2)^2. \quad (85)$$

and the damping threshold is

$$\gamma_0 = \frac{1}{4} k_{\perp} v_0 > \frac{1}{2} [v_1 \omega_1 / \omega_0 + v_2 \omega_0 - \omega_1 / \omega_0]. \quad (86)$$

where

$$\omega_i = \text{Re}(\omega) \approx \frac{\omega_0}{2} + \frac{c^2 k_{\perp}^2}{2 \omega_0}. \quad (87)$$

Thus we see that the linear-gradient results, using the local density-gradient scalelength at the matching point, apply until that point lies within a few wavelengths of the density maximum—that is, until $\mu \approx 1$. We anticipate that when sidescattering occurs at the maximum, the threshold is $\Gamma \approx 1$; we confirm and improve this estimate below.

Underdense Case. If $|\mu| \gg 1$ and $\mu < 0$, a different approximation scheme is available. For simplicity we neglect damping, putting $N_2 = 0$. We rewrite Eq. (74) as

$$\left[\frac{d^2}{d\xi^2} + \frac{(\xi^2 - A_-)(\xi^2 - A_+)}{(\xi^2 - B)^2} \right] v_1 = 0, \quad (88)$$

with

$$A_{\pm} = \mu - \rho \Omega \pm (\Gamma + \Omega^2)^{1/2}, \quad (89)$$

$$B = \mu + (1 - \rho) \Omega.$$

Above threshold, $A_{\pm} \ll A_{\pm}$ and $A_{\pm} \ll B$, and $A_{\pm} = 0$ corresponds to the

homogeneous dispersion relation. Thus we have, approximately,

$$\left[\frac{d^2}{d\zeta^2} + \frac{A}{B} (\zeta^2 - A_*) \right] v_1 = 0, \quad (90)$$

a parabolic cylinder equation whose eigenvalue condition is

$$A_* = -(2n + 1) i (B/A)^{1/2}. \quad (91)$$

Above threshold, the right hand side of Eq. (91) is small, and we can solve by iterating the equations

$$\Omega = \frac{-(\mu - \Delta)\rho - i[\Gamma(1 - \rho^2) - (\mu - \Delta)^2]^{1/2}}{1 - \rho^2}, \quad (92)$$

$$\Delta = -(2n + 1) i \times \left[\frac{\mu + (1 - \rho)\Omega}{\mu - \rho\Omega - i[\Gamma + \Omega^2]^{1/2}} \right]^{1/2}, \quad (93)$$

starting with $\Delta = 0$. Since Δ is of order 1, if $|\mu| \gg 1$ the threshold is essentially $\Gamma > \mu^2/(1 - \rho^2) \gg 1$, or

$$\frac{k_z v_{th}}{4} > \frac{\omega_0}{2} - \frac{c^2 k_z^2}{2\omega_0} - \omega_{pi}, \quad (94)$$

$$\left(\frac{v_{th}}{c} \right)^2 (k_z L)^2 > 1,$$

which requires that the excited plasma wave frequency corresponding to side-scattering should not exceed the plasma frequency at the density maximum by more than the homogeneous growth rate.

Sidescattering at the Density Maximum
 $\mu \approx 0$. Our analysis of the underdense and overdense cases indicated that the lowest threshold is for sidescattering right at the density maximum, with a threshold given by $\Gamma > F(\rho)$, where F is of order 1. The case $\rho = 0$ (where the maximum density of the plasma is quarter critical) can be readily treated by WKB methods, which yield a threshold

$$\Gamma > 0.31(2n + 1)^2 \quad (95)$$

and an approximate growth rate for $\Gamma \gg 1$ given by

$$\Omega = -i\Gamma^{1/2} \times [1 - 0.43 e^{-i\pi/3} \Gamma^{-2/3} (2n + 1)^{1/3}], \quad (96)$$

where $n = 0$ corresponds to the first and most unstable of a series of unstable eigenmodes. In dimensional units this becomes

$$\left(\frac{v_{th}}{c} \right)^2 \left(\frac{\omega_0 L}{c} \right)^2 > 1.7. \quad (97)$$

We have thus shown that one can expect enhanced SRS from regions of the target plasma where the local density gradient vanishes, and we have exhibited thresholds and growth rates. Figure 3-24 shows the

Fig. 3-24. Growth rate vs pump parameter for quarter-critical SRS at a density maximum.

normalized growth rate $\text{Im}(\Omega)$ as a function of pump parameter Γ for the case $\rho = 0$ (quarter critical density), for which the instability is truly absolute, for various values of the mismatch parameter μ . The agreement between the analytic approximations and the numerically determined eigenvalues is excellent. Figure 3-25 shows typical eigenfunctions and the corresponding effective potentials of the Schrödinger-like Eq. (74) for the underdense, overdense, and critical cases. The real part of the potential is essentially the negative of the density profile; it vanishes at the turning point of the light wave. The imaginary part of the potential arises from the unstable coupling with the plasma wave, and peaks at the matching point(s). Note that the eigenfunctions have, respectively, finite, small, and evanescent k-vectors at the density maximum in the underdense, critical, and overdense cases.

Because the curvature thresholds at plasma-density flat spots are so low, thresholds in these regions will, in practice, be primarily determined by damping. In this

sense, exploding-foil experiments simulate future reactor targets, for which scalelengths will be long enough that we must rely on collisional damping to keep SRS below dangerous levels.

Author: E. A. Williams

ZOHAR Simulations of Raman Sidescattering and Two-Plasmon Decay

In 1984 we resumed two-dimensional kinetic simulations of Raman scattering and two-plasmon ($2\omega_{pe}$) decay using the ZOHAR code.⁶⁹ Motivated by recent experiments, we typically use an electron temperature of 2 keV and illumination intensities corresponding to $I\lambda^2 \leq 10^{15}$ W/cm², where λ is the laser wavelength in microns. Spatial scalelengths are chosen to place the instability slightly above its density-gradient threshold. Our intent is to check linear threshold theory and nonlinear

Fig. 3-25. SRS eigenfunctions (right) and corresponding Schrödinger-like potentials (left) for $\Gamma = 1$ and $\mu = 0$. (a) $\mu = 0$ (underdense); (b) $\mu = 1$ (critical); (c) $\mu = 2$ (overdense). Arbitrary units.

saturation mechanisms and to learn why observed heating in the underdense plasma is commonly weak.

The plasma is initialized as a planar self-similar flow, shifted in velocity such that the positions of the $0.15n_i$ (for Raman scattering) or $0.25n_i$ (for two-plasmon decay) surfaces remain fixed. Plasma flows into the simulation region at the high-density end with speed $< c_s$, and out at the low-density end with speed $> c_s$. Because of this velocity gradient and the small values of $v_{Te}^2/v_{Ti}^2 \propto T_e/T_i$, Brillouin backscattering is below threshold.

Raman Scattering. The ZOHAR simulations support the Afeyan/Williams linear theory⁷⁰ for gradient thresholds for Raman sidescattering. Their analytic approximation to the growth rate is

$$\frac{\gamma}{\omega_0} = \frac{\alpha}{2} \left[1 - \frac{(\omega_p/\omega_0)^2}{(\omega_0 L_e/c) \alpha^2} \right] \quad (98)$$

where

$$\alpha = \frac{\omega_p}{\omega_0} \frac{v_0}{c} \left[2 - 2 \frac{\omega_p}{\omega_0} - \frac{\omega_p^2}{\omega_0^2} \right]^{1/2} \quad (99)$$

and where electron temperature modifications are neglected. The analytic results are supported by numerical solution of the coupled-mode eigenproblem in x space. For an exponential density profile with density (n), scalelength $L_n = 190c/v_0$, and periodic boundary conditions in the transverse (y) direction, corresponding to the theory (which has real-valued k_y), we find (at $0.15n_i$) good agreement between theoretical and simulation values of the growth rate γ (see Table 3-3). Reference 71 provides another theoretical calculation, which it incorrectly states is in disagreement with Ref. 70.

In these simulations, the laser is directed into the positive x direction (opposite to the density gradient) and polarized in the z direction. The simulation proceeds in the x - y plane, with periodic boundary conditions in the transverse (y) direction. For economical comparisons with theory, the ion motion is not calculated and the transverse dimension L_y is a single wavelength of the scattered light.

How large must the transverse dimension L_y be for reasonable duplication of linear growth in an aperiodic system? A simulation

Intensity, v_0/c	γ	
	Theory	Simulation
0.020	$\sim 4.6 \times 10^{-4}$	Stable
0.025	1.1×10^{-3}	1×10^{-3}
0.028	1.9×10^{-3}	1.6×10^{-3}

of duration T is surely valid if $L_y > v_g T$, where $v_g = c^2 k_y / \omega_0$ is the group velocity of the scattered light. Less conservatively, by analogy to other such problems, periodicity may be irrelevant if γ exceeds v_g/L_y , i.e., if a wave packet of scattered light exponentiates in less than its transit time. Periodicity has less nonphysical effect when density inhomogeneities disrupt phase coherence across L_y . These inhomogeneities can be due to ion waves associated with Brillouin sidescattering, ion waves driven by the plasmons of Raman sidescattering, or curvature of density contours.

The transverse dimension required for reasonable duplication of nonlinear saturation in an aperiodic system involves additional issues. The transverse wave vector k_y is discretized to multiples of $2\pi/L_y$. The sidescattering process for each mode is concentrated at the density for which $\omega_{pe}^2 = \omega_0^2 - (ck_y)^2$. Each mode extends over a range of densities, as can be observed in smaller systems. We would like to have the k_y spacing close enough that the modes overlap in space. This demands values of L_y higher than we have used yet (e.g., 10 wavelengths of the scattered light). Brillouin sidescattering, which coexists with the Raman, is even harder to model satisfactorily because the group velocity of its scattered light is higher.

A simulation with mobile ions was run to time $T = 4800\omega_0^{-1}$. In this simulation, the transverse dimension was $L_y = 53.6c/\omega_0$, or four wavelengths of the Raman-scattered light at $0.15n_i$. Although Brillouin backscattering is suppressed by the flow-velocity gradient, about half the light is sidescattered. The longitudinal field energy peaks at one-third its maximum value in the fixed-ion runs. Late in time, the absorption is less than 6% into a spectrum whose main body is Maxwellian ($T_{\text{hot}} \approx 40$ keV) but whose tail above 100 keV is lower than for such a Maxwellian.

Two-Plasmon-Decay Simulations. The $2\omega_{pe}$ decay simulations were similar, except that the laser was polarized in the plane of the simulation, the density range was $0.19n_i$ to $0.29n_i$, and the density scalelength was

Table 3-3. Growth rates for Raman sidescattering.

$L_{ii} = 200c/\omega_0$. Again, runs were made without and with ion motion. As in our earlier work,⁷²⁻⁷⁴ ion response is an essential aspect of the saturation of the $2\omega_{pe}$ instability. Fixed-ion runs reach much higher plasmon wave energy and T_{hot} .

With the ions frozen, the plasmon wave energy reaches 0.035 (in normalized units), then settles to 0.019. This burst and collapse are accompanied by a peaking of electron thermal energy and a temporary depletion of the light wave in and past the instability region. Later, absorption is 30% into a non-Maxwellian distribution that peaks at 100 keV and extends past 500 keV.

However, as we have reported before,⁷²⁻⁷⁴ ion response plays a dominant role in the saturation of $2\omega_{pe}$ decay. The beating of the broad spectrum of high-frequency decay waves drives up ion fluctuations, which then couple the electron plasma decay waves to shorter-wavelength, more strongly damped perturbations. On a much longer time scale, plasmon wave pressure produces a density step, whose gradient then stabilizes the $2\omega_{pe}$ instability. These features have been observed in several experiments using Thomson scattering. In a recent mobile-ion run, the longitudinal field energy peaked at 5.5×10^{-3} , and the absorption was less than 3%.

Although the theoretical thresholds are confirmed, the simulations show that exceeding the linear threshold need not result in strong absorption. These simulations will resume with additional diagnostics on the space and time evolution of the electron and ion temperatures and on the ion-density and longitudinal-field spectra. The relative timings of features in these diagnostics will clarify the nonlinear saturation processes that determine absorption and hot-electron temperature.

Authors: A. B. Langdon, W. L. Krueer, and B. F. Lasinski

Heating of Underdense Plasma at High $I\lambda^2$

We use a relativistic, electromagnetic, one-dimensional kinetic computer simulation to determine hot-electron temperature and absorption of laser light in underdense plasmas for $I\lambda^2$ approaching $10^{18} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$. Electrons are very rapidly heated to 0.1 to

1 MeV by a mixture of forward and backward Raman scattering. Higher intensities lead to preferential absorption by forward Raman scattering and hence to higher temperatures. The scattered light (both up- and down-shifted) appears to be strongly absorbed.

As $I\lambda^2$ approaches 10^{18} (as can easily occur with high-powered CO_2 lasers), the oscillatory velocity of electrons accelerated by the transverse electric field of the laser light wave approaches the speed of light, indicating enormous electromagnetic energy densities available for heating. (Here I is the laser intensity in W/cm^2 and λ_μ is the laser wavelength in microns.) At these intensities, the spatial gain length of Raman forward scattering⁷⁵ becomes less than a light wavelength. Recent experiments and two-dimensional simulations⁷⁶ have found high absorption ($\sim 80\%$) into very high electron temperatures (approaching 1 MeV); this article explains these results in terms of forward Raman scattering.

Stimulated back, side, and forward Raman scattering (SRS)⁷⁵⁻⁸⁵ in laser-produced plasmas can generate high-energy electrons, which can reduce the energy gain of laser fusion targets. SRS is the parametric decay of an incident photon into a scattered photon and a longitudinal electron plasma wave. The wave numbers k and frequencies ω obey the matching conditions characteristic of parametric processes, $k_i = \pm k_s + k_{epw}$, and $\omega_i = \omega_s + \omega_{epw}$ (the subscripts i , s , and epw refer to the pump, scattered, and electron plasma waves, respectively; the $-$ sign refers to backscattering). The growth rate for Raman backward and forward scattering is

$$\frac{\gamma_0}{\omega_0} = \frac{k_{epw} c}{4\omega_0} \frac{v_{osc}}{c} \left(\frac{\omega_p}{\omega_s} \frac{\omega_p}{\omega_{epw}} \right)^{1/2} \left(1 - \frac{n}{n_c} \right)^{1/4} \quad (100)$$

Here

$$\frac{v_{osc}}{c} = \frac{eE_0}{cm_e\omega_0} = 8.55 \times 10^{-10} I \lambda_\mu^2 [W/\text{cm}^2]^{1/2}, \quad (101)$$

where E_0 is the laser peak (vacuum) electric field, ω_p is the plasma frequency at density n , and n_c is the density at which the laser

frequency $\omega_0 = \omega_p$. The spatial growth rate k_i of the electron plasma wave can be approximated by $k_i = \gamma_0/(v_{gs}v_{gspw})^{1/2}$, where $v_{gs} = c(1 - \omega_p^2/\omega_0^2)^{1/2}$ and $v_{gspw} \approx 3k_{gspw}c^2/\omega_{gspw}$.

This rapid heating can be seen in the phase-space plot of Fig. 3-26. Laser light is incident from the left; the long-wavelength electron plasma wave has the correct wavelength for Raman forward scattering and lags behind the laser light wave by only about 10 free-space wavelengths. The noise in the forward-scattered wave grows as the wave moves from left to right. In the high-amplitude part of the plot, one can see the short-wavelength electron plasma wave (due to Raman backscattering) superimposed on the long-wavelength forward scattering. The noise in the Raman backscattering grows as the wave moves from right to left. The forward scattering is comparable to the backscattering, in spite of its lower spatial growth rate, probably because the electromagnetic turbulence from the laser-plasma interaction provides a large noise initialization.

The transmitted and reflected spectra are very broad, as is to be expected from the large imaginary part of the frequency due to the growth rate and to Landau damping.⁷⁷ Forward Raman scattering produces both up- and down-shifted light,⁷⁷ since both modes satisfy ω and k matching; backscattering does not upshift, since that is far from resonance.

The Manley-Rowe relations state that the pumping energy is partitioned into the scattered light wave and the electron plasma wave in the ratios ω_s/ω_0 and ω_{gspw}/ω_0 , respectively. If the fractions of light into plasma heating and external scattering are examined, one might conclude that the Manley-Rowe relations are not obeyed: they are locally obeyed, however, since the scattered light is so intense that it is subsequently absorbed. Even if the plasma density is greater than quarter critical for the scattered light ($\omega_0^2/\omega_s^2 > 0.25$), the light can still be rapidly absorbed by a forward Compton effect, as is shown by the points in Fig. 3-27 for $n/n_c = 0.5$. The k plots of the electromagnetic light for the $n/n_c = 0.5$ runs show that the light is internally upshifted to about $1.25\omega_0$ and is then rapidly absorbed. Note that the threshold for the absorption at $n/n_c = 0.5$ is about at $I\lambda_\mu^2 \approx (1 \text{ to } 3) \times 10^{17} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$. The heating

Fig. 3-26. Electron phase space in which the laser enters from the left. The long- and short-wavelength oscillations are Raman forward and back scattering, respectively.

Fig. 3-27. Heated electron temperature vs $I\lambda_\mu^2$ and n/n_c .

and frequency shifting mechanism may be a form of Compton scattering.⁸⁵

When characteristic runs were repeated with the code not relativistic, the heating was even faster. The time and space steps were reduced to 64 per laser period and laser wavelength without significant change in heating. To remove the oscillatory part of the heating, a wave packet was sent through the plasma and the heated electron temperature was measured after the electromagnetic wave had passed: again the heating

was about the same. Finally, the plasma length was doubled without leading to substantial changes to T_{hot} . Recent two-dimensional results⁸⁶ show that Raman sidescattering in the forward direction may be important.

In summary, Raman forward scattering may play a major role in the very strong underdense plasma heating observed in CO_2 experiments.

Author: K. G. Estabrook

2-D LASNEX Studies of Burn-Through by Small, Intense Spots

We used 2-D LASNEX to study the effects of foil thickness and laser spot diameter on the (burn-through) time required for laser light to ablate a CH foil enough that the light is transmitted. Throughout the study,

the laser intensity and wavelength were $3 \times 10^{13} \text{ W/cm}^2$ and $1.06 \mu\text{m}$, respectively. We find that $10\text{-}\mu\text{m}$ spots burn through much faster than larger spots: this is because spherical divergence allows the debris to convect away rapidly, and because a kind of thermal self focusing⁸⁷ allows the light to become quite concentrated in the foil hole. The burn-through time is about the same with diverging $f/4$ optics as with converging $f/4$ optics.

Figure 3-28 shows a representative calculation. Here a 100-J , 1-ns (FWHM) Gaussian pulse was focused into a $50\text{-}\mu\text{m}$ -diam spot on a $15\text{-}\mu\text{m}$ -thick CH foil. Light rays are just burning through 0.1 ns after the peak of the pulse. Other rays are reflecting off the funnel-like critical surface onto the central axis, where the peak intensity becomes $3 \times 10^{16} \text{ W/cm}^2$, an order of magnitude greater than the incident intensity (and probably limited by the numerical mesh of

Fig. 3-28. A 2-D ray-trace plot of a 100-J laser light pulse, incident from the right, focused to a $50\text{-}\mu\text{m}$ -diam spot on a $15\text{-}\mu\text{m}$ -thick CH film. The time is 0.1 ns after the peak of the 1-ns (FWHM) pulse.

the simulation). The density is $0.1n_e$ and the temperature is about 1.5 keV where the filament intensity is greatest.

Figure 3-29 gives burn-through times for CH foils 2.5, 5, 10, 15, and 20 μm thick with spots 50, 158, and 500 μm in diameter. The effect of the spherical divergence is small for very thin foils (e.g., 2.5 μm) and considerable for thick foils.

The implication of these results for laser fusion is that small spots simulate the hot spots of a much larger laser beam,⁴⁷ and could represent a failure mode in direct-illumination implosion experiments.

Author: K. G. Estabrook

Direct Implicit Electromagnetic Plasma Code

A direct implicit particle-in-cell (PIC) plasma simulation model with full electromagnetic (EM) effects has been implemented in two-dimensional Cartesian geometry [two spatial coordinates (x, y) and three velocity coordinates (V_x, V_y , and V_z)]. We expect to find this code useful for studies of electron heat flow and ion-beam propagation in ICF and other applications. In particular, we anticipate several applications related to the effects and generation of strong magnetic fields.

The model, implemented with the D_1 time differencing scheme,⁸⁸ has progressed through a succession of versions, starting with a one-dimensional electrostatic (ES) version.⁸⁹ The ES version was implemented first to gain experience with spatial differencing in forms suitable for extension to the full EM field in two dimensions.

The addition of EM effects entails a considerable extension of the ES model.⁹⁰ The most obvious additional requirement for the particle pusher is that particles must undergo a rotation; and, of course, EM codes require accumulation of particle current as well as charge density. Rotation poses no new difficulties, since it entails only a simple application of the algorithm used in explicit EM codes. Implicit codes also require the accumulation of coefficients with which to represent the part of the plasma charge and current densities that arises from the advanced electric field E . Unmagnetized ES codes need one scalar coefficient; EM codes require vector quantities, which become

Fig. 3-29. Burn-through time vs foil thickness for spots of 50, 158, and 500 μm diameter with an intensity of $3 \times 10^{15} \text{ W}/\text{cm}^2$.

tensors when the rotational contribution from the magnetic ($\mathbf{B}$) field is considered. The earlier ES work provided guidance for the accumulation of these coefficients; the tensors are easily obtained with vector quantities accumulated from the particles, as is done in explicit EM particle codes.

The implicit EM field solution is also considerably different from the implicit ES code. The $\mathbf{E}$ -field calculation now includes an inductive part in addition to the electrostatic, and of course the $\mathbf{B}$ field must be self-consistently advanced. Below we discuss the model and then describe numerical implementation of the field equations and some boundary condition considerations.

The EM Direct Implicit Model. The model with the D_1 difference scheme can be expressed as follows: The particle description is given by

$$\mathbf{v}_{n,1,2} = \mathbf{v}_{n,1,2} + \Delta t \times \left[\bar{\mathbf{a}}_n + \frac{1}{2} (\mathbf{v}_{n,1,2} + \mathbf{v}_{n,1,2}) \times \frac{q\bar{\mathbf{B}}_n}{mc} \times \mathbf{x}_n \right] \quad (102a)$$

$$\mathbf{x}_{n,1} = \mathbf{x}_n + \Delta t \bar{\mathbf{v}}_{n,1,2} \quad (102b)$$

$$\bar{\mathbf{a}}_n = \frac{1}{2} \left(\bar{\mathbf{a}}_{n,1} + \frac{q\mathbf{E}_{n,1}(\mathbf{x}_{n,1})}{m} \right) \quad (102c)$$

where the quantities $\mathbf{x}$, $\mathbf{v}$, and $\bar{\mathbf{a}}$ are considered to have index i for the i th particle. The

field equations are finite-differenced in time and are

$$\mathbf{E}_{n+1} - \mathbf{E}_n = c \Delta t \nabla \times \mathbf{B}_{n+1/2} + 4\pi \Delta t \mathbf{J}_{n+1/2} \quad (102d)$$

$$\mathbf{B}_{n+1/2} - \mathbf{B}_{n-1/2} = -c \Delta t \nabla \times \mathbf{E}_n \quad (102e)$$

$$\bar{\mathbf{B}}_{n+1} - \mathbf{B}_{n+1/2} = -c \frac{\Delta t}{2} \nabla \times \mathbf{E}_{n+1} \quad (102f)$$

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

The fields are coupled to the particles by

$$\rho_{n+1} = \sum_j q_j \sum_i S(\mathbf{X}_i - \mathbf{x}_{n+1}) \quad (102h)$$

$$\mathbf{J}_{n+1/2} = \sum_j q_j \sum_i \mathbf{v}_{n+1/2} \times \frac{1}{2} [S(\mathbf{X}_i - \mathbf{x}_n) + S(\mathbf{X}_i - \mathbf{x}_{n+1})] \quad (102i)$$

where the (vector) index j denotes the (j, k) grid location. The shape factor S relates each particle to the spatial grid at position $\mathbf{X}_i$.

This scheme is implemented in two steps. First we advance the particles to an intermediate level (the so-called $\sim$, or tilde, level) using everything but the advanced $\mathbf{E}$ field (a step called the "pre-push"):

$$\tilde{\mathbf{v}}_{n+1/2} = \mathbf{v}_{n+1/2} + \frac{\Delta t}{2} \times \left\{ \bar{\mathbf{a}}_{n+1} + \frac{1}{2} (\tilde{\mathbf{v}}_{n+1/2} + \mathbf{v}_{n+1/2}) \times \frac{q \bar{\mathbf{B}}_n}{mc} \mathbf{x}_n \right\} \quad (103a)$$

$$\tilde{\mathbf{x}}_{n+1} = \mathbf{x}_n + \Delta t \tilde{\mathbf{v}}_{n+1/2} \quad (103b)$$

Next we accumulate a charge and two current densities for each species by summing over particles:

$$\tilde{\rho}_{s,n+1}(\mathbf{X}_j) = q_s \sum_i S(\mathbf{X}_j - \tilde{\mathbf{x}}_{n+1}) \quad (104a)$$

$$\tilde{\mathbf{J}}_{s,n+1/2}(\mathbf{X}_j) = q_s \sum_i \tilde{\mathbf{v}}_{n+1/2} S(\mathbf{X}_j - \tilde{\mathbf{x}}_{n+1}) \quad (104b)$$

$$\tilde{\mathbf{J}}_{s,n+1/2}(\mathbf{X}_j) = \frac{1}{2} \tilde{\mathbf{J}}_{s,n+1/2}(\mathbf{X}_j) + \frac{1}{2} q_s \sum_i \tilde{\mathbf{v}}_{n+1/2} S(\mathbf{X}_j - \mathbf{x}_n) \quad (104c)$$

The source terms needed for the field solution are obtained by summing over species to obtain

$$\tilde{\rho}_n(\mathbf{X}_j) = \sum_s \tilde{\rho}_{s,n+1}(\mathbf{X}_j) \quad (105a)$$

$$\tilde{\mathbf{J}}_{n+1/2}(\mathbf{X}_j) = \sum_s \tilde{\mathbf{J}}_{s,n+1/2}(\mathbf{X}_j) \quad (105b)$$

and the tensors

$$\chi(\mathbf{X}_j) = \sum_s \frac{1}{4} \Delta t^2 \times \frac{q_s}{m_s} \tilde{\rho}_{s,n+1}(\mathbf{X}_j) [\mathbf{1} + \mathbf{R}_s(\mathbf{X}_j)] \quad (105c)$$

$$\zeta(\mathbf{X}_j) = \sum_s \frac{1}{4} \Delta t^3 \times \frac{q_s}{m_s} \tilde{\mathbf{J}}_{s,n+1/2}(\mathbf{X}_j) \times [\mathbf{1} + \mathbf{R}_s(\mathbf{X}_j)] \quad (105d)$$

are computed on the spatial mesh. The rotation tensor $\mathbf{R}_s$ is given by

$$\mathbf{R}_s(\mathbf{X}_j) = \frac{[1 - \theta^2] \mathbf{1} - 2\theta \times \mathbf{1} + 2\theta\theta}{(1 + \theta^2)} \quad (106a)$$

$$\theta = \frac{\Omega \Delta t}{2} = \frac{q B(\mathbf{X}_j) \Delta t}{2mc} \quad (106b)$$

the tensors χ and ζ are constructed so that the total charge and current densities ρ_{n+1} and $\mathbf{J}_{n+1/2}$ can be recovered from

$$\rho_{n+1/2}(\mathbf{X}_j) = \rho_{n+1/2}(\mathbf{X}_j) - \nabla \cdot \chi \cdot \mathbf{E}_{n+1} \quad (107a)$$

$$\Delta t \mathbf{J}_{n+1/2}(\mathbf{X}_j) = \Delta t \tilde{\mathbf{J}}_{n+1/2}(\mathbf{X}_j) + \chi \cdot \mathbf{E}_{n+1} - \nabla \times \zeta \cdot \mathbf{E}_{n+1} \quad (107b)$$

The field equations are solved by substituting the expressions for $\rho_{n+1/2}$ and $\mathbf{J}_{n+1/2}$ into Eqs. (102d) to (102g). Rearranging, we replace Eq. (102d) with

$$\mathbf{E} = A[\nabla^2 \mathbf{E} - \nabla(\nabla \cdot \mathbf{E})] + 4\pi(\chi \cdot \mathbf{E} + \nabla \times \zeta \cdot \mathbf{E}) = \mathbf{Q} \quad (108a)$$

$$\mathbf{Q} = \mathbf{E}_n - 4\pi \Delta t \tilde{\mathbf{J}} + \epsilon \Delta t \nabla \times \mathbf{B}_{n+1/2} - A \nabla \times \nabla \times \tilde{\mathbf{E}}_{n+1} \quad (108b)$$

$$A = \frac{c^2 \Delta t^2}{2} \quad (108c)$$

where the unsubscripted $\mathbf{E}$ is $\mathbf{E}_{n+1/2}$. The equations for $\mathbf{E}$ govern the evolution of both the electrostatic and inductive parts of the electric field. Taking the divergence of Eq. (108) and requiring that $\nabla \cdot \mathbf{E} = 4\pi\rho$, we obtain the following consistency condition on $\tilde{\mathbf{J}}_{n+1/2}$:

$$\nabla \cdot (\mathbf{1} + 4\pi\chi) \cdot \mathbf{E}_{n+1} = \nabla \cdot \mathbf{E}_n - 4\pi \Delta t \nabla \cdot \tilde{\mathbf{J}}_{n+1/2} \quad (109)$$

This is equivalent to requiring that the continuity equation

$$4\pi\delta\rho_{n+1/2} - \nabla \cdot \mathbf{E}_n + 4\pi \Delta t \nabla \cdot \tilde{\mathbf{J}}_{n+1/2} = 0 \quad (110)$$

be satisfied at the $\sim$ level.

It is possible to arrange the accumulation of $\mathbf{J}$ so that Eq. (110) is satisfied automatically. In practice, however, it is more convenient and less noisy to obtain each component of $\mathbf{J}$ in a manner analogous to that used for ρ and to correct the irrotational part of $\mathbf{J}$ so that Gauss's law is satisfied.³⁰ The correction is obtained in the form of the gradient of a scalar potential, added to $\tilde{\mathbf{J}}_{n+1/2}$. The scalar potential is governed by the Poisson equation that results from this substitution into Eq. (110).

Finally we finish the particle push (a step called the "final push") using the new $\mathbf{E}$ field:

$$\delta \mathbf{v} = \frac{\Delta t}{2} \left[\frac{q}{m} \mathbf{E}_{n+1}(\tilde{\mathbf{x}}_{n+1}) + \frac{\delta \mathbf{v}}{2} \times \frac{q \tilde{\mathbf{B}}_n}{mc}(\mathbf{x}_n) \right] \quad (111a)$$

$$\mathbf{x}_{n+1} = \tilde{\mathbf{x}}_{n+1} + \Delta t \delta \mathbf{v} \quad (111b)$$

$$\mathbf{v}_{n+1/2} = \tilde{\mathbf{v}}_{n+1/2} + \delta \mathbf{v} \quad (111c)$$

$$\tilde{\mathbf{x}}_n = \frac{1}{2} \left[\tilde{\mathbf{x}}_{n+1} + \frac{q}{m} \mathbf{E}_{n+1}(\tilde{\mathbf{x}}_{n+1}) \right] \quad (111d)$$

The consistency of the method depends on the details of the connection between the particles and the self-consistent fields. Inconsistencies can arise on several levels. We have found it unnecessary to iterate Eqs. (111), so that the new $\mathbf{E}$ for each particle is actually evaluated at the full new particle position, as the model calls for in Eqs. (102). This inconsistency can easily be reduced, if need be, by cycling through Eqs. (111a) and (111b) until $\mathbf{x}_{n+1}$ converges. Such consistency might be important for energetic particles ejected from a strong plasma-density gradient.

A higher level of consistency is required in the accumulation of χ and ζ to ensure that the correction of the $\tilde{\mathbf{J}}$ —which, in turn, is required to ensure that $\nabla \cdot \mathbf{E} = 4\pi\rho$ is in fact satisfied after the final push. The prescription given in Eqs. (105c) and (105d) is adequate and robust for many applications, and it is relatively fast. Other options are discussed by Langdon and Barnes.⁴⁰

This scheme appears to need two input/output passes through the particles, but one pass can be eliminated by combining the final push of step $n-1$ with the pre-push for step n for each buffer of particles. Both parts of the particle push can now be carried out at the same time, so that only one particle-buffering cycle is needed for each time step.

A minor inconvenience is that all particle quantities are at the $\sim$ level, except for a brief period in the middle of the reordered parts of the particle push. But these quantities are (conveniently) consistent for diagnostic purposes, because the particle and

current densities are also computed only at the $\sim$ level. In fact, all plasma source diagnostics of the simulation are representations obtained from $\sim$ -level quantities.

Numerical Implementation of Field Solution. The reduction of the field algorithm to a finite-difference procedure on a mesh has several pitfalls that can reduce the utility and robustness of the code. The relative positions on the mesh at which evaluations are made can significantly affect the numerical diffusion exhibited by the resulting algorithm. Almost as important, the positioning of the field and source components can determine the effort required to implement various choices for boundary conditions. Previous one-dimensional ES studies³⁹ provide guidance for the selection of locations for quantities likely to dominate when EM effects are less important. We use the storage locations shown in Fig. 3-30, where the superscript * implies all components and where

$$x_i = x_{\min} + j(x_{\max} - x_{\min})/N_x \quad (112)$$

$$y_k = y_{\min} + k(y_{\max} - y_{\min})/N_y$$

Quantities not specified in the figure correspond to the cell corners—that is, to integer (j, k) values.

All quantities in the particle pusher are considered to be at cell corners, so that the currents and various parts of χ and ζ must be averaged to the appropriate positions before the field solution, and the various components of $\mathbf{E}$ and $\mathbf{B}$ are averaged back to cell corners before returning to the particle advance.

The solution of Eq. (108a) is the most interesting part of the field solution. The

other field equations, Eqs. (108b) to (108d), can be evaluated point by point once $\mathbf{E}_{n+1}$ is known. Equation (108a) involves a coupling of $\mathbf{E}$ components and plasma source terms that is difficult to solve unless it is properly arranged. One-dimensional tests, for example, quickly showed the advantage of solving all components simultaneously with a linear system solver. Picard iteration, in which one solves for each component in turn using the most recent values for the other components, generally converged slowly and led to additional constraints on Δt . For $\omega_{ce} \Delta t \geq 2$, Picard iteration did not converge at all.

The one-dimensional results suggest that a simultaneous solution may be essential in two dimensions as well. Since a direct solution would require far too much storage, an iterative solution has been implemented. We have achieved good results with an alternating-direction implicit method that solves all three component equations simultaneously in each direction pass. As in one dimension, no new constraints on Δt emerge, and the convergence is remarkably rapid in those cases in which the χ coefficients dominate. Even if the plasma terms are small, convergence is still more than acceptable if the terms with A coefficients are comparable to other coefficients in Eq. (108a). For example, typical problems on a 32×32 mesh can easily be iterated to residue-to-field ratios of 10^{-4} for less than 10% of the particle-pushing cost, assuming 15 particles per cell.

The desired boundary conditions must be built in to the coefficients sent to the linear system solver. In general, we allow the field values needed at each boundary point to be unknowns and define boundary value equations for each point such that the boundary point and the nearest interior neighbor are coupled to the desired boundary condition. For example,

$$a_1 E_y(j = 0, k + 1/2) + a_2 E_y(j = 1, k + 1/2) = a_3(k + 1/2) \quad (113)$$

provides a Dirichlet condition on $E_y(x_{\min})$ for $a_1 = 1$, $a_2 = 0$, and a given $a_3(y)$ (the desired boundary value as a function of y at $x_{\min}$). Other choices allow easy imposition of Neumann conditions, and either type of boundary condition can be implemented for integer or half-integer mesh storage. (Periodic

Fig. 3-30. Mesh locations of field and source components used in the $\mathbf{E}$ -field calculation.

boundary conditions do not fit into this scheme. A separate linear system solver has been constructed to provide this capability.)

It is often useful to be able to run with TM wave boundary conditions. If we wish to consider a transverse-magnetic (TM) wave incident on the simulated plasma from the region to the left of $x_{\min}$, for example, we need a relation to connect E_y and B_z at $x_{\min}$ at the advanced time (because the field calculation is also implicit). The desired relation is

$$\frac{1}{2} (E_{y, 0, n+1} + E_{y, 0, n}) = \underbrace{B_{z, 0, n+1/2}}_{\text{right-traveling specific input}} - \underbrace{\left[\frac{1}{2} (B_{z, -1/2, n+1/2} + B_{z, -1/2, n+1/2}) - B_{z, 0, n+1/2} \right]}_{\text{left-traveling reflected rf}} \quad (114)$$

where $B_{z, -1/2, n+1/2}$, $B_{z, -1/2, n+1/2}$ and $E_{y, 0, n+1}$ have not yet been determined. The quantities $B_{z, 0}(t)$ and $B_{z, 0}(t)$ are the specified time-dependent rf drive input and bias fields, respectively.

To use Eq. (114) as a boundary condition for the E-field calculation Eq. (108a), we must eliminate these unknown B_z quantities with expressions for known quantities and earlier field quantities. We use Ampere's law, Eq. (102d), to eliminate $B_{z, -1/2, n+1/2}$ and Faraday's law, Eq. (102e), to eliminate $B_{z, -1/2, n+1/2}$. Then, assuming no plasma currents at $x_{\min}$, we have

$$\begin{aligned} (1 - v^{-1} - v) E_{y, 0, n+1} + v E_{y, 0, n} \\ = 2 B_{z, 0, n+1/2} - (1 + v^{-1}) E_{y, 0, n} \\ + 2 B_{z, 0, n+1/2} - v \bar{E}_{y, -1/2, n+1/2} - \bar{E}_{y, 0, n+1/2} \end{aligned} \quad (115)$$

where $v = c dt/dx$. This relation is in the form of Eq. (113), as our algorithm requires. An example of the use of this relation can be found in the article "Compression of Microwaves by a Moving Compact Torus," later in this section.

Filamentation Example. The problem of beam filamentation, which is of interest in studies of electron-beam propagation, currently serves as a test problem for our code. We start with a charge-neutral plasma

column of circular cross section. Two electron components, a low-density "beam" and a higher-density "target," are initialized to counterstream through the massive ion column. The filamentation instability results from the merging of low-density beam particles to form current filaments. These filaments coalesce in regions that now exclude the target electron component in order to preserve charge neutrality.

Figure 3-31 shows the disposition of beam electrons after several instability growth times. Later, as the instability nears saturation, the beam-species filaments have further coalesced into one current channel, as shown in Fig. 3-32. The target-electron distribution in Fig. 3-33 exhibits the expected exclusion from the regions now occupied by the beam filaments. Parameters are $\omega_{pe} \Delta t = 1$, $N_i = N_e = 32$; typically χ

Fig. 3-31. Early electron-beam filamentation.

Fig. 3-32. Coalesced electron-beam filaments.

Fig. 3-33. Reduction in target electron density at the location of the electron beam.

Fig. 3-34. Azimuthal magnetic field encircling coalesced electron beam.

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≈ 1.0 and $\zeta \approx 0.05$. Figure 3-34 shows the x and y components of the magnetic field that develop around the beam filaments.

This code should have wide application as it matures. We hope to use it to investigate the magnetic configurations that are believed to exist around laser focal spots on ICF targets. Other possibilities include the investigation of magnetic reconnection in ICF targets and of related phenomena in the earth's magnetotail. Eventually we expect that a cylindrical version of the code, able to treat various components of the plasma with a fluid representation, will be useful in investigations of particle-beam physics for ICF and other applications.

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Electron-Ion Collisions in Partially Ionized High-Z Plasmas

Laser energy absorption and subsequent heat transport in high-Z plasmas illuminated by short-wavelength, long-pulse laser beams of modest intensity are both controlled by electron-ion collisions. Commonly used expressions for the plasma electrical and thermal conductivities are in serious error for physical parameters of practical interest (as are laser absorption opacities, which are intimately connected to electrical conductivities). The improved low- and high-frequency electrical conductivities presented here are significantly lower than those given by the naive expressions, because we include ion-ion screening when calculating the scattering potential about an ion. Similarly, our improved thermal conductivity is significantly lower because we include the contributions of higher-energy electrons that penetrate the bound-electron screening cloud and are scattered by the nuclear charge, rather than by the ionic

charge. (We note in passing that all this physics is important in computing bremsstrahlung rates as well.)

Elementary considerations display the physics of interest. The collision cross section σ_{ei} for electrons colliding with stationary point ions of charge Z is roughly $\pi r_{ei}^2 = \pi (2Ze^2/m_e v_e^2)^2$, where r_{ei} is the interaction distance at which the potential and kinetic energies are equal. In fact, the cross section is considerably larger than this because of the long range of the screened Coulomb interaction, $(-Ze^2/r) \exp(-r/\lambda_D)$, which is only cut off by plasma screening at large impact parameters b : the correct expression is

$$\sigma_{ei} \approx \pi r_{ei}^2 \ln \Lambda_{ei} \quad (116)$$

Here $\ln \Lambda_{ei} = \ln(b_{\max}/b_{\min}) > 1$, where $b_{\max} \sim \lambda_D$ is the screening length and $b_{\min} \sim r_{ei}$ is fixed by the onset of large-angle scattering. The large-angle scattering contribution to the cross section is about πr_{ei}^2 , or a factor of $\ln \Lambda_{ei}$ smaller than the multiple-small-angle contribution of interest here. This article concerns the factors Z^2 and $\ln \Lambda$ in Eq. (116).

Essentially, σ_{ei} enters calculations of both the laser absorption and the electron heat conduction coefficients. In plasma science, the most commonly used expression for both λ_s and $b_{\max}$ is the electron Debye screening length $\lambda_{De} = v_e/\omega_{pe}$, where $m_e v_e^2 = T_e$ is the electron thermal energy. For laser absorption, the very high frequency of the laser field, $\omega_L > \omega_{pe}$, can introduce the modification $b_{\max} \sim v_e/\omega_L$, which reflects the effective reduction of the screening length because of the weak absorption of laser energy that occurs when an electron jitters rapidly while colliding with an ion. Collisions with larger impact parameters in the presence of the laser light tend toward adiabatic and thereby conserve energy, so that negligible absorption results.

Our primary improvement of the laser absorption opacity consists of recognizing an improvement to the above discussion that has already been made for the thermal transport coefficient.⁹¹ Screening of ions by other ions is dominant in our laser-produced plasmas. It is well known that the total screening length λ_s in a plasma is given by

$$\frac{1}{\lambda_s^2} = \frac{1}{\lambda_{De}^2} + \frac{1}{\lambda_{Di}^2} \quad (117)$$

where $\lambda_{Dn} = V_e/\omega_{pe}$ and where $\lambda_e \approx \lambda_{Dn} < \lambda_{Dn}$, whenever $Z^*T_e > T_e$, as is typically the case in the laser-produced plasmas already described. (Here Z^* is ionic charge.) A physical picture of this is that in establishing its screening cloud, an ion more readily displaces other ions from its vicinity than it attracts electrons.

The reduction of the effective screening length is often not as severe as indicated by Eq. (117), however; this fact also has already been recognized in the thermal transport coefficient calculation.⁴⁴ For the laser plasmas of interest r_0 , the ion sphere radius defined by $4\pi r_0^3 N_i/3 = 1$, comes into play. Our plasmas are typically not conventional plasmas at all—rather, they are dense plasmas, for which $4\pi r_{De}^3 N_e/3 = (\lambda_{De}/r_0)^3 \sim 1$. In this limit, Debye screening goes over to neutral-cell screening, which is characterized by the length r_0 . Thus we have

$$b_{\text{max}} \approx \min(r_0/\omega_e, \max(\lambda_e, r_0)) \quad (118)$$

Typically $\lambda_e \ll r_0 \ll r_0/\omega_e$, so that $b_{\text{max}} \sim r_0$ obtains, and the very-high frequency effect is negligible to occur in laser plasmas.

Our improvement of the heat transport coefficient arises from considering electron collisions with partially ionized atoms. A reduction of the electron thermal conductivity results from the increase in σ_{ei} for higher-energy electrons, for which b_{min} becomes smaller than r_0 , the radius that characterizes the bound-electron screening cloud about the nuclear charge Z_n . The energy-transporting electrons may collide with the ionic nuclei themselves; these collisions will have a cross section that is larger by roughly the factor $(Z_n/Z^*)^2$. Corrections associated with the Coulomb logarithm, which mitigate this effect, must be included in the calculation.

Electron collisions with the nuclear charge also increase the laser absorption strength, but only weakly because the modest-energy electrons that dominate this process are rarely energetic enough to penetrate the bound-electron cloud.

Quantum considerations also enter the calculation of collisions in plasmas, but they are rarely of significance for the systems of interest here. The classical minimum impact parameter b_{min} is determined by the onset of large-angle scattering [$\theta(b_{\text{min}}) = \pi$]; large-angle scattering may also be identified in the quantum treatment as $\theta_{\text{min}} \approx \pi$,

which emerges naturally from the same physics. The quantum analysis is logarithmically divergent in the limit of small scattering angles (as in the corresponding large-impact-parameter classical limit), but obtains its finite result from uncertainty physics.⁴² Quantum collisions on a screened Coulomb potential result in an effective minimum scattering angle $\theta_{\text{min}}^Q \approx \lambda_{De}/\lambda_e$, where the deBroglie wavelength $\lambda_{De} = h/m_e v_e$. Whenever the classical minimum scattering angle $\theta_{\text{min}} \approx \theta(b_{\text{min}})$ falls below the quantum value, the quantum value must be used. Thus we have $b_{\text{min}}^Q \approx b_{\text{min}}^C \min(1, \lambda_{De}/\lambda_e)$. For plasmas of interest this occurs only in the bremsstrahlung emission calculation: this is because the effective screening length of Eq. (118) is quite small, so that θ_{min}^Q usually greatly exceeds θ_{min}^C for those electrons that dominate the processes of laser absorption and heat transport.

Below, we analyze electron scattering on a partially ionized atom in a neutral cell. We obtain a cross section that includes contributions from the ionic and nuclear charges, and use it to calculate the laser absorption and electron heat conduction coefficients. We give some results for laboratory experimental conditions.

Model Scattering Cross Section.

Figure 3-35 shows an idealized neutral ion sphere in a dense, partially ionized plasma. The Z^* free electrons are modeled as a shell of charge at radius r_0 ; the $Z_n = Z^*$ bound

Fig. 3-35. (a) Model neutral ion sphere, with a shell of Z^* free electrons at r_0 and a shell of $Z_n = Z^*$ bound electrons at the atomic radius r_n ; the neutralizing ionic charge Z_n is at the origin. (b) The associated piecewise-Coulombic potential.

electrons are modeled as a shell of charge at radius r_a . Thus the charge density ρ in Poisson's equation for the electrostatic scattering potential, $\nabla^2 \phi = -4\pi\rho$, is

$$\rho = \frac{eZ_n \delta(r)}{4\pi r^2} = \frac{e(Z - Z^*) \delta(r - r_a)}{4\pi r_a^2} - \frac{eZ^* \delta(r - r_0)}{4\pi r_0^2} \quad (119)$$

Thomas-Fermi scaling and practical experience give $r_a \sim a_0/Z_n^{1/3}$ for highly ionized, high- Z atoms (we have not included the weaker dependence on Z^*); here a_0 is the Bohr radius. The model scattering potential is piecewise Coulombic and is readily manipulated in orbit calculations.

The momentum scattering cross section is

$$\sigma_{ei}(v) = \int d\sigma (1 - \cos \theta) \quad (120)$$

and the associated collision term on the right hand side of the free-electron Fokker-Planck transport equation is

$$\frac{df_{ei}}{dt} = \nabla_v \cdot \left[\frac{N_i \sigma_{ei} v}{2} (v^2 \mathbf{1} - \mathbf{v} \mathbf{v}) \cdot \nabla_v f_{ei} \right] \quad (121)$$

In Eq. (120), θ is the usual scattering angle and $d\sigma = 2\pi b db$, where (again) b is the collision impact parameter.

Straight-line orbit theory in the model potential gives

$$2 \sin \frac{\theta}{2} = \frac{Z^* e^2}{\epsilon b} \left(1 - \frac{b^2}{r_0^2} \right)^{1/2} + \frac{(Z - Z^*) e^2}{\epsilon b} \left(1 - \frac{b^2}{r_a^2} \right)^{1/2} \quad (122)$$

where $\epsilon = m_e v^2/2$ is the electron kinetic energy. If the argument of either square root is negative, the corresponding term is set equal to zero.

From Eq. (120), the momentum scattering cross section is found to be

$$\sigma_{ei} \approx 2\pi b_{\min}^2 + \frac{\pi e^4}{\epsilon^2} (Z^2 \ln \Lambda) \quad (123)$$

where $b_{\min}$ is that value of b in Eq. (122) for which $\theta = \pi$. The quantity $Z^2 \ln \Lambda$ represents the energy-dependent scattering strength of the model potential; it is given by

$$Z^2 \ln \Lambda \approx Z^2 \left[\ln \left(\frac{r_0}{b_{\min}} \right) - \frac{1}{2} + \frac{1}{2} \frac{b_{\min}^2}{r_0^2} \right] + (Z - Z^*)^2 \left[\ln \left(\frac{r_a}{b_{\min}} \right) - \frac{1}{2} + \frac{1}{2} \frac{b_{\min}^2}{r_a^2} \right] + Z^*(Z - Z^*) \cdot \left\{ \ln \left[\frac{1 + \left(1 - \frac{b_{\min}^2}{r_a^2} \right)^{1/2}}{1 - \left(1 - \frac{b_{\min}^2}{r_0^2} \right)^{1/2}} \right] + 2 \left(1 - \frac{b_{\min}^2}{r_0^2} \right)^{1/2} \right\} \quad (124)$$

To obtain Eq. (124) we have exploited the fact that $r_a < r_0$. To this same accuracy we have, for $b_{\min} > r_a$

$$b_{\min} \approx \frac{r_0 [Z^* e^2 / 2\epsilon]}{\left[r_0^2 + (Z^* e^2 / 4\epsilon^2) \right]^{1/2}} \quad (125a)$$

and, for $b_{\min} < r_a$

$$b_{\min} \approx \left\{ r_a^2 \frac{Z^* e^2}{2\epsilon} + r_a \frac{(Z - Z^*) e^2}{2\epsilon} \times \left[r_a^2 + \frac{(Z - Z^*)^2 e^4}{4\epsilon^2} - \frac{Z^* e^4}{4\epsilon^2} \right]^{1/2} \right\} \times \left[r_a^2 + \frac{(Z - Z^*)^2 e^4}{4\epsilon^2} \right]^{-1} \quad (125b)$$

Absorption of Laser Beams in Plasmas. A Wentzel-Kramers-Brillouin (WKB) treatment of laser-beam propagation in plasma is often sufficient. In one (spatially inhomogeneous) dimension, the intensity I_1 of the laser beam varies as

$$\frac{dI_1}{dx} = -\kappa_1 I_1 \quad (126)$$

where κ_1 is the plasma opacity. For an LTE plasma (one in which f_e has a Maxwell-Boltzmann distribution), the opacity is

n_e (cm $^{-3}$)	r_0 (cm)	I_0 (keV)	r_d (cm)	λ_L (μ m)	$(Z^2 \ln \Lambda)_L$	$Z^2 \ln (2.25 T_e / h \omega_L)$
1×10^{21} , 2.12×10^{-7}	9	1.23×10^{-9}	1.06	9.83×10^3	1.56×10^4	
1×10^{21} , 2.12×10^{-7}	9	0	1.06	8.60×10^3	1.56×10^4	
4×10^{21} , 1.34×10^{-7}	6	1.23×10^{-9}	0.53	7.83×10^3	1.39×10^4	
4×10^{21} , 1.34×10^{-7}	6	0	0.53	7.24×10^3	1.39×10^4	
9×10^{21} , 1.02×10^{-7}	3	1.23×10^{-9}	0.35	5.78×10^3	1.21×10^4	
9×10^{21} , 1.02×10^{-7}	3	0	0.35	5.74×10^3	1.21×10^4	

Table 3-4. Laser absorption coefficients for electron-ion collisions in the potential of Fig. 3-35 for gold with $Z^* = 40$. Scattering on the nuclear charge is not significant; neutral-cell screening reduces the opacity by about a factor of 2.

$$\kappa_1 = \frac{4}{3(2\pi)^{1/2}}$$

$$\times \left(\frac{T_e}{m_e v^2} \right)^{1/2} \frac{n_e}{n_i (1 - n_e/n_i)^{1/2}} \frac{1}{\lambda_i} \quad (127)$$

where n_e is the electron number density, n_i is the laser critical density ($4\pi e^2 n_i / m_e = \omega_i^2$), and λ_i is the opacity-averaged electron-ion collisional mean free path

$$\frac{1}{\lambda_i} = \int d^3v \left(\frac{m_e v^2}{2T_e} \right) \times \frac{4}{\lambda_{ei}} \frac{\partial}{\partial v} \left[\exp \left(- \frac{m_e v^2}{2T_e} \right) \right] \quad (128)$$

where $\lambda_{ei} = 2/N_i \sigma_{ei}$. In the following we neglect the subdominant term $2\pi b_{\min}^2$ in σ_{ei} ; it is straightforward to restore this term if desired.

In the integrand of Eq. (128), the factor $(m_e v^2 / 2T_e)^{1/2} 4/\lambda_{ei}$ varies slowly, through $(Z^2 \ln \Lambda)_i$, while the factor $(\partial/\partial v) \exp(-m_e v^2 / 2T_e)$ peaks at $m_e v^2 / 2T_e \sim 1/2$, so that the effective factor $(Z^2 \ln \Lambda)_i$ in $1/\lambda_i$ is

$$(Z^2 \ln \Lambda)_i \approx Z^2 \ln \Lambda_{i,1/2} \quad (129)$$

Here we mean to evaluate $Z^2 \ln \Lambda$ at the scaled energy given by the subscript: that is, using the $b_{\min}$ that corresponds to $m_e v^2 / 2T_e \sim 1/2$ from Eqs. (125a) or (125b) in Eq. (124). Numerical experience with the integral indicates that 0.6 is quite accurate.

Equation (129) emphasizes that electrons of modest energy ($m_e v^2 / 2T_e \sim 1/2$) determine the opacity of the plasma to a laser beam. Typically, scattering by the nuclear charge makes only weak contributions. The dominant opacity improvements come from

replacing naive expressions for $b_{\max}$ and $b_{\min}$ with those presented here.

Table 3-4 shows $Z^2 \ln \Lambda$ for parameters typical of laboratory experiments and target design calculations. For comparison, we include values obtained with the commonly used expression $Z^2 \ln (2.25 \lambda_{De} / \lambda_{Li})$.

Electron Thermal Conductivity. The momentum scattering cross section also enters the calculation of the electron thermal conductivity. In one spatially inhomogeneous dimension, we are concerned with energy transport in a quasineutral plasma carrying no current. In the spirit of the opacity calculation above, we can show that the heat flux is given by

$$Q_e = - \frac{64}{\sqrt{2\pi}} n_e \sqrt{\frac{T_e}{m_e}} \lambda_i \frac{\partial T_e}{\partial x} \quad (130)$$

where λ_i is the energy-transport-averaged mean free path. It is instructive to consider that as the higher-energy electrons move along the temperature gradient, an electric field is induced that drives a return current of lower-energy electrons to maintain charge neutrality. The effective factor $1/(Z^2 \ln \Lambda)_i$ in λ_i thus has two terms, which partially cancel:

$$\frac{1}{(Z^2 \ln \Lambda)_i} \approx \frac{5}{(Z^2 \ln \Lambda)_{i,1/2}} - \frac{4(Z^2 \ln \Lambda)_{i,1/2}}{(Z^2 \ln \Lambda)_{i,1/2}^2} \approx \frac{1}{(Z^2 \ln \Lambda)_{i,1/2}} \quad (131)$$

Equation (131) emphasizes that higher-energy electrons dominate the heat conduction process; these are, of course, the electrons most likely to suffer reduced mobility on account of scattering on the nuclear charge rather than the ionic charge.

Table 3-5. Electron heat conduction coefficients for electron-ion collisions in the model potential of Fig. 3-35 for gold with $Z^* = 40$. Scattering by the nuclear charge reduces the conductivity by a factor of about 2/3.

n_e (cm ⁻³), r_0 (cm)	T_e (keV)	r_{s1} (m)	$\frac{4!}{(Z^2 \ln \Lambda)_1}$	$\frac{4!}{Z^{*2} \ln [r_0/(Z^* r^2/3T_e)]}$
1×10^{21} , 2.12×10^{-7}	9	1.23×10^{-5}	1.32×10^{-3}	2.17×10^{-3}
1×10^{21} , 2.12×10^{-7}	9	0	2.08×10^{-3}	2.17×10^{-3}
4×10^{21} , 1.34×10^{-7}	6	1.23×10^{-9}	1.59×10^{-3}	2.49×10^{-3}
4×10^{21} , 1.34×10^{-7}	6	0	2.34×10^{-3}	2.49×10^{-3}
9×10^{21} , 1.02×10^{-7}	3	1.23×10^{-9}	2.13×10^{-3}	2.96×10^{-3}
9×10^{21} , 1.02×10^{-7}	3	0	2.75×10^{-3}	2.96×10^{-3}

Table 3-5 gives thermal conductivity coefficients for the same parameters as Table 3-4. For comparison with $1/(Z^2 \ln \Lambda)_1$, we include values given by the commonly used expression $1/[Z^{*2} \ln \{\lambda_{De}/(Z^* r^2/3T_e)\}]$.

For completeness, we give the expression for the electric field E that ensures quasi-neutral heat flow,

$$eI = \int dv v^2 v \lambda_{ei} \frac{\partial f}{\partial x} - \int dv v^2 v \frac{\lambda_{ei}}{m_e v} \frac{\partial f}{\partial v}, \quad (132)$$

and the heat flow expression itself,

$$Q_e = -\frac{4\pi}{3} \int dv v^2 v \frac{m_e v^2}{2} \times \frac{\lambda_{ei}}{2} \left(\frac{\partial}{\partial x} - \frac{eE}{m_e v} \frac{\partial}{\partial v} \right) f, \quad (133)$$

Using

$$f = \frac{n_e}{(2\pi I_e m_e)^{3/2}} \exp\left(-\frac{m_e v^2}{2I_e}\right) \quad (134)$$

in Eq. (132) and substituting the resulting expression for eI into Eq. (133), one may verify Eq. (130) and obtain the expression corresponding to Eq. (131).

We also record the dc electrical conductivity expression, $I = \sigma_1 E$, where

$$\sigma_1 = \frac{4\pi e^2}{6m_e} \int dv v^2 \lambda_{ei} \frac{\partial f}{\partial v} \\ \approx \frac{16}{(2\pi)^{3/2}} \frac{Z^* T_{e1} (T_e/m_e)^{1/2}}{e^2 Z^2 \ln \Lambda_{1/2}}, \quad (135)$$

Here again, the subscript 7/2 suggests that significant reductions may result from the scattering of current-carrying electrons by the nuclear charge in a hot plasma.

Author: J. R. Albritton

Compression of Microwaves by a Moving Compact Torus

It has been conjectured^{1,2} that a fast-moving plasma configuration known as a compact torus (CT) might be amenable to acceleration to high velocities. Should this conjecture prove correct, a large number of applications suggest themselves, including ICF drivers and other configurations that might be of interest in situations requiring compression of energy in time and space.

Another possibility^{3,4} is that moving CTs might be made to compress microwaves. The proposed mechanism is to prefill the coaxial vacuum region in front of a CT with a number of normal electromagnetic modes on which the CT impinges. A crucial assumption is that the CT excludes the microwaves and therefore compresses them. Performance would be degraded should the microwaves penetrate the CT to some degree. The degradation would range from diminished compression efficiency to significant CT heating. In the expected frequency regime^{5,6} and in similar (toroidal) geometries, studies of MFL applications have found electromagnetic radiation capable not only of penetrating but also of heating and driving currents in Tokamak devices. We report here a preliminary investigation of these effects using explicit (ZOHAR) and implicit particle-in-cell (PIC) codes in 1-D modes.

We used a much-simplified model that takes into account the relevant field

orientations but allows gradients only in the axial direction. Figure 3-36 shows the geometry and those parts of the configuration that we modeled. We consider a small region at the interface between the CT and a group of transverse electromagnetic (TE) modes.

Because of the small spatial scale the magnetic field of the CT can be considered to be homogeneous and static. The field of the TE mode is represented by a smaller rf component superimposed on, and in the same direction as, the zeroth-order CT field. Inhomogeneities in directions other than the axial direction are ignored, as are the nonzero gradients in the zeroth-order field. In the particular case reported here (Fig. 3-36), the bias field is $B_0 = 1$ (arbitrary units) for all x ; the density consists of piecewise linear segments with a peak density $N = 1$ for $25 \leq x \leq 35$, and the incident rf amplitude is a relatively large $0.3B_0$. For these choices the plasma electron and cyclotron electron frequencies are equal ($\omega_{ce}/\omega_c = 1$) and the lower-hybrid frequency is $\omega_{lh} = 0.07\omega_{ce}$ at peak density.

Even this simple model requires complicated analysis. Guidance is gained by considering linear dispersion analysis for this field orientation, assuming homogeneous density. Standard analysis¹⁶ gives the relation

$$k^2 = \frac{\omega^2}{c^2} - \sum_i \frac{\omega_i^2 \omega_c^2}{c^2(\omega^2 - \omega_i^2)} \quad (136)$$

$$= \frac{\sum_i \frac{\omega_i^2 \omega_c^2}{\omega^2 - \omega_i^2}}{\omega^2 - \omega_c^2} \quad (136)$$

for an applied electromagnetic (EM) wave of frequency ω , where the summations are over particle species i . We find $k^2 > 0$, indicating propagation of the applied EM wave, in two regions. At very low n , we have a vacuum-like EM mode; for high enough n , the lower hybrid mode propagates. This implies that there is a range of densities for which $k^2 > 0$ and in which no mode can propagate. We find that lower-hybrid

propagation can occur only if the density is high enough that $\omega_{lh} > \omega$, or

$$\omega_{lh} = \sqrt{\frac{n_e}{n_i} \frac{N}{1 + \frac{N}{B_0^2}}} > \omega \quad (137)$$

Fig. 3-36. Simplified model of compact-torus, transverse-electromagnetic (TE) mode geometry (all units arbitrary).

inhomogeneous density profiles, such as that shown in Fig. 3-36, were evaluated using explicit¹⁷ and implicit PIC models (see the article "Direct Implicit Electromagnetic Plasma Code,"¹⁸ earlier in this section). A right-traveling external rf mode with appropriate frequency and amplitude is specified by boundary conditions. The frequency is chosen so that the mode would propagate in the homogeneous peak density region shown in Fig. 3-36. Here we used $\omega = 0.05\omega_{ce}$ and $\omega_{lh} = 0.07\omega_{ce}$.

The inhomogeneous density region is almost entirely a forbidden region, through which a small-amplitude mode cannot propagate and from which it will therefore be reflected to the left. For nonzero rf amplitudes, nonlinear effects will become important; these effects, coupled with the complexities of the inhomogeneity, make it necessary to treat the problem by numerical simulation.

Figure 3-37 shows a typical result of nonlinear penetration of the CT by a large-amplitude ($0.3B_0$) wave. The penetration is most evident in the ion v_x - x phase space, which shows that the ions oscillate in resonance with the incoming wave. In the early stages, the steepening of the density ramp near the critical density gives some evidence of the rf reflection. At this amplitude

Fig. 3-37. Ion v_x - x phase space plot showing rf penetration of a TE mode.

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the phenomenon rapidly becomes non-linear, however, and it is not observable in Fig. 3-37.

In summary, the homogeneous linear model of Eq. (136) suggests that the rf energy might not penetrate if the external-field frequency ω were greater than the ω_{ih} in the interior. Our simulations verify this conjecture. The leading edge of the CT is significantly eroded by a high-amplitude incoming rf field. Thus the CT may be able to compress TE modes if the initial CT peak density gives rise to an ω_{ih} that is less than the initial TE-mode frequency ω . This conclusion depends strongly on our assumption that $\mathbf{k} \cdot \mathbf{B} = 0$, a condition that is probably not satisfied in more realistic multidimensional geometries.

We note a recent suggestion⁷⁴ concerning the nature of the magnetic field that precedes the moving CT and the effect this field may have on rf penetration. We expect that the level of B_0 in front of the CT, and therefore the field on which the rf is superimposed, should fall off to a small fraction of the interior magnetic field. Preliminary investigations verify that an inhomogeneous bias B_0 that falls to a tenth of the interior value on the TF side indeed changes and significantly reduces the penetration.

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LASNEX and Atomic Physics

Introduction

A number of projects discussed in the 1983 Laser Annual reached successful conclusions in 1984.

The new laser ray-tracing package was completed and is being widely used. This

package permits us to model problems in which the laser rays do not all pass through the symmetry axis. New macros for use in LASNEX generator decks model realistic ray bundles with a finite focal spot size. A random-scattering model has been added to study the scattering of laser light by small-scale plasma inhomogeneities.

The Thomas-Fermi-based analytic equation of state (EOS) package has been checked out on a wide variety of materials and temperature-density regimes, and has been improved where necessary. The package has been put into LASNEX and is being tried on real design problems. It yields thermodynamically consistent equations of state; true, separate electron and ion equations of state; correct EOS values for arbitrary materials and mixtures; and entropy values. Current EOS tables have none of these virtues.

The semi-Eulerian, continuous-rezoning hydrodynamics package has been further improved, and a number of laser fusion design problems are now being run in pure Eulerian mode with LASNEX. The Eulerian rezoner is fully second-order (when the mesh is orthogonal) and positivity-preserving (that is, physical quantities do not take on unphysical negative values at the new mesh points). The pure-Eulerian mode is very useful on problems with highly turbulent hydrodynamic motions. Fully Eulerian LASNEX performs very well against Woodward and Colella's⁶⁷ piecewise-parabolic method.

The programmers in the LASNEX group have also been active improving the tools that make LASNEX a responsive and versatile instrument for the user.

A new production-code-oriented graphics library, called PLOTLIB, will be the basis for all LASNEX graphics; it is being widely used in other codes. It includes a complete high-level facility for doing three-dimensional projections (both parallel and perspective).

A major effort in 1984 was directed toward the so-called GOAL version of LASNEX. This new user interface completely eliminates any restrictions on how a user may set or access data. Using a simple Fortran-like language, the user has complete access to every variable in the code and can set, edit, or plot any variable or any mathematical combination of variables.

This can be done at generation or (interactively) as the code is running. The first step, writing a GOAL library that performs these operations on data stored in Z (LASNEX dump) files, was completed in mid-1984, and the library is now being widely used. The next step will be to integrate this library into LASNEX so that all communications with LASNEX go through GOAL. This is being debugged and should be in production use by mid-1985. Conversion utilities are being provided so that LASNEX will still accept its old fixed-format commands.

Author: D. S. Kershaw

Second-Order Monotonic Advection in LASNEX

This article describes the scheme used in the LASNEX continuous rezoner to accomplish remapping ("advection") from an old, distorted mesh to a new one. Much of this scheme has been in place for some time, but wide, routine use began in 1984. Problems run with little or no user intervention and with accurate results; this makes life easier for users, and has led them to undertake problems that were previously too difficult to run at all.

It is common in ICF target design for purely Lagrangian hydrodynamic calculations (in which the fluid is fixed in the mesh) to lead to such severe mesh distortion that the calculation stops prematurely. This is particularly true for shear flows, which can produce very large distortion. The other (Eulerian) extreme of using a fixed mesh is inherently diffusive for material properties such as mass and energy.

To deal with such problems in fully or partially Eulerian calculations, mesh distortion must somehow be reduced without degrading the accuracy of the calculation. Fixing the problem after the calculation has stopped requires user intervention: this is diffusive, since it is first-order (as usually carried out), and it is onerous if it must be done often. With an automatic, continuous rezoner, on the other hand, one must be confident that the results are meaningful.

To give fully trustworthy results, an advection scheme should meet the following requirements:

1. It must be stable and conservative (i.e., not lose or gain any of the advected quantity).
2. It should be transportive (signals must go only in the direction of advection, not backwards).
3. It must have at least second-order accuracy (otherwise, diffusion will eventually spread the advected quantity everywhere).
4. It must be monotonic, not oscillatory. In particular, a negative density could cause an exponential instability in some other physical system.
5. It should be consistent (that is, it should preserve any significant constraints on or between the variables).

The extent to which the LASNEX advection scheme meets these requirements will become clear in the discussion below.

These requirements cannot all be satisfied simultaneously for non-smooth initial data. The essential conflict is between the second, third, and fourth requirements: high-order spatially centered schemes give high accuracy but violate the transportive and monotonicity requirements, while so-called upwind (or donor-cell) schemes are only first-order accurate.

To see how the LASNEX advection scheme works, we start with a first-order upwind scheme and improve the accuracy (this is the approach of van Leer¹⁰).

The donor-cell scheme is illustrated in Fig. 3-38, which shows that it is transportive: it is easy to show that it is monotone as well. Furthermore, since the algorithm does not depend on the quantity being advected, the scheme automatically preserves linear constraints, at least. But since we use one-sided differences we obtain only first-order spatial accuracy.

$$\frac{\Delta \xi_i}{\Delta t} = - \frac{u_R \xi_i - u_L \xi_L}{\Delta x}$$

$$\text{where } \xi_R = \xi_i \quad u_R > 0 \\ \xi_{i+1} \quad u_R < 0$$

$$\xi_L = \xi_{i-1} \quad u_L > 0 \\ \xi_i \quad u_L < 0$$

Fig. 3-38. Basic upwind algorithm.

Fig. 3-39. Advection model using nonzero slope.

Improved Upwind Method. The first-order scheme just described uses what we can call an "ice-block" model, in which a constant value is assumed for the density of the advected quantity throughout the zone. We can improve the accuracy by giving the quantity some structure within the zone. The simplest way to do this is to use a linear slope b derived from the zone-center values, as illustrated in Fig. 3-39; this makes the overall calculation second-order.

In this form, however, the scheme is not monotone; to avoid generating new maxima or minima, we impose a limit on b such that the interface values are bounded by the present maxima and minima:

$$\min(y_i) \leq \left(\bar{y} \pm b \frac{\Delta x}{2} \right) \leq \max(y_i) \quad (138)$$

$$j = i - 1/2, i + 1/2, i + 3/2 \quad .$$

In particular, this means that the slope will be zero next to a step discontinuity or at an extremum. To preserve monotonicity, we have thus given up formal accuracy at some places.

We have also given up on the consistency requirement. Consider linear constraints, for example: imagine a system composed of several materials, so that the mass is given by a sum over all ion types,

$$M = \sum_i x_i A_i \quad (139)$$

where x is the concentration and A the mass number. After we remap each x_i and M separately it will not, in general, be true that

$$M' = f(M) = \sum_i f(x_i) A_i = \sum_i x_i' A_i \quad (140)$$

so that consistency has been lost.

We can obtain consistency by identifying the f_i to which the greatest amount of slope-limiting has been applied, and using that f for all the constituents; this solution tends

toward the first-order scheme. The solution used by LASNEX is to define M' using the old relation on the new constituents:

$$M' = \sum_i x_i' A_i \quad (141)$$

Although this procedure does not guarantee monotonicity for M , it does guarantee positivity.

For nonlinear constraints, things are not so simple. As an example, for fixed total energy, the electron density and temperature in Saha equilibrium satisfy the nonlinear relation

$$\begin{aligned} n_e &= f(T_e, x_1, \dots, x_N) \\ T_e &= g(n_e) \end{aligned} \quad (142)$$

In this case, the only feasible choice is to map the constituent fractions and the temperature, and let the ionization routine determine the new electron density. (A similar remark applies to the total pressure.) To obtain consistency between the new temperature and energy, one should really determine the new energy and density and then invert the EOS function to obtain the desired temperature. Instead, LASNEX uses a single iteration pass to solve for the consistent temperature associated with the new energy.

For more complicated cases, such as determining the non-LTE populations, the remapping reverts to a first-order approach. The general philosophy is to avoid special-case knowledge in the remapping algorithms and to treat complicated mixing problems in the appropriate physics package; in the non-LTE case, the benefits of more careful mixing have not so far seemed worth the cost.

Momentum Remapping. An exception to our rule against special remapping algorithms is made for the velocities. In this case, we want to conserve two moments, the momentum ($\propto u$) and the kinetic energy ($\propto u^2$). Another example of this type is provided by a magnetic field remapping, in which one would like to conserve both flux and magnetic energy (although simple flux conservation is now used).

Since the velocities are point-centered, we temporarily form point-centered mass values,

$$M_i = \frac{1}{2} (M_{i-1/2} + M_{i+1/2}) \quad (143)$$

Now we write a transport equation for the nodal mass and momentum values:

$$\dot{M}_i = M_i - \delta M_{i-1/2} + \delta M_{i+1/2} \quad (144a)$$

$$\begin{aligned} \dot{M}_i u_i = M_i u_i - \delta M_{i-1/2} U_{i-1/2} \\ + \delta M_{i+1/2} U_{i+1/2} \end{aligned} \quad (144b)$$

where the nodal δM_i are derived from the zonal δM_i from an equation like Eq. (143), and where the effective velocity U associated with the mass advection to give the momentum change is as yet unspecified. (In what follows we call U the advection velocity, even though it is in fact a matter velocity.)

Since they are written in so-called conservation form, Eqs. (144) will conserve mass and momentum for an arbitrary advection velocity. To proceed, we multiply Eq. (144a) by another power of the velocity and rearrange to get an equation for the kinetic energy:

$$\begin{aligned} \frac{1}{2} \dot{M}_i u_i^2 = \frac{1}{2} M_i u_i^2 \\ - \delta M_{i-1/2} \left[\frac{\bar{u}_i^2}{2} - (\bar{u}_i)^2 + \bar{u}_i U_{i-1/2} \right] \\ + \delta M_{i+1/2} \left[\frac{\bar{u}_i^2}{2} - (\bar{u}_i)^2 + \bar{u}_i U_{i+1/2} \right], \end{aligned}$$

$$\bar{u}_i = \frac{1}{2} (u_i' + u_i) \quad (145)$$

We can now force kinetic energy conservation by an appropriate choice of the advection velocity:

$$\begin{aligned} 2U_{i-1/2} \bar{u}_{i-1} - u_{i-1}' u_{i-1} \\ = 2U_{i+1/2} \bar{u}_i - u_i' u_i \end{aligned} \quad (146)$$

or

$$\begin{aligned} U_{i-1/2} = \frac{1}{2} (\bar{u}_{i-1} + \bar{u}_i) \\ + \frac{1}{2} \frac{(\bar{u}_{i-1})^2 - \bar{u}_{i-1}'^2 - (\bar{u}_i)^2 + \bar{u}_i'^2}{\bar{u}_{i-1} - \bar{u}_i} \end{aligned} \quad (147)$$

In other words, the advection velocity is the average of the nodal velocities plus a correction term. Since the $\bar{u}_i$ involve the new nodal velocities, this is an implicit nonlinear equation, which we solve by a Newton iteration. In the two-dimensional case, each node is coupled to its eight nearest neighbors; the linearized equations for the Newton iteration are solved by a nonsymmetric conjugate-gradient method.²⁹

We have not imposed any monotonicity requirements on the velocities, so we check three simple cases for reasonable behavior.

1. Uniform Velocity Field. It is simple to show that

$$\begin{aligned} u_{i+1} = u_i = U \rightarrow U_{i+1/2} = U \\ \rightarrow u_{i+1}' = u_i' = U \end{aligned} \quad (148)$$

It is also possible to show that $\bar{u}_{i+1} = \bar{u}_i$ implies that

$$u_{i+1}' = u_i' = u_i = u_{i+1} \quad (149)$$

Thus a uniform velocity field remains uniform.

2. Velocity Reversal. Consider the case shown in Fig. 3-40, in which nodes A and B have equal masses but opposite velocities, so that the equations are simple. Then with a small mass transfer δ we have

$$\begin{aligned} (M + \delta) u_A + MU = \delta \cdot U_{ad} \\ (M - \delta) u_B - MU = -\delta \cdot U_{ad} \end{aligned} \quad (150)$$

The new velocity gradient ratio is

$$\frac{\Delta u'}{\Delta u} = \frac{u_B - u_A}{2U} = \frac{1 - \frac{\delta}{M} \frac{U_{ad}}{U}}{1 - \left(\frac{\delta}{M} \right)^2} \quad (151)$$

Fig. 3-40. Nomenclature for velocity-reversed example.

We could pick the advection velocity U in several ways. For the explicit guess

$$U_{ad}^{exp} = \frac{u_A + u_B}{2} = 0 \quad (152)$$

the change in the velocity gradient is second-order in the amount of advection.

$$\left(\frac{\Delta u}{\Delta x}\right)^{exp} = 1 + \left(\frac{\delta}{M}\right)^2 + \dots \quad (153)$$

For small δ/M , therefore, nothing unreasonable happens to the new velocities. Furthermore, this ratio is a function of $(\delta/M)^2$ only (that is, it is independent of the sign of δ). Both of these statements are also true for the exact U_{ad} :

$$U_{ad}^{exact} = \frac{\frac{1}{2} \frac{\delta}{M}}{\frac{1}{2} \left[1 + \sqrt{1 - \left(\frac{\delta}{M}\right)^2} \right]} \quad (154)$$

and

$$\left(\frac{\Delta u}{\Delta x}\right)^{exact} = 1 + \frac{1}{2} \left(\frac{\delta}{M}\right)^2 + \dots \quad (155)$$

Finally, if we use a backward Euler estimate for U_{ad} , then the velocity-gradient ratio is exactly unity:

$$U_{ad} = \frac{u_A + u_B}{2} - \left(\frac{\Delta u}{\Delta x}\right)^{he} = 1 \quad (156)$$

Fig. 3-41. Pure advection test case: (a) initial condition, (b) first-order result, (c) second-order result.

3. *Kinetic Energy Gradient.* We examine the kinetic energy for the velocity-reversed case of Fig. 3-40. For the sum of kinetic energies we have

$$\frac{(M + \delta)u_A'^2 + (M - \delta)u_B'^2}{2Mu^2} = \frac{\left(1 - \frac{\delta}{M} \frac{U_{ad}}{u}\right)^2}{1 - \left(\frac{\delta}{M}\right)^2} \quad (157)$$

so for any case in which $U_{ad} = O(\delta/M)$ at least [or $U_{ad} = 0$ for the explicit guess of Eq. (152)], we obtain exact momentum conservation as well as second-order kinetic energy conservation. The explicit guess is, in fact, used by LASNEX to initiate the Newton iteration; if no iteration is permitted, we still obtain second-order accuracy for the total kinetic energy.

In the case of the kinetic energy gradient, we have

$$\frac{(M + \delta)u_A'^2 - (M - \delta)u_B'^2}{2Mu^2} = -\frac{\delta}{M} \frac{\left(1 - \frac{\delta}{M} \frac{U_{ad}}{u}\right)^2}{1 - \left(\frac{\delta}{M}\right)^2} \quad (158)$$

so that even for the exact U_{ad} , for which the second factor on the right-hand side in Eq. (158) is unity, we have a first-order kinetic energy gradient ratio.

Sample Problems. To demonstrate the effect of the improved advection scheme, we show in Fig. 3-41 a simple problem of pure advection of an initially triangular density perturbation moving to the right at constant velocity. Figure 3-41(a) shows the initial situation, and Figs. 3-41(b) and 3-41(c) show the results for the first- and second-order schemes, respectively. The greater diffusion in the first-order scheme (which would continue indefinitely) is clearly evident.

Since the example of Fig. 3-41 had a uniform velocity field, it does not provide a serious test of the velocity remapping. Figure 3-42 shows the full-Eulerian solution and the corresponding exact results for a

Fig. 3-42. Full-Eulerian and exact results on a shock-tube problem with initial boundary at $x = 3$: (a) density, (b) pressure, (c) internal energy, and (d) velocity.

shock-tube problem: the agreement is quite good, and there are no spurious oscillations. (We assume that the higher propagation speed in the Eulerian calculation arises from residual diffusion, because use of higher resolution eliminates the effect.)

Extension to Two Dimensions. The advection scheme has been straightforwardly extended to two dimensions by means of operator splitting. We have not implemented corner coupling (i.e., coupling with diagonal neighbors): including corner-coupling terms (which are not necessary for second-order accuracy on an orthogonal mesh) would be straightforward given the coefficients, although twice as much work would be involved. Since the velocities are resolved into orthogonal components, the momentum and kinetic energy equations decouple into two scalar equations in two dimensions; again, this entails only twice as much work as in one dimension.

Author: D. S. Bailey

Statistical Model for Light Scattering by Density Fluctuations

Because of limitations on storage space and computational time, a macroscopic simulation program such as LASNEX must employ a computational mesh that is too coarse to resolve the steep, inhomogeneous plasma density structure near critical density. Thus, we must use phenomenological models to predict the effects of subgrid-scale density fluctuations that arise from such sources as uneven illumination and hydrodynamic instability. To this end we have developed a statistical model of the refractive scattering of laser light by random density fluctuations in the subcritical plasma, and have incorporated it into the LASNEX 3-D ray-trace package.¹⁰⁰

The existence of such fluctuations is suggested by disk experiments, which reveal a weaker dependence of absorption on angle than is implied by a simple theory based on

a smooth density profile. In addition, the nonabsorbed light is typically observed to emerge with an angular distribution too broad to be consistent with specular reflection.¹⁰¹⁻¹⁰⁴

In an earlier, simpler (and still-available) LASNEX model of subgrid effects, rays are treated as scattering only at their turning points; user parameters specify the probabilities of forward and backward scattering and the angular width of the scattered distribution. This model is useful, but it entails certain undesirable biases. First, normally incident rays (parallel to ∇n) cannot scatter at all, since their group velocity is zero at the turning point, and nearly normal rays scatter only a little. Second, the scattered-light distribution is biased toward the normal; at $x = \infty$, all scattered rays have $y < y_{\text{unscattered}}$ (see Fig. 3-43). This is because, at the turning point, the unscattered ray has the greatest possible velocity component perpendicular to the density gradient. Finally, the turning-point scattering model affects the inverse bremsstrahlung path length on the outbound leg only.

For these reasons, we have added to LASNEX a statistical model incorporating scattering along the entire ray path. In this model, a ray trajectory is computed by assuming that a point on the ray moves at the local group velocity of light waves, and

Fig. 3-43. Ray paths through a perturbed density profile. 1, unperturbed path; 2, 3, paths after scattering at turning point; 4, typical statistically scattered path.

advancing it through a sequence of time steps. The model perturbs the direction of the ray after each step. The development of the model consists primarily of a calculation of how much scattering to impose. Strictly speaking, our geometrical-optics approach means that we can consider only density fluctuations whose scale sizes are large compared with the laser wavelength. However, it might be expected that our model can treat other incoherent scattering processes with a suitable reinterpretation of the parameters.

The physical assumptions behind this model are similar to those used in a model developed at the University of Rochester Laboratory for Laser Energetics (UR-LLE).¹⁰⁵ Two approaches were followed in the UR-LLE work. The first was a Monte Carlo calculation in which a refracting ray path was computed through each member of a small ensemble of plasmas in which different realizations of the fluctuating density were superimposed. In the second, equations describing the spreading of the envelope of these rays were integrated along the unperturbed ray trajectory. Neither approach was suitable for our application, though both provide useful benchmarks for comparison.

For simplicity, we assume that the nominal scale size d of a density clump (or depression) is the same as the nominal spacing between clumps. Thus d is the correlation length of what we can call the scattering "force." In terms of the normalized zero-order density $N(x) = n^0(x)/n_c$, where n_c is the critical density, the acceleration of a ray point is given by¹⁰⁰

$$\ddot{\mathbf{x}} = -\frac{c^2}{2} \frac{\partial N}{\partial x} \hat{\mathbf{x}} + \mathbf{F} \quad (159)$$

where $\mathbf{F}$ is a (small) randomly directed force of magnitude

$$F = c^2 \frac{\delta n}{2d n_c} \quad (160)$$

and δn is a typical density fluctuation. The change in velocity (in the x direction, for example) due to $\mathbf{F}$ during a step from t_0 to t_1 is

$$\delta \dot{\mathbf{x}} = \int_{t_0}^{t_1} \mathbf{F}_x(t) dt \quad (161)$$

The mean velocity change $\langle \delta \dot{x} \rangle$ will be zero, but the mean-square velocity change $\langle \delta \dot{x}^2 \rangle$ will be nonzero:

$$\langle \delta \dot{x}^2 \rangle = \left\langle \left[\int F_x(t) dt \right] \left[\int F_x(t') dt' \right] \right\rangle. \quad (162)$$

The integrals can be brought outside the (ensemble) average, and we can assume that the magnitude of F is constant on the scale of a step. Averaging F_x^2 over solid angle to give $\langle F^2 \rangle / 3$, we have

$$\langle \delta \dot{x}^2 \rangle = \frac{\langle F^2 \rangle}{3} \int dt \int dt' G(|\mathbf{x}(t) - \mathbf{x}(t')|), \quad (163)$$

where for the correlation function we choose

$$G(\xi) = 0, \quad |\xi| \leq d \\ = 1, \quad |\xi| > d, \quad (164)$$

that is, the fluctuating force (from a given clump) extends over a distance d along the ray path. Other expressions for G lead to essentially the same result.

With this form for G , the integral in Eq. (163) reduces to a computation of the area in the t - t' plane shown shaded in Fig. 3-44:

$$t_0 \leq t \leq t_1,$$

$$t_0 \leq t' \leq t_1, \quad (165)$$

$$|\mathbf{x}(t) - \mathbf{x}(t')| \leq d.$$

Using the equation of motion (159), the unperturbed orbit is

$$x = -\frac{c^2}{4} \frac{\partial N}{\partial x} t^2, \\ y = \dot{y} t, \quad (166)$$

where after each step we choose the zero-point of time to correspond to that (past or future) time at which $x = 0$, i.e., the turning point of the instantaneous trajectory, which will vary slightly from step to step. The

Fig. 3-44. Area in the t - t' plane associated with the scattering integral (see text), for 10° incidence and $|L \cdot \nabla N| = 0.1$.

separation between points along the trajectory is

$$|\mathbf{x}(t) - \mathbf{x}(t')| = |t - t'| \\ \times \left[\dot{y}^2 + \frac{c^2}{4} \nabla N (t - t')^2 \right]^{1/2}. \quad (167)$$

With the change of variables

$$\tau = \frac{1}{2}(t + t'), \\ \sigma = t - t', \quad (168)$$

and the definitions

$$\tau_0 = \left| \frac{\dot{y}}{c^2 \nabla N / 2} \right|, \\ \sigma_0^2 = \left| \frac{d}{c^2 \nabla N / 2} \right|, \quad (169)$$

the integral becomes

$$\int dt \int dt' G = 2 \int_{t_0}^{t_1} \frac{\sigma_0^2}{\tau^2 + \tau_0^2} d\tau. \quad (170)$$

Fig. 3-45. Ray paths for 20° incidence. (a) No scattering; (b) 50 rays scattered using $\langle \delta n^2 \rangle / n_c^2 d = 5.32 \times 10^{-3}$.

We have used the fact that the correlation length is small compared with the distance traveled in a time step, so that the shaded area in Fig. 3-44 is long and thin, and thus the colored area is negligible. This breaks down for a small range of angles near normal incidence. We enforce a minimum τ_c of 10^{-10} .

To evaluate Eq. (163), we carry out the integration in Eq. (170), substitute for $\langle F^2 \rangle$ from Eq. (160) and for σ_z^2 from Eq. (169), and obtain

$$\langle \delta n^2 \rangle = \frac{c^2}{3} \frac{\delta n^2}{n_c^2} \frac{L}{d} \log \left(\frac{t_1 + \sqrt{t_1^2 + \tau_c^2}}{t_0 + \sqrt{t_0^2 + \tau_c^2}} \right), \quad (171)$$

where the zero-order scalelength $L = |\nabla n|^{-1}$. Identical expressions are obtained for $\langle \delta y^2 \rangle$ and $\langle \delta z^2 \rangle$.

Near grazing incidence, τ_c is large compared with t_0 and t_1 and the logarithm can be expanded to yield the rms angular deflection due to the step:

$$\frac{1}{c} \langle \delta x^2 \rangle^{1/2} \approx \frac{1}{\sqrt{6} n_c} \langle \delta n^2 \rangle^{1/2} \left[\frac{c}{d} (t_1 - t_0) \right]^{1/2}. \quad (172)$$

The deflection is proportional to the square root of the step size, as expected for a uniform random process.

Since the ray traverses many correlation lengths in a step, we can regard the deflections after each step as independent. We calculate t_0 and t_1 relative to the turning point, and use Eq. (171) to calculate $\langle \delta x^2 \rangle$ due to the step. A random velocity perturbation for each coordinate direction is calculated, with mean zero and variance $\langle \delta x^2 \rangle$. These are added vectorially to the end-of-step velocity, which is then rescaled to restore the speed to its original value (i.e., to v_{group} in the unperturbed medium). This scheme has the virtue of "failing gracefully" for very large perturbation velocities: in that limit, it just randomizes the ray direction.

Various options are available. The primary input parameter is $\langle \delta n^2 \rangle / n_c^2 d$. This can be interpreted as an inverse path length for randomization of the ray direction, for a ray moving in a low-density background with fluctuations characterized by $\langle \delta n^2 \rangle$ and d . Alternatively, one can specify $\langle \delta n^2 \rangle / n_c^2$. In this case, d is adjusted by the code to satisfy $d = L/10$. This option tends to increase the scattering during the peak of the pulse, when the path length through the steepened profile is short and the scattering would be minimal if d were fixed. In addition to either option, the user can also specify a dependence of density fluctuation amplitude on density, with $\langle \delta n^2 \rangle^{1/2} \propto [n(x)/n_c]^\gamma$ (where γ is an input parameter, whose default value is $\gamma = 1$), as well as a density cutoff below which scattering does not occur. The test problems described below were run with $\gamma = 0$.

A step-size control ensures that enough steps are taken for the trajectory to respond locally to velocity deflections. The control depends on a user-specified parameter a , via $\Delta t_{\text{max}} = aL/c$. Values of a between 0.04

and 0.4 are commonly used. In a linear profile, a normally incident ray would require 100 steps to complete its trip if $\alpha = 0.04$. This is not overly expensive, since in most zones the gradient is weak and the control is not active. It is important to take small steps when $\langle \delta n^2 \rangle$ varies strongly with position. In most other cases, the results seem to be insensitive to the value of α ; this is an important check of the model.

Figure 3-45(a) shows the trajectory of an unscattered ray incident at 20° . In Fig. 3-45(b) the fluctuation amplitude obeys $\langle \delta n^2 \rangle L / n_c^2 d = 5.32 \times 10^{-1}$. The step-size control allowed moderately large steps ($\alpha = 0.25$). Other runs with $\alpha = 0.04$ and 0.01 gave very similar results.

Figure 3-46(a) shows a run with $\langle \delta n^2 \rangle L / n_c^2 d = 8.51 \times 10^{-2}$, while Fig. 3-46(b) has $\langle \delta n^2 \rangle L / n_c^2 d = 0.532$. In both cases, $\alpha = 0.04$; if one assumes $L = 10d$, this corresponds to $\delta n_{rms} / n_c = 0.23$, which is at the limit of the model's validity.

When applied to a round-trip path from plasma entry to exit in a linear density profile, the expression for the deflection becomes

$$\begin{aligned} \delta \theta_{\text{exit}}^2 &\approx \frac{\delta n^2}{c^2} \\ &\approx -\frac{2}{3} \frac{\delta n^2 L}{n_c d} \log \tan \left(\frac{\theta_i}{2} \right) \quad (173) \end{aligned}$$

where the angle of incidence θ_i obeys $\cos \theta_i = N_{ip}^2$, and N_{ip} is the density at the turning point. At normal incidence, this expression diverges as the logarithm of θ_i , while at grazing incidence it approaches zero linearly. This calculation neglects the fact that the effect of scattering in the unperturbed plane of refraction is partially nullified by refractive focusing effects. A more careful calculation leads to a lower estimate of the in-plane scattering:

$$\begin{aligned} \delta \theta_{\text{exit in-plane}}^2 &\approx \frac{2}{3} \frac{\delta n^2}{n_c^2} \\ &\times \frac{L}{d} \left(\frac{\log \tan (\theta_i/2)}{\cos^2 \theta_i} + \frac{1}{\cos \theta_i} \right) \quad (174) \end{aligned}$$

At grazing incidence, the rms in-plane scattering angle on exit is ~ 0.577 times that

Fig. 3-46. Paths of 50 rays at 20° incidence: (a) $\langle \delta n^2 \rangle L / n_c^2 d \approx 8.51 \times 10^{-2}$; (b) $\langle \delta n^2 \rangle L / n_c^2 d = 0.532$.

out-of-plane, while at normal incidence (where there is no distinction between in-plane and out-of-plane) both diverge identically. Figure 3-47 shows the dependence of these rms scattering angles on θ_i . The figure also shows the results of test runs in which 1000 rays were traced using the model with $\langle \delta n^2 \rangle L / n_c^2 d = 10^{-1}$ and $\alpha = 0.01$. The agreement is excellent for $\theta_i \geq 30^\circ$ and is reasonable for smaller angles.

The angle dependence of inverse-bremsstrahlung absorption is reduced by the scattering. This accords with the experimental observation¹⁰¹⁻¹⁰⁴ that absorption in general has a weaker dependence on angle than expected for a smooth density profile. Figure 3-48 shows the calculated classical

Fig. 3-47. RMS scattering angles on exit vs angle of incidence (see text).

Fig. 3-48. Absorption fraction vs angle of incidence (see text).

absorption fraction vs θ_i , with an overall coefficient chosen to give 80% absorption at normal incidence. Results for three different fluctuation amplitudes are also shown. In all cases, $c \Delta t/L = 0.04$, and there were 250 rays per run. These results are in close agreement with those of the UR-LLE group.¹⁰⁵ (The value $\langle \delta n^2 \rangle L/n_c^d = 5.32$ is probably unreasonably large.)

We expect that our model will improve LASNEX ray-trace calculations in plasmas with rapid spatial density variations, and that forthcoming Nova experiments with tilted disks should provide good tests of its accuracy.

Authors: A. Friedman and E. A. Williams

Rapid Calculation of Properties of Plasma Atoms and Ions

Plasmas are a form of matter for which very little measured data exists. For a description of properties such as equations of

state, opacities, and conductivities, we often depend on calculations from theoretical atomic physics. Underlying much of the calculated atomic data are crude models that are far less accurate than the best models available but that continue in use because of their convenience and efficiency and because of the complexity of the better models. We describe here two fast, simple schemes for obtaining self-consistent field models of ions in plasmas. We expect such a scheme to replace, for example, the screened-hydrogenic model used in large computer programs such as LASNEX.

A plasma often contains atoms in a very large number of states of ionization and excitation. A self-consistent field calculation of any one of these ionic states is not hard (taking only a few seconds on our best computers), but dealing with many thousands of such calculations is a formidable task. Both of our schemes are based on the notion that a self-consistent field calculation for a typical ionic state can be adapted to other states that do not differ from it by too much. When a reasonable choice is made for the "typical" state, and nearby states will dominate in determining the important physical properties of the plasma. Economies of several orders of magnitude can be achieved by these schemes, which thereby make superior atomic physics models available for the calculation of the macroscopic plasma properties.

The interaction of radiation with matter depends critically on the frequencies of absorption lines and the energies of absorption edges. Both of these are obtained directly from the total energy of an ion in its initial and final states, so we concentrate here on the calculation of the total energy.

In the first scheme,¹⁰⁶ a self-consistent field calculation is carried out for a particular ionic configuration specified by orbital occupation numbers n_i^l ($i = 1, 2, \dots, N$). That is, for the prescribed n_i^l (which need not be integers) the variational principle is used to minimize the total ion energy with respect to variations in the orbital wave functions. A set of orbital functions $\phi_i^l(r)$, ($i = 1, 2, \dots, N$), is obtained as part of the calculation.

A nearby configuration is not computed directly; rather, the orbital occupation numbers are changed to n_i , and the orbital

functions are scaled to $\varphi_i(r) = \lambda_i^{3/2} \varphi_i^0(\lambda_i r)$. The total energy, which is a functional of the $\varphi_i(r)$, is then minimized with respect to the parameters λ_i , again using the variational principle. In carrying this out, the total energy is expressed as a Taylor series in the small parameters $\lambda_i - 1$, retaining only linear and quadratic terms (which means that the computational problem is reduced to linear algebra). Once the required coefficients have been computed, each calculation of a new configuration requires on the order of N^3 multiplications.

Details of this scale-parameter scheme are given in Ref. 106. In a series of calculations for argon, the largest error in the total energy was 0.6 eV out of 14 keV. The error is the difference between approximate and accurate self-consistent calculations; errors due to the neglect of such matters as relativity and correlations will be much larger. This scheme was developed for the nonrelativistic case only, and this is a significant limitation when high-Z atoms are of interest.

The second scheme again starts with a self-consistent field atomic calculation for a configuration specified by N occupation numbers n_i^0 . The total energy of a nearby configuration is expressed as a Taylor series:

$$L(n) = L(n^0) + \sum_i \epsilon_i n_i - n_i^0 \frac{\partial L}{\partial n_i^0} + \frac{1}{2} \sum_i \sum_j n_i - n_i^0, n_j - n_j^0 \frac{\partial^2 L}{\partial n_i^0 \partial n_j^0} + \dots \quad (175)$$

The partial derivatives of the total energy are easy to obtain, and we can use this series to evaluate the total energy for any configuration n that does not differ too much from n^0 . Even with just the few terms shown in Eq. (175), this series is remarkably accurate. Furthermore, it is just as easy to implement this scheme for the relativistic case as the nonrelativistic one, so it is particularly useful for high-Z atoms.

The first term in the Taylor series, $L(n^0)$, is a main result of the self-consistent field calculation. (Relativistic effects are included by use of the Dirac equation in the

self-consistent field calculation of the reference configuration n_i^0 .)

The first partial derivatives, $\partial E / \partial n_i^0$, are just the energy eigenvalues of the self-consistent field calculation:

$$\frac{\partial E}{\partial n_i^0} = \epsilon_i^0, \quad (k = 1, 2, \dots, N) \quad (176)$$

This well-known result is easily shown.

The second derivatives are numerically computed as the derivatives of the eigenvalues:

$$\frac{\partial^2 E}{\partial n_i^0 \partial n_l^0} = \frac{\partial \epsilon_i^0}{\partial n_l^0}, \quad (k, l = 1, 2, \dots, N) \quad (177)$$

To evaluate Eq. (177) we change one occupation number n_i^0 by a small amount to $n_i^0 + \Delta n_i$ and then carry out a new self-consistent field calculation, which enables us to calculate

$$\frac{\partial^2 E}{\partial n_i^0 \partial n_l^0} \approx \frac{\Delta \epsilon_i}{\Delta n_l} \quad (k, l = 1, 2, \dots, N) \quad (178)$$

This procedure is repeated for all the shells, ($l = 1, 2, \dots, N$), in the atom or ion. These additional self-consistent field calculations are much faster than the first one, because only a few cycles are needed to obtain convergence when Δn_i is small.

All the calculations needed to find the coefficients in the Taylor series for the relativistic self-consistent field evaluation of the total energy of a high-Z atom take less than 20 s on a fast computer. Once these coefficients are in hand, a single evaluation of the total energy requires only $N^2 + N$ multiplications. The resulting total energy is typically almost as good as the value obtained from a self-consistent calculation. In computing transition energies (the difference between two total energies), the relaxation energy is accounted for through the second-order term in the Taylor series.

The calculations for argon mentioned above and in Ref. 106 were repeated with the Taylor-series method. The largest error in the total energy was 2.1 eV, while this is greater than the error in the scale-parameter calculations; it is still more accurate than required for most plasma-physics needs.

These first tests of our new schemes for the rapid calculation of the properties of plasma ions encourage us to believe that a considerable improvement is possible over current screened-hydrogenic models. We expect that these methods will improve the calculation of the opacities and will prove useful in dealing with plasmas that are not in thermal equilibrium.

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J. R. Albritton

Electron-Ion Heat Exchange in Ionization and Recombination

LASNEX simulations of laser-produced plasmas usually predict unequal electron and ion temperatures T_e and T_i . We are working to improve our fundamental understanding of this interesting nonequilibrium thermodynamic state.¹⁰⁷⁻¹⁰⁹ Here we describe some recent advances in our knowledge of electron-ion heat exchange, which determines the ion temperature.¹¹⁰

In general, $T_e > T_i$ in the outer portions of an ICF target because laser energy absorbed by electrons is only slowly transferred to the ions. The situation is often reversed in the imploded fuel, where both hydrodynamic work and thermonuclear reactions heat the ions so that $T_i > T_e$; in this case, heat exchange with the electrons limits the ion temperature.

The ion temperature is important for three obvious practical reasons: first, the thermonuclear fusion reaction rate is very sensitive to T_i ; second, T_i controls a Coulomb logarithm that is important for laser absorption by inverse bremsstrahlung or electron thermal and electrical conduction (this is discussed in the article "Electron-Ion Collisions in Partially Ionized High-Z Plasmas," earlier in this section, and in Ref. 111); third, T_i determines the Doppler contribution to the width of spectral lines, and this is very important in connection with certain proposed designs for a soft x-ray laser.¹¹²

In this article we examine a novel mechanism of electron-ion heat exchange that occurs only in partially ionized plasmas and that is caused by sudden charge fluctuations resulting from ionization or recombination.

In the example examined below, this mechanism appears to set a lower limit on T_i .

Collisional Heating of Ions. The new effect operates in addition to the usual collisional coupling given by the Landau-Spitzer formula,

$$\frac{dE_i}{dt} = \frac{m_e}{M_i} (kT_e - kT_i) n_e \sigma_{eff} v_e, \quad (179)$$

where E_i is the energy per ion; m_e and M_i are the electron and ion masses; kT_e and kT_i are the electron and ion temperatures (in energy units); n_e is the electron density; $v_e = (kT_e/m_e)^{1/2}$ is the electron thermal velocity; and

$$\sigma_{eff} = 4\pi \frac{Qe^2}{kT_e} \log \Lambda \quad (180)$$

is an effective collision cross section (Qe is the ion charge, $-e$ is the electron charge, and $\log \Lambda$ is the Coulomb logarithm). The numerical coefficient in Eq. (180) is derived from a well-known formula¹¹³ for the electron-ion coupling time τ_{ei} . The physical process described by Eqs. (179) and (180) is energy transfer that occurs in electron-ion Coulomb collisions; a small fraction (approximately $m_e/M_i \ll 1$) of the kinetic energy of the hotter particle is transferred in each collision.

Ionization Heating of Ions. The new heat-exchange mechanism is also a consequence of electrical forces. Consider an ion of charge $+Qe$, surrounded by a neutralizing or screening cloud of electrons and ions. The electrical energy stored in this cloud is $\Delta E \approx Q^2 e^2 / R$, where R is the appropriate screening length. After electron-impact ionization, the screening cloud must adjust to the increased ion charge $Q' = Q + 1$. For times short compared with the ion plasma oscillation period $1/\omega_{pi}$, the extra unit charge is screened by electrons at a large distance $\sim D_e$, the electron Debye length. At the radius R_0 of the nearest-neighbor ions, this leaves an unscreened electric field that causes the neighbor ions to relax outward from the central ion. This outward motion is rapidly thermalized into ion kinetic energy. Thus in each ionization event, the ions

gain a thermal energy of order $d(\Delta E)/dQ \approx Qe^2/R$

This reasoning gives the rough estimate¹⁰⁹

$$\frac{dE_i}{dt} \approx \frac{Qe^2}{R} \phi n_e \sigma_{ii} v_i, \quad (181)$$

for the rate of ion heating associated with ionization processes. The rate of energy transfer is proportional to the impact ionization rate $n_e \sigma_{ii} v_i$; ϕ is a numerical factor discussed in Ref. 109. Essentially the same physical mechanism is suggested (from a rather different viewpoint) by Hagelstein,¹¹² who calls the phenomenon ion *proximity heating*.

It is important to observe that the energy transfer is not necessarily cancelled by a subsequent recombination. Indeed, if the time between ionization and recombination exceeds $1/\omega_{pe}$, and if T_e is low enough, the screening cloud around the ion of charge $Q + 1$ has time to relax to a low-energy configuration; when recombination occurs, the screening cloud must again readjust, and this transfers additional energy to the ion system.

Equations (179) and (181) indicate that energy transfer by the new mechanism can equal or exceed collisional energy transfer, although the ionization cross section is usually much smaller than the Coulomb cross section, the energy exchanged per ionization event ($\sim Qe^2/R$) can greatly exceed the small energy transfer in an electron-ion collision ($\sim m_e kT_e/M_i$).

The process described by Eq. (181) is very similar to a Raman effect that occurs in the absorption of visible light by impurities or defect centers in transparent solids.¹¹¹ In that case, absorption raises the defect center to an excited state that interacts with the environment with an altered potential; the defect and environment release energy (as phonons) in relaxing to a low-energy configuration in the changed potential.

After estimating the ion temperature rise produced by this mechanism in a rapidly ionizing plasma, we reconsider Eq. (181) to reconcile it with the principle of detailed balance.

Example: A Rapidly Ionizing Plasma.

Hagelstein's proposed design¹¹² for a soft x-ray laser based on resonant photoexcitation

illustrates the importance of the ionization heat-exchange mechanism (proximity heating). In this scheme, a low ion temperature ($T_i < 10$ eV) would have the favorable consequence of reducing the Doppler width of the pumping transition. In the proposed experiment, neon ions at a number density $\sim 10^{18} \text{ cm}^{-3}$ would be rapidly stripped to charge $Q = +8$ by photoelectric absorption of kilovolt x rays from an intense laser-plasma source. Three atomic scalelengths are important: the ion Debye length

$$D_i = \left(\frac{kT_i}{4\pi Q^2 e^2 n_i} \right)^{1/2} \approx 10 \text{ to } 20 \text{ \AA}, \quad (182a)$$

the ion neighbor separation

$$R_{ii} = \left(\frac{3}{4\pi n_i} \right)^{1/3} \approx 60 \text{ \AA}, \quad (182b)$$

and the electron Debye length

$$D_e = \left(\frac{kT_e}{4\pi Q^2 e^2 n_e} \right)^{1/2} \approx 200 \text{ to } 400 \text{ \AA}. \quad (182c)$$

The numbers refer to the proposed conditions $Q = +8$, $kT_i = 1$ to 5 eV, $T_e = 60$ to 200 eV. The situation is thus

$$D_i \ll R_{ii} \ll D_e. \quad (183)$$

The electrostatic interaction energy of adjacent ions is $\sim Q^2 e^2/R_{ii}$, much greater than kT_i ; this is the well-known condition of *strong coupling* of the ions, equivalent to $D_i \ll R_{ii}$. In this case, the ion of charge $+Qe$ is principally screened by ions at a range $\sim R_{ii}$.

In the ion-sphere model, the central ion is assumed to be surrounded by an electron gas that is nearly uniform for $0 \leq r \leq R_{ii}$; outside this range the other ions and electrons are taken to constitute a uniform, neutral plasma. This model predicts a screening energy $\Delta E = (9/10)Q^2 e^2/R_{ii}$, a value that agrees within a few percent with the results of elaborate Monte Carlo simulations of the strongly coupled plasma.

In this model each ion (of arbitrary charge Q') is surrounded by a cavity containing only electrons, whose radius is such

as to assure electrical neutrality. When we include the change of this cavity radius with ionization of the central ion we find

$$\frac{\partial \Delta E}{\partial Q'} = \frac{3}{2} \frac{Q' e^2}{R_0} \quad (184)$$

A minimum ion temperature can be calculated by integrating such energy changes with respect to Q' from $Q' = 0$ to the final charge state $Q' = Q$; assuming an ideal-gas specific heat, we find

$$kT_i \geq \frac{1}{2} \frac{Q^2 e^2}{R_0} \quad (185)$$

This mechanism alone thus predicts $kT_i \geq 7$ eV for the case considered. The actual ion temperature would be higher still because of the contribution of the collisional mechanism of Eq. (179) and because of cyclic ionization-recombination heating when steady ionization is achieved.

The limitation indicated in Eq. (185) applies to a plasma composed of a single substance (e.g., pure neon). If this material were diluted with a low- Z impurity (e.g., hydrogen gas) the same ion heating energy would be shared among more nuclei, resulting in a lower ion temperature.

The Problem of Detailed Balance. The formulation of electron-ion heat exchange based on Eq. (181) appears to predict a continuous heating of the ions, and thus would appear to predict that the ion temperature would continue to rise even if it exceeded T_e . This is obviously wrong; the heat-exchange mechanism must move the two species toward thermal equilibrium, and the net rate of heat transfer must approach zero as the temperature difference approaches zero.

Alternatively phrased, a proper view of the ionization-exchange mechanism must specifically include the possibility of transferring heat from (hotter) ions to electrons: near equilibrium, the two exchange rates must be equal, leading to detailed balance. This requirement is satisfied by the collisional mechanism described by Eq. (179).

To obtain detailed balancing for proximity heating, we examine ionization and

recombination that occur in the presence of a fluctuating electrostatic potential produced by plasma particles near a central ion.

Ionization and Recombination in a Potential. We consider an ion subject to an electrostatic potential Φ generated by neighboring charged particles. If the plasma density is low enough, the nearest-neighbor ion is often the main source of this potential, whose magnitude is therefore of order Qe/R_0 ; for definiteness, the reader may consider this case, although the reasoning is independent of the source of the potential.

We assume that Φ varies on a distance scale much larger than the radius of the ion core itself, and that it varies on the ion dynamical time scale (changing slowly compared with all electron transit times). An electron of energy ϵ_0 gains a kinetic energy $e\Phi$ as it approaches the central ion, and the impact-ionization cross section is therefore approximately $\sigma_{ii}(\epsilon_0 + e\Phi)$, where $\sigma_{ii}(\epsilon_0)$ was the ion's original impact-ionization cross section.

After impact ionization, there are two low-energy free electrons, one or both of which could be at low kinetic energies ($< e\Phi$), where they are actually loosely bound by the plasma potential fluctuation. However, electrons in this portion of phase space are extremely collisional and may therefore still be regarded as free.

The ionization rate is affected by a change in the distribution function for the incident free electron and by the probability of finding the target ion in its specified charge state. Assuming that the ion states are populated according to an equilibrium Boltzmann factor that contains the ion temperature, and that the electron distribution is a nondegenerate Maxwell distribution of the asymptotic energy ϵ_0 , the result is a modified ionization rate containing Boltzmann factors $\exp(e\Phi/kT_e) = (Qe\Phi/kT_e)$ and a cross section evaluated at the shifted energy $\epsilon_0 = \epsilon_0 + e\Phi$.

The three-body recombination that is the inverse of this ionization also has an altered rate, and the changes are somewhat different in the two-temperature case; the distribution functions contribute the Boltzmann factors

$$\exp \left[\frac{2e\Phi}{kT_e} - \frac{(Q+1)e\Phi}{kT_i} \right] \quad (186)$$

corresponding to two electrons incident on an ion of charge $Q \geq 1$. The three-body cross section is evaluated at modified energies determined by energy conservation. Again there is a question about the low-energy portion of phase space, where one or both of the incident electrons would be formally bound to the potential fluctuation; again it seems reasonable to assume that these electrons are in equilibrium with the adjacent positive-energy states.

If the ion densities are related by the (LTE) Saha equation, we find the ratio of ionization and recombination rates to be

$$\frac{\text{Ionization rate}}{\text{Recombination rate}}$$

$$= \exp\left(\frac{e\Phi}{kT_i} - \frac{e\Phi}{kT_e}\right). \quad (187)$$

Because this ratio is not unity, the Saha population ratios are altered when the ion temperature is not equal to the electron temperature. An effect of this character has been predicted.¹¹⁵

Using the corrections described we can readily determine the energy exchange rate. The result is

$$\frac{dE_i}{dt} = \left\langle e\Phi \left[1 - \exp\left(\frac{e\Phi}{kT_e} - \frac{e\Phi}{kT_i}\right) \right] \right\rangle \times \exp\left(\frac{e\Phi}{kT_e} - \frac{Qe\Phi}{kT_i}\right) n_e \sigma_{ii} v_e, \quad (188)$$

where the average is taken over all configurations of the local environment and over the distribution function of incident electrons. The recombination contribution has been related to the ionization rate as described above, using the LTE Saha detailed-balance condition. If we expand this rate assuming the potential $e\Phi$ to be small compared with kT_e and kT_i , we find

$$\frac{dE_i}{dt} = (e\Phi)^2 \times \left(\frac{1}{kT_i} - \frac{1}{kT_e} \right) (u_e \sigma_{ii} v_e). \quad (189)$$

This expression is the proper generalization of Eq. (181) for nearly equal electron and ion

temperatures: it predicts zero heat transfer in the equilibrium limit, and changes sign with the temperature difference.

In the other limit ($kT_i \ll kT_e$), Eq. (188) is qualitatively equivalent to Eq. (181), and so this more comprehensive approach resolves the difficulty associated with the requirement of detailed balance.

This entire discussion has referred to impact ionization and three-body recombination. An entirely similar coupling of electrons and ions (and photons) arises from the analogous processes of photoionization and radiative recombination.

Author: R. M. More

X-Ray Scattering by Bound Electrons and Plasma Fluctuations

In 1981 we developed an improved computational description of x-ray scattering from systems containing bound electrons (e.g., atoms or partially stripped ions). Our approach gives results that agree well with experimental data and extend existing models.

Among the motivations for this work is the fact that many laboratory experiments with hot plasmas involve the scattering of x rays as an essential diagnostic step. For example, the output of an x-ray laser can be measured by counting x rays scattered by a foil oriented at 45° to the x-ray flux (see Fig. 3-49).

Our special interest is in coherent scattering: for an atom containing Z electrons, the total soft x-ray scattering cross section is approximately $Z^2 \sigma_T$, where the Thomson cross

Fig. 3-49. If the source temperature greatly exceeds the foil temperature, x-ray scattering can dominate the observed x-ray signal (except for transmitted flux).

section $\sigma_T \approx 6.6 \times 10^{-25} \text{ cm}^2$. The proportionality to Z^2 (rather than to Z) arises from coherent scattering, and strongly enhances the scattering for high- Z atoms. Most computer codes used for simulating hot-plasma experiments neglect coherent scattering: this is valid only for fully ionized plasmas.

In certain cases, x-ray scattering is calculated from tabulated cross sections that include coherent scattering in a form that applies to cold (neutral) atoms; these cross sections greatly overestimate the scattering for hot plasmas. Our computational model describes both extreme cases.

Consider x-ray scattering from an isolated, neutral atom. In this case we can use a simple formula for the x-ray differential scattering cross section first derived by Heisenberg^{11b}:

$$\frac{d\sigma}{d\Omega} = S(q) \left(\frac{d\sigma}{d\Omega} \right)_T, \quad (190)$$

where $(d\sigma/d\Omega)_T$ is the Thomson differential cross section, proportional to $(1 + \cos^2 \theta)$, θ is the scattering angle, $q = (4\pi/\lambda) \sin(\theta/2)$ is the momentum transfer, λ is the x-ray wavelength, and $S(q)$ is a structure factor approximately given by

$$S(q) = S_{\text{coh}}(q) + S_{\text{incoh}}(q) \quad (191a)$$

$$S_{\text{coh}}(q) = F(q)^2 = \left| \sum_{i=1}^Z f_i(q) \right|^2, \quad (191b)$$

$$S_{\text{incoh}}(q) \sim Z \sum_{i=1}^Z f_i(q)^2. \quad (191c)$$

The $f_i(q)$ of Eqs. (191) are in turn determined by

$$f_i(q) = \int d^3r e^{i\mathbf{q} \cdot \mathbf{r}} \varphi_i(r)^2, \quad (192)$$

Here, $\varphi_i(r)$ denotes the i th normalized single-particle wave function entering into a Slater determinant used to describe the atomic electrons. In Eqs. (191) the sums over j run over all occupied orbitals. It can be shown that S_{coh} corresponds to elastic scattering; for a multi-electron atom, this is coherent scattering. The term S_{incoh} describes inelastic scattering, which is necessarily

incoherent. Equation (190) was derived under the assumption that the photon energy $h\nu$ greatly exceeds the k-shell binding energy E_k , but is less than the electron rest energy $mc^2 \approx 0.5 \text{ MeV}$.

Now consider the limiting behavior of $S(q)$. For low (but not too low) energies, $qr_a \ll 1$ for all θ , where r_a is an atomic radius. Using the normalization of the φ_i , Eq. (192) gives $S \propto S_{\text{coh}} \propto Z^2$, which implies a total cross section proportional to Z^2 , as noted earlier. For high energies, $qr_a \gg 1$ for sufficiently large θ ; in this case the $f_i(q)$ tend to zero, and we find $S \propto S_{\text{incoh}} \propto Z$. In this limit the atomic binding has no effect, and the scattering is entirely incoherent.

We have evaluated the x-ray scattering cross section over a large energy range for a variety of atoms with the following simple hydrogenic approximation for the wave functions:

$$\varphi_i(r) \propto e^{-r/r_n}, \quad (193)$$

where the orbital radius r_n depends only on the principal quantum number n of the state i . In our work r_n is determined by the screened-hydrogenic model^{11c}:

$$r_n = \frac{a_0 n^2}{Q_n}, \quad (194)$$

where Q_n is an effective charge for electrons in the n th shell. To treat high-energy photons, we replace the Thomson cross section in Eq. (190) by the relativistic Klein-Nishina differential cross section.

Figure 3-50 shows the resulting approximate model and more accurate tabular data^{11b}; the agreement is clearly adequate for practical purposes. The cross sections are very large for low-energy photons scattered by high- Z atoms. The present LASNEX plasma model is equivalent to the curve with $Z = 1$, so incorporating these results will make a large change in the calculated x-ray scattering.

In generalizing the model to treat scattering from a partially ionized plasma, we retain Eq. (190) but employ a more elaborate structure factor $S(q)$. For any many-electron system, the structure factor is given by

$$S(q) = \frac{1}{N} \langle \Psi | e^{i\mathbf{q} \cdot \mathbf{r}_1} \cdots e^{i\mathbf{q} \cdot \mathbf{r}_N} | \Psi \rangle, \quad (195)$$

where $|\Psi\rangle$ is the many-electron wave function and $\mathbf{r}_i$ is the position of the i th electron. In a partially ionized plasma, the structure factor includes contributions from both bound and free electrons.

The system of electrons and nuclei can be described by a set of adiabatic wave functions that incorporate the approximation that nuclear velocities are much smaller than electronic velocities. The result for $S(q)$ is

$$S(q) = S_n(q) + \sum_i F_i(q) S_{ni}(q) + \text{c.c.} \\ + \sum_i F_i(q)^2 S_{nn}(q) + n_i S_{nn}(q) \quad (196)$$

Here S_n , S_{ni} , and S_{nn} are the usual plasma structure factors describing correlations between pairs of free electrons, free electrons and nuclei, and pairs of nuclei. In Eq. (196), the quantity

$$F_i(q) = \sum_{j=1}^{N_{ij}} f_j(q) \quad (197)$$

is the ion form factor, as in Eqs. (191); in this case, however, the number of terms in the sum corresponds to the number of bound electrons N_{ij} . The final term $S_{nn}(q)$ is also a straightforward generalization of the atomic incoherent scattering factor of Eq. (191c). The coefficient n_i is the ion number density.

We have developed a computer code that evaluates x-ray scattering on the basis of the equations given above. The structure factors are obtained from the Debye-Hückel theory, modified to approximately describe the strong-coupling regime. The ionization state is calculated from the screened hydrogenic model. Figure 3-51 gives representative results for the total x-ray scattering cross section for a lead (Pb) plasma at 1 g/cm³ and 1, 10, 100, and 1000 eV. At low temperatures, where the majority of the electrons are bound, the cross section for soft photons is greatly increased by coherent scattering. At higher energies it does not matter whether the electrons are bound or free, since the x-ray wavelength is small compared to the interelectron spacing, and all curves coincide. Similar results are obtained for multicomponent plasmas.

Fig. 3-50. Photon scattering cross section vs energy (neutral atoms).

Fig. 3-51. Photon cross section (per unit atomic volume) vs photon energy.

However, the effect of partial ionization is much larger than the effect of ion-ion and ion-electron correlation in plasmas.

Authors: J. Oliva and R. M. More

Graphics-Oriented Application Language for LASNEX

GOAL, a graphics-oriented application language, was developed to help physicists understand the large amounts of data produced by LASNEX. GOAL combines many aspects of the old LASNEX language, computer graphics, and standard computer languages.

When LASNEX calculates a simulation, it creates "snapshot" and "history" data files. A typical simulation produces 100 files, each containing perhaps 300 000 floating-point numbers, or some 30 000 000

Fig. 3-52. Graph plots made with GOAL.

Fig. 3-53. LASNEX mesh plot made with GOAL.

floating-point numbers in all. The only means by which such a large quantity of data can be assimilated are graphical.

The Predecessor Language. The LASNEX graphics language grew by accretion from 1970 to 1978 as design physicists, code physicists, and computer scientists explored the capabilities of the CDC 7600s and graphics peripherals. In 1978, the language was redesigned to remove inconsistencies and to add more expression and data-selection capability. The language remained batch-oriented. The language was moved to the Cray-1s in 1979, and has changed little since then.

The old LASNEX graphics language had the following difficulties:

1. It did not permit users to create their own variables.
2. It allowed only limited algebraic expressions.
3. It allowed only limited data selection (a combination of subscript ranges and subscript functions).
4. It did not provide control structures, such as looping and conditional execution.

5. It had no inline macro facility (a macro is a way of naming text so that the text replaces the name whenever the name is used).
6. It was not "user friendly." For example, many error messages were vague or nonexistent.
7. It was batch-oriented and was clumsy when used interactively.

All these difficulties have been solved by the GOAL language. GOAL is available both as an interactive post-processor, which operates on LASNEX dump files, and as an integral part of LASNEX. In the latter form, it generates plots and evaluates expressions as a result of commands from the problem-generation file or the user's terminal.

Computer Graphics Concepts. GOAL has six ways to convert data into pictures. First and most widely used is the XY graph plot. X and Y are arrays of numbers $(X_1, X_2, \dots, X_N)$ and $(Y_1, Y_2, \dots, Y_N)$; a line is drawn from point (X_1, Y_1) to (X_2, Y_2) to ... to (X_N, Y_N) . Figure 3-52 is an XY graph plot that shows pressure as a function of the Z coordinate for three values of the radial zone index.

Another picture that requires X and Y arrays is a scatter plot, created by drawing a letter or dot at the points (X_1, Y_1) , (X_2, Y_2) , ..., (X_N, Y_N) .

GOAL lets the user create a picture of the LASNEX RZ Lagrangian mesh of quadrilaterals. Three arrays are needed: an array of radial coordinates $R(KMAX, LMAX)$, an array of axial coordinates $Z(KMAX, LMAX)$, and an array $XYMAP(KMAX, LMAX, 2)$ that tells which lines should be drawn. Note that there are twice as many lines as points. Each mesh quadrilateral (K, L) is associated with four vertex points: $(R(K, L), Z(K, L))$, $(R(K-1, L), Z(K-1, L))$, $(R(K-1, L-1), Z(K-1, L-1))$, and $(R(K, L-1), Z(K, L-1))$. Figure 3-53 shows a LASNEX mesh plotted with GOAL.

GOAL lets the user create a picture of the material boundaries of the mesh. Four arrays are needed: R, Z, and XYMAP to draw the mesh, and a fourth array of material region numbers $MREG(KMAX, LMAX)$. To each quadrilateral (K, L) there corresponds a material number $MREG(K, L)$.

GOAL lets the user create a picture of a field of velocity vectors. Four arrays are needed: R and Z for the location of the tail of each arrow, an array of radial displacement $U(KMAX, LMAX)$, and an array of

axial displacements $V(KMAX, LMAX)$. The velocity vectors are drawn from the point $(R(K,L), Z(K,L))$ to the point $(R(K,L) + DELTAT*U(K,L), Z(K,L) + DELTAT*V(K,L))$. Figure 3-54 shows velocity vectors plotted with GOAL.

GOAL lets the user create pictures of contour lines. Three arrays are needed: R and Z for the location of the mesh, and the elevation variable $F(KMAX, LMAX)$. To each point $(R(K,L), Z(K,L))$, there corresponds an elevation $F(K,L)$. The contour lines are lines of constant elevation. Figure 3-55 shows contour lines above a dotted mesh.

Computer Language Concepts. Physicists and computer scientists have used an extended Fortran, called LRLTRAN, during all of LASNEX development. LRLTRAN is rich in concepts: its features include data declaration, data assignment, mathematical operations, data selection, control structures, and macros. These concepts, combined with the old LASNEX graphics language, formed the nucleus for the creation of GOAL.

GOAL lets the user declare new variables. For example,

```
DECLARE REAL PRESSURE (KMAX,LMAX)
```

declares a floating-point mesh variable called PRESSURE. Other data types include INTEGER to store integers, LOGICAL to store true/false (0/1) logical values, SYMBOL to store variable names, and STRING to store ASCII strings.

Values can be assigned to variables. For example,

```
PRESSURE = 1.
```

assigns the floating-point value 1 to all of PRESSURE.

Mathematics can be performed on data. For example,

```
DENSITY = MASS/VOLUME
```

divides the MASS array by the VOLUME array and assigns the result to the DENSITY array.

Both scalar and array operations are available, including all the well-known mathematical operations: the usual operations (+, -, *, /, etc.), the Boolean operations (.I.E., .EQ., .AND., etc.), and the functions (SQRT, SIN, AMIN, etc.).

Fig. 3-54. Velocity-vector plot made with GOAL.

Fig. 3-55. Contour plot and dotted-mesh plot made with GOAL.

GOAL lets the user select data. For example,

```
MASSAVG(7) = MASS(7,AVG(2:LMAX))
```

selects the average MASS along K-line 7 for L from 2 to LMAX and assigns it to the seventh element of MASSAVG. Other subscript functions include minimum value (MIN), maximum index (MXX), peak-to-peak change in value (PTP), and partial sums (PSUM).

GOAL lets the user loop over commands and selectively execute commands. For example,

```
DO I = 1 : KMAX*LMAX
  IF (VOLUME(I) .NE. 0.0) THEN
    DENSITY(I) = MASS(I)/VOLUME(I)
  ELSE
    DENSITY(I) = 0.0
  ENDIF
ENDDO
```

GOAL lets the user define macros and macro arguments. For example, the following

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macro will plot a graph of any mesh-size quantity against the averaged position:

```
MACRO DEFINE PLGVSZT (K,A) =  
PLG A(K,2:LMAX), &  
    (ZT(K,1:LMAX - 1) + &  
    ZT(K,2:LMAX) + &  
    ZT(K - 1,1:LMAX - 1) + &  
    ZT(K - 1,2:LMAX))/4.  
MACRO END
```

To plot K-line 7 of ERAD against the average position using this macro type
PLGVSZT(7,ERAD).

An important feature for laser fusion designers is GOAL's ability to work with LASNEX data, experimental data, and analytical data.

GOAL is consistent. It was designed to avoid many of the inconsistencies found in the old LASNEX language.

GOAL is terse. Almost every example can be shortened by synonyms that are part of the language. For example, the command PLOT GRAPH can be shortened to PLG.

GOAL is "user friendly." In contrast to the old LASNEX language, GOAL reports all detected errors, issues useful error messages, and recovers without crashing. This makes errors easier to find and correct.

The Future. We plan to add three-dimensional pictures, cursor and movie commands, and additional macro capability to translate current languages into GOAL.

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Targets, such as those shown here, filled with a deuterium-tritium fuel mixture, are positioned in the center of the target chamber and illuminated by the laser. Parameters such as size, shape, and material of the target, the amount and the state of the fuel, the efficiency of laser-light absorption by the target, and the fuel's reaction to the laser energy are all important in producing uniform compression and ignition. Therefore, techniques have been developed to fabricate very small, uniform, hollow spheres from a variety of materials, and to fill them with DT fuel and other tracer gases. In the large photo, polymer spheres are magnified 134 times by a cross-polarized microscope. In the photo below, hundreds of the same polymer spheres floating in a water/soap solution illustrate a new chemical processing method, microencapsulation, being developed to fabricate these spheres. The inset photo at left is a low-temperature experimental chamber used for determining material properties at cryogenic temperatures. The inset photo at bottom left shows a resonant photoexcitation x-ray laser target whose components are made by anisotropic etching of silicon. This technique allows specific planes of the crystal lattice to be etched, leaving other planes unetched, thus producing precise geometric shapes.

Target Fabrication

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Target Fabrication

Introduction

I. F. Stowers

The Fusion Target Fabrication (FTF) Program researches, develops, and fabricates targets for the Inertial Confinement Fusion (ICF) Program. This year we developed targets for soft x-ray laser experiments and have begun fundamental studies into the formation of nuclear spin-polarized DT fuel. These and other major advances are highlighted below.

Novette. Using Novette's green output, LLNL performed the world's first laboratory demonstration of soft x-ray lasing action. For this demonstration, we developed the techniques to fabricate two new target designs: the resonant photoexcitation scheme and the collisional excitation scheme. We fabricated 23 targets of the first design and 145 of the second.

To aid in the study of laser plasma interactions at $0.53\text{ }\mu\text{m}$ and $0.26\text{ }\mu\text{m}$, we constructed a number of thin CH foil targets that contain a central spot doped with silicon or sulfur. Called multilayer polymer microspot targets, they are an improvement over the microspot targets reported on in the 1983 Laser Annual.¹

Nova. Because Nova will be operational as an experimental facility in 1985, we have begun to develop new concepts for advanced, high-gain targets. One of these concepts is nuclear spin-polarized DT fuel, a technique with high potential for increasing target gain. This process calls for the formation of pure molecular DT, and a determination of its spin-lattice relaxation time.

A series of experiments to demonstrate target concepts that are hydrodynamically equivalent to future high-gain targets is planned for the Nova laser. To prepare for this campaign we are designing and have begun constructing a cryogenic target support system, which has numerous advantages over the Novette system. And because some Nova targets may require condensed layers of fuel or diagnostic gases, work toward producing these layers uniformly has been continuing at the University of Illinois.

Finally, we have developed a method of forming ultrathin films of polystyrene, poly(butadiene-styrene), and polypropylene. Among these, polypropylene was found to have many potential uses in the construction of foil targets.

Novette Targets

Novette X-Ray Laser Target Fabrication

For the first time, the ICF Program at LLNL has produced amplified, coherent radiation in the soft x-ray region. X-ray lasers yielding this radiation used Novette's green output as their primary source of energy to indirectly pump a gain medium. The FTF Program developed the techniques to construct these x-ray laser targets. Related theory and experiments are discussed in "X-Ray Laser Experiments," Section 5.

We have investigated two principal x-ray laser schemes: the resonant photoexcitation design and the collisional excitation design. Like other lasers, x-ray lasers require a gain medium where energy can be stored in excited atoms or molecules. However, x-ray lasers, unlike visible lasers, require the higher energy transitions of atoms' inner-shell electrons for their short wavelength output. As a result, their gain medium is created using the highly ionized atoms of a lasing material. Once ionized, the medium is pumped to create the population inversion of excited states necessary for lasing.

The primary differences between the two schemes lie in their lasing materials, the structures used to support those materials,

and the methods of pumping the gain medium. While the resonant photoexcitation version has been the most difficult to design and fabricate, the collisional excitation version has demonstrated the amplification criteria for a laser.

Resonant Photoexcitation Scheme. Central to this design is the laser gain channel, a high-aspect-ratio, pressurized gas channel containing the lasing fluorine, in the form SF_6 . Figure 4-1 shows the laser channel, which is enclosed on two sides by input windows transmissive to pump x rays and on two sides by knife-edge shaped walls. Both ends of the channel are enclosed by x-ray transmissive output windows. Separated by $50\ \mu m$ from the input windows,

and parallel to them, are thin manganese foils used as the pump flashlamps. The lasing gas is supplied to the target while it is mounted in the target chamber. Prior to a shot, the laser gas pressure is monitored at these locations along the laser channel.

The population inversion is created by the resonant excitation of ionized lasing atoms.² As illustrated in Fig. 4-1, a high-power visible laser is focused onto a metal foil and a gaseous medium is pumped by the x-ray emission from this foil (the flashlamp). Broad-band x rays from the flashlamp ionize the cold lasing gas to a hydrogen-like state, while a specific x-ray line in the flashlamp emission excites the ionized fluorine atoms. From the excited

Fig. 4-1. Cross section of the resonant photoexcitation x-ray laser target. In the center of the target is the laser gain channel containing the lasing medium.

state, the atoms collisionally transfer to the upper laser level.

The cross section of the laser channel, shown in Fig. 4-1, requires a knife-edge geometry with a tip radius of curvature less than $1\text{ }\mu\text{m}$ to prevent possible trapping. However, theoretical uncertainties in the x-ray laser gain calculations require that the channel dimensions be adjustable. For example, the opposing knife-edge tips are parallel to $\pm 2\text{ }\mu\text{m}$, but the distance between them is variable, from 30 to $100\text{ }\mu\text{m}$ on several different channels. In addition, the channel width has to be adjustable, from 40 to $100\text{ }\mu\text{m}$, with a uniformity of $\pm 2\text{ }\mu\text{m}$ for any individual design. The channel dimensions have to remain within tolerance over their entire 14-mm length.

Side and end windows are necessary to contain the lasant gas within the channel at a pressure of 50 Torr. This requires low diffusivity and low pinhole density in the windows. The side (input) windows must be transmissive to both the hard, ionization x rays (a few keV), and the soft, photoexcitation x rays (a few hundred eV). The end (output) windows must transmit the soft x ray laser radiation ($\approx 153\text{ eV}$). In addition, calculations have shown that a $\pm 10\%$ pressure nonuniformity in the gas could significantly reduce the x-ray lasing action. To insure pressure uniformity, we measure the gas pressure along the laser channel. Under the worst conditions (all the gas leaking from one defect), a leak rate of 50 laser vol/s would produce a $\pm 10\%$

nonuniformity in the channel (one laser vol $\approx 5 \times 10^{-5}\text{ cm}^3$). Therefore, the leak rate of the structure has to be certified less than this value.

Outside of the laser channel, which is the central piece of the laser design, are flashlamps spaced approximately $50\text{ }\mu\text{m}$ from the input windows (Fig. 4-1). The flashlamp windows are 150-nm -thick silicon nitride (like the input windows) and are coated with 75 to 125 nm of manganese. A vacuum must exist between the input windows and the flashlamps to minimize pump radiation absorption losses. Finally we had to devise a way to fill the laser channel in situ with the fluorine lasant gas and verify the gas pressure during operation.

We anisotropically etched $\langle 100 \rangle$ crystal-line silicon wafers to produce all parts of the resonant photoexcitation laser, a useful technique that allows specific planes of the crystal lattice to be etched, leaving other planes unetched. Anisotropic etching micromachines structures with smooth and crystallographically defined surfaces. (The microfabrication process is illustrated in Fig. 4-2.)

To fabricate the flashlamp, we first thinned silicon wafers to $200\text{ }\mu\text{m}$ thick, and then coated them with 150 nm of silicon-rich silicon nitride. This form of silicon nitride has less stress and deformation than stoichiometric Si_3N_4 (Ref. 3). Our silicon nitride was prepared by low-pressure, chemical-vapor deposition with a 1:1 gas mixture of $\text{SiH}_2\text{Cl}_2:\text{NH}_3$. Auger analysis

- 1) Oxide/nitride deposition on a thinned wafer

- 2) Photolithography

- 3) Anisotropic etching

Fig. 4-2. The microfabrication process used to produce all parts of the resonant photoexcitation x-ray laser.

indicated that the silicon nitride film was composed of approximately equal atomic amounts of silicon and nitrogen. The coated wafer was then patterned using photolithography and the silicon was anisotropically etched in 77°C KOH, which does not attack the nitride. This yielded a 14-mm-long, 10-mm-wide silicon window frame with two 6-mm-long, 400- μ m-wide silicon-nitride windows separated by a 0.5-mm-wide silicon frame. In the center of the frame was a $250 \times 250 \mu\text{m}$ -square hole produced by double-sided anisotropic etching. This pump-out hole, shown in Fig. 4-3(a), was used to evacuate the region between the flashlamp and input windows. A 0.8-mm-long dead spot in the Novette line foci enabled us to split the flashlamp into two halves. The windows were sputter-coated with 50 to 150 nm of manganese, which formed the metal flashlamp foil. Calculations indicated that the 150-nm-thick, silicon-nitride layer would absorb only a small amount of the flashlamp pump x rays.

We fabricated the input windows similarly to the flashlamp wafers, except that we first thinned the wafers to 50 μm , then etched them to yield a 13-mm-long $\times$ 300- μm -wide, 150-nm-thick silicon-nitride window on each 14-mm-long, 10-mm-wide silicon wafer. The fragility of these thin wafers necessitated the use of special vacuum chucks and handling procedures. Assembly (or "glue" and "vacuum") holes [Fig. 4-3(b)], used to aid in the final assembly of the laser parts [Fig. 4-3(b)], were simultaneously etched into these wafers. Double-sided wafer patterning and etching of all parts were required to do this efficiently. We used a specially designed fixture

that sandwiches a photoresist-coated silicon wafer between two masks, allows the masks to be aligned to each other and to the flat of the silicon wafer (necessary for anisotropic etching), and then exposes both sides of the wafer.

The silicon, knife-edge laser channel [Fig. 4-3(c)], was fabricated by etching both sides of an appropriately thinned silicon wafer 40 to 100 μm thick, as shown in Fig. 4-2. By effectively designing the mask, we were able to control the interpoint spacing between the knife-edge tips to $\pm 2 \mu\text{m}$, with the spacing on an individual channel varying by less than $\pm 1 \mu\text{m}$. By carefully monitoring the etch, the radii of curvature of the knife-edge can be less than 50 nm. In addition, gas feed channels, pressure transducer lines, cleave lines and assembly holes were simultaneously etched into the channel wafer.

The gas feed channels were a set of 24 grooves, etched to a depth of half the wafer thickness and spaced 480 μm apart on alternating sides of the silicon wafer, as in Fig. 4-3(c). On the opposite side of the knife-edge from the feed lines, we etched three sets of pressure transducer lines. The three sets allowed us to tap the gas pressure at three distinct locations along the channel axis and measure the pressure with three externally mounted pressure transducers. Each set of transducer lines was made up of 10 grooves, spaced 100 μm apart on alternating sides of the wafer, and etched to a depth of half the wafer thickness. The wafer itself was 18 mm long and 10 mm high, with the knife-edge laser channel only 15 mm long. Two cleave lines symmetrically located and spaced 14 mm apart were etched approximately two-thirds of the way through

Fig. 4-3. Silicon wafer components for resonant photoexcitation x-ray laser: (a) flashlamp wafer, (b) input window, and (c) knife-edge laser channel wafer.

the wafer. After final assembly, we removed the excess length of the knife-edge wafer by snapping the wafer along the cleave lines.

This opened the knife-edge geometry, exposing the ends of the laser channel.

Finally, we fabricated the output window frames by anisotropically etching 100- μm -thick silicon wafers to produce a 1-mm-square structure, with a 150- μm -square opening in the center. Because of photon absorption at 153 eV, silicon nitride was an inappropriate choice for the output window material. Instead, a 100-nm-thick layer of Formvar was bonded to the silicon-end window frame.

After ETF designed each component wafer, the Electronic Research Division completed the photolithography-mask generation, wafer patterning, and anisotropic etching. The component wafers were then delivered to us for assembly into the target. We designed the component parts for the x-ray laser to make assembly an easy, modular procedure. Because particulate contamination of the wafer mating surfaces can cause wafer separation that results in failure of the target to perform as designed, we performed the assembly procedure in a clean room. A $\pm 2\text{-}\mu\text{m}$ flat assembly fixture with vacuum handling capability was machined for the procedure. We stacked two flashlamp wafers, two input window wafers, and one knife-edge wafer onto this fixture in appropriate order. Once positioned, we clamped the silicon together and, using 50- μm -diam pipettes, metered epoxy into the exposed assembly holes to bond the structure together. Once the epoxy cured, we turned the silicon "five-pack" of wafers over, and applied epoxy to the remaining assembly holes. At this point, we removed the excess length of the knife-edge wafer by cleaving, and positioned and epoxied the Formvar output windows into place. To insure leak-tightness, we also sealed any exposed edges of the laser with epoxy.

We then mounted and epoxied the five-pack into a metal frame, as shown in Fig. 4-4. The bottom of the aluminum structure held the five-pack in place and supported the pressure transducers, the laser gas supply, and a gas plenum. The top part of the structure was designed to transfer gas into the laser channel through gas feed lines etched into the knife-edge wafer.

After we assembled the five-pack and metal frame and completed an initial

leak-rate certification, we sputter-coated the manganese onto the flashlamp support windows. The flashlamp windows were not coated until this time because the opacity of the coating would interfere with the visual alignment of the laser channel and windows during assembly. We used sputter coating for this task because of its short turnaround time, easy thickness control, and low process temperature. Because considerable time and effort had been invested in the x-ray laser target prior to coating the flashlamp, we took stringent steps to minimize errors and accidents during the sputter process. The narrow gas-feed and pressure-transducer lines etched into the laser channel wafer provided a low conductivity gas path between the laser gain channel volume and the outside of the target. As a result, the normal, rapid pump-down and venting of the vacuum system caused a pressure differential as great as one atmosphere across the input windows. Additional control valves and a pressure monitor were installed into the vacuum system to manually control the rate of pressure change during pump-down and venting and reduce the risk of window breakage from excessive differential pressures. The rate of pressure change during the entire sputtering operation was kept below 10 Torr/s. Fifteen targets were successfully coated without a mishap.

To verify that the x-ray laser was assembled as designed, we performed several different tests. We used a Zygo large-area interferometer to verify the flatness of the

Fig. 4-4. The resonant photoexcitation x-ray laser in a metal support frame. The bottom metal piece holds the wafers in place, and supports the pressure transducers. The gas supply plenum is in the top piece.

x-ray laser parts. In the area of interest (14-mm-length $\times$ 1-mm-width along the channel) the 200- μ m-thick silicon flashlamp wafers were flat to $\pm 0.5 \mu$ m, and the 50- μ m-thick input window wafers and the x-ray laser knife-edge channel were flat to $\pm 1.5 \mu$ m. (Typically, the 10 $\times$ 14-mm-silicon parts were flat to $\pm 3 \mu$ m, but were much flatter in the smaller area of interest.) After assembly, we reexamined the silicon five-pack, and again interferometrically verified the channel to be flat within $\pm 2 \mu$ m along its length.

The silicon-nitride input windows were also required to bow less than 3 μ m under 50 Torr of pressure. When pressurized, the gas in the laser channel pushed the nitride sealing windows out. The extra volume created was outside the knife-edge. Lasing action in this region could not escape the output windows, and reduced the observable lasing action in the center. To measure this bowing, we evacuated one side of a 150-nm-thick, 13-mm-long, and 500- μ m-wide silicon-nitride film. Under the one-atmosphere pressure differential, the film bowed less than 2 μ m. In addition, this pressure differential was cycled several times to insure the integrity of the window.

After assembly, some of the five-packs were radiographed with 25-keV x rays. We analyzed the developed film with a microdensitometer to measure the accuracy of assembly and consistency of the interpoint spacing of the knife-edge tips. The resolution of the microdensitometer was approximately $\pm 1 \mu$ m, and no deviation in the interpoint spacing could be observed. In addition, the lateral placement accuracy of one wafer on top of another was typically $\pm 5 \mu$ m. All assembly was done using micromanipulators, viewing the samples under a stereomicroscope.

The final test for the x-ray laser assembly was to guarantee tightness against leaks and pressure uniformity. After assembly, we pressurized the silicon x-ray laser gain channel to 100 Torr above atmospheric pressure with helium gas. The helium leak rate was 10 to 15 x-ray laser vol/s. The leak rate for the larger lasant gas molecules would be less than this for either diffusion through the membranes or flow through gaps. The uniformity of the gas pressure

at 100 Torr inside the laser channel was measured with the transducers to be better than $\pm 10\%$.

The lasant gas was supplied to the resonant photoexcitation x-ray laser in the Novette target chamber through a gas-filling and monitoring system. Gas with greater than 0.995 purity was fed to the supply panel from a lecture bottle with the gas-supply network and the laser channel evacuated. Gas purity into and out of the channel was verified prior to shooting by means of a mass spectrometer. Readiness was determined when the channel pressures indicated an acceptable gas density profile and the spectrographs confirmed gas purity. Channel-filling transients were minimized by supplying gas to both the inlet and outlet of the laser assembly after purity confirmation. Gas purity for all shots was maintained above 0.995 and the pressure profiles were generally within the desired $\pm 10\%$ of the nominal 40-Torr value.

Collisional Excitation, Exploding Foil Scheme. Most of the x-ray lasers shot during this campaign were exploding foil targets using collisional excitation for the laser pumping. The lasant candidate is a 75-nm to 150-nm-thick foil of either selenium or yttrium coated onto a 150-nm-thick substrate of Formvar. The foil is supported on an aluminum frame that allows it to be illuminated along both sides of its length by the two green Novette beams.

Under the Novette radiation, the foil explodes while the lasant ionizes, creating a plasma with a high population density of neon-like-state ions. The ions collisionally excite to the upper laser level, creating the population inversion. The foil explosion forms a low-density column of the lasant. In the interior region of the column, the density gradients in the direction normal to the column axis are low. This minimizes refraction of the x-ray beam as it passes through the excited gain medium.

The selenium was coated by evaporation. We controlled the coating thickness by regulating the coating rate and time. The temperature of the selenium evaporation source had to be controlled to prevent radiative heating of the Formvar substrate. With the high vapor pressure of selenium, a hot coating and substrate would result in

re-evaporation of the selenium during coating and an uncertainty in its thickness. Because of the toxicity of selenium, we used a dedicated evaporation system.

The yttrium coatings were sputtered onto the Formvar substrates. The rapid oxidation kinetics of yttrium created a problem when entering the vacuum system. Normal venting rates allowed exposure of the yttrium to concentrations of oxygen high enough to cause the exothermic heat generated to ignite and burn the film. We implemented a staged venting procedure to force a reduction in the oxidation rate. This allowed the heat to dissipate, keeping the film temperature low. Once the yttrium was completely oxidized, the foils were exposed to air with no negative repercussions. The presence of oxygen in the lasant material did not degrade the targets.

Prior to each coating cycle during a trial run, we checked the film thickness. We also monitored the thickness in situ, and later confirmed this by a profilometer measurement on a witness slide coated during the same run. To experimentally verify that there was amplification along the plasma channel, we varied the length of the

illuminated foil. We could deduce gain from output intensity that was exponential with channel length. For this purpose, target lengths from 2.5 to 22 mm were needed.

The aluminum support frame was machined and lapped to have one flat side. The Formvar foil was floated from distilled water onto a wire hoop and then glued onto the flat side of the aluminum frame. We then coated the selenium or yttrium onto one side of the Formvar foil. We cut the ends of the lasant foil to the desired length for each particular shot using a 5- μm -diam glass fiber. A microscope with a micrometer translation stage having 2 μm resolution was used to measure the length of each target foil.

To keep background emission originating outside the laser gain channel from the primary diagnostics, 250- μm -apodizing pinholes were mounted on the ends of the frame. We fabricated the pinholes out of silicon by anisotropic etching as we did the component wafers used for the resonant photoexcitation x-ray laser targets. The holes, 250- μm square, were etched into 2-mm-square wafers. A line through the foil and the centers of the pinholes defined the

Fig. 4-5. The collisional excitation, exploding foil x-ray laser. The plasma gain channel is formed in the selenium or yttrium foils along a line between the output pinholes.

x-ray laser output beamline and the line that was to be illuminated by the Novette line foci.

After the pinholes were mounted, the frame was mounted to a target base and kinematic mount, the alignment mirror was mounted and aligned, the foil length was verified, and the target, as shown in Fig. 4-5, was delivered to Novette.

Because of the coherent output of the x-ray laser target, it was necessary to align the laser gain channel with two primary diagnostics located at opposing ends of the target chamber. The center of rotation of a goniometer was at the middle of the plasma channel. An alignment mirror mounted on one of the output ends of the x-ray laser was used to retroreflect from and insure the plasma gain channel was aligned to within 500 microradians of the primary diagnostics. The mirror was mounted to the target with a kinematic mount and, using a two-arm interferometer, was adjusted to be orthogonal to the expected output beam prior to delivery to the target chamber.

Summary. During the campaign, we fabricated 23 resonant photoexcitation and 145 collisional excitation x-ray lasers. Results for the resonant photoexcitation design are inconclusive; however, by developing this laser design and its silicon components we have demonstrated the utility of silicon and other microfabrication technology for future x-ray lasers and ICF targets. By developing the collisional excitation design we have demonstrated the criteria for amplification.

Fig. 4-6. Each multilayer microspot target was fabricated by coating a single-layer microspot target front and back with an additional layer of CH polymer.

Experimental results are described in detail in Section 5.

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Multilayer Polymer Microspot Targets

Last year we reported on the development of a seeded microspot x-ray diagnostic target.¹ This target consisted of a 300- μm -diam, 2- μm -thick disk of silicon or sulfur-seeded hydrocarbon polymer nested tightly in a hole in a 2- μm -thick film of pure hydrocarbon polymer. This year we extended our work on the microspot target, fully encapsulating the microspot in what we call the multilayer polymer microspot target.

The seeded polymer spot in the target is covered front and back with an additional CH polymer layer that further reduces temperature and density gradients for the plasma experiments conducted with this target. Several targets of this type were shot on Novette to measure the temperature and density of a laser-generated plasma and to study long-range plasma instabilities. "Exploding Foil Characterization," in Section 5, reports on experimental details and results.

The steps in fabricating a multilayer polymer microspot target are shown in Fig. 4-6. First a microspot target is prepared as described in the 1983 Laser Annual,² except that the film is left on the salt-coated substrate. Then, with all masks removed, the patterned polymer structure is overcoated with a layer of plasma polymer 2 μm thick.

Next we separate the desired area of the target from the excess film using a scalpel. A polyethylene support ring with a bead of adhesive on its face is then pressed against a defined area of the target. After the adhesive hardens the substrate is slowly immersed in water. The potassium chloride layer on the glass substrate dissolves, releasing the microspot and the support ring to float on the water surface.

The freestanding, two-layer film on the polyethylene support ring is then inverted

and coated with a second layer of plasma polymer. Thus, the microspot in the completed multilayer target is enclosed in CH polymer, both radially and longitudinally. Encapsulating the target in this way eliminates the undesirable temperature and density end effects that occurred in experiments using the single-layer microspot target.

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Nova Targets

Nova Cryogenic Target Support System

This year we completed the conceptual design of Nova's cryogenic target support system. The design is an evolutionary step beyond that of the Novette cryostat, described in the 1983 Laser Annual.¹ While the Novette cryostat was specifically designed to create a uniform layer of condensed DT in a glass microsphere, the Nova cryogenic system will be capable of supporting a variety of targets: those containing high density gas, spin-polarized fuel, condensed fuel, and diagnostic gases.

The Nova cryogenic system, illustrated in Fig. 4-7, provides a standard fixture for filling fuel capsules, in situ, to 200 atm; regulating target temperatures in the range of 20 to 300 K; and allowing the millimeter-size capsule to be consistently positioned in view of the target chamber alignment optics, despite the 5 m it must travel during insertion. Details of the target fixture are illustrated in Figs. 4-8 and 4-9.

One of our primary purposes for developing Nova's cryogenic system is to solve a problem that occurs when filling fuel capsules. IC² capsules are made of a permeable material so that fuel (usually D₂ or DT) can be diffused into their hollow interiors. Gas is diffused into the capsules to avoid the use of fill holes that might cause nonspherical implosions. Capsules made of materials that are highly permeable at room temperature can be filled in a matter of minutes. However, since most of the capsule materials are

essentially impermeable below 120 K, they can be filled at room temperature and then held in the cold state in an evacuated target chamber during alignment.

A cryogenic system also allows the capsules to be filled with higher density fuels without rupturing. This is significant since the desired fuel density will increase in proportion to the rise in laser power that we expect from Nova. Gas pressure in the capsule is greatly reduced by lowering the temperature to near the saturation point, between 20 and 40 K. This decrease in temperature also strengthens the capsule material. Capsules for spin-polarized and condensed fuel must be cooled to even lower temperatures, as discussed in "Spin-Polarized DT," later in this section.

Targets are mounted on a cryogenic utility post (CUP) and then attached to the Nova cryostat (Fig. 4-8). Most targets are assembled with radiation shields that protect certain diagnostics. These shields require that the target be positioned with an exact orientation in the target chamber. Correct orientation is guaranteed by the CUP, which is designed so that the target may be easily and securely attached to the subassembly under the cryostat end register (SAUCER), as shown in Fig. 4-8. Quasi-kinematic mounts orient the target and form a thermal switch between the SAUCER, which is a replaceable part of the cryostat, and the CUP. A thermal switch is activated by flowing helium into the annular space between concentric copper tubes in the CUP and SAUCER. Cold helium flowing through this space is blown onto the target to promote greater cooling capacity at a given temperature.

We diffuse gas into the target by gradually increasing the pressure in a high-pressure, stainless steel shroud, which is moved over the target by a two-axis servo. A large pneumatic cylinder generates sealing forces in excess of 50 kN between the shroud and SAUCER. These forces are maintained, despite thermal contraction (the high-pressure shroud is cooled by liquid nitrogen and liquid helium), by regulating pressure in the pneumatic cylinder. Following fill and cool-down, we evacuate the high-pressure shroud. The shroud and associated clamping equipment are then moved away from the cryostat.

Fig. 4-7. Nova's cryogenic system provides a means of filling a target above the target chamber, inserting the target, and holding it at cryogenic temperatures for up to 24 hours.

A gate valve below the cryogenic assembly chamber, shown in Fig. 4-7, can be opened to allow the cryostat to descend into the target chamber. We lower the cryostat 5 m to a motor and gear box using cables (attached to a winch with a vacuum feedthrough). Boil-off gas from the liquid helium Dewar of the cryostat is heated and exhausted through polymer tubes that follow the cryostat into the target chamber. The cryostat is then kinematically mounted in the target positioner to ensure that the target is correctly positioned and oriented.

The Nova cryogenic system is designed to support spin-polarized liquid fuel and high pressure, low-Z targets. The design has enabled us to create high pressure seals that operate at cryogenic temperatures, precisely position the target, and isolate high forces in the presence of delicate cryogenic structures.

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Top view of CUP

Fig. 4-8. The CUP is a convenient way to attach the assembled target to the cryostat at the SAUCER.

Fig. 4-9. After the high pressure shroud is removed, a thermal radiation shield is moved into position around the target.

Developments in Fuel Layering Techniques

We have continued to develop methods of forming uniform layers of condensed fuel on the interiors of laser fusion targets. As reported in the 1983 Laser Annual,⁶ we are investigating two techniques: the noncontact thermal gradient technique, aimed at providing uniform liquid layers, and the "free-fall" technique, aimed at providing solid layers. This article reports on experiments using both of these approaches. The experiments were carried out at the University of Illinois under the direction of Kyekyoon Kim.

Noncontact Thermal Gradient Method.

Workers at KMS Fusion, Inc. have shown that thermal nonuniformities can influence the distribution of liquefied hydrogen isotopes in the interiors of glass microspheres. When making these qualitative demonstrations, they induced the temperature field by cooling each sphere through an attached copper stalk while the sphere was exposed to thermal radiation.⁷ However, other methods may be used to induce the field. One that we have explored theoretically is to place the sphere between two plates at different cryogenic temperatures, and then

surround the sphere with low-density helium exchange gas. Ideally, the sphere should be suspended between the plates without contact to any support structures. In practice, of course, this is impossible, so instead we tried to minimize the thermal impact of any sphere supports. The model that we developed and reported on last year⁸ has allowed us to calculate the thermal gradients that must be impressed across the sphere to obtain uniform liquid layers.^{8,9} We have performed experiments to investigate the validity of the model, which we report here.

We constructed a target-cooling apparatus for testing the thermal gradient technique. The apparatus has two functions: to cool the sphere to liquid-hydrogen temperatures and to create a vertical temperature gradient across the sphere while it is cold. We mounted the sphere on a glass fiber and placed it in a stainless steel target-cooling cell, which had a vertical gas column and a line through which to supply ultra-pure helium as the exchange gas (Fig. 4-10). We attached the glass fiber to the end of a nylon screw, which allowed us to position the sphere in the cell. (Windows flanking the sphere enabled us to view it during experiments.) By running copper rods on the top and bottom of the column at different temperatures, we induced a roughly uniform vertical temperature gradient in the exchange gas column. The rods were in thermal contact with copper cold fingers, which in turn were in contact with a cryostat capable of supplying the differential temperatures. We were able to change the length of the gas column by changing the lengths of the copper rods. Since the glass mounting fiber has very low thermal conductivity at these temperatures, and since it is perpendicular to the temperature gradient, its impact on the temperature distribution in the sphere is minimal.

We have performed experiments with this apparatus on both D_2 -filled and HD-filled spheres. We formed uniform liquid layers for a variety of average sphere temperatures. The average temperature affects the amount of fuel condensed (and thus the layer thickness) and such liquefied fuel properties as surface tension, viscosity, and density. For each average temperature, we adjusted the temperatures of the copper rods so that the layers became uniformly distributed. We assessed the uniformity by

Fig. 4-10. Cryogenic cell for applying a thermal gradient to a stalk-mounted microsphere.

looking at shadowgrams and interferograms of the layers. The micrographs of a D_2 -filled sphere, shown in Fig. 4-11, clearly show that by applying the proper thermal gradient, the fuel can be "centered."

To compare our experimental results with the predictions of our theoretical model, we converted the temperature differences across the gas column to thermal gradients across the exterior of the sphere. We then compared these with the thermal gradients that the model predicts are required to support a uniform layer at any given temperature (and thus layer thickness). The results are summarized in Fig. 4-12 for a D_2 -filled sphere. All of the data fall on a common curve, independent of the height or pressure of the gas column. This indicates that our model and technique are generally valid, with agreement between our theoretical predictions and experimental results. The shapes of the predicted and experimental curves are similar, except for a relative shift of a few degrees between them. This shift may be due to systematic errors in thermometry and to uncertainties in the temperature-dependent properties of D_2 , but we have not resolved the discrepancy.

We also performed preliminary experiments on a 600- μm -diam sphere filled with equal amounts of H_2 and D_2 to a pressure of 75 atm; the gas-column pressure was 30 Torr. Again we adjusted the thermal gradient in the sphere to give uniform layers, as shown in Fig. 4-13. One very important difference between the results of this experiment and those of the experiments with D_2 -filled spheres is that the directions of the temperature gradients are opposite. As predicted by our model, the mixture-filled sphere is warmer at the top than at the bottom.

While there was general agreement between our theory and experimental results, that agreement was not as positive as in the D_2 case. We believe that a number of errors associated with the cell design contributed to this. In general, larger gap spacings were required for these experiments, making the thermal gradient calculations less accurate. Also, the required temperature differences were small, magnifying the effects of thermometry errors. We are building a new cell to address these issues.

Free-Fall Technique. Liquid layers on the interior of a sphere, even if uniformly

▲ Fig. 4-11. Shadow and interference micrographs of a uniform D_2 liquid layer held in steady state by an impressed thermal gradient. The glass microsphere has a diameter of 200 μm , a wall-thickness of 8.35 μm , and a D_2 fill of 133 atm.

◀ Fig. 4-12. Temperature gradient necessary to create a uniform liquid layer in a D_2 -filled microsphere.

thick, quickly form very nonuniform puddles in the bottom of the sphere as the liquid flows under the influence of gravity. One solution to this "layer slumping" problem is to eliminate gravity when condensing and freezing a fuel layer. Once

Fig. 4-13. Shadow and interference micrographs of a uniform $\text{H}_2\text{-O}_2$ liquid layer held by a temperature gradient. The glass microsphere has a diameter of $600 \mu\text{m}$, a wall thickness of $6.25 \mu\text{m}$ and a $\text{H}_2\text{-O}_2$ fill of $\sim 5 \text{ atm}$.

Fig. 4-14. Apparatus to investigate the formation of frozen layers in free-fall.

Fig. 4-15. Microsphere with a frozen fuel layer. The sphere has a $200\text{-}\mu\text{m}$ diameter, a $5\text{-}\mu\text{m}$ -thick wall, and a D_2 fill of 100 atm . (a) Interference micrograph before free-fall. (b) Interference micrograph after free-fall.

frozen, the layer will be stable. We have performed initial experiments with the free-fall apparatus described last year.⁶ An updated schematic is given in Fig. 4-14.

In a typical experiment, a microsphere is mounted inside the free-fall cell and cooled to below the triple point of the fill gas in the upper cryostat portion of the free-fall column. We examine the frozen layer, heat the frozen gas to above its condensation point with a laser pulse, and drop the cell

through the evacuated column. The gas reliquefies and freezes during free fall, and the cell is caught in the lower cryostat. We then examine the resulting layer. Results from the initial run of the apparatus with a 200- μm , D_2 -filled microsphere are shown in Fig. 4-15. Figure 4-15(a) shows an interference micrograph of the sphere taken at the top of the column, while Fig. 4-15(b) shows the corresponding micrograph at the bottom, after the free-fall refreeze.

While the layer formed during free fall may, on the whole, be more uniform than the original, this is not apparent from Fig. 4-15, which shows that in both cases a "crinkle" was present in the frozen layer. This crinkle pattern occurred consistently in several different microspheres. An analysis of the video tapes of the D_2 refreeze after vaporization showed that the pattern appeared to shift slowly, indicating that the fuel layer was not freezing solidly enough. We speculate that the freezing cell failed to reach a low enough temperature in the upper cryostat before free fall. This may have been caused by poor thermal contact between the cell and the cryostat, since the measured cryostat temperature was below 10 K. In many refreeze experiments in nonfree-fall apparatus, uniform fuel layers have been produced without crinkling. We are modifying the system to cool the cell more efficiently.

In summary, preliminary experiments to form uniform liquid layers using D_2 and H_2 - D_2 mixtures were successful. However, the results of the experiments with H_2 - D_2 mixtures were not consistent with our model. Experiments using the prototype free-fall apparatus to form solid layers proved disappointing. All parts of the system worked, but we suspect that cell temperatures were not low enough, and we are modifying the apparatus.

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Spin-Polarized DT

In 1982, Kulsrud et al.⁹ pointed out that the cross section for nuclear fusion of deuterium and tritium is increased by a factor of 3/2 if all of the nuclear spins are aligned along an external magnetic field. In

addition, calculations of the depolarizing forces in magnetic and inertial fusion reactors indicate that this polarization (or alignment) could survive long enough to enhance the performance of either type of reactor.^{9,10}

A simple analysis of ICF driver requirements¹¹ leads us to conclude that the driver energy required for a target burn of fixed gain is proportional to the inverse cube of the fusion cross section. Therefore, increasing the cross section by 3/2 reduces the required driver energy by more than a factor of 3. Because the use of spin-polarized DT fuel in fusion reactors could lead to a large saving in laser energy, we have begun an effort to polarize DT, focusing first on forming pure molecular DT and determining its spin-lattice relaxation time.

Polarization of nuclei refers to processes that give nuclear spins a net alignment in some direction. This direction is always defined by an externally applied magnetic field because in the presence of such a field an energy is associated with rotation of the spin. This field is the only "handle" by which the experimenter can orient the spins. The quantum mechanical energy for spins in a magnetic field, called the magnetic Zeeman energy, is given by

$$E_z = -\gamma \hbar m B, \quad (1)$$

where γ is the gyromagnetic ratio of the spin and m is a quantum number that takes on the values $-S, -S+1, \dots, S-1$, and S , where S is the intrinsic spin of the nucleus under consideration. Tritium has spin 1/2 and deuterium has spin 1. The value of m is proportional to the projection of the spin angular momentum of the nuclei along the applied field direction. The polarization p of a group of identical spins is simply the average of the quantum number m taken over the spins, or

$$p = \frac{1}{SN} \sum_{i=1}^N m N_{m_i}, \quad (2)$$

where N_m is the number of spins with quantum number m and N is the total number of spins. By definition $-1 \leq p \leq 1$.

Polarizing DT is advantageous for fusion because nearly all reactions $\text{D}(\text{Tn})^3\text{He}$ occur via an excited state of ^3He with total spin $|J| = 3/2$; J is the combined angular momentum of the interacting D and T nuclei. At

fusion reactor energies this reaction: only occurs with orbital angular momentum zero, so the angular momentum must come from the spins of the nuclei. If spins S_1 and S_2 are combined to give a total spin J , the allowed values of J are $S_1 + S_2$, $S_1 + S_2 - 1$, ..., $S_1 - S_2$. In the case of interacting T and D, $J = 1/2$ or $3/2$. For a given J value there are $2J + 1$ magnetic substates, with spin projections $m_J = m_1 + m_2$ taking on allowed values between $+J$ and $-J$. This means that there are four magnetic substates with $J = 3/2$ and two with $J = 1/2$. In an unpolarized plasma each of these six states is equally probable, and only those with $J = 3/2$ have a significant probability of producing fusion; therefore, only two-thirds of the interactions contribute to the fusion cross section σ . If the DT nuclei are polarized with $m_1 = 1/2$ and $m_2 = 1$, then all the interacting pairs have $J = 3/2$ and the cross section improves by the factor $1 + 2/3$, or $3/2$. Generally the polarizations will be less than unity, and the cross section is given by

$$\sigma = \sigma_0 \left(1 + \frac{p_T p_D}{2} \right) \quad (3)$$

where σ_0 is the unpolarized cross section and p_T and p_D are the polarizations of the T and D nuclei. Since the product of the polarizations determines the cross section improvement, it is essential to get both p_T and p_D close to unity to take advantage of the increase in cross section.

Another advantage of using polarized fuel in fusion reactors is that the angular distribution of the decay products is a function of the polarization, so some control over factors like neutron flux at the reactor wall is possible. This effect could also be used to study the effect of the implosion on the polarization of the fuel. The angular distribution of alpha particles and neutrons in $D(T,n)^4\text{He}$ is

$$\frac{d\sigma}{d\Omega} = \frac{3\sigma_0}{8\pi} \left\{ \frac{3}{4} (1 + p_T p_D - d_0) \sin^2 \theta + \left(1 + 3d_0 - p_T p_D \right) \left(\frac{1 + 3 \cos^2 \theta}{12} \right) \right\} \quad (4)$$

where θ is the angle between the direction of the emitted decay particle and the magnetic field and d_0 is the fraction of deuterons with $m = 0$. In the polarized case

with $p_T = p_D = 1$, $d\sigma/d\Omega$ is proportional to $\sin^2 \theta$, and the decay products come out of the reaction in a doughnut pattern perpendicular to the magnetic field. Another case of interest is $p_T = p_D = 0$ and $d_0 = 1$. Now $d\sigma/d\Omega$ is proportional to $1 + 3 \cos^2 \theta$, and the decay products are beamed along the field lines, although the total cross section is just σ_0 .

Polarization is difficult to achieve because of the relatively small value of the Zeeman energy in Eq. (1) compared to the thermal energy kT at accessible temperatures and magnetic fields. Table 4-1 gives the temperature at which the Zeeman energy is equal to the thermal energy for the listed spins in a 10-T magnetic field. The nuclei all have Zeeman energies equivalent to temperatures in the millikelvin range. Modern dilution refrigerators can reach these temperatures, but the heat produced by the decay of tritium in DT makes it impractical to cool DT to these temperatures. At higher temperatures the nuclear spins have nearly equal probabilities of populating the magnetic substates, and the polarization is nearly zero in thermal equilibrium.

The one spin that can be polarized in thermal equilibrium at temperatures to which tritium can be cooled is the electron spin. For instance, at 4.2 K in a magnetic field of 10 T the electronic polarization is 92%. Therefore, what we need is a method that will transfer the polarization of the electrons to the nuclei. Fortunately, techniques have been developed to supply accelerators with polarized proton targets. Called "dynamic nuclear polarization" or the "solid effect," these techniques involve the use of a microwave field to simultaneously flip an electronic and nuclear spin, leaving the nuclear spin with the electron's polarization. The electrons and nuclei in a solid are coupled to each other by mutual interaction of the two spins, or what is called dipole coupling.

The dipolar interaction mixes the magnetic substates of the electronic and nuclear spins in such a way that simultaneous spin flips are possible. These simultaneous spin flip-flops can occur as the spins interact with phonons in the solid or with an external microwave field that provides photons of the right energy. By interacting with the phonons in the solid, the spins come into thermal equilibrium with the environment,

The rate at which a group of spins comes to thermal equilibrium is given by a characteristic time, called the spin-lattice relaxation time. Because nuclei are magnetically weak compared to electrons, nuclei usually have long spin-lattice relaxation times T_n compared to the electron's relaxation time T_e .

The energy required to flip a spin is simply the difference between the Zeeman energies before and after the flip; therefore, if we want to flip-flop an electron-nuclear pair we must supply a photon of frequency $\omega = (\gamma - \gamma_n)B$.

Suppose a solid sample of electronic and nuclear spins is prepared at low temperatures in a high magnetic field so that the electronic polarization $p_e \approx 1$, but the nuclei have essentially zero polarization. This means that all the electron spins point in one direction, up, and the nuclei are equally likely to point up or down. We then use a microwave source to force flip-flops at a rate that is high compared to the rate of spin-lattice relaxation $1/T_n$. If the line width of the electronic transition is narrower than the nuclear transition frequency, then nuclear spins that are up cannot flip because all the electron spins are up as well, and the microwave frequency is wrong. Nuclear spins that are down can flip-flop with an electron, leaving the electron spin down. This electron could then undergo another flip-flop; however, because the electronic relaxation time is short, it is more likely to relax back to its equilibrium (up) orientation through a spin-lattice interaction, leaving the nucleus in the up orientation. The electron is then available to flip another nuclear spin, so one electron can polarize many nuclei if its relaxation time is much shorter than theirs. The nuclear polarization p_n that results from this process is given by

$$p_n = \frac{p_e}{1 + N_n T_n / N_e T_e} \quad (5)$$

where N_n/N_e is the ratio of the number of nuclei to free electrons in the sample.¹² Therefore, we want numerous electrons, a short electronic relaxation time, and a long spin-lattice relaxation time.

Several assumptions in this model require closer examination. First, where do the electrons come from? Electrons in hydrogen are

	Spin	$\gamma/2\pi$ (MHz/T)	$\gamma h \times 10^4/k$ (K)
Hydrogen	1/2	42.6	20.4×10^{-3}
Deuterium	1	6.5	3.1×10^{-3}
Tritium	1/2	45.4	21.7×10^{-3}
Electron	1/2	2.80×10^4	13.4

normally paired in molecules with opposing spins whose dipole couplings cancel. We need free electrons or electrons in atoms that are not bound in molecules. These atoms are usually made by irradiating a sample to break up molecules. In solid hydrogen, some of these atoms are trapped in the lattice and are unable to find a partner to recombine with. The electrons in these atoms are then available for polarizing nuclei. Samples with tritium manufacture their own free atoms because the decay of tritium produces an energetic electron that breaks up molecules. The concentration of free atoms that will be reached in a sample of DT is unknown and must be measured.

To derive Eq. (5), we assumed that the microwave field that forces flip-flops does not drive other transitions. This requires that the line width of the electronic transition $\Delta\omega_e$ be less than the nuclear transition frequency ω_n . This may not be the case in DT at low fields, but ω_n increases in proportion to B , so this tends to become a valid assumption at high fields. The line width $\Delta\omega_e$ must be measured to determine what field is necessary to make this assumption valid.

A second important condition for polarization is that $T_n \gg T_e$. In solid hydrogen, an important source of relaxation is the presence of rotationally excited homonuclear molecules. Rotationally excited molecules produce fluctuating electromagnetic fields that can reduce T_n by many orders of magnitude. The homonuclear molecules H_2 , D_2 , and T_2 are metastable in their rotationally excited states because the symmetry requirements imposed on the nuclear wave functions by the interchangeability of the two nuclei lead us to conclude that these molecules can only decay to their rotational ground state by changing the orientation of the nuclear spins. This requires a highly inhomogeneous magnetic field. As a result these molecules usually have long decay times.¹³

In samples containing tritium, rotationally excited molecules are produced by

Table 4-1. Temperature at which the Zeeman energy is equivalent to the thermal energy for listed spins in a 10-G magnetic field.

Cast Polymer Films

recombining atoms broken up by the decay electrons. But because the paramagnetism of atoms trapped in the lattice will cause them to decay more rapidly, their net effect must be measured. Normal DT consists of 25% D_2 , 25% T_2 , and 50% DT. If a normal sample is cooled, it will contain large concentrations of rotationally excited molecules and consequently short spin-lattice relaxation times. The ground state of T_2 also has antiparallel nuclear spins whose dipolar interactions cancel; therefore, they cannot be polarized.

For these reasons normal DT is unsuitable for spin polarization, and we must polarize pure molecular DT. We carried out the first bulk synthesis using the reaction $LiT + CH_3OD \rightarrow DT + LiOCH_3$. The product, analyzed by collision-induced infrared spectroscopy, had the following molecular composition: DT, 96; T_2 , 0.6; D_2 , 1.8; and HD, 1.6. Because of the tritium radioactivity, the DT began to decompose into D_2 and T_2 .

as soon as it was collected for analysis. A rough extrapolation of 3 h of data gave a decomposition $1/e$ time of 160 h.

During spin-polarization efforts, we will focus on measuring the quantities $T_{1\rho}$, $T_{2\rho}$, and N_1 at liquid helium temperatures in magnetic fields of approximately 1 T. These measurements will indicate how difficult spin polarization actually is, and how high a magnetic field is needed.

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Cast Polymer Films

Low-Z, oxygen-free thin films with structural integrity are desirable materials for x-ray laser substrates and fusion target capsules. For this reason, we have developed methods to obtain ultrathin films of polystyrene, a copolymer poly(butadiene styrene), and polypropylene. Using a casting procedure, we can produce these films to accurate and uniform thicknesses.

We conducted initial casting studies using polystyrene, preparing 1 to 10% solutions in dichloromethane and placing each solution into the casting apparatus shown in Fig. 4-16. Using the pressure bulb, we raised the solution level to point A. Then, using the metering valve, we lowered it at varying rates, depositing the polymer onto a glass slide. The solution concentration and the velocity at which the solution is

Fig. 4-16. Polymer casting apparatus.

Fig. 4-17. Experimental data showing the effect of liquid level velocity on film thickness for two concentrations of polypropylene.

lowered determine the thickness of the deposited polymer. A higher velocity leaves a thicker polymer than a slower velocity (Fig. 4-17). We released the film from the glass slide by floating the film on water, and transferred the film to a wire hoop for storage or transfer to a desired substrate. The quality and uniformity of these films depend on the surface quality of the substrate, in this case the glass slide.

Polystyrene was useful for initial casting studies but is not strong enough for many laser target applications. We cast a copolymer poly(butadiene-styrene) in the same manner as polystyrene, but what we derived also proved to lack structural integrity. Polypropylene, however, had both the strength and other necessary characteristics for our applications.

Until this time, polypropylene had been available only as thin as 3000 Å; this

thinness was obtained by heating and stretching. However, by dissolving the polymer in hot xylene, we have produced 100-Å-thin coatings of polypropylene.

We initiated thin-film polypropylene studies by modifying the coating apparatus to allow a constant-temperature solution (100°C). To study the effect of concentration on film thickness, we varied solutions from 0.5 to 6.0%. A higher concentration of polypropylene provided a thicker film. (Studies to quantify the relation between concentration and velocity are forthcoming.) We then fine-tuned to the desired thickness by varying the solvent-front velocity (Fig. 4-17). The structural integrity of polypropylene makes it a suitable candidate for use in fusion target fabrication.

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Carefully planned, successful experiments help determine the validity of scientific concepts. Equally important are the precise measurements with which the experimental data are verified. Using this double-pronged approach, experiments and diagnostics, target and fuel concepts, and laser components are tested. In addition, a complex computer system records and analyses the large volume of data obtained in graphic and numeric forms. One example of the specialized measuring devices is shown in the large photo at left; it shows a freestanding gold transmission grating used to diffract x rays for spectral measurement. The scanning electron micrograph shows the array of gold bars ($150\text{ }\mu\text{m}$ apart) that diffract the emerging x rays from a fusion burn and disperse them into spectrally separated images. The rough-surface edge is residue from etched silicon used as the base. The inset photo, top left, is an ellipsoidal mirror used with the x-ray diffraction grating. The photo below, a very thick zone-plate coded aperture, images x-ray emissions from a fusion implosion.

Laser Experiments and Advanced Diagnostics

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Section Editors: N. M. Ceglio and R. E. Turner

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Acknowledgments

This year marks a change in the Laser Annual process. Instead of assigning a specific person to be scientific editor, we chose to distribute the various aspects of this (sometimes thankless) job among several people.

To assure the continuity and consistency of the scientific content in the individual sections, we have always depended a great deal upon the section editors. This year, that task was to an even greater extent handed over to them, and was admirably executed. The responsibility for the overall technical and scientific continuity and consistency between individual sections was left almost entirely to the various program leaders, who also were responsible for the selection of articles.

The very often thankless part of the scientific editor's job, that of meeting deadlines, of "pushing" authors, of coordinating and juggling the schedules, and just, in general, getting the report out, was left almost entirely to our extremely capable Laser Annual staff headed by Louise Rufer and consisting of Peter Murphy and Sam Tindell. They may at times have made themselves unpopular among the authors and program leaders, but the quality and timeliness of this year's Laser Annual speaks for itself.

Publishing a book such as this would not be possible without the highly skilled technical editorial and production staffs of the *Technical Information Department*.

Finally, a special thanks goes to Lamar Coleman. Having been a scientific editor of the 1979 and 1980 Laser Program Annual Reports, he fortunately could not keep his hands off this one. His thoroughness and insistence upon changes and modifications greatly enhanced the coherence and technical quality of this year's Laser Annual.

Erik Storm

Laser Experiments and Advanced Diagnostics

Introduction

F. M. Campbell

In 1984, the Experiments and Diagnostics Program made substantial progress in the following areas: Novette laser experiments, advanced diagnostics development, computer support, and the construction of diagnostic systems for Nova.

The Novette laser, which was deactivated on September 1, 1984, proved to be the most successful experimental facility that the program has constructed to date. It was used to conduct a wide range of experiments, including laser-plasma interaction studies with 0.53 and 0.26- μm light, ablative compression of DT with 0.53- μm light, and soft x-ray laser studies. Highlights include:

- A demonstration that stimulated Raman scattering (SRS) is the dominant source of suprathermal electrons in long-scalelength coronas.
- Collisional stabilization of SRS with 0.26- μm light in high-Z plasmas.
- Substantially improved experimental and computational characterization of the corona conditions in a laser plasma.
- Measured x-ray conversion efficiencies as great as 80% with 0.26- μm light.
- Ablative compression of deuterium-tritium (DT) to inferred densities $>20 \text{ g/cm}^3$.
- The unambiguous demonstration of substantial amplification by stimulated emission at wavelengths of 153, 206, and 209 Å in neon-like yttrium and selenium.

We also renewed our efforts in the development of advanced diagnostics and associated technologies. To reflect this commitment, the program was renamed Laser Experiments and Advanced Diagnostics (LEAD). Since much experimental time on Nova will be devoted to hydrodynamics and implosion physics, a major part of the 1984 diagnostics development effort focused on these areas. New diagnostic instruments and methods include a multi-frame, sub-100-ps, x-ray framing camera and diagnostic techniques for fuel areal density measurements. We expect these diagnostics to be available for Nova experiments in late 1985 and 1986. In addition we are developing high temporal resolution (~ 50 ps) neutron diagnostics.

Expanding our resources for x-ray detection was another program priority. The need for calibrating and characterizing x-ray optics led to the joint development of x-ray charge-coupled device cameras with the Massachusetts Institute of Technology. We will soon be using these cameras to calibrate grazing- and normal-incidence x-ray mirrors and transmission gratings. Soft x-ray laser experiments involving laser radiation at wavelengths >100 Å call for measurements in this region of the spectrum. As part of the x-ray laser program, we developed a time-resolved spectrometer that operates in the desired spectral range. This instrument consists of a grazing-incidence ellipsoidal collecting mirror, a 2000-Å transmission grating, and an F1 NL soft x-ray streak camera. Its successful performance can be directly attributed to our ongoing research in state-of-the-art x-ray measurement and detection during the past 10 years.

We also made progress in the area of time-resolved instrumentation. In addition to the x-ray framing camera, we used an Austin switch to gate an extreme ultraviolet (XUV) spectrometer. This instrument was a principal

diagnostic in the 1984 x-ray laser experiments. Furthermore we upgraded LLNL streak cameras in preparation for their vital role on Nova. We increased the cameras' reliability by improving electronic components, and we incorporated a $7.35\text{-}\mu\text{m}$ optical-timing fiducial capability in all x-ray and optical cameras.

Proper laboratory facilities are essential in developing and testing advanced diagnostic technologies. As part of the advanced diagnostic effort, we established two new laboratories: the microfabrication and short-pulse laser facilities. They will play pivotal roles in developing the next generation of time-resolved instrumentation and x-ray optics.

Advanced computational facilities will be needed to process and archive the shot data from the sophisticated experiments planned for Nova. Accordingly, we improved both the hardware and the software used in data management and analysis. These enhancements include increased processing power and storage capacity for the Laser Experiments Analysis Facility (LEAF) VAX computer, the addition of an in-house microdensitometer directly linked to the VAX, and upgraded software for processing and archiving shot data.

When complete and activated, Nova will be the most sophisticated and versatile high-power laser experimental facility in the world. The target diagnostics initially installed as part of the Nova construction project will permit a wide range of experiments. We expect the first experiments with these diagnostics to begin in 1985.

Novette Experiments

Introduction

Experiments conducted in 1984 on the multi-terawatt, short-wavelength Novette laser were among the most productive ever completed by the LEAD Program. We made substantial progress in three areas of research: laser-plasma interaction, ablative compression, and soft x-ray laser studies. Experimental time on Novette was devoted about equally to inertial confinement fusion (ICF) related research and to the demonstration of a laboratory x-ray laser.

The primary objective of the laser-plasma interaction experiments was to study wavelength scaling in plasmas with dimensions of several hundred to 1000 laser wavelengths. Such large plasmas more closely resemble those expected from future high-gain targets. The short-wavelength lasers that preceded Novette had powers below $\sim 10^{11}\text{ W}$ and could not produce plasmas of this size. However, Novette produced such plasmas at powers in excess of 4.5 TW (with a $0.53\text{-}\mu\text{m}$ beam) and 1.5 TW (with a $0.26\text{-}\mu\text{m}$ beam), and with a near-Gaussian pulse of 10 ns, full width at half maximum (FWHM), irradiating both solid and exploding-foil targets.

In the case of solid targets, an overdense plasma ($n_e > n_{em}$) persisted throughout the laser pulse, and the plasma scalelength, which was not limited by the laser spot size in these experiments, scaled roughly as $c_s \tau$, where c_s is the ion sound speed and τ is the laser pulse length (FWHM). In the case of exploding-foil targets, the target thickness was less than the laser-ablation depth. This factor caused the target to decompress and produce a near-parabolic plasma density profile with an $n_e^{\text{max}}(t) \sim n_{em}$ that decreased with time. Such targets produce plasma scalelengths at $n_e < n_{em}$ that are several times greater than scalelengths from solid targets.

The laser-plasma interaction experiments included measurements of target absorption; suprathermal electron production; parametric instabilities, particularly SRS; and x-ray conversion efficiency. These experiments confirmed the more classical physical interactions we observed with submicron light obtained from smaller lasers. We also recognized the equally important role that the plasma scalelength plays in determining the mix of coupling processes. For example, the Novette experiments show clear evidence for stimulated Brillouin scattering, although there was no evidence for such processes in the Argus results. They are most likely due to the larger focal spots present in the Novette experiments. In addition,

lie above the curve labeled "absolute sidescattering." This curve shows the intensity and scalelength at which 90° Raman sidescattering would grow to saturation in a sufficiently large, plane parallel plasma irradiated by a planar laser beam.² The data did not show the systematic behavior one would expect for a planar plasma and suggest that transverse nonuniformities of the plasma and the laser beam can limit the growth of sidescattering.

Raman scattering at angles greater than 90° from the direction of the incident laser beam is less sensitive to transverse nonuniformity than 90° sidescattering. However, such scattering grows more slowly than 90° sidescattering (and grows over less distance in a planar plasma). In a planar plasma with a density gradient, Raman scattering grows most slowly at 180° . The curve in Fig. 5-3 labeled "convective backscattering" shows the intensity and scalelength at which 180° Raman backscattering¹ should produce a scattered light intensity that is 535 times the noise amplitude. The Novette experiments reach this curve at intensities of a few times 10^{15} W/cm². Thus these data should show significant amplification of the noise in the plasma produced by Raman scattering over some range of angle.

The data do, in fact, give evidence for such amplification over a wide range of angles. For example, Fig. 5-4 shows the observed Raman light amplitude as a function of average intensity for detectors located 169° from the direction of the incident laser beam. The observed scattering grows exponentially as a function of average intensity. At an average intensity of about 1.5×10^{15} W/cm² it reaches a level 535 times the apparent noise level. Considering that the target is actually irradiated by a distribution of intensities, this result is roughly consistent with the prediction shown in Fig. 5-3.

There are two other features of the data in Fig. 5-4 that we do not understand. First, the noise amplitude from which the emissions grow is quite large, about 10,000 times the blackbody emission amplitude. A recent publication has suggested a possible mechanism involving hot electrons from the $2\omega_{pe}$ instability to explain this.⁴ Secondly, the growth of the Raman scattering appears quite exponential at low intensities. This

result would be expected if the laser light incident on the plasma were a uniform plane wave. In the presence of the actual distribution of intensities irradiating the target, it is a surprising result.

Other Instabilities. In Novette experiments, we obtained evidence of several interaction processes besides Raman scattering. The Raman scattering discussed previously in this article occurred at densities from 0.1 to 0.2 times the critical density. Near quarter-critical density, the two-plasmon decay process is believed to dominate Raman scattering. We detected the two-plasmon decay instability by observing the emissions at the $3/2$ harmonic of the incident light frequency. (This corresponds to $0.35 \mu\text{m}$ for $0.53\text{-}\mu\text{m}$ light.) This $3\omega/2$ light is produced by the combination of plasma waves near quarter critical and an incident or scattered light wave. The spectrum of this radiation is similar to that observed in previous experiments.¹ The amount of $3\omega/2$ light was measured by an array of photodiodes,

Fig. 5-4. The fluence of the Raman light increased exponentially with the average laser intensity.

Fig. 5-5. The fraction of laser light scattered at $3\omega/2$ times the laser frequency as a function of average laser intensity.

as described in the articles "Summary of Absorption by Gold-Disk Targets" and "Stimulated Raman Scattering and Hot-Electron Production," later in this section. Figure 5-5 shows the fraction of the laser energy detected as $3\omega/2$ light as a function of average intensity. The $3\omega/2$ fraction increases linearly with intensity, although there is considerable scatter in the data. The linear increase may correspond to scattering of the incident laser light from a saturated level of plasma waves. The $3\omega/2$ fraction exceeds 0.01% at the higher intensities—more $3\omega/2$ light than previous experiments have reported. The exact $3\omega/2$ emission process and its dependence on the level and spectrum of plasma waves from two-plasmon decay is unknown. As a result, the small amount of $3\omega/2$ light does not guarantee that two-plasmon decay is unimportant. Nevertheless, we saw no evidence that two-plasmon decay produced significant amounts of suprathermal electrons.

A second plasma instability that may be important for ICF is the Brillouin instability. As discussed in the next article, "Summary of Absorption by Gold-Disk Targets," we observed strong, out-of-plane Brillouin sidescattering. The amount of sidescattering appeared to saturate at about 10% of the laser energy, at an intensity of about 10^{15} W/cm². We do not yet understand the saturation mechanism. We also observed that about 10% of the laser energy was more nearly backscattered, but could not determine whether Brillouin scattering or reflection from critical density was responsible.

Finally, filamentation may alter laser-plasma interactions by locally intensifying the laser light. The exact role played by filamentation in these results remains to be unraveled. We are still analyzing data from

Novette related to filamentation and are preparing to obtain better experimental evidence regarding this process in experiments on Nova.

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Summary of Absorption by Gold-Disk Targets

This article reports our measurements and analysis of the absorption of $0.53\text{-}\mu\text{m}$ light by gold-disk targets irradiated with 1-ns pulses by the Novette laser. These findings finalized preliminary results reported in 1983.⁸ Absorption is defined as the difference between incident laser energy and scattered light energy, divided by incident laser energy. For gold-disk targets, almost all the scattered light energy was at wavelengths near that of the incident green light. At most, a few percent was scattered at longer wavelengths by the Raman instability. (In contrast, the amount of Raman-scattered light from exploding-foil targets sometimes exceeded 10% of the incident laser energy.)

We measured the green light scattered from the target by means of two diagnostic systems. The green-light energy scattered into the focus lens was measured with a calibrated diode in the output sensor of the laser beam. The energy scattered throughout the chamber was measured with an array of filtered, calibrated photodiodes. One such set of measurements is shown in Fig. 5-6. The shot shown in the figure used 3.5 kJ of $0.53\text{-}\mu\text{m}$ light at an intensity of about 4×10^{15} W/cm². The energy scattered back into the focus lens was 65 J, and the total scattered light energy was 600 J. Various detectors measured the flux in plane, out of plane, and in between. The backscattering (defined here as scattering from 148 to 180°) showed azimuthal asymmetries of no more than a factor of 2. In contrast, the sidescattering (from 90 to 148°) was most often strongly out of plane. The backscattering was usually a few times more intense than the sidescattering.

These angular distributions allow us to calculate the total scattered light yield and the absorption. The diodes at each polar angle are used to define the integral of the fluence over azimuthal angle at that polar angle. Then these results are integrated over polar angle (with the $\sin \theta$ weighting to account for the variation in solid angle). The diode measurements closest to the lens are connected to the fluence scattered into the lens through a model distribution. The results were not strongly sensitive to the particular model used. The uncertainties in the diode measurements, the modeling, and the reflected-energy measurements combine to give a total uncertainty in the scattered light energy of ± 30 to 50%. For example, because the scattered light energy is typically about 25%, the resulting absorption is approximately $75\% \pm 8\%$.

Absorption results from a number of gold-disk experiments are shown in Fig. 5-7. The absorption approaches 90% at an incident laser intensity of 10^{14} W/cm² and declines gradually with increasing intensity to about 70% at 2×10^{16} W/cm². It is roughly 75% over a wide range of intensities. For comparison, Fig. 5-8 shows absorption measurements from experiments using the Argus laser. Here the absorption of 2ω light, typically about 80%, fell between that of 1ω and 3ω light. The absorption measured on Novette is similar to that from Argus for intensities below 2×10^{15} W/cm², but the physical mechanisms responsible for the scattering probably differ.

Strong out-of-plane sidescattering accounted for as much as half the scattered light in the Novette experiments. The wavelength and azimuthal distribution of the scattering on Novette suggest stimulated Brillouin scattering (SBS). In contrast, this

out-of-plane scattering was not observed in the Argus experiments at $0.53 \mu\text{m}$, in which small laser spots produced relatively steep velocity gradients. Presumably, this effect prevented stimulated Brillouin sidescattering from growing large enough to be observed. Similar results have been reported from experiments at $1.06 \mu\text{m}$.⁶

Although the larger plasmas produced by Novette allowed simulated Brillouin sidescattering to grow larger, they also decreased the amount of scattering by reflection of the laser light from the critical surface and by other processes occurring near critical density. Because the light must traverse more plasma to reach the critical surface and to escape the plasma after it is reflected, more light is absorbed by inverse

Fig. 5-6. The angular distribution of green light scattered from a gold-disk target in a Novette shot.

Fig. 5-7. Absorption measured in the Novette 1-ns gold-disk experiments.

Fig. 5-8. Absorption of 1ω , 2ω , and 3ω light by gold-disk targets, measured in experiments on Argus.

bremsstrahlung. Apparently, the reduction in this source of scattering compensated for the increasing SBS, so that both experiments produced the same absorption fraction.

At intensities above 2×10^{15} W/cm², the absorption measured when Novette was used proved to be much greater than that measured from Argus. At present, we believe this difference to be real, but we do not understand it. Two key differences between the two experiments are:

1. The Argus focal spot was half or less the size of the Novette focal spot ($70 \mu\text{m}$ vs $150 \mu\text{m}$ at best focus). This difference may have allowed more light reflected from near the critical density to escape the smaller plasma produced by Argus.
2. Only one potassium dihydrogen phosphate (KDP) crystal element was used in producing the Argus 2ω beam. In contrast, the 9 (or 25) crystals used on Novette divided the beam into segments with different relative phases. The interference of these segments may have altered the growth of the SBS instability.

For the lower-intensity Novette data, the plasmas produced are relatively planar. In these cases, the one-dimensional flow of ions away from the target determines the plasma density and velocity gradients. For the highest-intensity Novette data (at 2×10^{16} W/cm²), the laser spot size determines the gradients. The resulting two-dimensional plasma expansion and the unexpectedly high absorption remain to be explored in fully 2-D LASNEX modeling.

Modeling of the lower-intensity points is similar to that of the Shiva and Argus irradiations.^{6,7} The disk is modeled one-dimensionally as either the section of a sphere or as a slab with a fixed amount of lateral expansion. This "pseudo-two-dimensional" geometry is used for parameter and zoning studies. Fully two-dimensional calculations provide a better evaluation of the lateral extent of the plasma. They do not differ significantly from the various pseudo-two-dimensional calculations.

Some aspects of the modeling were discussed in earlier reports.⁶⁻⁸ Among these are the cross-section modeling of inverse bremsstrahlung, which is the dominant laser-light absorption mechanism. Non-LTE atomic physics is essential for this high-Z plasma with its steep temperature and density profiles in the vicinity of the critical surface.⁶ Flux-limited diffusion is used for modeling thermal electron transport; this procedure is by now routine for this type of calculation. Because suprathermal electrons have little effect on plasma evolution, they were not included in much of the modeling. This rationale is consistent with the size of the experimentally inferred hot-electron fraction (on the order of 1%), which is attributed to Raman scattering rather than to resonance absorption. (See the article "Stimulated Raman Scattering and Hot-Electron Production," later in this section, for more information.)

The chief improvement in the 1984 modeling effort over earlier work is in the

treatment of laser-ray propagation and scattering. The ray-propagation package is now fully three-dimensional (see the article "Statistical Model for Light Scattering by Density Fluctuations," in Sec. 3). The 1-D Brillouin-scattering package used earlier^{7,8} has been supplemented by 1-D and 2-D scattering packages that model turning-point scattering over a specified angular distribution and simulated scattering in underdense plasma by statistical methods.⁹

The chief constraints on modeling are the experimentally determined laser-light absorption fraction and the x-ray conversion efficiency. Matching both these quantities requires flux-inhibited transport with a flux limiter f of 0.03 to 0.05 and some provision for laser-light scattering. With a more classical transport model and/or no scattering model, calculated absorption approaches an unrealistically high 100%.

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Suprathermal Electron Production at Short Laser Wavelengths

Reducing the wavelength of laser light is expected to lower the levels of parametric instabilities. The adverse effects of these instabilities and their attendant suprathermal electrons have been discussed in earlier work.¹⁰ To verify the predictions of reduced suprathermals, we converted one beam of the Novette laser to 0.26 μm , delivering up to 1.6 kJ of ultraviolet energy on target with 0.9-ns pulses. The frequency conversion from 1.05 μm is described in Section 2 in the articles "Nova Frequency Conversion and Focusing Systems Performance" and "Operations Summary."

We used the filter-fluorescer diagnostic¹¹ to measure the absolutely calibrated x-ray spectrum from 9 keV to several hundred keV. (The upper energy limit is not more specific because the highest-energy instrument channels have poor spectral definition.) Only the SRS parametric instability was measured, by means of filtered and calibrated photodiodes. These were similar to

the photodiode array used in the green-light (0.53- μm) experiments described in the next article, "Stimulated Raman Scattering and Fio-Electron Production." The only difference was a change in filters to accommodate the shorter wavelengths. The Raman light was also spectrally and temporally resolved by a diffraction grating-streak camera combination.¹²

Gold-disk targets a minimum of 5 μm thick were irradiated by 0.26- μm light at intensities of 1×10^{14} to 2×10^{15} W/cm². These are average intensities calculated by dividing the ultraviolet (UV) energy by the pulse length (0.9 ns) and the spot area as inferred from x-ray microscope pictures. Typical energy on target was 1.5 kJ, and spot diameters ranged from 300 to 1500 μm . The incident energy was inferred from the 1.05- μm laser light energy and the conversion efficiency of the KDP crystals. The energy and pulse width of the UV light from the central crystal in the KDP array was monitored during some of the shots. Only on low-energy shots (150 J) was the pulse width noticeably shorter than that of the 1.05- μm light. These measurements are discussed in greater detail in the article "X-Ray Conversion Efficiency" later in this section. We did not attempt to adjust the quoted intensities for structure within the beam. Particularly at large spot sizes (1500 μm), the x-ray microscope pictures show considerable structure, so the effective intensities may be higher than the average ones.

The x-ray data fit a two-temperature Maxwellian spectrum with a thermal temperature of approximately 3 keV and a "hot" temperature of 30 keV. An example is shown in Fig. 5-9. The energy in hot

Fig. 5-9. X-ray spectrum from gold-disk target irradiated with 0.26- μm light at 2.3×10^{15} W/cm².

Fig. 5-10. Inferred fraction of laser energy in suprathermal electrons as a function of average laser intensity and laser wavelength (from gold-disk targets on Novette).

(suprathermal) electrons was deduced from a thick-target bremsstrahlung formula.^{13,14} This method neglects suprathermal energy lost in hydrodynamic expansion. The results are plotted in Fig. 5-10, along with data obtained during the gold-disk experiments on Novette using green light. The figure shows the expected reduction in suprathermal electrons at the shorter laser wavelength. In making this comparison, we must be careful to look only at wavelength effects and not at density-gradient thresholds for the various parametric instabilities. The green-light data shown were obtained with up to three times as much laser energy than was available at 0.26 μm , and the spot sizes were correspondingly larger. The density-gradient lengths scale roughly as the laser spot size in the transverse direction and, in these experiments, as the product of the sound speed and the pulse length in the direction normal to the target. Examining these factors, we see that the appropriately normalized gradient size (density-gradient scalelength divided by the laser wavelength) should be similar for both the 0.53- and 0.26- μm experiments.

Raman-scattered light was detected only during the shots at $2 \times 10^{15} \text{ W/cm}^2$. The SRS energy was about 7×10^{-4} of the laser energy, with an estimated accuracy of a factor of 3. Interestingly, this is the same order

of energy inferred to be present in the suprathermal electrons. A similar observation was made with higher accuracy and over a wider range of intensities during the green-light experiments, as described in the next article.

The 0.26- μm experiments lend further support to the role of parametric instabilities in producing suprathermal electrons and the control of these interactions by decreasing the laser wavelength.

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Stimulated Raman Scattering and Hot-Electron Production

High-intensity laser light can excite parametric instabilities that scatter or absorb it. One instability that can arise when laser light penetrates a plasma is sub-quarter-critical stimulated Raman (SQSR) scattering.^{15,16} It occurs below the quarter-critical density of the incident light and involves the decay of the incident light wave into a scattered light wave and an electron plasma wave. The scattered-light wavelength ranges from 1 to 2 times that of the incident light, depending on the plasma density and temperature. This article reports studies of SQSR scattering and hot-electron production in plasmas produced by irradiating thick gold targets with up to 4 kJ of 0.53- μm light in 1-ns (FWHM) pulses. These studies have important implications for laser fusion. Hot electrons attributed to the SQSR instability can increase the difficulty of achieving high-gain implosions by penetrating and preheating the fusion fuel.

From these studies came a number of results never observed in earlier work. First, we observed SQSR sidescattering of the laser light with intensities up to $0.015 \text{ J/(J} \cdot \text{sr)}$ and preferential SQSR scattering perpendicular to the plane of the electric vector of the laser beam. We attribute these observations to a combination of large axial-density-gradient scalelengths ($> 100 \mu\text{m}$), large irradiation spots (150 to 1880 μm), hot plasmas ($> 1 \text{ keV}$), and extensive diagnostics. Previous experiments have noted SQSR backscattering¹⁷⁻²⁰ or sidescattering that was orders of magnitude

less efficient.²¹ Other experiments in which 1- μm irradiation was used neither measured the angular distribution²² nor determined the scattering mechanism.²³

Second, our studies produced a total SQSR yield of up to several percent—at least three orders of magnitude greater than reported from previous experiments using solid targets.^{17,20} Efficient SQSR backscattering (but not sidescattering) has been reported from large, low- Z , underdense plasmas deliberately produced to minimize density gradients.^{19,21}

Third, this is the first report, among experiments with submicron laser light, of a large yield of suprathermal electrons (up to several percent) with a temperature higher than 15 keV. This result is important because short-wavelength lasers are being developed to reduce hot-electron production in laser fusion targets. The only previous hot-electron yields this high that were not attributed to resonance absorption were produced by 1- μm irradiation.²⁴ Low yields of very-high-energy electrons attributed to Raman forward scattering have been reported from 10.6- μm experiments.²⁴

Fourth, these are the first experiments in which the total yields of SQSR scattered light, $3\omega/2$ light, and hot electrons were simultaneously measured. The irradiation conditions were varied systematically, so we were able to attribute the hot-electron production to SQSR scattering. Thus, our results support and complement the previously reported interpretation of results at 10.6 μm (Refs. 19, 25) and 1.06 μm (Refs. 22, 26). They also are consistent with a model of SQSR scattering and hot-electron production based on computer simulations of Raman backscattering.¹⁰

The laser used in our experiments was Novette, which produced two opposed, linearly polarized, 74-cm-diam, 0.53- μm beams focused onto the target by $f/4$ doublet lenses. Output-sensor modules measured the energy and pulse history of these beams. (The residual intensity of 1.05- μm light at the target plane was about $3 \times 10^{10} \text{ W/cm}^2$.) For most of these experiments, only one beam was fired, and it struck the target while converging toward best focus. The gold targets were at least 1200 μm in diameter and 5 μm thick. X-ray micrographs (at x-ray energies of 1 to 3 keV) determined the focal spot size. The average intensity

varied from 10^{16} W/cm^2 to $2 \times 10^{16} \text{ W/cm}^2$. Based on images of the green beam obtained by the sensor package, the peak-to-average intensity variations at the target plane were no more than 10 to 1, and the intensity over half the area was within a factor of 2 of the average.

An array of filtered and absolutely calibrated photodiodes was used to measure the SQSR scattering from the target. These photodiodes were sensitive to wavelengths of 0.59 to 1.0 μm , excluding light near 1.05 μm . Time-integrated and time-resolved spectra of the Raman emissions confirmed that the spectrum was characteristic of SQSR. The Raman spectrum peaked near 900 nm, was less than 200 nm wide (FWHM), did not change markedly with irradiation conditions, and lasted for most of the laser pulse. The time-resolved spectra of the sidescattering and the backscattering were both similar to the backscattering spectra reported from a previous experiment.²⁰

Figure 5-11 shows a typical angular distribution of the SQSR light. Various detectors measured the fluence in plane, out of plane, and in between. The backscattering (defined here as scattering from 148 to 180°) showed azimuthal asymmetries of no more than a factor of 2. In contrast, the sidescattering (from 90 to 148°) was most often strongly out of plane. The sidescattering was intense, up to 0.015 J/(J·sr), and strongest at polar

Fig. 5-11. Angular distribution of Raman-scattered light.

angles of 120 to 135°. The backscattering was usually a few times more intense than the sidescattering.

We integrated the SQSR fluence over all solid angle to obtain the energy. There are enough detectors that the integrated energy was not strongly sensitive to the model used. For example, the extrapolation to $\theta = 0$ implied in Fig. 5-11 is not physical, but plausible physical models do not significantly change the integral. A few SQSR measurements obtained by detectors within the lens cone support this conclusion. The absorption of the SQSR light by inverse bremsstrahlung is small ($<10\%$). The absolute uncertainty in the SQSR yield is a factor of 2. The yield inferred from the data in Fig. 5-11 was 4.4%, with about 2/3 of the energy as backscattering. This yield was typical, but the ratio of backscattering to sidescattering varied considerably from shot

to shot. The total SQSR yield increased from about 10^{-4} at an average intensity of 10^{14} W/cm² to several percent at average intensities above 10^{15} W/cm². Comparisons of the measured yield to thresholds are difficult because the plasma conditions and the intensity distribution of the beam are not precisely known.

On an absolutely calibrated, multichannel x-ray spectrometer, we used scintillators with photomultiplier tubes to measure energetic x rays. Some channels were equipped with K-edge filters only, and others with a prefilter, an x-ray fluorescer, and a post-filter for higher resolution. To evaluate the x-ray spectrum, we used an iterative unfolding procedure that did not constrain the shape of the spectrum.²⁷ The absolute magnitude of the x-ray spectrum yields the total energy content of the hot electrons that produced the x rays to within a factor of about 2 (Ref. 28). One source of uncertainty is that some hot electrons may escape the target without producing x rays.

Figure 5-12 shows a characteristic x-ray spectrum. Each point is plotted at the energy at which that channel received 50% of its total signal. The spectrum shows one or more low-energy components with temperatures below 10 keV and an exponential decay with a temperature of 20 to 40 keV. These results correspond to a Maxwellian electron distribution of the same temperature. We never detected a more energetic "superhot" component of the spectrum. Any superhot component in the spectrum shown in Fig. 5-12 contains less than 0.003% of the laser energy. (The hot component represents 0.1%.)

The SQSR yield and the hot-electron yield (f_{hot}) increased strongly as the average intensity on target approached 10^{15} W/cm²; however, both quantities show significant scatter when compared with the average intensity and observed spot size. They correlate better with one another. The data in Fig. 5-13 show f_{hot} as a function of the yield of SQSR light. The line shows the energy in SQSR plasma waves for SQSR scattering at 900 nm (from conservation of energy). Within the uncertainties, f_{hot} equals the energy in plasma waves as both quantities increase by almost three orders of magnitude. Thus, the SQSR yield is sufficient to account for all the hot electrons, although the absolute uncertainties would allow another

Fig. 5-12. Hard x-ray spectrum from a gold disk.

Fig. 5-13. The hot-electron fraction is well correlated with the SQSR-scattered light fraction.

process of comparable magnitude and similar scaling to contribute. The fact that f_{hot} and the SQSR yield correlate better with each other than with intensity or spot size strengthens the case that the hot electrons result predominantly from SQSR scattering. In addition, the $3\omega/2$ emissions during these experiments correlate less well with intensity or f_{hot} than the SQSR emissions do. This observation suggests that instabilities at quarter-critical density are not significant sources of hot electrons in these experiments.

The observed hot-electron temperatures can also be caused by SQSR scattering. When plasma waves undergo Landau damping, they produce hot electrons with velocities comparable to the phase velocity of the plasma waves (which in these experiments is highly insensitive to the thermal-electron temperature). The phase velocities implied by the observed spectra and angular distributions correspond to electron energies of 14 to 40 keV, which are consistent with the observed temperatures of 20 to about 40 keV. The observed temperatures are not a strong enough function of intensity to result predominantly from resonance absorption.

These results have important implications for laser fusion. The hot electrons produced by the SQSR instability are of concern because, by penetrating and preheating the fusion fuel, they can make it more difficult to achieve high-gain implosions. This instability may be important because it occurs throughout large volumes of plasma and at low densities (and thus is not easily offset by inverse bremsstrahlung). The data from 10.6- μm , 1.06- μm , and 0.53- μm experiments (at low-to-moderate (λ^2)) indicate that long plasma scalelengths are necessary to produce efficient SQSR scattering from densities below quarter critical. With large plasmas, even solid targets irradiated with submicron light can produce significant scattering and quantities of hot electrons. The SQSR yields could be even greater for the larger plasmas in a laser fusion reactor, especially if the target produces a less-collisional, low-Z plasma. The important issue for laser fusion is the level at which SQSR scattering will saturate in very large plasmas. Collisional effects may help correct the problem by causing the instability to saturate more easily when even shorter-

wavelength irradiation is used. Further experiments are needed to learn the extent to which Raman scattering limits the intensity and wavelength of laser light and the collisionality of the plasmas needed for high-gain systems.

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Evidence for Collisional Damping in High-Energy Raman-Scattering Experiments at 0.26 μm

Experiments using 1.6 kJ of 0.26- μm wavelength light to irradiate thin burn-through targets show less SRS than similar experiments at 0.53 μm . At 0.26 μm , the SRS from high-Z gold targets is three orders of magnitude less than that from low-Z plastic (CH) targets irradiated at similar intensities. These results are the first direct observations of collisional damping of the Raman instability in high-Z targets as predicted by theory and simulations.

SRS²⁹⁻³⁵ in laser-produced plasmas can generate high-energy electrons, which can degrade laser fusion target gain. In nearly all previous laser-plasma experiments, SRS has been limited by steep density gradients, which restrict the phase-matching regions to relatively small lengths. However, the reactor-size targets envisioned for future laser fusion experiments will produce much longer density-gradient scalelengths. Here we report the observation of an important alternative limiting mechanism: electron-ion collisions.

In these experiments, we used UV laser light to irradiate thin films, causing them to expand rapidly. At the peak of the 1-ns laser pulse, the electron density was near resonance (quarter critical) for the Raman and two-plasmon-decay ($2\omega_{pe}$) instabilities.^{38,39} Other instabilities that can produce hot electrons occur near critical density [$n_c = 10^{21}(1.06 \mu\text{m}/\lambda_0)^2 \text{ e}^-/\text{cm}^3$], and are therefore mostly excluded from these observations.

With CH targets, the fraction of the incident energy scattered as Raman light was five times smaller: for irradiations with a laser wavelength λ_0 of 0.26 μm than with a 0.53- μm wavelength. When gold targets were used, the decrease was more than four orders of magnitude. By theory and simulation, we show that the suppression of SRS with increasing laser frequency or target Z was due to collisional damping of SRS.

These experiments used one arm of the 74-cm-aperture Novette laser.⁴⁰ An array of KDP crystals quadrupled the frequency of the 1.05- μm laser light to produce up to 1.6 kJ of light with a 0.26- μm wavelength in a near-Gaussian pulse of 0.9 to 1.2 ns (FWHM).⁴¹ A mask at the center of the lens and the chromatic aberration of the $f/4$ focusing lens caused the UV light to focus in a region free from residual 1.05- μm and 0.53- μm light.⁴⁰ The intensities quoted here are the average laser power divided by the illuminated area as determined from x-ray micrographs (made with x-ray energies of 1 to 1.5 keV). We define the beam diameter D as the FWHM of the x-ray intensity. Although the laser beam is known to be non-uniform,⁴² localized regions of higher intensities would be expected to enhance, not reduce, the growth of Raman scattering.

The targets were 6- μm -thick CH or 0.4- μm -thick gold foils supported on large washers. The thicknesses were chosen (by means of LASNEX calculations), so that the maximum electron density on-axis would be slightly less than quarter critical at the peak of the laser pulse. The parabolic density scalelengths²⁹ were calculated to be 500

to 1000 λ_0 in both the axial and radial directions for both CH and gold. The targets were irradiated at normal incidence, with the focus in front of the target. The focal spot diameters were 300 to 350 μm , resulting in average target-plane intensities of 1 to 2×10^{15} W/cm².

A streak camera and optical system, mounted orthogonal to the laser beam, provided time-resolved one-dimensional spatial images along the laser axis.⁴³ This instrument recorded the 0.31- μm bremsstrahlung emission from both the front and back: expanding plasmas. When the plasma was optically thin (as in the low- Z , green-light experiments), this emission was proportional to the line-averaged value of $n^2 T^{1/2}$. The simulation line integrals are in good agreement with the experimental observations. A time-resolved optical spectrometer measured the Raman light scattered at 135° to the laser beam. Calibrated photodiodes, filtered to look only at SRS from below quarter-critical density ($0.35 \mu\text{m} < \lambda_0 < 0.5 \mu\text{m}$), measured the absolute level of Raman-scattered light at various angles (between 105° and 169° from the laser). Spectroscopy verified that all Raman light emissions were within this range of wavelengths.

The results are summarized in Fig. 5-14. The error bars encompass all the data obtained during the experiments. We estimate the absolute accuracy to be within a factor of two. The principal source of error arises from the finite number of photodiodes (10) and the resulting uncertainty in the true angular distribution of the scattered light. The data analysis is identical to that discussed in Ref. 44. Most of the SRS energy was within 45° of direct backscattering, although no measurements were made back through the $f/4$ lens.

Figure 5-14 also shows the data from previous 0.53- μm experiments⁴⁴ conducted at similar average intensities and D/λ_0 , where λ_0 is the incident wavelength. The foils were thinner to allow the density to drop below quarter critical for the 0.53- μm light before the peak of the laser pulse. These green-light experiments produced very large amounts of SRS. The figure shows that at these intensities ($\approx 10^{15}$ W/cm²), the Z dependence in the level of Raman scattering was small in the 0.53- μm experiments. In contrast, the 0.26- μm experiments showed that Raman scattering decreases

Fig. 5-14. Energy in Raman-scattered light from low- Z and high- Z thin-foil targets. In all cases intensities are $\sim 10^{15}$ W/cm². The gold data point for 0.26 μm is an upper limit.

greatly with an increase in Z . With CH targets, the energy in the Raman-scattered light was approximately 2% of the incident laser energy. With gold targets, any SRS light present was below the minimum sensitivity of the diodes; therefore, the SRS energy could have been no more than about 10^{-5} of the laser energy. We attribute this large drop in SRS to collisional damping in the gold plasmas.

Figure 5-15 shows the Raman light spectrum from a CH target at two different times in the laser pulse. When scaled by the wavelength ratio, this spectral shape is remarkably similar to that obtained during the $0.53\text{-}\mu\text{m}$, thin-foil experiments.⁴⁴ The wavelength of the SRS implies that it is taking place at densities between 0.1 and $0.2n_c$. Assuming that the shortest-wavelength Raman light is limited by Landau damping of the corresponding plasma wave, we estimate the electron temperature to be $\sim 2\text{ keV}$. The decrease in the longest-wavelength Raman at later times is due to decreasing plasma density.

To clearly identify collisional effects, an experiment must exceed the competing gradient-scalelength threshold. Figure 5-16 is a contour plot from a 2-D LASNEX⁴⁵ hydrodynamic simulation of a CH foil, showing a large region of relatively constant-density plasma. The hydrodynamic calculations are consistent with several indirect density measurements: the time- and spatially-resolved bremsstrahlung⁴³ and the

time-resolved SRS spectrum. The calculations were further verified in the $0.53\text{-}\mu\text{m}$ experiments⁴⁴ by the timing of the emitted $3\omega/2$ light and of the burn-through of the laser light. The SRS was above the density-scalelength threshold, which is on the order of $4 \times 10^{13}\text{ W/cm}^2$ for sidescattering in a parabolic density profile.²⁹ Plasma turbulence, Brillouin sidescattering and backscattering, and light filamentation can change the thresholds, although these effects are difficult to quantify. Time-integrated x-ray

Fig. 5-15. Stimulated Raman scattering spectrum from a plastic target; t_1 is the peak time of the SRS pulse.

Simulation parameters:
 $\lambda_n = 0.26\text{ }\mu\text{m}$
 Target thickness = $6\text{ }\mu\text{m}$ (CH)
 Laser energy = 1400 J
 Laser spot diam = $385\text{ }\mu\text{m}$
 Time = peak laser pulse - 200 ps

— Electron density, n
 --- Electron temperature, T_e (keV)
 Laser intensity, I (W/cm^2)

Fig. 5-16. Contour plots of density, temperature, and laser intensity from a 2-D LASNEX simulation.

microscope pictures at 45°, 75°, and 135° from the laser beam show no structures that could be interpreted as filaments. The theoretical and computational studies here and in Ref. 46 indicate that thermal filamentation was stabilized. The diverging laser beam contributed to the stability.

Electron-ion collisions reduce SRS by: (1) directly damping the electron plasma waves, (2) absorbing the scattered-light wave that amplifies Raman, and (3) absorbing the incident laser light directly. To exceed the collisional threshold³³ in a plasma of constant electron density n requires:

$$\gamma^2 > (v_{ei}/2) (\omega_p/\omega_s)^2 [(v_{ei}/2) + v_{el}] \quad (1)$$

where γ is the Raman growth rate, ω_p and ω_s are the plasma and scattered wave frequencies, respectively; v_{el} is the electron Landau damping rate; and v_{ei} is the electron-ion collision frequency relevant to high-frequency waves³⁷:

$$v_{ei} \approx 9 \times 10^{-11} \bar{n} \ln(\Lambda) T_e^{-3/2} s^{-1} \quad (2)$$

where $\ln(\Lambda) \approx 31 - \ln(n^{1/2}/T_e)$, $\bar{n}$ is the average ionization state, T_e is the electron temperature (keV), and n is the electron density (cm⁻³).

The Raman backscattering growth rate³⁵⁻³⁷ is given by

$$\gamma \sim (k_{epw}/4) v_{osc} [(\omega_p/\omega_s) (c/v_{epw})]^{1/2} \times [1 - (n/n_c)]^{-1/4} \quad (3)$$

where ω_{epw} is the electron plasma wave frequency, k_{epw} is the wave vector of the electron plasma wave as shown below:

$$k_{epw} \sim \omega_0/c \left[1 - (n/n_c) \right]^{1/2} + \left[1 - 2(n/n_c)^{1/2} \right]^{1/2} \quad (4)$$

Table 5-2. Ratio of collision frequency to Raman growth rate, and corresponding Raman yield for various targets and laser wavelengths.

λ_0 (μm)	Target	v_{ei}/γ	Raman yield (E_{SRS}/E_{laser})
0.26	CH	0.2	0.02
0.26	Au	1.8	<10 ⁻⁵
0.53	CH	0.04	0.10
0.53	Au	0.2	0.04

for $k\lambda_D \ll 1$, and

$$\frac{v_{osc}}{c} = \left[\frac{I(W/cm^2) \left(\frac{\lambda_0}{1.06 \mu m} \right)^2}{1.2 \times 10^{18}} \right]^{1/2} \quad (5)$$

Equations (1) and (3) determine a collisional threshold intensity:

$$I(W/cm^2) > 3 \times 10^{-12} \left(\frac{k_{epw} c}{\omega_0} \right)^{-2} \left(\frac{\omega_{epw}}{\omega_s} \right) \times v_{ei} \left(\frac{v_{ei}}{2} + v_{el} \right) \alpha \lambda_0^{-3} \quad (6)$$

(The wavelength scaling is from the following analysis.)

Kinetic simulations³⁸ of Raman backscattering with collisions show that when n/n_c is about 0.2, SRS is severely reduced for values of v_{ei}/γ around 0.6 to 0.8. These values are lower than the collisional threshold condition of $v_{ei}/\gamma \sim 2$ from Fq. (1).

To calculate v_{ei} , one must know T_e and $\bar{n}$. We estimated the electron temperature of CH targets using the short-wavelength cut-off of the SRS spectrum and time-resolved x-ray spectroscopic data⁴⁴ from CH targets doped with a small amount of sulfur. Both measurements give $T_e \approx 2$ keV. LASNEX 2-D calculations for gold (similar to Fig. 5-16) show $T_e \sim 3$ keV and $\bar{n} \sim 56$. Figure 5-16 indicates that inverse bremsstrahlung absorption reduced the incident laser intensity by a factor of 2 to 3 before the light reached the density at which most of the Raman scattering occurred. (Similar effects were observed with the gold targets.) Table 5-2 shows the estimated v_{ei}/γ for $\lambda_0 = 0.26 \mu m$ and $\lambda_0 = 0.53 \mu m$ experiments. Note that for the 0.26-μm cases, $v_{ei}/\gamma \approx 1$ for gold; for CH, with $\bar{n}$ a factor of about 15 smaller, $v_{ei}/\gamma \ll 1$. This variation can account for the large difference in the observed SRS.

Finally, we performed LASNEX simulations^{48,49} of thick gold disks to roughly estimate the scaling of v_{ei}/γ . Least-squares fits of T_e and $\bar{n}$ from runs using 1-ns, 450-μm laser spot diameters are: $T_e \sim 1.2 \times 10^{-6} \mu m^{0.5} (\lambda_0/1.06 \mu m)^{0.85}$ and $\bar{n} \sim 4.48 \mu m^{0.06} (\lambda_0/1.06 \mu m)^{-0.06}$. These T_e and $\bar{n}$ values are consistent with the thin-film results shown in Fig. 5-16. T_e and $\bar{n}$ also

increase with the spot size and pulse length and are sensitive to other geometric factors. We used a heat-flux limit $f = 0.03$ for $\lambda_0 \approx 0.53 \mu\text{m}$ and $f = 0.1$ for $\lambda_0 = 0.26$. Solving for v_{ei}/γ from these fits and Eq. (2) gives

$$\frac{v_{ei}}{\gamma} \sim 0.04 \left(\frac{I}{10^{15} \text{ W/cm}^2} \right)^{-1.1} \left(\frac{\lambda_0}{1.06 \mu\text{m}} \right)^{-3.3}, \quad (7)$$

which is plotted in Fig. 5-17.

In summary, we have presented experimental data comparing reduced Raman scattering from low-Z targets irradiated with 0.26- μm laser light to those irradiated with 0.53- μm light. Raman scattering was not observed from long-scalelength, high-Z plasmas similarly irradiated with 0.26- μm light. These results are consistent with calculations and simulations, indicating that the suppression of SRS is due to collisions.

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Exploding Foil Characterization

Introduction. During 1983 and 1984, we conducted experiments with exploding-foil targets to produce and characterize large plasmas. The fabrication of these targets and preliminary results of the experiments were reported in 1984.^{50,51} This article presents the final results of experiments in which we used 0.53- μm light to irradiate thin CH foils. Gold foils were also used, and the results reported earlier.⁵¹ Similar experiments with 0.26- μm light produced evidence for collisional damping of Raman scattering, as discussed in the preceding article, "Evidence for Collisional Damping in High-Energy Raman-Scattering Experiments at 0.26 μm ."

We perform large-plasma experiments because the plasmas produced by future high-gain targets will be much larger and more

Fig. 5-17. Rough scaling of the collision frequency-to-Raman growth rate ratio (v_{ei}/γ) vs laser wavelength for a variety of laser intensities on gold disks; v_{ei} is calculated for quarter-critical density.

uniform than those produced by current disk or implosion targets. Because larger plasmas allow more growth of plasma instabilities, some processes that are currently unimportant might grow to damaging levels. This possibility is not a mere presumption, but has already been witnessed in the case of Raman scattering, as discussed in the article "Summary of Parametric Instabilities," earlier in this section. To guide our choice of laser wavelength, energy, and target design for future experiments, we must continue to study appropriately large plasmas and observe their instabilities. Research of this nature will proceed on Nova with the goal of obtaining a predictive understanding of plasma conditions and important plasma instabilities.

Figure 5-18 illustrates the setup for exploding-foil experiments on Novette. One of the $f/4$ beams struck the target while converging toward best focus. The target was a CH foil, 1 to 3 μm in thickness, which often included a 300- μm -diam microspot in the center. The microspot was doped with roughly 4 at.% sulfur⁵⁰ but was otherwise identical to the rest of the target. The laser energy was typically 3.5 kJ, with a spot diameter of 600 μm and pulse length of 1 ns. The average intensity was about $1.2 \times 10^{15} \text{ W/cm}^2$. Figure 5-18 shows the typical pulse shape of the incident beam, the light reflected back into the focus lens, and the light transmitted into the opposed focus lens. The relative and absolute timing of the reflected and transmitted light are only approximate for this shot. The foil exploded during the laser pulse, and as the plasma density decayed, the reflection of the laser

Fig. 5-18. In exploding-foil experiments, a Novette beam struck the thin CH target, with the pulse shown at the lower left. As the target expanded, reflection decreased and transmission increased.

light decreased and the transmission increased.

We varied the thickness of the CH foil and used the optical emissions from the plasma to diagnose the changes in the decay of the plasma density. Figure 5-19 shows data obtained by means of the optical/x-ray (OX) streak camera⁵² from foils with thicknesses of 1, 2.1, and 3 μm . The bottom row of plots in the figure shows the 1.05- μm laser pulses used to produce the 0.53- μm light. The width and center of each pulse are marked. The top row of plots shows the emission of light at the 3/2 harmonic of the incident 0.53- μm laser light. These emissions are a signature of the quarter-critical surface and disappear when the peak plasma density drops below quarter critical ($n_c/4$). The 1- and 2- μm -thick foils both dropped below quarter-critical density before the center of the laser pulse. In contrast, the 3- μm -thick foil remained above quarter-critical density much longer.

(Note that only part of the 3- μm foil need have been above quarter critical to produce the observed emissions.) Finally, the middle row of plots in Fig. 5-19 shows the Raman-scattered light, which was emitted from densities above 0.1 critical. These results indicate that if the Raman light ceases while the laser light is still intense, the density has dropped below 0.1 critical. The 1- μm foil decayed below 0.1 critical before the center of the pulse, but the 2- and 3- μm foils did not. In all these cases, the optical emissions allow us to diagnose the peak plasma density and to see the effects of changes in the experiment.

We also obtained density data from the time-resolved Raman spectrum and from the bremsstrahlung emissions. The bremsstrahlung data are discussed in the next article, "Optical Bremsstrahlung as an Electron-Density Diagnostic for Laser-Produced, Exploding-Foil Plasmas." The Raman spectrum provides density

Fig. 3-14. Timing of $3\omega/2$ and Raman emissions relative to the incident laser pulse for three foil thicknesses.

Experiment: $\lambda = 0.53 \mu\text{m}$, 1 ns (FWHM), 3 to 4 kJ, 0.6 to 1.4×10^{15} W/cm 2 ,
 1.8- to 2.1- μm CH foils

2-D LASNEX: $\lambda = 0.53 \mu\text{m}$, 1 ns (FWHM), 3 to 3.7 kJ, 0.6 to 1.3×10^{15} W/cm 2 ,
 2.0- to 2.1- μm CH foils

Fig. 5-20. A comparison of LASNEX predictions with experimentally measured plasma density.

information as a result of its long-wavelength cutoff, as indicated in the article "Raman Scattering in Exploding Foils" later in this section. When the plasma decays below a certain density, the Raman scattering at the wavelength corresponding to that density disappears. In the future, we hope to obtain additional density data from x-ray atomic-line spectroscopy.

Figure 5-20 compares the density data just described with predictions from two runs of the LASNEX computer code simulating 2- μm -thick plastic foils. The LASNEX calculations were two-dimensional, Lagrangian, hydrodynamic simulations with a flux limiter of 0.03. In the two calculations, the energy, intensity, and thickness parameters were: 3 kJ, 6×10^{14} W/cm 2 , and 2.0 μm ;

and 3.7 kJ, 1.3×10^{15} W/cm², and 2.1 μ m. The higher-intensity calculation produced a more rapid density decay.

The left half of Fig. 5-20 shows the simulated decay of the maximum plasma density with time, and the right half shows the density profile along the laser axis at the center of the laser pulse. Raman-scattering, 3/2 harmonic, and bremsstrahlung data are also shown. Although the shot-to-shot variations are significant, the data support the predictions within the uncertainty range. (A low-density plume, discussed in the next article, is an exception.) The LASNEX code correctly predicted the approximate shape and decay of the plasma density. It is notable that one of the LASNEX runs shown in the figure was performed before the experiments and was used to help determine the experimental conditions. Thus these experiments have given us increased confidence in our ability to calculate plasma production from exploding foils.

In addition to the density measurements, these experiments produced indications of plasma temperature. The short-wavelength cutoff of the Raman spectrum, which we attribute to Landau damping, implies a plasma temperature of about 2 keV, if a Maxwellian electron distribution is assumed. The temperature implied by the ratio of the x-ray lines emitted by the sulfur dopant is also near 2 keV. These results and their comparison to several LASNEX runs are discussed in the articles "X-Ray Line Spectroscopy on Exploding-Foil Plasmas" and "Raman Scattering in Exploding Foils," later in this section. The 2-D LASNEX calculations described above predict a temperature of 2.5 keV, and additional 1-D calculations predict a somewhat higher temperature. There are enough uncertainties and difficulties in the measurement, interpretation, and calculation of temperature that the actual value (or the equivalent in terms of mean energy) might differ from 2 keV by as much as a factor of 2.

Finally, we observed that large, well characterized plasmas produced significant absorption and scattering of the laser light by parametric instabilities. Raman scattering and stimulated Brillouin sidescattering each reflected roughly 10%. (Stimulated Brillouin sidescattering produced in these experiments was similar to that from disk targets described in "Summary of Absorption by

Gold-Disk Targets" earlier in this section.) The 3 ω /2 fraction, indicating the two-plasmon decay instability, exceeded 0.01%; this result was also similar to those produced with the disk targets.

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Optical Bremsstrahlung as an Electron-Density Diagnostic for Laser-Produced, Exploding-Foil Plasmas. This article describes a new diagnostic technique employing time-resolved photographic images of ultraviolet bremsstrahlung emission from hot, optically thin plasmas. This technique permits the calculation of the plasma electron density profile. It differs from traditional bremsstrahlung diagnostics, which are based on emission from dense, rather than thin, plasmas.

The new technique measures the line integral $\int n_e^2 \bar{Z} \ln \Lambda / \sqrt{T_e} dx$, where n_e is the electron density; $\bar{Z}$ is the average ionization state; $\ln \Lambda$ is the Spitzer factor or Coulomb logarithm (usually between 5 and 10, and relatively constant); T_e is the electron temperature; and x is the distance along the path of the line integral. If the temperature dependence of the emission is exceedingly weak, the line integral can be validly used to determine the plasma density profile. Furthermore, if $\bar{Z}$ and T_e do not vary greatly, we can deduce the density profile without including their variation.

The new diagnostic was demonstrated in exploding-foil experiments conducted to study instabilities in plasmas with long density scalelengths.⁵³⁻⁵⁶ We obtained a time-resolved, edge-on image of the 0.31- μ m (0.01- μ m FWHM) bremsstrahlung light emitted from plasmas produced by irradiating 2.1- μ m-thick CH foils. The foils were irradiated with 3.8 kJ of 0.53- μ m laser light at an intensity of 1.3×10^{15} W/cm² in 1-ns pulses. The foils exploded, producing a completely underdense plasma early in the pulse. The hydrodynamic measurements we obtained through this diagnostic technique are in good agreement with computer simulations. These measurements are relevant to

Fig. 5-21. Experimental setup. The streak-camera slit is parallel to the laser axis.

the laser fusion effort both because they test the validity of the LASNEX hydrodynamic simulation code and because they allow us to compare plasma instability measurements with theoretical predictions.

A schematic of the experiment is shown in Fig. 5-21. A major part of the apparatus is an inverse $10 \times f/4.7$ Cassegrainian telescope, which is nearly diffraction-limited. The second principal component is an LLNL streak camera that has a $500\text{-}\mu\text{m} \times 22\text{-mm}$ contact slit and is equipped with an RCA S20 streak tube with a sapphire window and an ITT 40-mm-diam microchannel plate (MCP) intensifier. The streak-tube grid, normally 2.3 kV positive with respect to the photocathode, was biased 2.7 kV negative about 20 ns after the end of the 118-ps/mm sweep. The bias was introduced to reduce background signals from the late-time target emission. The $500\text{-}\mu\text{m}$ -wide slit does not degrade the temporal resolution since the RCA streak tube is highly demagnifying along the time axis when the slit is positioned midway between the streak-tube grid wires.⁵⁷ Prior to these experiments, the focus voltage changed because of the failure of a power supply. This change degraded the temporal resolution to only 200 ps and the spatial resolution to only 200 μm at the target. Both types of resolution should have been 10 times better.

We achieved a calibration for the diagnostic instrumentation that we believe is accurate to within a factor of 2. The calibration was carried out as follows. The streak camera was first absolutely calibrated at

$0.53\text{ }\mu\text{m}$. A calibration wedge was identically exposed onto every film negative for both this calibration and the experiments. The film densities at any fixed position on the calibration wedges always correspond to the same film exposure. However, the light we imaged and time-resolved in these experiments was at $0.31\text{ }\mu\text{m}$. The calibration value at $0.31\text{ }\mu\text{m}$ was obtained from that at $0.53\text{ }\mu\text{m}$ by multiplying by the ratio of the dc sensitivities of the S20 photocathode at the respective wavelengths. The transmission or reflection coefficients of all the optics were measured at $0.31\text{ }\mu\text{m}$. This information, combined with the absolute calibration of the streak camera, the telescope solid angle and magnification, and the FWHM of the bandpass interference filter, gave us the overall calibration for the instrument. The accuracy of the calibration—to within a factor of 2—corresponds to a $\sqrt{2}$ uncertainty factor in the electron density n_e .

The streak record in Fig. 5-22 is the source of all the experimental data discussed in this article. The figure shows that the foil exploded symmetrically except for a low-density, high-velocity plume that jetted back toward the laser. This symmetry is to be expected since the foil plasma became underdense very early in the laser pulse, at about $t = -700\text{ ps}$ in the calculations. The dip in the brightness at the initial foil plane is due to obscuration by slow-moving cold plasma at large radii. The dip would have been much more pronounced except for the poor spatial resolution of the streak camera and the finite f number of the telescope.

Fig. 5-22. Streak record of the time-resolved ultraviolet thermal light from an exploding plastic foil. Experimental conditions: 2.0- μm CH foil, 5270 Å, 1 ns, 3.8 kJ, 600- μm spot diameter, $f/4.3$ converging beam, 1.3×10^{15} W/cm² peak intensity.

The CH foil was much larger than the laser spot size and was stretched across a thin washer. If the foil had been freestanding and exactly the same size as the laser spot the dip would have been absent. Although the streak record shown in Fig. 5-22 lacks a timing fiducial, we determined the absolute timing of the streak camera from other streak records for CH foils that had a 0.53- μm fiducial approximately coincident with the laser pulse. The fiducial was omitted in some cases because it reduces the length of slit usable for imaging the plasma.

Plasma brightness can be characterized by a Rayleigh-Jeans equivalent temperature T_{RJ} . This is approximately equal to the blackbody-equivalent temperature T_{be} only when the UV photon energy $h\nu$ is much less than kT_{RJ} , where k is the Boltzmann constant. The Rayleigh-Jeans brightness function $B_{\nu}^{RJ}(T)$ is

$$B_{\nu}^{RJ}(T) = \left(\frac{2\nu^2}{c^2} \right) kT \times [\text{erg}/(\text{s} \cdot \text{sr} \cdot \text{Hz} \cdot \text{cm}^2)] \quad (8)$$

Here ν is the UV photon frequency in Hz, and c is the speed of light. If κ_{ib}^* is the

reduced inverse bremsstrahlung absorption coefficient for the intensity, L is the length of the plasma, and T_e is the electron temperature, the Rayleigh-Jeans equivalent temperature is given by Eq. (9). It follows from detailed balance that the right-hand side of Eq. (9) is the brightness of the plasma:

$$B_{\nu}^{RJ}(T_{RJ}) = [1 - \exp(-\kappa_{ib}^* L)] B_{\nu}(T_e) \quad (9)$$

$B_{\nu}(T)$ is the blackbody brightness function:

$$B_{\nu}(T) = \left(\frac{2h\nu^3}{c^2} \right) \frac{1}{\exp(h\nu/kT) - 1} \times [\text{erg}/(\text{s} \cdot \text{sr} \cdot \text{Hz} \cdot \text{cm}^2)] \quad (10)$$

All these brightness functions are already summed over the two polarizations. The advantage of using T_{RJ} is that it is proportional to brightness and thus to the streak-camera film exposure.

The absorption coefficient κ_{ib}^* has a simple form when $h\nu \ll kT$:

$$\kappa_{ib}^* = \frac{4\sqrt{2}\pi}{3} \left(\frac{Zr_0^2}{\beta_e^2} \right) \times \ln \left(\frac{Q_{\max}}{Q_{\min}} \right) \left(\frac{n_e^2}{n_e \sqrt{1 - n_e/n_c}} \right) \text{cm}^{-1} \quad (11)$$

Here $Q_{\max}$ and $Q_{\min}$ are the maximum and minimum momentum transfer in electron-ion collisions; n_e is the electron density; n_c is the critical density for the observation wavelength (0.31 μm); r_0 is the classical electron radius (2.818×10^{-13} cm); $\beta_e = \sqrt{kT_e/(m_e c^2)}$, where m_e is the electron mass; and

$$\bar{Z} = \frac{\sum_i Z^2 n_i}{\sum_i Z n_i} \quad (12)$$

is the average charge state of the ion species i (n_i is the ion density). The argument

Fig. 5-23. Electron density profile vs distance from the initial foil plane at 0.5 ns after the peak of the laser pulse.

of the logarithm is⁵⁶

$$\frac{Q_{\max}}{Q_{\min}} = \frac{1.123 m_e c^2}{\alpha \lambda h \nu} \left(\frac{k T_e}{m_e c^2} \right)^{3/2} \quad (12)$$

when $\eta = Ze^2/(h\nu) \gg 1$, where v is electron velocity, and $\lambda \lesssim \sqrt{\pi}/(n_e r_D)$. Here α is the fine-structure constant $1/137.036$ and λ is the UV photon wavelength.

We assume that the plasma is optically thin, so that $1 - \exp(-\kappa_{ib}^* L) \simeq \kappa_{ib}^* L$. Then

$$\frac{T_{Rj}}{T_e} \simeq \kappa_{ib}^* L \quad (14)$$

when $h\nu \ll kT_e$. Notice that $\kappa_{ib}^* T_e$ has a $\bar{Z} \ln A/\sqrt{T_e}$ temperature dependence.

Figures 5-23 and 5-24 compare the electron density inferred from the UV bremsstrahlung emission to the actual electron density on axis. The path length L in Fig. 5-24 was derived from fits to simulations by the LASNEX 2-D hydrodynamics computer code^{59,60} and is given by

$$L = 1.2 \text{ mm} + \frac{|z|}{2}, \quad (15)$$

where z is the distance from the plane defined by the initial foil position.

A more meaningful comparison with the data was made from a simulated streak record produced by the TDG⁶¹ post-processor code operating on LASNEX-produced output files. The TDG simulation was treated just as the actual streak record data were, including the use of Eq. (15).

Fig. 5-24. Electron density profile vs distance from the initial foil plane at a distance of (a) $z = 0.5$ mm, and (b) $z = 1.0$ mm, from the initial foil position. Time $t = 0$ is the peak of the laser pulse.

TDG should have reproduced the actual streak record except that it did not simulate the poor resolution or finite f number of the experimental apparatus. Except for the plume, agreement is good. This plume is real unless the plasma is somehow parametrically generating $0.31\text{-}\mu\text{m}$ light. The TDG simulation of the streak record is shown in Fig. 5-25. The parameters $\Delta E = 0.5$ eV and $\Delta\Omega = 0.1$ sr are different from the experimental parameters, but these differences were canceled by normalization before Eqs. (8), (11), (13), (14), and (15) were applied to obtain the electron density on axis.

LASNEX performs Lagrangian, cylindrically symmetric, hydrodynamic simulations. It models laser light as a large number of rays that must be parallel to or intersect the Z axis; no angular momentum is allowed. Each ray is refracted according to geometric optics and absorbed according to inverse

Noveite Experiments

Fig. 5-25. Streak image as calculated from the hydrodynamic simulation.

bremsstrahlung and resonant absorption. The absorbed energy is transported both by radiation and by electrons with a flux limiter f equal to 0.1. About 700 ps before the peak of the laser pulse, the center of the foil becomes underdense to the $0.53\text{-}\mu\text{m}$ laser light. By the peak of the pulse, the plasma on the center line is quite underdense, as can be seen from Fig. 5-26. The peak density is $n_c/8$ (n_c is the critical density, at which the laser frequency equals the plasma frequency); $Z = 3.3$; and the temperature varies from 1.6 to 2 keV in the region of interest.

LASNEX saves the emissivity and opacity for each photon bin energy and for each zone at specified timesteps. TDG then calculates the line-integrated emission at a specified angle, which in this example is perpendicular to the laser beam axis.

Fig. 5-26. Contour plots of the LASNEX-calculated electron temperature T_e and electron density n_e .

In conclusion, we have compared the 0.31- μm spectral output from an exploding plastic foil with the results of a hydrodynamic simulation and found good agreement. UV bremsstrahlung emission appears to be a good electron-density diagnostic for low- Z plasmas.

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X-Ray Line Spectroscopy on Exploding-Foil Plasmas. X-ray line spectroscopy has been used to study the dynamics of exploding foils in underdense-plasma experiments.⁶² Time histories of the x-ray lines were modeled and correlate to experimental results within 100 ps. A hydrodynamic simulation from LASNEX⁶³ served as input to a code that solves time-dependent rate equations for atomic levels. The rate equations are used to calculate the ionization balance and atomic-level population of an ion species as it moves through the time-dependent temperature and density profiles of the plasma. The predicted line intensities of this model generally agree with the observed line intensities, although ionization during the early ablation phase of the experiment is underestimated. This discrepancy might indicate that the high-energy tail of the electron distribution penetrates the denser plasma.

X-ray spectroscopy on exploding low- Z foils was done by placing higher- Z tracer elements in the foil. A concentration of sulfur amounting to 4 at.% was placed in the center 300- μm diameter of the CH foil.⁶⁴ Since the spot size of the beam was about 600 μm , it overfilled the doped region. By limiting the sulfur to the center of the irradiation area, we minimized any effects the edge plasma might have on line intensities. The low sulfur concentration was expected to change the average charge state of the plasma by less than 20%. Simulations indicate that the sulfur produced a negligible perturbation of the foil hydrodynamics and caused minimal opacity in the spectral lines. Optical depths were estimated to be less than 1, even for the resonance lines of the spectra.

Fig. 5-27. Time-resolved x-ray spectrum from a thin, sulfur-doped CH foil.

The time-resolved spectra were measured by the x-ray crystal streak camera (XCSS) on Novette.⁶⁵ An example of the data is shown in Fig. 5-27. The data were taken with a potassium acid phthalate (KAP) crystal at a principal Bragg angle of 10.1°. The dispersion was 74 eV/mm at the slit, and the resolution was 7 eV. The $n = 2$ to $n = 1$ (α) and $n = 3$ to $n = 1$ (β) features from the He-like and Lyman series were recorded. The resolution was not sufficient to distinguish the He-like intercombination line (3P_1) from the resonance line (1P_1).

Time histories for the He α and Ly α lines are shown in Fig. 5-28, in which the time resolution is ~ 50 ps. Both time histories extend approximately 350 ps (FWHM)—considerably shorter than the 900-ps laser pulse. They show similar structure during the rising part of the pulse. On the trailing edge, however, the He α intensity decreases more rapidly than the Ly α , presumably when the foil burns through and becomes underdense to the 0.53- μm laser. A continuum emission appearing late in the pulse around the He α line was subtracted from the data. As a result, only an upper limit could be set for several data points after burn-through.

Calculated time histories are compared with experimental data in Fig. 5-28. These data have no absolute timing marker, so the time axis was normalized to the calculations at the half-intensity point on the leading edge of the He α line. The peak of the laser pulse is at $t = 0$. The overall agreement

Fig. 5-28. Time history of the He α and Ly α lines of sulfur compared with computer simulations.

— Simulations
• Experiment

Shot 93112105
CH foils, 2.1 μm thick
 $I = 1.3 \times 10^{15} \text{ W/cm}^2$
 $\lambda = 0.53 \mu\text{m}$
Peak of pulse = 2000 ps

Fig. 5-29. Time history of the Ly α -to-He α line ratios compared with calculations using two different flux multipliers.

between the data and the calculations is quite good; the widths of the time histories agree to within 100 ps. The Ly α time-history shows a leading edge that starts earlier than

calculations predict and decays more rapidly than predicted.

The Ly α -to-He α ratio is shown in Fig. 5-29. The large scatter in the data is due to the low signal levels and the inherent noise in the streak camera. The ratio is less than 1 early in the pulse and rises quickly to about 2. When the foil burns through, the ratio quickly increases to about 4 before the He α signal is lost in the continuum noise. If one uses a simple, collisional, steady-state model assuming uniform temperature, the ratio would indicate a plasma temperature around 1.5 keV at early times and 2 keV before burn-through. After burn-through, the inferred temperature would be 2.5 to 3.0 keV—in approximate agreement with temperature estimates derived from Raman spectra.

The calculated ratio is also shown in Fig. 5-29. The calculation underestimates the ratio at early times as shown by the early rise in the data for the Ly α line. After burn-through, the line ratio plateaus as the charge states are frozen into the underdense plasma. Comparisons were made for two flux multipliers in LASNEX, $f = 0.03$ and $f = 0.1$. The principal difference between the two results is higher coronal temperature in the underdense plasma as demonstrated by the higher ratio of He α to Ly α for $f = 0.03$. This value is in better agreement with the data, although the $f = 0.1$ curve is within the uncertainty range. The smaller flux multiplier also predicts a longer

time for burn-through in relation to the peak of the pulse. Since the XCSS has no absolute fiducial, the absolute burn-through time could not be measured. The predicted shapes of the time histories from both flux multipliers are similar, and both agree with the data as shown in Fig. 5-28.

Although the modeling predicts the overall time history of the line intensities and ratios, Ly α intensity increased more rapidly than calculated. This result may be evidence for nonlocal transport of electrons into the denser solid. The early presence of the Ly α signal is related to the ionization of the sulfur as it moves through the heat front. As the ion moves through the steep gradients, it requires a finite time to equilibrate with the local plasma temperature and density. This time is related to the collision frequency, which in turn is related to the ionization energies for He-like and H-like sulfur—3.2 and 3.5 keV, respectively. Electron energies two to three times these ionization energies are necessary to ionize the sulfur to a higher charge state.

Throughout the heat front, the model assumes a Maxwellian distribution for the thermal electrons at the local temperature. In the model, the electron temperature is greater than 1 keV only in the underdense corona, where the electrons have enough energy to efficiently ionize the K-shell; however, the plasma density decreases rapidly in the corona, and the ion quickly decouples from the plasma. On the other hand, electrons with energies greater than 6 keV can travel from the outlying corona through the heat front and ionize the sulfur. This movement increases the time during which the sulfur ions sample the energetic electrons and leads to faster equilibration. For example, the range of a 10-keV electron in cold matter is comparable to the total thickness of the foil. Penetration of such energetic electrons from the tail of the electron distribution can significantly preionize the sulfur. The possibility of this effect indicates that better nonlocal electron transport models are needed.

In summary, the time histories of the x-ray lines from sulfur confirm the hydrodynamic modeling of the thin-foil experiments. The line ratios favor a flux multiplier of 0.03, although a flux multiplier as high as 0.1 cannot be ruled out. The data show

evidence of nonlocal electron transport; better electron-transport modeling may be needed to predict ionization of the sulfur by energetic electrons traveling through the ablation region.

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Raman Scattering in Exploding Foils.

Exploding-foil experiments allowed us to study parametric instabilities in plasmas with a large-scale-length density gradient. This article reports time-resolved spectroscopic measurements of SRS obtained during irradiation of foil targets with 0.53- μ m light. The diagnostics we used were described previously.⁶⁶ Foil experiments with 0.26- μ m light are discussed in the article "Evidence for Collisional Damping in High-Energy Raman-Scattering Experiments at 0.26 μ m," earlier in this section.

Figure 5-30 shows the time-resolved SRS spectrum from a CH foil target irradiated with green light at 10^{15} W/cm². We analyzed its features as described below.

From the laser frequency and SRS-frequency measurements, we inferred the plasma wave frequency by using the SRS phase-matching relationships.⁶⁷ From the plasma wave frequency we estimated the density of the plasma at the density of the plasma at the site where scattering

Fig. 5-30. Time-resolved SRS spectrum from a CH foil target irradiated with green light at 10^{15} W/cm².

occurred. We consistently found that the highest-frequency (shortest-wavelength) SRS light from the CH targets was about $0.67\omega_p$. The corresponding electron density is $0.09n_c$, where ω_p is the laser frequency, and n_c is the critical density. The highest-frequency SRS observed from gold targets was $0.6\omega_p$, corresponding to a density of $0.12n_c$.

The phase-matching relations plus the experimental geometry (measurements were at 135° from the laser beam) enabled us to calculate the wave number k of the plasma wave associated with a particular scattered light wave. The highest-frequency SRS light was associated with the plasma wave that had the largest wave number. We would expect this plasma wave to be strongly Landau-damped if the Debye length λ_D , a function of electron temperature and density, were very large. Calculating the density as described above and assuming that higher-frequency SRS is not observed because of Landau damping, we set $(k\lambda_D)^2 = 0.1$, the Landau-damping limit. From this expression, we obtained an estimated upper limit for the plasma temperature T_e . For CH targets, T_e was 2 keV, and for gold targets, it was 3.5 keV. The CH result is in good agreement with temperatures inferred from the x-ray spectroscopy described in the preceding article, "X-Ray Line Spectroscopy on Exploding-Foil Plasmas." Both temperatures are in reasonable agreement with LASNEX calculations. Further confidence in this technique comes from the measurement of SRS from thin foils irradiated in the x-ray laser experiments. The Raman spectrum obtained

from these experiments predicted an 800-eV temperature limit, and LASNEX predictions were again in agreement.

The long-wavelength (low-frequency) limits of the SRS spectrum showed two interesting features. First, the lowest-frequency SRS observed was $0.54\omega_p$, corresponding to an electron density of $0.2n_c$. Even though the foil plasma was overdense initially, we observed no SRS light from the region just below and including the quarter-critical density. Similar observations have been reported from virtually all other SRS measurements; at present, there is no satisfactory explanation.

The second feature to note in the low-frequency end of the SRS spectrum is that the presence of the longest observed wavelength decreases with time. Because longer-wavelength SRS comes from higher plasma densities, the simplest explanation is that the plasma expands to lower densities and excludes longer-wavelength Raman processes. In Fig. 5-31, the maximum observed SRS wavelength is plotted versus time for three separate shots. Because the diagnostic lacked an absolute time reference, the time scales for each shot were adjusted to overlap. The right-hand vertical scale shows the densities that correspond to these peak wavelengths. The inferred expansion rate is in good agreement with the hydrodynamic simulations described in the article "Optical Bremsstrahlung as an Electron-Density Diagnostic for Laser-Produced, Exploding-Foil Plasmas," earlier in this section.

A question remains about the absence of SRS enhancement near the maximum density, where the gradient vanishes. There are at least two possible explanations. First, since we do observe convective SRS below the density plateau, we are obviously well over its threshold. Therefore, we must address the more difficult problem of determining the (nonlinear) saturation levels of Raman scattering at various densities. The second explanation relates to the finite time resolution of the instrument. Considering the hydrodynamic expansion rate, we estimate that a given frequency has high SRS gain in the plateau region for only a few picoseconds. The instrumental response, which is on the order of 50 ps, will smear out the SRS from the plateau region, so that this SRS signal can be noticeably stronger than the convective SRS below the

Fig. 5-31. Maximum observed SRS wavelength as a function of time for three shots irradiating CH targets.

plateau only if it is at least an order of magnitude stronger to begin with. Given the strength of the convective SRS signal (up to 10% of the laser energy) a much higher-strength signal from the plateau region seems unlikely.

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X-Ray Conversion Efficiency

In Novette experiments, we measured the conversion efficiencies for x rays emitted from gold-disk targets. We extended the experimental conditions to shorter wavelengths, larger spot sizes, and higher intensities than possible in earlier measurements with the Argus laser. X-ray conversion efficiency is defined as the energy E_x emitted by a target in thermal (0.1 to 1.5 keV) x rays normalized by the absorbed laser energy E_{abs} . The conversion efficiency is a sensitive measure of the coupling of laser energy to high-density matter. Modeling of the conversion efficiency is an integrated test of assumptions in laser absorption physics, hydrodynamics, and radiation physics.

X-ray conversion efficiency was studied extensively using Argus. Efficiencies were measured for incident-light intensities from 3×10^{13} to 1×10^{15} W/cm² at 1.06, 0.53, and 0.35 μm , in 600-ps pulses on targets with various atomic numbers.⁶⁸⁻⁷⁰ In general, the Argus results showed that the conversion efficiency increases with decreasing wavelength and with decreasing incident light intensity. For 0.53- μm and 0.35- μm light, the conversion efficiency decreased or remained constant when the intensity was below 1×10^{14} W/cm²; but with 1.06- μm light, the conversion efficiency continued to increase with decreasing intensity. Furthermore, at 3×10^{13} W/cm², the conversion efficiency for 0.53- μm light was lower than for 1.06- μm light. Simulations predict that the conversion efficiency would continue to increase with decreasing laser intensity (for the range covered in the experiments) and that shorter-wavelength light would always result in a higher conversion efficiency.

On Novette, we measured x-ray conversion efficiencies at 0.53 μm and 0.26 μm . The laser's higher power permitted larger focal spots and higher intensity than was possible with Argus. With larger spots, plasma expansion tends to be planar, and the plasma undergoes less lateral energy transport and spherical divergence. With smaller spots, plasma expansion is more spherical, resulting in steeper density gradients, and there is a larger proportion of relatively cooler edge plasma. Cooler plasma emits lower-energy thermal x rays. It can also absorb emission from hotter parts of the plasma if the cooler portion is along the observation line of sight. The Novette results at 0.53 μm agree with the Argus data for intensities 10^{14} W/cm² and greater. At 3×10^{13} W/cm², the Novette conversion efficiency is significantly higher than that on Argus. Novette conversion efficiency at 0.53 μm continues to increase with decreasing laser intensity, more in agreement than Argus data with simulations. Two-dimensional simulations are in progress to study the effect of spot size at low intensities. The 0.26- μm experiments on Novette extended the wavelength scaling study to shorter wavelengths. The conversion efficiency with 0.26- μm light was significantly greater than with 0.53- μm light for similar intensities. At 10^{14} W/cm², the x-ray conversion efficiency was over 80% for 0.26- μm light and 60% for 0.53- μm light. Above 10^{15} W/cm², the difference is not so great; we observed approximately 60% efficiency for 0.26- μm light compared to 45% for 0.53- μm light.

The x-ray conversion efficiency was measured by irradiating flat gold foils with one beam of the Novette laser.⁷¹ The foils, which were more than 5 μm thick, remained opaque to the laser for the irradiation conditions. The incident 1-ns pulse shape was nominally Gaussian before passing through the frequency-conversion crystals. The converted laser pulse $\sim$ both 0.53- μm and 0.26- μm light had a width of about 900 ps (FWHM). The peak of the 0.26- μm laser pulse appeared flattened, probably because of nonlinear effects *ir*, the conversion process at the highest intensities. For 0.53- μm light, the incident intensity varied from 3×10^{13} to 3×10^{15} W/cm². Data for 0.26- μm light were taken at two intensities: 1×10^{14} and

Fig. 5-32. Measured and calculated x-ray conversion efficiency from gold disks for laser wavelengths from 1.06 to 0.26 μm .

$2 \times 10^{15} \text{ W/cm}^2$, by fixing the energy at 1.5 kJ and changing the focal spot. (For more information, see the article "Operations Summary" in Sec. 2.) Each data point for 0.26- and 0.53- μm light, as shown in Fig. 5-32, represents the average of at least two measurements except at the highest intensity of 0.53- μm light.

The quoted incident intensity is the average intensity, defined as the total incident energy divided by the illuminated area and pulse width. Locally the intensity can vary significantly because of hot spots. We used x-ray microscope images of the target to make a qualitative assessment of the beam uniformity. Focal spots less than 500 μm in diameter showed little structure. Spots over 1000 μm in diameter showed near-field structure from obscurations due to the crystal array and the split in the disk amplifiers. In the 0.53- μm experiments, the area of the obscurations was small, so it is reasonable to use the outside diameter of the spot as an estimate of the illumination region. In the 0.26- μm experiments, the x-ray emission region was highly structured, and obscurations appeared to be much larger than at 0.53 μm . If the x-ray intensity is assumed to be proportional to the incident laser intensity, the average laser intensity is approximately

$1.5 \times 10^{14} \text{ W/cm}^2$. By using the outside diameter of the laser spot and roughly correcting for the unirradiated area from the x-ray images, we estimated the irradiation intensity to be $9 \times 10^{13} \text{ W/cm}^2$. The error limits for the intensity shown in Fig. 5-32 indicate the range of variation in the spot as well as errors in determining its average value.

The absorbed laser energy was determined for 0.53- μm light by subtracting the reflected energy from the incident energy. The incident energy was determined from 1.05- μm calorimetry and crystal conversion efficiency curves. The reflected energy was measured by an array of scattered light diodes, as described in the article "Summary of Parametric Instabilities," earlier in this section. Errors in the absorbed energy measurement are estimated to be about $\pm 10\%$. For 0.26- μm light, errors were larger because of the measurement technique. An incident-light calorimeter measured 0.26- μm light directly, but sampled only part of the beam as described in the article "Operations Summary" in Sec. 2. These errors are estimated at about 15%, resulting in larger error bars for the 0.26- μm data points in Fig. 5-32. Shots in which the entire beam was sampled were used to "calibrate" this procedure. The reflected energy for 0.26- μm light was not measured and is assumed to be small. Absorption of 0.26- μm light in short pulses ($< 200 \text{ ps}$) was greater than 90% at intensities between 3×10^{13} and $1 \times 10^{15} \text{ W/cm}^2$ with a variety of targets.⁷² Absorption is expected to be at least as high when pulses are longer. The assumption of 100% absorption should introduce errors smaller than those in the incident-beam calorimetry.

The x-ray spectrum was measured by means of a 10-channel Dante spectrometer.⁷³ Two systems, Dante A and Dante H, were used. Dante A was 75° from the west Novette beam, and Dante H was 90° from both beams. Most observations were made at 60° from the disk normal by rotating the disk normal 15° from the west beam toward Dante A or 30° from the east beam toward Dante H. For 0.53- μm light at an intensity of $3 \times 10^{15} \text{ W/cm}^2$, the disk was normal to the west beam, and Dante A viewed it at 75° from the disk normal. For

0.26- μm light, only the east beam and Dante H were used. X-ray spectra were unfolded from the Dante data by means of calibrated channel responses.⁷⁴ Some changes in the Dante x-ray diode response calibrated before and after the experiments were observed in portions of the spectral region. These changes in calibration resulted in a $\pm 10\%$ change in the integral of the unfolded spectrum and are within the uncertainties of the measurement.

The x-ray conversion efficiency from the Argus experiments is also shown in Fig. 5-32. The unfolded spectra were integrated from 0.1 to 1.5 keV. The Novette measurements were made at only one viewing angle, in contrast to the angular distribution measured on Argus. The angular dependence of the x-ray emission in the Novette experiments is assumed to be Lambertian, implying that the disk source is flat, so that the angular distribution varies as the cosine of the viewing angle with respect to the disk normal. Most of the data were taken at a 60° observation angle—the viewing angle shown by simulations and simple models to be the least sensitive to the plasma shape.⁷⁵ The x-ray measurement at 60° is the same whether the source is spherical or Lambertian.⁷⁵

The vertical error bars in Fig. 5-32 include errors due to uncertainties in the x-ray spectral measurements as well as the energy normalization. These include uncertainties in detector calibration and in the unfolding of the x-ray spectrum. Uncertainties due to integration of the Argus angular distribution are not included. Relative uncertainties among the Novette data points are less than indicated by the error bars by about $\sqrt{2}$ because the detection system was similar for the entire data set.

Conversion increased dramatically when the wavelength was changed from 0.53 μm to 0.26 μm . The increase is much greater than observed at Argus when the wavelength was changed from 0.53 μm to 0.35 μm . This increase implies much better coupling of energy to higher densities for 0.26- μm light. Critical electron density for 0.26- μm light is $1.6 \times 10^{22} \text{ cm}^{-3}$ compared with $4 \times 10^{21} \text{ cm}^{-3}$ for 0.53- μm light.

The x-ray conversion efficiency was modeled by LASNEX, and the results are also

shown in Fig. 5-32. This modeling was done in a pseudo-two-dimensional fashion with a transport-inhibiting flux limiter f of 0.03 (Ref. 70). Calculations for 1.06-, 0.53-, and 0.35- μm light were used to model the earlier Argus results,^{68,69} while the modeling for 0.26- μm light was done more recently. Although there are some differences in physics, the x-ray conversion efficiency for the two sets of modeled results should be similar.

The modeling is generally accurate in predicting trends in the data. The large increase in x-ray conversion efficiency with 0.26- μm light is correctly predicted. The x-ray conversion efficiency also agrees with observations and is insensitive to the flux limiter at $1 \times 10^{14} \text{ W/cm}^2$ when f is changed from the transport-inhibiting value of 0.03 to more classical transport. At $2 \times 10^{15} \text{ W/cm}^2$, the predicted x-ray conversion efficiency changes from 0.54 to 0.72 when the flux limiter changes from inhibited to classical transport. These calculate 1 values bracket the observed result of 0.62, and both are within the uncertainty range of the data. The simulations of the 0.26- μm irradiation are relatively insensitive to the flux limiter because most of the laser light is absorbed by inverse bremsstrahlung at relatively high electron densities near critical ($1.6 \times 10^{22} \text{ cm}^{-3}$), and little transport occurs. At intensities of 10^{14} W/cm^2 , the data agree with the simulation, since most of the x-rays are emitted and absorbed in the same region. Their limited transport makes the x-ray conversion efficiency insensitive to the flux limiter. At $2 \times 10^{15} \text{ W/cm}^2$, however, some light reaches critical density, and more emission comes from densities greater than critical; thus, the x-ray conversion efficiency becomes more sensitive to the flux limiter.

The measured x-ray spectra are compared with calculations in Fig. 5-33. The data points from individual detectors are plotted at the energy where 50% of the detected signal originates from lower energy and 50% comes from higher energy. The solid curve is a smoothed fit of the spectrum. For an intensity of $1 \times 10^{14} \text{ W/cm}^2$, the spectrum is nearly flat from 0.2 to 0.8 keV, while for $2 \times 10^{15} \text{ W/cm}^2$, the spectrum strongly peaks above 0.5 keV. The region

Fig. 5-33. Low-energy x-ray spectra from gold disks irradiated in 1-ns pulses with 0.26- μm light at two intensities.

above 0.5 keV is characterized by N-band emission from ionized gold. The ionization balance in the emission region greatly affects the intensity of the signal. LASNEX modeling predicts the overall spectral shape. It shows the peak at $2 \times 10^{15} \text{ W/cm}^2$ due to the N-band emission, although the peak is calculated at slightly higher energy than observed. The calculations seem to predict a higher temperature in the emission region than the experiment indicates. Such temperatures would cause the gold to strip to higher charge states and would shift the resulting emissions to higher energies.

The 0.53- μm data from Novette agree with the Argus results and with simulations for intensities greater than 10^{14} W/cm^2 . At $3 \times 10^{13} \text{ W/cm}^2$, the conversion efficiency measured on Novette continued to increase with decreasing intensity, but on Argus the data peaked around 10^{14} W/cm^2 and decreased at lower intensity. The Novette data are in better agreement with simulations. The principal difference between Novette and Argus is the energy available for irradiation. Novette had 100 times more energy than Argus, thus allowing irradiation diameters to be an order of magnitude larger. This difference in spot size may explain the difference in conversion efficiency at $3 \times 10^{13} \text{ W/cm}^2$. No spot-size dependence was observed at the higher intensity, even though the spot diameter increased by a factor of 5 from Argus to Novette. On Argus, the spot size was varied at an intensity of $3 \times 10^{13} \text{ W/cm}^2$, and some differences in conversion efficiency were observed; however, they were within the uncertainty of the measurements and were deemed not significant.

The spectra from Novette and Argus are plotted in Fig. 5-34 for intensities of 3×10^{13} and $3 \times 10^{14} \text{ W/cm}^2$ of 0.53- μm light. The data are normalized to the incident laser energy. The points represent data from the individual detectors, and the curves are smoothed to fit the data. At $2 \times 10^{14} \text{ W/cm}^2$, both sets of data have similar shapes. The Dante system on Argus had only one channel measuring the spectrum below 0.25 keV. A significant portion of its signal came from energies above 0.25 keV. The steep rollover in the Argus data is due to the signal from this one channel. The rollover is not significant because of the uncertainties in unfolding the response of this channel.

For an intensity of $3 \times 10^{13} \text{ W/cm}^2$, the spectra show large differences, especially between 0.7 and 0.9 keV. X-ray intensity increases in this region as the spot diameter increases. This is the region influenced greatly by N-band emission in gold. The spot size appears to affect ionization in the gold plasma. With smaller spots, the gold ionization from the emission region does not strongly excite these bands at such low irradiation intensities. With a larger spot, these bands are more strongly excited. This difference in band emission may be a result of changes in energy balance due to expansion. For large spots, plasma expansion is nearly planar, but for small spots, the expansion is more spherical. Spherical expansion allows more energy to be conducted laterally from the irradiation area. Although the energy loss is estimated not to be large, it may be sufficient to change the ionization balance and thus the emission of the gold. Another explanation for the variation in

Fig. 5-34. Low-energy x-ray spectra from gold disks irradiated with 0.53- μm light at (a) $3 \times 10^{14} \text{ W/cm}^2$ and (b) $3 \times 10^{13} \text{ W/cm}^2$ for different laser spot sizes.

band emission with spot size is that the radiation dynamics may depend on spot size. Optical depths and line trapping may depend on plasma size. In turn, they can change the ionization in the plasma. At the higher intensities, spot-size effects are not observed, perhaps because the gold is sufficiently ionized that small changes in the plasma dynamics do not affect the band emission. At higher intensities, the emission between 0.7 and 0.9 keV is high for both spot sizes, so energy loss due to lateral spreading or opacity effects may have little impact.

In summary, Novette experiments extended x-ray conversion efficiency measurements to higher intensities, larger spots, and shorter wavelengths. Conversion efficiency at 0.26 μm was significantly higher than at 0.53 μm for intensities below $3 \times 10^{13} \text{ W/cm}^2$. The difference in conversion efficiency between 0.26- μm light and 0.53- μm light on Novette was much greater than observed between 0.35- μm light and 0.53- μm light on Argus. The high conversion efficiency with 0.26- μm light was predicted by modeling and indicates an efficient coupling of short-wavelength laser light with high-density matter. Novette results with 0.53- μm light do not exhibit the rollover at low intensities observed on Argus. The

spectral shapes indicate that the conversion efficiency at low intensity is very sensitive to spot size, especially between 0.5 and 0.9 keV. Since this is the region at which gold N-band emission appears, this sensitivity indicates that spot size affects the ionization balance in the emission region. Simulations are in progress to examine this effect.

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Implosion Experiments

Ablatively driven, high-density implosions are essential for high-gain ICF targets.⁷⁶ The success of such implosions depends critically on the nature of the laser-plasma coupling, fluid instabilities, and drive symmetry. In recent years, numerous theoretical and experimental studies have demonstrated the strong dependence of laser-plasma coupling on laser wavelength.⁷⁷ Increases in target absorption,⁷⁸⁻⁸¹ x-ray conversion efficiency,^{31,82} mass ablation

rates and pressures,⁸³⁻⁸⁶ and reduction in fast-electron generation^{83,86} have been observed during target irradiation at submicron laser wavelengths. These improvements suggest that short-wavelength lasers can drive ablatively imploding DT fuel mixtures to thermonuclear ignition conditions. The minimum laser energy required for successfully driving high-gain targets to ignition will depend on the effects of asymmetry and fluid instabilities on the final fuel conditions.

Despite the improved coupling observed with submicron light, few systematic implosion experiments have been performed because of the relatively low laser energy and power available at submicron wavelengths ($E < 500$ J, $P < 0.5$ TW). Furthermore, for directly illuminated, electron-driven targets, inadequate symmetry due to the limited number of beams or the focusing geometry has also restricted the scope of the experimental implosion studies.

In this article, we report the results of target implosion experiments with the multi-kilojoule, 0.53- μ m Novette⁸⁷ laser. Comparison with the results from the 1.06- μ m Shiva experiments shows improved target compression with decreasing laser wavelength.⁸⁸

Indirect-drive (x-ray-driven) targets were used in these experiments. These targets consisted of DT fuel capsules mounted inside high-Z radiation cases, which were heated directly by the laser. The target configuration provided an alternative to direct illumination of the fuel pellet. Instead the focused laser energy was converted into soft x rays, which ablated the pellet's external surface. Successful production of a uniform radiation environment around the fuel capsule can result in a very symmetrical implosion, even in a two-beam laser system like Novette. Furthermore, since the pellet does not directly see the laser light, the optical focusing scheme used to deliver beam energy to the target does not necessarily play a pivotal role in the target performance. Thus, the targets used in the indirect-drive experiments on Shiva and Novette were configured in a nearly identical manner, although optical parameters differed greatly. (On Shiva, 1.06- μ m radiation was focused with two 10-beam clusters. The individual beams had an f number of 6,

and the effective f number of the clusters was 1.2. Novette was a two-beam, 0.53- μ m laser with an f number of 4.) To complete the reduction of independent variables to a single parameter, the laser wavelength, we also kept the laser output energy and the pulse length approximately constant during both series (~ 7 to 9 kJ and 0.7 to 1.1 ns FWHM, respectively). To optimize the conversion of focused laser energy into soft x rays, irradiation focal spots were set to produce beam intensities that optimized light absorption and x-ray conversion efficiency at each wavelength.

Capsule performance was characterized by measuring the 14.1-MeV neutron yield from DT fusion and the compressed pusher areal density integrated over the burn time. The neutron yield was measured with an absolutely calibrated Pb activation system. The pusher areal density $\rho\Delta r$ was obtained by detecting the activity of ²⁸Al created by neutron activation of a ²⁸Si tracer in the pusher via the ²⁸Si(n,p) ²⁸Al reaction.^{89,90}

The pusher areal density, sampled over the burn time (calculated to be < 100 ps), is related to the ²⁸Al activity by the expression⁹⁰

$$\rho\Delta r_p^b = 466 \times \left[\frac{N^*}{\eta_d \eta_c e^{-\lambda^* t_0} (1 - e^{-\lambda^* \Delta t}) Y} \right], \quad (16)$$

where $\rho\Delta r_p^b$ is the pusher areal density in (g/cm²), λ^* is the ²⁸Al decay constant (0.309 min⁻¹), η_d and η_c are the detector and target-debris collection efficiencies, respectively, N^* is the total number of nuclides that are counted in the experiment, Y is the neutron yield, t_0 is the delay time between the implosion and the start of counting (approximately 0.5 min), and Δt is the counting interval (5 min). The detector efficiency is measured by the β - γ coincidence counting of the ²⁸Al, and the target debris efficiency is determined by the ²⁴Na tracer technique.⁹¹

The pusher areal density measures the compression of the pusher and is related to the imploded fuel density (or areal density). For an ideal one-dimensional implosion, using arguments of mass conservation and continuity of pressure and temperature at the fuel-pusher interface at stagnation, and

assuming no mixing and uniform densities at burn time, one can show that⁹²

$$\frac{\rho_f^b}{\rho_f^i} = G(M_p, M_f, Z_p) \left[\frac{(\rho \Delta r)_p^b}{(\rho \Delta r)_p^i} \right]^{3/2}, \quad (17)$$

$$G(M_p, M_f, Z_p) = \frac{M_p^{3/2}}{3M_f(4\pi)^{1/2}} \times \left\{ \alpha \left(\frac{3M_f}{4\pi} \right)^{1/3} \left[\left(1 + \frac{M_p}{\alpha M_f} \right)^{1/3} - 1 \right] \right\}^{-3/2} \quad (18)$$

where M_p and M_f are the masses of the pusher and fuel respectively, Z_p is the average charge state of the pusher, ρ_f^i and ρ_f^b are the respective initial and burn-time fuel densities, $(\rho \Delta r)_p^i$ and $(\rho \Delta r)_p^b$ the pusher initial and burn-time areal densities, and α is the ratio of the pusher density to the fuel density, which is determined by the degree of ionization of the pusher. In our experiments, $\alpha \approx 2$. Realistic radial density profiles of the compressed target do not alter Eqs. (17) and (18) significantly. Equation (17) agrees well with models of gas-filled capsules, in which the entire hydrodynamic history of the implosion is simulated by means of the LASNEX code.⁹³ The relation between the measured pusher areal density and the inferred fuel density does depend upon the configuration of the imploded target. While one can show that departures from spherical symmetry do not significantly alter the interpretation of the results, pusher-fuel mixing can raise the $\rho \Delta r$ at a given fuel density.⁹² For example, if approximately 5% of the pusher mass is uniformly mixed with the fuel, then the actual fuel density is 40% lower than the value inferred from the neutron-activation measurement.

In Fig. 5-35, we summarize the results of compression experiments performed on both Novette (0.53 μm) and Shiva (1.06 μm). These data, represented by closed and open circles, include only experiments in which neutron activation was used. The figure also includes representative data from high-intensity, electron-driven compression experiments. These experiments are included to demonstrate the sensitivity of the pusher areal density to target performance.

Figure 5-35 shows that the 0.53- μm x-ray-driven target is reached higher final pusher areal density ratios than targets similarly imploded at a laser wavelength of 1 μm . For the data shown here, the average areal density ratio of the Novette targets is about 54, approximately 3 times the corresponding average from the highest 1.06- μm experiments. If we substitute the measured average pusher areal density ratio at each wavelength into Eq. (17), we obtain for the Novette experiments an average burn-time fuel density of approximately 120 times the density of liquid DT ($\rho \approx 0.2 \text{ g/cm}^3$). For the Shiva experiments, the result is 56 times the density of liquid DT. A more rigorous analysis, using LASNEX simulations, confirms that the 0.53- μm data shown in Fig. 5-35 are consistent with the 50-fold increase in $\rho \Delta r$ that one would expect from a target that is 100 times the density of liquid DT. The fuel densities inferred from the highest-compression Novette experiments are consistent with the measured ablation pressures and the fuel adiabat, which is determined by shock heating of the DT due to the rapid rise of the driver pulse. High-energy x-ray measurements show that target preheating by fast electrons is negligible. The fuel density inferred from Eqs. (17) and (18) is insensitive to modest levels of x-ray preheating of the pusher.

Figure 5-35 shows that the neutron yields of the high-density Novette experiments are comparable to those achieved in the highest-density Shiva implosions. LASNEX

Fig. 5-35. Ratios of compressed to initial target pusher areal densities. Experiments with x-ray and electron-driven targets are included.

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and a simple model for the neutron yield suggest a mass-averaged ion temperature of 500 to 600 eV for both sets of experiments.

The neutron yields obtained in the Novette experiments with the largest pusher areal density had yields that were a factor of 3 to 10 lower than our one-dimensional simulations predicted for the measured driving conditions. For ion temperatures <1 keV, the neutron yields scale approximately as T_i^2 . This sensitivity to T_i makes the yield dependent on many details of the experiment, such as implosion symmetry and pusher-fuel mix. None of these features were adequately monitored in the experiments. For example, computer simulations show that the ion temperature (and neutron yield) can be modeled by mixing a sufficient amount of pusher material with the DT. Future experiments are planned on the Nova laser to investigate such issues.

As mentioned previously, Fig. 5-35 also shows representative data obtained at Shiva from experiments with two types of directly illuminated targets. The first was a high-yield, thin-walled, exploding-pusher target with a wall thickness of $1.5\text{ }\mu\text{m}$ and a diameter of $300\text{ }\mu\text{m}$. It was filled with 2 mg/cm^3 of DT and illuminated with a high-intensity, 100-ps pulse.⁹⁰ The second type of target was an intermediate-density glass microballoon $140\text{ }\mu\text{m}$ in diameter and $5\text{ }\mu\text{m}$ thick.⁹⁴ The target was coated with $15\text{ }\mu\text{m}$ of $\text{CF}_{1.2}$ containing 10 mg/cm^3 of DT and was irradiated with 15 to 20 TW of $1.06\text{-}\mu\text{m}$ light in a 200-ps (FWHM) pulse. Neither type of target achieved high compression because both the pusher and fuel were imploded on a high adiabat. The implosion adiabat of these targets is set by fast-electron preheat resulting from the short-pulse, high- $I\lambda^2$ irradiation ($I\lambda^2 > 3 \times 10^{15}\text{ W/cm}^2\cdot\mu\text{m}^2$). In contrast to the x-ray-driven targets, the relatively small increase in pusher areal density observed in these experiments demonstrates the modest compressions that were achieved. The ratio of the compressed pusher areal density to its initial value was only ~ 1 for the exploding pusher target and ~ 5 for the thicker-walled targets.

In conclusion, we have presented laboratory data that show increased target compression at submicron laser wavelengths. The increase can be directly attributed to

improvements in laser-plasma coupling such as a reduction in hot-electron preheating associated with short laser wavelengths. The preheating levels observed on Novette are consistent with the requirements for the implosion of targets on the Nova laser that are hydrodynamically equivalent to high-gain targets. We believe that the low neutron yields in the submicron-wavelength experiments may have been the result of inadequate drive symmetry or hydrodynamic instabilities. The data suggest that problems of mix and asymmetry need to be resolved through further experiments.

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X-Ray Laser Experiments

Introduction. In mid-1984, after more than four years of effort, we succeeded in obtaining laser amplification of soft x rays. With Novette as the pump source, we conclusively demonstrated lasing action at XUV wavelengths. These experiments, which were prefaced by the experimental validation of computer code predictions at KMS Fusion, Inc., are described in the next article, "Collisional-Excitation Laser Scheme." We also tested a recombination scheme and a resonant-photopumping scheme, which are described in the subsequent articles, "Exploding Foil Recombination Laser" and "Resonant Photoexcitation Laser Experiments." Using the successful collisional-excitation scheme, we completed 112 shots. Only 11 shots were fired on the recombination scheme, and 10 on the resonant-photopumping scheme. Our work with recombination and photopumping, although not successful in demonstrating lasing, will improve our knowledge of physics and help us to design better tests in the future.

The success of these experiments and the important potential of this research for both civilian and military applications motivated us to add a two-beam target chamber to the recently completed Nova facility. The chamber (which is actually the existing chamber that was used on Novette) will utilize 2 of

Nova's 10 beams. It will be fully compatible with all of the x-ray laser diagnostics now on hand. Because this hardware is already developed, we will be able to resume our experiments by late summer 1985 with Nova. Although we plan eventually to use the full energy and power of the 10-beam Nova system, there are a number of near-term x-ray laser goals that we can most efficiently pursue with the 2-beam chamber. Some of these goals are:

1. Operate the selenium soft x-ray laser with a longer plasma so as to achieve gain saturation.
2. Characterize the coherence and divergence of this x-ray laser scheme.
3. Demonstrate applications.
4. Develop and demonstrate a cavity oscillator.
5. Continue to model and experimentally test schemes that produce shorter-wavelength lasers.
6. Develop techniques or designs to increase the efficiency of producing an x-ray laser in order to facilitate the transfer of this technology to the private sector.

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Collisional-Excitation Laser Scheme. For over a decade, researchers sought a way to produce amplification of XUV (or soft x-ray) radiation. Numerous papers review the history of these theoretical and experimental efforts.⁹⁵⁻¹⁰² Prior to 1985, the literature contains many investigations, but lacks conclusive evidence that such amplification was actually produced. In this article we review some proposed methods for producing XUV amplification and describe the first experiments in which amplification was conclusively demonstrated.

One proposed method of producing an XUV amplifier calls for a high-power optical laser to heat a plasma directly. Elton¹⁰³ first suggested that XUV lasing could be accomplished by the collisional excitation of $n = 3$ levels in carbon-like ions that are produced in a laser-heated plasma. Later Zherikhin, Koshelev, and Letokhov¹⁰⁴ described a mechanism for obtaining an inversion between the $2p^33p$ and $2p^33s$ levels in neon-like ions. A subsequent experiment on Ca XIX claimed lasing at 600 Å.¹⁰⁵ Vinogradov and coworkers have written a

series of papers^{106,107} refining the theoretical description of the neon-like scheme. According to these papers, the $n = 3$ excited levels are populated initially by electron-impact excitation from the ground state ($2p^6$) of the neon-like ions, which is itself produced in the plasma heated by an optical laser. The population inversion between the $2p^33p$ and $2p^33s$ levels develops because of the large difference between the radiative decay rates. The $2p^33p$ levels cannot decay to the ground state because of electric dipole selection rules, and the $3p-3s$ de-excitation rates are low compared to the high radiative rates out of ($2p^33s$) _{$j=1$} . Because the collisional excitation rate for the ($2p^6$) _{$j=0$} ($2p_{1/2}^5 3p_{1/2}$) _{$j=0$} monopole transition is apparently high, the greatest gain is predicted to occur for the ($2p_{1/2}^5 3p_{1/2}$) _{$j=0$} ($2p_{1/2}^5 3s$) _{$j=1$} transition,¹⁰⁶⁻¹¹³ although transitions from some of the other $2p^33p$ sublevels also experience significant amplification.

Hagelstein,¹⁰⁸ Feldman et al.,^{109,110} and Apruzese and Davis¹¹¹ applied this neon-like excitation scheme to higher-Z ions, in which the wavelengths of the lasing transitions approach the soft x-ray or XUV regime. Dahlbacka et al.,¹¹² Stewart,¹¹³ and Matthews et al.¹¹⁴ have reported experiments on neon-like krypton and selenium using a Z-pinch^{112,113} and a laser-produced plasma,¹¹⁴ respectively. In our previous work,¹¹⁴ the plasma had a steep electron-density distribution, which may have compromised the experiment because of the effects of refraction on the transport of soft x-ray photons through the gain medium. Our most recent efforts suggest an amplifier design intended to minimize the effects of refraction,¹¹⁵ and we have completed experiments to verify the attainment of the desired plasma conditions.

Recently, we reported the first results from a new set of experiments on a neon-like selenium collisional excitation scheme.¹¹⁶ We described the first conclusive demonstration of a macroscopic-sized gain medium that exhibits substantial amplification of soft x-ray transitions. We observed amplification of at least four $2p^33p-2p^33s$ transitions. The highest gain was observed for the $J = 2$ to 1 lines at 206.3 and 209.6 Å. The transition with the highest predicted gain, the $J = 0$ to $J = 1$ line at ~183 Å, has not been unambiguously identified. Moreover, no line in that wavelength region

exhibited significant amplification in comparison with the $J = 2$ to 1 transitions. These conclusions are based on reproducible observations in over 100 separate laser-irradiated target experiments. Our experimental method uses a variety of sophisticated x-ray spectroscopic instruments to measure the time-resolved absolute brightness of the amplified spontaneous emission (ASE) and to demonstrate the nonlinear (nearly exponential) increase of the emission with increasing length of the gain medium. The anisotropy of the ASE was also verified by comparing the relative intensity of amplified line emission from the end of the amplifier (on axis) with that from a line of sight far off that axis.

These experiments were conducted on the Novette laser-target irradiation facility. A schematic representation of the experimental arrangement is shown in Fig. 5-36. The special exploding-foil target design is a product of experiments done on long density-scaling plasmas produced in ICF research.¹¹⁵ The target was composed of a 750-Å layer of selenium, vapor-deposited on one side of a Formvar foil that is 1500 Å thick. The foil was stretched across a gap 0.08 cm wide and nominally 1.1 cm long, and was supported by an aluminum superstructure. The target foil was illuminated by a line focus (0.02×1.12 cm) of $2\omega_0$ light ($\lambda = 0.53 \mu\text{m}$) with a nominal pulse length of 450 ps and a typical irradiance of $5 \times 10^{13} \text{ W/cm}^2$.

Targets were irradiated in two different geometries: single-sided, in which the

vertical center of the foil was hit by only one laser beam, and double-sided, in which opposing laser beams irradiated a common target area. Using the single-sided geometry, we illuminated targets up to 2.4 cm long by displacing the two beams horizontally.

Target x-ray emission in the two axial directions was measured by means of a grazing-incidence spectrograph (GIS) and a transmission-grating streak spectrograph (TGSS-EM). The GIS has high spectral resolution ($\lambda/\Delta\lambda \approx 1800$) and a line-radiation-threshold detection limit of 6 erg/sr. It utilizes a microchannel plate contoured to the Rowland circle for detection of the x-ray spectrum. Gating of the microchannel plate detector by an Auston switch allowed viewing of the spectrum for a limited period of 250 to 750 ps, thereby discriminating against background spectral emissions that occurred late in the experiment. The TGSS-EM has a spectral resolution of $\lambda/\Delta\lambda \approx 200$ to 300, a threshold detection capability of $3.9 \times 10^{-3} \text{ W/Å}$ (radiated into a collection solid angle of $\sim 1.1 \times 10^{-4} \text{ sr}$), and a temporal resolution of 20 to 30 ps. This instrument contains an ellipsoidal mirror for efficient x-ray collection, an x-ray transmission grating with a 2000-Å period for dispersion, and a soft x-ray streak camera for time-resolved detection.

In addition to x-ray emissions along the target axis, we monitored isotropic emissions from the directly illuminated target foil. A second identical GIS viewed the entire target along an axis 77° off the target's longitudinal axis. An x-ray crystal streak spectrograph (XCSS) coupled to another x-ray streak camera viewed the target from 30° off axis and provided time-resolved spectra of the $n = 3$ to $n = 2$ transitions in sodium-like, fluorine-like, and neon-like selenium in the 7- to 9-Å region.

Our measurement strategy was to vary the selenium foil thickness (750 to 3000 Å), laser intensity (1.2×10^{13} to $2.5 \times 10^{14} \text{ W/cm}^2$), and pulse lengths (120 to 750 ps), and search for neon-like selenium transitions with certain characteristics. The parameter combinations we used were determined by computer modeling.¹¹⁵ We sought transitions that were brighter and of shorter duration than nearby spontaneous emission lines, and displayed strong angular anisotropy, as determined by comparison of

Fig. 5-36. Placement of the principal diagnostics used in measuring XUV transitions.

relative line intensities recorded by the two GIS diagnostics, one on and one off axis. Recognizing that the above criteria are only qualitative indications of amplification, we also measured the brightness of the transitions and brightness versus the length of the target. Both of these measurements can be used to roughly determine the amplifier gain. Our experimental method does not depend on any single criterion to demonstrate amplification.

In the course of more than 100 target shots, we found that 750-Å-thick selenium targets, pulse lengths of 450 ps or longer, and irradiances of $\sim 5 \times 10^{13}$ W/cm², in the

case of double-sided irradiation, and $\sim 1 \times 10^{14}$ W/cm², in the case of single-sided irradiation, are optimal for producing amplification. Because of the limited range of parameter space studied, we may not have witnessed the ultimate performance of these amplifiers.

We verified the strong anisotropy of certain lines emitted from these targets. Figure 5-37 illustrates spectral data from one of the time segments monitored by the two GISs. Note the presence of sodium-like selenium resonance transitions (not expected to be amplified) both on and off axis, whereas the strong transitions at 206.3 and

Fig. 5-37. (a) Film data from the time-gated GIS. (b) and (c) Corresponding spectra for less than maximum-gain target shots. A calculated spontaneous-emission spectrum of neon-like selenium is shown upside down in (c). (d) Spectrum obtained from a maximum-gain target 2.2 cm long.

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209.6 Å are seen only on axis. Based on *ab initio* calculations and isoelectronic extrapolation¹¹⁷ from lower Z, these lines are positively identified as the $(2p_{3/2}3p_{3/2})_{J=2} - (2p_{3/2}3s)_{J=1}$ and $(2p_{1/2}3p_{3/2})_{J=2} - (2p_{1/2}3s)_{J=1}$ transitions of neon-like selenium, respectively. With the exception of the $3d_{3/2} - 3p_{3/2}$ transition at 178.6 Å, the sodium-like lines are optically thick and appear brighter off axis because the spectrograph sees a larger surface area of the plasma in that direction. None of the significant lines observed in these spectra can be attributed to diffraction into second order.

Table 5-3 indicates that the wavelengths of the sodium-like transitions obtained in this work agree well with data published previously.^{118,119} Also identified, and listed in Table 5-4, are the strong neon-like lines caused by the transitions $2p^33p - 2p^33s$ and $2p^53d - 2p^53p$ with level degeneracy times oscillator strength (*gf* values) greater than 0.15 (Ref. 117).

Figure 5-37(c) shows a theoretical spontaneous-emission spectrum of neon-like transitions calculated under the

assumption of an optically thin plasma in steady state at $T_e = 1.0$ keV and $n_e = 10^{21}$ cm⁻³. In the experimental spectrum in the same figure, the 206.3- and 209.6-Å lines are much stronger than all of the other nonlasing neon-like transitions, even though some of the nonlasing transitions have higher *gf* values, e.g., the $2p^53d - 2p^53p$ lines at 191.5, 188.4, and 169.2 Å (Ref. 117). The two $J = 2$ to 1 lines dominating the spectrum were observed in about 100 laser shots. Furthermore, their intensity relative to other lines varied strongly, depending on the length of the foil target.

Figure 5-37(d) demonstrates the degree of amplification we can achieve and illustrates other $2p^33p - 2p^33s$ lasing lines at longer wavelengths. In this case the target was subjected to single-sided illumination, but with the beams displaced to give a total amplification length of 2.2 cm. There are few identifiable lines other than the neon-like $J = 2$ to 1 transitions at 206.3, 209.6, and 262.9 Å and the $J = 1$ to 1 transition at 220.2 Å. All of these are predicted to have substantial gain. It is important to note that we do not observe a strong line near 183 Å, the predicted position of the transition with largest calculated gain $[(3p_{1/2}2p_{1/2})_{J=0} - (3p_{1/2}2s)_{J=1}]$. If there were little or no amplification of that line, it would not be seen because of its relatively small *gf* value (0.1). We do see two weak lines in the axial GIS spectrum at 182.0 and 182.4 Å, but their assignment to Se XXV is not conclusive.¹¹⁷

Table 5-3. Present experimental wavelengths and previous results for the sodium-like transitions of selenium and yttrium.

Transition	Selenium (Å)		Yttrium (Å)	
	Present	Previous ^a	Present	Previous ^b
3s-3p _{1/2}	239.21	239.23	196.09	196.19
3s-3p _{3/2}	201.06	201.07	151.00	151.00
3p _{3/2} -3d _{3/2}	178.64	178.64	148.43	148.35
3p _{3/2} -3d _{5/2}	174.19	174.12	141.89	141.92
3p _{1/2} -3d _{3/2}	156.47	156.47	—	120.97

^aTaken from E. Y. Kononov et al., *Phys. Scr.* 19, 328 (1979).

^bTaken from B. Edlen, *Phys. Scr.* 17, 565 (1978).

Table 5-4. Identification of the strongest lines observed in the spectrum of neon-like selenium and yttrium.

Transition	J	J'	Selenium (Å)		Yttrium (Å)	
			Measured	Calculated	Measured	Calculated
2p _{3/2} ^{3s} - 2p _{3/2} ^{3p} _{1/2}	1	2	—	262.85	216.82	217.06
2p _{3/2} ^{3s} - 2p _{3/2} ^{3p} _{3/2}	1	1	220.20	220.22	165.37	165.35
2p _{1/2} ^{3s} - 2p _{1/2} ^{3p} _{3/2}	2	3	211.54	211.61	158.44	158.59
2p _{1/2} ^{3s} - 2p _{1/2} ^{3p} _{1/2}	1	2	209.64	209.74	157.08	157.19
2p _{1/2} ^{3s} - 2p _{1/2} ^{3p} _{3/2}	0	1	208.06	208.22	—	156.72
2p _{3/2} ^{3s} - 2p _{3/2} ^{3p} _{3/2}	1	2	206.30	206.40	154.90	155.03
2p _{3/2} ^{3s} - 2p _{3/2} ^{3d} _{5/2}	2	2	204.50	204.31	168.53	168.49
2p _{3/2} ^{3s} - 2p _{3/2} ^{3d} _{3/2}	2	2	196.95	196.87	148.43	148.55
2p _{3/2} ^{3s} - 2p _{3/2} ^{3d} _{5/2}	3	4	191.45	191.42	156.58	156.73
2p _{3/2} ^{3s} - 2p _{3/2} ^{3d} _{3/2}	1	2	188.93	188.95	154.25	154.43
2p _{1/2} ^{3s} - 2p _{1/2} ^{3d} _{5/2}	2	3	188.36	188.27	153.86	153.96
2p _{1/2} ^{3s} - 2p _{1/2} ^{3d} _{3/2}	2	3	187.96	187.88	153.49	153.60
2p _{1/2} ^{3s} - 2p _{1/2} ^{3d} _{3/2}	1	2	184.30	184.22	—	151.28
2p _{3/2} ^{3p} _{1/2} - 2p _{3/2} ^{3d} _{5/2}	3	3	176.19	176.14	144.32	144.10
2p _{3/2} ^{3p} _{1/2} - 2p _{3/2} ^{3d} _{3/2}	1	1	169.90	169.93	131.90	131.65
2p _{3/2} ^{3p} _{1/2} - 2p _{3/2} ^{3d} _{3/2}	2	3	169.18	169.17	130.70	130.45
2p _{3/2} ^{3p} _{1/2} - 2p _{3/2} ^{3d} _{3/2}	1	0	168.55	168.39	—	125.08
2p _{3/2} ^{3p} _{1/2} - 2p _{3/2} ^{3d} _{3/2}	1	2	167.41	167.58	—	129.19

Time-resolved spectra from the TGSS-EM are shown in Figs. 5-38 and 5-39. Figures 5-38(a) and 5-39(a) show raw data from two representative target shots; Fig. 5-38(a) shows a 1.0-cm-long target displaying moderate amplification and a rich x-ray spectrum, and Fig. 5-39(a) shows a 2.2-cm-long target for which the amplification is intense and the background spectrum is hardly observable because of the filtration required to keep the signal within the dynamic range of the instrument. In Fig. 5-38(b), the temporal profile of an amplified line at 209.6 Å is compared with that of the 201.1-Å transition in the sodium-like ion. Although both lines have a similar rise time (10 to 90% in 250 ps at 209.6 Å and 280 ps at 201.1 Å), the amplified emission is significantly contracted relative to the pulse width of the pump laser, whereas the spontaneous emission at 201.1 Å decays much more slowly ($\tau_{FWHM} = 450$ ps for the pump laser, 340 ps for the 209.6-Å emission, and 950 ps for the 201.1-Å emission). In general, we observed that, with increased amplification, the pulse widths of the lasing lines tend to decrease. Figure 5-39(b) is a spectral profile taken at the time of peak ASE for the 2.2-cm-long target. The emission at 206.3 Å has a peak power of ~360 W. The width of the line, which is source-size broadened, is $\Delta\lambda \approx 1.0$ Å, indicating a source region roughly 200 μm in diameter. As indicated above, the lasing transitions in this case have narrower pulse widths of approximately 160 ps.

The very intense emission shown at 206.3 Å and 209.6 Å in Fig. 5-39 is a strong indication of nonequilibrium spectral conditions in the plasma. If we assume that the line is Doppler-broadened with a width of 0.04 Å ($T_{ion} \approx 400$ eV) and that the source area has a 200- μm diameter, we obtain ~30 to 50 keV for the equivalent radiation temperature of the line at 206.3 Å. In contrast, the brightness temperature observed for the sodium-like nonlasing line at 201.1 Å is only ~0.1 keV. All evidence indicates that the maximum electron temperature in the laser-irradiated medium is approximately 1 keV.¹¹⁵

The final and most quantitative evidence for amplification—the line-energy vs target-length data—are shown in Fig. 5-40. These results come from the axial GIS and are

integrated over the line widths of the 206.3- and 209.6-Å transitions. The data from both the single-sided and double-sided irradiation geometries exhibit a very nonlinear dependence on the length of the amplifier. The output intensity (integrated over length, frequency, and solid angle) emitted by a one-dimensional homogeneous gain medium of length L under steady-state conditions is given by

$$I(L) = \frac{dA}{4\pi R^2} h\omega_0^{1/2} A_{21} \alpha(\omega_0)^{-3/2} L^{-1/2} \times \exp[\alpha(\omega_0)L] \frac{\text{power}}{\text{unit area}} \quad (19)$$

Fig. 5-38. Moderate-gain target shot. (a) Film data from the TGSS-EM. (b) Time history of power radiated by the neon-like $J = 2$ to 1 laser line at 209.6 Å and the nonlasing sodium-like resonance line at 201.1 Å.

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where dA is the detector area, ω_0 is the laser frequency, n_2 is the upper-laser-level population density, A_{21} is the transition

Fig. 5-39. Maximum-gain target shot. (a) Film data from the TGSE-EM. (b) Spectrum of power vs wavelength.

Fig. 5-40. Integrated intensity of $f = 2$ to 1 neon-like selenium transitions vs amplification length for (a) double-sided and (b) single-sided optical laser irradiation conditions.

probability, $\alpha(\omega_0)$ is the line-center gain coefficient, and R is the distance to the detector. Equation (19) is valid when $\alpha(\omega_0) L \geq 2$. Fitting the data from double-sided target shots to Eq. (19) yields a gain coefficient of $5.5 \pm 1.0 \text{ cm}^{-1}$ for both the 206.3- and 209.6-Å lines. Calculations of the length dependence shown by Eq. (19) for two different gains, 5 cm^{-1} and 6 cm^{-1} , are illustrated in Fig. 5-40(a). Although the data obtained with single-sided irradiation demonstrate significant amplification—up to 700 times spontaneous levels—the complexity of the target irradiation geometry precludes a simple application of Eq. (19). The rapid rise of the signals shown in Figs. 5-40(a) and 5-40(b) indicates that the amplifier was still below saturation.

In addition to verifying the gain, we have taken the first tentative steps toward second-generation x-ray laser experiments. Multilayer x-ray mirrors were used at normal incidence to attempt a double-pass amplifier—a precursor to an x-ray cavity oscillator. We have also conducted preliminary experiments to scale the observed amplification to shorter x-ray wavelengths. In particular, we shot yttrium targets to observe the analogous $2p^{5/2}p-2p^{5/2}s$ neon-like transitions. Figure 5-41 shows the spectrum taken with the GIS for an 800-Å yttrium/1500-Å Formvar foil irradiated at $\tau_L \approx 450 \text{ ps}$ and $I_L \approx 1.5 \times 10^{14} \text{ W/cm}^2$. The effects of time-gating the GIS are essential to simplify the spectrum and highlight the neon-like transitions. The wavelengths and the assignments of the observed $2p^{5/2}p-2p^{5/2}s$ and $2p^{5/2}d-2p^{5/2}p$ transitions of neon-like yttrium are also listed in Table 5-4. The $f = 0$ to 1 transitions were not seen, while the $f = 2$ to 1 transitions at 154.9 and

Fig. 5-41. Spectral representation of on-axis data for an yttrium/Formvar target. The top spectrum, taken with the time-gated GIS, shows two $J = 2$ to 1 lasing transitions absent in the bottom spectrum from an ungated GIS.

157.1 Å exhibited strong anisotropy and a pulse-width behavior similar to that seen for the selenium lines. We measured a brightness temperature of ~ 5 keV for the 154.9-Å transition. We conclude that these lines are also being amplified.

By using an optical laser to produce a population inversion in a simple exploding-foil plasma, we have demonstrated substantial amplification of spontaneous emission at soft x-ray wavelengths. The goals for future experiments are to optimize the gain, achieve saturation, measure coherence, spectroscopically characterize the kinetics, develop a cavity oscillator by using partially transmitting, multilayer mirrors, and extend our work with this and other inversion schemes to even shorter wavelengths.

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Exploding Foil Recombination Laser. In addition to experiments using the collisional-excitation approach to a soft

x-ray laser, we tested a recombination-pumped laser scheme. Intense ($I_L \sim 2 \times 10^{14}$ W/cm²) 0.53- μ m light of 100- to 200-ps duration from the Novette laser was used to fully ionize a magnesium exploding-foil target. After the peak of the laser pulse, the plasma cooled rapidly due to expansion, electron conduction, and radiation. In the regime of high electron density ($\sim 10^{20}$ cm⁻³) and low electron temperature (~ 100 eV), three-body recombination preferentially populates the upper levels of the hydrogen-like magnesium ion. The population of the lower levels is depleted by fast radiative decay. This process can result in a population inversion on the $n = 4$ to 3, 130-Å transition. The basic concept for this type of laser was originally proposed by Gudzenko and Shelepin^{120,121} in 1964. Since then, there have been numerous attempts to make an x-ray laser using this scheme,¹²²⁻¹²⁵ including an experiment at LLNL.¹²⁶

In our latest attempt, an improved target design was used to optimize the recombination pumping process. The target consisted of a thin foil of magnesium (1100 Å or 2200 Å) sandwiched between 1000-Å layers of Formvar. The Formvar layers performed several functions. First, they prevented the magnesium from being ionized during the rising portion of the optical-laser pulse

when the intensity was low because the laser had to burn through the Formvar first. Only when the laser pulse was near its peak intensity did the magnesium ionize. This effect maximized the density of the magnesium plasma and minimized the time required to ionize it. Secondly, the mass of the Formvar plasma inhibited the expansion of the magnesium plasma and kept the density high. Finally, the Formvar plasma allowed for electron-conduction cooling of the hotter magnesium plasma. As with the selenium collisionally pumped laser, the thin-foil target minimized refraction losses.

We modeled the target's performance by using the kinetic computer code RATION,¹²⁷ which calculates the ionization balance, level populations, and gain coefficients. The temperature and density evolution for these calculations were provided by the LASNEX¹²⁸ code. Small-signal gain coefficients as high as 3 cm^{-1} were predicted. Because of its relative simplicity compared with other x-ray laser schemes, the hydrogenic recombination laser was expected to provide a sensitive test of our modeling capability.

The experimental setup was similar to that used for the selenium collisionally pumped laser experiment described in the preceding article. A total of 11 shots were fired in the experimental sequence. No clear demonstration of lasing was observed on the $n = 4$ to 3, 130-Å line. However, x-ray spectra consistent with a population inversion were obtained. Figure 5-42 shows the time-resolved intensities of the $n = 4$ to 1 and $n = 3$ to 1 transitions in hydrogenic magnesium. These indicate that a population inversion ratio (n_{4g3}/n_{3g4}) of up to 2 ± 1 occurred after the peak of the laser pulse. The x-ray spectra also indicated that full ionization of the

plasma was never achieved; helium-like resonance lines were observed along with the hydrogen-like resonance lines. Substantial lasing would not be expected in this case.

Experiments are planned to improve the possibility of achieving lasing with this scheme. Current plans call for changing the target construction to correspond more closely with the model. This change involves using a narrow ribbon of magnesium surrounded on all sides with Formvar rather than having Formvar on only two sides. The new configuration would allow better cooling of the plasma after the peak of the laser pulse. In addition, a higher pump-laser intensity with a faster rise will be used to ensure that the magnesium plasma is fully ionized during the laser pulse.

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Resonant Photoexcitation Laser Experiments. This article summarizes work done in 1984 at KMS Fusion, Inc., Ann Arbor, Michigan, in support of resonant photoexcitation (RPE) x-ray laser experiments later conducted on Novette. The results of the Novette experiments are summarized in the article "Design of a Photoresonantly Pumped Fluorine X-Ray Laser," in Section 3.

Following a preliminary series of flash-lamp experiments on Shiva in 1981,^{129,130} we conducted a series of high-resolution spectroscopy experiments at KMS in 1982. Their purpose was to certify to $\pm 0.2 \text{ eV}$ the coincidence of candidate pump lines in a number of RPE schemes. Although the 1982 experiments were performed with high resolution, their interpretation in some cases was compromised by corrections for Doppler shifts and by artificial line broadening, which was due to target repositioning errors and to the need to integrate multiple shots to obtain detectable intensities.

In 1984, we repeated these experiments at KMS with improvements in the experimental method and the instrument sensitivity. Two sets of experiments were conducted. In

Fig. 5-42. Relative intensity as a function of time for the Ly γ $n = 4$ to 1 and Ly β $n = 3$ to 1 transitions. The dashed arrow indicates the time at which $n_{4g3}/n_{3g4} = 2$.

this article, they are called the "winter" series and the "spring" series. These experiments had the following goals for six potential RPE lasing line pairs:

1. To determine the energy coincidence of pump (flashlamp) lines with pumped lines to within 0.15 eV (1 to 2 mÅ).
2. To determine (to within 0.20 eV) the energy noncoincidence of pump lines that would reduce the population inversion by populating the lower state of the lasing transition. Such noncoincidence causes so-called "poison" transitions.
3. To determine Doppler shifts of the pumping spectral lines.
4. To determine the rate of collapse of flashlamp material into the laser channel (interferometric measurements).
5. To verify the production of the lasing spectral lines in an actual gas target when a line-pair coincidence is found and to make a time-resolved observation of the flashlamp being ablated by the coincident lasing line ("pillbox" measurements).

Line-Coincidence Measurements. We made the following improvements over the 1982 series of line-coincidence measurements: (1) the use of intensified spectrographs that eliminated the need for overlaying spectra for different shots and reduced associated problems; (2) the reduction of plasma emission or source size; (3) the use of an x-ray streak camera to measure the Doppler shifts; and (4) the use of spatial resolution to record spectra from adjacent metal comparison targets irradiated together in one shot.

We measured line coincidences to the specified precision using up to four flat-crystal spectrographs per shot. Two of these spectrographs were intensified, and one was time-resolved. See Fig. 5-43 for details of this setup. We used a fifth spectrograph employing a curved KAP crystal as a survey instrument. The two intensified spectrographs and the survey spectrograph viewed the target at right angles to the laser axis in order to minimize Doppler shifts. The time-resolved streak spectrograph viewed the target on axis for some shots and at right angles to the axis on other shots.

We examined six line-pairs for coincidence and their corresponding poison line

Fig. 5-43. Placement of spectrographs in target chamber.

pairs for noncoincidence. We used spectral lines from other elements for wavelength calibration. In all, 29 line-pair combinations were examined in 224 shots, as shown in Table 5-5.

We remeasured spectral lines of manganese and chromium that were previously found to coincide with the fluorine $X\ n = 3$ to 1 and fluorine $IX\ n = 3$ to 1 lines to within 3 to 5 mÅ.¹³¹ Our measurements showed these lines to be in coincidence at 12.643 and 14.458 Å, respectively, within ± 1.5 mÅ. This higher level of accuracy resulted not only from the experimental improvements mentioned above, but also from our extensive theoretical interpretation of boron- and beryllium-like transitions in manganese and chromium reported earlier.¹³¹ The corresponding fluorine $Ly\alpha$ line at 16.806 Å is devoid of manganese poison lines. The fluorine $He\alpha$ line at 14.988 Å is likewise devoid of any poison transitions of chromium.

An additional criterion for a line to be a pump candidate is its inherent strength. The manganese line appeared to have the required intensity, but the chromium line was extremely weak and therefore not considered for further testing as a pumping source.

Doppler Shift Measurements. In the coincidence measurements, we minimized Doppler shifts by recording spectra in directions at right angles to the laser and plasma

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Lines to be pumped	Pump wavelength (Å)	Poison wavelength (Å)	Pumping or poison candidates for wavelength comparison (candidate/substrate)
F (Heβ)	14.458		Fe/Formvar—CaF ₂ /Formvar Mn/Formvar—CaF ₂ /Formvar Mn/Formvar—Fe/Formvar Mn/Formvar—Cr/Formvar Cr/Formvar—Fe/Formvar
F (Heα)		16.806	Mn/Formvar—CaF ₂ /Formvar Mn/Formvar—Ti/Formvar
F (Lyβ)	12.643		Mn/Formvar—CaF ₂ /Formvar Mn/Formvar—Cr/Formvar
F (Lyβ)		14.988	Mn/Formvar—Fe/Formvar
Ne (Heβ)	11.547		Ni/Formvar—Fe/Formvar Ni/Formvar—Mn/Formvar Ni/Formvar—Ne/Be Mn/Formvar—Ne/Be
Ne (Heα)		13.447	Ni/Formvar—Ne/Be Ni/Formvar—Cr/Formvar
Ne (Lyβ)	10.239		Zn/Formvar—Co/Formvar Zn/Formvar—V/Formvar Zn/Formvar—Ne/Be
Ne (Lyα)		12.137	Zn/Formvar—Fe/Formvar Zn/Formvar—Ne/Be
O (Heβ)	18.628		V/Parylene—Formvar Ti/Parylene—Formvar
O (Heα)		21.601	V/Parylene—Formvar V/Parylene—Formvar
O (Lyβ)	16.006		V/Parylene—Formvar V/Parylene—Fe/Formvar
O (Lyα)		18.973	V/Parylene—Formvar Ti/Parylene—Formvar

Table 5-5. Line-coincidence measurements.

blowoff axis. By comparing spectra from front and back side radiation, we found the relative line shifts to be less than 1 mÅ.

We also measured Doppler shifts along the laser axis because they play an important role in determining the feasibility of the resonant photoexcitation scheme. The x-ray streak spectrograph was used to make these measurements. The shifts recorded in the winter series of experiments were of the same order as the distortions introduced by the streak camera's electron optics, and we attempted to correct these distortions. Our final results showed that the manganese lines were blue-shifted by about 2 mÅ over a time interval of 50 ps. Late in time they slowly reverted to a red shift.

In the spring series of experiments, we continued to analyze the Doppler shift of the manganese lines of interest by digital image-processing techniques. We found that the line shifts depended heavily on the accuracy of image registration when the film was digitized. To achieve the required accuracy, we devised a manual procedure employing precise-registration templates and tested it for repeatability. We measured

pincushion distortion effects on a calibration shot and removed them from subsequent images. We then used signal-enhancement and peak-analysis routines to trace out the spectral lines of interest.

Using Formvar substrates 2000 Å thick and silicon nitride (Si₃N₄) substrates 1500 Å thick, we observed blue shifts up to 6 mÅ in 100 ps during the laser pulse. By increasing the Formvar thickness to 8000 Å, we reduced the line shifts typically by 40%. Similarly, by increasing the Si₃N₄ thickness to 2300 Å, we reduced the shifts by 80%.

Interferometric Measurements. We measured the rate at which the flashlamp material expands into the laser channel by using holographic interferometry to record the following intervals: from the time the laser pulse strikes the flashlamp material to the time the plasma strikes the gas-cell wall; and from the latter time to the time the plasma appears beyond this wall.

Two types of targets were developed for these measurements. One was designed to view the vacuum region between the flashlamp material and the gas-cell wall, while the second was designed to view the inside gas-cell surface and the region beyond it.

Double-foil experiments performed during the winter series showed that with laser intensities of 8×10^{14} W/cm² incident on a first foil (with an areal density of 60 μg/cm²), the foil's plasma intruded into the second foil too early in the laser pulse. According to LASNEX simulations, this intrusion in an actual RPE-gain experiment would prevent the optimal buildup of the hydrogen-like state in the fluorine laserant gas by prematurely increasing its temperature above the allowed 3 eV.

In view of these findings, LASNEX predicted acceptable performance for an increased manganese flashlamp thickness of 87 μg/cm² on a Si₃N₄ substrate that also has an increased thickness of 119 μg/cm². Half of the Si₃N₄ target substrates provided by LLNL were coated with the appropriate thickness of manganese. They were then assembled with the uncoated substrates to form double-foil targets. The downstream side of the second foil was available to the holographic probe beam.

In the spring series, we irradiated the targets with 0.53-μm light on the manganese side. However, instead of an f/6 spherical lens, we used an f/1.8 parabolic mirror to focus the laser light onto the target.

The pulse length was maintained nominally at 160 ps, so that a spot-size diameter of 300 μm and laser energy of about 20 J produced a target irradiance of $1.8 \times 10^{14} \text{ W/cm}^2$.

In the spring series, new optics gave four interferograms per shot spaced 125 ps apart. Each was of 20-ps duration. The four frames allow the evolution of the target plasma to be followed. This capability reduces the number of shots required and the problems associated with shot-to-shot variations.

The first frame of a multiframe set is timed with respect to the peak of the irradiating laser pulse ($t \approx 0$). This timing was varied from -90 to -10 ps, providing probe exposures at -90, -50, -10, 35, 75, 115, 160, 200, 240, 285, 325, and 365 ps for the nine target shots.

A preliminary analysis of the data indicated that the added mass of the first foil did, indeed, delay the intrusion of its plasma into the second foil, compared with the earlier double-foil series. There was some indication that the second foil plasma began to appear at the peak of the target-irradiating laser pulse; however, it was moving very slowly with almost no appearance of fringes, indicating a density higher than the critical density. On the basis of these findings, we recommended for Novette experiments that the first foil be made thicker and/or that its distance from the second foil

be increased to make sure the interaction takes place much later in the laser pulse. However, the thicker substrates reduce the ionizing and pump radiation incident on the lasant gas.

Gas-Cell "Pillbox" Measurements. A special target (generically called a "pillbox" and shown in Fig. 5-44) was designed to contain a lasant gas (e.g., fluorine as SF_6). The gas was pumped by the x radiation emitted by the pumping element (e.g., manganese), which was coated onto Formvar and placed at a distance of 50 μm from the gas container.

Time-resolved measurements were made by using the streak spectrograph for two purposes: to observe the absorption of the manganese pump line (at 12.643 Å) by the lasant gas (hydrogen-like fluorine); and to record the fluorescence of the lasant gas lines produced by the x-ray emission of the pump element. Lasant gas pressures were varied from a few mTorr to 300 mTorr.

From a total of eight pillbox shots with the streak spectrograph set up to record absorption of the pump line, only one shot indicated that absorption occurred. Of four shots with the streak spectrograph arranged to record the lasant gas fluorescence, none showed any fluorescence. There are several possible reasons why the pillbox target experiments did not produce the expected results. For one, the metal

Fig. 5-44. Gas-cell target.

manganese and its Formvar substrate may have been too thin. This would lead to an earlier-than-expected interaction of the blowoff plasma with the gas container, which would heat the lasant gas to a higher temperature than required for appropriate level density populations. Another possibility is that the pump radiation flux was insufficient to pump the lasant gas.

In the spring series, we repeated the gas-cell pillbox experiments to examine the effect of the target changes on the spectroscopy of the RPE scheme. With the new double-foil target parameters, one of our principal objectives was to determine whether the manganese flashlamp would pump the fluorine gas. That is, would the beryllium-like manganese line at 12.643 Å preferentially excite the Lyβ line of fluorine in a pillbox target configuration? A secondary objective was to measure the absorption of the manganese pump line.

Target plans for the Novette experiments called for Si₃N₄ to contain the gas and to be used as the substrate for the flashlamp material. However, for the spring pillbox shots, the time available to complete the experiments was not sufficient for a major redesign and fabrication effort. The cell containing fluorine gas was made with Formvar rather than Si₃N₄. The manganese flashlamp was deposited on Si₃N₅ substrates (provided by LLNL) and placed on the pillbox 50 μm from the cell.

The laser-to-target irradiating optics were configured with a cylindrical lens to provide a line focus at the target flashlamp. The line was oriented to give a minimum source size as seen by the streak spectrograph.

The shot matrix included the following variations on nominal target and shot parameters:

1. A pulse length of 160 ps.
 2. A laser intensity of 2×10^{14} W/cm².
 3. Flashlamp thickness corresponding to an areal density of 87 μg/cm².
 4. A Si₃N₄ substrate with an areal density of 119 μg/cm².
 5. An SF₆ gas pressure of 50 Torr.
- Shot iterations included changes in flashlamp and substrate thickness, pulse length, and gas pressure.

We configured the streak spectrograph to view the side window of the pillbox targets

in order to record fluorescence of the helium- and hydrogen-like fluorine lines excited by the manganese flashlamp. To increase the resolution of the spectrometer, we used the second-order x-ray diffraction from a lead stearate crystal dispersion element. Because of the increased areal density of the pump foil compared with that in the winter experiments, the expected fluorescence was definitely (though weakly) observed, well ahead of the plasma emission from the manganese flashlamp. After reconfiguration to first order and reduction of the sweep speed to obtain greater sensitivity, the helium- and hydrogen-like lines were more clearly observed. A preliminary comparison of these line ratios showed the fluorine Lyβ line to be more intense than expected for a thermally emitting plasma. Such intensity indicates possible additional excitation of the line by the manganese pump line.

Two null shots were also performed to rule out other excitation processes. In one, the SF₆ gas was not introduced into the cell. In another more significant test, the manganese flashlamp was replaced with nickel. In both tests, the fluorine Lyβ line intensity was considerably reduced. These results support the inference that the fluorine Lyβ line was enhanced in the regular shots.

Attempts to measure the absorption of the manganese pump line by the fluorine ions in the pillbox were unsuccessful due to equipment malfunctions. The CsI photocathode sensitivity apparently was poisoned by high humidity.

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Diagnostic Development

Introduction

In developing diagnostic techniques for the LEAD Program, we traditionally maintain a strong commitment to advancing

technologies. In 1984, we strengthened our technical base with highly advanced facilities to aid the development of new measurement techniques. These additions to our laboratory are described in the article entitled "Facilities."

Our short-pulse laser facility provides precise measurement and testing capabilities in the 1- to 10-ps range. This laser will be used in the development of time-resolved instrumentation. A new microfabrication laboratory allows us to apply the rapidly advancing technology of microelectronics to x-ray optics. The laboratory is the only one of its kind at LLNL that is accessible to university and industrial research teams; thus it allows collaboration in this area of research. We also acquired a computer-controlled facility for reconstructing optical images, which allows rapid processing of coded image data. Finally, in cooperation with the MIT Center for Space Research, we developed a new charge-coupled x-ray camera system, which improves our capability to characterize reflected and transmitted x rays. This diagnostic is described in the article "X-Ray CCD Camera."

Much of our effort in 1984 was focused on developing new measurement capabilities for the first experiments on the Nova laser. Early compression experiments will benefit from our work in devising the measurement techniques described in the article on "Advanced Neutron Measurements." These techniques are designed to:

- Characterize the *pr* of the compressed fuel.
- Image the compressed core.
- Provide time-resolved information about the thermonuclear burn.

To accommodate experiments on the stability of compressed laser fusion targets, we are continuing work on appropriate diagnostics: the x-ray streak camera, described in the article "Improvements in Streak Camera Electronics," and the x-ray framing camera, described in "Development of an X-Ray Framing Camera." The goals of this work are to achieve more stable, absolutely timed streak measurements and to obtain sequences of two-dimensional images of dynamic implosion processes.

A satisfying example of the return on our investment in measurement technology is

the TGSS-EM. This instrument is described in the next article, "TGSS-EM: A Time-Resolved Imaging Spectrometer for the Measurement of Amplified Spontaneous Emission at Soft X-Ray Wavelengths." Consolidating our experience with grazing-incidence reflection optics, transmission x-ray optics, and x-ray streak cameras, the TGSS-EM provided the first time-resolved measurement of ASE at soft x-ray wavelengths in a series of $\sim$ -ray laser experiments on the Novette facility.

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TGSS-EM: A Time-Resolved Imaging Spectrometer for the Measurement of Amplified Spontaneous Emission at Soft X-Ray Wavelengths

In some of its most significant experiments of 1984, the LEAD Program achieved the production of amplified spontaneous emission (ASE) at soft x-ray wavelengths in laser-produced plasmas. Experiments of this nature are sufficiently different from conventional laser-plasma experiments to cause us to rethink our philosophy of x-ray measurement. To meet the unique challenge posed by ASE measurement, it is necessary to redesign instrumentation and to devise new techniques.

In particular, the distribution of ASE in angle, wavelength, time, and intensity have significant impact on diagnostic design and measurement philosophy. Because of the strong anisotropy of ASE, target-instrument coalignment becomes critically important. Likewise, anisotropy makes the solid angle of collection a signal-to-noise issue and suggests divergence measurement as a method of gain verification. The spectral characteristics of ASE are similarly consequential. The extremely narrow width of amplified lines puts a high premium on spectral resolution and makes instrument resolution a signal-to-noise and threshold-detection issue. The typically low energy of the ASE photons (relative to the mean electron energy in the parent plasma (i.e., $h\nu \ll kT_e$)) demands instrumental discrimination against higher spectral orders. Along with spectral

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discrimination, flexibility is an important requirement for future experiments. Our instrumentation must have a broad and easily adjustable spectral range to investigate a variety of population inversion schemes over a wide band of wavelengths. Finally, the short duration of the ASE demands time-resolved measurement, and the strong nonlinearity of these emissions requires measurement capability with a broad dynamic range.

Instrument Description. From these measurement issues evolved a transmission-grating streak spectrograph with an ellipsoidal mirror (TGSS-EM). The primary components of the TGSS-EM are its mirror, a transmission grating, and a streak camera. This instrument was used as a principal diagnostic on the recently completed Novette x-ray laser experiments. (See the article "Collisional-Excitation Laser Scheme," earlier in this section for more information on its use.) A schematic representation of this time-resolved spectrograph is presented in Fig. 5-45. Its ellipsoidal mirror reflects the output of the x-ray laser target (placed at

one focus) onto the photocathode of the x-ray streak camera (placed at the complementary focus). An x-ray transmission grating, placed in the path of the focused x-ray emission, disperses the spectrum along the slit of the streak camera. The streak camera can be moved along the dispersion direction in the detection plane to sample the desired portion of the spectrum. Table 5-6 summarizes the TGSS-EM performance specifications.

The ellipsoidal mirror is an annular section of an ellipsoid of revolution with major and minor axes of 122 and 8.69 cm, respectively, and an interfocus distance of 121.7 cm. The mean grazing-incidence reflection angle used in the x-ray laser experiment was $\sim 4.1^\circ$, providing a cutoff for x rays above ~ 700 eV. The mirror substrate is aluminum plated with electroless nickel. The figure was diamond turned to a slope-error specification of $<100 \mu\text{rad}$ and polished for 90 h to a low-scatter finish (with surface roughness $<25 \text{ \AA rms}$). The mirror was characterized by means of carbon K α x rays. It exhibited spatial resolution of 50

Fig. 5-45. Layout of the time-resolved imaging spectrometer, TGSS-EM.

to 70 μm , an image magnification of 1.35, and a reflectivity greater than 60%.

The dispersion element of the TGSS-EM is a free-standing gold, transmission grating. The grating period is 2000 $\text{\AA}$ with $\sim 1:1$ line-to-space ratio. The linear dispersion of the grating at the streak-camera detection plane is $\sim 3.2 \text{ \AA}/\text{mm}$. The grating is the limiting aperture of the TGSS-EM. Its active area is a 6-mm-diam circle, which provides a collection solid angle of $\sim 1.1 \times 10^{-4} \text{ sr}$.

The streak camera is a standard LLNL soft x-ray camera with a demountable photocathode. The cathode used in the x-ray laser experiments had slit dimensions of $500 \mu\text{m} \times 17 \text{ mm}$. It was 250- $\text{\AA}$ aluminum supported by 800- $\text{\AA}$ -thick Formvar. The spectral range covered by the 17-mm cathode slit dimension was 55 $\text{\AA}$, adjustable by simple camera movement over a wavelength band from 0 to 260 $\text{\AA}$. The 500- μm slit dimension allowed the entire output aperture of the x-ray laser target to be imaged onto the photocathode. As a result, there was a slight loss in temporal resolution from 20 to $\sim 40 \text{ ps}$.

Alignment Procedure. Alignment of the TGSS-EM with the x-ray laser target is a critical aspect of ASE measurements. The alignment procedure in the Novette x-ray laser experiments included three main operations:

1. Bench alignment of the target (Fig. 5-46).

In this procedure, a retroreflecting reference mirror is placed perpendicular to the laser-material channel. A ring

interferometer is then installed with a high-magnification telescope at one view port and an autoreflecting telescope at the other. Once the interferometer is balanced, the x-ray target is inserted into the interferometer's beam path. The attitude of the target is adjusted until the images of the two opposing ends of the laser channel overlap when viewed with the high-magnification telescope. The reference mirror is then adjusted to achieve retroreflection with the autoreflecting telescope.

2. Bench alignment of the TGSS-EM optical ensemble (Fig. 5-47). In this procedure, the target-chamber geometry is simulated by establishing a preferred line of sight between an autoreflecting telescope and a point light source (mock target). The axis of the ellipsoidal mirror is aligned parallel to this line of sight, and a reference mirror with a centering cross hair is aligned perpendicular to the same line of sight. The diffraction grating is centered on the image of the beam that originates from the point source and is reflected from the mirror.

3. Chamber alignment of the target-TGSS-EM system (Fig. 5-48). In this procedure, the target and the TGSS-EM ensemble

Table 5-6. Performance specifications for the TGSS-EM.

Temporal resolution	$\sim 40 \text{ ps}$	(source-size limited)
Spectral resolution	$\sim 1 \text{ \AA}$	(source-size limited)
Collection solid angle	$\sim 10^{-4} \text{ sr}$	(grating-aperture limited)
Dynamic range	$\geq 10^3$	(streak-camera limited)
Maximum sensitivity	50 W/sr	(photocathode dependent)
Spectral range	$\Delta\lambda_r \sim 55 \text{ \AA}$	(adjustable over range 0 to 260 $\text{\AA}$)

Fig. 5-46. Bench alignment of the x-ray laser target.

Fig. 5-47. Alignment of the TGSS-EM ensemble.

Fig. 5-48. Chamber alignment of the target and the TGSS-EM system.

are aligned along the preferred line of sight between the autoreflecting telescope and the center of the target chamber. The laser-material channel is first aligned along the preferred line of sight by means of a retroreflecting mirror. The TGSS-EM ensemble is then situated on the same line of sight by retroreflection in its reference mirror and centering on its cross hair. The location of the streak camera's photocathode along the dispersion line is verified by rotating the TGSS-EM

ensemble about its vertical axis (preset perpendicular to the ellipsoidal axis) and observing whether the image of the photocathode is centered on the cross hair of the reference mirror.

By using these methods and additional support procedures, we were able to align the spectrometer and the target laser channel with an accuracy better than ± 1 mrad.

TGSS-EM Data. A typical TGSS-EM data record is shown in Fig. 5-49. It was produced from a selenium x-ray laser target

illuminated by Novette. The strong ASE from neon-like selenium at 206 and 209 Å is clearly visible. A strong continuum emission occurs later at ~2 ns. Without temporal discrimination, the late-time continuum emission would have interfered with the detection and identification of the amplified lines. The record also shows two fiducials at the long-wavelength end of the spectrum. The bright edge is the second-order K-edge of a thin beryllium filter. It serves as a wavelength marker at 222 Å. The dark line adjacent to the beryllium edge is the shadow of a thin wire used as a spatial resolution check.

Two more data records are shown in Figs. 5-38 and 5-39 earlier in this section. Figure 5-38(a) shows the spectrum from a 1-cm-long selenium x-ray laser target. Easily seen are 206-Å, neon-like ASE lines (marked A), a sodium-like line at 201 Å (marked B), and other spectral features. Figures 5-38(a) and 5-38(b) show the contrast between the temporal shape of the ASE and the spectrum. The FWHM of the 209-Å line is ~250 ps, significantly shorter than the 450-ps Novette laser pulse, whereas the width of the 201-Å line is about 1 ns. Figure 5-39 shows the heavily filtered spectrum of a 2-cm-long selenium x-ray laser target. The photocathode was filtered to keep the ASE lines within the dynamic range of the streak camera. The peak power at 206 Å was measured at slightly greater than 3 MW/sr.

The TGSS-EM is a synthesis of our capabilities in time-resolved x-ray detection, grazing-incidence reflection optics, and microfabrication. It shows a generous return on our long-term investment in x-ray measurement technology.

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X-Ray CCD Camera

A balanced development program in x-ray optics must include goals for improving x-ray detection. Newly developed charge-coupled devices (CCDs) make it possible to

Fig. 5-49. Typical TGSS-EM data record from a selenium x-ray laser target.

detect x rays electronically with low noise, high sensitivity, good spatial resolution, real-time readout, and extreme positional stability. Such devices provide a significant advantage over x-ray film, with its combined problems of quality control, high grain noise, limited sensitivity, and slow data turnaround.

We have developed a direct x-ray-imaging CCD camera designed to test and characterize reflection and transmission x-ray optics. It covers a spectral range from ~0.5 to >15 keV, and its high sensitivity allows the use of low-brightness laboratory x-ray sources. The front-illuminated virtual-phase CCD, made by Texas Instruments, Inc., has a 390- × 584-pixel array. Since each pixel is 22 μm square, the total active detector area is 8.6 × 12.8 mm. The pixel capacity is 4 × 10⁵ electrons. The CCD operates in a temperature range from -60 to -100°C. It has an equivalent noise of ~30 electrons (rms) and exhibits a dark current less than 0.05 nA/cm².

The camera system is a stand-alone unit with three components: the camera head, the front-end electronics, and the main computer. The camera head contains a CCD

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Fig. 5-50. Schematic layout of the CCD camera system including the camera head, the front-end control electronics, and the main computer with peripherals.

chip, which is the active imaging element, and the CCD pre-amplifier board, all in a vacuum environment cooled by a liquid nitrogen reservoir. The Dewar design of the cooling apparatus allows for universal orientation of the head. The cooling apparatus can hold the temperature below -100°C

for more than 16 h. The front-end electronic components are the heart of the camera system. They monitor and control the CCD temperature, drive the clocks that shift the charge out of the CCD, perform low-noise amplification, do analog signal processing, convert the output signal from analog to

digital, and transmit the digital data to the main computer. Eight circuit boards make up the front-end electronics. They are listed below and represented schematically in Fig. 5-50.

1. 1802 Cosmac microprocessor.
2. Cosmac memory.
3. Digital-to-analog converter.
4. Temperature control.
5. Programmable sequencer.
6. High-speed clock driver.
7. Analog signal processor.
8. High-speed serial link.

The main computer, a Pacific Microcomputer Model 68000, is also represented with its peripherals in Fig. 5-50. Image data transferred over the high-speed serial link can be analyzed by this computer, displayed by the AED 512 color graphics terminal, and stored on the 20-megabyte Winchester disk system or on 1/2-in. magnetic tape by means of the Cypher tape drive.

Major improvements to the CCD camera are under way. By early 1986, we expect to extend CCD sensitivity to longer x-ray wavelengths (carbon $K\alpha \approx 280$ eV). Furthermore, a new camera is being built with an 800×800 -pixel array and pixel dimensions improved to $15 \mu\text{m}$ square. Finally, we are designing a compact, thermoelectrically cooled camera head that will be easier to handle and require less cooldown time.

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Development of an X-Ray Framing Camera

This article describes the development of a framing camera capable of recording an x-ray image with a time resolution of ~ 100 ps and a spatial resolution of $\sim 25 \mu\text{m}$. The potential applications of such an instrument have been described previously.¹³² The x-ray framing camera (XFC) will allow the direct observation of instabilities and asymmetries that can degrade the performance of fuel-pellet implosions. In a typical experiment, the imploding fuel pellet will be backlit by an intense laser-induced

x-ray source, and the framing camera will record a radiograph of the implosion event through an x-ray microscope. Since the frame duration of 100 ps is much shorter than the duration of the implosion (> 1 ns), the image will be a "snapshot" of the pellet at a particular stage of compression. Hence the XFC should be an ideal tool for diagnosing the dynamics of the implosion process.

Performance Requirements. If the XFC is to be a useful experimental instrument, it must meet severe requirements imposed by the laser fusion processes that are to be studied. In a fuel-pellet implosion, the radial velocities can be as great as $\sim 3 \times 10^7$ cm/s. Nonuniformities of the pellet surface are expected to grow because of Rayleigh-Taylor instabilities, with maximum growth at a wavelength approximately equal to the shell thickness. Since it is important to resolve these fine-scale nonuniformities, the spatial resolution of the XFC should be several microns at the source. In order that the image blur associated with the rapid implosion not exceed a spatial resolution of several microns, the frame duration should be no longer than 100 ps.

The limited intensity of available x-ray backlighters imposes further constraints. For instance, an aluminum backlighter (1.4 keV) used in conjunction with an $8 \times$ Kirkpatrick-Baez (K-B) x-ray microscope (viewing a solid angle of $\sim 10^{-7}$ sr) can provide on the order of 10^8 x-ray photons to the input of the camera during a 100-ps interval. If an image consists of $\sim 100 \times 100 = 10^4$ pixels, an exposure of 100 photons per pixel (which implies a statistical noise level of 10%) would require the XFC to have a detection efficiency of at least 1%. In this regime of signal-limited imaging, it is essential that the XFC be able to detect single x rays with high efficiency.

Another important requirement for the XFC is an adequate signal-extinction ratio—the ratio of the detection efficiency in the "open-shutter" condition to that in the "closed-shutter" condition. Any backlighter illumination or target self-emission recorded outside the frame interval must not reduce the dynamic range of the framed image. If significant x-ray emission exists for 1 ns and a dynamic range of 100:1 is desired, the extinction ratio must be better than 1000:1.

Finally, the field of view of the XFC must be large enough to conveniently image reactor-size fuel pellets of up to 1 mm in diameter. An x-ray optic well matched to the XFC is the 8× K-B microscope. This component satisfies all optical requirements and is simple to operate. The 8× magnification factor requires that the field of view of the camera be $\sim 1 \times 1$ cm. The requirement that the XFC have a spatial resolution of several microns at the source translates to a spatial resolution at the image plane of at least 50 μm . Thus the 1-cm² field of view implies an image array of $\sim 200 \times 200$ pixels.

In view of the considerations presented above, the following are realistic performance requirements for the XFC:

- Temporal resolution: ~ 100 ps.
- Spatial resolution: $< 50 \mu\text{m}$.
- Dynamic range: $> 100:1$.
- Detection efficiency: $> 1\%$ in the energy range 1 to 5 keV.
- Detection sensitivity: single x rays.
- Extinction ratio: $> 1000:1$.
- Image size: $\sim 1 \times 1$ cm.

Framing Camera Design. The design of the camera we are fabricating is shown in Fig. 5-51. The x-ray image is incident on a transmission photocathode composed of CsI deposited on a $\sim 25\text{-}\mu\text{m}$ beryllium window. The photocathode is incorporated into the upper half of a strip transmission line. The conducting strip is composed of a film of gold ink deposited on a ceramic substrate. The width of the strip is 25 mm, and the

transmission-line gap is 1 mm; these dimensions yield a characteristic impedance of 15 Ω . In the closed-shutter condition, a reverse-dc bias of 500 V prevents the photoelectrons from leaving the surface of the photocathode. To record an image frame, a 5-kV gate pulse is generated by a GaAs photoconductive switch and propagated across the photocathode. The GaAs switch is activated with a short light pulse of moderate intensity (~ 0.25 mJ) coupled in through a fiber-optic light guide. The width of the gate pulse is twice the length of the charge line. The rise time of the gate pulse is largely determined by the temporal width of the trigger light pulse and by broadening due to transmission-line dispersion. The gate pulse travels along the transmission line and is eventually dissipated in a thick-film terminating resistor. The width of the gap underneath the resistor is decreased to maintain a constant line impedance, which minimizes reflections.

Throughout the time that the gate pulse travels across the photocathode, photoelectrons are accelerated toward the ground plane. A thin microchannel plate (MCP) is built into the ground plane directly below the photocathode. Computer simulations indicate that the electrons are deflected in the transmission-line gap by magnetic fields associated with the propagating electromagnetic wave. The angle of the channels (3.8° with respect to normal) is chosen to allow

Fig. 5-51. The x-ray framing camera design. Framing is achieved by gating a photocathode built into the strip transmission line.

most of the electrons to pass through the MCP. Meanwhile the background signal of x rays that are transmitted through the nearly transparent photocathode impinges on the MCP at normal incidence, and the x rays are absorbed without being detected.

The photoelectron image is directly detected by an RCA SID-501 CCD placed within 1 mm of the bottom of the MCP. The image area consists of a 320×512 array of $30\text{-}\mu\text{m}$ -square pixels. Because the CCD is thinned for back-side optical illumination, it is particularly suitable for direct electron bombardment. Previous studies of the direct electron bombardment of CCDs have demonstrated that the devices are very linear and have high sensitivity for detecting electrons with sufficient energy to penetrate the dead layer.^{133,134}

The performance of this XFC was studied by means of a computer simulation that incorporated all the basic camera components. Monte Carlo calculations of individual electron trajectories were used to generate simulated images. The calculations predict that the width of the point-spread function will be within a single CCD pixel, i.e., $<30\text{ }\mu\text{m}$. The simulation also shows that XFC efficiency is limited to the quantum efficiency of the CsI photocathode, which is 1 to 5%.

Evaluation Facilities. Conceptually, the XFC is a simple instrument. However, there are several critical aspects of the design that have not yet been tested. One assumption crucial to the successful operation of the camera is that a high-voltage pulse composed of purely transverse electromagnetic (TEM) modes can be generated and propagated a short distance along the transmission line, even though the wave contains frequency components well above the transverse electric (TE) and transverse magnetic (TM) cutoff frequencies. Another important consideration is the response of the CCD to direct bombardment by relatively low-energy (4.5-keV) electrons. In order to optimize the performance of the XFC, it is necessary to investigate the behavior of the transmission line and the CCD thoroughly. We are setting up two laboratory facilities to perform these evaluations.

The transmission line will be characterized on an electrooptical sampling system, based on a technique originally developed by Valdmanis et al.¹³⁵ This sampling system

is shown schematically in Fig. 5-52. A small birefringent crystal (LiTaO_3) is suspended inside the transmission line gap. A light pulse lasting 3 to 5 ps (available from the short-pulse laser facility) is used to trigger the photoconductive switch and probe the crystal after a variable delay. As the linearly polarized beam passes through the crystal, the Pockels effect rotates the plane of polarization by an amount proportional to the instantaneous strength of the electric field. By analyzing the polarization state of the transmitted beam, the electric field strength can be measured with a temporal resolution equal to the width of the light pulse ($<5\text{ ps}$) and with a spatial resolution equal to the size of the crystal ($\sim 0.5\text{ mm}$). The delay of the probe beam and the position of the crystal are controlled by a computer that can rapidly map temporal and spatial profiles of the propagating voltage pulse.

The response of the CCD to direct electron bombardment will be studied using a scanning electron microscope (SEM) as the electron source. We are installing digital-beam-scanning and high-speed beam-blanking capabilities on a Nanolab 2100 SEM. The microscope is being modified to allow for the mounting and operation of a standard CCD package inside the sample chamber. This facility will make it possible to bombard the CCD with an electron beam of variable flux and duration with energy varying from 750 to $30\,000\text{ eV}$ and a spatial resolution of better than $1\text{ }\mu\text{m}$. The system will be used to characterize important effects such as the magnitude and uniformity of the dead layer and the extent of signal spreading inside the CCD.

Fig. 5-52. The configuration of the electrooptical sampling system, capable of measuring voltage waveforms inside the transmission line with picosecond resolution.

These facilities will allow us to study the critical components of the XFC separately. Such investigations will provide a thorough understanding of XFC operation, which is essential to the development of an effective and versatile instrument.

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Advanced Neutron Measurements

Advanced Neutron Diagnostics for the Nova Laser Facility. Implosion experiments performed on Nova are expected to produce an increased yield of thermonuclear neutrons compared with that of earlier ICF experiments. This yield will make feasible a number of neutron-based measurements heretofore not possible.

Laser fusion neutron diagnostics can be divided into two categories: invasive and noninvasive. Invasive techniques require the placement of a tracer material in an interesting region of the target to be activated by the thermonuclear neutrons. The degree of activation indicates the location or the areal density (ρr) of the tracer material at the time of neutron emission, which in turn is indicative of some aspect of target performance. This technique is attractive because it is passive and can be made quite sensitive. Disadvantages stem from the requirement for specially fabricated targets and possible effects on target performance. The compressed ρr of various target components and values related to pusher-fuel mixing are the quantities that can be expected to be measured by these techniques.

Noninvasive techniques involve the energy, spatial, or temporal analysis of the neutrons emitted from the target. These techniques take advantage of the large mean free path of the neutrons, which allows the easy transport of information out of the targets. Noninvasive techniques require no specially prepared targets, and the measurement does not affect target performance. The target parameters measured include the total yield, compressed fuel ρr , burn history, core symmetry, ion temperature, and interval time between the laser and neutron peaks.

After examining a host of diagnostic options from both categories for Nova, we have decided to pursue the following ideas:

1. We are planning to develop two diagnostics—an invasive radiochemical technique and a noninvasive measurement of the neutron energy spectrum to obtain the compressed fuel ρr . These diagnostics will apply to a range of Nova target designs. They will have an approximate ρr -yield threshold of 10^9 neutrons·g/cm². The next article on the facing page, "Measurement of the Fuel ρr of Nova Targets," discusses these techniques.
2. We are also designing a streak-camera-based instrument capable of measuring neutron production as a function of time with a resolution of 15 to 20 ps. This instrument will also be able to measure the ion temperature and the interval time between the laser and neutron peaks. We expect the detection threshold for such an instrument to be in the range of 10^9 to 10^{10} neutrons. These measurements will be made with an optical streak camera coupled to a stack of uranium foils interleaved between thin layers of plastic scintillator. It is crucial that the rise time of the scintillator be less than the desired resolution of 15 to 20 ps. The article "Fast-Scintillator Measurements" describes the rise-time measurements of several candidate scintillator materials as well as the instrument itself.
3. With yields of 10^{11} and greater, we can begin to consider neutron imaging. Imaging the thermonuclear-neutron source region with a spatial resolution of 5 to 10 μ m will require the fabrication of a very high-aspect-ratio aperture. Such a structure will have dimensions or tolerances on the order of the desired resolution in one direction and dimensions comparable to a significant fraction of the neutron mean free path in the other direction. Furthermore, to operate at these yields, the aperture must have a larger collection area than a simple pinhole, the only aperture considered thus far. We are performing simulations to examine a number of devices and techniques, including high-aspect-ratio zone plates, uniformly redundant arrays, and penumbral imaging. This last technique is described in the article "Penumbral Imaging for Neutrons."

Author: S. M. Lane

Measurement of the Fuel (ρr) of Nova Targets. In 1984, we investigated two methods for determining the compressed fuel (ρr) of Nova targets:

- Neutron activation of a tracer material placed in a target.
- Measurement of the neutron energy spectrum.

The neutron-activation method requires that a special tracer material be placed in a relevant region of the target. This is an established radiochemical technique that has the advantage of being passive; that is, it is insensitive to target-produced x rays, electrical noise, and gamma rays. Because the amount of tracer can be adjusted, this method can be made quite sensitive.

The neutron-energy method takes advantage of the fact that the predominant interactions of the fusion neutrons before they leave the target are with the fuel. Thus, a direct determination of the fuel (ρr) can be made by measuring the neutron energy spectrum. This technique is not straightforward because of the low fuel areal density ($\rho r < 300 \text{ mg/cm}^2$) expected from Nova targets where only a few percent of the neutrons interact with the fuel. Therefore they must be measured in the presence of the large number of noninteracting neutrons.

We studied a number of methods of measuring the neutron spectrum and concluded that a time-of-flight technique employing a detector made of a large number of independent cells was the most attractive. Alternative techniques we considered include:

1. Threshold detectors.
2. The use of plastic nuclear track detectors or a magnetic analyzer to measure the energy and direction of charged recoil particles produced when neutrons impinge on a radiator material.
3. Time-of-flight measurements by means of a neutron streak camera, a neutron-sensitive gated transmission line, gated microchannel-plate detectors, or a single large scintillator detector.

These techniques were eliminated because of low sensitivity, low energy resolution, or complex data reduction requirements (in the case of track detectors).

Neutron Activation. Neutron activation of silicon in the pusher is a standard LEAD Program (ρr) diagnostic.¹³⁶⁻¹³⁸ In general, our neutron-activation techniques employ a tracer element that is activated by fusion

neutrons. The tracer can be mixed uniformly with the fuel, placed in a thin layer between the fuel and the pusher, or (least desirably) mixed with the pusher. The latter two tracer locations yield an indirect measurement of the fuel (ρr) and require computer modeling support. The number of tracer atoms that are activated is proportional to the neutron yield (which must be measured independently) and to the burn-averaged (ρr) of the tracer. After the implosion, a known fraction of the activated atoms must be collected and then counted. Counting can be done by a detector system in the chamber or in a remote, shielded enclosure. The fraction of collected debris is determined by including a known amount of a second radioactive tracer in the target.

Using an excessive amount of the tracers can cause problems such as hydrodynamic instabilities, fuel preheating, or quenching of the burn. The degree of these problems can be partially quantified by measuring neutron yield as a function of tracer concentration. Another difficulty is the need to develop techniques for the collection of tracer debris and the determination of the fraction collected.

There are several advantages to neutron-activation techniques. First, they can be comparatively sensitive. For instance, a ^{80}Kr tracer, enriched to 90% in an amount equal to 0.1 at.% of the fuel and mixed uniformly with the fuel, requires a yield of only 2×10^8 neutrons to measure a (ρr) of 0.05 g/cm^2 when 1% of the ^{80}Kr is collected following the implosion. ^{80}Kr is an attractive tracer candidate because of its anticipated solubility in liquid DT. We expect that a uniform mixture of ^{80}Kr and DT can be achieved in directly driven cryogenic targets containing liquid fuel layers. If neutron yields are high enough, it is possible that different tracers could be placed in separate regions of the target to obtain simultaneous fuel and pusher (ρr) measurements or even mix measurements.

Neutron Energy Spectrum Measurements. The dominant interactions of fusion neutrons with the target are elastic scattering and (n,n) reactions with the DT fuel. These interactions cause the neutrons to lose some of their energy. An example of the resultant energy distribution is given in Fig. 5-53. The intensities of the fusion neutrons in the figure are normalized according to the yields given in Table 5-7. The total intensity from elastic scattering and (n,n) reactions in the

Fig. 5-53. Neutron spectral distribution in an equimolar fuel mixture.

Fusion reaction	Rate (relative)	Weight
DT	1.00	1
DD	4.14×10^{-3}	1
TT	2.87×10^{-3}	2

Table 5-7. Characteristics of fusion reactions assumed in spectral calculations. (Mean plasma temperature is 5 keV.)

Neutron reaction in fuel	Cross section at $E_n = 14$ MeV (b)	Weight
D(n,n)D	0.638	1
D(n,2n)D	0.177	2
T(n,n)T	0.932	1
T(n,2n)T	0.048	2

fuel was assumed to be 1% of the DT neutron yield. The fraction of neutrons f interacting in this manner is given by

$$f = 0.22 \langle \rho \rangle (\text{g/cm}^2) \quad (20)$$

for a target having a 1:1 DT fuel mixture. Nova targets are expected to achieve $\langle \rho r \rangle$ values sufficient to cause 1 to 10% of the neutrons to interact. If the energy spectrum of neutrons emitted from the target were measured, the fraction of neutrons interacting with the fuel, and thus the fuel $\langle \rho r \rangle$, could be estimated.

The tritium-tritium TT reaction is a potential source of interfering neutrons. (The TT yield is roughly 0.6% of the DT yield at an ion temperature of 5 keV.) With an optimal fuel mixture—about 90% D to 10% T—the total yield will be reduced by about a factor of 3, whereas the TT yield will be reduced by a factor of 25.

Besides the TT yield, there are other potential sources of interference. One consists of scattered neutrons resulting from interactions with structures other than the fuel. Structures inside the target chamber, such as the target manipulator and the cryogenic shroud, are capable of scattering vastly more neutrons than the fuel. Although these neutrons may be partially shielded or discriminated against on the basis of time of arrival at the detector, it is critical that they be well characterized to avoid confusion with fuel-scattered neutrons. Because of the complex geometry of the target chamber and associated hardware, the non-fuel-scattered neutrons must be characterized by experiment as well as by computer simulations. Such experiments could be performed by using targets with high yields and very low fuel densities, such as exploding-pusher targets. Alternatively, a small neutron generator could be used to produce neutrons at the chamber center.

Another type of interference is caused by full-energy neutrons—those that have not interacted with the target or have interacted but have lost very little energy. The detector must have sufficient energy resolution to distinguish between the large quantity of these neutrons and the relatively few neutrons that have lost a substantial amount of energy in the fuel. For $f = 1\%$, Doppler broadening alone limits the observable portion of the spectrum to below 13.5 MeV, if we assume a Gaussian line shape having a width appropriate for a 5-keV ion temperature. Instrument broadening will reduce this region even further. Although computer predictions confirm that the Gaussian line shape is a good approximation even at 0.001 of the peak value, non-Gaussian shapes may arise in the course of experiments.

Because the emission of neutrons from the source is essentially a delta function in time and space, the energy spectrum can be obtained by measuring the relative time of arrival of neutrons at a detector. Such a detector might usefully be placed anywhere from 10 cm to 10 m from the target. The closer detectors are placed to the target, the shorter is the required resolving time, and the more sophisticated the detection system must be. For reference, the arrival times of

14-, 13-, 8-, and 3-MeV neutrons at a detector 1.0 m from the target are 19.4, 20.1, 25.7, and 41.9 ns, respectively.

We selected a 1.0 m flight path inside the 5-m-diam Nova target chamber because a short target-to-detector distance makes it possible to discriminate against some non-fuel-scattered neutrons on the basis of arrival time. For instance, a detector placed about midway between the target and the chamber wall will detect a 3-MeV neutron before a 14-MeV neutron has time to bounce off the nearest wall and return to the detector.

A major difficulty with a time-of-flight system is detector saturation. The detector must not be saturated by the burst of 14.1-MeV neutrons that arrives ahead of the slower neutrons scattered in the fuel. The detector system used in this application must be able to recover rapidly or remain active immediately following the 14-MeV burst.

By using a detector made of a large number of identical cells, we can increase the ratio of downshifted to full-energy neutrons observed by each cell, even though the ratio emitted by the target may be very low. This ratio increases with the number of cells to the point at which, on average, a single cell detects either 0 or 1 downshifted neutron and some number of full-energy neutrons. Cells that detect no downshifted neutrons can be ignored; thus the ratio of downshifted to full-energy neutrons is increased in the remaining cells. We call this approach the multi-hit technique. By increasing the number of cells even more, a limit can be reached when no full-energy neutrons are detected in some cells; thus they are available to detect one downshifted neutron. This is the single-hit technique. With either of these techniques, each cell typically detects neutrons in a pulse mode. This allows measurement of each neutron's energy individually rather than counting an undetermined number of neutrons at once.

The single-hit approach requires a large number of cells in the detector array. For an estimate of this number, consider a case in which the downshifted neutrons make up 1% of the neutron spectrum. If the neutron yield and single-cell detection efficiency are

such that an average of one neutron per cell will be detected, about 60% of the cells in the array will be hit by one or more full-energy neutrons and become "dead." This leaves about 40% of the cells "alive" for the incoming 1% component of downshifted neutrons. Thus there would be a 4×10^{-3} probability that a downshifted neutron is detected in a "live" cell. In order to construct an energy spectrum, it is desirable to detect at least 100 downshifted neutrons; therefore, 25 000 cells are required.

The multi-hit configuration takes advantage of the fact that the downshifted neutrons do not all arrive simultaneously, but over a 20-ns period following the full-energy burst. This delay allows sufficient time for many cells to recover. By using the multi-hit technique, the number of required cells is reduced to about 1000.

Monte Carlo computer simulations using realistic neutron spectra and realistic cell-response functions yield substantially the same detection efficiencies as described above. Single-hit signal processing is simpler and easier to implement than multi-hit systems, but its cost may be prohibitive. We are trying to determine the optimized cost per channel for both configurations.

Conclusion. To measure the fuel (*pr*) on Nova targets, we will focus on a time-of-flight system consisting of an array of small detector cells placed approximately halfway between the target and the chamber wall. This system can be made quite sensitive, it will use existing technology, it requires no special target, and it can provide a rapid and direct measure of the fuel (*pr*). The major risk lies with x rays, electrical noise from the target, and neutrons scattered by structures inside the target chamber other than the fuel. All of these might tend to overwhelm the fuel-scattered neutron signal. Simple analysis of the emitted neutron spectrum suggests that a 90% D to 10% T fuel mixture is optimal. An appropriate detector system might consist of an array of silicon detectors placed on a chip using integrated circuit technology.

We also plan to implement some form of neutron-activation technique. This method is established, passive, and potentially very sensitive. We are analyzing methods of incorporating appropriate tracer materials into

the target. Experiments are being conducted at LLNL and at the University of Rochester's Omega Laser Facility to develop debris-collection techniques and to determine the tracer's effect on target performance.

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Fast-Scintillator Measurements. We are investigating scintillators because their fast timing properties may be applied to the development of neutron diagnostics. Measuring the history of a target burn by direct observation of DT neutrons requires a time resolution of 20 ps. An instrument designed to measure the plasma ion temperature by neutron time of flight, when the flight path is ≤ 1 m, requires a detector system with resolution of 60 to 100 ps.

Fast plastic scintillators like NE111, BC-422, and SG180 typically have decay constants of about 1400 ps. With quenching, the decay constant can be decreased to about 700 ps—still too slow for the instruments that we would like to build. One yet-unexploited property of fast scintillators is

their rise time. If a scintillator's rise time is short and its decay time is long relative to the particle time distribution being measured, then the distribution of particle arrival times is approximately the time derivative of the light output. Accurate rise-time information for fast scintillators has been unavailable in the literature. Using a 200-ps (FWHM) burst of soft x rays, Cheng et al. estimated the 10 to 90% rise time of their samples to be about 100 ps.¹³⁹ Other experimenters suggested that the rise times of fast plastic scintillators are greater than 100 ps.¹⁴⁰ Previous scintillator studies were designed to measure sample decay times greater than several hundred picoseconds. In each study, either the excitation pulse or the detector-system response made it difficult to determine a sample rise time less than 100 ps.

Rise-Time Measurements. In 1984, we began experiments designed to measure scintillator rise times. For our application—the measurement of target burn histories—we are especially concerned with the temporal width of the sample excitation pulse, the temporal resolution of our measurement system, and the need to characterize the excitation pulse and the scintillator output simultaneously. Our initial work was done at the LLNL electron linear accelerator (linac), which provides a beam of 120-MeV electrons to excite the scintillator sample. A beam pulse consists of 1 to 100 micropulses spaced at intervals of 350 ps. The width of a micropulse is 30 to 40 ps.

Figure 5-54 shows the experimental configuration we used to measure scintillator time response as a function of spectral wavelength. The electron beam from the linac is focused onto a 1-mm-thick scintillator sample at a point several centimeters

Fig. 5-54. Experimental configuration used at the LLNL linac to measure scintillator rise times.

Fig. 5-55. (a) Sample of streak-camera data recorded with the instrument configuration shown in Fig. 5-54. (b) Temporal response of scintillator (SG180) taken at $\lambda = 0.38 \mu\text{m}$.

past the end of the linac vacuum pipe. Approximately 0.5 MeV of energy per electron is lost as the beam travels through a 1-mm-thick sample. Cerenkov light generated in the sample and sample scintillation are collected with a spherical lens and relayed by a turning mirror to a cylindrical lens, which focuses the light onto the input slit of a spectrometer. The spectrometer output is recorded by an LLNL ultrafast optical streak camera with an S20 photocathode. The temporal response of this arrangement is about 25 ps (FWHM). The spectrometer has a 20-ps (FWHM) response due to its 1-cm-wide, 600-line/mm diffraction grating, and the streak camera has less than a 10-ps (FWHM) response.

Figure 5-55(a) shows streak-camera data recorded with this arrangement using a sample of SG180. The long-lived scintillator output around $0.38\text{ }\mu\text{m}$ is easily distinguished from the fast, broadband Cerenkov background pulses. A Cerenkov pulse is a measure of the sample-excitation pulse, since it is formed only as the electrons pass through the sample. Figure 5-55(b) shows the temporal response of the scintillator at the wavelength of peak light output with the Cerenkov background subtracted. Preliminary analysis indicates that the rise times of several samples, including BC-422, are less than 50 ps. Our measurements appear to be limited by the response of the measuring system and the width of the excitation pulse. Our work in this area is continuing.

Application to the Measurement of Target Burn Histories. A fast plastic scintillator is an important component in the design of an efficient neutron streak camera. The design calls for a stack of many $15\text{-}\mu\text{m} \times 1\text{-mm} \times 1\text{-mm}$ sheets of plastic scintillator, each separated by a $10\text{-}\mu\text{m}$ layer of a uranium oxide. Dimensions are approximate and must be optimized for maximum detector efficiency. The scintillator stack is placed against the slit of an optical streak camera, so that light leaving the edge of the plastic layers couples into the camera. Neutrons traveling through the uranium oxide layer cause fission. Fission products leave the oxide layer and deposit much of their energy in the adjacent scintillator layer. This design provides high energy magnification and efficient energy collection for each absorbed neutron. The time for energy transfer from

the fission product to the scintillator is estimated to be several picoseconds. As a pulse of neutrons passes through the scintillator stack, the light output for periods shorter than the decay time of the scintillator will increase to a level proportional to the total number of neutrons passing through the detector. If the scintillator response is fast relative to the temporal width of the burst of neutrons, the derivative of the leading edge of the streak-camera output will be the shape of the neutron pulse.

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Penumbra Imaging for Neutrons. In early laser fusion experiments, we were able to image the fuel burn region by means of α particles emitted from the target.^{141,142} However, as the targets became larger and were driven to higher density by more powerful lasers, the α -particles no longer escaped the target and could not be imaged. On the other hand, thermonuclear neutrons are penetrating enough to image the burn region without significant attenuation or distortion.

We first studied the possibility of using a simple pinhole-camera geometry for neutron imaging. With this device, we concluded that a source intensity of 10^{14} neutrons would be needed to obtain images of useful quality and $10\text{-}\mu\text{m}$ resolution.¹⁴³ Since the neutron flux expected from a Nova target is at most $\sim 10^{14}$, it is important to develop a new technique to image neutrons at lower fluxes.

We considered two approaches to this goal: (1) using multiple pinholes, and (2) applying to neutrons the penumbral imaging technique¹⁴⁴ developed for x-rays by Nugent and Luther-Davies at the Australian National University Laser Physics Laboratory. By means of large-aperture penumbral imaging, these authors obtained good-quality, high-energy x-ray images that could not be obtained using a conventional pinhole camera because of low x-ray fluxes. Thus they have already solved the x-ray analog of our neutron problem. In spite of differences from the x-ray case arising from the large mean free path of the neutrons and from neutron scattering in the pinhole

Fig. 5-56. Simulation geometry for pinhole-aperture neutron penumbral imaging.

Fig. 5-57. Simulations using (a) the pinhole-aperture geometry of Fig. 5-56 and a source intensity of 10^{13} neutrons, and (b) the cylindrical-aperture geometry of Fig. 5-58 and a source intensity of 5×10^{12} neutrons.

collimator material, we decided that the most efficient way to test the concept of neutron penumbral imaging was to apply a simulation code developed by the Australian researchers. At our request, they conducted a simulation and concluded that 10- μ m resolution can, in fact, be obtained with a source intensity of $\sim 10^{13}$ neutrons.^{145,146} They also found that the threshold can be lowered to $\sim 10^{12}$ neutrons if the apertures are arranged in arrays of 3×3 . These findings are important enough to warrant further study.

Because the Australian code takes into account only neutron absorption and not scattering in the collimator material, we carried out Monte Carlo calculations with the neutron transport code TART.¹⁴⁷ Once these simulated images are deconvolved with the Australian code, we should be able to assess the effect of the scattering on the spatial resolution and the quality of the image.

Figure 5-56 shows the geometry used in the simulation. The neutron source, consisting of two point sources separated by a variable distance of 10 to 50 μ m on the x axis, has its center of gravity at the origin. A quantity of 800 000 neutrons is generated within a cone of 0.03° with respect to the z axis. This corresponds to $\sim 10^{13}$ neutrons emitted isotropically from the source. The tapered aperture 10 cm from the source casts the penumbral shadow of the neutrons on the detectors 20 m away. The detectors form a one-dimensional array in the x direction and are assumed to be 100% efficient in detecting 14-MeV neutrons. The simulated one-dimensional penumbral images of the two point sources are shown in Fig. 5-57(a). The change in the images as a

Fig. 5-58. Simulation geometry for cylindrical-aperture neutron penumbral imaging.

function of source separation can clearly be seen. These images will be deconvolved by means of the Australian code.

In γ -ray shadowgraphy, the edge of a large solid cylinder is used to provide maximum one-dimensional isoplanatism in the penumbral images, i.e., the aperture shape in one direction appears the same regardless of the source position in that direction.¹⁴⁸ We used this principle in a second simulation, the geometry of which is illustrated in Fig. 5-58. Two point sources separated by a variable distance of 10 to 50 μm are placed at or near the level of the top edge of a large cylinder, which casts penumbral shadows at the detector 20 m away. A quantity of 400 000 neutrons is generated in a small cone, corresponding to 5×10^{11} neutrons emitted isotropically from the sources. The result of this simulation is presented in Fig. 5-57(b).

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Ultrafast Spectrometer Gating Using a Photoconductive Switch

Recent x-ray laser experiments conducted on the Novette facility employed instrumentation new to the LEAD Program. Two of these instruments—a microchannel-plate, grazing-incidence spectrometer (McPIGS) and a slitless, intensified-readout, gated

spectrometer (SIRGAS)—had been used elsewhere, but either ungated or with relatively long gates. In order to detect the short burst of laser x rays in the presence of a long-lived background, these instruments must be gated shortly after the laser x rays arrive. A fast-gating capability was developed and implemented on Novette,¹⁴⁹ providing a gate-off time less than 200 ps. Two McPIGS and one SIRGAS were used in the x-ray laser experiments.

McPIGS and SIRGAS have very similar detectors: open-faced MCPs with microstrip electrodes plated onto the front surface. There are three electrodes on each MCP, each having a characteristic impedance of 25 Ω . Whenever the MCP electrode is biased to about -1 kV, photoelectrons freed by incident x rays are multiplied along a channel and accelerated into a phosphor screen in fiber-optic contact with film. Reducing the bias voltage sufficiently (below about 500 V) results in gating the MCP "off."

The gating system consists of three principal parts:

1. An optical relay system over 175 m long.
2. A "slow" pulsed bias supply, producing a negative 1-kV, 1- μs pulse.
3. A silicon photoconductive switch and wideband pulse-distribution network.

The photoconductive switch illumination is derived from a laser amplifier near the beginning of the Novette laser chain. An optical relay system was constructed to delay the beam so that it arrives properly timed to the x rays and maintains a good beam profile. The relay system is shown schematically in Fig. 5-59. Multiple reflections

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between plane mirrors fold 120 m of beam path length into only 25 m of space. A pair of lenses with a 6-m focal length provides image relaying. Near the target chamber, the beam is split into five separate beams, one each for the three spectrometers and two for purposes of monitoring. Optical trombones provide timing adjustment for each beam path.

The leakage current of a silicon photoconductive switch will result in thermal runaway if it is dc-biased to 1 kV. To prevent this, a pulsed bias supply that produces ~ 1 kV in a few hundred nanoseconds was developed. The switch bias is present for less than $1 \mu\text{s}$, gating the spectrometer MCP "on" until the switch shorts out the bias voltage. The bias pulse is removed well before the photoconductive switch recovers—in about $100 \mu\text{s}$.

Fig. 5-59. Five-pass trigger beamline relay and folding system.

A silicon photoconductive switch¹⁵⁰ was chosen for three reasons:

1. It provides switching speed essentially limited only by the incident-light pulse width.
2. It can readily switch 1 kV.
3. It can be implemented in a low-impedance geometry.

Figure 5-60 is a schematic of the fast-gate circuit. Each spectrometer consists of three $25\text{-}\Omega$ transmission-line electrodes. A fourth $25\text{-}\Omega$ output was implemented as a monitor. The resulting configuration is a $6.25\text{-}\Omega$ (6-mm wide) strip line, fanned out into four $25\text{-}\Omega$ strip lines. The switch functions by rapidly shorting the pulsed bias voltage to ground.

The three gating signals are carried to the spectrometer over semirigid $25\text{-}\Omega$ cable. Turn-off times staggered by 250 ps provide some time resolution. Similar cables are used to carry the gating signals from the spectrometer to a terminator/monitor module. This module minimizes electrical reflections and provides a means for monitoring the continuity and quality of each signal path.

This gating scheme provides a gate-off transition of less than 200 ps. It has proven to be extremely stable and reliable, and has significantly improved the detection capability of these spectrometers.

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Fig. 5-60. Fast-gate circuitry utilizing photoconductive switch.

Improvements in Streak Camera Electronics

We improved the performance of the avalanche-transistor sweep circuitry in the streak cameras that are used for target and laser diagnostics by reducing trigger-timing drift, adding a circuit to monitor trigger delay, replacing the trigger transistor with a more reliable type, and adding a plug-in option that increases the sweep time to 50 ns. These changes make the camera data more reliable and extend the recording period to cover longer laser pulse and fusion experiment durations.

Avalanche Transistors.¹⁵¹ In the past, trigger-to-sweep-output delay timing drift and timing jitter were excessive in some of the streak cameras. Our investigation showed that this delay drift^{151,152} is a function of the unstable breakdown voltage of some avalanche transistors. Plots of the trigger drift in optical camera No. 12 with a new set of 18 avalanche transistors show ~1 ns of drift during the first day after installation and an initial 1-h warm-up period [Fig. 5-61(a)], and additional drift of 3 ns by the end of the first week [Fig. 5-61(b)]. In contrast, the well burned-in set of transistors normally used in this camera had been stable within a fraction of a nanosecond for over two years. We determined that the timing drift in camera No. 12 was caused by breakdown voltage instability in an avalanche-transistor collector. Then we measured the change in breakdown voltage with time in more than 150 samples (in sets of 20 at a time) of the transistor we have used for over 10 years, the National Semiconductor 2N3700. The breakdown voltage

in 2 to 3% of the transistors dropped by 20 to 70% in the first two weeks of testing and those transistors were discarded. The voltage drift in the remaining samples was between +1.8 and +5.2% per month. A voltage change of +5% results in a change in trigger time of <200 ps; this is acceptable, and thus the +5% per month rate of drift in voltage appears to be tolerable for most of our applications.

We instituted a regular, long-term burn-in procedure in 1984 to select only the most stable transistors for the sweep circuit. Burn-in tests now eliminate transistors with breakdown voltage instability before they are installed in the streak cameras. We also discard the transistors that have more than 600 mV of collector-noise voltage (peak-to-peak), and we bias the remainder (~90%) into breakdown at 90 μ A for about two months. The breakdown voltage is monitored every few days for the first week, and then weekly until the test is terminated. Transistors that drift more than 10% in the first week or 10% per month thereafter are eliminated. Our confidence is high that this approach is adequate because we have several streak cameras with sets of transistors that have been in use for 2 to 5 years. Their triggering stability is excellent: <500 ps change in timing occurs over months of operation after a 0.5-h warmup period. However, burn-in testing is manpower-intensive. We designed a fully automated burn-in tester operated by a small computer to reduce labor time; this is now under construction.

Trigger-Timing Monitor. The monitor is used primarily to set up and check the timing of streak cameras in the Nova electrical triggering system. A circuit was installed to

Fig. 5-61. Trigger drift for optical streak camera No. 12 with a new set of avalanche transistors: (a) drift in the first 8 h after a 1-h warm-up and (b) total drift in the first week after installation.

monitor the time difference between electrical trigger input to the streak camera and the actual time the avalanche transistors fire. In normal operation, this delay is ~ 19 ns measured at the monitor output and jitter is less than ± 100 ps. Any deviation is an indication that the camera is not working properly.

A capacitive pickoff consisting of two 1.6-pF printed circuit capacitors was constructed, and added to the back side of the sweep board across the four transistors in the middle of the avalanche string. The output is cabled to a small insertion box containing a discriminator and an amplifier that we designed, and the box is mounted on the bottom plate of the streak camera. The amplifier produces a -6.8 -V pulse with a 1.8-ns fall time and a duration of ~ 300 ns. The pulse, shown in Fig. 5-62, is cabled to a 50- Ω SMA connector mounted on the right door of the camera.

Trigger-Transistor Current Regulator. The electrical trigger to the streak camera fires the base of an avalanche transistor, which in turn triggers the collectors of the avalanche transistors in the sweep string. Selecting reliable trigger transistors was time-consuming because an unacceptably low fraction of the 2N3700s avalanche when the base is triggered. We switched to the Raytheon RS3800 (TO-18 version of the R33500) because they can be triggered at the base. The 2N3700s used a bias of $90 \mu\text{A}$ and the replacement transistors require a bias of $\sim 15 \mu\text{A}$. We designed a series pass-current regulator on a small printed circuit card, and installed it on the back of the sweep card to reduce the current to the new trigger transistor to $15 \mu\text{A}$.

Fig. 5-62. Monitor output pulse measured at the right door of the streak camera; scope sensitivity of 1 V/div. and 5 ns/div. gives a fall time of 1.8 ns and an amplitude of 6.8 V.

After installation of the new trigger transistors, we also discovered that timing jitter, which varied from ± 50 to ± 300 ps with the 2N3700s, was consistently less than ± 100 ps with the RS3800s; this is within our stated camera specifications.

Extending Camera Sweep Time. Longer laser pulses and the consequent longer duration of laser fusion processes require longer sweep times for the streak camera. Complete design equations for the plug-in time bases of the streak camera, previously unavailable, were developed¹⁵¹ and used to construct a 50-ns-sweep plug-in called the X-30. Designs for other sweep times can now be executed easily and accurately.

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Facilities

Introduction. The increasing need in fusion research for highly advanced diagnostics led to the construction of on-site facilities for developing such instrumentation. In 1984, we made substantial progress on two facilities: the short-pulse laser laboratory and the microfabrication facility.

The short-pulse laser laboratory, described in the next article, will allow us to produce optical pulses up to 30 Hz, 3 to 6 ps, and ~ 3 to 4 mJ over a wide range of wavelengths in the IR, visible, and near-UV. When fully operational, this ND:YAG pumped dye laser will enable us to test and characterize ultrafast optical switches and time-resolved diagnostics such as the x-ray framing camera now under development.

The microfabrication facility, containing a wide variety of state-of-the-art systems for producing micron and submicron structures, is also nearing completion. As discussed in the article "Microfabrication Facility," later in this section, we will soon be able to produce x-ray optical components such as transmission gratings, zone plates, multilayer mirrors, photocathodes for x-ray streak cameras, and miniature targets for a variety of laser experiments.

An important advantage of these facilities is their location in unclassified areas of

LLNL. Their open access affords us the benefit of working on site with colleagues from outside organizations.

Author: E. M. Campbell

Short-Pulse Laser. The Nova experimental program requires diagnostics that can make time-resolved measurements with a resolution of 10 to 100 ps. The development of many such instruments requires a stable source of optical pulses with energy ~ 1 mJ and duration < 50 ps. We have installed an amplified, synchronously pumped, mode-locked dye laser to serve as this source. This system—the short-pulse facility (SPF)—was first activated and tested in 1984. Its specific applications include:

1. Framing camera development (optical and x ray).
2. Measurement of the temporal resolution of streak cameras.
3. Temporal calibration of fast x-ray and optical diodes.
4. Generation of short electrical pulses.
5. Development of radiation-to-light converters.
6. Short-pulse electrooptical probing.

7. Digitization of fast transients.

8. Neutron streak camera development.

All equipment in the SPF was purchased from Spectra Physics. The facility's current configuration is depicted in Fig. 5-63. A train of 80-ps pulses originates in a Model 3000 mode-locked yttrium-aluminum-garnet (YAG) laser at a pulse-repetition frequency of 82 MHz. Mode locking is accomplished by acousto-optical modulation of cavity gain in the laser. The generated beam consists of $1.064\text{-}\mu\text{m}$ light in a TEM_{00} mode. Its frequency is doubled in a KTP crystal, and the residual $1.064\text{-}\mu\text{m}$ light is dumped. The average output power of this doubled system is approximately 900 mW. The resulting $0.532\text{-}\mu\text{m}$ beam is used to synchronously pump a Model 375 dye laser. The beam from the dye laser, linearly polarized and operating in the TEM_{00} mode, is roughly 150 mW of $0.606\text{-}\mu\text{m}$ light when the dye being pumped is Rhodamine 6G. When the cavity length of the dye laser is adjusted for transform-limited operation, 3-ps pulses were measured at its output by an autocorrelator. Typical SPF performance characteristics are listed in Table 5-8.

Fig. 5-63. Schematic of LLNL short-pulse laser facility.

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Pulse-repetition frequency	Wavelength (μm)	Power level, cw (mW)	Pulse width (ps, FWHM)
82 MHz	1.064	10 000	~ 80
82 MHz	0.532	900	~ 80
82 MHz	0.606	240	~ 3
10 Hz	0.606	35	~ 6

Table 5-8. Output characteristics of short-pulse laser facility.

Pulses from the Model 375 dye laser may be used directly for sampling applications or may be amplified in a four-stage pulsed dye amplifier if high-power operation is required. The dye amplifier, when operated with sulforhodamine dye, has a measured amplification of 3×10^6 at a pulse repetition frequency of 10 Hz. Interstage isolation to suppress ASE is accomplished by spatial filtering except for the final stage, which is isolated by a stream of malachite green absorber. The absorption effectiveness of the malachite is 18 dB for residual unamplified pulses. The suppression of amplified pulses is roughly 3 dB. The width of the amplified pulse, measured with a scanning autocorrelator, is 6 ps (FWHM). Typical energies per pulse are approximately 3.5 mJ.

In summary, we activated and tested a short-pulse laser facility. Work on this facility will continue in order to make it a more versatile tool in the development of Nova diagnostics in the picosecond-measurement domain.

Author: W. R. Graves

Microfabrication Facility. The ability to produce precise microstructure patterns has become an increasingly important part of the LEAD Program's diagnostic effort. Sub-micron fabrication techniques have been applied to produce x-ray optical components such as x-ray gratings, focusing zone plates, and normal-incidence mirrors.¹⁵³⁻¹⁵⁶ All of these applications require fabrication control of a critical dimension as small as the x-ray wavelength. Technological advances in fabricating submicron structures also lend themselves to the production of "macrostructures" such as thick gold zone plates for coded imaging.¹⁵⁷ The micro-fabrication processes applied to our work are based on techniques developed in the integrated-circuit industry. Our micro-fabrication facility was built to reproduce such techniques.

The microfabrication facility consists of an inner clean-room area and an outer fabrication area. The clean room (18×30 ft) consists of conventional laboratory space with several class-100 clean hoods and benches installed over the equipment and work areas. The clean room houses a chemical exhaust hood for work with toxic chemicals, a wafer-cleaning station (where solvents and 13-M Ω distilled water are used), a spin-coating and baking station, and an rf plasma-ashing station. Also in the clean room are facilities for optical lithography. They consist of a contact-photolithography exposure station, a Kasper contact-mask aligner, and an Olympus microscope with several parfocal objective lenses for optical inspection of submicron features. We also have two class-100 laminar-flow hoods to accommodate the assembly of diagnostic instruments.

In the outer area are two 8- $\times$ 10-ft, class-100, portable, overhead clean hoods. In one of them is a thermal evaporator with a crystal thickness monitor, capable of evaporating metals such as aluminum, gold, chromium, and silver. Also in the same hood is an x-ray lithography station, which will be capable of replicating sub-1000-Å lines in thick films. The assembly and testing of this lithographic system should be finished by May 1985. The second clean hood in the outer area houses a customized reactive-ion etching (RIE) system. Using up to five gases, the RIE system will be able to anisotropically etch materials such as silicon, silicon dioxide, polymers, and aluminum. This system should be operational by March 1985. In the same hood will be a broad-area (4-in.-diam) ion-beam milling system. This unit can operate with either argon or oxygen gas. It will produce ion beams with accelerating voltages up to 500 V and current densities up to 5 mA/cm². Delivery of the milling system is due in March 1985.

Using the equipment already installed, we accomplished the following in 1984:

1. Optically replicated structures with a line width of 1 μm .
2. Produced thick ($>5\text{-}\mu\text{m}$), gold, x-ray streak camera resolution plates.
3. Manufactured aluminum witness plates for drive-measurement experiments.

4. Produced freestanding gold x-ray transmission gratings for the TGSS-EM diagnostic used in the x-ray laser experiments of 1984. These were fabricated in cooperation with the MIT Submicron Structures Laboratory.

Within the next year, all of the micro-fabrication systems described above should be operational. We plan to produce submicron-period gratings and zone plates using x-ray lithography, assist in the fabrication of x-ray laser and laser fusion target and diagnostic components, and play an important role in the fabrication of neutron and x-ray diagnostic instruments.

Authors: A. M. Hawryluk and G. F. Stone

Eikonix CID-1 Automated Image Reconstruction System. The advanced diagnostics and x-ray optics group in the LEAD Program has acquired an automated image reconstruction system that reduces the time required to process zone plate coded image (ZPCI)¹⁵⁸ data from laser fusion experiments. ZPCI is used extensively as an ICF diagnostic for x-ray and charged-particle imaging. The advantage of the ZPCI technique are a significantly greater collection solid angle than pinhole camera imaging at equivalent spatial resolution and the ability to image higher energy x rays than the grazing incidence reflection (GIR) x-ray microscopes currently in use. One disadvantage of the ZPCI technique has been the number of steps required to process the data when compared with other x-ray imaging techniques. The new Eikonix image reconstruction system greatly reduces the processing time and permits faster data turnaround.

ZPCI itself is a two-step image technique that uses simple shadowcasting followed by optical reconstruction to reproduce an image of the original source distribution. In the method used here previously,¹⁵⁸ the reconstructed image was recorded directly onto a high-resolution photographic glass plate. This plate was processed and the image inspected with a microscope. If the image had no defects, it was labeled by hand and sent to be digitized with a microdensitometer. Then the digital image was transferred to the LLNL Octopus computer

system for processing. In this old method, shot information and magnification and identification data were recorded manually, introducing the possibility of error.

The new image reconstruction system incorporates a more direct electronic method to record the image data and the shot information. It reduces the number of steps and the time required to process images from laser implosion experiments. The new technique uses the same method as before to record the coded image data and to perform the optical reconstruction, but the reconstructed image is recorded and analyzed by a novel method. The reconstructed image is relayed by a microscope objective onto a linear photodiode array (PDA) with 2048 picture elements (pixels). The linear array is driven through the image in 2048 discrete steps, producing a 2-D data array of intensity values that represent the reconstructed image. The magnification and identification data are recorded by the computer that controls the digitizer and stores a permanent record of all the pertinent information. The direct digitization of the image eliminates chemical processing of the photographic plates, digitization of the data at a remote site, processing the data on an external computer system, and manual recording of data.

The Eikonix CID-1 automated image reconstruction system is shown in Fig. 5-64. The components that were also used in the previous method are a He-Ne laser, spatial filter, collimating lens, coded image holder, and reconstructed image holder. Components that have been added to the new system are a closed-loop laser intensity controller, stepper motor and controller for coded image positioning, image relaying microscope optics, the Eikonix model 780 digitizer camera head, an image display system, a digital position scale, and an LSI-11/23 computer system with RL02 storage disks.

The He-Ne laser light is first passed through the closed-loop laser intensity controller. This unit reads a small fraction of the laser light and monitors the output with a photovoltaic cell. The laser light output is kept within $\pm 1\%$ by the servo-controller during the scan time of the image (typically 30 to 180 s). In the previous technique, the

Fig. 5-64. Eikonix CID-1 automated image reconstruction system.

image was recorded on the photographic plate with a single exposure of $t \leq 1$ s.

The coded image transparency plate is held in a frame attached to a translation stage that is fastened to the optical bench. The plate holder is translated along the axis of the laser beam with a lead screw driven by a stepper motor. A digital scale with a 3-m range is used to obtain the positional accuracy needed for this technique. A magnetically coded strip on this scale has a sliding sensor that is attached to the coded image plate. The sensor moves along the scale coincident with the coded image plane, and reads its position in relation to the reconstructed image plane. This scale can precisely locate the position of the coded image within $\pm 5 \mu\text{m}$ of the reference point over the full 3-m travel path.

Another new system component is the microscope objective assembly; it magnifies and relays the reconstructed image onto the PDA detector in the digitizer head. The reconstructed images are typically 150 to $750 \mu\text{m}$ square. They are magnified from 5 to 20 times, depending on the microscope objective used, before they are relayed to the PDA array. A disadvantage of microscope optics is that optical aberrations may be introduced into the image (e.g., dust particles on the optics could cause diffraction rings in the image plane that would confuse the interpretation of the results).

The heart of the new system is the computer-controlled image digitizer, the

Eikonix model 780 camera head. This image scanner uses a 2048-pixel linear PDA. Each pixel is $15 \mu\text{m}$ wide, spaced on $15\text{-}\mu\text{m}$ centers. The array is stepped through 2048 discrete positions by a fine lead screw connected to a stepper motor.

The Eikonix digitizer has a dynamic range of 255 grey levels, the result of an 8-bit analog-to-digital converter. The top 7 bits contain picture grey-level (image intensity) data and the last bit (bit 8) contains graphics information. The 7 grey-level bits describe 128 levels of image intensity. In those instances where the dynamic range of the image exceeds that of the digitizer, the previous technique of recording the image on film is still available to us.

A DEC LSI-11/23 computer controls the entire digitizer assembly; it has two 10-Mbyte RL02 cartridge disks for operating system and image storage. The image from the digitizer can be sent directly to the disk or to the 512×512 -pixel image display unit, which can show as many as 128 grey levels in addition to graphics labeling information. A software package, EPPL (Eikonix Picture Processing Language), handles data transfer and graphics generation. The types of graphics produced by the image system include image identification data, data intensity ranges, 2-D and 3-D data plots, and 1-D lineouts of intensity versus position. Pseudo-color representations of the intensity data can also be constructed with EPPL. The graphics data can be plotted

in black and white on the Versatec plotter, and images from the display can be recorded on 35-mm film or on 8- x 10-in. color positives or negatives with a color unit attached to the image system. For more sophisticated image processing of the data files, we can transfer images digitized by the CID-1 system onto a DEC RL02 disk cartridge and transfer them to the LEAD Program VAX 11/780.

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Data Management and Analysis

Introduction

During 1984 we continued to make improvements in the computer system for data management and analysis in the LEAD Program. Both hardware and software were substantially upgraded.

In the hardware category, we increased the storage capacity of the LEAF VAX-11/780 computer by over 4 gigabytes with the addition of four Winchester disks. We also added two new high-speed tape drives. These improvements, which are described in the article "System Update," have increased overall system reliability. They will allow easy handling of the increased volume of data expected as a result of Nova experiments. A number of peripherals were also acquired to increase our graphics and hard-copy capabilities and to facilitate file transfer between the VAX and an IBM Personal Computer (another recent addition). We are planning to upgrade the VAX-11/780 to an 11/785 in the summer of 1985. This upgrade will increase our processing capacity by approximately 50%.

The major hardware acquisition in 1984 was a Perkin-Elmer PDS 1010GM microdensitometer, described in the article "PDS Capability." This instrument provides the capability to rapidly digitize film data from Nova. The microdensitometer is controlled by a dedicated PDP-11/34, which has been appropriately upgraded to handle the task.

This computerized photodigitizing system is connected directly to the VAX via DECNET, the Digital Equipment Corporation (DEC) inter-computer network. DECNET will permit rapid transfer of image files from PDP-11/34 storage disks to the VAX for immediate processing. The system will be operational by March 1985.

The need for archiving and correlating shot data and the availability of commercial data base systems have led to the development of an "electronic shotbook." Work on this data management software system, now in its second year, is described in the article "Electronic Shotbook." Basic components are now in place, and we are developing the means to transfer, store, and retrieve Nova diagnostics data. To accomplish this task, we have begun to link the output files from diagnostics data-processing codes directly to the SHOTBOOK system, and we have installed new software to provide graphic displays of data from a number of diagnostics instruments. Examples include streak camera records in which the axis orthogonal to the time direction contains spatial or spectral information. We expect to use SHOTBOOK for the first experimental series on Nova.

The article "Analysis Software" describes other enhanced software capabilities. New and updated programs were developed in-house and acquired from commercial sources. We installed the latest versions of the DISSPLA graphics package and the ORACLE relational data base management system. These new versions have demonstrated increased speed and reduced execution times. The upgrades enhance SHOTBOOK and many of the processing codes that have graphics output. In-house software development has focused on linking data-processing codes to SHOTBOOK and on upgrading the algorithms for streak camera data reduction and calculation of x-ray-detector response functions.

Finally, a secure fiber-optic interface connecting the Nova control VAX system to the classified LEAF VAX has been installed. This communication link is similar to that used successfully in the Novette system. It is one-way, so that unprocessed shot data can only be sent to the LEAF VAX for processing. Communication of data in the

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opposite direction is inhibited. This interface is described in the article "Nova Communications and Data Transfer."

Author: E. M. Campbell

System Update

The Laser Experiments Analysis Facility consists of two DEC computers—a PDP-11/34 and a VAX/VMS-11/780; and a Perkin-Elmer 1010GM microdensitometer—a photodigitizing system (PDS). The PDP-11/34 is a dedicated machine used to control the PDS, and the VAX is used for storage and analysis of ICF data and for general-purpose computing. A one-way, fiber-optic link connects the LEAF VAX to the Nova

VAX for transfer of shot data. Figure 5-65 shows a block diagram of the system components and their logical connections.

Over the past year, major improvements were made in the hardware, system software, and applications codes. The VAX hardware was modernized by the addition of two high-speed tape drives and four Winchester disks. The disks have a combined capacity of over 4 gigabytes. A number of special peripherals were also added: an array processor, an IBM Personal Computer (PC), a Versatec Random Element Processor (REP), a high-resolution color monitor, and two 1000-line-per-minute Versatec printer/plotters. Currently on order is an upgrade to the VAX central processing unit (CPU) that will increase its computing speed by about 50%.

Fig. 5-65. A block diagram of the system components and their logical connections.

Acquisition of the Winchester disk drives was dictated by two immediate requirements: increased storage capacity for Nova shot data and higher reliability than could be achieved with the old removable pack drives. Two of these old drives were accepted as trade-ins on the purchase of the four new ones. Similarly the old tape drives, which had a maximum capacity of 1600 bytes per inch (bpi), were traded in for two new ones with a 6250-bpi capacity and higher data-recording reliability. The IBM PC was installed primarily to allow transfer of files from 5-1/4-in. diskettes to the VAX, since many of the staff utilize IBM PCs in their work. The REP is a dedicated microprocessor used to scan-convert plot files for the Versatec printer/plotters. Off-loading this process-intensive task from the VAX to the REP allows plots to be produced at much higher speeds. The new printer/plotters were installed for the increased reliability and higher-quality output they provide. The VAX CPU upgrade, to be installed in the summer of 1985, will ensure the computing power needed when Nova is fully operational. DEC will replace the CPU, cache memory, floating-point accelerator, and backplane with faster components.

Like the VAX, the PDP-11/34 system has been improved by the addition of larger-capacity disk drives. In addition, several old, unneeded peripherals were removed. These changes were necessary to provide adequate storage capacity for digitized film data. Communication circuit boards for controlling both the photodigitizer and a DECNET connection to the VAX were also added.

The operating systems of both the VAX and the PDP have been upgraded to the latest version, and software has been added to control the DECNET connection. This network will allow transfer of digitized film data from the PDP to the VAX for storage and analysis.

Author: J. E. Randolph

PDS Capability

The LEAD Program has acquired a Perkin-Elmer PDS 1010GM microdensitometer to digitize film data produced from Nova

experiments. The digitizer will provide, for the first time, an in-house capability to rapidly photodigitize film up to 25 × 25 cm in area, in pixel sizes as small as 5 μm, and with optical densities up to 5.11 D (optical transmission $\tau \leq 10^{-5.11}$). Prior to this, data were scanned outside the immediate program facility on the PDS 1010A in the laboratory's L Division. Magnetic tapes containing the results were then hand-carried to the LEAF VAX to be read for storage. Since the data on the tapes were not in VAX format, they had to be read record by record and assembled into a correctly formatted file. A large file required several minutes of VAX time. This former data-transfer process is shown in Fig. 5-66.

The new PDS, consisting of the 1010GM microdensitometer controlled by a DEC PDP-11/34 computer, is shown in Fig. 5-67. Located within the LEAD computer facility, the PDS is connected to the LEAF VAX via DECNET, which permits rapid transfer of image files from the PDP storage disks to the VAX storage disks. As a backup for the network, magnetic tapes can be used to transfer files from the PDP-11/34 to the VAX. Coupled with the VAX, the PDS 1010GM will greatly enhance the data analysis capabilities of the LEAD Program.

The PDS 1010GM is a computer-controlled, flatbed, scanning microdensitometer. The 2300-pound machine consists of scanning optics and a stage carriage held on granite blocks. The entire assembly rests on a metal spaceframe support stand, which is bolted to the building to isolate it from vibrations. A granite center C-frame bears the y-axis carriage, and the x-axis carriage rides on two pieces of granite, one on each side of the center frame. A console at the right of the base houses the control electronics, x- and y-axis stepper-motor controls, photomultiplier (PMT) control electronics, a digital readout indicating the relative stage position, and the optical-density readout.

During the scanning process, the light from an incandescent source is imaged, by means of microscope optics, through an aperture onto the emulsion of the film being digitized. The light passing through the illuminated area of the film is then relayed by similar microscope optics and apertures onto the face of an S-11 response PMT. The electrical signal from the PMT is proportional

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Fig. 5-66. The method of handling the transfer of digitized data to the LEAF VAX prior to the installation of the new system.

Fig. 5-67. Schematic diagram of the LEAF photodigitizing system including the PDS 1010GM, PDP-11/34 computer system, and the inter-computer network DECNET.

to the optical transmission of light through the film. The transmission data are collected by raster-scanning an area on the film to produce a "numerical image" of the original film record. The transmission values $t(x, y)$ are converted to optical density $D(x, y)$, and the data are recorded onto a storage disk for subsequent processing. This numerical image consists of data points representing the individual picture elements (or pixels) that make up the image. During the digitizing, the stage position is controlled to $\pm 5 \mu\text{m}$ over the full 25 cm of travel by means of linear optical encoders, with a repeatability of $\pm 1 \mu\text{m}$. The numerical image of the film record is then reformatted in the PDP-11/34 and stored on disk for eventual transfer to the LEAF VAX.

The PDS can scan film samples at speeds up to 20 cm/s in the x direction, but the y -axis scanning speed is limited to 2 cm/s because of the large granite C-frame holding the y axis and optics. The sample can be digitized with a spatial resolution at the film plane of 5 to 250 μm for circular or square apertures. For rectangular apertures, spatial resolution can range from $2.5 \times 500 \mu\text{m}$ to $100 \times 1600 \mu\text{m}$. The individual pixels can be digitized or sampled in discrete step sizes ranging from $\Delta x = \Delta y = 1 \mu\text{m}$ to the full 25-cm travel of the stage. The PMT assembly and associated electronics are capable of measuring optical densities up to 5.11 D in 1024 discrete levels, translating to a density resolution of 0.005 D. When a single scan of the data on a film record will suffice, 1-D plots of density vs position can be obtained from a strip-chart recorder. The PDS is situated under a Class-100, laminar-flow, clean hood to ensure a dust-free environment during scanning.

The digitizing of data is continuous while the sample stage is in motion. The area being scanned and the scan parameters are specified by a master control program, SCANSALOT, which resides on the PDP-11/34 computer. SCANSALOT directly controls the many functions of the digitizer. The program downloads instructions from the PDP-11/34 to the X-10, an MC6800 microprocessor in the digitizer, which controls the digitization process. Parameters specified by SCANSALOT include the scan direction and step size on both the x and y axes, the origin of the scan, the number of

raster scans per film, and identification lines for each scan. SCANSALOT also performs general-housekeeping and diagnostic functions needed to maintain proper operation of the system.

When a film sample is digitized, SCANSALOT collects the data and writes it to a disk file on the PDP-11/34. The data are repacked with one header per file and with each record representing a single scan line. This file format is compatible with existing analysis codes on the VAX. Once a file is created, it is sent over DECNET to the VAX and stored—available for the experimenter to manipulate. The average time to scan a film record and transfer the image file to the VAX is about 10 min.

To provide PDS images in the Octopus system¹⁵⁹ data format (which is different from the VAX file format), a computer code will be written to produce tapes identical to those made by the L-Division PDS. This code will allow any image files in the LEAF PDS format to be written onto a tape and transported for image processing to the Octopus computer network or to other LLNL computer systems.

Authors: G. F. Stone and J. M. Auerbach

Electronic Shotbook

SHOTBOOK¹⁶⁰ is a utility software package developed in-house with structured programming techniques.^{161,162} The software is designed to provide LEAD Program staff members with rapid, easy access to final experimental results. It is supported on the LEAF VAX-11/780 computer—an integral part of the LEAD Program. SHOTBOOK is based on the commercial relational data base management system ORACLE.¹⁶³ The application programs of SHOTBOOK utilize DEC's Forms Management System (FMS)¹⁶⁴ and the graphics library DISSPLA.¹⁶⁵

SHOTBOOK will accommodate a unique feature of the Nova laser: the fact that Nova will have at least two target chambers. Because multiple chambers imply a need to account separately for data coming from each chamber, SHOTBOOK has been programmed to manage this task. The system will keep track of the origin of all data without ambiguity.

SHOTBOOK offers access to final experimental results through a series of menu selections. Menus eliminate the need for a sophisticated knowledge of the computer or the data base. The entire operating system of SHOTBOOK is now complete and operational. Additional display sections are under development, with the goal of providing:

- Consistency, readability, and ease of access to final experimental results.
- Software-controlled entry of data into the data base.
- Close administrative control of data-modification privileges (e.g., updating, deletions, or insertions) to ensure data integrity.
- Standard display formats.

There are two prerequisites for implementing a new display. First, the display software itself must be programmed and available to SHOTBOOK as a menu option; and secondly, appropriate processing software must be available to place related data in the LEAF data base. The display software then draws from the data. New display software will continue to be graphical in order to maximize uniform communication of experimental results.

Figure 5-68 shows the relation between SHOTBOOK and its environment. Raw data (film images, digital data, and instrument

configuration data) are generated at the laser facility when an experiment (shot) is conducted. Following a shot, the relevant digital data are sent, via a one-way optical-fiber link, to the LEAF VAX. In addition, films from the optical diagnostics are digitized, and the resulting files are sent to LEAF. These raw data are stored on the VAX disk, with the shot number as the key identifier. Next, each type of raw data is analyzed by a particular processing code to obtain the final calibrated experimental results. These codes generate printed data summaries for the experimenters and place the experimental results into the LEAF data base. The electronic shotbook is the tool experimental staff members use for easy, rapid access to these stored results.

Using SHOTBOOK. Figures 5-69 to 5-71 provide an impression of the environment encountered by a user of SHOTBOOK. Figure 5-69 shows menus that lead the user to select a particular experimental series. In turn, the menus in Fig. 5-70 allow a choice of subseries, shot number, and display type for the specified shot. Figure 5-71(a) shows one type of display the user can select (22 × streak microscope). With preknowledge of the shot number, the user can reach the desired display more directly and bypass this stepwise assistance. Each display provides

Fig. 5-68. Relation between the electronic shotbook and data collection, analysis, and storage.

Select: "Examine Shot Results"

Select: "List Shot Series"

Select a series

Fig. 5-69. Selecting a particular experiment series.

Select a subseries

Select a shot

Select a display

Fig. 5-70. Selecting a particular display.

(a) View Display

(b) Select frame with cross hairs

(c) Zoom and take horizontal lineout

Fig. 5-71. Examining data.

clear and complete identification of the exhibited data. For this reason, heavy emphasis is placed on the use of graphics.

Users may just wish to "page" through a display section and examine the data; however, they can also enter "cross hair mode," illustrated by Fig. 5-71(b), and readily reformat a display. They can shift the ordinate or abscissa scales of a plot, switch between linear and logarithmic scaling, make vertical or horizontal lineouts of contour plots, or zoom in on a selected plot area, i.e., display a full-screen view as in Fig. 5-71(c) with any of the above-mentioned enhancements. While such actions can produce drastic changes in the appearance of data displays, these changes remain unique to the user who invokes them. No changes occur either in the data base or to displays viewed by other SHOTBOOK users. As users review

shot data, they can obtain permanent records of interesting displays. A hard copy of any screen can be printed by pressing a single key.

Access Levels. It is the task of SHOTBOOK to support two groups of users: those who are responsible for inserting, deleting, and updating data associated with a specific shot or diagnostic; and those who merely need access to final experimental results. For example, experimental physicists must define in SHOTBOOK such details as the experimental series, subseries, and desired shot diagnostics. Such definition provides a necessary link to the results developed by the shot data-processing codes. On the other hand, diagnostic physicists must be able to review particular results obtained from their diagnostics as well as to insert or revise experimental comments

regarding these results for later assistance to the casual user. Anyone without specific data responsibility on a given shot is considered to be a casual user, who is given access to shot results but not the privilege to insert, update, or delete results. SHOTBOOK access is designed so that the highly privileged user on one experiment becomes a casual user with respect to another's work. It is basic to the design of SHOTBOOK's security scheme that a clear and relatively transparent separation be maintained between the various classes of users. Data base integrity is maintained by assigning update, insert, and delete privileges for specific data to those who are responsible for it.

Although simulated data were used to develop new displays, testing is now done with actual results from the processing codes. As users continue to drive SHOTBOOK with real data, data-protection measures will be implemented.

Displays and Data Analysis. Software to produce three new types of displays has been completed: SCL, CONTOUR, and LINEOUT. SCL gives line plots for time slices from a streak-camera image, plus a summary plot. The CONTOUR package plots any two-dimensional array up to 512×512 . It reduces the original array to 30×30 to accommodate the DISSPLA graphics library contour routine. LINEOUT superimposes an elevation profile reflecting a selected plane of a contour plot. It uses the original data array (up to 512×512) and overlays a lineout across the contour. It also draws a dashed line to indicate the position across the contour from which the lineout was taken. Both the lineout and the dashed line are drawn by using calls to the user's terminal to obtain a very fast screen refresh. Either vertical or horizontal lineouts are allowed.

Three new graphic displays have been implemented: OX, BEAMTIME, and $22 \times$. The OX display plots data from the optical/x-ray streak camera by using the SCL package. The BEAMTIME display plots both incident- and reflected-beam temporal profiles. It has the capability to display all 10 beams from Nova. The $22 \times$ display plots data from the $22 \times$ x-ray microscope by using both the CONTOUR and LINEOUT packages. The data, axis labels,

and title information needed for these displays all come from the data base and related data files.

Authors: R. A. Lerche, E. W. McCauley, and J. G. Shaw

Analysis Software

During the past year, software development focused on:

- Processing codes to analyze data from diagnostics.
- An interface between processing codes and the electronic shotbook.
- Increased capability to calculate x-ray response functions.

To improve our analysis capability, we installed the latest versions of our two most-utilized commercial software packages: the DISSPLA graphics library and the ORACLE relational data base management system. ORACLE, Version 4.0, has made data base queries three times faster than the previous version. In certain applications, revised calls in the ORACLE-Fortran interface also result in dramatically decreased execution time. In previous ORACLE versions, these calls had to be executed as many times as the number of quantities desired; now multiple quantities can be retrieved with only one call. The new ORACLE also allows quicker execution of the electronic shotbook code (SHOTBOOK). The updated DISSPLA library provides a faster interface for Tektronix terminals. This improvement has increased the speed of processing codes that produce graphic displays on these terminals.

New Diagnostics Processing Software. Upgraded codes for processing streak-camera-based diagnostics data are now compatible with Ramtek color display systems that have trackball cursors. These graphic systems display color-enhanced contour plots of digitized data from streak-camera film. The user can move the cursor to a selected position on a color contour plot and transmit the screen position to the computer. The system then relates that site to an index in the data array, which in turn can be used to produce profile displays or simply to retrieve values of the data at that point.

Another improvement in the streak-camera processing software is the capability to use the entire film-density data array in unfolding the diagnostic system response. In previous years, only lineouts of density data could be unfolded. The primary unfolding operation consists of converting film density to relative exposure by means of a D -log E curve, which provides the film density corresponding to a given exposure. The exposure corresponds to light energy from the streak-camera phosphor on the film. To relate this energy to intensity at the target, the response of the diagnostic must be determined. The front end of a streak-camera diagnostic contains components for wavelength and spatial resolution. The diagnostic records this quantity as a function of the camera's sweep direction (time). Since most streak cameras are not absolutely calibrated, the processing software calculates relative source intensity as a function of position or wavelength in one direction and as a function of time in the other direction. This new software methodology was incorporated into an existing processing code for space-time diagnostics such as the $22\times$ streak x-ray microscope and the spatially discriminating streak spectrometer (SDSS). These diagnostics provide space-time information on target x-ray emission over a relatively narrow energy band. New data processing codes have been developed for the TGSS-EM, which provides a time history of x-ray spectra in the sub-keV range, and the streak optical spectrometer (SPEX), which provides a time history of Raman-scattered light emitted by the target.

The SPEX is a combination reflection grating spectrometer and streak camera. It records scattered light from the target, ranging from approximately the laser wavelength to half the laser wavelength. The corresponding analysis program receives as input: (1) a file created by digitizing the film streak record; (2) a file containing the D -log E response of the film; and (3) a file containing the spectral response of the photocathode (sensitivity vs wavelength), the spectral response of various filter glasses used between the target and the spectrometer (transmission vs wavelength), and the dispersion of the spectrometer. The SPEX is configured so that a given position along the streak-camera slit corresponds to a

Fig. 5-72. Screen display of film-density data from a SPEX streak record.

Fig. 5-73. Color-enhanced screen display of SPEX intensity data showing a lineout taken in the horizontal direction.

particular wavelength. This design allows the user viewing a color display of the streak to align the trackball cursor on a known fiducial feature and enter a wavelength value to initialize the dispersion relationship.

Figure 5-72 shows a contour plot of a digitized SPEX streak record of film density.

Figure 5-73 shows a contour plot of relative intensity after spectrometer, filter, and film response are unfolded from the density data. The horizontal direction of the contour plot corresponds to wavelength, and the vertical direction to time. Using the trackball cursor, the user can select spectra or time profiles from this color display.

Figure 5-74 shows the spectrum corresponding to the horizontal lineout of Fig. 5-73. The entire intensity distribution is stored by the code in a disk file, which can

Fig. 5-74. Detailed plot of the intensity lineout shown in Fig. 5-73.

Fig. 5-75. Screen display of film-density data from an x-ray pinhole streak record.

subsequently be read by other software such as SHOTBOOK and by general image-processing codes.

The $22 \times$ streak x-ray microscope and the SDSS fall into a class of diagnostics called $x-t$ (or space-time) diagnostics. Their output is a film record produced by a streak camera. The film shows the time history of the x-ray emission from the section of the target seen through the streak camera slit. For example, in the case of a spherical target, the $22 \times$ microscope produces a time history of the x-ray emission along the diameter

of the sphere. The processing code for $x-t$ diagnostics uses as input the magnification, camera sweep speed, and $D\text{-log } E$ film response to convert the film density data of the streak record to the intensity distribution at the target as a function of space and time. Viewing a color display of this distribution, the user can select horizontal and vertical lineouts by means of the trackball cursor. In turn, these lineouts can be converted to spatial and time profiles of x-ray emission. Each data point in the streak record data array is converted to relative exposure. Figures 5-75 to 5-77 are displays created by the data processing software.

Figure 5-75 shows a contour plot of the film-density distribution from a $22 \times$ microscope. Figure 5-76 shows a contour plot of the intensity distribution unfolded from the film-density data by means of a $D\text{-log } E$ film-response data set. Superimposed on this last display is a horizontal lineout (space profile) generated by the software. Figure 5-77 represents a different type of display showing a detailed plot of this spatial profile. As with the SPEX software, the $x-t$ processing code stores the intensity distribution in a file for later use.

Another new diagnostic this year is the pinhole microchannel plate (MCP) camera which consists of a pinhole camera in which

x-ray spatial-emission patterns from the target are recorded by a combination of MCP, scintillator, and film. Processing software unfolds the film-density data to relative x-ray intensity, allows the integrals of selected sub-regions of the distribution to be calculated, and offers a lineout option. The local integration results are stored in a table of the ORACLE data base and can be used for comparative analysis of x-ray emission from different regions of each target. As with the streak-camera software, the user can display full-color or gray-scale contour plots of the calculated relative-intensity distributions.

Electronic Shotbook-Processing Code Interface. A major software development effort this year focused on the interface between processing codes and SHOTBOOK. The interface consists of a data base that holds the processing-code output and the associated software to read and write it. Three categories of data have been identified for display by SHOTBOOK:

1. Scalar data: single values or small data sets, each element of which has significance by itself.
 2. Profile data: large sets of x, y pairs corresponding to such entities as laser power histories and x-ray spectra.
 3. 3-D data: large sets of x, y, z data values.
- These data sets correspond to digitized

film images unfolded to intensity distributions.

There is a general format for storing each type of data that is the standard for each processing code. Data items in category 1 are sorted into tables in the ORACLE data base. They are read from or written into the data base by means of a generalized ORACLE-Fortran interface (ORACLE HLI). Data in categories 2 and 3 are maintained in files, which also contain titles for each profile or set of contour data, labels for plot axes, dimensions of x, y , and z quantities, the number of data pairs in each profile, and the row

Fig. 5-76. Color-enhanced screen display image of the relative intensity distribution on a streak record from an x-ray pinhole. A horizontal lineout is superimposed on the display.

Space profile 81062206 - 22X T - 1.4630E + 00 nsec

Fig. 5-77. Detailed plot of the intensity lineout shown in Fig. 5-76.

Fig. 5-78. The interface between the diagnostic processing software and SHOTBOOK.

and column dimensions of the 3-D data array. Four subroutines were developed to read and write the file data. The processing codes use subroutines WRT_PROFILE_FILE and WRT_CONTUR_FILE, and SHOTBOOK uses subroutines RD_PROFILE_FILE and RD_CONTUR_FILE. Figure 5-78 shows a data-flow diagram illustrating the processing code-SHOTBOOK interface.

X-Ray Detector Response-Function

Calculations. We made significant improvements in the software for calculating the response functions of x-ray detector channels. First, we extended the Henke library¹⁶⁶ of f_1 and f_2 atomic scattering factors from an upper energy limit of 2 keV to a new upper energy limit of 10 keV. This was the first such effort worldwide, and the extended

library is now used in many laboratories in the United States. This library allows us to rapidly calculate reflectivities for x-ray mirrors and efficiencies for transmission gratings in x-ray detector channels. We also upgraded the software for calculating response functions for the Dante soft x-ray spectrometer. X-ray-mirror responses are now included in these calculations. This software was crucial in understanding spectra from the many Novette shots that included Dante spectrometers with several mirror channels.

To support x-ray-streak camera data processing, we incorporated the Henke data library into software for calculating photocathode response. The quantum efficiency of a photocathode is proportional to the

scattering factor f_2 for a given material. So with the f_2 library, quantum efficiency calculations became a fast and simple task. This last type of calculation has been applied to processing data from the TGSS-EM and the XCSS.

Author: J. M. Auerbach

Nova Communications and Data Transfer

A typical target experiment on Nova will produce many megabytes of digitized data that must be processed and analyzed for proper interpretation of the experiment. This article describes the form of the raw data produced at Nova for ultimate transfer to LEAF and the means for transfer.

Complete interpretation of an experiment requires processing and analysis of data from the Nova target and laser diagnostic systems. Data from these systems are recorded by electronic sensors and transferred to film or to digital recording systems, i.e., CCDs and CAMAC-based electronics. The digitized electronic data are then sent to the Nova control VAX computer system for transport to the LEAF VAX. Films containing data are developed and then digitized by means of the LEAF PDS 1010GM microdensitometer.

Data files created at Nova contain the following:

1. Electronic sensor data from the target diagnostic system.
2. Laser energy data from the laser diagnostic system.
3. Data from CCDs recording output of optical streak cameras in the laser diagnostic system.

One-Way Link. After a shot, all files are transmitted to the LEAF VAX disks via a one-way fiber-optic link. This special inter-computer link was built to satisfy the security requirements of the LEAF VAX, which is a classified computer system. No data other than a simple yes/no (0 or 1) can be sent from it to an unclassified computer system such as the Nova control system of four VAX computers. The hardware comprising the link is identical to that used at the Novette laser facility. A diagram of the link is shown in Fig. 5-79.

The link consists of NOVALINK network cards installed in the Nova north VAX and in the LEAF VAX. The two computers communicate with each other via a communication controller consisting of an LSI-11 microprocessor containing two more NOVALINK cards (one each on the LEAF side and on the NOVA side). The LSI-11 contains a resident programmable read-only memory (PROM) program that allows full data transfer only from Nova to the LEAF VAX. When transfer in the other direction is attempted, the LSI-11 program masks out the left-most 15 bits of each 16-bit data word and limits the transfer rate of the

Fig. 5-79. Diagram of one-way link from Nova to the LEAF VAX showing the principal hardware components.

remaining bit to once per second. This rate-encoding process prohibits data transfer while allowing simple yes/no signals for synchronization and control.

Data files are transferred from Nova to the LEAF VAX by means of specialized file-handling programs: a file transmitter on the Nova VAX, and a file receiver on the LEAF VAX. Both programs use the Novanet driver for communication through the NOVALINK cards. The file-transmitter program opens the file to be transferred, reads its characteristics from the file header, and sends those parameters over the link to the file receiver. The receiver then uses the parameters to open a file in the LEAF computer system with the same characteristics as the one on Nova. As soon as the file is open, the file receiver sends a "yes" reply to Nova indicating it is ready to receive the data contents of the file. Data are then sent from Nova to the LEAF computer by block transfer, which is independent of the record structure of the file.

In block transfer, consecutive 512-byte sections of the file are sent over the link. Block transfer is made possible by the "block I/O" feature of the VAX VMS operating system. Note that the operation of the LSI-11 is completely independent of any file structure on the VAX. The LSI-11 simply passes buffers of file data originally broken down by the file-transmitter program. The file-receiver program on the LEAF VAX is always active, waiting for a "ready to send file" signal from Nova. Sending of files at Nova is initiated on a file-by-file basis by the operator. An 800-block file requires about 30 seconds to transfer. The transfer rate is limited by the operation of the LSI-11. The software is designed to cause re-transmission whenever an error is detected.

Target Diagnostic Electronic Data.

Electronic sensors in the target diagnostics are the primary source of raw data used by the LEAF data-processing software. Each file contains all the raw data collected by a CAMAC-based data acquisition system plus mapping indicating the path between the output of a detector in a diagnostic system and either: (1) the channel of a recording module (i.e., charge integrator or waveform digitizer), or (2) the channel of a programmable power supply or programmable amplifier.

Raw data in the file are labeled by the name and identification number of the associated recording module, programmable power supply, or amplifier. This form of labeling is most efficient for a data acquisition system but is of no use to experimenters, who must see data labeled and grouped by diagnostic and detector number when analyzing data for physical correlations. Hence, the mapping is crucial in that it allows automated data processing programs to relabel and regroup the data as desired.

The mapping is used in processing all CAMAC-based diagnostic data. It is taken from a Nova-originated file once, preprocessed, and inserted into an ORACLE data base, where it can be accessed much more quickly and easily than if kept in the original file. Moreover, once in the data base, the mapping provides an on-line archive of the diagnostic setup for all shots.¹⁶⁷ Each time data from a particular diagnostic is processed, the respective program pulls only the raw data from the target diagnostic data file. The required mapping parameters are taken from the ORACLE data base.

Laser Energy Data. The laser diagnostic system produces files containing laser energy data recorded by diodes from positions along the laser amplifier chain, and from the incident- and reflected-beam modules that deliver laser energy to the target. The data used by experimenters are a small subset of all the recorded data, consisting of the laser energy incident on the target and the laser energy reflected back into focusing lenses. After the laser-energy data file is transferred to the LEAF VAX via the one-way link, the incident- and reflected-energy data are extracted from the file, and the results stored in the ORACLE data base where they are indexed by shot number and beam number.

Laser energy data from all shots remain on-line in this data base for access by processing codes that calculate, for example, laser energy absorbed by the target.

Laser Temporal Data. The largest data files sent over the one-way link hold digitized data from CCDs. These record the output of streak cameras measuring the temporal histories of the incident laser light on the target and that reflected back through the lenses. In processing the CCD data files, lineouts of the center region of each streak pulse pattern are obtained in

terms of CCD counts versus position. Then laser-energy and streak-camera-speed data are used to convert these lineouts to a time history of incident or reflected laser energy.

Film Data. After being digitized, film data constitute the bulk of all data from Nova laser and target diagnostics. Laser-diagnostic film data include near-field and far-field beam images. These images will be digitized by means of the LEAF PDS 1010GM microdensitometer and analyzed on the LEAF VAX. The results provide data on the spatial uniformity of the beam and the intensity of the pulse at the target. After analysis, the digitized film data will be archived in magnetic tape libraries. Film data from target diagnostics include streak records from optical and x-ray streak cameras and x-ray imaging devices (i.e., x-ray microscopes and pinhole cameras).

From the start, Nova will have a highly automated means of transferring raw experimental data to the LEAF VAX for interpretation. Most significantly, the transfer of digitized film data will be completely automated, since film digitization will be done by means of the LEAF PDS.

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Nova Target Diagnostics

Introduction

The Nova laser facility at LLNL is the most sophisticated high energy laser experimental facility in the world. When it is completely activated in 1985, a wide and flexible range of laser energies, pulse widths, pulse shapes, and wavelengths can be delivered onto a target. To fully utilize these capabilities, a number of diagnostic instruments were built and installed as part of the Nova construction project. The phase 1 instruments described in the next article, "Nova Diagnostics Summary," measure a number of optical and x-ray emissions from the laser-produced plasma. While similar to diagnostics in use on earlier laser systems,

these instruments have been upgraded and improved for use on Nova. The optical and electrical fiducial systems discussed later in this section in the articles "Nova Optical Fiducial System" and "Electronic Instrumentation Timing System" are designed for use with existing and future diagnostics. These systems now allow us to cross-time the diagnostic instruments and to time them absolutely with respect to the laser pulses from Nova. We are proceeding with the design and testing of advanced diagnostic instruments that will be essential in order to conduct the symmetry, hydrodynamic stability, and implosion experiments planned for the first few years of Nova operation, and to move eventually into the hydrodynamically equivalent target campaign.

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Nova Diagnostics Summary

We intend that Nova be the best diagnosed ICF research facility in operation today. Our experience in providing advanced diagnostics for previous laser systems will be extended at Nova, and will be challenged by the development of new instrumentation to diagnose the more advanced targets made possible by this powerful laser. Previous experience has shown that to understand target performance, we must have as complete a set of diagnostics as possible.

The Nova diagnostics are divided into two sets: the basic set required for the initial Nova experiments and the more advanced set for later, generally more complex, experiments. The basic set will be operational for the first Nova shots; it was a Nova line item funded with Nova construction money. The instruments were redesigned to be compatible with the Nova target chamber configuration. The electronic and vacuum systems were upgraded and automated, and computer controls were implemented. The more advanced set of diagnostics, examples of which are presented in the series of articles earlier in this section entitled "Diagnostic Development," will be completed by the LEAD Program during 1985 and 1986.

The 5-9 lists and describes the basic set of diagnostics. These instruments represent the culmination of our efforts to improve

Table 5-9. Basic diagnostics instruments.

Name and number of instruments	Detectors	Recording on	Description	Purpose
Dante 2	Diodes	Oscilloscopes	Dante is a 10-channel, absolutely calibrated, soft x-ray spectrometer that uses K-edge and L-edge filters and mirror casettes to define energy bins.	Dante measures the soft x-ray spectrum from 0.1–1.5 keV with a resolution of $E/\Delta E \approx 2-8$.
Filter fluorescer (FFLEX) 1	Fluor-photo-multiplier	Charge integrator	FFLEX is a 20-channel, hard x-ray spectrometer that uses K-edge filters and fluorescers to define energy bins.	FFLEX measures the suprathermal x-ray spectrum from 3–300 keV with a resolution of $E/\Delta E \approx 2-5$. This measurement allows us to determine the hot electron fraction and temperature.
Energy balance modules (EBM) 130	Photodiodes	Charge integrator	EBMs are absolutely calibrated optical energy bins. They are mounted at many positions at the chamber wall.	EBMs measure the reflected light, allowing us to calculate laser light absorption. By measuring the scattered light at several frequencies from $\omega/2$ to $3\omega/2$, we identify the processes responsible for reduced absorption and increased hot electron production.
Spatially discriminating streak spectrograph (SDSS) 2	X-ray streak camera	Film	SDSS is a pinhole-streak camera combination that provides 1-D space (100 μm) and time (20 ps) x-ray information simultaneously. This diagnostic can also be configured as a multichannel spectrometer with up to 8 energy bins defined by grazing incidence mirror and K & L edge filters. Optical fiducials provide correlation with the optical delivery pulse.	SDSS is used to measure x-ray emission in the soft x-ray region from specified points on the target as a function of time. It is also used to measure the x-ray spectrum from 200–1500 eV with a resolution of 100–300 eV.
Streak crystal spectrograph (SCS) 1	X-ray streak camera	Film	SCS is a high sensitivity, high energy and time resolution crystal-streak camera spectrograph. This diagnostic is reentrant.	SCS measures the x-ray spectrum from 0.6 to >10 keV with a resolution of 20 ps. It allows us to characterize temperatures and densities in a wide range of plasma conditions using relative x-ray line intensities and spectral line shapes.
Streak optical pyrometer (SOP) 1	Optical streak camera	Film	SOP images visible and UV light with a Cassegrainian telescope onto an S-20 streak camera.	SOP images 2500–7000-Å light from the target with a spatial resolution of 50 μm and a temporal resolution of 20 ps. It is used to measure preheat and shock breakout times.
Optical x-ray streak camera (OX) 1	Optical-x-ray streak camera	Film	OX is a combination optical and hard x-ray streak camera. Light filters are used to define optical energy bins, and K x-ray filters (and photocathode response) define x-ray energy bins. The signals are absolutely timed with an optical fiducial.	OX measures the time correlation of hard x rays (>3 keV) and optical signals from laser-plasma instabilities. This allows us to identify mechanisms for the production of hot electrons.
Streak optical spectrograph (SOS) 2	Optical streak camera	Film	SOS is a grating spectrograph-streak camera combination.	SOS measures scattered light from $\omega/2$ to $3\omega/2$ with a temporal resolution of 20 ps. From these measurements we identify processes responsible for reduced absorption and increased hot electron production.
8 $\times$ microscope 2	Film	Film	The 8 $\times$ microscopes use high resolution Kirkpatrick-Baez x-ray mirrors to image soft x-ray emissions onto film with a magnification of 8. Each microscope produces 4 images that are time-integrated.	These instruments measure x-ray images with a resolution of less than 10 μm in energy bands from 1.1–5 keV. They provide an excellent guide to where the laser energy is deposited.
Zone plate camera (ZPC) 1	Film	Film	The ZPC is a high sensitivity ($10^{-1}-10^{-2}$ sr) x-ray coded aperture imaging system featuring high spatial resolution ($\sim 5 \mu\text{m}$). Energy definition (3–100 keV) is achieved using x-ray filters. Images are time-integrated.	ZPC allows us to characterize the geometry of the compressed target capsule and thereby address questions of symmetry and mix. We can also characterize the spatial distribution of high energy x-ray emission and address questions of preheat and drive symmetry.
Neutron copper activation (Cu activation) 1	NaI photo-multipliers	Multichannel analyzer	The Cu activation system is a Cu disk placed near the target in a reentrant tube. The Cu is activated by 14.1-MeV neutrons. After the shot, the disk is transported to an automatic counting system consisting of a pair of NaI-photomultiplier gamma ray detectors.	The Cu activation system measures the number of fusion reactions that occurred in the implosion.
Radiochemistry (RadChem) 1	NaI photo-multiplier; intrinsic Ge	Multichannel analyzer	RadChem consists of a cylindrical collector placed near the target to collect target debris activated by neutrons. An automated system transports the collector to the counting laboratory in <30 s.	The radiochemistry diagnostic measures 14.1-MeV neutron activation products to determine areal density of various target components.

previously developed diagnostics with the most current technology available and to tailor them to our needs for Nova.

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Nova Optical Fiducial System

The Nova optical fiducial system provides reference timing information to the fusion diagnostic streak cameras mounted on Nova. Timing signals are distributed in the form of optical pulses at $0.35\ \mu\text{m}$. Our tests have determined that all diagnostic streak cameras, both optical and x-ray, are sensitive to the 3ω fiducial pulse, and a usable signal is generated in each camera.¹⁶⁸ Thus, we will have a single, absolute time base that relates all diagnostics. The modifications necessary to do so will be incorporated into the diagnostic software packages.

Fiducial timing pulses with $>500\ \text{nj}$ of energy and $<250\ \text{ps}$ long make the fiducial system compatible with the sweep rates of all Nova streak cameras. Because the fiducial system can be fired independently of the main Nova laser, it can be used to check streak camera operation without a target shot.

The Nova optical fiducial system is composed of two subsystems: an optical fiducial laser supplying a short 3ω light pulse and a fiber-optic distribution subsystem that delivers the light to various Nova diagnostics. The interface of the optical fiber to an x-ray camera is different from the interface to an optical streak camera. For x-ray streak cameras, the tip of the optical fiber is placed near the photocathode, where it serves as the electron source. For optical cameras, the light transmitted by the optical fiber illuminates the 5-1 or 5-20 photocathode. Tests of fiber-optic photocathodes showed that we needed to stringently specify the temporal shape, wavelength, repetition frequency, and energy of the fiducial pulse. We determined the specification for the fiducial pulse at the interface between the fiducial generator and the fiber-optic distribution subsystems by adding the effect of the losses of the fiber-optic distribution system. System specifications were also evaluated for compatibility with the Nova electronic trigger generator.

Previous tests established that only $100\ \text{nj}$ of 3ω light are required to generate a suitable signal to an x-ray camera. A fiber output energy of $500\ \text{nj}$ per pulse was established as a system requirement.

The temporal characteristics of streak cameras require that the fiducial pulse be no longer than $250\ \text{ps}$ with a well defined peak and that its separation from the main Nova pulse exhibit a maximum jitter of $50\ \text{ps}$. This standard yields absolute and cross-timing of streak cameras to an accuracy of $\leq 100\ \text{ps}$.

Three synchronization requirements are imposed on the optical fiducial system:

- Matching the Nova and fiducial pulse times.
- Coordinating the timing of the various cameras with one another.
- Synchronizing the electronic trigger with the optical fiducial.

The timing of the fiducial with respect to the Nova pulse is specified to be adjustable over $\pm 10\ \text{ns}$ from nominal Nova timing in steps of $50\ \text{ps}$ or less. The various cameras are cross timed with one another by timing them to the optical fiducial.

We also plan to synchronize the electronic trigger by timing it to the optical fiducial. This will allow tests of the trigger without firing the Nova laser.

Fiducial Laser Requirements. The optical fiducial design specifications that have a primary influence on the final beamline configuration are the fixed, short pulse length and timing, and sufficient 1ω energy to generate $5\ \text{mj}$ of 3ω energy. There is no point in the Nova baseline configuration where these requirements can all be satisfied. As a result, we have designed an independent beamline, essentially duplicating the short-pulse beamline configuration of the master oscillator room (MOR). That is, we will add a mode-locked, Q-switched oscillator and a Pockels cell for single pulse selection, followed by two 1-cm rod amplifiers and a frequency conversion assembly.

Timing the fiducial pulse within $50\ \text{ps}$ of the arrival of the target x-ray signal at the diagnostic streak cameras on the target chamber requires that we synchronize the optical fiducial oscillator with the main short-pulse oscillator. This is done by driving the two oscillators from the same rf signal generator. To permit this coupling, the two oscillators will stand side by side in the

Fig. 5-80. Nova optical fiducial layout.

MOR, followed by their respective switch-out Pockels cell chassises (Fig. 5-80).

The limitations on the transmission of 3ω light through optical fibers require that the fiber insertion point be no further from the north (10-beam) target chamber than the south wall of the target bay. We relay a low-energy (100-mJ) 1ω pulse from the MOR to this point, then amplify and convert it to 3ω .

A four-lens relay configuration is designed to place a $5\times$ image of the oscillator output window at the frequency conversion crystal near the target bay. The frequency converter itself will be mounted within the 4-ft-thick wall separating the switchyard from the target bay in a hole intended for beam No. 20 of the original 20-beam Nova design. This lends both mechanical and thermal stability to the frequency converter.

The optical fiducial amplifier contains several components required by the physical layout. (See Fig. 5-80.) A standard Nova chain input sensor will align the MOR beam through the amplifiers to the

frequency converter. Two rod amplifiers, separated by a Pockels cell, are necessary to provide between 20 and 100 mJ of 1ω energy required for the 5-mJ, 3ω specification. The third rod amplifier shown is included in the design to supply 1ω trigger pulses to x-ray diagnostics on the north target chamber wall.

After frequency conversion to 3ω , the UV light is separated from the residual 1ω and 2ω light with a fused silica prism. For north chamber shots, the 3ω pulse is inserted into a 61-fiber optical bundle for transmission to the diagnostic instruments. The other harmonics are presently unused in north chamber operation, but in the future, the unconverted light may drive photoconductive switches on advanced diagnostics.

Fiducial Output Coupling and Distribution. The fiducial distribution system (Fig. 5-81) is composed of a fiber-optic array that splits the fiducial beam into 61 independent optical fibers, each of which may be assigned to a diagnostic instrument. A down-collimating telescope presents a

Fig. 5-81. Nova optical fiducial distribution system.

Fig. 5-82. Fiber-optic array.

parallel beam of optimum size from the laser to the fiber-optic array. The diagnostic streak cameras are coupled to this array with 20 m of fiber. The length of each fiber is adjusted to deliver a properly timed fiducial to each streak camera.

The total attenuation from the fiber array to the streak camera cathode is determined by measuring the attenuation of each individual component. A multimode, 100- μm semi-graded index fiber was selected based upon the relatively high power levels required at the entry to the fiber array and on the measurements¹⁶⁹ of the attenuation, dispersion, and wavelength behavior of a variety of optical fibers. Attenuation data were obtained by removing 3 m of cable in six steps and recording the input and output characteristics at each step. Four such runs made using various lengths of cable show the attenuation to be less than 8 dB/18 m. This is an extremely high value compared to the 7 dB/km at 0.8 μm advertised by the manufacturer. This difference is explained by the poor transmission of 3ω light in fused silica.

A second measurement of attenuation was performed on two sets of fiber arrays (Fig. 5-82). We uniformly illuminated the front surface of the 61-fiber bundle with a pulse of 3ω light. The output from 3 of the 61 fibers was recorded and plotted against

the input energy, and the array attenuation was calculated from the slope of the input-output plot. The test was repeated using only four of the eight fibers, and the slope of the input-output plot increased by a factor of 2. This verified that the array was uniformly illuminated.

The last system components evaluated were the fiber-optic connectors. The manufacturer specifies the connector loss at 1 dB/connector; however, our measurements

show the attenuation to be 1.3 to 1.5 dB per connector.

Summing the attenuation values of the fiber-optic array, 20 m of fiber and 3 connectors, we arrive at a total attenuation of 41.5 dB.

We expect to obtain a 350-nJ output pulse at the streak camera cathode with a 5-mJ input beam. But to meet the 500-nJ output energy design, we must increase the input intensity at the fiber array, either by supplying more than 5 mJ of input energy or by further reducing the diameter (5 mm) of the input beam and eliminating some of the 61 fibers.

Dispersion measurements made with 0.606 μm light pulses through 20 m of fused-silica fiber indicate that a 250-ps pulse will broaden by only 70 ps. Tests are now in progress to evaluate the dispersion at 3ω .

From the results of the attenuation and dispersion measurements, we will be able to determine accurately the peak of the fiducial pulse. Then we can compare temporal data taken from separate diagnostics on the Nova experiments with a high degree of precision.

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Electronic Instrumentation Timing System

An electronic timing system that serves two important functions has been implemented on Nova. First, pulses can be generated at precise, controlled times relative to the Nova system trigger. These pulses simulate experimental signals without the need for laser shots. Transient recorder and streak camera triggers can now be timed independently of the laser system, greatly reducing setup time.

The second function of the Nova electronic timing system is to accurately cross-time signals on transient recorders. In past experiments, signals were cross-timed assuming common breakout times. Now, fiducial signals are common to all transient recorders allow them to be independently cross-timed to an accuracy of better than 50 ps.

The heart of the electronic timing system is a signal simulator combined with a high-resolution time interval meter (TIM). The signal simulator is a pulse generator with controllable amplitude, polarity, and trigger delay. A signal simulator monitor signal, independent of the output pulse amplitude and polarity, is used by the TIM to measure the delay to a precision of better than 10 ps. Software has been developed that determines the correct signal simulator delay for each signal channel, and produces that delay by a successive approximation technique. The correct delay is based on time-of-flight and jumper-cable delays that are part of a pre-loaded data file. The software includes an interface to a portable controller that operators use to control the signal simulator.

During instrumentation setup, the timing system causes the signal simulator pulse to arrive at the transient recorder at the correct time relative to the Nova system trigger. The operator then adjusts individual trigger times to obtain the desired display. The same procedure is used to set up streak cameras, except that a TIM reading is used instead of a displayed waveform. The TIM is started by a sweep monitor pulse from the streak camera and stopped by the signal simulator pulse. The streak camera is being triggered correctly when the TIM reading equals a calculated, pre-determined reading.

Cross-timing is achieved by mixing a common fiducial pulse with the signal on each transient recorder. This pulse is distributed via passive, jitter-free components. In order to account for different cable lengths, the fiducial pulses must be accurately timed at each recorder. The simulator and precision TIM are used to do this. The operator selects a signal channel and adjusts the signal simulator until the leading edges of the simulator output and the fiducial pulse are aligned. A TIM reading is taken and stored in a data file. This read operation is repeated for each channel to measure the relative timing of the various fiducial pulses. These readings are then used to accurately correlate transient recorder records.

The system includes a method to monitor the temporal performance of a streak camera. For each streaked diagnostic, a time-to-digital converter channel is used to record the time differences between the diagnostic trigger and the sweep monitor pulse from the streak camera. These measurements will

be routinely made on system dry runs and on target shots to reveal timing shifts that may occur.

The Nova electronic timing system provides a more efficient use of the laser facility by allowing most setup operations to be accomplished independently of the laser. For the first time, it makes possible accurate cross-timing of experimental data. And it is now possible, also for the first time, to detect a timing error or sweep malfunction in a streaked diagnostic before the target shot.

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Microwave Experiments

Microwave simulation experiments were performed at the University of California at Davis to study hot electron production due to parametric instabilities. In the experiments, s-polarized microwaves with a frequency $\omega_0/2\pi = 1.2$ GHz are incident^[70,171] on an inhomogeneous plasma in a special section of an L-band waveguide (8.25 × 16.5 cm). Therefore, a well-defined TE₁₀ mode of the microwaves interacts with the inhomogeneous plasma near the critical surface. The plasma is essentially collisionless because the electron mean free path is about 1 m, which is greater than the device dimensions. The plasma has an electron temperature $T_{e0} = 3$ eV, an electron-to-ion temperature ratio $T_{e0}/T_{i0} = 10$, and a critical density $n_c = 1.8 \times 10^{10}$ cm⁻³.

When the microwave power is above a well-defined threshold value, the simultaneous excitation of an ion wave ω_i and an electron plasma wave ω_e is observed. This process is identified as the ion acoustic decay instability. For microwave powers well above threshold, the electron plasma waves have a turbulent spectrum. The broad spectral width, shown in Fig. 5-83, is

$$\Delta\omega/\omega_i \geq 3\% \gg \omega_{pi}/\omega_i, \quad (21)$$

where the ion plasma frequency $\omega_{pi} = 4\pi n_i e^2/m_i$, e is the electron charge, and m_i is the ion mass. For high-power microwaves, the density profile near the critical surface is strongly steepened^[70,171] as a result of the ponderomotive force of the

plasma waves, and the ion waves also become turbulent.

Langmuir probe measurements are shown in Fig. 5-84. A straight line on this electron current-voltage characteristic indicates a Maxwellian energy distribution. The current of high-energy electrons increases with increasing microwave power. The suprathermal electron distribution has a cut-off at a relatively low energy (shown by the arrow on the 50-W curve) for weak microwave powers. The cutoff velocity

$$v_{cut}/v_{ph} \approx (v_{cut}/v_{e0}) k\lambda_{De} \approx 1.2 \quad (22)$$

Fig. 5-83. High-frequency spectrum of electron plasma waves with microwave power well above threshold.

Fig. 5-84. Electron current vs energy measured with Langmuir probe at four levels of microwave power. All electrons with energy above the abscissa value are collected.

Microwave Experiments

is on the order of the phase velocity of the most unstable waves.

The measured steady-state hot electron temperature is given approximately by $T_H = 66 P_0^{1/2}$ eV, where P_0 is the incident microwave power in kilowatts and the area is $\sim 64 \text{ cm}^2$. Thermal electrons are also strongly heated (Fig. 5-85). The thermal electron temperature is given approximately by $T_e = 9 P_0^{0.3}$ eV.

The measured ratio of hot to thermal electron temperature $T_H/T_e \approx 7.3 P_0^{0.1}$ is therefore a very slow function of microwave power. Computer simulation calculations (ZOHAR) give $T_H \approx 4.4 \times 10^{-4} (\lambda_0^2)^{0.3}$ keV for $10^{15} \leq \lambda_0^2 \leq 10^{18} \text{ W/cm}^2$, (the parameter regime of the experiment) and $T_e = 4 \text{ keV}$; λ_0 is the transverse wavelength in microns. Thermal electrons are only weakly heated in the computer calculations.

Fig. 5-85. Thermal electron heating T_e vs incident microwave power from three sets of measurements with the same plasma parameters.

Fig. 5-86. Ratio of hot and thermal electron temperatures vs $f_H \lambda_0^2 / T_e^{3/2}$.

An approximate theory¹⁷⁰ indicates that T_H/T_e should be only a function of the variable $f_H \lambda_0^2 / T_e^{3/2}$ where f_H is the fraction of energy into fast electrons. Figure 5-86 illustrates the microwave simulation results, ZOHAR computer simulation results for laser parameters, and the theoretical results¹⁷⁰

$$T_H/T_e \approx A \times 10^{-4} (f_H \lambda_0^2 / T_e^{3/2})^{1/3}, \quad (23)$$

where $A \approx 1.4$ for $T_e/T_i \approx 1$ and $A = 1.0$ for $T_e/T_i \gg 1$ with l expressed in W/cm^2 , λ_0 in μm , and T_e in keV. The results should be compared to theory for $A \approx 1.4$ in the microwave experiments with strong ion heating, and they should be compared to theory for $A \approx 1$ in the computer simulations with a large temperature ratio $T_e/T_i \approx 16$. We see good agreement (within $\pm 20\%$) between theory, microwave measurements, and simulation calculations for laser parameters. As expected, these results only apply to cases where the density profile is self-consistently determined by the ponderomotive force. Simulation calculations with shallow density profiles that are not in equilibrium give quite different results at the same laser intensities.

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Advanced research and development for each laser component ensures that current laser systems perform efficiently and that future laser system designs meet ICF goals. Continued research in amplifier designs and in optical components will increase the efficiency of the Nd:glass laser and will develop advanced solid state laser technology. Early high-average-power lasers will feature glass amplifiers, therefore, strength tests of glass are being conducted, strength being an important parameter in an amplifier's ability to withstand high-average-power laser operation. The large photo at left shows a side view of a high-energy fractured glass sample. The two inset photos at far left show a biaxial test jig (top) and a bend test jig (bottom), two devices used to characterize defect populations. Inset photo at page top shows a fracture initiating from a surface defect. The photo below, a gadolinium gallium garnet crystal, is, when doped with neodymium and chromium, one of the most efficient of any known solid state laser material. It is a candidate for high-average-power laser amplifiers.

Advanced Laser Development

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Section Editor: J. R. Murray

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Advanced Laser Development

Introduction

J. E. Hozzichter and W. H. Lowdermilk

The responsibility of the Advanced Laser Development Program is to develop and test concepts, components, and materials for present and future laser systems. These activities have three goals and objectives.

- To support the experimental program on Nova. During 1984, construction of the laser system was completed. Laser activation and preparations for the first target experiments began at the end of the year. A major portion of the Advanced Laser Development Program resources were dedicated to bringing the Nova laser to operational status. These activities included finishing the technology development of the harmonic-conversion arrays and of the high-power antireflective (AR) coatings used on all transmitting optics.
- To develop the technology and cost base for an affordable 5- to 10-MJ high-gain research facility to demonstrate and optimize high-gain ($\approx 100\times$) targets. Wavelength and temporal pulse-shape flexibility will also be required (see Table 6-1).
- To invent an efficient, repetitively pulsed high-average-power driver for various future civilian and military applications of ICF. This objective places particularly stringent demands on laser technology, as shown in Table 6-1. It is important to note that no demonstrated concept can meet this final objective. However, we have identified advances in solid-state laser technology that can permit 1- to 10-Hz laser systems to operate at $>10\%$ efficiency.^{1,2} We continue to support free-electron laser development by the Beam Research Group at LLNL because that technology also shows promise of high efficiency and repetition rate even though it is relatively undeveloped at present.

The 5- to 10-MJ, short-wavelength laser system must be both affordable and acceptable to a variety of users in the weapons, fusion, and general physics communities. The system must be flexible, applicable to a wide variety of target shapes, sizes, and pulse requirements, and should cost no more than $\sim \$250$ million. We have investigated both gaseous laser systems (primarily the KrF excimer laser) and harmonically converted Nd:glass laser systems for this application. We conclude that the most flexible and economical approach at present is to use harmonically converted Nd:glass technology.

Table 6-1. Future ICF Program laser-irradiation facilities.

Characteristics	High-gain ($\sim 100\times$) experimental facility	Reactor driver (1 GW _e)
Pulse energy	5 to 10 MJ	2 to 6 MJ
Peak power	500 to 1000 TW	200-600 TW
Pulse duration	0.030 to 100 ns	~ 10 ns
Pulse shape	Diverse, flexible	Single, fixed
Pulse repetition rate	Single shot	Few Hz
Efficiency	Unimportant except for cost impact	$\geq 10\%$
Wavelength	0.25 to 1.00 μm	Fixed λ , $\leq 0.50 \mu\text{m}$
Cost goal	25 to 50 $\$/\text{j}$	30 to 60 $\$/\text{W}_{\text{ave}}$ 100 to 150 $\$/\text{j}$

However, the construction of a 5- to 10-MJ, short-wavelength driver using Nd:glass at an acceptable price requires cost reduction in every area of laser technology.

Two major areas of cost reduction involve the amplifiers and the optical train. Compared to Nova, we believe that amplifier costs on a per-joule basis can be reduced by over a factor of 10. Savings include reducing laser-glass costs by a factor of 4 to 8, doubling extraction efficiency, increasing amplifier storage efficiency (stored laser energy relative to bank energy) from 1 to 4%, and reducing the cost of electrical energy storage from \$0.40 to \$0.10/J. Our objectives are to make better use of the laser optics (and thereby to reduce their cost per joule) by increasing the laser fluence from present levels of 5 J/cm^2 (an average at a 1-ns pulse duration) to approximately 20 J/cm^2 (an average at a 10-ns pulse duration). The goals of our materials program are to meet these objectives by using new materials such as sol-gels for AR coatings, as well as for high reflectivity and dichroic coatings, and by developing new techniques for more economical production of glasses and harmonic crystals.

Over the past decade, we have invented a variety of gaseous laser media directed at the high-repetition-rate, high-efficiency application. We are continuing this work; however, all short-wavelength gaseous laser media to date have failed to meet our requirements of either $>10\%$ system efficiency or acceptable cost per joule. For example, the most well developed short-wavelength gas laser uses the rare-gas excimer KrF. We have worked since its invention in 1975 to optimize the KrF excimer laser for the ICF application. However, it requires expensive pulse power (\$1 to \$2/J) based on megavolt electron-beam pumping, its characteristic wavelength of $0.249 \mu\text{m}$ limits the optical fluence to $\sim 5 \text{ J/cm}^2$, and the nonstorage character in the laser medium requires the use of pulse-compression techniques to shorten the $>100\text{-ns}$ amplifier extraction pulse to the 10-ns pulse needed for target implosions. These operating constraints preclude KrF laser technology from being cost competitive for a single-shot research facility when compared to harmonically converted Nd:glass laser technology optimized for minimum capital cost. In addition, we estimate that the ultimate system efficiency of KrF is about 5%, which unfortunately also precludes this system from being used for direct generation of electricity with expected $\sim 100\times$ target gains.

Free-electron laser technology holds promise of high efficiency (~ 15 to 20%) and very high average power at short wavelengths. However, the technology is too poorly developed to be a serious candidate for the next fusion facility. The development is much too expensive for the ICF Program to support, given current priorities. Fortunately, free-electron laser development is being supported in the Beam Research Group at LLNL by the Department of Defense (DOD). We have provided support during the project's beginning, and we continue to follow their work closely.

Our failure to invent new fusion laser systems caused us to reconsider the use of solid-state lasers as fusion drivers and, more generally, for military and commercial applications.^{1,2} While the success that we have experienced using the Nd:glass laser for target research has been gratifying, the extension of this technology into these new operating regions is indeed exciting. Our conclusions are that

1. Solid-state laser systems have as much potential for scaling to high average power (megawatts) as do known gas lasers.
2. Waste heat can be removed in a straightforward fashion by using gas-flow-cooled, thin-disk amplifier structures (at the MW/m^2 rate).

3. Solid-state lasers suitable for fuel-breeding ICF applications (few percent efficiency, few Hz repetition rate) can be constructed with modest extrapolation from existing Nd:glass laser technology.
4. Solid-state laser systems with efficiencies in excess of 10% appear to be achievable using new pumping sources (electron-beam-excited phosphors or semiconductor lasers) and advanced solid-state materials.

As a result of these studies, we have begun a technology program to ascertain the efficiency and average power capability of solid-state laser systems. Because of the immediate applicability of this technology to DOD problems, the Defense Advanced Research Project Agency (DARPA) has begun a program at LLNL to develop advanced solid-state laser technology for their specific applications. These activities will be reported in a separate document. The synergism between ICF and DOD applications is enhancing progress in both areas.

Summary of Activities

During 1984, we performed experiments directed at the 5- to 10-MJ high-gain, research laser goal. We increased our understanding of the important variables that affect the flashlamp pumping efficiency of amplifiers. The variables we studied included various reflector shapes, flashlamp packing fraction, and single- vs. double-sided pumping. We continued studies on flashlamp performance as a function of gas fill and bore diameter and on using enhanced feedback to the lamps of light with wavelengths not directly absorbed by the Nd-doped glass. Combining these factors resulted in a substantial increase in overall energy-storage efficiency of amplifiers ($\sim 4\%$ at 0.2 J/cm^3).

This year, we successfully completed development of sol-gel AR coatings for surfaces of optical elements exposed to $0.35\text{-}\mu\text{m}$ (third-harmonic or 3ω) light. Damage thresholds of sol-gel AR coatings deposited on fused-silica focusing lenses were equal to the damage threshold of uncoated fused silica. This damage threshold is 3 to 4 times higher than thresholds of vacuum-deposited, multilayer dielectric AR coatings. In addition, we developed a sol-gel AR coating that can be applied directly to the potassium dihydrogen phosphate (KDP) crystals used for harmonic frequency conversion on Nova. These coatings eliminate the need for cells filled with index matching fluid (IMF) to contain the KDP crystals. Eliminating such cells was desirable because they presented a higher risk of damage through laser-induced breakdown of the fluid.

We have continued development of improved Nd-doped laser glass. This year, we formulated and optimized new glass compositions for long fluorescence lifetime and low concentration quenching. In addition, we measured the effects of impurities on the Nd excited-state lifetime.

Large regenerative-amplifier systems are expected to have cost and maintenance advantages over current master-oscillator, power-amplifier (MOPA) laser designs. However, such systems require a large-aperture optical switch for which there were no previously known practical concepts. We have continued development of our plasma-electrode Pockel's cell and successfully demonstrated highly uniform switching of a $27\text{-} \times 27\text{-cm}$ cell with rise times of less than 50 ns.

We continued our systematic search for improved crystals for harmonic generation with high efficiency and wide dynamic range. We established experimental capabilities to screen materials for high nonlinear coefficients and noncritical phase matching. New solid solutions of KDP isomorphs have shown promise. In addition, we have examined the conditions under which currently used KDP crystals could be grown more rapidly to reduce cost. We built an experimental growth station of new design and achieved growth rates that are 10 times higher than those currently achieved in commercial practice.

We have continued our search for candidate laser systems for both the low-cost, single-pulse research system and for the eventual ICF commercial-power-generation system. New crystal-line-laser host materials have been grown, and their spectroscopic

Solid-State Amplifiers

and laser properties were measured. With one material, neodymium and chromium co-doped into crystals of gadolinium scandium gallium garnet (GSGG:Nd,Cr), we achieved an overall lasing efficiency of 5%, which is approximately 2 times higher than that of YAG:Nd.

In research that is separately funded by the Department of Energy (DOE). Office of Basic Energy Sciences, we have continued the development of computational methods to understand transition-metal ion impurities in crystalline host lattices. This program is directed toward developing the ability to predict systems with good properties as high-efficiency lasers. We have also continued our studies of laser-induced damage to surfaces and have successfully identified impurities that we believe initiate surface damage.

Solid-State Amplifiers

Single-Segment Amplifier: Experimental Results

Introduction. As discussed in the 1983 Laser Annual,³ we have constructed an extremely flexible Nd:YAG disk amplifier, which we call the single-segment amplifier (SSA). The SSA has rectangular laser disks and could represent a single segment of a

large segmented-array disk amplifier. However, its primary purpose is to provide a testbed for exploring amplifier designs and to produce fundamental data that can be compared to results from modeling. The experiments discussed in "Flashlamps," later in this section, and in previous Laser Annuals^{4,5} provide important input to such models.

Our major emphasis in the SSA is to maximize energy-storage efficiency, the ratio of laser energy stored in the glass at peak gain to the energy delivered to the lamps by the capacitor bank. In contrast, the Nova amplifiers were designed to maximize the peak gain coefficient and thereby minimize the effect of nonlinear self-focusing when used in a system. A megajoule-class laser driver will operate at much longer pulses so that amplifier efficiency also becomes an important parameter for maximizing the performance-to-cost ratio.

The main features of the SSA are illustrated in Fig. 6-1. The laser disks are rectangular to maximize the area of glass available for pumping and to be compatible with a segmented geometry. The nominal aperture of the amplifier is 32 cm on each side. We have used 4% Nd-doped silicate glass disks (LG-660 from Schott Glass Technologies, Inc.), which are 4.3 cm thick. The disks have a relatively large doping-thickness product to enhance the absorption of pump light. The baseline laser disks in the Nova 31.5- and 45-cm amplifiers are

Fig. 6-1. The SSA with black end panels and side panels used for measurements of storage efficiency.

also 4.3 cm thick; however, the Nd doping is 2%. The edge cladding, which provides absorption of 1- μm light, is constructed of KG-3 filter glass plates (Schott Glass) imbedded in clear silicone rubber with an optical index that matches both the laser glass and the absorbing plates.⁶ The internal reflectors in the SSA are plated with vacuum-deposited silver. The reflectivity of these reflectors is 90 to 93% at 0.53 μm compared to 85 to 90% for the electroplated silver coatings used in the Nova amplifiers. The SSA is four disks long to minimize end losses, and the arc length of the flashlamps is 234 cm.

A selection of end and side panels that can be used with the SSA are also shown in Fig. 6-1. The choice of end panels includes the standard bead-blasted aluminum ends (used in Nova testing), flat silver reflectors, or black anodized-aluminum boxes, as shown in the figure. The black boxes, which are about 30 cm deep, act as a light trap and provide a large area for resisting damage from the intense pump light in the amplifier. The number of lamps per side and the type of reflector behind the lamps can be varied. These silver-plated reflectors can have special shapes or can be flat. In some cases, we use black-absorbing side panels, which are 30-cm-deep troughs of black anodized aluminum. The black troughs can be used with or without flashlamps on a given side. Finally, the amplifier can be assembled with or without Pilkington glass blast shields between the laser disks and the flashlamps.

The three flashlamp reflectors we studied in 1984 are the 5-lamp Hagen, the 5-lamp Winston, and the 8-lamp cylindrical reflectors, as shown in Fig. 6-2. In cross section, the 5-lamp Hagen and 8-lamp cylindrical reflectors have a constant radius, and the curves extend to the front surfaces of the lamps (preventing the lamps from directly illuminating each other from side to side). The minimum clearance between the back of the lamps and the reflector surface is ≈ 0.1 cm. In the Winston design, any light ray leaving the surface of the lamp exits the reflector without passing back through the lamp, and vice versa (all incoming light rays pass through the lamps).⁷ We conclude that the Winston design should be ideal in terms of using light from all sides of the

lamps, directing light toward the disks, and feeding unused pump light back into the lamp plasma after passing through a disk. Our conclusion is, however, tempered by the fact that Winston reflectors spread light over large angles (up to $\pm 90^\circ$), and the light may require several reflections before striking a disk. Moreover, the efficiency of optical recycling (the conversion by flashlamps of absorbed light to optical output) is somewhat uncertain. The relative benefit of absorbing unused light in the plasma, rather than passing light back through the disk, remains unclear.

For the Winston shape, the width of the reflector channel is πd , where d is the diameter of the lamp. Hence, with the packing fraction of flashlamps associated with that shape ($\pi^{-1} = 0.32$), the area of the surface of the lamps exactly equals the area of the projected plane of the disks through the center of the amplifier. From conservation of brightness, we expect this lamp-packing fraction to be optimum with regard to amplifier efficiency. The brightness (in $\text{W} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1}$) of light that can be achieved on the pump plane by any reflector system cannot exceed that which is emitted from the surface of the lamps, given blackbody behavior. With equal areas for emitter and absorber together with an appropriately designated reflector, we expect to produce the maximum pumping at the maximum efficiency.

We have compared, under similar conditions, the performance of the SSA with that of the best Nova amplifiers. At energy-storage densities of interest ($\sim 0.2 \text{ J/cm}^3$), the efficiency of the SSA is greater by about a factor of 2. We have studied the efficiency of the SSA, operated with five lamps per side, as a function of electrical pulse width.

Fig. 6-2. Cross sections of three lamp reflectors used on the SSA.

(a) 5-lamp Hagen reflector

(b) 5-lamp Winston reflector

(c) 8-lamp cylindrical reflector

We have also measured the efficiency of single-pass pumping to provide a simple case for comparison to theory. The difference in pumping efficiency from the two sides has given us information on the angular distribution of pump light and on the loss in pumping caused by end loss. Moreover, we have measured the benefit of two-pass pumping relative to one-pass pumping and have examined the mutual effect when lamps are used on both sides of the amplifier. By comparing amplifier single-pass pumping efficiency with a given reflector shape to the case in which a black side panel is used, we have obtained a factor representing reflector enhancement for the various shapes. The effect of reabsorption of pump light by the lamps after reflection is clearly evident in these measurements. Finally, we have measured the effect of the blast shields on pumping efficiency for several conditions.

Measurement Techniques. We determined the performance of the SSA by measuring gain as a function of time using a cw Nd:YAG probe laser. The measured gains were converted to stored-energy densities and to efficiencies. We have emphasized measurement accuracy and repeatability to evaluate similar amplifier configurations and to reduce the need to repeat measurements.

Figure 6-3 shows a schematic of the setup for gain measurements, which is similar to that described previously.³ An etalon internal to the Nd:YAG resonator allowed us to choose the Nd:YAG transition at $1.061 \mu\text{m}$, which is near the peak cross section of

LG-660 laser glass at $1.057 \mu\text{m}$. The diffraction grating in the reference arm provided a continuous check on the output of the probe laser at the chosen wavelength. We monitored the power of the probe laser before transmission through the amplifier (reference signal) as well as after (transmitted signal) to account for fluctuations in the probe intensity. A diffraction grating was used in the transmitted-signal beamline to increase spectral discrimination against flashlamp light. The combined spectral and spatial filtering in the transmitted-signal beamline reduced flashlamp light leakage to less than 0.1% of the gain signal. For the reference beamline, spectral filtering alone reduced flashlamp background to less than this level. A TV camera was imaged on the spatial-filter pinhole, the smallest alignment point in the system, to provide continuous confirmation of alignment.

For accuracy and reproducibility, the reference and transmitted signals were digitized by Tektronix 5223 oscilloscopes. Using an LSI-11/23 computer, we divided the transmitted waveform by the reference waveform, normalized the result, and smoothed around its peak using the Savitsky-Golay method⁸ to find the peak gain. Figure 6-4 shows the standard deviations for peak amplifier gains obtained in this manner and plotted as a function of gain. The standard deviations were obtained from 6 sets of gain measurements taken on the SSA over a 9-mo period using our standard amplifier configuration: two 5-lamp Hagen reflectors, blast shields,

Fig. 6-3. Schematic of the setup for gain measurements shows an exaggerated top view of the four disks in the SSA to illustrate orientation and define the plus (+) and minus (-) sides of the amplifier.

bead-blasted aluminum end plates, and a flashlamp pulse width of $3T = 655 \mu\text{s}$. The least-squares fit to the data is also shown. Since we typically do not take enough shots at one condition during a run to determine errors accurately, we determined the error in the (average) measured gain by dividing the standard deviations shown in Fig. 6-4 by the square root of the number of shots taken for that data point.

The standard deviation remains finite as the gain approaches unity, presumably because it is dominated in that case by digitizing errors. For small gains, we passed the probe beam through the amplifier twice, thereby squaring the gain and reducing the relative error. Typical relative errors in gain coefficients or storage efficiencies in our data varied from $\pm 3\%$ at a gain of 1.04 to $\pm 0.3\%$ at a gain of 3.

We have seen small but consistent decreases in gain as the Nova amplifiers and the SSA heat up. We monitored the surface temperature of an end disk in the SSA with a thermocouple and found that heating the amplifier externally with heater tapes produced an effect comparable to firing the flashlamps. For LG-660 glass in the SSA, the gain coefficient decreased by 0.2% per $^{\circ}\text{C}$ referenced to the surface temperature of the end disk. Because a typical set of gain measurements taken over one day increased the temperature by only a few degrees, the temperature corrections were usually less than or comparable to our error bars. Therefore, we did not make temperature corrections except in those cases for which many shots at large energies were necessary or when a data set spanned more than one day.

We periodically repeated gain measurements with our standard configuration to test for aging effects in the amplifier. Remarkably little change occurred over a 10-mo period. Storage efficiencies, calculated from data sets separated by over 1500 shots, remained within error bars for inputs between 100 and 300 kJ, and they decreased only about 3% for lower inputs. These results indirectly demonstrate the survivability of the silicone-rubber-based edge claddings developed specifically for the SSA.

The most important quantities determined from our measurements are the peak

Fig. 6-4. Experimentally measured standard deviations from gain measurements of the SSA in the standard configuration.

gain coefficient α and the peak storage efficiency η . The peak storage efficiency is the total energy stored in the amplifier at peak gain divided by the total energy initially stored in the capacitor bank. We determined these two quantities from centerline gain G as follows:

$$\alpha = \frac{\ln G}{mt(1 + 1/n^2)^{1/2}}, \quad (1)$$

and

$$\eta = \frac{\alpha h \nu C_s V}{\sigma_g E_p}, \quad (2)$$

where

- m = number of disks in the amplifier,
- t = disk thickness,
- n = refractive index of laser glass,
- h = Planck's constant,
- ν = frequency of peak gain,
- C_s = spatial-uniformity correction factor to account for the difference between the centerline gain coefficient and the average gain coefficient over the full clear aperture of the amplifier,
- V = total volume of laser glass accessible to a laser beam,
- σ_g = gain cross section of the glass at the measurement wavelength, and
- E_p = initial energy stored on the capacitor bank.

For the SSA with LG-660 disks, $m = 4$, $t = 4.3 \text{ cm}$, $V = 34.9 \text{ liters}$, $n = 1.52$, $\nu = c/1.057 \mu\text{m}$, and $\sigma_g = 1.81 \times 10^{-20} \text{ cm}^2$. To determine σ_g we used the cross section at the center of the line ($1.93 \times 10^{-20} \text{ cm}^2$ at

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1.057 μm) and assumed that it scales with wavelength as $\lambda^{5/3}(\lambda)$, where $I(\lambda)$ is the measured fluorescence line shape.⁹

The uniformity correction C_u was determined from gain measurements at three spatial positions. The SSA disks rested on edge at Brewster's angle with respect to the horizontally polarized input beam, as shown in Fig. 6-3. We typically measured gain at the +95 and -95% points of the clear aperture in the horizontal plane as well as along the centerline. The plus and minus sides of the amplifier are the high- and low-gain sides, respectively, as discussed previously.³ Detailed measurements of the gain profile of the SSA showed flat profiles in the vertical direction and asymmetric parabolic profiles in the horizontal direction, similar to those for the Nova amplifiers.³ We assumed the same profile was maintained over the full range of bank energies and calculated C_u for a given data set from the parabolic profile determined by the three gain values at a charge of 20 kV. Values for C_u varied from 0.96 to 1.00.

Experimental Results. To obtain a direct comparison with the amplifier performance

Fig. 6-5. Storage efficiency of the SSA and the two largest Nova amplifiers as a function of stored energy density.

Fig. 6-6. Measured ASE derating factors for the SSA and the 20.8- and 31.5-cm Nova amplifiers compared to a GAINPK6 calculation.

of Nova, the SSA was first tested in our standard configuration, with blast shields and bead-blasted aluminum end plates.¹⁰ The 5-lamp Winston reflectors were employed, and the lamps were excited by a 655- μs ($3T = 3\sqrt{LC}$) circuit, with $C = 106\ \mu\text{F}$ and $L = 450\ \mu\text{H}$. Figure 6-5 shows the measured peak storage efficiency vs stored energy density for the SSA and the Nova 46- and 31.5-cm amplifiers. Such comparisons based on stored energy density are appropriate because all three amplifiers have 4.3-cm-thick disks. The SSA is roughly twice as efficient at a given energy density as the best Nova amplifier, and its fall-off in efficiency with loading is much reduced. We attribute these improvements to the larger disk area (rectangular vs elliptical), greater Nd doping, optimized lamp-packing fraction, higher cavity reflectivity, reduced end losses, and smaller loss in gain caused by amplified spontaneous emission (ASE).

We measured the loss in peak gain caused by ASE in the same manner as for Nova amplifiers—by using a peak-fluorescence detector (PFD).³ Our PFD technique uses pump light from inside the cavity to excite a small sample of 4%-doped LG-660 amplifier glass located outside the laser cavity. The effect of ASE depumping of the small external sample is negligible. Thus, α/PFD (peak gain coefficient α divided by peak fluorescence signal) gives a direct measure of the loss in gain caused by ASE depumping. To provide an angle-averaged sample of pump light to the PFD, a small quart bulb filled with quartz shards scattered a sample of pump light into a fiber-optic bundle. The bundle, encased in a stainless steel tube, and the quartz bulb were inserted into the amplifier to provide light that excited the PFD, which was located outside the amplifier.

Figure 6-6 shows the measured α/PFD vs αD_{eff} for the SSA and for the 31.5- and 20.8-cm Nova amplifiers.³ The three data sets were normalized individually by least-squares fitting of each to a second-order polynomial and scaling α/PFD to make the fit equal unity at $\alpha D_{\text{eff}} = 0$. We finetuned each normalization by as much as the error in the zero intercept of the fit to obtain the best fit of data sets to the

Fig. 6-7. SSA storage efficiency for four flashlamp pulse widths with circuit parameters chosen to give equal damping factors at equal input energies.

GAINPK6 curve in Fig. 6-6. We used an effective elliptical disk size of $1.12 \times$ the long SSA disk dimension for D_{eff} to compare SSA data (rectangular disks) to the data for elliptical disks of the two Nova amplifiers.¹¹ Figure 6-6 shows that ASE caused, at most, an 11% drop in stored energy in the SSA, compared to the 33% drop in the 31.5-cm amplifier. The solid curve in Fig. 6-5 was subsequently used to remove the effect of ASE when many of our amplifier measurements were analyzed.

To optimize SSA performance, we measured its gain with four different circuits having $3T$ values of 274, 436, 655, and 922 μ s (Rel. 12). Values for inductance and capacitance were chosen to give the same circuit-damping factor at the same bank energy. Figure 6-7 shows that shorter pulses substantially improve the SSA performance at low energy densities but help little at stored energy densities of interest (~ 0.2 J/cm²). This behavior follows from the change in opacity of the flashlamps with input energy. At low stored energies, for which the lamps are lightly loaded and relatively transparent, the shorter pulse lengths obtained with smaller $3T$ values compete favorably against spontaneous decay. At higher stored energies, for which the lamps are driven harder, their increased opacity at shorter pulse lengths begins absorbing useful pump photons, thus reducing the improvement.

Single-Pass Pumping. Single-pass pumping is the simplest SSA experiment to model using fundamental flashlamp data.¹³ In the ideal case, we pass all pump photons

emitted by the flashlamps through the disks once. This simplified geometry eliminates the modeling difficulties associated with multipass pumping and feedback of light to the flashlamps. To approach the ideal case, we use flashlamps and a reflector on one side, a black-absorbing panel on the opposite side, and black end plates. To minimize Fresnel reflection, we removed the blast shields; however, laser disks themselves are Fresnel reflectors. Moreover, some of the reflectors are not ideal in that some reflected light passes back through the lamps and can be reabsorbed.

Figure 6-8 shows the resulting storage efficiency vs bank energy for single-pass pumping. The flashlamps with 5-lamp Hagen reflectors were mounted either on the minus or the plus side of the amplifier. Similar curves are shown for the black reflector. In this case, black absorbers were placed on both sides and on the ends to create an "all-black" amplifier.¹⁴ In general, higher efficiency is obtained for pumping on the minus side because the light must pass through the disks before reaching an end. For pumping on the plus side, some fraction of the light can pass directly out the ends. As discussed below, single-pass pumping efficiencies with reflectors are about 70% of ideal predictions. The difference is primarily due to losses in the silver reflectors. Finally, the presence of reflectors behind the flashlamps enhances pumping by the expected factor of about 2.

The difference in single-pass pumping efficiency for lamps either on the plus or the minus side provides a convenient method

Fig. 6-8. Single-pass storage efficiencies with 5-lamp Hagen reflectors or black reflectors showing the differences between pumping from the minus and plus sides of the SSA.

Fig. 6-9. Measured end loss for the 5-lamp Hagen, Winston, and black reflectors vs initial bank energy to one side of the SSA. ▶

Fig. 6-10. Spatial uniformity for single-pass pumping from (a) the plus and (b) the minus side of the SSA using 5-lamp Hagen, Winston, and black reflectors.

for measuring the end loss. The ratio of efficiencies produced by pumping from the two sides provides a direct measure of end loss for pump light on its first pass across the amplifier. Figure 6-9 shows the average end loss for two-sided pumping. This average loss factor is represented by

$$\frac{1}{2} \left(\frac{M}{\bar{M}} + \frac{\bar{M}}{M} \right) \quad (3)$$

where M represents the orientation of the four disks, and the bar identifies the side where the flashlamps were mounted.

Figure 6-9 shows that the end-loss is somewhat different for Winston, Hagen, and black reflectors, all with five lamps per side. While the end-loss factor is roughly constant with loading for all three reflector shapes, the loss for Winston and black reflectors is about 8%; that for the Hagen reflector is about 5%. This difference suggests that Hagen reflectors produce some collimation, directing light toward the disks. This difference in collimation can be visualized in two dimensions, looking down at the amplifier, and with the disks arranged as shown in Fig. 6-3. We refer to the

angular distribution of pump light seen from this view as the projected longitudinal distribution. Only the projection of a light ray, as seen from above, determines whether the ray hits a disk or exits the end. The end-loss experiments suggest that the shape of lamp reflectors changes the projected distribution.

The transverse (horizontal) distribution of gain for single-pass pumping from the two sides and with different lamp reflectors supports the qualitative difference just described. Figure 6-10 shows the transverse gain distribution for pumping from the plus [Fig. 6-10(a)] or minus side [Fig. 6-10(b)] with the three types of reflectors. Greater variation across the aperture is seen for pumping from the plus side because the short end sections of lamps used with plus-side pumping radiate directly to the ends, and the gain falls off more rapidly with distance. The fall-off across the aperture is greater for the Winston and black reflectors than for the Hagen reflector. Once again, this outcome suggests that the Hagen reflector produces greater collimation in the projected view.

The reduction in peak storage efficiency caused by blast shields is plotted in Fig. 6-11 for the 5-lamp Hagen reflectors. The effect of blast shields is greater on the plus side, where some of the reflected light is apparently coupled out the ends. We find only a weak indication that the loss due to blast shields increases with loading. A roughly constant blast-shield loss of 6% was also observed with two-sided pumping.

The magnitude of the effect from blast shields and its constancy with loading are difficult to explain. We have calculated that there should be a 16% reflection of light by the blast shield, assuming a Lambertian input and two interfaces. That this reflection produces only a net loss of 6% suggests either flashlamp recycling of light is very efficient or a majority of the reflected light is not absorbed in the lamps. There is also a slightly beneficial effect of the blast shields, which helps to explain our results. Blast shields refract light rays toward the normal during passage, and the lamps appear to be somewhat closer to the disks.

A factor representing reflector enhancement can be derived by dividing the one-sided pumping efficiency observed with a given reflector by the efficiency observed with a black reflector, as shown in Fig. 6-8. Figure 6-12 shows this enhancement factor for the case of minus-side pumping (to eliminate end-loss changes). The top curve

Fig. 6-11. Measured influence of blast shields on storage efficiency for single-pass pumping of the SSA using 5-lamp Hagen reflectors.

Fig. 6-12. Reflector-enhancement factors for minus-side pumping without blast shields.

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(see curve for Winston without ASE) shows the effect of removing ASE. Corrections of the same order of magnitude would apply to the remaining curves. The enhancement factor is roughly 2, as expected from turning around one-half of the pump light. The observed decline with input energy is presumably caused by the reabsorption of pump photons into the lamp after reflection. On this basis, the enhancement factor for the Winston reflector should be constant. Moreover, the substantial decline for the 8-lamp cylindrical reflector illustrates the problem of a large lamp-packing fraction.

Multipass Pumping. To study two-pass pumping, we used 5-lamp Hagen reflectors on one side of the SSA and a flat, silver reflector on the other side (instead of the black absorber used in the one-pass case). In fact, the pumping is *not* simply two-pass because further passes can occur following reflection from the side with lamps. Figure 6-13 shows the ratio of two-pass to one-pass pumping with Hagen reflectors, black ends, and no blast shields. This ratio is about 1.18 and varies little with input energy. Given that the lamps are quite transparent at low loading and absorbing at high loading, the increased efficiency is caused

Fig. 6-13. Ratio of two-pass to one-pass pumping from the plus and minus sides of the SSA using 5-lamp Hagen reflectors.

Fig. 6-14. Comparison of one-sided and two-sided pumping with and without ASE using 5-lamp Hagen reflectors.

Fig. 6-15. Comparison of SSA performance for two-sided pumping without blast shields.

by multiple-pass effects at low loadings and by multipass effects combined with lamp absorption and reemission at high loadings. The effects apparently trade off nearly equivalently for the SSA because the improvement factor is roughly constant.

When lamps are fired on both sides of the amplifier, as in actual laser systems, the situation becomes complex and difficult to analyze. The effects of multiple-pass pumping vs lamp absorption and reemission are illustrated in Fig. 6-14. These data were taken with flashlamps and 5-lamp Hagen reflectors on both sides, black ends, and the 655- μ s circuit.

The curve for one-sided pumping corresponds to the case of firing lamps on only one side and is the average of the plus- and minus-side pumping cases. The curve for two-sided pumping corresponds to the case of firing lamps on both sides. The two-sided case is closely related to the two-pass case discussed previously, except that it is influenced by lamp absorption and reemission after each pass instead of every other pass. The curves in Fig. 6-14 are essentially identical at low loadings, as they should be when lamp absorption is low. At the greatest loadings, the results for one-sided pumping are greater (even after ASE is removed), suggesting that the second pass through our disks enhances the pumping more than the effect of lamp absorption and reemission. In the region of intermediate loading, the opposite relation appears to be true. However, such conclusions are quite sensitive to corrections for ASE. Nonetheless, we conclude that in the SSA, the benefit in efficiency of additional passes through the glass is closely matched to the benefit of lamp absorption and reemission after a single pass.

Figure 6-15 summarizes the performance for two-sided pumping of the SSA with our three reflectors. These data were taken with black ends, a 3T value of 655 μ s, but without blast shields. As in Fig. 6-5, the efficiency values were computed relative to stored bank energy and are plotted vs peak stored-energy density. The Winston reflector produces the best storage efficiency at high energy density,¹⁵ while the 5-lamp Hagen reflector provides the highest storage efficiency at lower energy densities. The efficiency with 8-lamp cylindrical reflectors is substantially lower at all energy densities, confirming that its lamp packing is larger than optimum. The differences in the Hagen and Winston reflector data are very small. The measured axial gain coefficients of the Hagen reflector are, in general, higher than for the Winston reflector, especially with blast shields. Thus, the performance differences are actually contained within the differences in the correction factor C_c for pump uniformity.

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Single-Segment Amplifier: Models and Calculations

Neodymium Absorption Cross Section.

Neodymium (Nd) absorption cross sections in various glasses are generally determined from the optical density of 5-mm-thick samples.¹⁶ For such thin samples and at typical values for Nd ion densities (10^{21} to 10^{23} cm⁻³), the absorption in the wings of the major bands is weak. Consequently, the relative error in the calculated cross section

Fig. 6-16. Predicted absorption cross section for LG-660 glass.

due to baseline subtraction becomes significant. The error is amplified when data on cross sections are used to predict the spectral absorption in thicker pieces of glass (e.g., 4.3 cm in the SSA) and can lead to inaccuracies in the calculated pumping efficiency. To eliminate such problems, we have analyzed transmission data for both 5-mm and 5-cm glass samples.

Figure 6-16 shows the spectral dependence of the absorption cross section of LG-660 glass, inferred from the transmission spectrum of a 5-cm-thick, 4%-doped sample (Nd-ion doping-thickness product = $1.90 \times 10^{21} \text{ cm}^{-2}$), and that inferred from the transmission of a 5-mm-thick, 2%-doped sample (doping-thickness product = $9.52 \times 10^{19} \text{ cm}^{-2}$). Just as the optical properties of a thin piece of glass give no information about the absorption in the wings of the bands, it is clear from Fig. 6-16 that a thick sample provides no values for the magnitude of the cross section around the peaks of the spectrum. To predict the absorption properties accurately over a broad range of glass thicknesses, we must use an absorption cross section that is a weighted average of both sets of data.

We determined best-fit values for the cross section over the entire spectral range by a least-squares fit to the transmissions of

both the thin and thick glass samples. Assuming a constant uncertainty in the transmission measurement, we determined the cross section σ at each wavelength by minimizing the error

$$\Delta^2 = [T_1 - \exp(-\sigma\rho_1d_1)]^2 + [T_2 - \exp(-\sigma\rho_2d_2)]^2, \quad (4)$$

where T_1 is the transmission of the thin sample, and ρ_1d_1 is its doping-thickness product. The values of T_2 and ρ_2d_2 are defined similarly for the thick sample. Requiring that $d(\Delta^2)/d\sigma = 0$ results in a transcendental equation, which we solved numerically for the least-squares prediction of the cross section at each wavelength.

The best-fit cross section obtained by this technique is plotted as a function of wavelength in Fig. 6-16. In the limit of saturation of the thick sample ($T_2 \rightarrow 0$), the calculated cross section approaches $-\ln T_1/\rho_1d_1$; conversely, in the regions where the thin sample is transparent ($T_1 \rightarrow 1$), the predicted cross section approaches $-\ln T_2/\rho_2d_2$. We estimate that our cross-section values are accurate to within $\pm 3\%$ over a dynamic range of 100.

Single-Pass Pumping. We used the fitted values of the cross sections to generate the absorption spectrum at normal

incidence of the 4.3-cm-thick, 4%-doped LG-660 slabs (doping-thickness product = 1.64×10^{21} ions cm^{-2}) used in the SSA. The instantaneous pumping efficiency for Nd:glass by 1.5-cm-bore xenon flashlamps was obtained by convolving the predicted absorption with the previously measured and absolutely calibrated lamp-emission spectra at the peak of the output for a wide range of lamp electrical inputs.¹⁷ In Fig. 6-17, the Nd*-pumping efficiency is defined as the ratio of extractable laser power at $1 \mu\text{m}$ to electrical power deposited in the lamp. We assumed that each photon passed through the glass once and that the quantum yield for producing the laser upper state ${}^4F_{3/2}$ was unity.

The time dependence of the deposition of electrical power into the amplifier flashlamps was modeled using the circuit equations of Markiewicz and Emmett.¹⁸ Specifically, we assumed that the lamp current I and voltage V_l are determined by

$$LI + RI + KI^{1/2} + \frac{1}{C} \int I dt = V_b, \quad (5)$$

and

$$V_l = KI^{1/2}, \quad (6)$$

where L is the circuit inductance, C is the capacitance, R is the resistance, and V_b is the initial voltage across the discharge capacitor. The lamp parameter K is given by

$$K = 1.3 \frac{l}{d} \left(\frac{P}{450} \right)^{0.2} \text{ VA}^{-1/2}, \quad (7)$$

where l is the arc length, d is the bore diameter, and P is the xenon fill pressure in Torr. The circuit parameters for the xenon flashlamps in the SSA were $C = 106 \mu\text{F}$, $L = 450 \mu\text{H}$, $R = 0.4 \Omega$, $V_b = 6$ to 24 kV , $l = 234 \text{ cm}$, $d = 1.5 \text{ cm}$, and $P = 300 \text{ Torr}$.

Once the lamp electrical power deposition $P_l(I, V_l)$ and the instantaneous Nd:glass pumping efficiency $\eta(P_l)$ are known, the determination of E_s , the amount of energy stored in the excited Nd ions per flashlamp, requires the solution of the differential equation

$$\dot{E}_s = \eta(P_l)P_l(t) - \Gamma_{\text{decay}}E_s, \quad (8)$$

where $\Gamma_{\text{decay}}(t)$ is the effective decay rate for the stored energy. A three-component exponential fit, which accurately tracked the fluorescence decay of a $5 \times 5 \times 0.5\text{-cm}^3$ sample of our LG-660 laser glass, was used with branching ratios of 0.0738, 0.5400, and

Fig. 6-17. Calculated Nd*-pumping efficiency vs lamp input for xenon lamp pumping (1.5-cm-bore, 300-Torr xenon, 400-nm cutoff) of LG-660 laser glass.

Fig. 6-18. Ratios of observed to predicted storage efficiencies for single-pass pumping using the 5-lamp Hagen and 5-lamp Winston reflectors.

0.3862 and lifetimes of 41, 476, and $232 \mu\text{s}$, respectively.

Equations (6) through (8), together with the results shown in Fig. 6-17, complete the calculations necessary to predict the amplifier performance for single-pass pumping. The differential equations for the lamp current and stored energy were integrated using a 4th-order Runge-Kutta numerical approximation. The peak storage efficiency was then calculated by finding the maximum value of $E_s(t)$ and dividing that value by the energy initially stored in the capacitor bank.

The ratio of measured SSA performance, using either a 5-lamp Hagen or 5-lamp Winston reflector in the single-sided illumination geometry (see the previous article), to predictions for single-pass pumping are shown in Fig. 6-18. In both cases, the ratios of experiment to theory start near 0.7 at low energy and fall off slightly with increasing bank energy. This falloff, at least with Winston reflectors, may be an artifact of errors in the correction for ASE that were applied to the experimental data.

These calculations represent an upper limit on possible amplifier performance because we assume that the transport efficiency of pump radiation from the lamps to the disk is unity. The true transport efficiency is lower because of the nonperfect reflectivity of silver reflectors and because

laser-disk holders intercept some portion of the pump light. Ray-tracing calculations based on measured reflectivity indicate that the net transport efficiency should be about 0.8. This value is in good agreement with the value of 0.7 suggested by Fig. 6-18, particularly when we consider the uncertainty (on the order of 10%) in the predictions due to experimental uncertainties in the input flashlamp data.

Multipass Pumping. In SSA experiments, a two-pass geometry was approximated by placing a flat, silver reflector on the side of the amplifier opposite the flashlamps. This setup differs from the ideal case of two-pass pumping in two opposing ways. The transport efficiency of radiation from the lamps to the laser glass (and from the glass to the opposite wall) is certainly less than unity. In addition, the photons that survive two trips through the disks have some probability of making additional passes through the glass. The remaining power may be absorbed in the flashlamps and reemitted over a broad spectral range. Such complications greatly increase modeling difficulties.

At very low input energies, the flashlamps are presumably transparent, and we can calculate in more direct fashion the enhancement for two-pass pumping (actually n -pass pumping for this condition) over single-pass pumping. Using the absorption cross section shown in Fig. 6-16, we can calculate the effect of two or more passes on pumping efficiency simply by using the curves for pumping efficiency associated with successively larger doping-thickness products. The improvement factors over single-pass pumping calculated in this manner are shown in Fig. 6-19(a). Obviously, the predicted enhancement is considerably greater than that observed (~ 1.18)

when the effective number of passes becomes large.

Much better agreement with the two-pass experiment is obtained when we allow a finite transport efficiency from lamps to disks, from disk to reflector, and so forth. Calculations were done in essentially the same manner, assuming a fixed decrement on each transport step. As shown in Fig. 6-19(b), if the transport efficiency is 0.8, then the n -pass improvement factor approximates that from experimental observation. Moreover, the value of 0.8 is in good agreement with our estimate of the actual transport efficiency.

Finally, we considered the case of two-sided pumping at reasonably large flashlamp loadings, for which lamp absorption becomes significant. The model is again idealized in that we assumed that all the unused infrared (IR) and visible radiation that reaches the lamp walls is redeposited in the lamps. We also assumed that the deposited optical energy affects the radiative output of the lamp in the same manner as would an identical increase in electrical power input.

Using the assumptions above, we incorporated energy conservation to calculate the electrical power input P_e required to produce a net power input P_i . Then

$$P_e = P_i [1 - T_{LG} T_{GL} \eta_l (1 - f_{abs} - f_{uv})] \quad (9)$$

where T_{LG} and T_{GL} are the radiation-transfer efficiencies from lamp to glass and glass to lamp, respectively, η_l is the flashlamp radiative efficiency, f_{abs} is the fraction of the total lamp output absorbed in the laser glass in a single pass, and f_{uv} is the fraction of the lamp output that lies in the ultraviolet region and is lost to absorption in either the lamp or reflector walls. Due to the dependence of the radiative characteristics

Fig. 6-19. Ratio of predicted efficiencies for multipass and single-pass pumping.

Fig. 6-20. Predicted peak Nd* storage efficiencies.

of the lamp on the input power, η_r , f_{abs} , and f_{uv} are all implicit functions of P_i .

From Eq. (9), the net effect of energy recirculation is to increase by a multiplicative factor the total power deposited in the lamp. As a result, the effective and instantaneous Nd-glass pumping efficiency relative to the electrical energy input is given by

$$\eta_{\text{eff}} = \frac{P_i}{P_e} \eta(P_i) \quad (10)$$

The ratio P_i/P_e and, thereby, $\eta_{\text{eff}}(P_e)$ was calculated using our analysis of the radiation from bare, 1.5-cm-bore xenon flashlamps.¹⁷ Once the effective pumping efficiency is known, it is straightforward to calculate the theoretical performance of the SSA with recycling using the same method as described for our single-pass analysis.

Figure 6-20(a) shows the results of our calculation in the idealized case, in which the transfer efficiencies of the radiation in the amplifier are assumed to be unity. For comparison, we have also plotted the results of single-pass pumping (with a transfer efficiency of unity). As discussed above, both T_{GL} and T_{LG} probably equal about 0.8 in the SSA. Therefore, as a more accurate model, we have repeated the analysis in the case where $T_{\text{LG}} = T_{\text{GL}} = 0.8$, as shown in Fig. 6-20(b).

With respect to the observed performance of the SSA under two-sided illumination, Fig. 6-20(b) represents an overestimate of the effects of the recycling of radiation by flashlamps. At the lowest lamp loadings, the lamps are essentially transparent to their own radiation, absorption in the lamps is negligible, and the efficiency of the lamps should lie closer to the two-pass value. However, for recycling with a finite transport efficiency, the improvement factor is very similar (~ 1.20) to the n -pass case.

Thus, the two effects trade off nearly equally, and the prediction in Fig. 6-20(b) is in relatively good agreement with experimental observations.

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Radiation Transport

Introduction. We have used ray tracing to calculate radiation transport in a pump cavity for various lamp opacities and Nd absorptions. The purpose of our evaluations was to improve our understanding of amplifier performance and to provide guidance for improved amplifier designs. We reduced the complexity of our simulations by replacing individual components of the pump cavity with equivalent emitter and absorber surfaces that had identical angular emission, absorption, and reflection characteristics averaged in space. The ray-tracing calculations were two-dimensional (2-D), while three-dimensional (3-D) corrections were introduced.

Ray Tracing. For the cylindrical reflector shapes used in our amplifiers, it is sufficient to trace the rays only in the transverse plane that is perpendicular to the flashlamp axis because the axial component of the radiation is not changed by the reflector. The calculated 2-D angular radiation profiles transverse to the lamps can then be combined with longitudinal radiation profiles from flashlamps to simulate amplifier characteristics. Our representation of a 3-D pump cavity is exact for opaque lamps. For transparent lamps, however, the representation is only approximate because the transverse profiles can change with the axial ray component. To reduce the deviations, the axial components were taken into account

by using effective emission and absorption lengths in the 2-D simulations.

We determine the angular radiation profiles from the flashlamp-reflector combination (radiator) with a ray-tracing code in which parallel rays enter the radiator at various angles, are reflected, absorbed, deflected, and then leave the radiator. The power of the transmitted rays is collected in bins for the incoming and outgoing directions, with angles ϕ_i and ϕ_o , respectively, relative to the normal of the radiator. Our code also allows us to calculate the emission into the initial ray direction using a backward accounting technique. Hence, the code also provides the emission profiles of the radiator, as discussed in detail in Ref. 19.

Flashlamp Emission. The spectral radiance at the surface of a uniform transparent plasma at an angle ψ relative to the surface normal, in units of $\text{W}/\text{cm}^2 \cdot \text{sr} \cdot \text{nm}$, is

$$R_{\lambda\psi} = \frac{W_{\lambda B} \epsilon \cos \psi}{\pi}, \quad (11)$$

where $W_{\lambda B}$ is the spectral emittance of a blackbody radiator at wavelength λ . The spectral emissivity is

$$\epsilon = 1 - \exp(-aL), \quad (12)$$

where a is the spectral absorption coefficient, and L is the 3-D path length in a cylindrical plasma. We selected a coordinate system such that θ is the polar angle with respect to the lamp axis, and β is the azimuthal angle defined with respect to the surface normal, as shown in Fig. 6-21.

The spectral power emitted from a surface element dA_o on the outside of the flashlamp into the solid angle $d\Omega_o = \sin \theta \, d\theta \, d\beta_o$ is then

$$\begin{aligned} dP_o &= TR_{\lambda\psi} d\Omega_o dA_o \\ &= \frac{TW_{\lambda B} \epsilon \sin^2 \theta \cos \beta_o \, d\theta \, d\beta_o \, dA_o}{\pi}, \end{aligned} \quad (13)$$

where the effective transmission of the lamp envelope, including multiple internal reflections in the lamp, is determined by

$$T = \frac{1 - R_W}{1 - (1 - \epsilon)R_W}, \quad (14)$$

The reflectivity of the quartz wall, including multiple reflections within the wall, is

$$R_{W} = 1 - \frac{(1 - R_i)(1 - R_o)}{1 - R_i R_o}, \quad (15)$$

where R_i and R_o are the reflectivities of the inside and outside surfaces, respectively.

Because of refraction, the angle of emission on the outside surface of the flashlamp is limited to $\beta_o = \arcsin \chi \approx 48.6^\circ$ for our standard lamps with χ (the ratio of inner-to outer-wall radius) equal to 0.75. Because of this angular limitation, the apparent source diameter of the flashlamp is equal to the plasma diameter and not to the outside diameter of the tubing.

The transverse emission profiles for an internal and external surface element of a standard flashlamp are shown in Fig. 6-21 for various lamp opacities. The lines for external emission with $T \approx 1$ show the effects of refraction alone, while the lines for external emission with $T < 1$ show the effects of both refraction and reflection on the distributions. Without reflection, the radiance $R_{\lambda o}$ at normal incidence is identical at the inside and outside surfaces of the lamp. Moreover, for opaque lamps with emissivity $\epsilon \approx 1$ and $T = 1$, the radiance of the outside surface is identical to that of the inside surface up to the cutoff angle β_o . For low emissivities, the effective transmission T of the lamp approaches unity because of vanishing lamp absorption.

The complexity of full 3-D simulations can be reduced with an angle-averaged approximation of Eq. (13) so that

$$d\bar{P}_o \approx \frac{\bar{T} \tau W_{\lambda B} \epsilon \cdot \beta_o \, d\beta_o \, dA_o}{\pi}, \quad (16)$$

where averaging over the polar angle θ yields

$$\bar{T} \tau = \frac{4}{\pi} \int_0^{\pi/2} T \epsilon \sin^2 \theta \, d\theta = \bar{T} (1 - e^{-aL}). \quad (17)$$

The effective emission length and transmission can be approximated by

$$T \approx \left(1 + \frac{0.27}{\sqrt{1 + 0.8\chi aL + aL}} \right), \quad (18)$$

Fig. 6-21. Transverse-emission profiles for cylindrical plasmas with and without quartz tubing.

and

$$\bar{T} \approx \frac{1 - R_{\beta}}{1 - R_{\beta}(1 - \epsilon_v)} \quad (19)$$

where ϵ_v is the emissivity for $\theta = 90^\circ$, and

$$R_{\beta} \approx 0.08 + 0.92 \left(\frac{\mu_i^2}{90^2} \right)^n \quad (20)$$

is the θ -averaged and polarization-averaged reflectivity of the quartz tubing for a Lambertian emitter. The path length

through the plasma of a ray at angle β_i , as seen in two dimensions, is designated as l .

For radiation that is absorbed by the flashlamp, an effective absorption length can be obtained in a similar manner. Because the axial component of the emitted radiation is essentially maintained in the pump cavity, the absorbed power in the axial plane is proportional to $Te dP_v$. This leads to a θ -averaged value of

$$\bar{T}\epsilon_v = \frac{4}{\pi} \int_0^{\pi/2} \frac{(Te)^2 \sin^2 \theta d\theta}{\bar{T}\epsilon} \approx \bar{T}(1 - e^{-\alpha l}) \quad (21)$$

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The effective absorption length can be approximated by

$$l_a \approx l \left(1 + \frac{0.56}{1 + 2.2\sqrt{a l}} \right) \quad (22)$$

The approximations of the effective emission and absorption lengths are within 1% of the integrated values. For typical operating conditions of lamp opacity $ad \approx 1$, the effective lengths are only about 20% greater than those seen in a 2-D view, which means that the correction terms from 2-D to 3-D are also only of this magnitude. Applying the effective length according to Eqs. (18) and (22) in modified 2-D simulations reduces the remaining higher-order correction terms to insignificant deviations relative to 3-D simulations.

Radiation Profiles. The calculated emission profiles, as a function of the transverse angle ϕ_o , of the Hagen reflector (radius of curvature = 36 mm, average reflectivity = 90%), are illustrated in Fig. 6-22(a) for absorption coefficients ranging from 0 to 10 cm^{-1} . A source diameter equal to the effective plasma diameter $d \approx 2r/(1 + 1 \text{ mm}/r)$ was applied in this analysis, which takes into account the transition layer of the plasma at the flashlamp walls. The emission profile

produced by the Hagen reflector is well collimated, with most of the energy emitted within $\pm 20^\circ$. For high opacities, the collimation is somewhat reduced because some of the reflected light is reabsorbed, which also reduces the reflector efficiency from 94 to 82% relative to the total emission of the flashlamp. The intensity-weighted number of reflections is 0.6 for this geometry.

The transmission profiles for radiation entering the Hagen reflector at angle ϕ_i are shown in Fig. 6-22(b). The transmission is lowest at normal incidence because most of the light is reflected toward the lamp. The angle-averaged transmission of the Hagen radiator for an external Lambertian source decreases with increasing opacity from 79 to 40%.

The absorption profiles as a function of input angle ϕ_i , shown in Fig. 6-22(c) are normalized to unity for normal incidence. The absorption profiles widen only slightly with opacity and exhibit similarities to the emission profiles. In other words, a well-collimated emitter is also a good absorber for collimated light.

Since the ray directions in a box amplifier from a given flashlamp panel to the opposite panel are essentially preserved, the angular distributions of radiation on the

Fig. 6-22. Angular (a) emission, (b) transmission, (c) absorption, and (d) reflection profiles for the Hagen reflector and various lamp opacities.

receiver panel are essentially equal to the emission profiles. The resulting angular distributions of reflected light are shown in Fig. 6-22(d) for various opacities. The distribution is plotted as a function of the output angle and assumes the input given in Fig. 6-22(a). Notice that the reflected radiation is distributed over various output angles ϕ_o for a given input angle ϕ_i . For low lamp opacities, the reflected distribution is approximately Lambertian and has a nearly circular emission profile. For high lamp opacities, the reflected distribution is hollowed out at normal incidence due to high lamp absorption. The angle-averaged reflectivity of the Hagen reflector for its own emission distribution decreases with increasing opacity from 77 to 35% and, hence, is only slightly lower than that for a Lambertian illuminator.

The same sequence of radiation profiles for a Winston reflector is shown in Fig. 6-23. The emission profile is independent of lamp opacity by design, but the emission reaches a maximum at $\pm 35^\circ$ from the normal. The efficiency of the Winston reflector is 93%. This figure translates into a value of 0.7 for the intensity-weighted number of reflections. The transmission profile exhibits preferred transmission at normal incidence,

which correlates well with the absorption profile. Once again, a broad-angle emitter is a good broad-angle absorber.

The angle-averaged transmission of the Winston case for an external Lambertian source is essentially the same as for its own emission distribution, and that transmission decreases with increasing opacity from 75 to 27%. The reflection profile of the Winston reflector for its own radiation distribution is, again, approximately Lambertian for low opacities. It appears that the radiation profiles, after one round trip in the pump cavity, approach Lambertian distributions because directions of high emission also have high absorption, and vice versa.

Pump Cavity. The calculated radiation profiles can be used to simulate amplifier performance by transmitting the emitted radiation through a blast shield, reflecting portions on the side reflectors, absorbing a specified amount in the Nd segment, and recirculating the remainder in the opposite radiator. This process is then carried out for various lamp opacities and laser-glass absorptivities. Given both the spectral absorptivity of the glass and the emissivity of the flashlamp, it is then possible to calculate the spectrally integrated pump efficiency for various lamp loadings.

Fig. 6-23. Angular (a) emission, (b) transmission, (c) absorption, and (d) reflection profiles for the Winston reflector and various lamp opacities.

While the simulations of pump uniformity and end losses require 3-D calculations, with the segments tilted at the Brewster angle it is valid to simplify the geometry for pumping calculations. The pump cavity was represented by an emitting surface, the Nd absorber, an equivalent reflecting surface, and two flat-sided reflectors.

In our simplified geometry, the effective number of side reflections necessary to reach the midplane of the amplifier is $N = (0.5X/Y) \tan \phi$, where X is the spacing of the side reflectors, and Y is the separation of the radiating surfaces. This leads to effective side reflectivities of $R_s = R_M^N$, where R_M is the reflectivity of the flat mirrors. Fresnel reflections from the blast shield and segment surfaces were approximated by θ -averaged and polarization-averaged values for Lambertian emitters of $R_F \approx 0.05 + 0.95(\phi/90^\circ)^6$ per surface. The performance of the SSA with blast shields was limited by evaluating the transfer efficiency for 2 and 3 parallel surfaces with corresponding effective reflectivities of $R_2 = 2R_F/(1 + R_F)$ and $R_3 = 3R_F/(1 + 2R_F)$.

The transfer efficiency of photons from the lamps to the Nd ions is then

determined by evaluating the ratio

$$\eta_T = \frac{P_A}{P_N} = \frac{P_A}{P_T - \eta_L P_L} \quad (23)$$

where P_A is the total absorbed light in the segment, P_N and P_T are the net and total emission from the lamps, respectively, P_L is the total reabsorbed light in the flashlamps, and η_L is the reemission efficiency. The absorbed light in the segment, including recirculation within the pump cavity, was approximated for monochromatic light by

$$P_A = \int \frac{p_E \eta_{\Lambda\phi} R_s}{1 - R_s^2 R_\phi (1 - \eta_{\Lambda\phi})} d\phi \quad (24)$$

where p_E is the emitted power from the radiator in the direction ϕ , R_ϕ is the effective reflectivity of the radiator for light entering at the angle ϕ , and

$$\eta_{\Lambda\phi} \approx \frac{\eta_\Lambda (1 - R_3)}{1 - R_3 (1 - \eta_\Lambda)} \quad (25)$$

is the effective absorptivity of the segment, including multiple Fresnel reflections. The

Fig. 6-24. (a) Single-pass transfer efficiency and (b) efficiency ratios of the SSA with blast shields for Hagen and Winston reflectors.

absorptivity of the segment depends on the effective path length in the segment, which is about 14% greater than its thickness t for external Lambertian illumination at normal incidence. For an internal rectangular disk pumped by a Lambertian radiator, the angle-averaged path length is only about 12% greater than the segment thickness. These corrections are relatively small because the maximum path length within a disk is limited to $t(1 - 1/n^2)^{-1/2} = 1.33t$ for $n = 1.52$. For the modified 2-D simulations of this analysis, a θ -averaged thickness of

$$t_\phi \cong t \left[1.07 + 0.31 \sin^2 \left(\frac{\phi}{n} \right) \right] \quad (26)$$

was applied, which accounts for the essentially Lambertian longitudinal emission of the flashlamps.

Amplifier Simulations. The results for one-sided pumping of opaque disks are illustrated in Fig. 6-24 for the Winston and Hagen reflectors. The transfer efficiencies in Fig. 6-24(a) and efficiency ratios in Fig. 6-24(b) are plotted as a function of plasma absorption for various values of lamp reemission efficiency. Blast shields were included, and mirror reflectivities of

90% were assumed for these examples. With these assumptions, the Hagen reflector exhibits the highest transfer efficiency of 84% at low lamp opacities and high reemission efficiencies. However, for high lamp opacities and low reemission efficiencies, the transfer efficiency falls off to 68%. The Winston reflector is somewhat less efficient at low opacities but is also less sensitive to variations in lamp opacity and reemission efficiency. The transfer efficiencies of the two reflector shapes are identical at plasma absorptions of 0.4, 0.6, and 1 cm^{-1} for reemission efficiencies of 0, 0.25, and 0.5, respectively. At higher reemission efficiencies, the Hagen reflector would be more efficient than the Winston reflector at all lamp opacities. The measured ratios of the axial gain coefficients with Hagen reflectors relative to Winston reflectors shown in Fig. 6-24(b) represent averages of 160 measurements. The effective plasma absorption for these data was scaled with the square of the current density, as observed experimentally, and normalized at 1 cm^{-1} for the highest input energy. The simulations are in reasonable agreement with experiments if reemission efficiency is about 50%. The differences in performance of the Winston and

Fig. 6-25. (a) Transfer efficiency and (b) cavity-enhancement ratio of the SSA with blast shields and Hagen reflectors for various lamp opacities and Nd absorptivities.

Hagen reflectors, considering axial gain coefficients or storage efficiency as shown in Fig. 6-15, are due to differences in the correction factor for pump uniformity.

The transfer efficiency for two-sided pumping of segments with angle-averaged absorptivities η_A ranging from 0.1 to 1 is illustrated in Fig. 6-25(a) for Hagen reflectors. For high absorptivities, $\eta_A = 1$, the results are identical to those for one-sided pumping. At low Nd absorptivities and low lamp opacities, recirculation of pump light leads to significant enhancements (by factors of 2 to 3) in effective Nd absorption, as shown in Fig. 6-25(b). For high lamp opacities, the enhancement is obviously reduced and leads to the observed stronger decrease in

amplifier efficiency with pump energy than that for one-sided or open-lamp pumping.

These simulations represent the behavior of specific wavelengths within a pump cavity. While lamp opacity generally increases monotonically from the ultraviolet to the infrared, except for line emission, the absorptivity of the Nd pump bands varies drastically. The spectrally averaged absorptivity for typical disk amplifiers is about 0.4, while the lamp opacities at the dominating pump bands are about unity for typical operating conditions. In other words, the spectrally integrated results will be comparable to those at $\eta_A \approx 0.4$ and $ad \approx 1.0$.

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Fig. 6-26. Coordinate system to predict end losses and transverse gain profiles based on the angular distribution projected into the x - z plane.

Pump Uniformity and End Effects

Performance of the SSA varies with the choice of flashlamp reflector, as discussed in the previous three articles. Both the end loss, which affects overall amplifier energy-storage efficiency, and the transverse variation of gain across the amplifier depend on the angular distribution of pump radiation. We have developed a model to calculate the angular-emission profiles from various flashlamp-reflector combinations and to estimate the resulting pump distributions. We have found qualitative and, in some cases, quantitative agreement with experimental observations.²⁰ A better understanding of pump uniformity has been motivated following the observation of optical distortions caused by nonuniform pumping.²¹

To simplify the numerical analysis, we idealized the SSA geometry, as described in the previous article. We assumed that the amplifier glass slabs are opaque. The flashlamps and reflectors were replaced with equivalent plane-emitting surfaces. Losses caused by a finite value of reflectivity at the top and bottom walls were determined from the average number of reflections necessary to reach the slab. The model is essentially two-dimensional and therefore presumes a flat gain profile in the vertical direction, which is consistent with experimental observations.

Our calculation of radiation transport is further simplified by choosing the geometry shown in Fig. 6-26(a). To completely specify the angular distribution of the source, we

chose a polar coordinate system, with the z axis in the direction of the lamp axis and the x axis normal to the lamp-reflector plane. The equivalent radiating surface is presumed to lie in the y - z plane at the exit plane of the reflector. In this coordinate system, the angle θ that a ray makes with the z axis is unaffected by the lamp reflectors. The reflectors simply change the angle ϕ projected into the x - y plane. Moreover, the ray path projected into the x - y plane simply follows the 2-D laws of reflection. Because the angular distribution of a bare lamp at the y - z plane $I(\theta, \phi)$ (in $\text{W} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1}$) can be described by $f(\theta)g(\phi)$, where $f(\theta) \sim \sin \theta$ for lamp, the presence of a reflector simply changes this distribution to $f(\theta)g'(\phi)$. Examples of the distributions $g'(\phi)$ for 5-lamp Hagen and Winston reflectors, are shown in Figs. 6-22(a) and 6-23(a), respectively, in the previous article. Hence, 2-D ray tracing gives the critical information for 3-D modeling.

Further simplification results from the fact that complete specification of the angular profile is not required for determination of either the direct loss of radiation out the ends of the amplifier or the transverse variation in gain. Both quantities depend only on $I(\theta_p)$, the angular distribution projected into the x - z plane, as shown in Fig. 6-26(b). For any source, we calculate this reduced profile from the relation

$$I(\theta_p) = \int_{-\pi/2}^{\pi/2} I(\theta, \phi) \sin \theta \frac{d\theta}{d\theta_p} d\phi, \quad (27)$$

where the absolute magnitudes of $I(\theta, \phi)$ and $I(\theta_p)$ are set by normalizing their integrals to unity over the allowed range of emission angles. In this integration over ϕ , the angle θ is implicitly a function of θ_p and ϕ according to the relation $\cot \theta_p = \tan \theta \cos \phi$ from Fig. 6-26(a).

Figure 6-27 shows the projected angular distributions calculated for several limiting cases of the amplifier-lamp geometry for which the required integrals can be calculated analytically. The cylindrical Lambertian distribution corresponds to a line source placed in the emitting plane $I(\theta, \phi) \sim \sin \theta$, while the plane Lambertian case corresponds to a surface source $I(\theta, \phi) \sim \sin \theta \cos \phi$. The $\phi = 0$ curve corresponds to the case of placing a lamp at the focus of an infinitely deep 2-D parabola,

which would channel all radiation from the cylindrical Lambertian emitter toward the laser glass, with $\phi = 0$.

Figure 6-27 clearly shows that significant changes in the projected distribution, and thus the end losses and gain profiles, can be caused by modifying the transverse dependence of pump emission. The rather surprising result that the cylinder source has a finite value at 90° derives from rays that are directed up or down with considerable intensity but with a projection along the lamp axis. The Lambertian plane simply projects to a $\cos \theta_p$ distribution, and the fully collimated case ($\phi = 0$) gives $\cos^2 \theta_p$.

Figure 6-28 shows the projected angular-emission profiles predicted for the Hagen and Winston reflectors used in the SSA. The ϕ distributions are given in the preceding article, and the θ variation was simply assumed to be $\sin \theta$. The integration for $I(\theta_p)$ was carried out numerically. For both reflectors, the projected emission curve is essentially equal to that predicted for a plane Lambertian source. The projected distribution with the Hagen reflector is slightly more forward-directed than the Winston case, but the difference is small.

Calculation of End Loss. The major source of end loss for the SSA is from rays that are emitted directly out the ends of the amplifier without passing through laser glass. This contribution can be calculated directly from the projected angular distribution as an integral over the allowed range

Fig. 6-27. Projected angular distributions based on a Lambertian cylinder, a Lambertian plane, and an emitter having a collimated ϕ distribution.

Fig. 6-28. Projected angular distributions predicted for the 5-lamp Winston and Hagen reflectors.

of exit angles. This fractional loss is given by

$$f_{\text{direct}} = \frac{\int dz \int_{\theta_{\text{max}}}^{\pi/2} d\theta_p I(\theta_p)}{\int dz}, \quad (28)$$

where θ_{max} is an implicit function of the spatial-integration variable z as well as of the exact geometry of the ends and disk.

Using the angular profiles shown in Figs. 6-27 and 6-28, the predictions for fractional end loss of the total pump energy are 5.5% for both the Hagen and Winston reflectors, and 8% for an amplifier with black reflectors as modeled by a cylindrical Lambertian distribution. The fractional loss for those lamp segments that directly radiate to the ends is actually 4 times greater because those segments represent only one-fourth of the pump length. Introduction of a finite lamp diameter and $\sim 90\%$ reflectivity for the top and bottom walls has a negligible effect on the results for Hagen and Winston reflectors. However, these nonideal effects, particularly the addition of a gap, reduce the fractional loss in pumping to 6% for the black reflector by depressing the projected distribution at large values of θ_p .

In addition to the direct loss, a significant portion of pump radiation is reflected out the ends of the amplifier due to Fresnel reflection from the laser disks. For each ray emitted by the source, the angles θ and ϕ uniquely define its angle of incidence with the glass surface and therefore specify the (polarization averaged) Fresnel reflection coefficient. Using this reflection coefficient, we calculated the total amount of energy reflected out the end by integrating over the allowed range of emission angles for each surface element. The predicted reflection

end loss for the black amplifier, calculated from the corrected cylindrical Lambertian emission profile, is 1.5% of the radiation delivered to the disks. Essentially the same reflected energy loss was found for both the Hagen and Winston reflectors.

Summing the direct and reflected contributions, our model predicts a total end loss of 7.5% of the pump radiation using either Hagen or Winston reflectors in the amplifier and a loss of 8% using black reflectors. The results for the Winston and black reflectors are in good agreement with the experimentally measured loss of 8% in both cases. The model overestimates the experimentally measured loss of 5% for the Hagen reflector. This outcome indicates that the Hagen reflector produces a greater degree of collimation toward the laser glass than we have calculated.

Calculations of Transverse Gain Profile. The transverse variation in gain across the aperture of the SSA is a somewhat more sensitive measure of the pump profiles in the amplifier than is the end loss. The latter quantity samples only the integrals of the emitted distribution. Comparisons of gain profiles produced from single-sided pumping and from either the plus or minus side of the amplifier represents a much more serious test of amplifier modeling. The transverse gain profile is determined in the same manner as for calculation of the direct end loss: by integration over the projected angular-emission distribution. Specifically, the local pumping experienced by an infinitesimal glass element at distance x from the radiating surface is given by

$$P(x) = \int dz I(\theta_p) \frac{d\theta_p}{dx}, \quad (29)$$

where θ_p depends on z , and the integration limits depend on the disk element under consideration. The transverse distribution in the SSA is specified by summing over the eight surface elements (including two non-equivalent elements) that contribute to the gain in our model.

In Fig. 6-29 we show the results of our model calculation for the transverse gain variations in the SSA with black reflectors along with the experimentally observed gain. The calculations were normalized to give a best fit to the data. The simple model (idealized curve) using the projected

Fig. 6-29. Transverse gain profile of the SSA predicted for the cases of idealized and nonidealized black reflectors.

Fig. 6-30. Transverse gain profile of the 55A predicted for the (a) 5-lamp Winston and (b) 5-lamp Hagen reflectors.

angular distribution of the Lambertian cylinder clearly overestimates the side-to-side variations as a result of the large values of emission at large angles for θ_{p} . This overestimate is reduced by taking into account the finite lamp size, the finite reflectivity, and the gap between the top and bottom reflector walls and the lamps for the black reflector. The results of our refined calculation are shown by the corrected curve, which is in good agreement with the experimental observations.

Figure 6-30 shows the results of similar calculations for the Winston and Hagen reflectors together with experimental data. Reflectivity of 90% for the top and bottom amplifier walls as well as the gap were included in the calculation. These corrections, however, caused a negligible change in results when compared to those obtained with the idealized profiles given in Fig. 6-28. The predicted profiles and observed values are in good agreement for the Winston reflector, as was the case with our comparison of predicted and observed end losses. For the Hagen reflector, the experimental results once again indicate that the true angular distribution is more strongly directed toward the laser glass than the model predicts.

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Flashlamps

Dependence of Flashlamp Performance on Gas Fill and Bore Size

Our flashlamp experiments are part of a continuing study on the feasibility of building affordable, large Nd-glass lasers with

improved energy-conversion efficiencies. We have investigated the radiant and electrical properties of 1.5- × 112-cm and 0.55- × 112-cm refillable, clear-fused-quartz flashlamps with gas fills of xenon, krypton, and argon. We have made measurements of the lamp impedance parameter, overall radiant efficiency, and relative Nd-pumping efficiency over a range of input energies and gas fill pressures.

Our present work with a 0.55-cm-bore lamp extends the range of bore diameters previously studied.²²⁻²⁵ It further confirms that xenon-filled flashlamps of different bore sizes (with optimum fill pressures) have nearly the same overall radiant efficiencies and Nd-pumping efficiencies when excited at a given explosion fraction f_x (input energy/explosion energy) for a fixed pulse length. We also measured the decline in overall lamp radiant output with lighter rare gases. The relative single-pass Nd-pumping efficiencies of the 1.5-cm-bore lamp with xenon, krypton, and argon at $f_x = 0.2$ were 1.0, 0.93, and 0.83 for optimized fill pressures of 300, 500, and 750 Torr, respectively. The Nd-pumping efficiencies obtained in the 0.55-cm lamp at a fixed gas fill pressure of 750 Torr for xenon, krypton, and argon were 1.02, 0.98, and 0.91, respectively, relative to the 1.5-cm-bore results for xenon. Estimated errors in our lamp-to-lamp and gas-to-gas comparisons are about 3%. These measurements were made in an open lamp configuration. Performance in an amplifier will also depend on lamp plasma absorption and reradiation efficiency.

Experiments. The experimental apparatus for our measurements has been described in previous Laser Annuals.^{24,25} The lamps were mounted in a brass frame on plastic insulating supports allowing for clear viewing of lamp radiation and for symmetric current return. All optical measurements

were made using light scattered from a plate packed with highly reflective BaSO_4 powder. The scatter plate provides for spatial averaging of discharge inhomogeneities and angle averaging of light from the flashlamp. The optical instrumentation included a spectrally flat pyroelectric detector, a Nd-fluorescence detector, and a Tektronix J-20 spectrometer with a rotating-wheel shutter (20- μs gate) to record spectra at the time of peak current. The aperture of each detector was tight to detect light from only the scatter plate, without direct contributions of lamp light. The voltage, current, pyroelectric, Nd-fluorescence, and J-20 outputs were attached to digitizing oscilloscopes and interfaced with an LSI-11/23 computer. The standardized computer-assisted data acquisition system (SCADAS) was used to acquire data.

Estimates of the absolute, overall lamp radiant efficiencies were obtained from measurements of the temperature rise of a 15-cm-diam black-anodized copper can of known mass, which surrounded the flashlamp. Nine thermocouples were placed at equal-mass intervals along the can to provide measurement of spatially averaged temperature. The heat energy initially deposited in the copper can gives a lower limit on lamp radiant efficiency because of incomplete absorption and light losses through apertures in the end plates.

It is convenient to compare²⁶ the radiant properties of lamps with different diameters (for purposes of amplifier pumping) at energy inputs corresponding to the same fractions of the explosion energy, which is given by²⁷

$$E_x = 20\,000\,ld\tau^{1/2} \quad (30)$$

where E_x is the empirically determined single-shot explosion energy, l and d are the flashlamp length and bore diameter (in cm), respectively, and the circuit time constant $\tau = (LC)^{1/2}$ is determined from the inductance L and capacitance C . The Nd-pumping efficiency was determined by dividing the peak Nd-fluorescence signal by the electrical energy deposited in the lamp during the entire pulse. Because Nd-fluorescence signals are convolutions of pumping rates and decay during the pulse,

accurate comparisons of Nd-pumping efficiencies for lamps with different sizes or different gases may be made only if the pulse shapes are constant. Lamp pulse shapes are characterized²⁸ by the circuit parameter

$$\alpha = K_0(V_0 Z_0)^{-1/2} \quad (31)$$

where V_0 is the capacitor charging voltage, and Z_0 is the circuit impedance $(L/C)^{1/2}$. The impedance parameter K_0 has been found empirically²⁹ to describe the relation between lamp voltage and current

$$V = K_0 I^{1/2} = k_0 \frac{1}{d} I^{1/2} \quad (32)$$

where $k_0 = 1.3(P/450)^{0.2}$, and P is the xenon fill pressure in Torr. Different values for the circuit L and C would be required for each input energy to maintain the same pulse shape through the range of energy inputs we explored. However, another way of comparing different lamps is to allow the usual change in pulse shape with explosion fraction (while keeping L and C fixed) and to choose appropriate L and C values for different lamps to achieve the same pulse shapes at given explosion fractions. With changes in lamp length or diameter, this requirement is met²⁶ if the capacitance remains proportional to $d^{5/3}/k_0^{4/3}$. To obtain a selected explosion fraction, the charging voltage would then scale as $k_0^{2/3}(l/d)^{1/2}$. The inductance is determined by holding τ constant.

These rules were used to choose test circuits²⁴ starting with $C = 87\,\mu\text{F}$, $L = 450\,\mu\text{H}$, and with two 1.5-cm lamps in series (electrically equivalent to one lamp of twice the length) or with $C = 174\,\mu\text{F}$, $L = 225\,\mu\text{H}$, and with one 1.5-cm lamp in the circuit ($3T = 600\,\mu\text{s}$). The former circuit was advantageous because the higher voltages used to obtain the same explosion fraction made it easier to achieve breakdown of the lamp without the use of a preionization pulse. We have found that preionization has a beneficial effect ($\sim 3\%$) on lamp output efficiency. The high explosion fractions not accessible with this circuit due to charge-voltage limitations of the capacitors were obtained by using the latter circuit.

Gas	0.55-cm-bore lamp		1.5-cm-bore lamp	
	A (in $V/A^{1/2}$)	q	A (in $V/A^{1/2}$)	q
Xenon	1.42 ± 0.17	0.196 ± 0.004	1.30 ± 0.16	0.184 ± 0.016
Krypton	1.24 ± 0.10	0.185 ± 0.011	1.01 ± 0.14	0.163 ± 0.017
Argon	1.14 ± 0.11	0.200 ± 0.012	0.95 ± 0.09	0.154 ± 0.005

Table 6-2. Least-squares fitting with statistical uncertainties for parameters A and q [defined in Eq. (33)] in the 0.55- and 1.5-cm-bore lamps.

When required for lamp breakdown, the preionization pulse was delivered 200 μ s before the main pulse.

The 0.55-cm-bore lamp was investigated in a test circuit ideal for a 0.5-cm-bore lamp, with $C = 21.75 \mu F$ and $L = 1850 \mu H$. The discrepancy in bore diameter from the nominal 0.5-cm value was discovered after the tests, when the lamp was cut for bore measurement. Use of a nonideal circuit resulted in slightly shorter pulse widths than desired for comparison of the Nd-pumping efficiencies with the 1.5-cm-bore lamp results. The effects of differences in pulse width were small enough (typically less than 10%) that accurate correction factors could be generated by modeling the Nd fluorescence.

Pressure Variations. A set of measurements was made at the fixed explosion fraction of 0.2, while varying the xenon, krypton, and argon fill pressures in each lamp. The lamp-impedance parameter K_0 , evaluated at the time of peak electrical power, is plotted vs pressure in Fig. 6-31 for the three gases. The impedance parameter was roughly constant in time but was always somewhat larger early in the pulse, before the arc expands to fill the lamp bore. The impedance parameter was also found to drop near the end of the pulse, a result probably due to the reduction in gas density³⁰ as hot gas expands into the dead volume behind the electrodes. The variation in impedance during the pulse was less pronounced for the smaller-bore lamp perhaps due, in part, to faster filling. Another contributing factor may be the smaller clearances between the lamp wall and the electrodes in the 0.55-cm-bore lamp, which would restrict gas flow to the dead volumes at low fill pressures. Ablation of the fused-silica envelope contributes to the lamp vapor pressure, so that the impedance parameter reaches a limiting value as fill pressure is reduced. This is best seen in the xenon data taken with the 0.55-cm-bore lamp.

A least-squares fit was made to the data at the higher fill pressures, where wall ablation was not important ($P \geq 100$ Torr for the 0.55-cm lamp; $P \geq 50$ Torr for the 1.5-cm lamp). We assumed that the impedance parameter for each gas could be described by

$$K_0 = A \left(\frac{P}{450} \right)^q \frac{l}{d} \quad (33)$$

Fig. 6-31. Lamp impedance parameter (measured at the time of peak input power) vs fill pressure for $f_x = 0.2$.

The fitting results for the three gases in each lamp are listed in Table 6-2. Lower lamp impedances and higher conductivities occur with the lighter rare gases.

The relative overall radiant efficiencies (integrated pyroelectric-detector signals divided by input energy) are plotted as a function of gas fill pressures in Fig. 6-32. The results have been normalized to the values obtained with 300-Torr xenon in the

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Fig. 6-32. Relative lamp radiant efficiency at $f_{\lambda} = 0.2$ determined from pyroelectric detector data vs gas fill pressures in the 0.55- and 1.5-cm-bore lamps.

Fig. 6-33. Relative Nd-pumping efficiencies vs gas fill pressures in the 0.55- and 1.5-cm-bore lamps.

1.5-cm-bore lamp. The efficiencies for each gas increase with fill pressure up to an optimal pressure, beyond which radiant efficiency is insensitive to increases in pressure. Radiant efficiencies are lower for the lighter rare gases at each fill pressure. Comparison of the data for each gas in the two lamps clearly shows that the rise in efficiency with rare-gas pressure occurs at a higher pressure in the smaller-bore lamp. This

finding suggests the importance of diffusion losses to the wall for the smaller-bore lamp. A common curve is generated when the radiant efficiencies measured previously with a 2.7-cm-bore lamp are plotted as a function of $Pd^{2/3}$. We have no explanation for the $Pd^{2/3}$ dependence.

The Nd-pumping efficiencies (peak Nd-fluorescence signal divided by input energy) are plotted vs pressures for each of the

Fig. 6-34. Overall lamp radiant efficiency from calorimetry data vs f_x in the 0.55-cm-bore lamp.

Fig. 6-35. Overall lamp efficiency from calorimetry data vs f_x in the 1.5-cm-bore lamp.

three rare gases in both lamps in Fig. 6-33. Only slight differences in the shapes of these curves distinguish them from the corresponding plots of overall radiant efficiencies in Fig. 6-32, showing that most of the variation in Nd-pumping efficiency with pressure is due to changes in overall lamp radiant efficiency. However, the Nd-pumping efficiencies increase slightly faster than the overall radiant efficiencies with pressure. A slight peak in Nd-pumping efficiency occurs near the optimized pressures at which point the radiant efficiency becomes flat. The slight decline beyond this point occurs because of changes in the lamp spectra with pressure (line emission becomes

less prominent, and there is an enhancement at shorter wavelengths) and because of changes in the pulse width with pressure (pulse widths become greater).

Variations in Input Energy. The Nd-pumping and radiant efficiencies are plotted vs explosion fraction for the 0.55- and 1.5-cm lamps in Figs. 6-34 through 6-37. Data were taken with gas fills of 750 Torr for the three gases in the 0.55-cm lamp and with fills of 300-Torr xenon, 550-Torr krypton, and 750-Torr argon in the 1.5-cm lamp at a 3T pulse length of 600 μ s. The Nd-pumping efficiencies have been slightly corrected for pulse shape so that at a given explosion fraction the standard pulse shape

Fig. 6-36. Relative Nd-pumping efficiencies vs f_x in the 0.55-cm-bore lamp.

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Fig. 6-37. Relative Nd-pumping efficiencies vs f_x in the 1.5-cm-bore lamp.

is that obtained with the 1.5-cm lamp with 300-Torr xenon. Corrections were, at most, 10% of those for the krypton and argon data with the 0.55-cm-bore lamp. To obtain this correction, we simply assumed that Nd-pumping is proportional to lamp electrical input, with single exponential decay of Nd fluorescence at an experimentally determined rate. The absolute sensitivities of the Nd-fluorescence and pyroelectric detectors to lamp radiation were found to change slightly with time. Three or four shots in each data set for each gas were repeated in the course of a few days so that data could be normalized to these round-robin shots. The standard deviations of the normalization

factors were $<1\%$ for each gas in each lamp. The Nd-pumping efficiencies were normalized again to the values obtained with 300-Torr xenon in the 1.5-cm lamp at $f_x = 0.2$.

Spline fits were made through the data so that the radiant characteristics of the lamp could be evaluated at $f_x = 0.2$. The relative single-pass Nd-pumping efficiencies and the relative radiant efficiencies from the pyroelectric-detector data and the calorimetry results were normalized to the values obtained with xenon in the 1.5-cm lamp. These values are given in Table 6-3. The Nd-pumping efficiency is proportional to overall lamp radiant efficiency multiplied

Gas	Nd-pumping efficiency	Relative radiant efficiency		
		Pyroelectric	Can calorimetry	Nd*/f pyro
1.5-cm-bore lamp				
Xenon	1.0	1.0	1.0	1.0
Krypton	0.93	0.91	0.93	1.02
Argon	0.83	0.86	0.85	0.97
0.55-cm-bore lamp				
Xenon	1.02	1.01	0.99	1.01
Krypton	0.98	0.95	0.93	1.03
Argon	0.91	0.89	0.88	1.02

Table 6-3. Nd-pumping and overall radiant efficiencies for two lamps with optimized gas-fill pressures. Spline fits made through the data were evaluated at $f_x = 0.2$.

Fig. 6-38. Radiant output fractions in three spectral regions vs f_x for 1.5-cm-bore, 300-Torr and 0.55-cm-bore, 750-Torr xenon flashlamps.

by a factor that takes into account the overlap of lamp spectra and Nd-pump bands.^{24,31} The ratio of Nd-pumping efficiency to overall radiant efficiency (Nd^*/f_{pyro}) is therefore a parameter that depends on the lamp spectral distribution only. The fact that nearly the same ratios are obtained for xenon and for krypton in our two lamps suggests that the spectral distributions do not change with lamp bore at $f_x = 0.2$. With argon, however, there is a slight difference in the ratio of the different lamps.

More detailed measurements of spectral distributions were made by recording pyroelectric signals using long-pass filters with cutoffs near 0.4 and 1.0 μm , along with unfiltered pyroelectric-detector data. The curves for filter transmission were convolved with the spectra recorded with the J-20 spectrometer at each explosion fraction to determine the fractions of radiant output

for $\lambda < 0.4 \mu m$, $0.4 < \lambda < 1.0 \mu m$, and $\lambda > 1.0 \mu m$. The fractional distributions in these three spectral regions are plotted in Fig. 6-38 vs explosion fraction for the output of the 0.55-cm lamp filled with 750-Torr xenon. For comparison, similar results are plotted for the 1.5-cm lamp filled with 300-Torr xenon. Little change in spectral distribution occurs between lamps at a given explosion fraction.

Spectra. The J-20 spectrometer uses a silicon Vidicon detector, which has a nonlinear response to pulsed sources. Corrections must be made in software to obtain quantitative spectra. Calibration of the spectrometer, which was described previously,²⁵ has now been done somewhat differently. Spectra of either flashlamps or cw tungsten lamps (with a 2- to 8-ms shutter) were recorded both with and without neutral-density filters of known transmittance. The

Fig. 6-39. Spectral output distributions with old and new J-20 calibration factors. Data were recorded with 300-Torr xenon at $f_x = 0.19$.

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ratios of the signals obtained at each element ($R = S/S_{ND}$, unattenuated signal/attenuated signal) were plotted vs unattenuated signals. The natural variation in intensity over the wavelength range of the 512 elements was exploited to obtain a dynamic range of signal. The curve of ratio vs signal was then modeled by choosing the parameters x , A , B , and C in the formula used to relate the relative intensity I and the measured unattenuated signal S such that

$$I = S^x (A + BS + CS^2) \quad (34)$$

Equation (34) was used because it allows the ratio of the unattenuated to attenuated signal to approach any power law as the signal becomes small (determining the x parameter).

After correcting the J-20 signals for nonlinearity of response with input intensity, a shuttered tungsten source was used as a standard to determine the response of the instrument as a function of wavelength. Long-pass filters had previously been used to eliminate second-order contributions in each of the three spectral regions used: blue (0.3 to 0.7 μm), yellow (0.5 to 0.9 μm), and red (0.7 to 1.1 μm). Because tungsten sources are relatively weak below 0.5 μm , the measured spectrum of the tungsten lamp was susceptible to scattered light contributions in the blue range. This problem was solved by replacing the blue-range long-pass filter with a Corning 4-96 filter (0.38- to 0.56- μm transmittance) to reduce

the contributions of scattered light from long wavelengths. Spectra of the 1.5-cm lamp filled with 300-Torr xenon at $f_x = 0.2$ using the previously generated corrections and the new corrections are plotted in Fig. 6-39. The corrected spectra now show larger emissions at shorter wavelengths.

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Chemical Flashlamps

We have characterized the optical output and Nd:glass pumping performance of large-scale (120-cm-long, 1.2-cm-inner-diam), metal-oxidizer chemical flashlamps supplied to us by G.T.E. Sylvania. The experimental results were obtained on the same testbed that was used to study xenon electrical flashlamps, as described in "Dependence of Flashlamp Performance on Gas Fill and Bore Size," earlier in this section. The peak Nd^{3+} inversion levels produced by the chemical lamps were $\leq 10\%$ of those generated by a xenon lamp of similar size and energy loading. The peak Nd^{3+} levels are in good agreement with predictions for the pumping rates in Nd:glass by a blackbody³² at the color temperatures of 3000° to 5000°C, which we have measured during the burn of the pyrotechnic lamp.

The chemical lamps provided by G.T.E. Sylvania consisted of either a quartz or an

alumino-silicate glass tube (1.2 cm inner diam, 1.6 cm outer diam) surrounded by a polycarbonate envelope (1.9 cm outer diam). The envelope was filled with approximately 820 mg of $\sim 25\text{-}\mu\text{m}$ -wide (i.e., 1-mil) shred cut from $25\text{-}\mu\text{m}$ -thick zirconium foil. The purpose of the polycarbonate envelope was to contain the glass, which usually fragments on firing. Copper electrodes (1.6 cm long, 0.64 cm diam) inserted into the ends of the tube provided contact with the lamp ignition circuit. The ends of the lamp were sealed mechanically and connected to the gas-handling system for evacuation and fill with a stoichiometric (1230-Torr) atmosphere of either O_2 or OF_2 . Lamp ignition was established by a 10- to 20-kV trigger pulse, which deposited 50 to 100 J of electrical energy into the lamp in less than 1 ms.

The time-dependent spectral output of the lamps was measured on a coarse scale using filtered ($>1\text{-}\mu\text{m}$, long-pass filter) and unfiltered, spectrally flat pyroelectric detectors. The ratio of the signals from the two detectors was used to obtain data on temperature vs time, assuming the output to be blackbody. The detailed shape of the time-integrated spectrum, obtained

using a Tektronix J-20 optical spectrometer, confirmed the blackbody character of the output.

An ED-2 Nd:glass, peak-fluorescence detector was used to evaluate the excited-state level produced by the lamp output. As was the case for studies on the xenon lamp, the detectors viewed only the light scattered from a barium sulfate (packed-powder) scatter plate, which provided an angle-averaged spectrum. All data, including the voltage and current characteristics of the lamp ignition pulse, were fed into an LSI-11 computer for analysis. The pyroelectric detectors were calibrated on an absolute basis using the signals produced by a 1.5-cm-bore xenon flashlamp as a reference.

Figure 6-40 shows typical results for the time-dependent optical output of the lamps (as evidenced by the signals from the unfiltered pyroelectric detector) for O_2 and OF_2 oxidizers. For comparison, the observed output power from a 1.5-cm-bore xenon flashlamp is displayed at a ratio of input to explosion energy of 0.2. In both cases, the pyrotechnic lamp triggers were delivered at a time of 200 μs . The several-millisecond delays between trigger pulse and onset of

Fig. 6-40. Comparison of radiant output from Zr/O_2 and Zr/OF_2 lamps with output from a xenon flashlamp.

the output from the chemical lamps was found to be typical of their operation. An inverse correlation between the length of the delay and the amount of energy deposited by the trigger pulse was observed, but the typical jitter at a given trigger input was about 1 ms.

Integration of the output pulses for the chemical lamps yields a radiation efficiency, as measured by the pyroelectric detector (output energy/stored energy), of $20 \pm 2\%$ and $7 \pm 1\%$ for the Zr/O_2 and Zr/OF_2 systems, respectively. The primary cause of the lower yield using OF_2 as the oxidizer is a premature termination of the output pulse, which is apparently caused by the deposition of an opaque layer on the inner wall of the lamp. Based on the energy released in the reaction and the peak measured temperatures, we would expect a pulse length of several milliseconds. This premature turn-off of the lamp output was not observed in shorter lamps (length < 20 cm), which were examined at G.T.E. Sylvania

Fig. 6-41. Time dependence of color temperatures for Zr/O_2 and Zr/OF_2 lamps.

Fig. 6-42. Predicted ratio of peak Nd^* pumping efficiency from a blackbody source to that produced by a 1.5-cm-bore xenon flashlamp ($f_s = 0.2, 3, LC = 600 \mu s$).

and which also exhibited radiant efficiencies of 20 to 30% with OF_2 as the oxidizer. We postulate that short lamps do not become as hot because of their greater end losses. The opaque material apparently results from the reaction of hot ZrF_4 with the silica wall. In contrast, Zr/O_2 lamps run cooler and produce a translucent layer of ZrO_2 on the wall.

To ensure that a nonuniform ignition was not responsible for the decreased efficiency of the longer lamps, a framing camera was used to acquire multiple exposures of the entire lamp length during its output. Occasional propagation of the burn front starting at either of the electrodes was observed, with less than 0.5 ms required for the front to travel the entire length of the lamp. In several cases, uniform ignition was established by the addition of fine copper whiskers to the electrodes or by the introduction of a thin tungsten trigger wire in the middle of the lamp bore. However, we observed no substantial increase in duration of lamp output or decrease in darkening of the walls.

Figure 6-41 shows the inferred color temperatures during the burn of the Zr/O_2 and Zr/OF_2 lamps, which were calculated from the ratio of the filtered and unfiltered pyroelectric signals. At the peak of the output, the color temperatures of the O_2 and OF_2 lamps are ~ 3500 and ~ 4500 K, respectively. These values are in good agreement with the brightness temperatures of ~ 3700 and 4600 K, which are indicated by the radiative power from the lamps determined from the magnitude of the pyroelectric-detector signals. The close correspondence between color and brightness temperatures is a strong indication of the blackbody nature of the pyrotechnic lamp output.

The signals generated in the Nd:glass fluorescence detector by the chemical lamps were observed to follow the pump pulses with an integration time roughly equal to the 350- μs spontaneous decay time of the excitation in ED-2. Peak excitation levels, relative to those produced by the standard output pulse of a 1.5-cm-bore xenon flashlamp, were $4 \pm 0.5\%$ and $8 \pm 1\%$ for the Zr/O_2 and Zr/OF_2 lamps, respectively. Since the Nd fluorescence signals follow the lamp output with little delay, the peak signal is not sensitive to the lamp efficiency but, rather, to its instantaneous power.

In Fig. 6-42, we show the calculated values for the ratio of peak Nd^* pumping

efficiency to that of a xenon flashlamp. The values result from our previously described model for the blackbody pumping of Nd:glass.³² The calculation assumed a steady-state, blackbody excitation pulse and a 350- μ s lifetime for the Nd* excitation. In addition, the calculation incorporates our experimentally determined values for the pumping efficiencies of Nd:glass by the xenon flashlamp (600- μ s pulse length). At 3500 and 4500 K, the predicted values for the blackbody performance relative to the xenon lamp are 3 and 10%, respectively. The agreement between the experimental and predicted values for this ratio, based on measured temperatures, provides a good consistency check. Moreover, we conclude that a significant improvement in the pumping of Nd:glass by a chemical lamp could be achieved only through a dramatic increase in its radiant temperature. Such an increase seems unlikely.

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Electron-Beam-Pumped Phosphors

Electron-beam excitation of solid-state scintillators, or phosphors, can result in efficient generation of visible light confined to relatively narrow regions of the spectrum.^{33,34} The conversion efficiency can exceed 20%, and, with proper choice of phosphors, radiation can be obtained anywhere from the near infrared (IR) to the near ultraviolet (UV). These properties qualify the phosphors as a potentially useful pump source for new solid-state lasers.

A demonstration of laser pumping using a conventional television phosphor was performed fairly early in the history of laser research.³⁵ Conventional television phosphors consist of glass plates covered with crystalline powder and organic binders. These phosphors have low heat capacity and poor thermal conductivity. Thus, they can be operated only at very low energy and average power densities before thermal quenching or physical destruction of the coating occurs (i.e., for current densities ≤ 50 mA/cm² at 20 kV).

New phosphors are being developed for high-brightness television tubes that are

capable of higher power dissipation.

Here, an epitaxial film of fluorescent material is grown on a crystalline substrate with good thermal properties. For example, researchers at North American Phillips Laboratories have developed a cerium-doped yttrium aluminum garnet (YAG) grown on a YAG substrate, which has operated at 1 A/cm² at 20 kV without observed thermal quenching.³⁶ The input power is higher by almost two orders of magnitude than that which can be tolerated by a conventional television phosphor.

The temperature rise in the phosphor is determined by heat capacity, thermal conductivity, and depth of penetration by the electrons. Assuming a flat-topped excitation pulse of duration t and uniform energy deposition over depth l , the peak temperature rise at the end of the pulse³⁷ is

$$T = \frac{E}{\rho c_p} \left[1 - 4i^2 \operatorname{erf} \left(\frac{l}{2\alpha\sqrt{t}} \right) \right], \quad (35)$$

where E is the deposited energy per cm², ρ is density, c_p is heat capacity, and α is the coefficient of heat conduction. For example, for YAG, $\alpha = 0.22$ cm²/s, and $\rho c_p = 2.7$ J/cm³°C. The function $i^2 \operatorname{erf}$ is the twice-integrated error function compiled in Ref. 37.

Assuming an excitation pulse of 100 μ s and $l = 2 \cdot 10^{-4}$ cm (for ~ 20 -kV electrons), the temperature rise is

$$T = 0.1 \frac{E}{\rho c_p}. \quad (36)$$

A temperature rise of 300°C reduces the fluorescence of the Ce:YAG by one-half,³⁸ and this energy density occurs with an input beam fluence of ~ 1.6 J/cm². Increasing the operating voltage to 200 kV, which is the maximum tolerable voltage before serious material damage occurs, increases the depth of electron penetration to $\sim 5 \times 10^{-3}$ cm and the allowed e-beam fluence to 4 J/cm². With an efficient phosphor, this system could become a useful pump source for a solid-state laser.

We constructed the experimental setup shown in Fig. 6-43 to test the new phosphors. A glow-discharge electron beam^{39,40} was used to generate >100 - μ s pulses with currents up to 0.2 A at ~ 25 kV, focusable to a spot of 30×3 mm on a phosphor. The

Fig. 6-43. Experimental setup for studying solid-state laser excitation using e-beam-pumped phosphors.

electron gun worked best at a pressure of 10^{-7} Torr of helium. Philips Laboratories provided phosphor samples. The cerium-doped YAG layer was 10- μ m thick and was overcoated with aluminum. We attached the laser crystal (for example, Nd:YLF or Nd:YAG) to the phosphor substrate using optical cement.

We performed some phosphor tests without the laser crystal. Fluorescence efficiency of the phosphor did not roll over up to the maximum power inputs possible with this e-beam, but the absolute fluorescence efficiency measured with a calibrated pyroelectric detector was below the expected 9% by at least a factor of 2.

This system did not have enough gain for self-sustained laser oscillations. To estimate the gain, we added a xenon flashlamp, which also pumped the laser. We could easily observe the change in laser threshold due to the e-beam-pumped phosphor and deduced a net single-pass gain of 0.5% from our preliminary measurements in Nd:YLF crystals. The cavity loss in the laser was 4%; thus the laser was below threshold with phosphor pumping alone. It should be possible to build a higher-Q laser cavity to reach oscillation threshold even with this simple geometry at somewhat higher e-beam fluence.

It would be better, however, to epitaxially grow phosphors directly on the crystal itself. We are considering phosphors thought to be more efficient and that match laser ions and crystalline hosts.

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Optical Materials, Surfaces, and Coatings

Introduction

We spent most of 1984 completing the development of antireflective (AR) coatings for fused-silica lenses in Nova and developing a recently discovered process for applying AR coatings to water-soluble KDP crystals.⁴¹ Initially, we evaluated coatings made of a polymeric porous silica that was first produced at Westinghouse.⁴² We fabricated thin (60- to 90-nm) films with good resistance to damage induced by 0.35- μ m laser pulses. Thicker layers of this material contained carbon and damaged at low fluences.⁴³ Thus, we investigated alternative porous-silica coatings made from sols containing silica particles with diameters from 5 to 20 nm.

Coatings with high damage thresholds at 0.35 μ m were readily applied to fused-silica substrates from these sols, regardless of whether the carrier was water or alcohol. We found that the alcohol solutions provided damage-resistant AR coating on water-soluble KDP crystals. The latter discovery led to a change in the baseline design of Nova, from harmonic converters containing KDP crystals immersed in index-matching liquid to converters with dry AR-coated surfaces.

We also completed a few tests of vacuum-deposited, highly reflective (HR) coatings and evaluated the new sol-gel coatings for use at 1.06 μ m. Finally, we began two new efforts: evaluation of damage in crystals of interest for high-average-power applications and a program to evaluate sol-gel coatings for U.S. Air Force (USAF) applications.

Sol-Gel Coatings

Sol-Gel Antireflective Coatings. A new porous-silica, narrowband AR coating was developed for the Nova silica and KDP optics.⁴¹ This coating was prepared from a silica sol in ethanol, obtained by the ammonia catalyzed hydrolysis of tetraethyl silicate in ethanol, as follows:

The solution was prepared at room temperature by mixing the silicate in ethanol with a concentrated aqueous ammonia solution. About three days are required to complete the formation of the sol. The resulting product consisted of a suspension of silica particles in ethanol. The particles are substantially spherical with diameters greater than 200 Å. We found that a silica concentration of 3% was the most useful for coating fabrication.

Optics were coated by dip, drain, or spin. After air drying, no further treatment was required. The coating consisted of layers of particles packed to yield an average porosity of about 50%. This resulted in a coating refractive index near the value of 1.22 that is required to give a quarterwave narrowband coating on substrates with an index of about 1.5.

The coating thickness could be varied either by changing the silica concentration in the sol or by changing the withdrawal rate or spin speed of the optic during the coating process. Multiple coats could also be applied. The transmittance spectra of coatings on 5-cm-diam silica disks and 5- × 5-cm KDP crystals are shown in Figs. 6-14 and 6-45. Figure 6-44 also shows the effect of multicoats. Each of three coating applications produced a layer with an optical thickness of about 46 nm.

We scaled up this process to coat Nova optics. For silica lenses and debris shields, we used a 200-gal coating tank with a drain-coating process in which the coating mixture was drained at a constant rate from the tank containing the optic. We had no problem with scaling up the sol preparation to 200 gal. However, to attain maximum purity, we found it necessary to fractionally distill the tetraethyl silicate starting material

Fig. 6-44. Two-surface transmittance of a fused-silica window with sol AR coatings.

prior to use. We also found that filtration of the coating mixture through a 0.2-μm filter was essential to maintain cleanliness in the system. The silica optics were coated at a thickness to give 100% transmission at 0.42 μm. This wavelength was selected to enable us to use the products at both the second (2ω) and third (3ω) harmonics. At both 2ω and 3ω, the theoretical two-surface transmittance of this coating is 99.3%, and this value was obtained for several of the optics that were coated. For others, transmittance was lower by a few tenths of one percent.

Fig. 6-45. Two-surface transmittance of a KDP crystal with sol AR coatings.

Surface effects caused problems in the dip coating of large KDP samples. Therefore, we switched to a spin-coating process. Spinning has two principal advantages over dipping. In spinning, the two surfaces of a crystal are individually coated so that the coating applied to a given surface of the crystal can be tailored for the light that will pass through that surface. In addition, far less material is required for coating by spin than that for either dip or drain.

Sol-Gel Highly-Reflective Coatings.

Oxide coatings with a high refractive index have been investigated as potential components for damage-resistant, dielectric HR coatings. These were prepared by application of solutions of metal chloride alkoxides to silica substrates. The derivatives, after coating, react with atmospheric moisture to give the relevant oxide, such that

In contrast to the silica system, we found that a heat treatment to at least 150°C was required to attain high-density (high-index) coatings. In this manner, Ta₂O₅, Nb₂O₅,

ZrO₂, and TiO₂ coatings were prepared, each having an index greater than 2.0. Damage thresholds of 4.5 and 3.5 J/cm² were obtained for Ta₂O₅ and ZrO₂, respectively, at 0.346 μm; Nb₂O₅ and TiO₂ absorbed strongly at this wavelength.

Multilayer reflective coatings were prepared using Ta₂O₅, with SiO₂ as the low-index component. Some problems were encountered with crazing due to stress; however, it was possible to produce a good coating with five layers of Ta₂O₅ and four layers of SiO₂. This multilayer coating had a reflection of 93% at 0.4 μm.

Laser Damage Studies

Tests of Sol-Gel AR Coatings with 0.6-ns, 0.355-μm Pulses. We tested about 100 sol-gel coatings, which can be grouped into four sets:

- Siloxane AR coatings on silica.
- Sol AR coatings on silica.
- Sol AR coatings on KDP.
- Coatings containing layers of progressively lower refractive index.

The latter coatings were made by mixing the particulate silica in sols with the siloxane silica.

We made little progress in fabricating damage-resistant siloxane coatings for silica. However, we made several attempts to incorporate oxidizers into these films to eliminate the carbon believed to be responsible for their low thresholds. Careful attention to all steps in fabrication of thin (60- to 90-nm) siloxane films usually produced coatings with 0.355-μm, 0.6-ns thresholds of 5 to 9 J/cm². We also found these coatings to be useful as 0.248-μm AR coatings, but thicker siloxane coatings usually had 0.355-μm thresholds below 2 J/cm². Fabrication of damage-resistant siloxane coatings continues to be an important task because this material is needed in several designs for broadband AR coatings.

The distribution of thresholds measured on the newer sol AR coatings for silica is shown in Fig. 6-46. Both the maximum threshold and the range of thresholds are approximately the same as those measured in thin siloxane coatings, but thresholds of the sol coatings were independent of coating thickness. Our most recent results suggest that coatings with thresholds exceeding

5 J/cm^2 can be routinely fabricated by making the coating solutions from distilled chemicals, filtering finished solutions to remove large particles caused by clustering of the sol particles, and then carefully cleaning the substrates.

It was difficult to unambiguously establish a $0.355\text{-}\mu\text{m}$ threshold for the sol coating on either diamond-turned or polished surfaces of KDP crystals because the bulk of the KDP usually damaged at a fluence below the apparent threshold of the AR coating. However, in measurement of third-harmonic conversion efficiency using two AR coated 27-cm Nova crystals, we subjected two 1-cm-diam sites on the crystals to 1.3-ns, $1.053\text{-}\mu\text{m}$ pulses at fluences up to 5 J/cm^2 for about 240 shots. The converted fluences at $0.351 \mu\text{m}$ were as high as 3.3 J/cm^2 . Neither of these two test sites suffered damage. The distribution of fluences recorded in these tests is shown in Fig. 6-47. The outcome serves as an adequate characterization of the minimum damage resistance of sol coatings on KDP surfaces.

Finally, we also tested AR coatings consisting of one to four layers of varied refractive index made by using varied mixtures of the sol and siloxane solutions. Thresholds measured on some single-layer films of this type are shown in Table 6-4. Some of these thresholds were as high as 10 J/cm^2 , which is the median threshold for bare, polished silica. However, the multilayer films usually damaged at less than 2 J/cm^2 . We currently believe that this may be the best coating system for fabrication of broadband AR coatings, although problems remain in developing a successful technique for combining the acidic siloxane solutions with the basic sols.

Dependence of Damage in Sol-Gel Films on Pulse Duration. In the Novette laser, lenses and windows that transmitted only the fundamental $1.053\text{-}\mu\text{m}$ beam were made of BK-7 glass with AR surfaces formed by neutral-solution processing (NSP). At positions in the chain where fluence was greatest, these lenses damaged due to the presence of a low-volume density of platinum inclusions. It may be necessary to replace some of the BK-7 components on Nova with components made of fused silica and coated with sol-gel AR films to avoid this problem. Therefore, we used pulses with durations of 1, 5.5 and 9 ns to measure the $1.064\text{-}\mu\text{m}$ damage

Fig. 6-46. Damage thresholds measured with $0.355\text{-}\mu\text{m}$, 0.6-ns pulses in sol AR coatings on fused-silica substrates.

Test	Number of layers	Laser-damage threshold (J/cm^2)
L-710	1	9.7 ± 1.6
L-719	1	10 ± 2
L-748	2	2.9 ± 0.7
L-770	1	9.7 ± 1.5
L-779	1	<2.0
L-808	4	<2.1
L-810	4	<2.2
L-812	4	<2.1
L-813	3	2.2 ± 0.4
L-814	4	<3.0
L-823	4	<2.4

▲ Fig. 6-47. Comparison of input $1.053\text{-}\mu\text{m}$ fluences with converted $0.351\text{-}\mu\text{m}$ fluences used to measure harmonic-conversion efficiency in AR-coated KDP crystals.

Table 6-4. Laser-damage thresholds measured with $0.355\text{-}\mu\text{m}$, 0.6-ns pulses on coatings made of sol-siloxane mixtures.

thresholds of 10 silica samples with sol-gel AR films.⁴⁴

The test coatings using a sol solution were prepared by the spin process described in "Sol-Gel Coatings," earlier in this article. We used a solution that deposited an 88-nm-thick porous-silica layer. Three applications provided coatings with thickness almost equal to the $0.266\text{-}\mu\text{m}$ quarterwave thickness for the $1.064\text{-}\mu\text{m}$ wavelength. The substrates were Corning 7940 fused silica

Optical Materials, Surfaces, and Coatings

Table 6-5. Laser-damage thresholds at 1.064 μm for sol-gel AR coatings on fused-silica substrates.

Sample	Preparation group	Coating type (sol)	Laser-damage thresholds (J/cm^2)					
			1-ns pulse		5.5-ns pulse		9-ns pulse	
			Damage	Light	Damage	Light	Damage	Light
3811	1	Ethanol	10 ± 1.5	<8	27 ± 4	19 ± 3	6 ± 5	32 ± 11
3815	1	Ethanol	13 ± 2	5.8 ± 6	26 ± 4	21 ± 2	41 ± 6	34 ± 5
3810	1	Methanol	9.4 ± 1.4	8.0 ± 2	16 ± 2.4	11 ± 1	17 ± 2.5	15 ± 3
3817	1	Methanol	7.7 ± 2	13 ± 3	16 ± 4.5	>31	11 ± 2	>15
3746	2	Methanol	3.4 ± 3	8 ± 2	7.6 ± 1	14 ± 1	6.7 ± 1	>11
3744	2	Ethanol	7 ± 1.2	7 ± 1.2	19 ± 3.8	18 ± 1.8	28 ± 4	16 ± 1.6
3815	3	Methanol	18 ± 2.7	12 ± 1	46 ± 7	30 ± 1	61 ± 9	37 ± 15
3817	3	Methanol	17 ± 2.6	12 ± 2	42 ± 6	33 ± 6	70 ± 10	50 ± 15
3810	3	Ethanol	11 ± 1.7	11 ± 2	45 ± 7	43 ± 4	62 ± 10	53 ± 6
3811	3	Ethanol	13 ± 2	13 ± 1	48 ± 8	32 ± 3	40 ± 4	32 ± 5

Fig. 6-48. Damage thresholds at 1.064 μm on porous-silica coatings vs pulse duration.

Fig. 6-49. Damage thresholds measured with 0.355-nm, 0.6-ns pulses on vacuum-deposited HR coatings. ▼

that had been polished by Zygo, Inc. The two-surface transmittance of the coated samples at 1.064 μm was 99.5%.

Damage thresholds for the test samples are given in Table 6-5. The samples were prepared in three sets. The thresholds measured on coatings in the first two sets indicated that coatings made from methanol solutions were less resistant to damage than were coatings prepared from ethanol solutions. Since both solutions should have yielded coatings of comparable purity, we used HF to remove the coatings made in the first set and then recoated the substrates, being careful with the cleanliness of the coating process. Thresholds measured on coatings in the third set were independent of precursor solution.

The damage thresholds are plotted as a function of pulse duration in Fig. 6-48. The median thresholds at each pulse duration scaled as $T = k\tau^{0.56}$, where $k = 10.8 \text{ J}/\text{cm}^2$, and τ is pulse duration. This scaling is approximately the same at the $\sqrt{\tau}$ scaling that we usually observe in tests of carefully polished bare-silica substrates.

The median of the 1-ns, 1.064- μm thresholds measured in these sol-gel coatings was $10.8 \text{ J}/\text{cm}^2$, the same as the corresponding median threshold for NSP surfaces on BK-7 glass. However, the spread is much greater in the thresholds of sol-gel coatings. Therefore, while the sol-gel films should be usable in Nova applications, the probability of making an unsatisfactory coating is greater than that of making an unsatisfactory NSP surface.

Vacuum-Deposited HR Coatings. We evaluated several sets of vacuum-deposited HR coatings. The first set was fabricated at Optical Coating Laboratory, Inc. to provide

a second test of the usefulness of overcoat layers in improving thresholds of 0.355- μm reflectors. These reflectors were deposited on substrates held at either 150° or 250°C, and each contained 25 quarterwave-thick layers alternating between MgF_2 and Sc_2O_3 . Some had halfwave-thick overcoat layers of either SiO_2 or of MgF_2 . The substrates were made of BK-7 glass.

Figure 6-49 shows the damage thresholds measured on these reflectors with 0.6-ns, 0.355- μm pulses. Overcoats did not provide an increase in thresholds. This outcome is in agreement with results of previous 0.355- μm tests but contradicts results obtained at both 0.248 and 1.064 μm . Thresholds were greatest for coatings deposited at 250°C, but the median threshold for such coatings, 2.9 J/cm², is not significantly greater than the median for coatings tested in 1983. Because continued and careful practice of the art of vacuum deposition yielded only a slight improvement, it seems likely that a major alteration of deposition technique is now needed.

A coating technique under current evaluation is ion-assisted vacuum deposition in which energetic (50- to 1000-eV) ions are applied to a film during its growth. Collisions between the ions and the atoms being adsorbed from the vapor increase the surface energy of the atoms and allow fabrication of denser films. Two general results were obtained in evaluating several sets of coatings grown under bombardment by either oxygen or argon ions. First, when the ion stream was not electrically neutralized, surface charging at the substrate caused arcing between the tooling and the substrate, which resulted in deposition of metal in the film. Second, coatings bombarded with electrically neutral ion beams had good cosmetic quality, but thresholds were usually less than or equal to those of beams grown without ion-assisted deposition. While it is too early to accurately forecast the full potential of ion-assisted deposition, this technique has failed to quickly provide threshold improvement.

We also measured thresholds well below 2 J/cm² on films grown by ion-beam sputtering. This technique is one in which the coating vapor is made by ion bombardment of a target.

Float-Polished Fused Silica. A float-polishing procedure has been recently demonstrated.⁴⁵ In this technique, the substrate is either mechanically carried so that it is not in direct contact with the polishing lap or is moved so rapidly that it hydroplanes in the coating slurry. Loss of contact between the lap and the substrate lessens the surface fracture caused by large particles embedded in the lap. In some instances, a very smooth surface seems to have been produced largely by chemical erosion rather than by abrasion. Float-polished samples of fused silica were provided to LLNL by J. R. McNeil of the University of New Mexico, where a capability for float polishing had been refined during an extended visit by the developer of this procedure.⁴⁵

Nomarski microscopy showed that these surfaces consist of smooth areas bounded by large scratches believed to be residual grinding damage. Etching with mild HF did not reveal significant polishing damage in the smooth areas. At the large scratches, damage induced by 1-ns, 1.064- μm pulses occurred at fluences of 12 to 16 J/cm², which is slightly less than the 16 J/cm² median threshold for silica surfaces fabricated by standard abrasive polishing. The smooth areas had thresholds as high as 23 J/cm². Therefore, either extended application of float polishing to eliminate residual grinding damage or use of float polishing as a final step on carefully polished surfaces might be a useful technique for producing damage-resistant surfaces.

Crystals. We tested several crystals with nearly collimated beams that were approximately 2 mm in diameter. In such tests, damage may occur at either a front- or rear-polished surface, or at isolated inclusions in the bulk of the crystal. Only the threshold for the weakest part of a sample can

Table 6-6. Laser-damage thresholds measured with 1.064- μm , 10-ns pulses on selected crystals.

Material	Sample	Laser-damage thresholds (J/cm ²)		
		Front surface	Rear surface	Bulk
GSGG	9953	19 $\pm$ 2	16 $\pm$ 2	19 $\pm$ 2
Cr:Ni:GSGG	9954	16 $\pm$ 3	16 $\pm$ 3	>28
LiF	9952	>15	>15	8.8 $\pm$ 1.6
LiF	9957	14.6 $\pm$ 1.5	>24	12.4 $\pm$ 1.9
CaF ₂	9956	>12	>12	8.4 $\pm$ 0.8
CaF ₂	9951	>13	>13	10.7 $\pm$ 2.6
BaF ₂	9955	21 $\pm$ 4	21 $\pm$ 4	28 $\pm$ 3
BaF ₂	9950	22 $\pm$ 2	22 $\pm$ 2	>30

be measured unambiguously. If the bulk threshold is low, a surface threshold measured at higher fluences is of questionable significance. Similarly, scattering of laser light from surface damage might shield inclusions in the bulk.

Table 6-6 presents a representative set of thresholds measured with 10-ns, 1.064- μ m pulses. For a few crystals, surface and bulk thresholds were sufficiently alike that all the thresholds could be measured, but the limiting threshold for most of the crystals was that for bulk damage at inclusions. This has been a consistent observation in our tests during the last 6 years in more than 300 crystals of a wide variety,

Table 6-7. Desired properties for new Nd-doped laser glasses in future laser systems.

Megajoule system	
Fluorescence lifetime (μ s)	>280
Nd concentration (ions/cm ³)	$\sim 5 \times 10^{20}$
Emission cross section (cm ²)	<3
Thermal expansion (K ⁻¹)	<90 $\times 10^{-7}$
High-average-power system	
Fluorescence lifetime (μ s)	>180
Nd concentration (ions/cm ³)	$\sim 10^{21}$
Emission cross section (cm ²)	~ 4
Thermal figure-of-merit, ^a R_T (W/m)	>100

^a $R_T = K(1 - \nu)/\alpha E$, where K is thermal conductivity, ν is Poisson's ratio, S_T is the stress at material fracture, α is the coefficient of thermal expansion, and E is Young's modulus.

including both oxide and fluoride materials. Therefore, it is our opinion that elimination of bulk inclusions in materials selected for use in high-power lasers is a task of high priority.

UV Damage Studies for the USAF. Because coatings developed for Nova are potentially useful in several USAF applications, funds were provided to LLNL to establish two new test facilities for evaluating coatings produced at the Laboratory. The first facility will employ a 150-pps XeF laser, which emits 0.35- μ m, 20-ns pulses. This laser has been installed in Building 165, and a damage experiment is being assembled around the laser. The second laboratory will employ a 10-pps, flashlamp-pumped dye laser, which generates 500-ns, 0.7-J pulses at a frequency between 0.35 and 0.4 μ m. This laser will be installed in January 1985. With funding supplied by the USAF, we will fabricate sol-gel coatings on CaF₂ or silica substrates and measure the damage thresholds of these coatings as a function of pulse duration and repetition rate.

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Fig. 6-50. Phosphate glass systems studied by (a) Hoya Optics, Inc. and (b) Schott Glass Technologies, Inc. during 1983 and 1984.

Laser Glass

New Laser-Glass Compositions

In collaboration with Schott Glass Technologies, Inc. and Hoya Optics, Inc., we continued development of new laser-glass compositions in 1984. Our two main applications are for ICF lasers and for high-average-power (HAP) lasers. The two applications set somewhat different requirements for laser glasses, which are listed in Table 6-7.

Megajoule-class laser for fusion research require inexpensive glasses with good energy storage. Of primary importance in achieving good energy storage is a long fluorescence lifetime for high pumping efficiency and a stimulated-emission cross section of about 3 pm^2 or less to minimize the loss from amplified spontaneous emission (ASE). To lower production costs of glass for a megajoule laser, the glass should be produced in a system that incorporates continuous melting and forming rather than the discontinuous process used for Nova laser glass. The glass companies are designing such systems for production of phosphate glass.

HAP lasers also require glasses with a long fluorescence lifetime, but at higher Nd-doping levels than that for a megajoule laser glass. In addition, HAP laser glasses must have good thermomechanical characteristics, high mechanical strength, and low thermo-optic and stress-optic coefficients. We have focused on obtaining glasses with low concentration quenching and a high thermal figure of merit R_T , which is defined in Table 6-7.

The glass compositions studied by Hoya and Schott during 1984 are shown in Fig. 6-50. For comparison, we also show the compositions studied during 1983. Both companies discovered that silico-phosphate glasses have excellent characteristics for both megajoule and HAP lasers. The thermomechanical and spectroscopic properties of selected silico-phosphates, ultraphosphates, and some commercial laser glasses are listed in Table 6-8. The thermal figure of merit for all glasses investigated, except for LG-670, is below the desired 100 W/m for a HAP laser. However, we conservatively estimated the strength based on the measured fracture toughness and an assumed maximum surface-defect size of $100 \mu\text{m}$. We predict a possible increase in strength, by at least a factor of 2, for surfaces with smaller defects due to improved polishing

Table 6-8. Observed properties for laser glasses.

Glass	Emission cross section (pm^2)	Fluorescence lifetime at $10^{21} \text{ ions/cm}^3$ (μs)	Thermal conductivity ($\text{W/m}\cdot\text{K}$)	Thermal expansion ($10^{-7}/\text{K}$)	Young's modulus (GPa)	Poisson's ratio, ν	Tensile strength (MPa)	Thermal figure of merit (W/m)
Phosphates								
UP-16	4.0	210	0.49	131	45	0.270	37 ^a	22
UP-91	3.3	188	0.65	90	61	0.247	60 ^a	54
UP-148	3.9	182	0.63	99	54	0.247	— ^b	49
UP-120	3.0	—	0.77	93	64	0.242	— ^b	54
UP-136	3.4	133	0.91	64	76	0.212	— ^b	81
Silico-phosphates								
UP-160	3.8	186	0.69	80	56	0.250	— ^b	64
UP-137	2.9	155	0.91	64	73	0.210	— ^b	84
Z-85	3.2	184	0.84	93	69	0.231	55 ^a	55
H-310	3.9	198	0.80	80	67	0.230	55 ^a	63
Commercial phosphates								
LHG-5	4.1	135	0.77	98	64	0.184	90	90 (190) ^c
LHG-8	4.0	150	0.58	112	52	0.260	67	49 (102) ^c
LG-750	4.0	137	0.62	114	52	0.256	55	43
Q-100	3.8	180	0.73	96	73	0.24	65	59
Commercial silicates								
LG-670	2.7	45	1.35	93	94	0.242	100	117
LG-660	2.0	43	1.05	107	69	0.233	—	—

^aEstimates based on measurements of fracture toughness.

^bAssumed 55-MPa tensile strength.

^cStrengthened glasses demonstrated by Hoya Optics, Inc.

Laser Glass

techniques, chemical polishing, or surface coatings. Hoya has already demonstrated strengthened LHG-5 and LHG-8 glasses, as listed in Table 6-8.

Schott investigated silicate glasses for megajoule-laser applications. Such glasses are not as useful as phosphates because, at the Nd concentrations needed for efficient pumping, the Nd-fluorescence lifetimes are too short. Silicate glasses also have low emission cross sections, which makes energy extraction difficult because of the high laser fluences required.

A summary of the conclusions reached by Schott and Hoya is as follows:

- The obtainable fluorescence lifetime at the Nd doping of 10^{21} ions/cm³ is about 180 to 200 μ s for practical phosphate glasses. The longest lifetime found was 210 μ s for UP-16, which has poor durability and strength.

Fig. 6-51. Water-induced quenching in the phosphate glass UP-91.

Table 6-9. Increase in Nd-fluorescence decay rates due to OH⁻, Fe²⁺, and Cu²⁺ impurities in LG-660, UP-91, and UP-16.

Impurity	Glass	Increase in decay rate (s ⁻¹)		
		0.5×10^{20} ions cm ⁻³	5×10^{20} ions cm ⁻³	10×10^{20} ions cm ⁻³
OH ⁻ (see footnote a)	LG-660	70 (3.4%)	162 (4.2%)	319 (2.4%)
	UP-91	61 (2.3%)	163 (4.6%)	287 (5.1%)
	UP-16	59 (2.3%)	154 (4.4%)	303 (5.9%)
Fe ²⁺ (see footnote b)	LG-660	29 (1.4%)	43 (1.0%)	67 (0.5%)
	UP-91	51 (1.9%)	75 (1.9%)	125 (2.1%)
	UP-16	71 (2.7%)	120 (3.2%)	184 (3.2%)
Cu ²⁺ (see footnote c)	LG-660	25 (1.2%)	27 (0.7%)	37 (0.3%)
	UP-91	24 (0.9%)	36 (0.9%)	41 (0.6%)
	UP-16	30 (1.1%)	54 (1.6%)	69 (1.2%)

^aIncrease in decay rate (s⁻¹) per absorption coefficient (cm⁻¹) at the OH⁻ IR band peak. Peak = 2880, 2750, and 2750 cm⁻¹ for LG-660, UP-91, and UP-16, respectively.

^bPer 10⁻² cm⁻¹ absorption coefficient at 1000, 1100, and 1100 nm for LG-60, UP-91, and UP-16, respectively.

^cPer 10⁻² cm⁻¹ absorption coefficient at 960 nm.

- The obtainable fluorescence lifetime at the Nd doping of 5×10^{20} ions/cm³ is about 280 to 300 μ s.
- The presence of SiO₂ in phosphate glass lowers thermal expansion but decreases chemical durability. Al₂O₃ can be added to increase durability.
- The emission cross section can be controlled by varying the content of alkali, Al₂O₃, and MgO. However, concentration quenching becomes more severe for the lower-cross-section glasses with concentrations above 5×10^{20} Nd ions/cm³.

Author: S. E. Stokowski

Impurity Quenching

In 1984 we measured the Nd-fluorescence quenching rates induced by OH⁻, iron, and copper impurities, thus expanding our previous studies on self-quenching by Nd ions.⁴⁶ Schott Glass Technologies, Inc. prepared samples of the silicate glass LG-660 and two phosphate glasses UP-91 and UP-16 with these impurities added. Nominally, 30, 100, and 300 ppm of iron and copper were added to the melt. However, we determined the actual concentrations in the glass samples by measuring the optical absorption coefficients at the absorption band peaks.

The additional decay induced by the impurities appears to increase linearly with increasing impurity concentration. We show an example of the dependence in Fig. 6-51. Table 6-9 lists the additional decay rate induced per absorption coefficient (cm⁻¹) of the impurity at the given wavelength. The

Table 6-10. Spectroscopic properties of impurity ions in LG-660, UP-91, and UP-16.

Impurity	Glass	Peak wavelength (nm)	Absorption coefficient at peak per ppm impurity added to melt (cm^{-1})	Absorption coefficient at 1.053 μm per ppm impurity added to melt (cm^{-1})
Fe^{2+}	LG-660	990	1.1×10^{-3}	1.0×10^{-3}
	UP-91	1080	7.5×10^{-4}	7.3×10^{-4}
	UP-16	1060	5.9×10^{-4}	5.8×10^{-4}
Cu^{2+}	LG-660	750	6.7×10^{-4}	3.0×10^{-4}
	UP-91	900	2.8×10^{-3}	2.3×10^{-3}
	UP-16	920	2.5×10^{-3}	2.1×10^{-3}

table also lists the percentage increase in Nd decay rate over the decay rate extrapolated to the pure sample. All results in Table 6-9 are based on linear fits, which have regression coefficients greater than 0.98.

The OH^- impurity-induced decay rate is nearly independent of glass type. Between 5 and 10 wt% of Nd_2O_3 , the OH^- impurity-induced decay rate increases by a factor of 2; however, the rate only increases by a factor of about 2.5 between 0.5 and 5 wt%. Thus, the quenching rate is not strictly proportional to the Nd concentration. At Nd concentrations used in practical laser glass, water content is a critical factor in the pump efficiency of a laser. Most laser glass produced in a large furnace has an OH^- attenuation coefficient between 1 and 2 cm^{-1} . For maximum efficiency, we must keep this coefficient below 1 cm^{-1} .

Iron and copper are the impurities with the largest expected effect on Nd quenching because their energy levels match the $^4\text{F}_{3/2}$ initial-state energy of the Nd fluorescence. Table 6-10 lists the spectroscopic properties of Fe^{2+} and Cu^{2+} in LG-660, UP-91, and UP-16. Of particular significance is the large difference between silicate and phosphate glasses in the attenuation coefficient of Cu^{2+} per ppm of added copper. This difference could be due to the presence of some Cu^{1+} in the silicate. We have also found Fe^{3+} in both the silicate and phosphate glasses; however, its fraction of the total iron content is small, particularly in the more highly doped samples.

Both copper and iron are less effective in quenching the Nd fluorescence in LG-660 than in LG-750. This effect could be due to the 50%-smaller Nd transition strength in the silicate, which, according to the dipole-dipole quenching theory, should result in less quenching. For phosphate glasses, the percent increase in decay rate due to added iron or copper is approximately independent

of Nd concentration, which is a somewhat unexpected, but fortunate result.

In addition to their detrimental effect on the Nd decay rate, copper and iron also contribute to passive loss at the laser wavelength of 1.053 μm . The absorption coefficients of these ions at 1.053 μm are listed in Table 6-10 in terms of ppm impurity added to the melt. (The actual impurity concentrations in the glass samples are not known.) To keep the passive loss below 10^{-3} cm^{-1} , copper concentrations in laser glass melts must remain below 0.4 ppm and those for iron must remain below 1 to 2 ppm.

Author: S. E. Stokowski

Nd Spectra

During 1984 we observed the absorption spectra of the Nd laser transition $^4\text{F}_{3/2} \rightarrow ^4\text{F}_{5/2}$ in two laser glasses. This absorption is weak because the lower state is approximately 1800 cm^{-1} above the ground state and has a small population. The three reasons for our interest in this absorption are as follows:

- The attenuation coefficient due to this transition limits the maximum power available in a zig-zag laser, particularly when the glass temperature increases.
- The emission cross section for this transition can be determined from the absorption strength and its temperature dependence.
- A comparison of the fluorescence and absorption spectra yields information about the inhomogeneities of the Nd spectra.

We measured the absorption spectra in the phosphate glass LG-750 and the silicate glass LG-660. When the absorption spectra are compared to fluorescence spectra in Fig. 6-52, it is clear that the absorption and fluorescence peaks are at the same wavelength. The absorption cross section is

Fig. 6-52. Absorption and fluorescence spectra of the $^4F_{3/2} - ^4F_{5/2}$ transition of Nd^{3+} in the (a) phosphate glass LG-750 and (b) silicate glass LG-660 at 298 K.

Table 6-11. 1PA and 2PA coefficients in several transparent oxide glasses.

Glass	$\lambda_{\text{edge}} (\mu\text{m})^a$	α at 0.5525 μm (10^{-3} cm^{-1})	β at 0.5525 μm (cm/TW)	β at 0.532 μm (cm/TW)	β at 0.351 μm (Ref. 48) (cm/TW)
LG-650	0.318	1.2 ± 0.2	2.0 ± 1.0	—	—
LG-660	0.317	0.4 ± 0.1	0.3 ± 0.3	—	—
LG-670	0.376	1.1 ± 0.2	0.4 ± 0.4	—	—
LSG-91H	0.296	0.7 ± 0.1	0	—	—
BK-3	0.280	0.3 ± 0.1	≤ 0.4	0.6 ± 0.6	5.5
BK-7	0.312	0.6 ± 0.1	2.7 ± 1.5	2.9 ± 1.5	30
BK-10	0.275	0.6 ± 0.1	≤ 0.5	0.4 ± 0.4	15
LG-750	0.320	—	≤ 0.4	—	—

^a λ_{edge} is the wavelength at which the linear absorption coefficient α equals 1 cm^{-1} .

larger than the fluorescence cross section for shorter wavelengths; however, the opposite relationship holds for longer wavelengths. This expected behavior is due to thermalization among the substate: within the initial-state manifolds. The attenuation coefficient is $2 \times 10^{-4} \text{ cm}^{-1}$ in LG-750 and about 10^{-4} cm^{-1} in LG-660 per $10^{20} \text{ Nd ions} \cdot \text{cm}^{-3}$ at 298 K. Thus, at high Nd concentrations, the absorption due to Nd will become significant. The shoulder at $1.09 \mu\text{m}$ in LG-660 appears in both the fluorescence and absorption spectra. The shoulder is due to site-to-site inhomogeneities and not to the splitting of the $^4F_{3/2}$ state manifold.

Author: S. E. Stokowski

Measurements of Weak One- and Two-Photon Absorption in Oxide Glasses

As part of our research into the basic optical properties of laser glasses, we measured the two-photon absorption (2PA) coefficient

β of some highly transparent oxide glasses. We tested several alkali silicate compositions (LG-650, LG-660, LG-670, and LSG-91H), several borosilicates (BK-3, BK-7, and BK-10), and one phosphate (LG-750). Other than for fused silica⁴⁷ and BK-3, BK-7, and BK-10 (Refs. 47 and 48), no previous data exist for 2PA in laser glasses.

We used a previously described technique of dual-beam thermal lensing.⁴⁹ A 10-Hz pulsed dye laser [$\lambda = 0.5525 \mu\text{m}$, $\tau = 5 \text{ ns}$, full width at half maximum (FWHM)] was the pump beam, and a continuous-wave He-Ne laser was the probe. The peak intensity of the pump beam was about 2 GW/cm^2 in the samples, and each sample was about 1 cm thick. The samples exhibited no obvious absorption when measured in a Cary 17 spectrophotometer, but a weak one-photon absorption (1PA) coefficient α was detected with the thermal lensing technique.

Table 6-11 shows the measured values of the coefficients α and β . The magnitudes of the linear and nonlinear absorptions appear to be independent. The fact that nearly all of the nonlinear coefficients are within one

order of magnitude of each other suggests that some common, unidentified mechanisms may account for 2PA in the samples. For comparison, Table 6-11 also shows the values of β measured at 0.351 and 0.532 μm in borosilicate glasses as part of previous studies of two-photon-induced solarization.^{47,48} In the borosilicates, the values of β change little between 0.5525 and 0.532 μm , but they increase by about an order of magnitude at 0.351 μm .

Author: W. T. White, III

Nonlinear and Electrooptic Systems

Introduction

The engineering development of frequency-conversion and optical-switch technology is a high priority for the activation of Nova and for future ICF laser development. In 1984, we continued research on the technology of frequency conversion and studied the performance of the Nova frequency-conversion arrays. We also constructed a demonstration model of a 27- $\times$ 27-cm Pockel's cell using a technology scalable to future laser systems.

High-Dynamic-Range Quadrature Doubling Experiments

The quadrature second harmonic generation (SHG) scheme was designed to give high-

dynamic-range frequency conversion; that is, high conversion efficiency over a wide range of drive intensities. In this scheme, fully described in the *Laser Program Annual Report 83* (Ref. 50), the planes formed by the beam direction and the optic axes of two KD*P crystals of unequal thickness are arranged to be orthogonal (see Fig. 6-53). Thus, the green light from the first crystal passes through the second without back-conversion. The second crystal converts only the red drive left unconverted from the first crystal. The overall conversion efficiency is

$$\eta = \eta_1 + \eta_2(1 - \eta_1) \quad (39)$$

The quadrature concept was tested this year using the Cyclops laser to verify the high dynamic range. The experiment confirmed that quadrature schemes are capable of very high conversion efficiencies (>95%) over a wide range of drive intensities and that this efficiency can be attained using pulses that are Gaussian in space and time.

Two crystals of KD*P, with thicknesses of 12 and 44 mm, were used in the experiment. The deuteration level was measured, using nuclear magnetic resonance (NMR), to be 90.6%, and the optical absorption at 1.053 μm was measured to be 0.6%/cm for light appropriately polarized for Type II operation. The optical absorption at 0.532 μm is essentially zero. The Fresnel loss from the surfaces is 3.75%. Thus, the losses lie between 15 and 18.4%, depending on the conversion efficiency; the higher loss corresponds to lower efficiency. The 1.053- μm drive had a Gaussian profile in space and time with a pulse width (FWHM) of 1.2 ns and a diameter of 3 cm. It was close to

Fig. 6-53. The quadrature scheme for 2ω .

diffraction limited. The measured conversion efficiency, including all the losses, is shown as a function of the average drive intensity in Fig. 6-54. The efficiency rises above 70% at $\sim 200 \text{ MW/cm}^2$. In the region from 1 to 3 GW/cm^2 the efficiency was nearly 80%, which is very close to the maximum achievable. This experiment demonstrates the high dynamic range (~ 45 before surface losses are subtracted out) and the high efficiency of quadrature schemes. Using the AR coatings developed recently at LLNL, there

Fig. 6-54. Experimental data on a 12/44 quadrature system shows the high dynamic range and conversion efficiency this scheme permits.

Fig. 6-55. A thin-thick polarization-switched quadrature scheme for 4ω .

is every reason to expect that system efficiencies approaching 95% are possible.

Quadrature schemes are robust in that their sensitivity to beam profile, polarization errors, laser bandwidth, and angular and thermal misadjustments is lower than that of single crystal schemes. They will find many programmatic uses in the future.

A four-crystal quadrature scheme has been developed to give high dynamic range at 4ω . This is shown in Fig. 6-55. As yet, high dynamic range at 3ω appears to be impossible, although some four-crystal 3ω schemes do show a slight increase over the two-crystal doubler/mixer approach. However, their dynamic range is still less than 10.

Author: D. Eimerl

Nova Frequency-Conversion-Array Development

We have performed detailed second-harmonic generation (SHG) and third-harmonic generation (THG) measurements on the Cyclops laser system to support development of the Nova frequency-conversion array, as described in "Nova Frequency Conversion and Focusing Systems," in Section 2. Recent measurements using Nova production sol-gel coated $27 \times 27 \times 1\text{-cm}$ KDP crystals have resulted in 80% energy conversion to the

third harmonic as measured by calorimetry. Our results for frequency conversion are in reasonable agreement with computer-code calculations and indicate the high sensitivity of cascade THG to fundamental-beam polarization angle and to the alignment of the third-harmonic crystal.

We performed our measurements with the Cyclops laser system, which operates at the 1.053- μm Nova wavelength and generates a nominally 1.3-ns Gaussian laser pulse. We used a $2\times$ beam-reducing telescope and fundamental energies of 2.5 to 3.0 J to obtain peak intensities up to 5 GW/cm^2 in a nominally 1-cm-diam beam. Laser damage to 1.05- μm , 45-deg HR turning mirrors limited most measurements to intensities below 3.0 GW/cm^2 , but this

intensity was sufficient to cover the range of interest on Nova.

For our SHG experiments, shown in Fig. 6-56, we used a single 1.0-cm-thick Nova production crystal. SHG experiments were useful in verifying the accuracy of our calorimetry and crystal-alignment procedure with a vidicon camera system. For our THG experiments, shown in Fig. 6-57, we added the second 1.0-cm-thick Nova production crystal and performed sum-frequency mixing to generate the third-harmonic wavelength. Both $27\times 27\times 1\text{-cm}$ crystals were mounted in a precision aluminum egg crate designed specifically for these experiments. The crystals were held in the egg crate with absorbing-glass-backed buttons used on Nova, and the egg crate was mounted

Fig. 6-56. Experimental arrangement for single-crystal Type II SHG.

- A_1 - A_3 Apertures
 BG-18 Filter glass
 BS_1 - BS_3 Wedged beam splitters
 HR Highly reflecting mirror
 ND Neutral-density filter
 RG-1000 Filter glass

Fig. 6-57. Experimental arrangement for Type II cascade THG.

- A_1 - A_3 Apertures
 BS_1 - BS_2 Wedged beam splitters
 HR Highly reflecting mirror
 ND Neutral-density filter
 UG-11 Filter glass

inside a 25-in.-diam Aerotech mirror mount. Precision stepper-motor-driven micrometer screws, controlled with an LSI-11 minicomputer, were used to make fine adjustments in crystal angles. The egg crate was oriented in the mirror mount so that the crystal faces containing the optical axes would be parallel to the horizontal and vertical rotation axes of the mirror mount.

The crystal arrangement and polarization directions are shown in Fig. 6-58 (Ref. 51). The doubling crystal was oriented with its optic axis in the horizontal plane so that rotation about the vertical axis established the phase-matched condition. The mixing crystal had its optic axis oriented in the vertical plane so that horizontal-axis rotations were required to obtain phase matching. The angular rotation of the mirror mount was calibrated by reflecting a He-Ne laser off the surface of the first crystal at near-normal incidence and noting the beam deflection as a function of motor steps. Calibration was $0.74 \pm 0.01 \mu\text{rad}/\text{step}$.

The crystal alignment was accomplished by placing a corner cube in the beamline following both crystals and retroreflecting the 10-pps oscillator beam. Approximately 5% of the retroreflected beam was picked off with an uncoated 3-deg beam-splitting wedge (BS₁) and directed into a lens with a 4-m focal length. A vidicon camera was

placed in the focal plane of this lens, and the far field of the retroreflected oscillator beam was displayed on a TV monitor. With the corner cube blocked, the crystals were rotated so that the reflection from the first crystal would be coincident with the position previously established on the monitor screen by the corner cube. This established the "normal" position for the crystal pair. The egg crate was purposely wedged so that the retroreflection from the second crystal surfaces were displaced from the reflections of the first crystal by approximately $2200 \mu\text{rad}$ in the horizontal plane. This was done to minimize etalon effects between crystals and to allow a clear identification of each crystal. The etalon effects internal to each crystal were substantially eliminated by the sol-gel AR coatings. Internal etalon effects were observed in reflection as a function of crystal tuning angle, but they were not observed to degrade the frequency-conversion process.

The field of view of our vidicon camera system was approximately $900 \times 1125 \mu\text{rad}$. Changes in crystal angle as small as $20 \mu\text{rad}$ could be clearly detected. We estimate our precision in determining the normal position of the crystal at $\pm 50 \mu\text{rad}$. With an additional uncertainty of $\pm 50 \mu\text{rad}$ at the phase-matched position as a result of uncertainty in stepper-motor calibration, our

Fig. 6-58. Crystal arrangement and polarization directions for Type II cascade THG.

total alignment accuracy was estimated at $\pm 100 \mu\text{rad}$. The far-field of the Cyclops oscillator beam, as seen by the vidicon system, had a full angular divergence of approximately $125 \mu\text{rad}$.

Single-crystal SHG experiments were performed first. The crystal was nominally aligned for phase matching by rotation from the normal position about the vertical axis of the mirror mount by $-2380 \mu\text{rad}$. A negative angle of rotation in this case is a decrease in the angle of beam propagation to the optical axis of the doubling crystal. As seen in Fig. 6-56, three calorimeters were used with 5-mm-diam apertures and appropriate filter glasses to separate the second-harmonic and fundamental wavelengths. The input 1ω energy, output 2ω energy, and residual 1ω energy were measured on every shot, and the corresponding beam energies and fundamental input intensity were determined from the calculated ratios from the beam splitters. We estimate the accuracy

of the calorimetry at $\pm 2\%$ per station, and therefore the energy conversion ratio of second-harmonic to fundamental wavelengths could be determined to $\pm 4\%$.

A 2-in.-diam S-1 vacuum photodiode coupled to a TEK 7104 oscilloscope were used on occasional shots to monitor the temporal shape of the Cyclops laser pulse. Figure 6-59 shows a typical example of the pulse. From previous experience with

Fig. 6-59. Temporal shape of Cyclops laser pulse.

Fig. 6-60. Single-crystal Type II SHG conversion efficiency as a function of crystal angle from the normal position.

mode-locked pulses, we know that the TEK 7104 scope and vacuum photodiode combination exhibit a rise time of 0.5 ns. Assuming the detection system rise time adds to the laser-pulse full width in a root-mean-squared fashion, we estimated the laser-pulse full width to be 1.3 ± 0.1 ns. Therefore, in computing the fundamental input intensity (averaged over the calorimeter aperture), an uncertainty of $\pm 10\%$ has been assumed (including a calorimetry error of 2%).

The SHG conversion efficiencies data for our single-crystal experiments are shown in Fig. 6-60. The output $2\omega/1\omega$ energy-conversion ratio, the fraction of residual to input 1ω energy, and the fraction of total 1ω and 2ω to input 1ω energy, are plotted vs input 1ω intensity. Each data point is accompanied by a letter denoting the angular tuning of the doubling crystal from the normal position, described earlier. Data points lying on a vertical line are from the same laser shot. The data points are plotted against theoretical curves that are computed assuming a Gaussian temporal pulse shape and a crystal detuning of $100 \mu\text{rad}$ from perfect phase matching. The data were obtained in the same run and with a crystal temperature of $21.5 \pm 0.3^\circ\text{C}$.

Data points that range in angular tuning from -2315 to $-1515 \mu\text{rad}$ (points C to G)

are at or above the computed conversion curve. The points with the highest conversion were taken with angular tunings of -2115 and $-1915 \mu\text{rad}$. An angle from the normal of $-2371 \mu\text{rad}$ at 20°C (i.e., an internal angle of $-1587 \mu\text{rad}$) corresponds to the optimum phase-matched condition for the Nova KDP crystals. Applying a temperature correction of $+176 \mu\text{rad}/^\circ\text{C}$ gives a theoretical optimum angular tuning of $-2107 \mu\text{rad}$, which is in good agreement with our data.

For THG experiments, the mixing crystal was mounted to the egg crate following the doubler, with its optic axis in the vertical plane (orthogonal to the doubler). With the doubling crystal nominally aligned for SHG, the mixing crystal and doubler were rotated together about the horizontal axis by $+2371 \mu\text{rad}$. A positive rotation angle represents an increase in the angle of the propagation direction to the optic axis of the mixing crystal. This alignment procedure is equivalent to that performed on Nova.⁵¹ To obtain the optimum 35.3-deg , 1ω input polarization angle ψ (measured from the ordinary axis of the doubling crystal, i.e., the vertical axis), a half-wave plate was inserted before the first beam-splitting wedge and appropriately adjusted.

The calorimeter arrangement used for THG measurements was simplified somewhat from the earlier scheme. In addition to the input 1ω calorimeter, one calorimeter was placed after the crystals to measure the generated 3ω energy, and one calorimeter measured the total energy at all three frequencies. This arrangement allowed us to measure the $3\omega/1\omega$ energy-conversion ratio and the total ratio of transmitted energy to input 1ω energy, but not the residual 1ω and 2ω fractions. At high 1ω energies, the total-energy calorimeter had to be blocked to prevent damage. At $2 \text{ GW}/\text{cm}^2$ and below, the total-energy data indicated good agreement with the computed curves; thus, we were confident with respect to the accuracy of the calorimetry.

Two typical data runs are plotted in Fig. 6-61 as a function of input 1ω intensity. The theory curves in these two figures are based on $100\text{-}\mu\text{rad}$ angular detuning from perfect phase-matching of both the second-harmonic and mixing processes. All the data points are 10 to 20% above the computed conversion curve. Peak energy

Fig. 6-61. Type II cascade THG conversion efficiency with optimum crystal tuning angles from the normal position.

conversion of $80\% \pm 5\%$ occurs at an input 1ω intensity of approximately 2.25 GW/cm^2 and greater. Coating damage to the 45-deg turning mirrors prevented reliable measurements at higher input 1ω intensity, and this prevented us from seeing the expected rollover in conversion above 4 GW/cm^2 . The optimum angular tuning of the mixing crystal was predetermined by taking five to seven data points at 2 GW/cm^2 while "rocking" the crystal angle in steps of $100 \mu\text{rad}$. At temperatures slightly above 20°C , the optimum tuning for the mixing crystal appeared to be $+2580$ to $+2680 \mu\text{rad}$, which is greater than the theoretical 20°C value of $+2371 \mu\text{rad}$. The THG experiments were performed at two places on the crystal aperture: at 6 cm from the midpoint of one edge and at 17 cm from the midpoint. No differences in conversion efficiency or angular tuning were observed to exceed our stated error estimates.

The sensitivity of the cascade process to angular tuning of the mixing crystal for input 1ω intensities from 1.0 to 3.0 GW/cm^2 is shown in Fig. 6-62. Similar to the data in Fig. 6-61, the optimum mixing-crystal tuning corresponded to $+2680 \mu\text{rad}$. However, the input polarization angle for this data run

was inadvertently set at 37.7° rather than the optimum value of 35.3° . This probably accounts for the premature rollover in conversion efficiency after 2.0 GW/cm^2 . The data points for the $+2680 \mu\text{rad}$ detuning are nevertheless still above the theoretical curve in Fig. 6-62 based on perfect phase matching and a 37.7° -deg polarization.

The conversion efficiencies for the $+2380$ - and $+2480 \mu\text{rad}$ data points appear to fall off much more rapidly than those for the $+2828$ - and $+2925 \mu\text{rad}$ data points as a function of detuning from $+2680 \mu\text{rad}$. Assuming a temperature correction for the THG process of $+196 \mu\text{rad}/^\circ\text{C}$, the optimum angular tuning should be $+2567 \mu\text{rad}$ for a crystal temperature of 21°C . Much of our THG data indicates a slightly larger tuning as optimum for the mixing crystal. However, our estimated alignment accuracy of $\pm 100 \mu\text{rad}$ and uncertainty in temperature preclude making a definitive judgment based on the present data. The THG conversion efficiency was also briefly investigated for polarization angles less than 35.3° . The conversion efficiency for polarization angles of 33.3 and 34.3° deg is substantially degraded at input 1ω

Fig. 6-62. Type II cascade THG conversion efficiency as a function of second (mixing) crystal tuning from the normal position.

intensities above 2 GW/cm² compared to the 35.3-deg data.

In summary, our experiments indicate that sol-gel AR-coated Nova production crystals are capable of 80% energy conversion from the fundamental to the third-harmonic wavelength. Our data were obtained with a spatially and temporally well controlled laser beam and with crystal-alignment procedures that are functionally equivalent to those implemented on Nova. Our measurements also indicate the relatively high angular sensitivity of THG compared to single-crystal Type II SHG and demonstrate the necessity of having tight control of polarization angle for THG.

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Plasma-Electrode Pockels Cell

Large-aperture optical switches are of great potential importance for the realization of the next generation of high-energy

ICF lasers. Large-aperture switches may be used for the isolation of amplifier stages in master-oscillator power-amplifier (MOPA) chains, for injection and extraction from regenerative amplifier cavities, and for isolation from target retroreflections. Conventional technologies for optical switches are not scalable to the large apertures (1 to 5 m) required for the next generation of ICF lasers.⁵²

We proposed a new approach to the problem of designing large-area electro-optical switches and recently tested small-scale devices.⁵³⁻⁵⁵ We used ionized gas as large-area conducting electrodes for longitudinal electrooptic switches. Using plasma electrodes, we built a Pockels cell with an open aperture of 26 × 26 cm and a Z-cut KDP crystal with dimensions of 27 × 27 × 1.1 cm (Ref. 56). The device is the largest high-performance optical switch ever built and can now be developed for integration into new laser systems with large architectures.

Figure 6-63 is a photograph of the Pockels cell. The cell body was machined from Lucite and consists of two separate discharge chambers separated by the center section, which holds the electrooptic crystal. The crystal can be mounted "bare," with plasma directly in contact with its surfaces, or in a "sandwich" geometry, with the crystal located between two glass plates.⁵⁴ The bare-crystal Pockels cell has the advantage of mechanical simplicity and contains the smallest number of optical surfaces. However, the bare crystal requires development of high-quality AR coatings for KDP. Fortunately, recent developments in sol-gel coating technology now allow high-damage-threshold AR coatings to be applied to large-aperture KDP crystals. In addition, the bare crystal has high capacitance and a relatively low half-wave voltage. Thus, significantly slower turn-on times are obtained when compared to those for the sandwich configuration. In our 26- × 26-cm Pockels cell, we tested both the bare and sandwich configurations, and the results confirmed our expectations.⁵⁶

Figure 6-64 shows the four-electrode geometry for creating the plasma electrodes and the electrical circuit for applying the switching voltage. Pairs of electrodes on the same side of the crystal are used to establish plasma in both chambers, with minimum voltage applied across the crystal.

Fig. 6-63. 26- × 26-cm-aperture Pockels cell mounted on a large tilt-rotation table.

Following preionization, the main switching pulse is applied to a pair of electrodes on the opposing sides of the crystal. The voltage pulse charges the crystal faces using the plasma as the conducting electrode. A 1- μ s preionization pulse of nominally 30 kV breaks down the gas to form the plasma and, after a 200-ns delay, triggers the main switching pulser. A 50-ns transmission line at 25- Ω impedance (two RG 217 cables) couples the main pulser to the Pockels cell. The driving line is terminated at the cell with a 25- Ω load for the sandwich configuration and with a 10- Ω load for the bare-

crystal configuration. The ballast resistance for each preionization electrode is 250 Ω . Prior to operation, the discharge chambers are evacuated and filled with helium gas. Typical operating pressures are in the range of 5 to 10 Torr.

The Pockels cell was routinely operated at approximately 1 pps for many hours during the course of experiments, and no degradation of performance was observed. An occasional pumping out and refilling of the cell with helium was required due to a slow buildup of impurities in the gas from outgassing of the Lucite. Exposure to the

Fig. 6-64. Electrical circuit (four-electrode geometry) for driving the plasma-electrode Pockels cell.

Fig. 6-65. Optical layout for testing the switching performance of a Pockels cell at points A and B across the aperture.

plasma discharge caused no observable deterioration of the bare KDP-crystal surfaces.

Experiments were conducted with our previously developed 5-cm-aperture Pockels cell⁵⁴ to determine the effect of the plasma-electrode glow discharge on the colloidal silica sol-gel, AR coating devised recently for KDP crystals. We investigated both helium and argon discharges at pressures from 2 to 10 Torr. A soft sol-gel coating was deposited on two-thirds of the aperture of a 5- × 5-cm KDP crystal, and the remaining one-third was left uncoated to serve as a reference. The crystal was coated in the same manner on both sides and optimized for maximum transmission at 0.633 μm . A He-Ne laser beam was modulated with a chopper and split into two equal-intensity beams, one reflected from the coated surface and one from the uncoated surface. The KDP crystal was polished with parallel faces, and reflections from both were monitored. Phase-sensitive detection with lock-in amplifiers was used to minimize noise pickup from the discharge. A 20-kV, 10-Hz, 15- μs -long voltage pulse was applied to the KDP crystal. The leading edge of the pulse broke down the gas forming the glow-discharge plasma that constitutes the plasma electrodes of the Pockels cell.

With the virgin AR coating in the helium discharge at 10 Torr, the AR properties of the coating dramatically improved over ~ 120 min. The coating reflectivity was initially $\sim 0.6\%$ per surface and was reduced to less than $\sim 0.1\%$ per surface by the action of the discharge. With the discharge removed, the coating was placed under vacuum for several hours during which the

reflectivity gradually increased to its initial value. This process was repeated several times in helium and then in argon gas, and similar results were obtained. Eventually, however, the degree of improvement decreased, and the reflectivity of the coating increased to $\sim 2\%$ per surface. Repeated discharge in argon eventually caused coating deterioration after $> 3 \times 10^5$ shots. Over the course of these measurements, the forward-scattered light from the coated surfaces did not increase. The physical mechanisms responsible for these observations are not yet understood. However, the soft sol-gel AR coating appears to have good resistance to sputtering from the plasma-electrode discharge and would probably work successfully in a large-aperture Pockels cell.

The optical performance of the 26- × 26-cm Pockels cell was monitored with a single-mode continuous-wave (cw), argon-ion laser of nominally 1-W output at 0.514 μm to probe discrete points on the aperture, as illustrated in Fig. 6-65. A vacuum photodiode with an S-20 cathode and a subnanosecond response time was used as a detector. Polarizers were placed at both the laser and the detector and were oriented for minimum transmission. The argon-ion laser was aligned to propagate along the z axis of the crystal by monitoring the isogyre pattern at each probed point. The optical signals were recorded on a 7104 oscilloscope, and the photographs were digitized using a HP-9874 digitizer. From the optical transmission, the voltage at each probed point was calculated using the electrooptic coefficient for KDP.

Figure 6-66 shows the voltage measured directly on the electrodes for the bare-crystal configuration and the voltage deduced from optical transmission at two points on the aperture. Point A was midway up the crystal on the side nearest the main switching electrodes; point B was on the side nearest the preionization electrodes. A voltage of 1.5 times the bare-crystal half-wave voltage of 8 kV (V_{π}) was applied in this example. The voltages at points A and B rose to nearly the same value 70 ns after application of the main pulser voltage.

We define the optical risetime of the Pockels cell as the interval from the

Fig. 6-66. Voltage pulse V_P measured by a high-voltage probe across the switching electrodes. Voltages V_A and V_B were calculated from transmission data at two extreme points on the bare crystal.

beginning of transmission at point A to 95% transmission at point B. At 8 kV, corresponding to a half-wave retardation at $0.5145 \mu\text{m}$, the rise time of the bare-crystal Pockels cell was ~ 90 ns by this criterion. In this operating regime, the Pockels cell resembles a resistive transmission line, and the switching time is determined by the diffusion of charge across the aperture. The total preionizing current in these experiments was ~ 100 A. Further increase in the preionizing current as well as increase of the crystal thickness, which reduces the capacitance, will result in shorter switching times for the bare-crystal configuration.

The sandwich geometry was tested with ~ 6 -mm-thick glass plates mounted on each side of the KDP crystal. A 1-mm gap between each plate and the KDP crystal was filled with Koolase, which is an index-matching fluid for harmonic-generation crystals.²⁷ After a few days, we observed a chemical reaction between the Koolase and the Lucite walls of the Pockels cell. Even though this did not affect our measurements, we do not recommend this index-matching fluid for future work. It is also important that the pressure above the index-matching liquid be kept nearly the same as that in the discharge chamber. If the index-matching liquid is at a pressure

higher than that in the chamber, the glass plates will be placed under stress, thus reducing the optical figure of the Pockels cell. The plates might break if a sufficient difference in pressure occurs.

The results obtained with the sandwich configuration are shown in Fig. 6-67 for an applied voltage of 36 kV equal to the external half-wave voltage. The half-wave voltage increased by a factor of 4 compared to that for the bare crystal, and the switching time became significantly shorter. The turn-on time was limited by the high-voltage pulser rise time, which was ~ 20 ns for the $25\text{-}\Omega$ load used with the sandwich. A switching speed comparable to the electromagnetic propagation time across the crystal (~ 5 ns) should be obtainable by further lowering the high-voltage driver impedance and using a crystal that is 2 to 3 times thicker.

The field uniformity across the aperture was monitored by firing the pulser at 1.5 times the half-wave voltage and monitoring the optical transmission at many points. The results for the sandwich configuration show that after ~ 20 ns, all points on the aperture were charged to the same voltage within 10%, the size of our experimental error. To confirm this outcome, a diagnostic technique was developed

Nonlinear and Electrooptic Systems

Fig. 6-67. Voltage pulses V_P , V_A , and V_B for a sandwich geometry.

Fig. 6-68. Arrangement for the framing camera using a pulsed dye laser for full-aperture diagnostics.

to freeze-frame photograph the entire aperture during the turn-on of the optical transmission.

A framing camera was devised using a 5-ns pulse from a small nitrogen-laser-pumped dye laser operating at $0.550 \mu\text{m}$. The dye-laser pulse was synchronized with the high-voltage switching pulse and was incident on a ground-glass scatter plate placed in the nominal focal plane of a large ($70 \times 95\text{-cm}$) Fresnel lens. The arrangement is shown in Fig. 6-68. The scattered light was collimated by the Fresnel lens and overfilled the $26 \times 26\text{-cm}$ aperture of the Pockels cell. After the dye-laser light passed through the Pockels cell, a second large Fresnel lens imaged it to the input lens of a photographic film camera. A Polaroid sheet

polarizer was placed in the focal plane of a second Fresnel lens with sufficient angular acceptance to act as an analyzer. The timing of the dye-laser probe was adjusted relative to the high-voltage switching pulse by using various delay cables. In this way, the optical transmission of the entire aperture was freeze-framed at different times during the turn-on of the cell. Three photographs on Plus-X film are shown in Fig. 6-69. The data show excellent switching uniformity of the plasma electrodes after an $\sim 20\text{-ns}$ voltage propagation time across the aperture. The nonuniformities visible in the photographs are due to imperfections in the Fresnel lenses.

In conclusion, we have demonstrated that the plasma-electrode concept is scalable to

Fig. 6-69. Photographs from the framing camera show the turn-on of optical transmission across the 26- x 26-cm-aperture Pockels cell.

the large apertures that are of interest for new fusion lasers. The aperture can be further increased either by growing larger crystals or by segmentation. The switching times of our Pockels cell are short enough for most fusion-laser applications. Further engineering work should result in even faster large-area switches based on this concept.

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Nonlinear and Electrooptic Materials

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Frequency conversion and electrooptic switching are an important aspect of modern laser technology. Most practical lasers are not tunable, and the frequencies directly available from a lasing ion are discrete. Frequency conversion through harmonic generation, or sum and difference frequency generation, provides wavelengths in otherwise inaccessible regions of the spectrum. Electrooptic devices permit switching, optical isolation, and pulse-shape control. The cost and performance of a fusion laser is critically dependent on the properties of its nonlinear components; optimizing the nonlinear materials can increase design flexibility and reduce cost.

This year the program launched a new project in nonlinear materials. Its purpose is to identify materials whose properties (nonlinear and otherwise) permit reduced cost

and improved laser performance. The program's primary focus is on harmonic generation and electrooptic switching for fusion lasers; the secondary focus is on high-average-power lasers. Our scientific aim is to lay a solid basis for optimizing nonlinear materials for high-power-laser applications.

We had already analyzed one of the performance requirements for fusion lasers, high-dynamic-range frequency conversion.⁵⁸ The two-crystal quadrature system, invented last year, was predicted to have a dynamic range of ~ 50 for 100- μ rad laser pulses in potassium dideuterium phosphate (KD*P) and to be capable of increased wavelength, divergence, and polarization bandwidth. This year these predictions were confirmed experimentally on Cyclops; a dynamic range of ~ 70 with a peak conversion efficiency of over 95% was obtained with the 44/12 mm quadrature system described in "High-Dynamic-Range Quadrature Doubling Experiments," earlier in this section. The analysis was continued this year. We derived the range of material constants over which high performance can be achieved, and determined the relation of crystallographic symmetry to these requirements. The most important parameters for single-pulse systems are damage threshold and the behavior of the linear index as a function of wavelength, angle, and temperature. It is likely that there are many materials with sufficient nonlinearity to perform well; this is not considered to be the most critical parameter. It is the damage threshold and linear index that determine the utility of the material.

This year we began a survey of all known acentric crystals with useful transparency in the visible, UV, and near-IR ranges. A number of potentially interesting

new materials were identified. (Refer to "Crystal Symmetry and the Electrooptic Effect" in the next article.)

We also developed the concept of index engineering to obtain favorable phase-matching properties. This was tested successfully in the KDP group of isomorphs as reported in "New Harmonic Materials: Index Engineering" later in this section. With this series of homologs, we can do noncritical phase-matching (NCPM) at room temperature and at wavelengths ranging from 0.5 to 0.7 μm . This is expected to be a very important tool for materials engineering. Work on other crystal systems is planned.

Authors: D. Eimerl and S. P. Velsko

Crystal Symmetry and the Electrooptic Effect

The search for new electrooptic materials to improve laser performance and reduce cost requires assessing the potential of candidate materials for a desired application. If a set of criteria is used to set up screening tests, the number of candidates to evaluate can be reduced. One such criterion is the well-known requirement that there be no inversion center in the crystal structure. Of the 32 crystal classes only 18 are acentric. The use of von Neumann's principle further reduces the number of independent components the electrooptic tensor may have.⁵⁹

From knowledge of the crystal class alone, a considerable amount of information may be inferred. There are several other criteria that may be used, based on the intended application. Recent fusion research, reported in "Plasma Electrode Pockels Cell," earlier in this section, has led to a breakthrough in large-aperture optical-switch technology. The nonlinear material used in this device must satisfy several criteria, including:

- A high threshold for high-intensity laser-induced damage.
- A moderately large electrooptic coefficient.
- Near collinearity of the applied electric field with the direction of propagation of the laser beam.
- No depolarization of the laser beam during operation of the switch.

The first two criteria are probably not correlated strongly with the crystal class, but the last two are essentially symmetry requirements on the effective electrooptic tensor, and they might be expected to carry certain implications for the crystal symmetry of the electrooptic material.

The objective of the present study was to determine the relative merits of the acentric crystal classes for use in large-aperture electrooptic devices. This objective imposes two new requirements: the applied field must be nearly collinear with the beam direction, and there must be no depolarization of the beam.

We considered plane electromagnetic waves propagating in an arbitrary direction in an electrooptic crystal. We examined in detail the index changes and the rotation of the propagation eigenmodes that result when an electric field with an arbitrary direction is applied. The study extended to all 18 electrooptic classes, including the biaxial classes 2, *m*, 222, and *mm*2. We showed that there are indeed advantages to using crystals of certain classes; those advantages are dependent upon the geometry used and upon the performance requirements.

First, the optical properties of biaxial media were fully analyzed, and the perturbation theory for the distortion of the index ellipsoid was developed. We then treated the general problem in uniaxial media and in biaxial media. Finally, we analyzed the case of propagation along an optic axis and the role of permutation symmetry.

The index ellipsoid experiences both a rotation and a change in shape or birefringence as a result of the applied field. The tensor character of the effect is insufficient on its own to lead to any general conclusions about these perturbations. However, as the crystal symmetry increases, some simple geometric relations become evident. In particular, if the crystal symmetry is high enough, the ellipsoid rotation and the induced birefringence become orthogonal in the sense that they are proportional to orthogonal components of the applied field.

We begin with the crystals of lowest symmetry, the triclinic class 1, and consider successively higher crystal symmetries, ending with 432. For propagation in a general direction, there is insufficient symmetry in

any of the biaxial classes to draw any universal conclusions about the index ellipsoid. It appears that for the biaxial classes in general, the electrooptic effect depends on the relative magnitudes of the electrooptic coefficients and the linear refractive indices.

Consider first the orthogonality of the perturbations, in the sense described above. The eigenmode rotation χ and the birefringence Δn are linear in the applied field. That is,

$$\chi = \mathbf{Z} \cdot \mathbf{E} \quad (40)$$

and

$$\Delta n = \mathbf{W} \cdot \mathbf{E} \quad (41)$$

where $\mathbf{E}$ is the applied field. If the applied field is chosen to maximize the induced birefringence with no ellipsoid rotation, then the effective electrooptic coefficient is proportional to the sine of the angle between $\mathbf{Z}$ and $\mathbf{W}$. If $\mathbf{Z} \cdot \mathbf{W} = 0$, this geometric factor is maximal. Orthogonality, therefore, addresses the size of the effective electrooptic coefficient. Now, all classes except $\bar{4}$, $42m$, 23 , and $43m$ necessarily have $\mathbf{Z} \cdot \mathbf{W} = 0$.

Moreover, these classes also have orthogonal perturbations if the propagation direction is chosen appropriately. Thus, uniaxial crystals can definitely be configured to be orthogonal, whereas orthogonality for biaxial crystals depends on the electrooptic coefficients and the linear refractive indices.

If the direction of the applied field is arbitrary, then it is always possible to make the ellipse rotation zero by an appropriate choice of applied field direction. We require only that $\mathbf{W} \cdot \mathbf{E} = 0$. Then orthogonality guarantees that the induced birefringence is both finite and maximal. If the applied field is constrained to be collinear with the beam direction, then it is not always possible to have finite induced birefringence and zero rotation simultaneously. Again, for the biaxial classes, this depends on the electrooptic coefficients and the linear refractive indices. However, the uniaxial crystals divide into three categories:

- Those for which the rotation is identically zero ($4mm$, $6mm$).
- Those for which the rotation can be made to be zero for an appropriate choice of

beam direction ($\bar{4}$, 6 , $42m$, 3 , 32 , $3m$, $\bar{6}2m$, 23 , $43m$).

- Those that are unsuitable because the induced birefringence is zero if the rotation is zero (4 , 6 , 422 , 622).

Finally, we considered the case of propagation along an optic axis. To avoid significant beam depolarization, the eigenmodes must not rotate with the applied field. For some crystal classes, the eigenmodes are locked in place by virtue of the crystal symmetry. The eigenmode rotation is described in Tables 6-12 and 6-13. Locking occurs in the biaxial class m if the crystallographic mirror plane contains the optic axis. This also occurs in the biaxial class $mm2$, except when the optic axis is perpendicular to the two-fold crystallographic axis. For the other biaxial classes, 1 , 2 , and 222 , locking does not occur in general. Among the uniaxial classes, locking occurs for all the tetragonal and cubic classes and for the hexagonal classes 6 , 622 , and $6mm$. The trigonal classes, 3 , $3m$, and 32 , and the remaining hexagonal classes, $\bar{6}$ and $\bar{6}2m$, are unsuitable in this configuration.

If both the beam direction and the applied field are parallel to the optic axis, then the only uniaxial crystals that can generate birefringence are $\bar{4}$, $42m$, 23 , and $43m$. All members of these four classes exhibit eigenmode locking. For biaxial crystals the formulas change, but the conclusions relating to crystal class are the same as for general applied fields.

Crystal class	Tan 2α
4 , 422 , $4mm$, 6 , 622 , $6mm$	0
$\bar{4}$	$-r_{63}/r_{41}$
$42m$, 23 , $43m$	1
3 , $\bar{6}$	$\left(\frac{r_{11}E_2 + r_{22}E_1}{r_{22}E_2 - r_{11}E_1} \right)$
$3m$, $\bar{6}2m$	E_2/E_1
32	$-E_2/E_1$

Table 6-12. Eigenmode rotation for uniaxial crystals.

Plane	Class				
	1	2	m	222	$mm2$
$x-z$	y^a	y^b	$Yy c$	y	$Yy $
$x-y$	Y	Y	Y	y	y
$y-z$	Y	Y	Y	y	$Yy $

^a y indicates birefringence with $\mathbf{E} \parallel \mathbf{s}$.

^b y indicates the $\mathbf{Z}$ and $\mathbf{W}$ are orthogonal.

^c $|$ indicates that the eigenmodes do not rotate for $\mathbf{E} \parallel \mathbf{s} \parallel \mathbf{z}$, where $\mathbf{z}$ is the optic axis.

Table 6-13. Summary of electrooptic behavior in dielectric planes.

Permutation symmetry was not used to derive any of these conclusions, and it does not alter them in any way. The large-aperture requirements on electrooptic materials reduce the number of crystal classes from which a material may be chosen. This information will be useful in researching new electrooptic materials.

Author: D. Eimerl

Crystal Symmetry and Noncritically Phase-Matched Harmonic Conversion

The nonlinear material used for frequency conversion of fusion lasers must satisfy several criteria, including:

- A high threshold for damage induced by high-intensity lasers.
- A moderately large nonlinear coupling coefficient.
- Comparative insensitivity to crystal orientation relative to the laser beam.

The first two criteria are probably not correlated strongly with the crystal class, but the last is essentially a requirement on the crystal orientation. Together with the requirement that the effective tensor coupling does not vanish, this requirement has certain implications for the crystal symmetry and linear refractive indices of the electrooptic material.

The importance of NCPM in frequency conversion is well known. In general, efficient mixing or conversion requires that the wave velocities be precisely matched. Such matching can be accomplished using birefringent crystals, but the crystal must be oriented correctly. Usually the orientation must be controlled to within ~ 1 mrad. Clearly, the divergence of the laser beam must be limited to significantly less than the tolerance in the orientation angle. Under NCPM, the restrictions on the crystal orientation and the beam divergence are very forgiving; they may be as large as 100 mrad. This allows very efficient mixing, especially for beams with spatial modulations. Achieving NCPM requires special relations between the refractive indices of the crystal at the relevant frequencies.

On the other hand, the nonlinear coupling between the waves depends on the orientation of the crystal and on its crystallographic symmetry. At the crystal

orientations involved in NCPM, the nonlinear coupling between the waves is zero for some crystal classes.

We studied the conditions on the linear indices and crystal symmetry under which NCPM is possible with a nonzero nonlinear coupling. The goal of this study was to understand the limitations of crystal symmetry in NCPM three-wave mixing. A thorough analysis of phase-matching in biaxial media showed that the only way to achieve NCPM is to arrange for propagation along one of the dielectric axes. In uniaxial media, NCPM is achieved by propagation in any direction in the plane perpendicular to the polar axis. Our analysis of the coupling coefficients for these orientations is given below.

Crystal Symmetry and NCPM. Consider a biaxial crystal with normal dispersion and three dielectric axes labeled in order of their indices,

$$n_a > n_b > n_c, \quad (42)$$

which are assumed to be unequal. There are three dielectric axes and two types of phase-matching for each axis of propagation. The relevant couplings are summarized in Table 6-14.

In Type I phase-matching, the two lowest-frequency waves have the same polarization. In Type II phase-matching, the lowest-frequency waves have different polarizations. There are two kinds of Type II processes; in one, the lowest-frequency wave has the higher index polarization, and in the other, the lowest-frequency wave has the lower index.

Biaxial Crystals. There is no fixed relation between the crystallographic axes (abc) and the dielectric axes ($\alpha\beta\gamma$) in biaxial media. Therefore, it is not always possible to draw conclusions about the relation between crystal symmetry and NCPM.

NCPM requires that there be a coupling with two indices the same and the third index different. It is immediately clear that the crystal class 222 is not capable of NCPM.

For crystals of class 2, noncritical coupling is not possible for waves propagating along the b axis. For a -axis propagation, Type I phase-matching is possible only if $n_b < n_c$. For c -axis propagation, we require $n_b < n_a$. Thus, if the 2 axis is also the high-index axis, no noncritical processes are allowed.

Noncritical coupling is also possible for crystals of class *m*. For waves propagating perpendicular to the mirror plane, both Type I and Type II processes are possible regardless of whether $n_a > n_c$ or $n_a < n_c$. For propagation in the mirror plane, both types are again possible. If $n_a < n_b$, then the waves must propagate along the *c* axis. If $n_a > n_b$, they must go along the *a* axis. Given a sufficiently high birefringence, it is always possible to find a coefficient allowed by crystal symmetry that generates either a Type I or a Type II process.

For *mm2* crystals, there are no noncritical processes for waves propagating parallel to the 2 axis. If $n_c < n_b$, then both Type I and II processes exist for *a*-axis propagation. For *b*-axis propagation, a Type II process exists regardless of the relative magnitude of n_a and n_c , whereas a Type I process exists only if $n_c < n_a$. Given sufficiently high birefringence, it is always possible to find a coefficient allowed by crystal symmetry that generates either a Type I or a Type II process for *m* crystals.

These results for biaxial crystals are summarized in Table 5-15, which shows the combinations of birefringence Δn and coupling constants d_{eff} required to achieve NCPM for Type I and Type II processes, with the axis of propagation and biaxial crystal class. This table applies only to crystals for which three indices are clearly distinct. Biaxial crystals that are nearly uniaxial are discussed below.

Uniaxial Crystals. Two of the indices are equal in uniaxial crystals. If $n_y = n_z$, then the crystal is positively birefringent, and the *a* axis is the polar axis (*c* axis) of the crystal. If $n_y = n_x$, then the crystal is negatively birefringent, and the *y* axis is the polar axis of the crystal. Recall that NCPM is achieved for propagation in any direction perpendicular to the polar axis. The waves therefore propagate in the *a*-*b* plane at an angle ϕ to the *c* axis.

For Type I processes, the effective coupling for positively birefringent crystals is

$$d_{\text{eff}} = -Sd_{13} + Cd_{23} \quad (43)$$

for negatively birefringent crystals, it is

$$d_{\text{eff}} = S^2d_{31} - 2SCd_{36} + C^2d_{32} \quad (44)$$

where

$$C = \cos \phi, \text{ and } S = \sin \phi$$

Type	Axis		
	<i>a</i>	<i>b</i>	<i>c</i>
I	d_{cbb}	d_{caa}	d_{baa}
II	d_{cbc}	d_{cac}	d_{bab}

Table 6-14. Tensor couplings required to achieve NCPM in biaxial crystals.

		Crystal class			
		Axis	2	<i>m</i>	<i>mm2</i>
Type I processes					
d_{12}	$n_a < n_b$	<i>c</i>		x	
d_{13}	$n_a < n_c$	<i>b</i>		x	
d_{21}	$n_b < n_a$	<i>c</i>	x		
d_{23}	$n_b < n_c$	<i>a</i>	x		
d_{31}	$n_c < n_a$	<i>b</i>		x	x
d_{32}	$n_c < n_b$	<i>a</i>		x	x
Type II processes					
d_{15}	$n_a < n_c$	<i>b</i>		x	x
d_{16}	$n_a < n_b$	<i>c</i>	x		
d_{24}	$n_b < n_c$	<i>a</i>		x	x
d_{26}	$n_b < n_a$	<i>c</i>		x	
d_{34}	$n_c < n_b$	<i>a</i>	x		
d_{35}	$n_c < n_a$	<i>b</i>		x	

For Type II processes, the effective coupling for positive crystals is

$$d_{\text{eff}} = S^2d_{15} - SC(d_{14} + d_{25}) + C^2d_{24} \quad (45)$$

and for negative crystals, it is

$$d_{\text{eff}} = -Sd_{35} + Cd_{34} \quad (46)$$

The restrictions placed by crystal symmetry on the tensor coupling have profound consequences for NCPM in uniaxial crystals. First, there are no uniaxial crystals for which any of the couplings d_{13} , d_{23} , d_{35} , or d_{34} are allowed. Therefore, we can immediately conclude that Type II processes for negative crystals, and Type I processes for positive crystals, cannot exhibit NCPM. Second, the relevant coefficients (or combinations) for the remaining two cases, Type II positive and Type I negative, are all zero for crystal classes 422, 32, 6, 622, and 6mm2.

Therefore, NCPM is not possible for these classes regardless of the sign of their birefringence, and only eight uniaxial crystal classes are capable of supporting NCPM: 4, $\bar{4}$, 4mm, 4m2, 3, 3m, 6, and 6mm. Only positively birefringent crystals can support Type II NCPM, and only negatively birefringent crystals can support Type I NCPM. These are severe restrictions on the crystal class and birefringence.

Finally, the effect of Kleinman symmetry, the symmetry of the tensor coupling under interchange of any of its indices, is to set $d_{14} = d_{25} = 0$ in the classes 4, 422, 3, 32, 6, and 622. This does not alter any of the

Table 6-15. Coupling constants d_{eff} for all combinations of birefringence Δn and crystal class that support NCPM in biaxial crystals.

conclusions drawn about the effect of crystal symmetry on NCPM, although it may possibly affect the magnitude of d_{eff} . We conclude that permutation symmetry is not relevant to this issue.

NCPM in Nearly Uniaxial Crystals. If the optic angle of a biaxial crystal is close to 0 or 180 deg, then two indices are nearly equal, and the angular sensitivity of the phase-matching is small for waves propagating in any direction in the plane spanned by the two, nearly equal directions. For example, in a crystal where $n_b \sim n_c$, the angular sensitivity is small for propagation perpendicular to the a axis, that is, in the b - c plane itself. This leads us to re-examine the couplings for biaxial classes with this extra degree of freedom in the direction of propagation.

If the a and b axes are the nearly equal index axes, then the polar axis is the c axis, and the same formulas derived above for uniaxial crystals apply to this case. The couplings in the cases where the a or b axis is the polar axis are obtained by cyclic permutation of the tensor components in these formulas.

The extra feature of the nearly uniaxial case is that by propagating in a plane away from the dielectric axes, it is possible to couple the waves through the tensor components d_{14} , d_{25} , and d_{36} . These three tensor components can support Type I NCPM for negative crystals and Type II NCPM for positive crystals, and they are present in all biaxial crystals. Thus, besides the processes summarized in Table 6-15, these additional processes are possible for biaxial crystals whose optic angle is close to 0 or 180 deg.

This class of materials is quite large and is of particular interest as a way to reduce the angular sensitivity of frequency-conversion crystals. Refer to Zumsteg, Bierlein, and Gier⁴⁰ for a discussion of KTiOPO_4 , one of the materials in this class.

Author: D. Eimerl

Electrooptics in the Biaxial Crystal BaMgF_4

In a longitudinal electrooptic switch that utilizes a uniaxial crystal such as KDP, a laser beam is propagated along the optic axis, which is also the direction of the applied electric field. In the absence of an

electric field, no birefringence is observed along the optic axis. When a voltage is applied to transparent electrodes or ring electrodes to establish the electric field, the uniaxial crystal exhibits birefringence along the optic axis, and the induced phase retardation between "fast" and "slow" rays is independent of crystal thickness. For this reason the crystal thickness need not be rigorously constant across the aperture, and we also find no first-order sensitivity to misalignment for small angular deviations. However, when a biaxial crystal is used in an electrooptical switch, only one of these two properties can be obtained.

A biaxial crystal has two optic axes, both contained in one of the three dielectric planes, and it exhibits a natural birefringence when a laser beam is propagated down any one of the three dielectric axes. If the direction of propagation is along an optic axis, we find considerable sensitivity to misalignment, but the crystal thickness is not critical. If, however, the direction of propagation is along one dielectric axis, then the crystal thickness must be precisely controlled, but we find a reduced sensitivity to angular misalignment. There is a considerable difference between the voltages required to induce complete switching in these two configurations.

We considered several orthorhombic, biaxial crystals as materials for use in electrooptic switches, and we investigated their angular sensitivity for propagation along an optic axis and their thickness tolerances for propagation along a dielectric axis. In particular, we measured the half-wave voltages for longitudinal and transverse configurations in BaMgF_4 and compared the measurements with theoretical estimates.

The BaMgF_4 crystal we used was grown by H. Guggenheim⁶¹ at the Bell Laboratories and oriented, cut, and polished at LLNL. A sample with dimensions of $8.6 \times 8.0 \times 7.5$ mm was used for the measurements reported here. A low-power He-Ne laser at $6328 \text{ \AA}$ served as the probe beam.

Placing the crystal between crossed polarizers and propagating rays close to one of the optic axes produced the isogyre pattern characteristic of a biaxial crystal that is shown in Fig. 6-70. For comparison, the isogyre pattern for a 10-mm-thick uniaxial KDP crystal is shown in Fig. 6-71. As we can see from these pictures, the acceptance

angle for KDP is one order of magnitude greater than that for BaMgF₄ in the sensitive direction. For propagation in the x - y plane of BaMgF₄ that contains the optic axes (see Fig. 6-72), the two indices along the eigen directions are determined by

$$n_o = n_z, \quad (47)$$

and

$$\left(\frac{1}{n_c}\right)^2 = \frac{\cos^2 \theta}{n_x^2} + \frac{\sin^2 \theta}{n_y^2}. \quad (48)$$

We specify the ordinary polarization n_o as that along the z axis and the extraordinary polarization n_e as that contained in the x - y plane. θ is the angle of the propagation direction with respect to the y axis of the crystal. For BaMgF₄ at 1.06 μm (Ref. 61),

$$n_x = 1.4674, \quad (49)$$

$$n_y = 1.4436, \quad (50)$$

and

$$n_z = 1.4604. \quad (51)$$

Thus, the indices for both polarizations are equal at the angle $\theta = 34$ deg, which specifies the optic-axis direction.

The acceptance angle for small deviations from one of the optic-axis directions can be calculated from the angular sensitivity of the index of refraction. By differentiating Eq. (48), we find

$$\frac{\partial n}{\partial \theta} = n_z^3 \sin \theta \cos \theta \left(\frac{1}{n_x^2} - \frac{1}{n_y^2} \right). \quad (52)$$

For an angular deviation $\Delta\theta$ in the x - y plane from the optic axis, with path length along the optic axis l and wavelength λ , the induced phase difference is

$$\Delta\phi = \frac{2\pi\Delta\theta}{\lambda} \left(\frac{\partial n}{\partial \theta} \right), \quad (53)$$

and the transmission with the crystal placed between crossed polarizers is given by

$$T = \sin^2 \left(\frac{\Delta\phi}{2} \right). \quad (54)$$

Fig. 6-70. Isogyre pattern (transmission through crossed polarizers) of BaMgF₄ for propagation along one of the optic axes.

Fig. 6-71. Isogyre pattern of KDP for propagation along the optic axis.

Fig. 6-72. Dielectric and crystallographic axes of BaMgF₄ (point group-orthorhombic $mmm2$). The optic axes are in the x - y plane at $\theta = 34$ deg and 1.064 μm .

Assuming we are willing to tolerate a leakage transmission of $T \leq 1\%$, we obtain $\Delta\phi \leq 0.2$ rad for a 1-cm crystal and an angular acceptance of $\Delta\theta \leq 0.15$ mrad from Eq. (53). A similar calculation for a 1-cm-thick KDP crystal gives an angular acceptance of $\Delta\phi \leq 10$ mrad for $\leq 1\%$ transmission between crossed polarizers.

The calculation of the electrooptic effect for propagation along the optic axis is straightforward. Using first-order

perturbation theory, we can write expressions for the indices of refraction in the presence of an external electric field E :

$$\frac{1}{n_o^2} = \frac{1}{n_z^2} + M_{33} \quad , \quad (55)$$

and

$$\frac{1}{n_e^2} = \frac{1}{n_z^2} + (S_x^2 M_{22} + S_y^2 M_{11}) - 2S_x S_y M_{12} \quad , \quad (56)$$

where

$$M_{11} = \Delta \left(\frac{1}{n^2} \right)_1 = r_{11} E_x + r_{12} E_y + r_{13} E_z \quad (57)$$

$$M_{22} = \Delta \left(\frac{1}{n^2} \right)_2 = r_{21} E_x + r_{22} E_y + r_{23} E_z \quad (58)$$

$$M_{33} = \Delta \left(\frac{1}{n^2} \right)_3 = r_{31} E_x + r_{32} E_y + r_{33} E_z \quad (59)$$

$$M_{12} = \Delta \left(\frac{1}{n^2} \right)_6 = r_{61} E_x + r_{62} E_y + r_{63} E_z \quad (60)$$

and the appropriate direction cosines are $S_x = \sin \theta$ and $S_y = \cos \theta$.

BaMgF₄ belongs to the $mm2-C_{2v}$ point group for which the electrooptic tensor is given as

$$\begin{array}{ccc} 0 & 0 & r_{13} \\ 0 & 0 & r_{23} \\ 0 & 0 & r_{33} \\ 0 & r_{32} & 0 \\ r_{51} & 0 & 0 \\ 0 & 0 & 0 \end{array} \quad (61)$$

This sets most of the terms in Eqs. (57) through (60) to zero and simplifies Eqs. (55) and (56) to:

$$\frac{1}{n_o^2} = \frac{1}{n_z^2} + r_{33} E_z \quad , \quad (62)$$

and

$$\frac{1}{n_e^2} = \frac{1}{n_z^2} + (S_x^2 r_{23} + S_y^2 r_{13}) E_z \quad . \quad (63)$$

Only the z component of the applied electric field induces phase differences. Thus, a nonlinear electrooptic effect occurs with the electric field along the optic axis. The crystal must be tilted in such a way that there is a finite projection of the field on the z axis. From Eq. (63) we obtain

$$\Delta n = n_o - n_e \approx -\frac{n_z^3}{2} \times E_z (r_{33} - r_{23} S_x^2 - r_{13} S_y^2) \quad . \quad (64)$$

Writing the path length along the optic axis as l and the thickness of the crystal along the z direction as d , we can calculate the half-wave voltage,

$$V_\pi = \frac{d}{l} \frac{\lambda}{n_z^3 (r_{33} - r_{23} S_x^2 - r_{13} S_y^2)} \quad (65)$$

applied along the z direction. The r_{ij} coefficients have not been measured for BaMgF₄. However, a range of values, estimated from the known SHG coefficients using formulas based on wave-mixing and on anharmonic oscillator treatments of the electrooptic effect are given as

$$r_{13} = -0.076 \text{ to } -0.20 \times 10^{-12} \text{ m/V} \quad , \quad (66)$$

$$r_{23} = +0.150 \text{ to } +0.40 \times 10^{-12} \text{ m/V} \quad , \quad (67)$$

$$r_{33} = +0.056 \text{ to } +0.15 \times 10^{-12} \text{ m/V} \quad , \quad (68)$$

$$r_{51} = +0.120 \text{ to } +0.37 \times 10^{-12} \text{ m/V} \quad , \quad (69)$$

and

$$r_{42} = +0.080 \text{ to } +0.22 \times 10^{-12} \text{ m/V} \quad . \quad (70)$$

Assuming $l/d \approx 1$, we calculate from Eq. (65) a half-wave voltage V_π ranging from 2.13 to 5.71 MV at $\lambda = 1 \mu\text{m}$, and 1.27 to 3.4 MV for He-Ne lasers.

Because V_π depends on the difference between several electrooptic coefficients, and because there is some uncertainty about relative signs and values of these coefficients, V_π can significantly differ from these estimates. Thus, experimental measurement is essential.

The experimental setup is shown in Fig. 6-73. The voltage is applied along the z axis, and the laser is propagated along the

Fig. 6-73. Experimental configuration for measurement of half-wave voltage V_{π} ; electric field applied along the z direction and propagating down the optic axis in BaMgF_4 .

optic axis. The incident polarization was aligned at 45 deg to the two polarization eigen directions of the crystal, and a $\lambda/4$ plate was inserted before the crystal for maximum sensitivity. The optical transmission with the crystal placed between crossed polarizers is given by

$$T = \sin^2 \left(\frac{\pi}{2} \frac{V}{V_{\pi}} + \phi \right), \quad (71)$$

where the $\lambda/4$ plate biases the transmission to the 50% point ($\phi = \pi/4$). An 8.5-kV pulse of ~ 1 -ms duration produced a modulation $f = 0.5\%$. The half-wave voltage from Eq. (71), when $V \ll V_{\pi}$, is given by

$$V_{\pi} = \frac{\pi V}{2 f} \approx 2.7 \text{ MV}. \quad (72)$$

This value is in the mid-range estimated for the He-Ne laser.

Now consider a different geometry for a Pockels cell with the beam propagating down the z axis and the applied field along the z direction. Because the material is birefringent, the thickness variation will give a fringe pattern when the crystal is placed between the crossed polarizers. For 1% leakage transmission, we require

$$\sin^2 \left(\frac{\pi \Delta l (n_x - n_y)}{\lambda} \right) \leq 10^{-2}, \quad (73)$$

which gives $\Delta l \leq 1.2 \mu\text{m}$. Because we are presently fabricating KDP crystals for Nova with this degree of precision,⁶² fabrication of other crystals should also be possible.

For this longitudinal arrangement, the alignment tolerance is sensitive to tilts only

in the second order. For a 1-cm-thick BaMgF_4 crystal, the alignment sensitivity is close to 10 mrad, a value comparable to that for KDP crystals.

The half-wave voltage for this geometry is given by

$$V_{\pi} = \frac{\lambda}{r_{23}n_y^3 - r_{13}n_x^3}. \quad (74)$$

For $\lambda = 1 \mu\text{m}$, Eq. (74) yields 560 kV to 1.52 MV for V_{π} . We measured V_{π} as ≈ 800 kV at $0.63 \mu\text{m}$. When this experimental value is scaled up to 1.3 MV for a $1\text{-}\mu\text{m}$ laser, the half-wave voltage is near the top end of the estimated range.

A somewhat lower V_{π} was observed for a transverse configuration with the laser propagating down the x axis and the applied field along the z direction. In this configuration, we measured V_{π} as ~ 600 kV. When extrapolated for $\lambda = 1 \mu\text{m}$, this measurement becomes 1.0 MV, somewhat less than our estimated values of 1.5 to 3.7 MV.

The transverse electrooptic effect can be used with large laser beams by tilting a thin plate of the electrooptic material so that a component of the electric field (applied with plasma electrodes) is perpendicular to the direction of propagation inside the material. We term this novel geometry the "quasi-transverse" configuration. This is shown in Fig. 6-74 for BaMgF_4 . This configuration will also work with transverse electrooptic crystals (linear and quadratic effects), isotropic Kerr-effect glasses, transparent ceramics, and liquids. For BaMgF_4 , however, a large-aperture application in the quasi-transverse geometry increases the applied voltage requirement by a factor of

Fig. 6-74. Quasi-transverse, large-aperture Pockels cell; applied field E (transparent plasma electrodes) normal to crystal surface; component E_z orthogonal to propagation direction λ inside BaMgF₄ crystal.

Crystal	Direction of propagation			
	Optic axis		z axis	
	V_{π} (1.06 μm)	$\Delta\theta$ (mrad)	V_{π} (1.06 μm)	$\Delta\epsilon$ (μm)
BaMgF ₄	4.5 MV ^a	0.15	1.3 MV ^a	1.2
BaZnF ₄	9.1 MV ^b	0.12	240 kV ^b	1.3
Li(COOH)·H ₂ O	58 kV ^c	0.026	79 kV ^c	0.3
a-Gd(MoO ₄) ₃	26 kV ^b	0.3	49 kV ^b	80.0

^aMeasured at 6328 Å and scaled to 1.06 μm .

^bCalculated from r_{ij} coefficients scaled from SHG coefficients.

^cCalculated from known coefficients.

Table 6-16. Electro-optic switching with biaxial materials.

$\sim \cot 30 \text{ deg} = 1.73$. This requires that $V_{\pi} > 1.7 \text{ MV}$ for a 1- μm laser, therefore this approach has no advantage over the longitudinal configuration along the z axis. Although clean measurements were not possible with the present technique, the V_{π} for the other configurations were much higher.

We have concluded that BaMgF₄ is not a good candidate for an electrooptic switch because the voltages required are one order of magnitude too high for practical applications. However, there are other biaxial crystalline materials with the same symmetry that may be more useful. Table 6-16 lists some of the more promising candidates and the estimated switching voltages. The half-wave voltages, angular sensitivities, and thickness tolerances of these crystals must be confirmed experimentally before the materials can become serious candidates for electrooptic switches using plasma electrodes.

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New Harmonic Materials: Index Engineering

The goal of advanced nonlinear materials research is to develop the capability of matching material parameters to a desired application. Although it is conceivable that enough materials could be found to cover all applications, this is considered unlikely. Rather, we expect to identify a suite of pure materials that sparsely cover a broad range of operating conditions and to develop the skills to engineer or fine-tune nonlinear materials to specific applications. Our survey work to find a suite of base materials is in progress, and we have developed the technique of index engineering, successfully applying it to the KDP series of homologs.

Index engineering is a method of adjusting the refractive indices of a material to obtain NCPM over a range of wavelengths. The index is adjusted either by doping the crystal or by making solid solutions. The index adjustments affect the lattice constants and the UV absorption edge, changing both the birefringence and the wavelength dispersion. By an adroit choice of materials, some decoupling between the changes in the birefringence and the dispersion can be arranged.

The concept was tested using the group of KDP homologs. Table 6-17 lists the wavelengths at which NCPM was achieved for Type I doubling. (The coupling for Type II doubling vanishes at NCPM in the KDP group.) The following substitutions and solid solutions were attempted:

- $\text{Rb}_x\text{K}_{1-x}\text{H}_2\text{PO}_4$,
- $\text{KH}_{2(1-x)}\text{D}_{2x}\text{PO}_4$,
- $\text{KH}_2(\text{PO}_4)_{1-x}(\text{AsO}_4)_x$,
- $\text{RbH}_2(\text{PO}_4)_{1-x}(\text{AsO}_4)_x$, and
- $\text{NH}_4\text{H}_2(\text{PO}_4)_{1-x}(\text{AsO}_4)_x$.

Some work on D-H and K-Rb substitution has been reported,⁶³⁻⁶⁵ but none has been reported on the P-As substitution. If the full range of solid solutions exists, then presumably NCPM is possible anywhere between 0.515 to 0.699 μm .

RDP-KDP. Crystals were grown from aqueous solution containing redissolved, pure end-member crystals. The solutions and crystals were analyzed to determine segregation coefficient and phase purity. The proportion in solid solution is quite small, $0 < x < 0.15$, and at the upper end of this range the segregation coefficient for rubidium depends strongly on the concentration of rubidium in the solution. For $x = 0.11$, the segregation coefficient is 7, but for rubidium concentrations in solution greater than $C_s(\text{Rb}) = 70\%$, the concentration in the crystal appeared to be inhomogeneous and some cracking occurred. For $C_s(\text{Rb}) < 70\%$, the crystals appeared to be homogeneous. The homogeneity was tested by observing the dependence of the doubling efficiency to 0.266 μm on the crystal orientation. The lobes and zeros of the $\sin^2 \theta$ distribution were clearly observable in a 1-cm crystal, indicating that the phase matching was not disturbed. At the highest concentrations, the NCPM wavelength was 0.530 μm . Thus, Rb-K substitution gives about the same variation in NCPM as D-H substitution, 0.515 to 0.530 μm .

Arsenate-Phosphate Series. Crystals were grown from aqueous solution containing redissolved, pure end-member crystals. The atomic concentrations in the solutions were 0, 0.2, 0.4, 0.6, 0.8, and 1. Attempts to grow them using arsenic acid rather than the alkali salt were unsuccessful. The solutions and crystals were analyzed chemically and with x-ray diffraction to determine the segregation coefficients and phase purity. The complete solid solution exists for both Rb and K salts, and the segregation coefficient for As in the phosphate crystal is 0.8; that is, there is proportionately more As in the crystals than in the solution.

Specific crystals were grown in an attempt to fabricate a crystal that exhibits

Crystal	λ (μm)
KDP	0.515
KD*P (90%)	0.527
ADP	0.528
RDP	0.623
KDA	0.609
ADA	0.587
RDA	0.695

Table 6-17. NCPM wavelengths in KDP homologs.

NCPM at room temperature for $\lambda = 0.532 \mu\text{m}$. If the NCPM wavelength varies linearly with the As concentration, the appropriate solid solutions are

- $\text{KH}_2\text{P}_{0.82}\text{As}_{0.18}\text{O}_4$ and
- $\text{NH}_4\text{H}_2\text{P}_{0.97}\text{As}_{0.03}\text{O}_4$.

The crystals were grown by the temperature-lowering technique from pure KDP or ADP seeds in a solution with the appropriate molar fractions of the end-member compounds. The As concentration in the ADP/A crystal was 4%, and the crystal was not phase-matchable at 0.532 μm , in agreement with the assumption that NCPM is linear with As concentration. The As concentration in the KDP crystal was 21%, and its phase-matching angle at 0.532 μm was 81 deg. Apparently, NCPM in KDP/A is not linear in the As concentration.

The homogeneity of the KDP/A crystal was tested by observing the $\sin^2 \theta$ dependence of the doubling efficiency. The first lobe and zero of the distribution were only barely observable in a 1-cm crystal, indicating that the phase matching was disturbed by the crystal inhomogeneity. Clearly, some further work on the As-P solid solutions is necessary.

Author: D. Eimerl

Rapid Growth of KDP Crystals

Introduction. Our immediate objective during 1984 was to increase the growth rate of KDP single crystals from the industrial rate of 1.5 mm/day to about 15 mm/day. The higher rate would reduce the time for bulk growth from one year to one month and effect a roughly proportionate decrease in production cost. In addition, the broader objective of this research is to establish a scientific and engineering basis for the development of economical processes for the growth of any optical crystal from solutions or fluxes under analogous conditions of mixed kinetic and mass-transport control.

Fig. 6-75. Three-zone crystallizer used to define system mass and heat transfer requirements for high-rate crystal growth at constant temperature and supersaturation.

In the first year of this project, which was initiated in 1984, we identified three fundamental limitations of current industrial practice:

1. Bulk transport of dissolved KDP into the growth chamber is insufficient to support high growth rates.
2. The rate of solute transport across the diffusion boundary layer adjacent to the growing crystal is insufficient.
3. The industry does not continuously measure or control supersaturation, which is the driving force for crystal growth.

Accordingly, we have performed a heat and mass-transport analysis of a canonical growth system with separate chambers for crystal growth, stock dissolution, and superheating. This analysis of multizone crystallizers yielded the transfer rates required for high-rate growth. We developed a semi-empirical model using hydrodynamics, mass-transport, and growth kinetics to predict the operating conditions necessary for high-rate growth in three candidate cell configurations. This model has been verified for 20- to 25-mm crystals grown at rates as high as 1 mm/h using a turbine to enhance solute flux to the pyramidal faces. Finally, we identified three in situ, non-destructive techniques for the determination of supersaturation that are based on differential measurements of solution density, buoyancy, and electrical conductivity.

Analysis of Multizone Crystallizers.

Phillips Laboratories was subcontracted to

develop a growth process⁶⁶ with separate vessels for crystal growth, feed stock dissolution, and superheating. The theoretical groundwork for the design and operation of multizone crystallizers was done here.⁶⁷ The three-zone system in Fig. 6-75 was chosen to illustrate the theoretical findings, which are more broadly applicable. The mathematical procedure we developed relates process-control parameters (flow rates and zone temperatures) to the growth of individual crystal faces at a fixed temperature and supersaturation. Heat and mass transport requirements for high-rate growth were determined in terms of generalized kinetic and hydrodynamic models under conditions of 3-D crystal growth. Maintenance of constant temperature and supersaturation at the crystal surface requires time-dependent control over the rate of feed of solution to the growth chamber. Requirements for heat removal from the chamber (a function of feed rate) can be offset by increases in solute feed concentration and corresponding increases in solution temperature at the inlet to the growth chamber. Procedures were formulated to optimize feed rate, temperature, and heat rejection, but no cell configuration was identified for the optimization of these factors.

Growth of Crystals at High Rates Under Mixed Kinetic and Mass Transport Control. We determined the operating conditions (flow rates, temperature, and supersaturation) required for the growth of KDP

crystals in three candidate cell designs at specified rates.⁶⁸ First, a kinetic equation for KDP growth in the [001] direction was derived from published rate data for 3-mm crystals grown in a solution of specified bulk concentration under impinging flow.⁶⁹ Concentration at the crystal surface was determined by application of mass-transport theory, in which transport by convective-diffusion to the right-angle pyramid was approximated by equations for a right-angle cone.⁷⁰ Resultant values of surface concentration c_s , absolute temperature T , and the rate of KDP growth in the [001] direction g_z were reduced by linear regression to a kinetic equation of the following form:

$$g_z = k_R [\exp(-H_R/RT)] (c_s - c_{sat})^n \quad (75)$$

where c_{sat} is the solubility of KDP in units, kg-KDP/kg-H₂O. The regression analysis yielded a rate constant, $k_R = 3492$ m/s; an activation energy, $H_R = 41,399$ kJ/mol; and an effective reaction order, $n = 2.277$. The required average flux to the pyramid is then given by

$$j = g_z d_{KDP} \frac{1}{2} \quad (76)$$

where d_{KDP} is crystal density.

Next, transport models were chosen for three configurations of potential interest to the crystal grower:

1. Pyramidal growth under axial flow.
2. Growth of sector disks or blocks cut from the (101) plane with perpendicularly impinging flow.
3. Single crystal surfaces positioned in one wall of a narrow, rectangular duct-flow channel in the laminar and turbulent regimes.

To calculate the average flux to the pyramid in the first configuration, we use the correlation for cones, taking the pyramid slant height as the characteristic dimension X in lieu of the cone generatrix:

$$j = 0.8 D \left(d_b \frac{c_b}{1 + c_b} - d_s \frac{c_s}{1 + c_s} \right) \times \left(\frac{U}{X\nu} \right)^{0.5} \left(\frac{\nu}{D} \right)^{0.33} \quad (77)$$

where U is the potential flow rate, D is the diffusion coefficient, d_b , and d_s are bulk

and surface solution densities, respectively, c_b is the bulk concentration, and ν is kinematic viscosity.

We determined the solubility and density of KDP solutions over the operating range of potential interest ($t = 30$ to 90°C) and reduced the data by linear regression to equations of the form:

$$c_{sat} = k_1 + k_2 t + k_3 t^2 \quad (78)$$

and

$$d = d_1 + d_2 c + d_3 c^2 + d_4 t + d_5 t^2 \quad (79)$$

Finally, Eqs. (75) to (79) were solved numerically using correlations for the kinematic viscosity and diffusion coefficient⁶⁹ to yield a functional dependence G of growth rate on temperature, relative bulk supersaturation $[\sigma = (c_b - c_{sat})/c_{sat}]$, and a hydrodynamic parameter U/X :

$$g_z = G(T, \sigma, U/X) \quad (80)$$

Curves of constant growth rate are plotted on axes T and σ in Fig. 6-76, which also shows the metastable limit of the solution

Fig. 6-76. Constant growth rate g_z as a function of temperature and bulk relative supersaturation for a fixed hydrodynamic parameter $U/X = 30 \text{ s}^{-1}$. The broken line indicates the reported metastable limit of the solution.

Fig. 6-77. Lines of constant growth rate ($g_z = 15 \text{ mm/d}$) for various values of U/X vs temperature and relative bulk supersaturation.

according to Loiacono.⁶⁶ The effect of the hydrodynamic parameter U/X is shown in Fig. 6-77. Analogous results are obtained for the flow channel in either turbulent or laminar flow, when appropriate transport correlations are used in place of Eq. (77).

In summary, growth rates of 15 mm/day or greater are predicted for relative

supersaturations below 0.1 and at temperatures above 50°C in both impinging and parallel flow systems. At higher temperatures, the required supersaturation appears to decrease relative to the metastable limit of the solution. The ratio of velocity to crystal dimension is a measure of the degree of kinetic control; for fixed growth rates, the growth process is controlled completely by kinetics as U/X approaches infinity, and by mass-transport as U/X approaches zero.

Experimental Studies of Crystal

Growth. Crystal boules 20 to 25 mm thick ($X = 14$ to 18 mm) were grown at rates up to 25 mm/day using a novel turbine (shown in Fig. 6-78) to enhance mass transport. The crystal seed was attached to the turbine so that the axis of rotation coincided with the z axis of the crystal.

Figure 6-79 shows the growth rate g_z , computed on the basis of Eq. (80), the integrated rate (crystal tip displacement), and the experimental displacement. The agreement between model and experiment is excellent in view of the tenfold increase in crystal size from that of the rate data⁶⁹ on which the model is based, extrapolation to higher temperatures, and the approximations involved in the transport model. Growth rates of 5 to 25 mm/day were achieved.

Preliminary measurements of laser damage to samples grown at rates of 5 to 20 mm/day indicate damage thresholds that

Fig. 6-78. A turbine with peripheral stators used in studies of the growth of crystals at high rates (≤ 1 mm/h).

Fig. 6-79. Comparison of the computed growth rate g_z , the integrated growth rate (i.e., computed crystal tip displacement), and the experimental displacement for a typical growth experiment.

are comparable to or greater than the thresholds for industrial crystals grown at 1 to 1.5 mm/day.

The Measurement of Supersaturation.

The driving force for crystal growth is related to the difference between surface concentration and solubility. The narrow range of accessible metastable concentrations requires that bulk concentration and solubility each be measured with great precision (at least to 0.1%) if the growth rate is to be accurately controlled (to 5%). This can be done by measuring the difference between a property of the growth solution (e.g., buoyancy, electrical conductivity, refractive index, or density) and that of a reference solution of known concentration. We used differential hydrometry to detect changes in concentration to levels below $1 \text{ g} \cdot \text{KDP/kg} \cdot \text{H}_2\text{O}$. Systems based on differential buoyancy or electrical conductivity that are appropriate for industrial use have been designed.

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Advanced Crystalline Laser Materials

Introduction

Recent analysis shows the potential for solid-state laser media to produce the high average output power at high efficiency required for eventual operation of a fusion power plant.^{71,72} For such high-average-power operation, the solid-state gain medium is configured as face pump disks, with heat removed by flowing gas across the disk faces. Achievement of this system requires an integrated and self-consistent program of research and development on improved pumping sources, sensitizing and activator ions, and solid-state host materials. Our current analysis of each of these areas suggests the feasibility of overall laser efficiencies in excess of 10%. The solid-state media under consideration include both crystalline materials and composite glass structures consisting of a thin layer of active material with layers of nonabsorbing glass for mechanical support.

In the following three articles, we summarize progress in the development of

advanced crystalline laser materials. The scope of this project includes spectroscopic analysis of microcrystals, synthesis and laser testing of larger ($\sim 1\text{-cm}$), single crystals and the ultimate deployment of large-disk-geometry crystal slabs (tens of centimeters) in high-average-power laser configurations.

Authors: J. A. Caird and M. D. Shinn

Spectroscopic and Laser Studies of Gadolinium Gallium Garnets

Our current work on gadolinium scandium gallium garnet (GSGG) is motivated by several factors. It is the first crystal to be identified that can be grown in sufficient size and with an optical quality consistent with the cross section of the Nd active ion and the ASE and parasitic limits to permit validation of the design concepts for high-average-power lasers. It is important to recognize that the nonlinear refractive index of GSGG is too large for this material to be considered for ultimate application in a fusion power plant. Nonetheless, in addition to serving as an appropriate surrogate for design-concept studies, it has other potentially useful applications.

The base crystal, gadolinium gallium garnet (GGG), with the chemical formula $\text{Gd}_3\text{Ga}_5\text{O}_{12}$, was developed by the electronics industry for magnetic bubble-memory substrates. For this application, 10-cm-diam crystals are routinely grown by the Czochralski technique. An example of a typical commercial GGG single crystal is shown in Fig. 6-80.

Both GGG and several substitutional variants of GGG have been used as host crystals, which can be readily doped with

Fig. 6-80. Production boule of gadolinium gallium garnet (GGG) of 8-cm aperture and 25-cm length (Material Progress, Inc., Santa Resa, Calif.).

Fig. 6-81. Absorption spectrum of GSGG:Cr,Nd showing broad absorption bands due to Cr^{3+} and sharp lines due to Nd^{3+} . The emission spectrum of Cr^{3+} is also shown.

Fig. 6-82. GSGG:Nd $^4\text{F}_{3/2} \rightarrow ^3\text{I}_{11/2}$ fully resolved fluorescence spectrum was used to determine the line width $\Delta\nu_s$ and the inter-Stark component branching ratio β_s for the laser transition $R_2 \rightarrow Y_3$.

active laser ions. GSGG [chemical formula $\text{Gd}_2(\text{Sc,Ga})_2\text{O}_{12}$] doped with chromium and neodymium ions is a particularly attractive laser material for the aforementioned reasons.⁷³ The absorption spectrum of a typical GSGG:Cr,Nd crystal is shown in Fig. 6-81. The principal advantage of this crystal is its efficient absorption of flashlamp light in the visible portion of the spectrum. Light absorbed in the broad Cr^{3+} absorption bands is efficiently transferred to Nd^{3+} ions, from which it is extracted as laser emission. Efficient energy transfer occurs through nonradiative dipole-dipole interaction due to the overlap of the Cr^{3+} emission band with

strong Nd^{3+} absorption lines, as shown in Fig. 6-81.

To evaluate the potential of GSGG:Cr,Nd as a laser material, we have made several spectroscopic and laser measurements. Thermomechanical properties were also studied and will be published in a forthcoming report. The peak stimulated-emission cross section for Nd^{3+} in GSGG was determined through analysis of its absorption and fluorescence spectra. The emission cross section is given by

$$\sigma = \frac{\lambda^2 \beta_m \beta_s}{4\pi^2 n^2 c \Delta\nu_s \tau_r} \quad (81)$$

where τ_r is the radiative lifetime of the upper laser level, β_m is the inter-manifold branching ratio for the $^4\text{F}_{3/2} \rightarrow ^3\text{I}_{11/2}$ group of transitions, β_s is the branching ratio of the transition between the individual Stark-split components of the upper- and lower-energy manifolds, $\Delta\nu_s$ is the line width (FWHM) of the homogeneously broadened Stark component of Lorentzian line shape, λ is the laser wavelength, and n is the refractive index of the material at that wavelength.

A radiative lifetime of 278 μs was measured in low-concentration samples. An intermanifold branching ratio, $\beta_m = 0.51$ was calculated using the Judd-Ofelt theory,^{74,75} with intensity parameters determined by analysis of the Nd^{3+} absorption spectrum.

The Stark component line width, $\Delta\nu_s = 11.5 \text{ cm}^{-1}$, and branching ratio, $\beta_s = 0.35$, were obtained by analysis of the fluorescence spectrum shown in Fig. 6-82. The index of refraction, $n = 1.942$, at the laser wavelength, $\lambda = 1.061 \mu\text{m}$, was measured at the Hughes Research Laboratories⁷⁶ using an undoped sample of GSGG supplied by LLNL. Equation (81) thus gives an emission cross section of $\sigma = 1.41 \times 10^{-19} \text{ cm}^2$. This value is lower than that for YAG:Nd by about a factor of 2 but is 3 times higher than typical values for Nd doped glasses. Perhaps the most significant impact of this result lies in its impact on the control of ASE in slab-geometry lasers. At a constant level of stored energy per unit volume, a GSGG:Cr,Nd slab can have twice the linear dimensions of a YAG:Nd slab.

We performed a series of experiments to determine the level of loss per unit length in GSGG:Cr,Nd laser rods and its effects on laser performance. We measured losses in the laser rods with a modified version of the technique first proposed by Findlay and Clay.⁷⁷ In their method, the threshold for laser oscillation is measured in a laser resonator as a function of the reflectivity of the output coupler, which is varied by changing mirrors. In a free-running laser resonator, the threshold condition for laser action can be expressed as

$$R_{\text{out}} \cdot \exp(2\sigma N_{\text{th}} - \delta L) = 1 \quad (82)$$

where R_{out} is the reflectivity of the output coupler, N_{th} is the population density in the

excited state at threshold, δ is the loss per unit length, and L is the length of the laser rod. Taking the natural log of Eq. (82) yields

$$2\sigma N_{\text{th}} L - 2\delta L = -\ln R_{\text{out}} \quad (83)$$

For the Findlay-Clay analysis, it is not necessary to know the inversion density in absolute terms. Any signal or quantity that is proportional to N_{th} can be used. Findlay and Clay proposed the use of preshot energy in the capacitor bank of the flashlamp pulse-forming network (PFN) as a quantity that is directly proportional to the inversion density. We attempted to use this method, but we found that the assumption of proportionality between capacitor energy and inversion density led to systematic errors, especially at low input energy, where the data are most important. The discrepancies appear to arise when a reasonable fraction of the stored electrical energy is consumed in the lamp-ignition process.

To avoid the aforementioned systematic errors, we decided to use the $1.3\text{-}\mu\text{m}$ fluorescence from the laser rod to produce a signal proportional to the excited-state density. The experimental arrangement is shown in Fig. 6-83. M_1 is the output coupler with reflectivity R_{out} , which is varied by changing out this mirror. M_2 is the rear mirror of the laser resonator, with $\sim 100\%$ reflectivity at $1.06 \mu\text{m}$ but very high transmission ($T > 90\%$) at $1.3 \mu\text{m}$. The fluorescence is imaged through a field stop onto an InAs detector cooled to liquid-nitrogen temperature. A narrow-band filter is also

Fig. 6-83. Apparatus for measurement of $1.3\text{-}\mu\text{m}$ fluorescence intensity from a laser rod simultaneous with detection of laser action at $1.06 \mu\text{m}$.

used to filter out $1.06\text{ }\mu\text{m}$ signals and to reduce any effect of scattered flashlamp light. A small correction factor is used to account for signal variations caused by changes in the $1.3\text{-}\mu\text{m}$ reflectivity of the output coupler M_1 .

The results of the loss measurements on four 0.25-in.-diam $\times$ 3-in.-long GSGG:Cr,Nd rods obtained from four different vendors and one YAG:Nd rod of the same size are shown in Fig. 6-84. From Eq. (83), it is evident that the intercept of the straight lines in the figure with the horizontal axis is equal to $-2\delta L$ or the round-trip loss factor for the resonator. The loss per unit length of laser crystal is obtained by dividing this number by twice the length of the rod. The losses measured for these and a few other rods are given in Table 6-18.

The lowest loss was obtained for the YAG:Nd rod, with a round-trip loss of 5%, or 0.3 cm^{-1} . The two best GSGG:Cr,Nd

rods measured to date yielded a round-trip loss of approximately 12%, or about 0.8 cm^{-1} . Also listed in Table 6-18 are the attenuation coefficients expected in these materials due to absorption from the thermally preplated lower laser level at the laser wavelength and room temperature. The insertion loss in the best GSGG:Cr,Nd rods is several times the expected intrinsic attenuation and indicates that there is room for significant improvement in material quality.

Several GSGG:Cr,Nd laser rods exhibited loss levels $>2\text{ cm}^{-1}$. Several different sources of loss have been observed. These sources include (1) distortion due to embedded strain in the crystal; (2) scattering from visible defects in the crystals, such as dislocations, inclusions and bubbles; and (3) linear absorption due to impurities. The magnitude of the contribution from each of these mechanisms to the total measured loss varies from rod to rod. Improving crystal perfection is a major goal of our near-term efforts with respect to crystal growth.

We also measured the lasing efficiency of the free-running oscillators for each rod with a number of output coupling mirrors. The best results for each rod are listed in Table 6-18. Even for moderately lossy rods, the GSGG rods all performed better than YAG:Nd. The lowest loss GSGG rods showed almost twice the efficiency of the YAG rod in the same laser head, despite higher losses in the GSGG rods. The output energy vs input energy for the best GSGG rod tested to date is shown in Fig. 6-85. The YAG:Nd results for the same resonator are plotted for comparison.

The value for the stimulated-emission cross section of Nd^{3+} in GSGG, determined

Fig. 6-84. Comparison of loss measurements on GSGG:Cr,Nd and YAG:Nd.

Table 6-18. Loss and efficiency measurements on 0.25-in.-diam $\times$ 3-in.-long laser rods.

Rod	Host crystal	Output efficiency (slope)	Output efficiency (absolute)	Insertion loss ^a (10^{-2} cm^{-1})	$\alpha(\text{Nd})$ at $1.06\text{ }\mu\text{m}$ (10^{-2} cm^{-1})	Nd concentration 10^{20} cm^{-3}	Cr concentration 10^{20} cm^{-3}
43833	YAG	3.8	2.8	0.30	0.2 ^b	1.4 ^b	0
9B	GSGG	5.0	3.5	0.85	0.19	3.0	0.29
52-1	GSGG	6.4	4.9	1.29	0.08	1.3	0.92
55-1	GSGG	6.5	4.9	0.79	0.16	2.5	0.92
57-1	GSGG	4.9	3.0	3.31	0.13	2.0	0.92
004-2	GSGG	5.9	3.9	2.13	0.15	2.3	2.0
77A4	GSGG	5.1	3.3	3.10	0.08	1.3	1.0
60-1	GSGG	7.0	5.0	0.78	0.13	2.0	2.0

^aFindlay-Clay plot.

^bEstimate based on assumed 1% doping.

above, places a constraint on the aperture of the active components because of ASE. This, in turn, limits the maximum power output per active laser component. If the Nd^{3+} cross section can be reduced without compromising other properties, scaling to larger apertures is possible.

A host material that may satisfy this criterion is GGG with a structure modified by calcium and zirconium ions [GGG(Ca,Zr)]. This modified version of GGG was also originally synthesized for use as a magnetic bubble memory substrate; hence, knowledge of how to grow large-diameter crystals already exists.⁷⁸

Crystals of GGG(Ca,Zr):Cr, GGG(Ca,Zr):Nd, and GGG(Ca,Zr):Cr,Nd were grown by F. Bruni at Material Progress Corp., using the Czochralski method. The composition of the latter two crystals was determined using x-ray fluorescence.

The fluorescence due to the ${}^4\text{F}_{3/2} \rightarrow {}^4\text{I}_{11/2}$ transition of Nd^{3+} in GGG(Ca,Zr) is shown in Fig. 6-86. When compared to GSGG:Nd (see Fig. 6-82), a large degree of inhomogeneous broadening exists. We believe this inhomogeneous broadening is due to non-equivalent Nd^{3+} sites, which are the result of the many different positions the charge-compensated ion pairs (Ca-Zr) may assume in the lattice. Using radiative rates and branching ratios calculated from the Judd-Ofelt theory^{74,75} and the measured fluorescence spectrum, the peak cross section, referenced to the ${}^4\text{F}_{3/2}$ manifold population, was $7.1 \times 10^{-20} \text{ cm}^2$ using the method of Caird et al.⁷⁹ This value is less than the equivalently defined value for Nd^{3+} in GSGG by almost a factor of 2. Another consequence of the increased inhomogeneous broadening is a reduction in the degree of concentration quenching, compared to that in either GSGG:Nd or YAG:Nd.

The sensitization of Nd^{3+} fluorescence by Cr^{3+} ions was also studied and found to be comparable to that observed for GSGG:Cr,Nd. A summary of some of the spectroscopic properties of GGG(Ca,Zr):Cr,Nd is given in Table 6-19.

The laser performance of a 0.25-in.-diam $\times$ 3-in.-long rod of GGG(Ca,Zr):Cr,Nd was measured in the free-running mode. The best results are shown in Fig. 6-87. The best performance of a YAG:Nd laser rod in the same cavity is also shown. Even though

Fig. 6-85. Laser output energy vs input energy (capacitor bank) for GSGG:Cr,Nd and YAG:Nd in the same free-running laser resonator. Output-coupler reflectivity was 49%.

Fig. 6-86. Emission spectrum due to the ${}^4\text{F}_{3/2} \rightarrow {}^4\text{I}_{11/2}$ transitions of Nd^{3+} in GGG(Ca,Zr) at room temperature.

Table 6-19. Spectroscopic properties of ∇ GGG(Ca,Zr):Cr,Nd.

Judd-Ofelt Parameters (10^{-20} cm^2)	
Ω_2	0.27
Ω_4	3.03
Ω_6	3.54
${}^4\text{F}_{3/2}$ Judd-Ofelt calculated radiative lifetime (μs)	244
Measured Nd lifetime (μs)	246
${}^4\text{F}_{3/2}$ Stark splitting (cm^{-1})	57
${}^4\text{F}_{3/2} \rightarrow {}^4\text{I}_{11/2}$ peak wavelength (μm)	1.061
${}^4\text{F}_{3/2} \rightarrow {}^4\text{I}_{11/2}$ peak cross section (10^{-20} cm^2)	7.1
${}^4\text{A}_2 \rightarrow {}^4\text{T}_2$ peak wavelength (μm)	0.64
${}^4\text{T}_2$ fluorescence lifetime (μs)	113
Nd^{3+} segregation coefficient	0.56
Cr^{3+} segregation coefficient	2.5

Fig. 6-87. Laser performance of GGG(Ca,Zr):Cr,Nd and of YAG:Nd in the same cavity.

the former rod was of poor optical quality, with observable banding of the Cr^{3+} density and internal strain, it still performed better than YAG:Nd. We believe that better performance can be expected if material of higher optical quality can be grown.

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Studies of Other Crystal Laser Candidates

Our interest in other crystal systems is motivated by material requirements for ultimate laser fusion applications. For such applications, GSGG:Cr:Nd and other oxide gallium garnets are excluded by their large nonlinear refractive indices. In addition, the linear refractive index of GSGG is large ($n \approx 1.94$), which leads to a reduced fluorescence lifetime. Furthermore, the emission cross section is too large for fusion applications.

One promising alternate concept is Cr,Nd co-doping of a fluoride host crystal. The lower crystal field in fluoride hosts facilitates efficient energy transfer to the Nd ion by broadening Cr emission bands

and increasing their overlap of the Nd absorption bands.

$\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}:\text{Cr}$. The fluoride garnet crystal $\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}:\text{Cr}$ is isomorphic to oxide garnets, such as YAG. Our crystals were obtained from H. Guggenheim at AT&T Bell Laboratories, Murray Hill, N.J., who grew them in a horizontal zone furnace. The crystals contained few cracks and had no observable change in color along their lengths, suggesting that the segregation coefficient of Cr^{3+} in this material is nearly unity.

The absorption and emission spectra for $\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}:\text{Cr}$ are shown in Fig. 6-88. The emission is entirely from the $^4\text{T}_2$ excited state to the $^6\text{A}_1$ ground state as indicated by its bandlike character. The fluorescence exhibited a single exponential decay with a time constant of 330 μs .

Laser emission from a $\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}:\text{Cr}$ (4 at.% Cr) crystal was obtained using the experimental setup shown in Fig. 6-89. Both resonator mirrors had high reflectivity (>99%) in the region from 0.76 μm to beyond 0.85 μm . The 0.51-cm-thick laser crystal was cut and polished at LLNL. Surface reflection losses were minimized by placing the crystal between quartz windows, which

Fig. 6-88. Absorption and emission spectra of Cr^{3+} in $\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}$ at room temperature.

were AR coated on the outside surface only. Index-matching fluid was held by capillary action between the crystal and the inner surface of the windows. A flashlamp-pumped dye laser (Candela LFDL-6) operating with a 1- μ s pulse width at a wavelength of 0.64 μ m was used to produce a population inversion in the crystal. It was difficult to accurately determine the threshold for laser action because crystal laser action took place almost immediately after the dye laser reached its oscillation threshold. We estimated that the absorbed energy flux at threshold was about 0.5 J/cm². The lasing wavelength was 0.776 μ m, which is very close to the peak of the fluorescence spectrum (see Fig. 6-88). To our knowledge, this is the first time laser action has been obtained from this crystal.

It appears that this material could be a useful, tunable laser in the near infrared.

More work is needed on the thermomechanical properties of the material, and further tests are required to see if it can be flashlamp pumped.

KZnF₃:Cr. The cubic perovskite KZnF₃, doped with Cr³⁺, has been found to lase in the 0.82- μ m region by Brauch and Dürr.⁸⁰ This is an interesting result because most Cr³⁺-based lasers are derived from hosts that have a trivalent ion site, such as sapphire or alexandrite. To better understand the optical properties of this material, we have had a series of KZnF₃ crystals, doped with various amounts of chromium, grown by H. Newkirk at LLNL using the horizontal Bridgman method. H. Guggenheim has also provided one crystal. Two of the LLNL-grown crystals are shown in Fig. 6-90.

Emission and absorption spectra taken at room temperature are in general agreement with Brauch and Dürr. The fluorescence

Fig. 6-89. Experimental arrangement for production of laser action in a Na₃Ga₂Li₃F₁₂:Cr fluoride garnet crystal.

Fig. 6-90. KZnF₃ crystals grown at LLNL. Upper boule was grown from a melt with 1-wt% CrF₃ doping level. Lower, undoped boule lies in a graphite boat used in the crystal-growth process.

decay was single exponential, with a lifetime of 176 μs , in contrast to Brauch and Dürr's value of about 85 μs . It is possible that the discrepancy lies in the preparation of the material. Their crystals were grown from purified fluorides, as received from the manufacturer. Our fluorides are hydrofluorinated to ensure that oxide contaminants are removed. Dürr believes, based on ESR data, that an appreciable amount of charge compensation in his crystals is due to oxygen ions.⁸⁰ This should result in optical properties different than the intrinsic charge-compensation K^+ vacancies.⁸¹ The charge compensation is apparently local enough to reduce the site symmetry at the Cr^{3+} site, which, in turn, enhances the stimulated-emission cross section. It also gives rise to a segregation coefficient of about 0.1, as determined by x-ray fluorescence measurements. This leads to a visible gradient of Cr^{3+} concentration in the crystal (as is apparent in Fig. 6-90). This result also appears to limit the maximum amount of Cr^{3+} in KZnF_3 to less than 1 at.%, otherwise the strain is high enough that numerous cracks occur. Further attention to the growth may solve this problem.

Laser action from a $\text{KZnF}_3\text{:Cr}$ crystal was produced with an experimental arrangement almost identical to that shown in Fig. 6-87. The only change was that a thicker region of index-matching fluid (~ 1 cm) was used between the windows and the crystal. The $\text{KZnF}_3\text{:Cr}$ crystal used in the laser experiment was grown by H. Newkirk and W. Wade at LLNL by the horizontal zone method in a graphite boat. The melt from which it was grown had

0.5 wt% Cr. A clear, 4.1-cm-long section was cut and polished at the LLNL optics shop. The pump energy absorbed by the crystal at threshold was estimated to be 5 J/cm² at the center of the pump beam. Lasing occurred primarily at a wavelength of 0.799 μm , with a weaker sideband at 0.8165 μm . This laser experiment was the first to be performed at LLNL using an in-house grown, cut, and polished crystal.

Improvements in the performance of $\text{KZnF}_3\text{:Cr}$ as a laser crystal could be achieved by using higher Cr^{3+} doping levels. The strong segregation of Cr^{3+} out of the crystal might be reduced through charge compensation by another dopant, such as Li^+ .

An attempt to co-dope $\text{KZnF}_3\text{:Cr}$ with Nd failed due to the very low segregation coefficient for Nd. In addition, no crystalline material was obtained in attempted growth of $\text{RbCaF}_3\text{:Cr}$. However, preliminary spectra show that $\text{RbCdF}_3\text{:Cr}$ is a broadband emitter with short fluorescence lifetime (≈ 170 μs). Further examination of this material is in progress.

$\text{SrGaF}_5\text{:Cr}$. Strontium gallium pentafluoride has a tetragonal crystal structure, with two types of octahedrally coordinated Ga^{3+} sites.⁸² Several crystals, grown with 2 at.% chromium in the melt, were obtained from H. Guggenheim. The optical quality of the Bridgman-grown crystal was poor, with many cracks and bubbles. A Czochralski-grown crystal appears to be of higher optical quality.

The absorption and emission spectra of $\text{SrGaF}_5\text{:Cr}$ are shown in Fig. 6-91. Compared to $\text{Na}_3\text{Ga}_2\text{Li}_3\text{F}_{12}\text{:Cr}$ (see Fig. 6-88), the $^4\text{T}_2$ level is at lower energy in both absorption and emission. The fluorescence decay could be resolved into the sum of two exponentials, with lifetimes of 9.5 and 41 μs . The observation of a dual lifetime is consistent with occupation of both Ga^{3+} sites by Cr^{3+} ions. These lifetimes are short compared to Cr^{3+} in other fluoride media. At this time, it is not known whether the observed lifetimes are radiative. Experiments are planned to answer this question. If the lifetimes are radiative, calculations indicate that the emission cross section will be high (4 to 6×10^{-20} cm²) relative to those for other Cr-doped systems and in the region where efficient extraction of stored energy can be achieved in a short (i.e., 10-ns) pulse.

Fig. 6-91. Absorption and emission spectra of Cr^{3+} in SrGaF_5 at room temperature.

LaMgAl₁₁O₁₉:Cr. Lanthanum magnesium hexaluminate (LMAO), doped with Nd³⁺ ions, is being considered as a replacement for YAG:Nd in commercial systems by the French⁸³ and Soviets.⁸⁴ Due to the various positions of the Mg²⁺ ion relative to the La site, inhomogeneities broadening of the Nd³⁺ transitions occurs. Such broadening reduces the effect of concentration quenching of the Nd³⁺ fluorescence, permits higher doping levels, and, in turn, enhances the pumping efficiency. The thermomechanical and optical properties of LMAO appear to lie in the range that would make it useful as a fusion driver.⁸⁵

The LMAO lattice also contains octahedral Al³⁺ sites, which can accept transition metal ions such as Cr³⁺. Preliminary results on a Cr-doped LMAO crystal grown at LLNL indicate that the absorption is similar to that of alexandrite, with the ²E level clearly at lower energy than the ⁴T₂ level.⁸⁶ Fluorescence measurements indicate that the ⁴T₂ level shares population in thermal equilibrium with the ²E level, since both line and band emission are observed (as in alexandrite). The fluorescence decay is the sum of at least two exponentials, and the slower decaying component has a lifetime of about 3 ms. Thus, Cr³⁺ will probably not sensitize Nd³⁺ fluorescence efficiently in this host crystal.⁸⁷

It is possible that the gallium analog of LMAO may exhibit better Cr³⁺ photonic properties because the crystal field should be lower. Other sensitization schemes for LMAO are possible, and deserve further study because divalent and trivalent ion sites are available.

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Crystal Growth

As a part of our continuing effort to rapidly grow a wide variety of crystals for evaluation, several additional crystal-growth stations were constructed and placed into operation during 1984. This equipment is located in Building 241, Chemistry and Materials Science Department, together with the hydrofluorination systems already operating. The stations permit the preparative

chemistry of high-purity starting materials and the single crystal growth of oxide and fluoride compounds to 2000°C.

A resistance-heated furnace is on line and has the capability to control crystal diameter using the crucible weighing technique. The furnace has a crucible capacity permitting the growth of boules to a diameter of 45 mm. The furnace may be operated as a Czochralski crystal-growth station or as a controlled-atmosphere heat-treating furnace. Processing control by a programmable microprocessor can be interfaced to a personal computer for more centralized control. Another independently operating Czochralski growth station can control crystal diameter by weighing the boule. Processing control by a microprocessor is interfaced to a personal computer. It may also be operated in the vertical Bridgman growth mode for crystals up to a diameter of 25 mm.

When needed, rf power can be transferred to an adjacent station and used to operate a horizontal zone melter or refiner. A resistance-heated zone melter and refiner, using two narrow zone furnaces in tandem, provides two refinements in a single pass. Replacing the two furnaces with a longer single zone furnace provides the capability for horizontal freezing operations. This unit will be replaced by a new unit that will double the crystal-production capacity. These stations will be operated as zone melters in horizontal or vertical positions and as Czochralski pullers with controlled atmospheres. A flux growth station using a resistance-heated furnace with three independently controlled zones is operated by programmable microprocessors. The furnace is mounted in a vertical position on a

Table 6-20. Crystal boules grown at LLNL in 1984.

From the melt	RbCdF ₆ : 0.5% Cr ³⁺
KZnF ₃ : 1% Cr ³⁺	RbCaF ₃ : 0.5% Cr ³⁺
KZnF ₃ : 5% Cr ³⁺	LaMgAl ₁₁ O ₁₉ : 0.5% Cr ³⁺
KZnF ₃ : 0.5% Cr ³⁺	
KZnF ₃ : 0.2% Cr ³⁺	
KZnF ₃ undoped	
From the flux	
Rb ₂ NaYF ₆ + 0.1% Yb ³⁺	Rb ₂ NaYF ₆ + 0.1% Nd ³⁺
Rb ₂ NaGdF ₆ + 1% Yb ³⁺	Rb ₂ NaGdF ₆ + 1% Nd ³⁺
Rb ₂ NaGdF ₆ + 5% Cr ³⁺	Rb ₂ NaYF ₆ + 5% Yb ³⁺
Rb ₂ NaYF ₆ + 1% Nd ³⁺	Rb ₂ NaYF ₆ + 5% Cr ³⁺
Rb ₂ NaGdF ₆ + 5% Yb ³⁺	Rb ₂ NaGdF ₆ + 5% Nd ³⁺
Rb ₂ NaYF ₆ + 1% Yb ³⁺	Rb ₂ NaYF ₆ + 5% Nd ³⁺
Rb ₂ NaNdF ₆	Rb ₂ NaGdF ₆ + 5% Nd ³⁺
Rb ₂ NaGd _{0.75} NdF ₆ Cr _{0.25}	Rb ₂ NaGdF ₆ + 5% Nd ³⁺
Ca ₂ Nd ₂ La _{0.75} NdF ₆ Cr _{0.25}	Ca ₂ Nd ₂ La _{0.75} NdF ₆ Cr _{0.25}
KMgF ₃ + 1% Fe ²⁺	KFeF ₃
LiAlF ₃ + 1% V ²⁺	LiAlF ₃ + 1% Eu ²⁺

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gimbal stand to permit rotation of the furnace, while hot, to drain the contents of the crucibles. A metallic envelope surrounding the crucibles permits operation in controlled atmospheres to 1200°C.

Using this equipment, several large single-crystal boules were grown for measurement of mechanical properties, spectroscopic evaluation, and laser testing. The crystal boules are listed in Table 6-20.

Author: H. W. Newkirk

Multimegajoule Laser Project

Introduction

The objective of the multimegajoule laser project is to provide a low-cost, single-shot, 5- to 10-MJ, short-wavelength laser driver to demonstrate high-gain inertial-confinement fusion. This project is still in the conceptual stage and has led to the invention of a variety of new laser architectures, materials, and components. In the following three articles, we present an analysis of the basic differences between multipass systems and our traditional preamplifier/amplifier laser systems. We also discuss our cost objectives

for a simplified multimegajoule laser design and present the simplest, lowest-cost laser concept that we have yet analyzed.

Author: J. F. Holzrichter

Development Objectives

To meet our objective of constructing a 5- to 10-MJ, short-wavelength laser for \$200 to 300 million (overnight construction cost), we must reduce the recurring costs of the laser hardware by at least tenfold, compared to those costs for the Nova laser project. In addition, we plan to use manpower (design and assembly), diagnostics, and target-system costs that are similar to those on Nova. Table 6-21 shows the cost breakdown of the Nova project and a proposed breakdown for a conceptual multipass system that we invented in 1984 and describe below. To meet the cost objectives shown in Table 6-21, we have instituted research programs to reduce costs in all areas of technical activity. We have made considerable progress in these areas.

Our technical approach in accomplishing our objectives is to reduce the cost of individual parts and to increase the effectiveness (on a J/\$ basis) in using these parts. Tables 6-22 and 6-23 summarize the techniques that we have developed to meet the cost objectives shown in Table 6-21. In Table 6-21, the techniques that we are using to reduce the cost of the amplifier are reviewed. In any of the compact laser systems, amplifier cost dominates the optical-system cost and accounts for up to two-thirds of the system cost. The amplifier is assembled from materials in separate cost centers: laser glass, pulse power, and mechanical hardware. Table 6-22 shows how we are working to reduce the costs in each

Table 6-21. Cost breakdown for conceptual multipass system vs. Nova.

Budget category	Megajoule laser (\$ million)	Nova (\$ million)
Laser chain hardware	165	45
Labor		
Design	20	20
Assembly	30	40
Oscillators, diagnostics, and control	15	14
Target system	20	9
Tooling	10	12
Subtotal	260	140
Buildings	—	36
Total cost	260	176

Table 6-22. Techniques to reduce amplifier cost.

Cost center	Technique to reach 10 x cost reduction	Expected improvement factor	Progress during 1984
Laser medium	Improved extraction	~2 x	2 x
	Reduced glass cost	~8 x	~4 x
Pulse power	Improved amplifier storage efficiency	~4 x	4 x
	Reduced pulse-power cost	~4 x	~3 x
Mechanical hardware	Architecture	~2 x	~2 x
	Surface/volume	~3 x	—

of these amplifier centers to meet our objective of one-tenth the Nova cost. Significant cost reduction, for example, can be achieved by better using the laser glass. This is possible because improved extraction occurs when the amplifier medium (glass) is operated at a higher level of saturation (higher fluence). We plan to extract twice the energy stored in the laser glass of the amplifier, compared to that extracted in Nova.

In the previous articles on laser technology (see "Laser Glass," earlier in this section), considerable progress in reduced materials costs was described. Lower-cost glass, based on continuous production technology, should reduce the cost from the Nova average of $\$4/\text{cm}^3$ to a cost as low as $\$0.50/\text{cm}^3$. Improved amplifier efficiency, from 1 to 4%, reduces pulse power costs from $\$0.40$ to $\$0.10/\text{J}$, and large, compact (log-pile) mechanical structures reduce the cost for the flashlamp amplifier structures.⁸⁸

Table 6-23 shows the cost benefits that can accrue by using the optical materials in the laser chain at a higher fluence and by taking advantage of simpler system architectures. Such architectures lead to a reduction in the number of optical elements in the laser system. For example, we have made significant progress in increasing the threshold at which AR-coated optics damage, from $5 \text{ J} \cdot \text{cm}^{-2}$ (1-ns Nova operation) to $30 \text{ J} \cdot \text{cm}^{-2}$ (20-ns megajoule operation). We are pursuing several similar approaches to enhancing HR surfaces. One of the advantages of operating the laser amplifier systems at $1 \mu\text{m}$ is that higher damage thresholds appear to be attainable at that wavelength than at longer or shorter wavelengths. Subsequent harmonic conversion to nw reduces the fluence (by the efficiency of

conversion) but places an increased strain on the focusing optics, which have lower damage thresholds, depending on whether the final wavelength is 2ω , 3ω , or 4ω .

In summary, our objective is to obtain successful performance of a 5- to 10-MJ laser-target system by using a simplified architecture, by reducing the recurring cost of the laser components, and by more effectively using the laser materials.

Author: J. F. Holzhrichter

New Compact Multipass Laser Design

Figure 6-92 shows a simple laser design that has the fewest laser components of all fusion systems that we have studied and that packs closely, thus minimizing space requirements. The laser operation begins with the reflection of a spherical wave from a sacrificial mirror target located $\sim 1 \text{ cm}$ from the fusion target. The sacrificial mirror reflects light into a large, solid angle that includes many beamlines, such as the one shown. Approximately $1 \mu\text{J}$ of energy enters the laser chain through the partial reflecting output mirror. Energy is amplified in the backward direction and relayed down to a

Table 6-23. Techniques to reduce optical train costs.

Cost-reduction technique	Improvement factor	Progress during 1984
Increase average optical fluence $1 \mu\text{m}$, $5 \text{ J}/\text{cm}^2$ — $20 \text{ J}/\text{cm}^2$	} $4 \times$	} $(20 \text{ J}/\text{cm}^2)$ $(4 \text{ J}/\text{cm}^2)$
$0.25 \mu\text{m}$, $1 \text{ J}/\text{cm}^2$ — $4 \text{ J}/\text{cm}^2$		
Reduce number of surfaces	$\sim 2 \times$	$2 \times$
Reduce cost/ cm^2 of mirrors, windows, and harmonic crystals	~ 2 to $4 \times$	2 to $4 \times$
Total improvement	16 to $30 \times$	$10 \times$

Fig. 6-92. Multipass laser chain injected from a target reflector, harmonically converted, and focused onto a fusion target.

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Fig. 6-93. Typical amplifier element and estimated dimensions. Each unit could produce 1.2 kJ of short-wavelength light.

1- to 5-cm square beam, where it continues to propagate to the preamplifier-phase-conjugator system. Upon reaching the phase conjugator (or retroreflector), the beam phase (direction) is reversed. The beam then propagates out through the laser chain, with phase precorrected to both point the beam exactly at the sacrificial target and to precorrect the optical phase aberration in the amplifier medium. As the beam reaches the harmonic converter, it attains enough energy and power density (~ 17 kJ total at several GW/cm^2) to

Fig. 6-94. Thirty-six amplifier units are stacked into a log-pile configuration for efficient use of space. Flashlamps between each amplifier element pump several elements at once. All amplifiers are injected simultaneously by the same seed pulse.

Fig. 6-95. Twelve packages of each 36-unit amplifier structure can be arranged to occupy a space similar to that of the present Nova laser.

efficiently convert to a shorter wavelength. The harmonically converted beam passes through the dichroic output mirror and is deflected by the wedge, causing it to be focused onto the fusion target.

Each laser beam, operating under these optimal conditions, can produce ~ 12 kJ of short-wavelength light. To provide 5 or 10 MJ of short-wavelength light (see Figs. 6-93 through 6-96), 450 to 900 total beams (or 12 to 24 clusters of 36 beams each) are required. The identical nature of each beamline enhances the opportunities for cost reduction and permits straightforward testing of one such architecture.

Summary. The Advanced Laser Program objectives are determined by the requirements of the subsystems. The requirements consist of the following elements: high damage thresholds on reflectors; AR layers and dichroic coatings; high-efficiency amplifiers; low-cost production of laser glass, pulse power, and optical elements; and special optical elements, such as an effective phase conjugator and isolator. The combination of a compact architecture and lower-cost, higher-performance components can lead to significant reduction in overall system cost.

Author: J. F. Holzhrichter

Comparison of Architectures

Introduction. Within the next decade, we must construct a new laser system if ICF experiments are to be extended to higher energies. An economically attractive laser system will require improvements in both performance and reductions in unit costs by one order of magnitude compared with present systems. Significant improvements in the unit costs of construction could be achieved by changing the optical configuration or architecture of the laser system. We have explored several configurations in detail during 1984.

Currently, we envision a departure from the linear-chain approach for laser systems to a design that uses multiple passes through a large laser amplifier. From our experience gained during the construction of Shiva and Nova, we found that

preamplifier components constitute an inordinately large fraction of the hardware cost of a laser system, compared to the large, final amplifier stages. In addition, high extraction efficiency can be achieved at relatively low fluences in multipass systems.

If we select a multipass system over the linear chain, we have concluded that a change in architecture alone will not provide the magnitude of cost reduction necessary to make a new laser system more economically attractive. This does not mean that a change of architecture would be disadvantageous. Rather, the differences in performance and cost of various architectures are relatively small compared to the magnitude of improvement in cost effectiveness that will be required for future systems. To reach our cost goals, we will have to focus on reducing the cost of individual laser components by improvements in manufacturing processes and by designing systems that take maximum advantage of economy of scale.

Types of Architectures. We use the term architecture to refer to the specific arrangement of optical components that form a functional system. Examples of five configurations are shown in Fig. 6-97. Each architecture can be conveniently categorized by the number of passes the beam makes through an element of the amplifier medium.

Figure 6-97(a) is the familiar linear-chain system or MOPA. The beam passes only once through a given amplifier. Efficient use of the gain medium requires that the intensity of the beam be increased either by increasing the size of the beam or by subdividing it into multiple beams. In

Fig. 6-96. Opposing beams from each array of twelve 36-unit amplifier packages are arrayed around the target. The amount of space is similar to the present Nova and Shiva buildings.

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Fig. 6-97. Representative selection of architectures for laser systems.

(a) Linear-chain, single-pass system

(b) Cassegrainian, double-pass system with single-pass preamplifier

(c) Regenerative, multipass amplifier with polarization switch

(d) Regenerative, multipass amplifier with frequency-conversion switch

(e) Regenerative amplifier with noncollinear frequency-conversion switch

conventional MOPA construction, individual amplifiers are matched to the size of the beam; however, in principle, the beam could be folded appropriately to pass through different parts of a single amplifier.

Multiple passes through the same volume of active medium require some means for separating input and output beams. Angular separation can be used if the field of view of the amplifier is large enough. In the system shown in Fig. 6-97(b), a Cassegrainian telescope is used to expand the beam that passes twice through the outer annulus of the amplifier. The center of the amplifier can be used as a convenient preamplifier.

In Fig. 6-97(c), rotation of the angle of polarization is used to separate the input and output beams. Two passes through the amplifier can be obtained by use of a passive rotator. More passes require an active switch that must open and close within some fraction of the transit time through the cavity.

Figure 6-97(d) shows how frequency conversion can be used as a switch. Such switches use the dependence on nonlinear intensity or the narrow angular acceptance angle of the frequency-conversion process to distinguish between multiple passes.

More esoteric configurations using frequency conversion can be conceived. The architecture shown in Fig. 6-97(e), is based on the fact that a frequency-converted beam need not be collinear with the input drive beams but requires the presence of both beams to convert efficiently.

Each of the five architectures has advantages and disadvantages, and each one differs with respect to performance and cost. Nevertheless, our study has shown that the similarities in these architectures far outweigh the differences. Most of the components are the same, and each architecture requires large final amplifiers. Each system requires spatial filters to limit longitudinal ASE, to control the growth of spatial noise, and to relay the beam image. In addition, each system has elements for frequency conversion and focusing on target, and each one needs a driver of some sort. The primary differences among systems arise from the number of required preamplifiers and the presence or absence of a switch.

Transverse Scaling. Transverse scaling and associated issues are insensitive to

the choice of architecture. An obvious exception to this general rule is the case of an architecture-dependent component, such as a switch. At the energies of interest, a complete system must consist of many beams organized into assemblies at a scale set by structural and economic considerations. The fundamental size of an individual beam is that of an ASE-limited disk in the amplifier, and all large amplifiers for any architecture must consist of multiple disks arranged in longitudinal and transverse arrays of various dimensions. Spatial filters can be constructed with focusing elements matched either to the aperture of a single laser disk or to a large number of apertures.

Frequency-conversion crystals are matched to some unit number of beams. Our experience with the Nova amplifiers and crystal arrays suggests that assemblies of multiple beams into larger structural units will result in substantial cost savings. At present, however, we have insufficient data to establish the scale of the minimum-cost beamline. However, the modularization will be similar for any of the architectures we might choose.

Our ultimate choice of an architecture, therefore, will be based on the performance and cost of a single beam. Because our study shows that the largest contrast is between a linear MOPA [Fig. 6-97(a)] and a regenerative multipass amplifier [Fig. 6-97(e)], we have selected these two configurations to illustrate specific issues.

Extraction Efficiency. One potential difference between multipass and single-pass systems arises from the fraction of energy stored in the amplifier that can be extracted. The extraction efficiency reaches a maximum at a fluence determined by the gain, the loss, and the number of passes. However, maximum extraction efficiency for an idealized amplifier is independent of the number of passes through the amplifier. Figure 6-98 shows extraction efficiencies for an amplifier with a gain of 20 and fixed losses per pass as a function of output fluence for 1, 2, and 4 passes. Each curve reaches the same maximum extraction efficiency, although the fluence required to reach maximum extraction efficiency for a multipass amplifier decreases as the number of passes increases. Most of the fluence difference in a multipass amplifier is

Fig. 6-98. Extraction efficiency vs output fluence for single, double, and multiple passes through the amplifier.

obtained between 1 and 2 passes through the amplifier.

An important question is whether the difference in fluence required to obtain a specific extraction efficiency represents an advantage for one approach over another. The answer depends on several factors. If the fluence is limited by damage to the laser medium in the amplifier, then a multipass system can extract up to 30% more energy from the amplifier than a single-pass system, as indicated in Fig. 6-98. If damage is not an issue and extraction efficiency is limited only by internal losses, then the single-pass system can reach the same extraction efficiency at an output fluence that is higher by a factor of 2.

In practice, the differences between systems are smaller than this simple analysis suggests. Nonlinear effects and pulse distortion increase as the point of maximum extraction efficiency is approached. A design using some fraction of the maximum extraction efficiency (perhaps 90%) is more desirable, and the difference in fluence in this case would only be about 30%.

Laser Glass. Extraction efficiency and output fluence also depend on the saturation fluence of the laser glass. In principle, the extraction efficiencies of multipass and single-pass amplifiers can be made identical at the same output fluence by choosing two different laser glasses with appropriate saturation fluences. A comparison of the performance of systems with different

glasses reveals a significant difference in the extraction characteristics of multipass and single-pass amplifiers. For a given amount of energy stored in an amplifier, a multipass system extracts a fraction of this energy at a fluence that is independent of the saturation fluence. In contrast, a single-pass amplifier requires a fluence that increases with saturation fluence to obtain the same extraction efficiency. If the output fluence is limited by some mechanism, such as optical damage, and if for some other reason, such as material cost, larger apertures, or improved excitation efficiency, it were appropriate to choose a high-saturation-fluence glass, multipass systems would be clearly advantageous.

A closer examination reveals that other factors reduce this potential advantage. For an amplifier with constant stored energy, gain varies inversely with saturation fluence. The internal losses in the laser medium tend to be relatively independent of saturation fluence. Figure 6-99 shows the extraction efficiency as a function of saturation fluence for an amplifier with constant stored energy and losses. The upper curve is the maximum local extraction efficiency set by the losses. The extraction efficiency decreases monotonically with increasing saturation fluence of the glass. If the fluence is limited by damage, extraction efficiency is reduced, as shown in the lower curve. The damage fluence is assumed to be equal to the fluence required for maximum extraction efficiency in a glass with a saturation fluence of 65 cm^2 . These two cases correspond closely to multipass (upper curve) and single-pass (lower curve) systems. The difference between the two cases, which could be attributed to a difference in architecture, is relatively small. The larger variation is associated with the general difficulty of using high-saturation-fluence glasses without a corresponding reduction in the losses. In fact, there is some evidence of improved storage efficiency in some high-saturation-fluence glasses. Development of such glass, with improved characteristics sufficient to overcome the degradation due to losses, will influence the choice of architecture and render multipass systems more attractive.

Pulse Distortion. Efficient extraction of energy stored in an amplifier requires operation in the saturated regime. As

fluence increases, gain decreases. This relationship can produce severe distortion of the amplified pulse. A significant difference in pulse distortion of the linear-chain system and a multipass amplifier arises under saturating conditions.

A convenient measure of pulse response is the ratio of final to initial power P_f/P_i , required at the input to maintain constant intensity (a square pulse) at the output.

Figure 6-100 shows a comparison of pulse distortion, defined in this manner, as a function of output fluence for a single pass and for multiple passes through the same amplifier. As the output fluence increases, the square-pulse distortion increases monotonically. At the output fluence giving maximum extraction efficiency, the pulse distortion for a single-pass or for a multipass amplifier is on the order of the single-pass gain. However, optimal designs of a linear chain consist of several stages of saturated amplification to obtain the same overall gain as the multipass system, and the pulse distortions of individual stages are multiplied. Thus, the square-pulse distortion for the linear chain is significantly higher than that for the multipass system.

As noted previously, the relative output fluences are also different. If the output of the linear chain were reduced to the same value as that for the multipass system, the square-pulse distortion would be reduced significantly. Whether this difference in the input pulse shape is significant depends on the penalties exacted by the controls required to generate the input pulse. If a contrast ratio of 1000:1 is technically viable, then there is no reason to downrate the linear chain because of its higher pulse distortion. If the distortion must be reduced in the linear chain, the output fluence could be reduced or the saturation in the preamplifiers could be reduced by increasing their size and number. However, both of these solutions tend to increase the relative cost of the linear chain and may give multipass systems a significant advantage.

Optical Damage. Optical damage is the most important and perhaps the most uncertain design issue for any new laser system. At present, damage fluences for bare glass surfaces are expected to be approximately 60 J/cm^2 at 10 ns. If we assume a reduction by a factor of 2 to account for intensity structure on the beam, then the

Fig. 6-99. Peak extraction efficiency vs saturation fluence of a laser medium for constant stored energy and losses.

design fluence is about 30 J/cm^2 . Operation at this fluence with available glasses will require relatively high-gain amplifiers for which the differences in performance between multipass and linear-chain systems are small if optical damage is not an issue.

Unfortunately, the damage fluences of other components in the beam, such as the frequency-conversion crystals, mirrors, and polarizers, tend to be lower by a factor of 2 than those for bare or AR-coated surfaces. Assuming that no progress is made in solving this problem, two design approaches are possible. First, the beam can be expanded to match the damage fluence of the vulnerable components. This is the Nova design, but it makes close packing difficult. Second, the laser can be designed and operated at lower fluences, so that a multipass system would have the advantages described above.

Fig. 6-100. Square-pulse distortion vs output energy for single and multiple passes through the same amplifier.

Cost. The differences in performance and architecture are ultimately reduced to the issue of cost. Many of the cost issues are related to components and materials that are common to all architectures. To separate the issues associated with architecture, and to make the comparison of costs more explicit, consider again a linear-chain and a multipass system [see Figs. 6-97(a) and 6-97(e)] designed using the same type and number of amplifiers in the final stage. Each system stores 25 kJ. Both systems require spatial filters, frequency-conversion crystals, and optics to form a complete system. The linear chain requires a preamplifier, while the multipass system uses a polarization switch.

Table 6-24 shows a comparison of the costs of the various subsystems for the two architectures. The costs used for the comparison of subsystems are based on our procurement experience with Nova hardware. Hardware costs represent about one-third of the total Nova construction costs; other cost centers are labor and buildings, including support facilities. These costs do not include any projections for improvements in performance or cost and, therefore, overestimate the ultimate cost of any new laser system. Our experience shows that the smaller preamplifiers tend to be more expensive than the final amplifiers, and this difference is reflected in the cost structure. The costs of optics and optical subsystems are assumed to scale with

surface area. Specialized components and optical assemblies cost more than conventional components. The polarization switch is assumed to cost the same as a Faraday isolator plus an additional polarizer. This is probably an underestimate of the actual cost of a fast switch.

Table 6-24 also compares the costs of a linear-chain and a multipass-amplifier system constructed with components of the same aperture. The amplifiers make the largest single contribution to cost, but the contributions of other components to the overall cost of the system are significant. Most of the major components are the same for both systems, and the relative costs reflect this fact. The primary differences between the two systems are in the cost of the preamplifiers and isolators for the linear chain and the cost of the switch for the multipass system.

The multipass system costs approximately 25% less than the linear chain. This advantage in cost must be weighed by the relative output of the two systems because the same total output is desired in each case. There are two outputs, depending upon whether or not the system is fluence limited. This comparison is shown in Table 6-25. As noted in the discussion of Fig. 6-98, the linear chain is capable of producing approximately 30% higher output from the same amplifier if it is not fluence limited. The net cost effectiveness would be approximately the same for the two systems. However, if the fluence is fixed by damage, then the multipass system is the more cost effective.

In summary, there are differences in the cost effectiveness of various architectures. However, changes in architecture alone will not provide sufficient improvements in cost effectiveness for the next system. Multipass systems are being investigated because they have the potential for cost savings, particularly when operating in the regime of saturated pulse extraction. The development of a low-cost, high-saturation-fluence glass would require a multipass system for effective operation. Multipass systems may have significant advantages when overall system size and components are included. Evaluation of the magnitude of these savings is, however, limited by uncertainty in the real performance characteristics of

Table 6-24. Costs of laser components.

Sub-system	Cost for comparison of subsystems	MOPA (\$ thousand)	Multipass system (\$ thousand)
Amplifier	\$40 per J stored	1000	1000
Preamplifiers	\$60 per J stored	400	—
Spatial filters	400 ^a	250	150
Mirrors	100 ^a	—	40
Isolators	300 ^a	200	20
Switch	400 ^a	—	200
Frequency conversion	500 ^a per color	100	100
Focusing	1000 ^a	100	100
Total		2040	1610

^aIn \$ thousand/m².

Table 6-25. Relative cost effectiveness of multipass and linear-chain systems.

	MOPA (fluence limit)	MOPA (no fluence limit)	Multipass system
Relative cost	1.00 to 1.25	~1.25	1.0
Relative output	1.0	~1.3	1.0
Relative cost per unit output	1.00 to 1.25	~0.95	1.0

multipass systems. The magnitude of the potential cost savings is sufficient to warrant an experimental program to resolve such uncertainties.

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Other Laser Systems

Introduction

We continued small-scale studies of innovative approaches to high-efficiency, short-wavelength, low-cost lasers during 1984. Our studies covered several potential lasers that are based on complexes of Group II metals and rare gases. We developed a novel technique for generating a high concentration of such metal atoms in a cool background gas. In addition, we funded studies of electronic energy-transfer processes that might be useful in electronic chemical lasers.

Flashlamp Evaporation of Metal Dusts

We have investigated a novel method of producing metal vapors of Group IIa and IIb elements.^{1,2} These metals have several characteristics that make them attractive as possible storage media in laser systems. They have s² ground-state configurations and therefore do not form strongly bound dimers; they have low-lying metastable states with excitation energies of 2 to 5 eV, and they have relatively high vapor pressures.

We predict that this method is capable of yielding high, uniform concentrations of metal atoms in a cool, but reasonably high-pressure, background gas. Our proposed method would use an optical technique for evaporation of metal dusts, which has the advantage that the energy is used efficiently to produce vapor atoms. The medium remains relatively cool because photons are absorbed only by the metal particles and not

by the buffer gas. We find that the light intensity produced by standard flashlamps is adequate to evaporate Group IIa and IIb metals if the particle diameter is on the order of a few microns. Metal dusts with particles of this size apparently are readily available and are ideal for dispersion in a gas medium.

The residence time of these particles in the vapor is determined by their downward velocity. This velocity in (cm/s) is given by

$$v = \frac{g\rho d^2}{18\mu} \quad (84)$$

where g is the acceleration of gravity, ρ is the metal density, d is the particle diameter, and μ is the gas viscosity. The viscosity of xenon is 227 μ P at room temperature and is independent of pressure. Assuming that the maximum tolerable downward velocity is 0.1 cm/s, then the maximum particle diameters for the metals of interest are $d_{Mg} = 4.9$, $d_{Sr} = 4.0$, $d_{Ba} = 2.4$, $d_{Ca} = 2.2$, and $d_{Hg} = 1.8$ μ m. Our calculations show that by using dusts with particles of these sizes, high-concentration metal vapors can be produced by flashlamps.

A theory of the evaporation process has been developed over the past few years to investigate the laser evaporation of various aerosols. We have used the theory of Belyaev et al.,^{3,4} which assumes the evaporation rate is diffusion limited. The thermal-balance equation is

$$\frac{dT_m}{dt} = \frac{3}{C_p \rho_s a} \times \left[K_{ab} - L_p \frac{da}{dt} - (\epsilon \sigma T_m^4 - T_c^4) - I_1 \right] \quad (85)$$

where

- T_m = temperature of the metal,
- C_p = heat capacity of the metal,
- ρ_s = density of the metal,
- a = radius of the metal sphere (time dependent),
- I_1 = incident intensity of flashlamp radiation,
- K_{ab} = absorption coefficient of the metal,
- L = heat of vaporization of the metal,
- ϵ = emissivity of the metal.

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σ = Stefan-Boltzmann constant,
 T_x = temperature of the buffer gas, and

$$j_l = -\lambda \left. \frac{\partial T}{\partial r} \right|_{r=a} \quad (86)$$

is the thermal flux to the gas, with λ the thermal conductivity of the buffer.

The Knudsen equation, which may be expressed as

$$\rho_s \frac{da}{dt} = v \rho_s (T_s) - \alpha v \rho_s^l \quad (87)$$

with

α = sticking coefficient (unity for these metals),

$\rho_s(T_s)$ = equilibrium vapor density of the metal at T_s ,

ρ^l = actual density of metal vapor at $r = a$, and

$v = (kT_s/2\pi m)^{1/2}$ is the rate per unit area at which a Maxwell distribution of atoms strikes a surface.

Assuming that the evaporation rate is actually controlled by the rate of diffusion of

the metal atoms away from the surface boundary, we determine the mass flux from

$$j_m = (\rho_s^l - \rho_s) \frac{da}{dt} = -D \left. \frac{\partial \rho_s}{\partial r} \right|_{r=a} \quad (88)$$

where D is the diffusion coefficient.

To illustrate the complete solution of these coupled equations for the evaporation process, we present detailed results for magnesium in Fig. 6-101. The time evolution is shown for some of the significant parameters for magnesium particles in 1 atm of xenon, with initial radii of $5 \mu\text{m}$, and with an absorbed power (IK_{ab}) assumed to be 8 kW/cm^2 .

Figures 6-101(a) and 6-101(b) show the evolution of the radius and the temperature of the sphere. (In all cases the metal particles are rapidly heated above the melting point so that it is valid to treat them as spheres.) During the initial phase ($\sim 100 \mu\text{s}$), the spheres are rapidly heated until diffusion-limited evaporation can occur. After the heating period, evaporation proceeds under these conditions for $\sim 0.5 \text{ ms}$ until the radius becomes small

Fig. 6-101 Time-dependent evaporation of $5\text{-}\mu\text{m}$ Mg particles in 1-atm xenon buffer for an absorption rate of 8 kW/cm^2 .

enough that conduction of heat to the buffer becomes competitive.

Figure 6-101(c) shows q_{lamp} , the rate of energy absorbed (in W) by the particle, Fig. 6-101(d) shows q_{cond} , the energy conducted to the buffer, and Fig. 6-101(e) shows the fraction of energy being absorbed that is lost due to conduction. From the total energy absorbed, given by the integral of q_{lamp} , and the total conductive loss, given by the integral of q_{cond} , we find that the conductive loss is only 6% of the total energy absorbed. An additional contribution to heating the background gas is due to the heat carried by the metal atoms. Figure 6-101(e) shows the sudden onset of the dominance of the conductive loss after about 0.8 ms. At this time however, the radius has been reduced enough that the energy being absorbed is negligibly small, and the volume vaporized is a most 100%. The final radius is $0.08 \mu\text{m}$, and the fraction of the atoms left in the particle is $\sim 4 \times 10^{-6}$ of the initial mass.

From such calculations, we have shown that high densities (1 to 10 Torr) of metal atoms can be produced in a background gas, which remains approximately at room temperature by the flashlamp vaporization of metal dust particles. The evaporation can be reasonably efficient, particularly if absorption of radiation by the walls is not large. The energy required is on the order of a few eV per vapor atom, and we predict that a quiescent and homogeneous vapor will be established on a time scale comparable to the flashlamp pulse duration (~ 1 ms). The residual particles will remain small enough that scattering of laser light and recondensation of the vapor will be insignificant.

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Resonant Charge-Exchange Excimer System

We continue to explore new media that have the potential for converting electrical energy to light more efficiently than current pump sources. One system that has been examined previously is a class of heteronuclear ions with emission bands corresponding to

radiative charge-transfer transitions of the type

The transition energy corresponds roughly to the difference in the ionization potentials of X and Y. For example, the spectra of the heteronuclear rare-gas ions were observed by Tanaka et al.⁹¹ More recently, the structure and kinetics of these systems under e-beam excitation were studied theoretically and experimentally by Szoke, Winter, and Powell.⁹²

A problem with these systems is that the rare gas with the higher ionization potential has the lower stopping power. Hence, conditions that favor population inversion correspond to large overpressures of the lighter rare gases. Under such conditions, formation of the homonuclear diatomic ion via three-body recombination

is a major loss mechanism.

We have carried out a theoretical study of a rare-gas alkaline-earth ionic system $(\text{MgXe})^+$. The kinetics of such a system may be significantly influenced by the fact that the charge-exchange reaction

is nearly resonant ($\Delta E = 0.05$ eV) and produces excited magnesium ions in the $3p^2P^o$ state. As expected, our calculations showed that there is a XeMg^+ potential-energy curve dissociating to $\text{Xe} + \text{Mg}^+(3p)$, which is below by more than 1 eV. Three-body collisions might well be expected to produce $\text{Mg}^+(3p)\text{Xe}$ on a time scale that competes with three-body formation of Xe_2^+ if the Mg density can be made large enough. We would then expect to observe fluorescence in a red-shifted $3p \rightarrow 3s$ band of Mg^+ from

At large internuclear separations, the dominant interaction between an ion and a closed-shell atom is charge-induced polarization, which we expect to play a major

Other Laser Systems

role in the molecular binding of $(\text{MgXe})^+$. Therefore, we attempted to ensure that the atomic basis sets used in our calculations were flexible enough to describe the atomic dipole polarizabilities. In the case of Xe, we first derived a relativistic core potential to replace the inner-shell electrons of atomic Xe, which contribute negligibly to the polarizability of the atom and are not important for chemical bonding. In this approach,³³ an effective local potential, based on a relativistic Hartree-Fock calculation for the atom, is defined which, by construction, reproduces the energy and outer shape of the valence orbitals. Because the inner electrons are not treated explicitly, the use of these potentials can greatly reduce the time spent in calculating two-electron integrals.

We used the configuration interaction (CI) method to generate the potential-energy curves for $(\text{MgXe})^+$ in LS-coupling. In choosing the list of configurations to be included in our calculations, we tried to meet two limited objectives. First, we wanted to equally describe all states correlating with $\text{Xe}(^5\text{S}) + \text{Mg}^*(nl)$, where subshell $nl = 3s, 3p, 4s, 3d$, or $4p$. This description is important because excited-state absorption from the proposed laser upper state $[\text{XeMg}^+(3p)]$ to higher lying excited states is an important issue. Second, we wanted to account for the interaction of the single Mg^+ electron with the $5s$ and $5p$ electrons in Xe to obtain reasonably accurate dissociation energies for those XeMg^+ states that are bound.

Figure 6-102 shows the potential-energy curves of the XeMg^+ and MgXe^+ states in LS-coupling. The potential-energy curves of all low-lying states correlating with $\text{Xe} + \text{Mg}^+$ are attractive, whereas those correlating with $\text{Mg} + \text{Xe}^+$ are repulsive for $R \leq 7 a_0$, where a_0 is the Bohr radius. For simplicity, the repulsive curves dissociating to $\text{Mg}(3s3p, {}^1\text{P}) + \text{Xe}^+$ are not shown. The potential-energy curves of the "charge-transfer" states (those of MgXe^+) are not realistic because we made no attempt to calculate accurately the ionization potential of Xe or the spin-orbit splitting of the ${}^3\text{P}_{1,2}$ and ${}^3\text{P}_0$ levels of Xe^+ . However, we believe that the calculated potential-energy curves dissociating to $\text{Xe} + \text{Mg}^+$ are reasonable.

The equilibrium internuclear separations of the $X^2\Sigma^+$, $A^1\Pi$, and $C^2\Sigma^+$ states are estimated at 5.63, 5.15, and 5.01 a_0 , respectively, and the calculated dissociation energies are 0.52, 1.16, and 1.93 eV, respectively. The $B^2\Sigma^+$ state appears to have a shallow minimum near 6.5 a_0 . If the curves are shifted to reproduce the experimental splittings of the atomic states, then near the equilibrium geometry of the A state the vertical electronic transition energy is 3.70 eV for the A-X transition and 3.57 eV for the A-C transition. Near $R = 6.5 a_0$ the vertical transition energy of the B-X transition is 4.60 eV.

The calculated electric-dipole transition moments for the X-A, X-B, and A-C transitions are large and approximately independent of internuclear distance because they correspond to the atomic 3s-3p and 3p-4s transitions, respectively, of Mg⁺. The transition moment for the X-C transition

Fig. 6-102. Calculated potential-energy curves of XeMg^+ and MgXe^+ .

decreases to zero at large internuclear distances because the atomic 3s-4s transition is forbidden.

Our study illustrates two potential difficulties with $(\text{XeMg})^+$ as a laser candidate. First, the strong A-X transition is a "bound-bound" transition. Thus, rapid vibrational relaxation of the $X^2\Sigma^+$ state is required to prevent bottlenecking. Second, the photons emitted by the A-X transition can be absorbed by the equally strong A-C transition. However, the latter is also a bound-bound transition, and, therefore, a coincidence between the vibrational bands of the A-X and A-C transitions is by no means a certainty.

It may be worthwhile to consider a potential laser based on the strongly allowed B-X transition with an energy of ~ 4.6 eV. Our calculations predict a shallow minimum in the potential-energy curve of the $B^2\Sigma^+$ state near $6.5 a_0$. In this case, excited-state absorption is still a potential loss mechanism because the calculations predict that two $^2\Sigma^+$ states, correlating with $\text{Xe} + \text{Mg}^+(3d)$ and $\text{Xe} + \text{Mg}^+(4p)$, respectively, lie approximately 4.4 eV above the B-state near $6.5 a_0$ (see Fig. 6-102). In addition, the transition moment connecting the B-state to the $\text{XeMg}^+(3d)$ state is especially large ($\sim 2.0 a_0$).

It is interesting to speculate whether we could observe the emission due to the charge-transfer transition $\text{MgXe}^+ \rightarrow X^2\Sigma^+ \text{XeMg}^+$ in the 4.4-eV region near the resonance line of Mg^+ . Based on our calculations, the $\Omega = 1/2$ state dissociating to $\text{Mg} + \text{Xe}^+(^2P_{1/2})$ has a reasonably large transition moment (~ 0.7 to $1.0 a_0$) to the ground state at internuclear distances between 6 and $8 a_0$.

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Group II-Rare-Gas Excimer Systems

Group II metal atoms all have closed-shell (n^2s^2) electronic configurations. The ground electronic states of Group II-rare-gas systems are therefore van der Waals molecules, as are those of interacting rare-gas atoms, with predominantly repulsive interactions. The low-lying excited states of these systems, which correlate with the n^1P and n^3P excited metal atoms, are only weakly

bound. Thus, the excimer bands would be on the wings of the metal-vapor lines. They would be characterized by low gain coefficients per excited atom, homogeneous broadening across the excimer band, and the necessity for high rare-gas densities to ensure adequate concentrations of transient excimer molecules.

Group II-rare-gas systems have several characteristics that make them more attractive as potential fusion laser candidates than the more thoroughly studied alkali-rare-gas excimers.⁹² Group II metal atoms, unlike the alkalis, do not readily form stable dimers. Absorption of excimer radiation by metal dimers, which has proven to be debilitating in the case of the alkalis, may not be an issue in the case of the Group II atoms. Moreover, the lowest excited states of the Group II atoms are metastable (n^3P), and, in the absence of any significant excited-state absorptions, they may be viable storage media.

The n^3P_1 wavelengths and radiative lifetimes are listed in Table 6-26. The strongly allowed $n^1P \rightarrow n^3S$ atomic-line radiation will be trapped, but the far-line wings will escape. Excited-state absorption should be negligible for these allowed transitions, but rapid, high-power pumping would be necessary. For some systems, the collision-induced $(n-1)^1D \rightarrow n^3S$ transitions may also be worth considering.

With the exception of mercury, elevated temperatures are necessary to obtain adequate equilibrium metal-vapor pressures for reasonable gain coefficients. Indeed, the difficulties associated with operating high-pressure systems at high temperatures is undoubtedly responsible for the paucity of experimental data presently available. However, novel techniques for flashlamp vaporization of metal dust currently being investigated (see "Flashlamp Evaporation of Metal Dusts," earlier in this section) point to the possibility of obtaining high metal-atom

Table 6-26. Group II n^3P_1 wavelengths and radiative lifetimes.

Transition	Wavelength, λ (nm)	Radiative lifetime, τ (μ s)
Mg ($3s^2 \rightarrow 3s3p$ 3P_1)	457	232 ^a
Ca ($4s^2 \rightarrow 4s4p$ 3P_1)	657	385
Zn ($4s^2 \rightarrow 4s4p$ 3P_1)	309	20
Sr ($5s^2 \rightarrow 5s5p$ 3P_1)	689	16
Cd ($5s^2 \rightarrow 5s5p$ 3P_1)	326	1.7
Ba ($6s^2 \rightarrow 6s6p$ 3P_1)	791	3.2
Hg ($6s^2 \rightarrow 6s6p$ 3P_1)	254	0.11

densities at low operating temperatures and have given us added incentive to undertake further study of these systems.

During 1984 we focused on the metal-rare-gas excimers that correlate with metastable (n^3P_1) metal atoms. It is important to realize that the excimer states that derive from the n^3P_1 levels are not expected to have lifetimes that depend strongly on internuclear distance or rare-gas partner. These "spin-forbidden" transitions cannot be collision induced in the usual sense. They derive their intensity from spin-orbit mixings with the n^1P_1 excimer levels, the lifetimes of which vary only slightly from the asymptotic atomic values. Excited-state absorption may be a problem in some of these systems. The most important excited-state absorption for the metastable levels

Fig. 6-103. MgXe potential-energy curves.

Fig. 6-104. SrXe potential-energy curves.

will be from bands associated with the $n^3P \rightarrow (n+1)^3S$ atomic transition in the case of the Group IIa atoms (Mg, Ca, Sr, and Ba) and the $n^3P \rightarrow n^3D$ transitions in the case of the Group IIb atoms (Zn, Cd, and Hg).

Absorption of some Group II-rare-gas mixtures has been measured and used to deduce information about the potential-energy curves.⁹⁵ These measurements are of no help in predicting excimer bands because they relate only to long-range interactions. We therefore carried out *ab initio* calculations of the low-lying electronic states of MgXe and SrXe. In the cases of Sr and Xe, we used effective relativistic core potentials⁹⁶ to replace the inner-shell electrons that contribute negligibly to the polarizability of the atoms and that are unimportant for chemical bonding. The use of these potentials greatly reduces the computer time for calculating two-electron integrals. A list of reference configurations was chosen consisting of those states dissociating to ground-state Xe plus a metal atom in either the ground 1S or excited 3P and 1P states. We then carried out CI calculations, including a restricted set of single and double excitations. We defined these excitations by the valence metal and rare-gas orbitals and the list of reference configurations (referred to as "first-order" or POL CI). This prescription resulted in ~ 2600 to 3000 configurations for each symmetry considered.

Figures 6-103 and 6-104 show the potential-energy curves of MgXe and SrXe in LS coupling. The curves have been shifted to reproduce the correct splittings of the asymptotic atomic states. These calculations show that the excited states arising from the $^1S + ^3P$ and $^1S + ^1P$ asymptotes behave similarly to the states arising from a 3P alkali atom colliding with a rare-gas atom. The $^3\Pi$ and $^1\Pi$ states are weakly attractive, with well depths of ~ 700 and ~ 1600 cm^{-1} , respectively. The $^3\Sigma$ and $^1\Sigma$ curves arising from the $^1^3P + ^1S$ asymptotes, on the other hand, are unbound and are dominated by the repulsive overlap between a directed $p\sigma$ metal orbital and a rare-gas atom.

We have also plotted a higher-lying $^3\Sigma$ curve, which dissociates to a rare-gas atom and a Group II atom in the 3S state. This state has a large transition moment to the lowest $^3\Pi$ curve and can be a

potentially damaging source of excited-state absorption of photons emitted from the $^3\Pi \rightarrow ^1\Sigma$ transition. Our calculated curves show that for MgXe, this state is energetically accessible to $^3\Pi$ excimer photons. For SrXe, however, $^3\Sigma$ lies too high to be an excited-state absorber. SrXe seems to be a promising candidate.

Spin-orbit coupling is relatively weak for Mg and Sr. Such coupling has the effect of inducing only small fine-structure splittings in the calculated LS potential curves. However, spin-orbit coupling will have a significant effect on the radiative properties of the 3P excimers.^{97,98} Consider the isolated Group II atom. The spin-orbit interaction can be characterized by the two parameters λ and λ' . The parameter λ refers to coupling within the 3P manifold that splits the level into three nondegenerate $^3P_{0,1,2}$ states. The parameter λ' refers to coupling between the 3P_1 and 1P_1 states. Based on the magnitude of these spin-orbit parameters, the 3P_1 state will acquire some admixture of 1P_1 and thus attain a nonzero radiative rate to the ground state. By using the known 1,3P -level spacings, we can determine the parameters λ and λ' and the coefficient of 1P_1 in the 3P_1 wave function.

To calculate molecular spin-orbit curves, we use an approximate treatment. The atomic spin-orbit parameters are first used to determine the matrix elements of the spin-orbit Hamiltonian at infinite separation. These matrix elements are assumed to be independent of internuclear distance. The total molecular Hamiltonian is then $H(R) = H_{el}(R) + H_{so}$, where H_{el} is the electronic Hamiltonian, and H_{so} is the spin-orbit Hamiltonian. The functions that diagonalize $H_{el}(R)$ are simply the calculated LS configuration-interaction wave functions. The eigenvalues of $H(R)$ at the calculated R values are the spin-orbit (Hund's case-c) curves.

Figure 6-105 shows a schematic of the $^1,3P + ^1S$ potential-energy curves with and without spin-orbit coupling. The $^3\Sigma$ and $^3\Pi$ states are split into a family of spin-orbit states that correlate with the three 3P_j asymptotes. The 0^+ and 1 states dissociating to 3P_1 are optically connected to the ground state. The 0^+ state is essentially $^3\Pi$ plus a small admixture of $^1\Pi$; the 1 state is $^3\Pi$ plus a small admixture of $^1\Sigma$. The 1 state dissociating to 3P_2 , which is $^3\Sigma$ plus a small admixture of $^1\Pi$, can also radiate. The 0^- and 2 states are metastable. Because the $^1\Pi$

Fig. 6-105. Schematic of potential-energy curves (a) without and (b) with spin-orbit coupling.

and $^3\Sigma$ transition moments have little R dependence, their admixture into the $^3\Pi$ and $^3\Sigma$ states varies only slightly with internuclear separation as well. Thus, our model predicts that $^3\Pi$ and $^3\Sigma$ excimer lifetimes are essentially equal to the atomic 3P_1 radiative lifetimes (2.3 ms for Mg and 16 μ s for Sr).

If the atomic 3P_1 excited states are statistically populated, then the metastable population residing in states that correlate with 3P_0 and 3P_2 is not directly accessible to laser action. In mercury, for example, where the spin-orbit splitting is large (~ 4000 cm^{-1}), there may be a reservoir of metastable states that is not readily deactivated. In the lighter atoms, however, where the spin-orbit splitting is smaller, we expect rapid cross relaxation at high rare-gas pressures to minimize this problem.

The weakly bound character of the Group II-rare-gas excimers gives rise to low gain coefficients and homogeneous broadening, with the possibility of high power levels and laser-pulse energies. Excimer states associated with the metastable 3P atomic levels may be viable candidates, provided that excited-state-absorption bands do not exist in the same wavelength region. Our calculations show that SrXe is particularly favorable in this regard. The CdXe system, which we have not yet studied, may also be interesting. Based on the atomic 3P_1 lifetime, CdXe should have a gain coefficient larger than that of SrXe by one order of magnitude, although $5s5p\ ^3P \rightarrow 5s5d\ ^3D$ excited-state absorption may be a problem. Detailed emission spectra of these systems should be measured. Theoretical calculations such as these can only provide a qualitative picture of what may be expected. It is particularly important to obtain more information about the radiative lifetimes of the triplet excimer states.

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Electronic Chemical Laser

During 1984 we renewed our interest in gas storage media and investigated several potential candidates that might lead to a large-scale fusion laser system. Much of our theoretical work is discussed elsewhere in this section, but, in addition, we funded two outside experimental efforts involving the

metastable $N_2(A^3\Sigma_u^+)$ state. This state is attractive because of its 2-s lifetime, its 6.2-eV excitation energy, and its stability with respect to quenching, except by atoms or molecules that can undergo a spin-changing excitation in the process. Therefore, the $N_2(A^3\Sigma_u^+)$ state is similar to the $O_2(a^1\Delta_g)$ state, which has been used successfully in the oxygen-iodine excitation-transfer laser, except that the $N_2(A^3\Sigma_u^+)$ state has considerably more energy to yield.

The first experimental study involved searching for new methods to produce the $N_2(A)$ state, either through photodissociation or chemical reaction. The second study involved measuring the excitation-transfer rates to metastable levels of selected molecules with shorter radiative lifetimes (shorter lifetimes make such molecules better candidates for laser media).

Generation of $N_2(A)$. Two methods for generating $N_2(A)$ states were tested by C. Wittig and his group at the University of Southern California. These researchers investigated the photolysis of NCNO for a range of photon energies from 0.266 to 0.900 μm , hoping to induce the following process:

Although this photolysis reaction is highly exothermic, and there appeared to be a reasonable likelihood of producing $N_2(A)$, the only products observed were NO and CN (Ref. 79). It appears that the substantial nuclear rearrangement necessary to obtain the desired products does not occur quickly enough to compete with the more direct dissociation.

Next, a study of the possible reaction

was undertaken. This reaction to ground-state products is exothermic by 6.44 eV so that there is sufficient energy to produce the $N_2(A)$ state. Because it might be necessary to overcome a reaction barrier to obtain the excited products, high translational energies were considered desirable. To obtain this translational energy, BrCN was dissociated by 0.193- μm photons. This approach produced vibrationally cold, but translationally and rotationally hot, CN fragments. Rotational energies were as high as $10\,800\ \text{cm}^{-1}$. The rapidly rotating

molecules are not expected to be very reactive and require thousands of collisions to cool. Unfortunately, after many collisions, the translational temperature is also thermalized, and no reaction was observed. Therefore the results were inconclusive, and a new source of CN fragments must be considered to pursue this reaction.

Reactions of $N_2(A)$ with SO and S_2 . D. Setser and his group at Kansas State University studied the excitation transfer from the $N_2(A)$ state to SO and S_2 , where

and

The SO(A) states have lifetimes on the order of 10 to 35 μ s and emit photons with wavelengths in the range from 0.240 to 0.270 μ m (Ref. 100). These parameters make this molecule an attractive candidate for a storage laser medium. Cao and Setser¹⁰¹ found that the total quenching-rate constant of $N_2(A)$ by SO is $(1.2 \pm 0.2) \times 10^{10}$ cm³/s. The rate constant for producing the SO(A) state [see Eq. (95)] is $(0.5 \pm 0.2) \times 10^{10}$ cm³/s. The $S_2(B)$ state has a lifetime of 30 to 45 ns (Ref. 102) and emits between 0.280 and 0.440 μ m (Refs. 101 to 103). The total quenching-rate constant of $N_2(A)$ by $S_2(X)$ was $\sim 6 \times 10^{11}$ cm³/s, with a branching fraction of 0.3 for formation of the B-state.

These rate constants were measured in a flowing afterglow reactor. The $N_2(A)$ state was generated via transfer from Ar($^3P_{0,2}$) atoms, which are created by flowing purified Ar through a low-power hollow-cathode discharge. The SO molecules were produced by microwave discharges in dilute SO₂/Ar and S₂ from S₂Cl₂/Ar mixtures. The SO₂ and S₂Cl₂ gases were added to the flow reactor from radially distributed nozzles. The quenching rates were measured by directly observing the rate of decay of the $N_2(A)$ concentration. The SO(A $\rightarrow$ X) emission from the reactor was measured under high resolution, and the relative vibrational populations of the lowest seven levels of the excited molecule were determined. The $S_2(B\rightarrow X)$ emission was also

observed and found to be one order of magnitude weaker than the SO(A $\rightarrow$ X) emission under the same conditions. The vibrational distribution of $S_2(B)$ was also measured.

These experiments show that the excitation-transfer reactions in Eqs. (95) and (96) are relatively efficient, with branching fractions for SO(A) and $S_2(B)$ in the range of 0.5 and 0.3, respectively. The larger total quenching-rate constant for SO(X) makes it the more interesting candidate for efficient energy conversion from $N_2(A)$ and therefore for an excitation-transfer ultraviolet laser.

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Atomic and Molecular Theory

Introduction

Accurate theoretical models of atomic and molecular energy levels and interactions are important in many areas of the ICF Laser Program. During 1984, we extended our understanding of collision-induced absorption and emission, reactive atom-diatom collisions, photoionization, *ab initio* calculations of energy surfaces of molecules containing heavy atoms, and coherent excitation.

Far-Wing Absorption in Na-Ar Collision

Collision-induced absorption and emission at wavelengths well removed from line center play important roles in many atomic and molecular processes. We have developed the theory and computer codes to calculate exact quantum mechanical cross sections for these optical and radiative collisions between atoms.¹⁰⁴ We also have produced a quasi-classical model that can efficiently generate accurate absorption cross sections. This model cannot, however, give branching ratios for the final-state populations. Our codes and model can be used to study the propagation of nearly resonant light through gaseous media and to calculate accurate gain and absorption cross sections for the far wings of atomic transitions.

We have used our theory to study the collision-induced absorption by sodium in argon for wavelengths in the vicinity of the resonance lines D_1 and D_2 . Figure 6-106 shows the potential-energy curves for the ground and first three excited electronic states of the Na-Ar system. Red-wing absorption excites primarily the $A^2\Pi_{3/2}$ and $A^2\Pi_{1/2}$ states, and blue-wing absorption involves the $B^2\Sigma_{1/2}$ molecular state. Following absorption, the collision pair dissociates to produce excited sodium atoms, which subsequently fluoresce. During the dissociation, fine-structure transitions between the

Fig. 6-106. Electronic potential-energy terms for the Na-Ar system.

Fig. 6-107. Spectrum of the partial excitation cross sections per collision pair vs detuning from the $2P_{1/2}$ state.

excited molecular states occur, which determine the final populations of the $2P_{1/2}$ and $2P_{3/2}$ states of sodium.

The collision-induced absorption is actually a three-body process involving two atoms and a photon. Therefore, we define η_p , the partial excitation "cross section" per collision pair, by

$$\frac{d[2P_{1/2}]}{dt} = \eta_p(\Delta\omega) c n_\phi [Na][Ar] \quad (97)$$

where $\Delta\omega$ is the detuning of the laser wavelength from line center, c is the speed of light, n_ϕ is the number density of photons, and $[Na]$ and $[Ar]$ denotes the concentrations of sodium and argon, respectively. Similarly, we can define the usual optical collision cross section for a given laser intensity (photon density) that is related to η by

$$\sigma_l(\Delta\omega) = \frac{\eta_p(\Delta\omega) c n_\phi}{v} \quad (98)$$

where v is the relative velocity of the collision pair.

Figure 6-107 shows two partial cross sections for a large range of both red and blue detunings. The scale on the right-hand side of Fig. 6-107 gives the optical collision cross section for a laser intensity of 1 MW/cm^2 .

From the quantum mechanical results, we also can predict the polarization of the fluorescence of the excited atoms.

Our quasi-classical theory can be shown to be equivalent to the usual quasi-static theory of line broadening for spectra outside the impact region (i.e., the far-wing region). Classical approaches do not include an accurate treatment of the fine-structure transitions between the excited molecular states, so that we cannot obtain partial cross sections. For our system, however, we obtained excellent agreement between the quasi-classical and quantum mechanical total absorption cross section $\Sigma_i \sigma_i(\Delta\omega)$. Therefore, for cases in which only the total cross sections are of interest, the much more efficient quasi-classical method is adequate.

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Resonances in Atom-Diatom Reactions

A model collinear, symmetric reactive scattering system ($A-B-A$), which exhibits a highly oscillatory reaction probability as a function of collision energy, has received considerable recent attention.^{105,106} Resonances in molecular systems are important because they can change reaction cross sections by orders of magnitude over small energy ranges. Resonances can be studied via their absorption spectra, which can become highly complicated. For example, the recently reported¹⁰⁷ infrared predissociation spectrum of H_3^+ contained approximately 27 000 lines across only 222 cm^{-1} . Many of the lines represent transitions between two states, both of which contained energies above the dissociation threshold (i.e., molecular resonances). The existence of long-lived vibrational resonances with at least 5 eV of internal energy in this molecular ion is indeed surprising. To understand such complex spectra, we must develop methods for characterizing the resonant states of this and other molecules.

The model system we have studied obeys the Hamiltonian

$$H = \frac{1}{2\mu} p_1^2 + \frac{1}{2\mu} p_2^2 + \frac{1}{m_B} p_1 p_2 + V(r_1) + V(r_2) \quad (99)$$

where V is a Morse potential, r_1 and r_2 are bond distances, p_1 and p_2 the corresponding momenta, and the reduced mass μ is given by $m_A m_B / (m_A + m_B)$. From scattering calculations, we determined that for $m_A = 35/31$ and $m_B = 210/31$ the reaction probability is very small, except in the vicinity of narrow resonances, where the probability increases to unity.

The usual method for determining resonance wave functions is to diagonalize the total Hamiltonian in a basis set made up of products of square functions with finite norms. Some of the eigenvalues will correspond to resonances. If a resonance is well represented (i.e., overlaps strongly) with a single product function, it is possible to assign quantum numbers to this state.¹⁰⁸

In performing these calculations, the choice of coordinate system strongly influences the results obtained. We cannot use "normal" coordinates that are typically employed for lower vibrational levels because they cease to be well defined for the configurations accessible in these higher energy states. Therefore, we tested two different coordinate systems in our preliminary studies. We used "local" coordinates or bond distances r_1 and r_2 . We also used hyperspherical coordinates r and ϕ defined by

$$r_1 = r \sin(\phi_m - \phi) \quad (100)$$

and

$$r_2 = r \sin \phi \quad (101)$$

where

$$\phi_m = \arctan \left[\frac{m_B(m_A + m_B + m_A)}{m_A^2} \right]^{1/2} \quad (102)$$

We constructed product states of Morse functions for the local-mode representation and vibrational adiabatic hyperspherical product states for the hyperspherical modes. We then performed a variational calculation by diagonalizing the molecular Hamiltonian in each basis set. Both representations resulted in eigenfunctions with energies closely approximating the resonance energies we had obtained from full scattering calculations.

Neither of our two coordinate systems was optimal for compactly representing all resonant states. Some states were clearly local-mode resonances, and others were hyperspherical. Some resonances were a sum of several product states in either basis.

In effect, our calculations show that unraveling a complex spectrum, such as that of H_2^+ , will be a difficult problem. We expect that no zero-order, separable Hamiltonian can be found, that no quantum numbers can be assigned to the resonant states, and that no simple selection rules exist for transition strengths between the states. Even for our relatively simple model, we have found that different resonances can have very different characters.

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Atomic Photoionization by Complex Basis-Function Expansion

Gaseous photoionization processes can play a critical role in determining the efficiency of short-wavelength pump sources and the ultimate gain possible in many gas lasers. Although significant improvements in the measurement of ground-state photoionization cross sections have been made in the last decade, measuring absolute cross sections for excited-state photoionization remains extremely difficult.

We have developed a practical theoretical technique for computing photoionization cross sections based on the method of complex basis functions. The method has proven to be a practical tool for characterizing resonance (autoionizing) states of many-electron atoms and molecules.¹⁰⁹ More recently, we have extended this technique to study nonresonant scattering as well.^{110,111} The technique of complex scaling provides a formal method for "rotating" the continuous spectrum of the Hamiltonian into the complex energy plane. Rendering the continuum discrete via an expansion in a finite set of L^2 functions can then provide a convenient method for calculating, at real energy values, matrix elements of the resolvent from which the photoionization cross section can be obtained. Complex scaling moves the singularities of the resolvent to a

set of isolated points in the complex energy plane and, hence, renders them innocuous.

The photoionization cross section for an atom in state ψ_0 with energy E_0 is related to the negative frequency component of the dipole polarizability

$$\alpha^-(\omega) = \lim_{\epsilon \rightarrow 0} \left\langle \psi_0 \left| \mu \frac{1}{H - E_0 - \omega - i\epsilon} \mu \right| \psi_0 \right\rangle \quad (103)$$

through the relation¹¹²

$$\sigma(\omega) = \frac{4\pi\omega}{c} \text{Im} [\alpha^-(\omega)] \quad (104)$$

The method of complex basis functions involves making a projection of H onto a finite set of (complex) orthonormal basis states $|\chi_i\rangle$ and yields an approximation to the photoionization cross section in the form

$$\sigma(\omega) \approx \frac{4\pi\omega}{c} \text{Im} \sum_i \left(\frac{\langle \psi_0 | \mu | \chi_i \rangle}{E_i - E_0 - \omega} \right)^2 \quad (105)$$

where E_i and χ_i are the complex eigenvalues and eigenfunctions of the matrix H , and μ is the dipole operator. The square, not the modulus square, of the dipole matrix element appears in Eq. (105) (Ref. 111).

We have carried out the first practical applications of the method of complex basis functions to photoionization in a many-electron atom.¹¹³ We have computed cross sections, including the effects of target-state correlation and final-state channel coupling, for photoionization of magnesium in its ground and first excited 3P states. We have considered the two ionization processes

and

The initial bound states and final continuum states were represented by linear superpositions of many-electron configurations. The configurations were constructed from a set of orthonormal atomic orbitals, which were expanded as linear combinations of real and/or complex Slater orbitals.

In the ground-state calculations, we investigated the effects of including correlation in the initial state and channel coupling in the final state. The calculations of excited-state photoionization employed a self-consistent field (SCF) target state and channel coupling in the final states. Calculations were performed using both the length and velocity form of the dipole transition moment in Eq. (105).

The ground-state cross section of magnesium is shown in Fig. 6-108. One striking result is the very broad resonance feature in the cross section, which arises from the inclusion of the $(3p4s)$ configuration in the final-state wave functions. The doubly excited $(3p4s)^1P$ state is the first member of an infinite set of autoionizing states. We found the position and width of the $(3p4s)^1P$ resonance to be relatively insensitive to changes in the basis set parameters; this was not true of higher members of this series, which we chose to project out of the calculations. The cross section rises somewhat sharply and then falls off more slowly on the high-energy side of the resonance. This result is in qualitative accord with the far-ultraviolet absorption spectrum of magnesium measured by Balloffet and Esteve¹¹⁴ and the recent cross-section measurements of Preses et al.¹¹⁵

In the case of photoionization from the excited 3P state, optical selection rules allow for the construction of final-state wave functions of 3S , 3P , and 3D symmetry that are dipole connected to the initial state. We can label the final channels as

$$(3kp), ^3S + (3ks), ^3S \quad (108)$$

$$(3kp), ^3P \quad (109)$$

and

$$(3kp), ^3D + (3kd), ^3D \quad (110)$$

The $(3pnp), ^3S$ and $(3pnp), ^3D$ configurations will again lead to a series of autoionizing

Fig. 6-108. Photoionization cross section of ground-state Mg.

Fig. 6-109. Total photoionization cross section of Mg $(3s3p)^3P$.

states converging to the $(3p)^2P$ state of Mg^+ . The $(3p^2)^3P$ state of Mg lies below the ground state of Mg^+ and is not autoionizing. Our calculations included the first $(3p4p)^3S$ and 3D autoionizing states. The total cross section, which is dominated by the 3D component, is plotted in Fig. 6-109. To our knowledge, no other results are available for comparison.

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Relativistic Calculations on Polyatomic Systems

Ab initio calculations on polyatomic molecules and solids have progressed to the extent that complexes containing any atom in the periodic table can now be treated. For meaningful comparison with experiments, the relativistic terms in the Hamiltonian can no longer be ignored. These terms contribute to bond shortening, core contraction for s and p electrons, core expansion for d and f electrons, and energy shifts in the valence

shells due to spin-orbit coupling.^{116,117} These effects are currently included by replacing the core electrons on the heavy atoms with a relativistic effective potential (REP) and then carrying out either a Dirac-Fock (DF) calculation¹¹⁸⁻¹²⁰ or a nonrelativistic valence electron calculation,¹²¹ followed by a perturbation calculation of the spin-orbit splittings. The first (DF) approach is currently feasible only for diatomics. The nonrelativistic method may not be appropriate for heavy atoms with p, d, and f valence electrons, where the spin-orbit splitting is comparable to the electron repulsion integrals.

We have developed a new program to carry out configuration interaction (CI) calculations over spin orbitals that have been determined using an REP. The CI Hamiltonian contains an *ab initio* one-electron spin-orbit operator determined from the REP. Therefore, the CI calculation not only introduces electron correlation, but allows for orbital relaxation and energy-level shifts due to spin-orbit coupling.

The complete spin-orbit operator can be written as

$$V_{so} = \frac{1}{2} \alpha^2 \left[\sum_i \xi(r) \mathbf{l}_i \cdot \mathbf{s}_i - \sum_{i,j} (\mathbf{r}_i - \mathbf{r}_j) \times \mathbf{p}_i \frac{\mathbf{s}_i + 2\mathbf{s}_j}{r_{ij}} \right] \quad (111)$$

The two-electron terms for core-core electron pairs are nearly constant, and the terms for valence-valence pairs are small. The core-valence pairs act approximately as shielding terms. In our work, only the first term, $V_{so} = 1/2 \alpha^2 \sum_i \xi(r) \mathbf{l}_i \cdot \mathbf{s}_i$, is considered.

Following the notation of Ermiler et al.,¹²² the REP can be written as

$$U^{\text{REP}} = U_{lj}(r) + \sum_l \sum_j [U_l(r) - U_{lj}(r)] \times \sum_{m_l} |l m_l\rangle \langle l m_l| \quad (112)$$

where L and J are the smallest angular momentum quantum numbers not present in

the core. It is possible to rewrite the REP in the form¹²²

$$U^{\text{REP}} = U^{\text{AREP}} + U_{so} \quad (113)$$

where U^{AREP} is the angular momentum averaged REP given by

$$U^{\text{AREP}} = U_{lj}(r) + \sum_l [U_l(r) - U_{lj}(r)] \times \sum_{m_l} |l m_l\rangle \langle l m_l| \quad (114)$$

with

$$U_l = \frac{(l+1)U_{l,l+1/2} + lU_{l,l-1/2}}{2l+1} \quad (115)$$

The spin-orbit operator for use with pseudo-orbitals is simply the difference of the REP for $j = l + 1/2$ and $j = l - 1/2$, and

$$\Delta U_l = \frac{2(U_{l,l+1/2} - U_{l,l-1/2})}{2l+1} \quad (116)$$

is multiplied by the appropriate projection operator. By using the values of the matrix elements of $\mathbf{l} \cdot \mathbf{s}$ over the projection operators and evaluating the summations, where possible, the effective spin-orbit operator can be derived from

$$U_{so} = \mathbf{s} \cdot \sum_l \Delta U_l \times \sum_m \sum_{m'} |l m\rangle \langle l m| \mathbf{l} |l m'\rangle \langle l m'| \quad (117)$$

By adding the first term of Eq. (113) to the Hartree-Fock (HF) operator and solving for the one-electron orbitals, we obtain a set of HF orbitals that include all the relativistic effects in the DF equations, except for spin-orbit coupling. The second term of Eq. (113) is then added to the nonrelativistic Hamiltonian, which is diagonalized in the basis of HF spin orbitals to incorporate both electron correlation and spin-orbit interaction. In this approach, the spin-orbit correction is not added as a perturbation, and the spin-orbit matrix elements are not estimated

semiempirically or adjusted to reproduce atomic splittings. Additional advantages of this method are that the integrals can be computed and stored over orbitals rather than spin-orbitals, and, if sufficient symmetry is present, the CI coefficients can be identified in advance as either pure real or pure imaginary. This information can be used to make the Hamiltonian matrix pure real.¹²³ It is also possible to carry out the calculations using double-group-adapted functions.

During 1984, we have written programs to evaluate the integrals over the U^{AREP} and U_{so} operators and rewritten the CI program to diagonalize the Hamiltonian over spin orbitals. The CI program is now being tested on the complexes UF_6 , NpF_6 , and PuF_6 .

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Modeling Noise in Strong Laser-Atom Interactions

The response of an atom to near-resonant radiation changes dramatically as the radiation intensities increase. The weak-field rate processes of traditional low-intensity spectroscopy give way to Rabi-type population oscillations in strong fields (such as those that occur in laser-induced isotopically selective excitation). In turn, the Rabi oscillations influence atomic spectral response by modifying weak-field absorption and emission-line profiles to include sidebands.

The theory of such strong-field spectral features has been known for many years, under the simplifying assumption of an isolated atom that is excited by monochromatic radiation. In practice, the excitation is never monochromatic, and an atom is never truly free from influence by random encounters with neighboring atoms. A vast literature treats the effect of such a fluctuating environment on weak-field spectral lines, and many authors have reported theoretical extension into the strongly excited regime. In their simplest form, these theoretical approaches introduce effective relaxation times or bandwidths to model random fluctuations of the field or of the environment. The more elaborate theories

treat the details of fluctuations as stochastic Markovian processes.

Previous approaches to this subject fall into two broad categories. The first category uses the theory of Gaussian random process. This approach, which permits an exact solution in an interesting variety of cases, ultimately rests on the assumption that some variable represented by white noise underlies the interaction.

The second category is less well known and uses the theory of random-interruption (or jump) processes. This model was first applied to laser excitation by Burshtein and coworkers.¹²⁴ It assumes that excitation dynamics can be characterized by a set of parameter values (such as the values of instantaneous laser frequency and intensity). These values remain constant, except at random, discrete times, when one or more parameters may assume new values as a result of an infinitesimal interruption.

In collaboration with colleagues at the University of Rochester and the Polish Academy of Sciences, we examined examples of interruption-process effects on observable properties of atomic excitation. We published accounts of our work as a series of four papers.¹²⁵⁻¹²⁸

In Ref. 125, we described a practical scheme for computing time-dependent properties, such as excitation probabilities and dipole correlation functions, of a multi-level atom strongly excited by a laser in the presence of noise. The method, based on the notion of randomly interrupted and multivalued jump processes, was shown to be applicable to noise in laser phase, amplitude, and/or frequency and to noise in the microfield environment. The method was shown to provide analytic expressions, based on eigenvalues and eigenvectors of a finite-dimensional matrix, for time dependence and spectral dependence, and to permit computation of time-dependent fluorescence emitted from atoms excited by a noisy laser.

In Ref. 126, we applied a specialized version of the theory to examples of phase noise modeled as a superposition of random telegraph signals. We gave examples of laser spectra computed by direct calculation, and we derived a third-order differential equation for the average inversion of a

resonantly excited two-state atom. We discussed the general features of the interaction, such as the optical analog of motion narrowing. We showed that the theory of generalized Poisson processes allows laser-atom interactions, in the presence of random telegraph noise, to be treated systematically by means of a master equation first used in the context of quantum optics by Burshtein. We used this equation to obtain an exact expression for the two-state atom, steady-state resonance fluorescence spectrum when the exciting laser exhibits phase telegraph noise. We then compared our results with those from other noise models.

In Ref. 127, we extended the general Markov-chain approach to the theory of noisy laser-atom interactions. We considered the special case, a chain composed of n independent two-state jump processes, which is an example of a non-Gaussian stochastic process. We showed that the n -component Markov chain has the property of pre-Gaussian noise because, in the limit of $n \rightarrow \infty$, it converges to a Gaussian stochastic process. We gave solutions for the field spectrum and the atomic-fluorescence spectrum for phase, frequency, and amplitude pre-Gaussian fluctuations of the driving laser light.

In Ref. 128, we extended the preceding results of phase noise to the case of frequency noise. Such noise can be used either as a model of collisional effects, in which the atomic frequency randomly jumps, or as a model of finite laser bandwidth effects, in which the laser frequency exhibits random jumps. We discussed how these two types of frequency noise can be distinguished in light-scattering spectra. Finally, we discussed examples of both temporal and spectral motion narrowing, of nonexponential correlations, and of non-Lorentzian spectra.

Author: B. W. Shore

Gating of Population Flow in Resonant Multiphoton Excitation

The behavior of n -state atoms excited coherently by near-monochromatic radiation is known to be predictable from solutions to the time-dependent Schrödinger equation in the rotating-wave approximation. The application of Schrödinger dynamics to multistep excitation in isotope separation is well known. However, realistic modeling of an actual atom, including hyperfine structure, magnetic-sublevel degeneracy, Doppler shifts, inhomogeneities, and excitation pulse characteristics, often inject such complexity that it is difficult to interpret the results of either realistic modeling or experiments. For this reason, the results for elementary models (such as monochromatic illumination of stationary, nondegenerate multistate atoms) can be useful to serve as touchstones for interpreting more realistic models. These models of time-dependent dynamics, often admitting exact analytic solution, are the analogs of the hydrogen atom as an element of molecular structure. By examining and understanding simplified models, we gain insights that aid in the interpretation of complicated time dependence.

As a part of an ongoing effort, we examined a model in which excitation probability flows along a multistate chain of energy levels driven by coherent photoexcitation.¹²⁹ We examined the gated interruption of this probability flow caused by a strongly coupled side branch on the excitation chain. We found that, if the side branch is itself a chain rather than a single transition, then flow inhibition occurs if the number of branch transitions is odd. No inhibition occurs if the number is even. Such inhibition provides a limiting bound on the intensity of ASE.

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protects the reactor vessel and, by stopping the fusion neutrons, provides heat from which electricity can be produced. The large photo at left illustrates granular behavior studies with which some of the granular blanket parameters were defined. These studies, combined with other calculations, showed that the rotation speed and the shape of the chamber could maintain an evenly moving flow of granular material on the inside wall of the chamber, thus protecting it from damage and excessive wear. In the photo below, a rotating cone examined qualitative flow characteristics. In the inset photo left, an inclined chute qualitatively determined the velocity profile that would exist in Cascade. Top inset photo shows that portion of Cascade where stationary scoops carry heated granules to the heat exchanger.

Applications of Inertial Confinement Fusion

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Applications of Inertial Confinement Fusion

Introduction

W. J. Hogan

We have long recognized that there are many potential applications, both military and civilian, for inertial confinement fusion (ICF) technology. Specific proposals have included:

- Studying the physics of materials at high temperature and density for nuclear weapons design applications.
- Using the x rays and neutrons emitted by a fusion target to study nuclear weapons vulnerability, lethality, and effects.
- Producing nuclear materials (e.g., Pu, T₂, or ²³³U) for weapons or civilian power application.
- Producing electricity.
- Propelling space vehicles.

What is not generally recognized is that the ICF Program has already begun to realize some of these potential applications. While high fusion gains have not yet been obtained, the quest for high pellet compressions has already produced temperatures and pressures not attainable anywhere else on earth except in nuclear explosions. Thus, some applications are under way.

In this fast-changing field, it is natural that there has been some shift of focus in the ICF Applications Group. In the past, a great deal of emphasis was placed on first-generation power reactor technology. We required that all materials and technology used in the design (except those of the driver) be available today. Then if a driver were available, it would follow that a practical power plant could be built. The best example of a reactor designed within these precepts is the high-yield lithium-injection fusion energy reactor, HYLIFE, in which steel walls are protected by falling columns of liquid lithium that also breed tritium and transfer the heat to a secondary loop.

We are shifting our attention to include studies applicable to two other time periods:

- The near-term bridge between present research and applications that are realizable in the near future.
 - The ICF power plants of the far future that must be economically competitive with the cost of projected future fission and coal power plants.
- This year's articles describe progress on the Cascade and Pulse*Star power-reactor concepts, and include some general contributions to ICF-applications technology.

The greatest fraction of our effort this year was devoted to Cascade studies. We see this reactor concept as a most promising one for improving the safety and economy of ICF power relative to previous designs. In the spirit of looking beyond first-generation power plants, we have allowed ourselves to use more exotic materials and technologies than might be generally assumed available today. At the same time, we do not invoke any "magic" solutions to problems that would exist. We use only solutions that have been proposed and that appear to be sound, even though they may not be in commercial application today. In particular, we have switched to a chamber made of silicon carbide (SiC) tiles held together in compression by a web of tendons that are a composite of SiC fibers in an aluminum matrix.

To obtain the highest energy-conversion efficiency, we redesigned the cascading granular blanket and heat exchanger. Granular-flow experiments and analyses convinced us that stratified flow in a two-layer, two-material blanket is feasible. The new granular blanket has an inner layer of beryllium oxide (BeO) and an outer layer of lithium aluminate (LiAlO_2) with a mixed-mean outlet temperature of almost 1700 K, compared with 1200 K in the previous design. This should lead to higher conversion efficiency, although the design of the power system is not yet final. In addition, the granule transport method chosen uses no auxiliary energy: granules fly from the reaction chamber to the heat exchanger using the kinetic energy they obtain from the chamber rotation needed to keep them against the wall.

We have estimated the recondensation time in the chamber for the new blanket and for the lower pellet yield made possible by the high thermal-to-electric conversion efficiency. It appears that resultant chamber-clearing times will allow us to select pulse rates at least as high as the desired 5 Hz.

Work on Pulse*Star slowed this year to allow work to start in other areas. Progress was made, however, on two issues raised last year—tritium containment and polonium production. Progress reports are given on these issues, and a new version of Pulse*Star is presented that avoids both issues by substituting liquid lithium for the lithium-lead (LiPb) coolant.

Cascade and Pulse*Star differ from HYLIFE in several important respects. They operate at lower gains and yields and at higher pulse rates. The high temperature of Cascade and the pool geometry of Pulse*Star are both aimed at better power production economics. The use of ceramic granules, LiPb coolant, and SiC walls is aimed at increasing safety and minimizing environmental impact. These design trends are illustrated quantitatively in the next article.

We developed an economic model to help identify technological developments that could make power from ICF reactors fully competitive with that from other sources. We first used this model to look at heavy-ion fusion, and found that smaller plant sizes and lower pulse rates appear more attractive than previously thought.

This year, we present a new reactor concept, called Starfish. We have been considering how to increase the efficiency of converting fusion energy to electricity. In previous designs, the fusion energy is converted to heat, the heat is transferred to steam, and the usual 39% efficiency of the steam cycle is accepted. Turbines themselves, however, can be very efficient machines. Some hydroelectric turbines achieve efficiencies of more than 80%. In the Starfish reactor, fusion energy is first converted to fluid kinetic energy by placing a mass of lithium very close to the pellet. The fluid kinetic energy is converted to electricity directly in high-efficiency turbines, and the residual thermal energy is used in a conventional steam cycle. While we do not minimize the problems and objections that can (and will) be raised to this concept, we believe that the possible energy conversion efficiency of over 60% makes it worth examining.

As a first step in designing the bridge from present research to near-term applications, we have begun to examine reaction-chamber concepts useful for a facility in which a 5- to 10-MJ driver produces single fusion pulses of 100 to 1000 MJ. Our objective is to keep the cost of such a chamber to a minimum. Work will be needed to consider the design impacts of making the reactor a multipurpose facility.

In sum, the age of practical use of ICF technology has begun. We have responded by shifting our vision from first-generation power plants further toward the future of fusion power and toward the bridge between present

research and near-term applications. The work reported here reflects the continuation of past work and the first impacts of that shift.

Power Reactor Design Trends

Cascade and Pulse*Star represent new trends in ICF power reactor design that have emerged in the last few years. The most recent embodiments of these two concepts, and that of the HYLIFE design with which we will compare them, are shown in Figs. 7-1(a), 7-1(b), and 7-1(c). All three reactors depend upon protecting structural elements from neutrons, x rays and debris by injecting massive amounts of shielding material inside the reaction chamber. However, Cascade and Pulse*Star introduce new

ideas to improve the economics, safety, and environmental impact of ICF reactors. They also pose different development issues and thus represent technological alternatives to HYLIFE.

Economic Issues

The operating regimes of these three reactors are shown in Table 7-1. As can be seen, the design yields are lower and the pulse rates higher for the two new designs. Recent economic studies indicate that the cost of electricity from an ICF plant is likely to decrease as a function of pulse rate (at constant net power). The rate of decrease

Fig. 7-1. (a) Cascade's high operating temperature obtains a high thermal-to-electric conversion. (b) The Pulse*Star reactor's compact, pool-type geometry maintains a very high power density. (c) The high-yield lithium-injection fusion reactor, HYLIFE, a first-generation power reactor.

Table 7-1. Operating regimes of three LLNL power reactor concepts.

Power Reactor Design Trends

	HYLIFE	Pulse*Star	Cascade
Driver energy (MJ)	4.5	2	1.5
Target yield (MJ)	1800	600	300
Repetition rate (Hz)	1.5	4.3	5
Blanket outlet temperature (K)	770	770	1670
Gross thermal power (MW _t)	3170	3300	1670
Net electrical power (MW _e)	1010	1115	800

is high between 1 and 4 Hz but diminishes rapidly thereafter, so we have chosen pulse rates at the knee of this curve. Table 7-1 also shows that the driver energy and target gains are smaller for Cascade and Pulse*Star than for HYLIFE. These differences reflect both economic motives and a desire to be less dependent upon the achievement of very high gains. Finally, both new designs take additional measures to improve their economics. Pulse*Star uses a pool-type geometry to eliminate many costly pipes, valves, and other equipment used in a loop-type reactor. This compact design has a very high power density inside the containment structure. On the other hand, Cascade uses ceramic materials to push the operating temperature to a very high value so that the thermal-to-electric conversion efficiency can be significantly raised.

Safety and Environmental Issues

Since radioactive materials will exist in any fusion reactor, attention must be paid to the amounts of such materials that might be released in an accident and the amounts that must be safely disposed of during the lifetime of the plant. One factor that is important in accidental-release scenarios is the quantity of stored energy that is present in a plant at any time. The HYLIFE reactor has about 10^{13} J of stored chemical energy in its liquid lithium (about the same as a liquid-metal fast breeder reactor). For comparison, the chemical energy in a railcar full of propane is about 10^{12} J. Even though a lithium fire releases energy more slowly than a gasoline fire, the large potential energy gives rise to the perception of significant hazard. The Pulse*Star and Cascade reactors use blanket materials that have no significant fire hazard. The only identifiable

stored energy in these reactors is the 10^7 J of electrical energy for the driver. This number can be compared to the 10^{12} J of stored thermal energy in any reactor of this size.

An indication of the size of the radioactivity problem is shown in Table 7-2 for five reactors. In addition to the induced activity of the three ICF reactors discussed in this article, the total radioactivity in a typical pressurized water reactor (PWR) and the induced activity in the Starfire magnetic fusion reactor are shown for comparison. The Cascade reactor (a LiO₂ blanket was used for these calculations) not only has the smallest amount of activity at shutdown, but its activity is short lived. In fact, at 50 y, the activated material from Cascade would satisfy U.S. government standards for shallow land disposal without dilution. It is the only one of the five reactors for which this is true.

Decay afterheat, shown in Table 7-3 is an indicator of the need for emergency cooling to prevent meltdown. If values are smaller than a few megawatts, it is likely that radiative cooling would be sufficient and no active cooling system would be required. Both Cascade and HYLIFE satisfy this criterion at shutdown, and again Cascade's value falls rapidly after shutdown.

Finally, the inhalation biological hazard potential (BHP), an indicator of the potential risk due to release of debris to the biosphere, is shown in Table 7-4. The BHP is the volume of air that must be mixed with the radioactive material in order to make it safe for humans to breathe. Of course, any given accident would release only a small fraction of the material but, again, the total amount is used as a relative indicator of the magnitude of the problem. Cascade again appears to have the lowest potential risk.

Two facts stand out in these comparisons:

1. All ICF reactors appear to be better than either a PWR fission plant or the Starfire magnetic fusion design for these figures of merit.
2. The Cascade reactor is superior to the other two ICF designs.

Development Issues

The three ICF designs have several development issues in common. All must, of

course, have a suitably high pulse rate and high energy drivers. All three could probably be made to work with either lasers or heavy-ion beams. Cascade requires illumination from two directions only and thus could not be used with many-sided illumination (direct drive targets). The other two could probably be made to work with more than two beams (but HYLIFE becomes more awkward as the number of beams increases). All three reactor designs must solve the problem of assured clearing of the beam path at high pulse rates. Finally, they all must address radiation damage from pulsed sources.

Some development issues are different for the three designs. HYLIFE must establish the engineering feasibility and reliability of high liquid-metal pumping rates. For the Pulse*Star design, the main issues are the lifetime of the inner screen, tritium containment, and the effect of ^{210}Po from the Pb activation. Cascade must address engineering questions about the practicality of the high operating temperatures and the lifetime of the ceramic granules.

In conclusion, Cascade and Pulse*Star represent new trends in ICF reactor design. The design ideas they embody offer the possibility of greater economic competitiveness and safety and reduced environmental impact when compared with earlier designs. They also represent alternative technological pathways to ICF power.

Author: W. J. Hogan

Cascade Reactor

Introduction

Cascade (Fig. 7-2) is a concept for an ultrasafe, highly efficient, easily built reactor to convert inertial-confinement fusion energy into electrical power. The Cascade design includes a rotating double-cone-shaped chamber in which a moving, 1-m-thick ceramic granular blanket is held against the reactor wall by centrifugal action. The granular material absorbs energy from the fusion reactions.

Accomplishments this year associated with Cascade included improvements to

Reactor type	Time after shutdown				
	0	1 d	1 m	30 y	100 y
PWR	1.7×10^7	— ^a	— ^a	— ^a	— ^a
Starfire	6×10^6	— ^a	— ^a	2400	292
HYLIFE	3×10^5	2×10^5	1.5×10^5	90	33
Pulse*Star	4×10^6	9×10^5	7×10^5	660	220
Cascade	2×10^5	7×10^3	6×10^3	11	8

^aNot available.

Reactor type	Time after shutdown	
	0	1 d
PWR	1.8×10^5	— ^a
Starfire	7×10^4	— ^a
HYLIFE	2×10^3	2×10^2
Pulse*Star	1.8×10^4	1.4×10^3
Cascade	8×10^3	11

^aNot available.

Reactor type	Time after shutdown	
	0	1 d
PWR	1.2×10^{10}	— ^a
Starfire	1.7×10^{10}	— ^a
HYLIFE	5×10^7	3×10^7
Pulse*Star	7×10^8	6×10^8
Cascade	8×10^6	1.3×10^6

^aNot available.

Table 7-2. Total neutron-induced activity (kCi) for a 1000-MW_e reactor operating for 30 y.

Table 7-3. Decay heat of induced activity (kW) for a 1000-MW_e reactor operating for 30 y.

Table 7-4. Inhalation BHP (km³) for 1000-MW_e reactor operating for 30 y.

simplify chamber design and lower activation. We switched from a steel chamber wall to one made from silicon-carbide (SiC), panels held in compression by SiC-fiber/Al-composite tendons that gird the chamber both circumferentially and axially (Fig. 7-2). We studied a number of heat-exchanger designs and selected a gravity-flow cascade design with a vacuum on the primary side. This design allows granules leaving the chamber to be transported to the heat exchangers using their own peripheral speed. The granules transfer their thermal energy and return to the chamber gravitationally; no vacuum locks or conveyors are needed.

The Cascade rotating-bed reaction-chamber concept¹—two conical frustums attached at their major radii—offers simplicity in its mechanical design and a number of favorable stress characteristics in its structural design. First, only about 1% of the neutron flux contacts the wall of the chamber. This minimizes swelling-induced stresses and radiation embrittlement. Second, with a 50% porosity, solid-granule blanket, the pressure pulse that follows each fusion pulse imposes only modest forces on the structural wall.² Third, the

Fig. 7-2. (a) The Cascade reaction chamber. (b) Detail of the 2-cm-thick SiC panels. (c) Detail of the SiC/Al composite ring girder.

blanket weight and centrifugal forces can be carried by support rollers located outside the reaction chamber.

The granule transport system connects the reaction chamber and the power conversion system. For Cascade, we use a centrifugal-throw system that relies on the peripheral speed of the granules to transport them to the power conversion system: no additional equipment is required. The centrifugal-throw system is, however, limited to heights determined by the chamber rotational speed.

We conducted experiments to determine the flow and thermal characteristics of the granular blanket. We found that a two-layer, two-material blanket is feasible. The inner layer and the inner portion of the outer layer of the blanket are composed of beryllium-oxide (BeO) granules that reach exit temperatures up to 2300 K and act as a

neutron multiplier. The remainder of the outer layer is composed of lithium aluminate (LiAlO_2) granules that breed tritium. About one-third of the fusion energy is deposited in a thin, fast-moving layer of the inner-blanket surface. The average blanket exit temperature is about 1700 K, which permits thermal efficiencies of up to 60%.

In subsequent articles, we discuss the structural design of the Cascade reaction chamber, optimization of various blanket designs and their respective thermal and neutronic characteristics, report on methods of granule fabrication, and include a study that shows that the BeO granules on the inside surface of the blanket can be recondensed before the next pulse (0.2 s later). Our selection of a BeO-LiAlO₂ blanket was somewhat arbitrary because each blanket has its design advantages. However, our choice gave a high average blanket exit

temperature, a tritium-breeding ratio greater than unity, and BeO granules, which have good resistance to thermal shock, on the inside blanket surface.

Cascade continues to be a particularly promising concept. We hope to complete our basic study of Cascade and report the final design in the next calendar year.

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Chamber Structural Design

The low-activation chamber concept for Cascade, shown in Fig. 7-2(a), is constructed of 2-cm-thick, box-shaped panels of SiC that are each about 50 cm on a side [Fig. 7-2(b)]. The chamber is held together in compression by an axial and circumferential net of 3-cm-diam SiC-fiber/Al-composite tendons. Because SiC can operate at high temperatures, we avoid the need for external cooling and minimize neutron activation. The bottoms of the panels form the chamber wall. The sides of each panel are 10 cm high and form faces that butt against one another. The joints between panels include a compliant bond, possibly of an uncured ceramic compound that is fired during reactor operation to form a permanent joint. The corners of each panel are 15 cm high—5 cm higher than the sides—and form a corner post that supports the net of tendons. This net can be prestressed so that the SiC panels are always in compression. There is also thermal insulation between the chamber and the tendons to maintain the tendon's temperature within their usable range.

The chamber is ringed circumferentially with a pair of 40-cm-deep Al-composite girders that serve both to support the chamber and to provide a surface for chamber rotation [Fig. 7-2(c)]. The girders are mounted on 24 supporting wheels—six axes, carrying two wheels each, on each ring. Four of the 12 axes are motor driven.

Modular, prestressed SiC panels offer some very useful advantages for the chamber. First, the configuration lends itself to prestressing techniques. In an application similar to Cascade that included proof-testing techniques,³ α -SiC was found to have a design tensile-stress allowable of 200 MPa, and a compressive-stress allowable four to five times greater. The maximum

prestress required is 60 MPa. Young's modulus for α -SiC is 360 GPa, and the buckling stress for α -SiC is 8 GPa, assuming a value for Poisson's ratio of 0.3. The rupture stress is greater than 400 MPa. Consequently, the design has very high allowable prestress and applied working stress. Second, SiC has very high abrasion resistance. This is a very useful attribute, particularly in the present application in which the granules are expected to slide along the SiC wall. Third, SiC's activation is low, even with presently achievable impurity levels, which allows some hands-on maintenance, high inherent public safety, and permits shallow land burial of wastes only 50 years after shutdown.^{4,5} Finally, by using properly prestressed and insulated SiC, the structure and prestressing tendons can be radiatively cooled. The resultant temperatures and temperature gradients produce negligible thermal stresses and avoid the need for active cooling systems.

The prestressing material is an Avco SiC-fiber/Al composite.⁶ At a 45 to 90% fiber-volume fraction in Al, this material has an axial tensile strength of about 970 MPa at 700 K. A composite with 6063 Al alloy is preferred in the present application because it has only low-activation alloying elements.

The insulation material is fibrous, like Fiber-frax⁷ or Min-K.⁸ Foil-type insulators,⁹ although superior to the fibrous type, are not recommended because they commonly use thorium, zirconia, and other metals that would result in higher activation levels.

Chamber Stress Analysis. The sources contributing to the total stress in the Cascade reaction chamber were identified and calculated for a metallic chamber design in Refs. 1, 2, and 10. Following the methodology and equations presented in the references, the stresses were calculated for a SiC chamber with a 2-cm-thick wall, and a 1-m-thick Li₂O blanket rotating at 40 rpm.

Mechanical stress from centrifugal action was calculated using

$$\Gamma_c = \frac{m_b \omega^2 R^2}{A t \cos \theta} \quad (1)$$

where

Γ_c = hoop stress in the chamber wall = 20 MPa,

Cascade Reactor

- m_H = blanket mass = 200 000 kg,
 ω = chamber rotational speed = 40 rpm,
 R = maximum chamber wall radius = 5 m,
 A = chamber surface area = 250 m²,
 l = chamber panel thickness = 2 cm, and
 θ = chamber wall half angle = 35°.

Thermal stress was calculated using

$$\Gamma_{th} = \frac{\alpha E \Delta T_w}{2(1 - \nu)} \quad (2)$$

where

- Γ_{th} = thermal stress in the chamber wall = 5 MPa,
 α = SiC coefficient of thermal expansion = $4.9 \times 10^{-6} \text{ K}^{-1}$,
 E = Young's modulus = 360 GPa,
 ΔT_w = temperature difference across the chamber wall = 4 K, and
 ν = Poisson's ratio = 0.3 (assumed).

The stresses due to pressure from vaporized material, and shock due to x rays, debris, etc., were scaled from Ref. 10 to the present chamber-wall thickness and material properties. The results are summarized in Table 7-5.

The maximum stress of 60 MPa is acceptable in comparison to the 200 MPa tensile strength expected in the SiC tiles that make up the reaction chamber. In fact, the tiles

would be preloaded in compression to this maximum stress level and thus not experience any tensile loading.

The shear stresses on the panels were calculated by assuming the load between the ring girders and chamber wall occurs on 3 m of chamber-wall circumference in each of four regions around the horizontal (two regions per ring girder). The maximum shear stresses on the 0.02-m panels are 4 MPa, well below both the 400 MPa rupture stress and 8 GPa critical buckling stress. Therefore, the panels are not expected either to rupture or buckle. Because considerable flexibility exists in the structural design of the panels, the 2-cm-thick wall is chosen on the basis of ease of fabrication.

The tendon diameter is determined by the maximum 60-MPa preload. Assuming single Si-fiber/Al-composite tendons at 700 K prestressed to 80% of their tensile strength over the entire chamber, the required tendon diameter is 3.0 cm.

Because of its large size, the stresses in the reaction-chamber wall are comparatively low; however, the support requirement generates large forces virtually normal to the wall. To distribute these local forces, a pair of 40-cm-deep, SiC-fiber/Al-composite ring girders surround the chamber at two positions, providing a track for the rotation equipment. These girders distribute support reaction into the prestressed panels in the form of shear that is easily carried in the panels at a stress of 4 MPa.

The stress analysis of the ring girders shows that the rotating chamber can be supported by two rings of I-shaped cross section, mounted on a set of wheels. Assuming the reaction loads are distributed over a 60° angle, the maximum moment (M_{max}) and shear (V_{max}) along the circumference of the ring are $3.3 \times 10^5 \text{ N} \cdot \text{m}$ and $2.5 \times 10^5 \text{ N}$, respectively. The yield stress of the composite was assumed to be two thirds of the tensile strength at 700 K. A girder composed of 23- × 2.5-cm flanges and a 1.3- × 36-cm web satisfies the requirements on stress allowables and geometric constraints of the American Institute of Steel Construction (AISC).¹¹ Calculated bending stress in the flanges is 140 MPa, and the shear stress in the web is 55 MPa.

A support-wheel load of around 80 000 N is used, which requires 24 supporting wheels. Support-wheel diameter is approximately 0.75 m. Six axles, each carrying

Table 7-5. Summary of reaction chamber hoop stresses.

Source of stress	Stress (MPa)
Centrifugal action ^a	21
Thermal stress ^b	5
Pressure from vaporized granules ^{a,c,d}	18
Shock caused by x rays, debris, and neutrons (400-MJ, 5-Hz chamber) ^{b,e}	34
Maximum stress	44 to 60

^aPrimary stress.

^bSecondary stress.

^cStress caused by pressure from ablated material, assuming no heat transfer. The value given is twice the static value, to account for dynamic loading.

^dPeak stress caused by pressure from ablated material occurs later than the peak stress caused by shock interaction; therefore, pressure and shock stresses are not additive. Totals shown include one or the other of these stresses.

^eShock-caused stresses are calculated at cylindrical limit. Our value for Cascade should be lower.

two wheels, are required on each ring girder. The axes will tilt and will be adjustably sprung to give control of the force per wheel and to compensate for any out-of-roundness in the ring girders. We can cool the support-wheel bearings by circulating gas or water through the bearing housings; no elaborate mechanical sealing is necessary.

The present configuration for the drive system has the entire power input for the chamber, with its granule blanket, through 4 of the 12 axes. The direct power required to lift a blanket circulation of 2700 kg/s to a height of 17 m is 450 kW, but considerable internal friction is anticipated; thus, 225 kW on each of the four shafts is a reasonable installed requirement. The shaft speed is 320 rpm. Because the motors are external to the vacuum chamber, the drive-shaft penetrations would incorporate radiation-shielded, oil-flooded vacuum seals. The SiC-fiberglass composite drive shafts are less than 8 cm in diameter.

Granule Transport and Systems Integration. The granules slide along the chamber wall under centrifugal action, and exit from the chamber through an equatorial slot onto an exposed shelf outside the chamber. The granules rest on the shelf at the angle of repose with respect to the centrifugal field. A controllable scoop

removes the granules from the shelf as uniformly through the cross section as possible, thereby providing the primary reactor flow rate. Uniform removal of granules minimizes collisions between granules and results in a homogeneous and efficient granule trajectory. There will be no convergent flows until the line of the first-released-particle's trajectory starts to reenter the stream.

The reaction-chamber heat-exchanger design is shown from two different angles in Fig. 7-3. The granule stream: ejected from the chamber shelf by the scoop is directed along a throw duct and is then caught by the heat exchanger. The heat-exchanger tubes are completely submerged below the granules. As granules are removed from the chamber, an equivalent amount of granules automatically flows through the heat exchanger and the hopper, and back to the reaction chamber.

We examined the physical space requirements and configuration of the components that make up the primary heat transport system of Cascade. For vacuum heat transfer, the total height requirement on the throw system is that due to the heat exchanger, and the distribution and collection plena above and below the heat exchanger. The plena heights are established by the 500 m² heat-exchanger frontal surface area

Fig. 7-3. Two views of the proposed Cascade heat-exchanger design.

requirement and the angle of repose of the granules.

For the cascade gravity-flow vacuum heat exchanger, assuming frontal dimensions of 30-m width and 17-m depth, the plena heights required for adequate distribution within the heat exchanger and transport back to the reaction-chamber inlets are approximately 5 m each. These are combined with a 4-m heat-exchanger active height and 3-m clearance for the particle trajectories for a total throw height requirement of 17 m. This can be provided by a chamber rotation speed of 40 rpm or greater. The total vacuum boundary height is thus 23 m, assuming a 5-m-radius reaction chamber and an additional 1 m for the granule exit shelf and ground clearance. These dimensions give a good indication of the size requirements of the Cascade reactor plant.

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Granule Chute Flow Experiments

We conducted a series of experiments with the inclined rectangular chute shown in Fig. 7-4 to determine the flow characteristics of the Cascade granular blanket. These

Fig. 7-4. Test apparatus for the granule chute flow experiments.

experiments, combined with the analysis described in the next article, were used to establish the 35° half-angle of the Cascade double-cone-shaped chamber. Our experiments show that a two-layer moving granular blanket is feasible. The top layer is thin, has high speed, and has almost complete mixing. The rest of the blanket moves more slowly, with a velocity that decreases with radius in a parabolic fashion. Two-layer flow enhances the thermal properties of the blanket because about one-third of the fusion energy is deposited in the thin, fast-moving layer. Hence, the temperature rise in this layer can be kept closer to that in the remainder of the blanket and a higher average blanket exit temperature is obtained.

Experimental Description. The chute, shown in Fig. 7-4, is 250-mm wide and inclined at about 35°. We used type-O, Monterey-crystal, amber-color sand in the experiments. This sand has an angle of repose of about 34° and a grain size of about 0.5 mm. We noted, however, that the slope of the sand on the chute after experiments was closer to 32°, which indicates that dynamic conditions may lower the effective angle of repose.

We conducted the experiments by filling the upper container (positioned on a forklift), with sand (Fig. 7-4). The upper container was placed so that a clear plastic pipe extending from the bottom of the container ended in an inlet hopper that led to the entrance of the inclined chute. When a plate in the bottom of the upper container was opened, sand would fill the inlet hopper to the plastic pipe exit. As sand moved out of the hopper through a bottom-opening weir and down the chute, replacement sand from the upper container would automatically refill the hopper.

Sand was prevented from flowing out of the chute until the start of a test by a removable plastic barrier just upstream of the bottom end of the chute. When the plastic barrier was removed, sand would flow to the bottom end, out one or more of five exit weirs and into a second container (see Fig. 7-5). The second container was located on a scale so we could determine the flow rate. We adjusted the flow rate out of the various weirs by changing the angle of individual weir exit plates. If an exit plate was close to horizontal, then no sand would flow out of that particular exit weir. We

Fig. 7-5. Time exposure of two-layered flow on a chute inclined at 36.6° . Vertical markers are 2 cm apart and horizontal markers are 10-cm apart.

conducted tests with the middle three weirs closed—with flow only out of the top and bottom weirs. Under these conditions, sand from a thin, fast-moving layer on the chute was removed solely by the top weir; the rest of the sand was removed by the bottom weir. We did not close the bottom weir because this would establish a zone upstream where the sand was stationary. If stationary sand was present in the Cascade blanket, it would continue to heat up until it melted, which would destroy the granular nature of the blanket. The middle three weirs were not used: using the bottom weir only was more effective and produced similar results. In later tests, we used two exit containers and two scales so we could establish independently the flow rate from the top and bottom weirs.

The sand was seeded with a small amount of white-colored alumina that had both a lower density and a larger average grain size than the sand (~ 3 mm). Hence, the alumina granules remained on the top layer of sand as it flowed down the chute, and acted as a marker for motion pictures. We positioned flat mirrors at an angle so that when viewed from the side, the top,

side, and bottom of the sand could be seen simultaneously with a motion picture camera (Fig. 7-5).

In this test, our data consisted of geometrical measurements to establish the chute inclination, sand flow rates, and motion pictures taken at 50 frames/s. By viewing the motion pictures frame by frame, we determined the velocity of the sand as a function of position from the bottom of the chute. In some cases we took time-exposure photographs, an example of which is shown in Fig. 7-5. In this figure, the chute is inclined at 36.6° .

Results. Referring to the experiment shown in Fig. 7-5, the sand entered the chute through a 2-cm-high, bottom-opening weir in the inlet hopper. Sand at the upper right of the chute is flowing down at supercritical speed. Because of the time exposure of the photograph, the sand is seen as a blur rather than as individual grains. In the center of the photograph, the flow changes abruptly with only a thin layer at the top flowing at supercritical speed down the chute and out the top exit weir. The bottom sand layer is stationary and individual sand grains can be seen. No flow was

permitted out the bottom four exit weirs for this particular test. The center of the photograph shows the equivalent of a hydraulic jump with water flow. Note that the top mirror shows supercritical flow over the entire width of the chute. The bottom mirror shows stationary flow separated from supercritical flow by a curved line.

Had both the bottom and top exit weirs been opened, a region of the desired two-layer flow would have been produced. The region of stationary sand would have been moving at subcritical speed, the boundary between subcritical flow and supercritical flow seen on the bottom mirror would have moved upstream, and the curvature of the boundary would have been substantially reduced. The effects of the sides of the chute that would decrease the velocity of the granules would still have been present but would have been negligible. Further, the depth of sand in the region where there was a thin supercritical top layer flowing over a thick subcritical flow layer would have been nearly uniform rather than increasing to the lower left as in Fig. 7-5.

The boundary between two-layered flow and complete supercritical flow moves even further upstream if the inclination of the chute is reduced to a value closer to the sand's angle of repose. A small region of complete supercritical flow is always present just downstream of the bottom-opening weir in the inlet hopper, but it is only a few centimeters long when the inclination angle is between 35° and 36° .

Figure 7-6 shows the velocity profile when two-layered (supercritical/subcritical)

flow exists on a chute inclined to 35.9° .

These data were obtained from motion pictures, in which individual grains of sand on the side of the chute were followed from frame to frame. Flow rate was measured by weighing the amount of sand flowing into the exit container over a period of ~ 30 s. The flow rate so measured was $940 \text{ cm}^3/\text{s}$. Flow rate calculated by integrating the velocity profile was about 20% lower. The two values are in reasonable agreement and give confidence in the accuracy of both flow rate and velocity measurements. Further, the calculated flow rate is lower, which is expected because sand grains at the side of the chute (where the velocity measurements were made) travel somewhat slower than at the center of the chute.

Data reduction of tests where flow rate in both layers was measured individually has yet to be completed. However, the velocity profile would not change substantially from that shown in Fig. 7-6.

Conclusions. We found that two-layered granular flow existed on a chute that was inclined slightly above the angle of repose. The top layer is thin and moves with high speed (supercritical flow). The remainder of the flow has a parabolic (subcritical) velocity distribution, with the velocity near the chute roughly one-tenth that at the interface of the thin top layer.

The two-layered flow is advantageous for Cascade because about one-third of the fusion energy is deposited in the thin, fast-moving layer of the blanket. The high speed permits a reasonable change in temperature as the granules pass through the blanket. The velocity distribution determined from these tests allows design of a thermally efficient blanket for the Cascade chamber.

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Motion of Single Frictional Granules on Rotating Cylinders and Cones

Laboratory tests last year verified that we could achieve a granular layer flowing continuously on the inside of a rotating

Fig. 7-6. Velocity profile in two-layered flow down a chute inclined at 35.9° .

horizontal cone.¹² This year, we obtained multilayer flow on inclined chutes. We think that such a flow pattern can also be obtained on the inside of rotating cones with proper design of the inlet and exit regions. The motion of the entire flowing granular blanket in a Cascade chamber is not simple to predict in detail. Many aspects of the flow behavior of granular solids are not well understood. Lack of verified constitutive relations and flow equations prevent us from putting together complete numerical models of the granular flow. Instead, as a means of evaluating Cascade design concepts, we have used various physical models¹² and analyses of certain limiting cases.^{13,14}

It is necessary to maintain a clear, open space in the central region of the chamber for injection of fusion pellets and the laser or ion beams into the Cascade reactor. Also, to determine the final chamber shape and size, it is necessary to know what cone half-angle the inner granular surface is likely to form. We analyzed the motion of single frictional granules on rotating cylinders and cones to address these two questions.

Granules on a Rotating Cylinder. Intuitively we know that if a cylinder containing granules is spun fast enough, the granules will stay on its inner surface. However, in laboratory tests of granules in a rotating cylinder,¹² we observed some granules to be raining into the open central region, even when the centrifugal force exceeded gravity by more than 20%. To determine what minimum rotation rate would maintain the granules in continuous contact with the surface, we considered the behavior of individual frictional granules on the interior surface of a rotating cylinder.

The motion of a single granule can include circumferential sliding as the cylinder rotates. An examination of the forces acting on a granule moving with a rapidly rotating cylinder shows that the point 45° above the

horizontal is the place where sliding is most likely to begin as rotational speeds are reduced. The ratio of the tangential forces to normal forces acting on a granule rotating with the cylinder is given by

$$\frac{\text{tangential force}}{\text{normal force}} = \frac{mg \sin \theta}{mg \cos \theta + m\omega^2 R} \quad (3)$$

where m is the mass of the granule, g is the acceleration due to gravity, θ is the angular position of the granule measured from the bottom of the cylinder, ω is the angular velocity of the cylinder, and R is the cylinder radius.

This ratio reaches a maximum when $\cos \theta = -\sin \theta$, or when θ has a value of $3\pi/4$. No sliding will occur at this point if the tangential force is less than the coefficient of friction, μ , times the normal force. Hence, for all rotational speeds greater than $\omega_{\text{no slip}}$, given by

$$\omega_{\text{no slip}} = \left[\frac{g}{R} \left(\frac{1}{\mu} \sin \frac{3\pi}{4} - \cos \frac{3\pi}{4} \right) \right]^{1/2} \\ = \left[\frac{g}{R} \left(1 + \frac{1}{\mu} \right) \right]^{1/2} \quad (4)$$

no circumferential sliding is possible for the granule as it moves with the cylinder wall. The values for this minimum upper-bound rotational speed are given in Table 7-6, and shown in Fig. 7-7. The dimensionless rotational speed (shown on the ordinate in Fig. 7-7) equals unity when the centrifugal force acting on a granule equals the granule's weight. At rotational speeds below the $\omega_{\text{no slip}}$ value of Eq. (4), it is possible for the particle to move relative to the cylinder wall; however, circumferential sliding does not mean that the granule will lose contact with the wall. Contact with the wall will be lost near the top only if the effective angular velocity of the granule about the cone

Radius R (m)	Coefficient of friction (μ)	Minimum rotational speed $\omega_{\min}$ (rad/s)	Dimensionless rotational speed (R/g) ^{1/2} $\omega_{\min}$	Upper bound (R/g) ^{1/2} $\omega_{\text{no slip}}$ Eq. (4)	Empirical fit (R/g) ^{1/2} $\omega_{\min}$ Eq. (6)
1.0	$\tan 24^\circ$	4.38-4.39	1.399-1.402	1.515	1.395
1.0	$\tan 30^\circ$	4.08-4.10	1.305-1.310	1.390	1.309
1.0	$\tan 35^\circ$	3.90-3.91	1.246-1.249	1.310	1.251
0.045 ± 0.001	$\tan 24^\circ$	20.8-20.9	1.499-1.416 + 0.015	1.515	1.395
0.045 ± 0.001	$\tan 35^\circ$	18.5-18.6	1.254-1.260 + 0.015	1.310	1.251

Table 7-6. Calculated minimum rotational speed required for a single frictional granule to maintain continuous contact with the inside of a rotating cylinder wall.

Fig. 7-7. Minimum rotational speed required to maintain a frictional granule in continuous contact with the wall of a rotating cylinder as a function of the coefficient of friction between the granule and the wall.

axis ω_p is low enough for the centrifugal force to be less than the normal component of gravity mg_n at that location:

$$m\omega_p^2 R < mg_n \quad (5)$$

We numerically integrated the equations of motion for various frictional granules on the inner surface of horizontal rotating cylinders—including the effects of friction, gravity, and centrifugal forces—to determine the minimum rotational speed $\omega_{\min}$ necessary to keep the granules in continuous contact with the cylinder wall. We used a radius of 0.045 m in some of our calculations to correspond with the radius of the laboratory tests with rotating cylinders.¹² In the other calculations we used a radius of 1.0 m to correspond with the inner-surface radius near the inlet region of the Cascade chamber. The results of our calculations are shown in Table 7-6 and Fig. 7-7.

The calculated values of $\omega_{\min}$ correspond almost exactly to the upper-limit curve shifted to the left by 0.125. Thus we obtain an empirical equation that is a reasonable approximation to the calculated results for $\omega_{\min}$ by adding 0.125 to μ in Eq. (4):

$$\omega_{\min} = \left[\frac{g}{R\sqrt{2}} \left(1 + \frac{1}{\mu + 0.125} \right) \right]^{1/2} \quad (6)$$

The values for $\omega_{\min}$ calculated in Eq. (6) are also given in Table 7-6 and plotted in Fig. 7-7.

In the calculations of individual-granule motion at rotational speeds higher than the upper-limit curve in Fig. 7-7, no circumferential movement of the granule relative to the rotating cylinder wall was observed. As the rotational speed was reduced, intermittent circumferential sliding of the granule on the cylinder wall occurred. As the rotational speed approached the minimum value required to maintain continuous contact, the amplitude of the sliding increased. Figure 7-8 shows the calculated angular displacement vs time for a granule on a cylinder rotating at two speeds. The top curve is a rotational speed just below the no-slip, or upper-limit curve of Fig. 7-7; the lower curve is for a rotational speed about 1% faster than $\omega_{\min}$. The calculated maximum amplitude of the circumferential sliding that occurred without a granule leaving the surface depended on the coefficient of friction, with much larger displacements being calculated for low coefficients of friction (the cylinder also rotated faster in those cases).

In laboratory tests,¹² most of the granules stayed on the cylinder wall ($R = 0.045$ m) at 172 rpm, which corresponds to 1.22 in the dimensionless rotational-speed units of Fig. 7-7. The measured coefficient of wall friction μ_{wall} acting between the granules and the Plexiglas had a value of $\tan 24^\circ$, or approximately 0.445. A point corresponding to $\mu = 0.445$ at a dimensionless rotational speed of 1.22 would fall well below the $\omega_{\min}$ curve on Fig. 7-6, and thus should correspond to material falling off the wall of the cylinder. Yet, at this speed, most of the granules stayed on the wall. This apparent contradiction is explained by taking into account the effects of the remainder of the granular material in the cylinder. Individual granules attempting to slide down the wall exert forces on the material below them. If this force is not sufficient to deform the granular layer, no significant movement occurs. The shear strength of a non-cohesive granular material increases as the confining lithostatic load on the granular layer in the bottom half of the cylinder. This load gives the granular layer sufficient strength to

resist the small deformation forces due to the tangential forces that are slightly higher than the friction forces acting on granules in the upper part of the cylinder. If the granular layer does not deform, then there will not be any circumferential sliding because there are symmetrical tangential forces acting on the entire hoop of granular material (with equal total forces in the positive and negative circumferential directions). Because of the granular layer's material-strength property, most of the granules remained on the cylinder wall whenever the rotational speed was high enough to provide the shear strength for the granular layer to resist deforming. Because we do not know quantitatively how the shear strength of the granular material used in the laboratory tests depends on confining load, we could not predict the exact rotational speed that would cause this to happen; however, it is apparently well below the rotational speed required to keep a single granule on the wall in this small laboratory-scale apparatus.

The granules that make up the innermost surface of the rotating layer do not experience any confining load from material on top of them. For the innermost-surface granules, our single-granule analysis should therefore be valid. However, these granules are on a layer of granular material with an effective coefficient of surface friction equal to the tangent of the angle of repose for that material. For the granules in the rotating-cylinder tests, the angle of repose was near 35° (there is evidence that under dynamic conditions the effective angle of repose for granules may decrease by a few degrees as we discussed in "Granule Chute Flow Experiments" earlier in this section).¹² The horizontal line through the data point near the bottom of Fig. 7-7 represents possible interpretations of the granule tests in the small, rotating cylinder. The left end of this line corresponds to the measured wall-friction angle of 24° ; the marked point corresponds to the measured angle of repose, θ_r , of $35^\circ \pm 2^\circ$, which we think more closely approximates the effective angle of friction for a granule on the inner surface of the granular layer. Because the entire line is below our $\omega_{\min}$, it is not surprising that some surface granules rained into the central region. We expect that rotational speeds

Fig. 7-8. Circumferential displacement vs time for a frictional granule inside a 1-m-radius cylinder. The upper curve represents a speed just below the upper-limit curve of Fig. 7-7; the lower curve represents a speed about 1% higher than the calculated $\omega_{\min}$ curve of Fig. 7-7.

10% higher than the maximum value of 172 rpm obtained in the small cylinder tests would eliminate all raining particles. However, no laboratory verification of this prediction has been performed yet.

Granules on a Rotating Cone. The motion of individual granules on the inner surface of a rotating cone can be analyzed similarly to granules in a cylinder. In a cone, the analysis is slightly more involved because the frictional force and the other tangential forces (gravity and the Coriolis force) are not always collinear. The tangential, or surface force can be considered to be composed of two orthogonal components: one purely circumferential, and the other in the direction of a projection of the cone axis on the surface. We call this the axial surface direction. The Coriolis force always acts in a circumferential direction with a magnitude that depends on the rate of change of the radius (i.e., on the axial surface velocity). Gravity, acting in the axial surface direction, is constant independent of the granule's position or velocity: gravity in the circumferential direction is exactly the same as in the case of a cylinder. The maximum magnitude of the frictional force depends on the current normal force and the coefficient of friction, while the direction of the friction force depends on the direction of the current surface velocity (because friction acts to oppose current sliding motion). In our numerical integration of the equations of motion for a single frictional granule, we used the following expressions for the three components of the granule's acceleration in the normal, circumferential, and axial surface directions when it was sliding on the cone's surface:

$$a_n = \left[\left(\frac{v_\theta}{r} + \omega \right)^2 r + g \cos(\omega t + \theta) \right] \cos \alpha \quad (7)$$

$$a_{\theta} = -2 \left(\frac{v_{\theta}}{r} + \omega \right) v_a \sin \alpha - g \sin (\omega t + \theta) - \mu a_n \left(\frac{v_{\theta}}{|v_i|} \right), \quad (8)$$

$$a_a = \left[\left(\frac{v_{\theta}}{r} + \omega \right)^2 r + g \cos (\omega t + \theta) \right] \sin \alpha - \mu a_n \left(\frac{v_a}{|v_i|} \right), \quad (9)$$

where

- ω = rotational speed of the cone,
- t = time,
- θ = angle of granule relative to its starting position,
- r = radius of cone at current position of granule,
- v_{θ} = circumferential velocity of granule relative to point at same radius rotating with cone,
- v_a = axial surface velocity of granule on cone surface,
- α = half angle of cone,
- g = acceleration due to gravity,
- μ = coefficient of friction between granule and wall,
- $v_i = (v_{\theta}^2 + v_a^2)^{1/2}$ = total tangential (surface) velocity.

The normal-direction acceleration a_n of Eq. (7) is opposed by the cone surface, so no net acceleration occurs in this direction. We numerically integrated Eqs. (8) and (9) to obtain the displacement-time history for individual granules moving on the interior of a rotating cone.

We performed these calculations for several rotational speeds, effective surface friction coefficients, and cone half-angles. Initially, we assumed that the actual inner surface of a granular blanket would be a cone with a half-angle equal to the angle of repose for the material (independent of the choice of the shape for the outer chamber wall). We further assumed that the effective friction coefficient for particles on the surface would be the tangent of the angle of repose. We determined the minimum rotational speed required to keep a granule in continuous contact with the cone surface. In our calculation, the tangent of the cone half-angle α was equal to the effective

surface-friction coefficient μ , and the granule was initially moving with the cone at the bottom of the inlet region. We found that the minimum rotational speed was about 50 rpm for cone half-angles from 24° to 35°. (The inner-inlet radius for these calculations was 1 m.) We also found that in all cases, the calculated axial motion of the granule involved considerable acceleration, and the resulting residence time was significantly less than what we would extrapolate from our laboratory experiments.¹³ We interpret this result to indicate that the actual inner surface of the flowing granular layer will be at a cone half-angle somewhat below the angle of repose (i.e., $\tan \alpha$ will be somewhat less than the effective coefficient of surface friction).

To determine the probable shape of the inner blanket surface, we reasoned that a frictional granule would not accelerate without limit on a surface whose effective slope was at or below the angle of repose. But this effective slope needs to be determined from the effects of all forces acting on the granule. With a fixed effective-surface-friction coefficient $\mu \approx \tan 35^\circ \approx 0.700$, we examined various cone surface angles to see under what circumstances a granule would undergo unlimited acceleration, come to rest, or lose contact with the wall.

For each assumed inner surface half-angle, we obtained a different $\omega_{\min}$ (for a cone with an initial radius of 1 m). The values varied from the cylindrical value of near 37 rpm for a half-angle of zero to 50 rpm when the cone half-angle became equal to the arc tangent of the assumed coefficient of friction. However, reduction of the cone half-angle only a few degrees below the value corresponding to $\tan^{-1} \mu$ eliminated the unrealistic unlimited acceleration discussed above. Above 45 rpm, all sliding relative to the cone surface ceased after the first meter of length down the cone. This implies that the inner surface cone half-angle will probably be very close to the angle of repose for the material, with perhaps a small region near the inlet with a cone half-angle a few degrees smaller than the angle of repose.

We also considered the motion of a single granule simulating the movement of the entire blanket in a Cascade chamber. To do this, we added a constraint to the maximum axial velocity of the granule that corresponds

to the total expected mass flow rate of granules through the chamber, assuming a constant-thickness granular blanket. With this restriction on v_n , we examined the motion of a granule with $u \approx \tan 24^\circ \approx 0.445$ on a 35° cone with an initial radius of 2.0 m, corresponding to a granule in contact with the chamber wall near the inlet region. We looked at rotational speeds ranging from 40 to 51 rpm and found that such a granule moved steadily down the cone according to the constraint condition, undergoing as it went very small, almost sinusoidal, circumferential oscillations in its trajectory. The amplitude of these oscillations varied from ± 1.5 cm at 40 rpm to less than ± 0.5 cm at 51 rpm. Our calculations indicate that circumferential oscillations of the flowing granular layer, if they occur at all in a full-sized Cascade chamber, will be very small and will have a negligible effect on the granular blanket thickness at the top of the chamber.

We conclude that the Cascade chamber operating at 50 rpm will have an open central region with no granules raining from the inside surface of the blanket and that the inner blanket surface will have a half angle close to the angle of repose of the granular material.

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High-Temperature Blanket Design Options

A unique characteristic of DT-fueled inertial confinement fusion is that a large fraction of the energy yield, typically $\sim 30\%$, is in the form of a surface heat flux from charged particles, reflected laser light, and x rays. The balance of the energy yield is in the form of volume heating from neutron absorption. The ideal ICF reactor could convert a high percentage of this energy to electricity. In the Cascade reaction chamber concept,^{15,16} the two-layer blanket configuration promises high energy conversion. The two-layer blanket combines a thin, fast-moving surface layer of high-temperature (>2000 K) material to absorb and transport the surface heat flux, with a thicker, slower-moving layer of high-temperature tritium-breeding material

that can be used to absorb and transport the neutron energy. The high-temperature material can be used to achieve very high power conversion efficiency (up to 60%) while still allowing for adequate tritium breeding.

In this article, we present a number of high-temperature blanket-design options available to Cascade. We examine the high-temperature material options for the surface layer and include analyses of mechanical properties for thermal stress resistance and irradiation stability. Material temperature-limit suggestions are based on compatibility of the surface layer with materials used for tritium breeding, the chamber, and the heat exchanger. Finally, we calculate the overall thermodynamic, and thus power-producing, potential of the selected blanket designs, and arrive at a reference design for subsequent, secondary system design and tritium analyses.

High-Temperature Material Options.

Ceramic materials, such as SiC and BeO, are usable at temperatures of ~ 2000 K and higher. A high-temperature surface layer of ceramic material can significantly increase the power-conversion efficiency of a Cascade reactor by permitting higher temperatures than the typical 1400 to 1500 K maximum of solid-ceramic tritium breeders such as Li_2O and LiAlO_2 . Table 7-7 presents mechanical properties for surface layer materials at elevated temperatures. We used these properties to evaluate granule resistance to the thermal stress induced by the surface-heat flux and to estimate irradiation lifetimes. The survey was restricted to low-activation materials that would maintain the safety and environmental attractiveness of Cascade.

Data for metallic-beryllium compounds, Li_2O , and LiAlO_2 are shown in Table 7-7 for comparison. Data from Ref. 17 are for high-density polycrystalline materials prepared by hot pressing, except for SiC, where they refer to chemical vapor deposited (CVD) β -SiC. The thermal stress parameter shown in Table 7-7 is calculated according to

$$R = \frac{S(1 - \nu)}{\alpha E} \quad (10)$$

where R is the thermal-stress parameter (a measure of the material strength relative to the thermal stress), S is the flexural

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Material	Melting point (K)	Density (Mg/m ³)	Thermal expan. (10 ⁻⁶ K)	Thermal conduct. (W/m·K)	Young's modulus (GPa)	Poisson's ratio	Flex. strength (MPa)	Specific heat (J/kg·K)	Thermal stress parameter	References
Al ₂ O ₃	2320	3.99	9.0	7	340	>0.45 ^d	300	b	54	16,17
MgAl ₂ O ₄ (spinel)	2408	3.59	9.0	6	206	b	250	b	100	16,17
MgO	3100	3.58	15.7	9	275	b	250	b	43	16,17
SiC	>3000 ^c	3.21	4.9	50	360	0.2	500-1000	1300 ^d	230-450	16,17
Neutron-Multipliers										
Be ^e	1558	1.85	11.5	150	290	b	210-410 ^{e,f}	1883	47-92	16-19
Be-38Al ^e	918	2.08	16.2	210	190	b	210-280 ^{e,f}	b	51-68	16-19
Be ₂ C	>2400 ^c	2.24	10.8	23 ^e	214	0.1	100 ^g	1400 ^c	39	16-18
BeO	2800	2.90	9.1	29	360	0.22	390	2000	93	h
Tritium-Breeders										
LiAlO ₂	1883	2.55 ^e	12.4	2.2	b	b	b	1500	31 ⁱ	16-20
Li ₂ O	1706	2.00	35.9	3.6	54	0.25	28	2890	11	16

^aAt 1773 K unless otherwise noted.

^bNot available.

^cDecomposes.

^dRefs. 16-18, at 1800 K.

^eAt room temperature.

^fYield strength depends on form.

^gModulus of rupture at room temperature.

^hSee Tables 7-8 and 7-9 below.

ⁱAssuming flexural strength, Poisson's ratio, and Young's modulus of Li₂O.

Table 7-7. Typical properties of ceramic materials.^a

strength, ν is Poisson's ratio, α is the coefficient of thermal expansion, and E is Young's modulus. Poisson's ratio was not available for certain materials. For these materials, we used the Li₂O value in the calculation of R .

Table 7-7 shows that, among non-neutron multipliers, SiC is clearly superior in terms of thermal stress resistance, followed by the MgAl₂O₄ spinel. Among the beryllium neutron-multipliers, BeO exhibits the highest thermal stress resistance at elevated temperature, comparable to that of MgAl₂O₄. As shown below, Be₂C must be coated because of compatibility concerns, and thus its thermal stress resistance is less important. In all cases, however, these high-temperature ceramics exhibit thermal stress resistances that are at least an order of magnitude greater than that of Li₂O. Thus, they are more likely to survive the thermal shock induced by the surface heat load without experiencing the cracking and chipping predicted for Li₂O.

The surface layer of the Cascade blanket will also experience the highest neutron fluence. We examined two materials: SiC and BeO. The properties of single-phase β -SiC after neutron irradiation in a fission spectrum are documented in Ref. 19. In summary, at fluences up to $5 \times 10^{26} \text{ n/m}^2$ at 823 K, specimens of β -SiC showed no structural deterioration and a linear swelling of 0.4%. At 1773 K, the same linear swelling of 0.4% was reached in a fluence of

10^{26} n/m^2 . The unirradiated dimensions can be recovered by annealing. Because of the cyclic nature of ICF neutron loading and the recirculation of the blanket for energy transport, this annealing process may be inherent to the Cascade concept. Furthermore, SiC's flexural strength, Young's modulus, and thermal expansivity are insensitive to irradiation up to $5 \times 10^{26} \text{ n/m}^2$. The neutron fluence expected at the surface layer at a radius of 4 m for 3000 MW fusion is $1.5 \times 10^{26} \text{ n/m}^2 \text{ y}$, and thus an irradiation lifetime of 4 to 5 y appears possible for surface granules of SiC.

Data on radiation damage to BeO are available in Refs. 20 through 24. For the best irradiation stability, BeO should have a grain size ($\leq 10 \mu\text{m}$). At irradiation temperatures above 1300 K, the volume swelling is likely to be less than 3% for fast neutron fluence up to $5 \times 10^{25} \text{ n/m}^2$. Ultimately, microcracking will occur and the material will weaken. The limiting fluence is not well defined, but some fine-grained specimens have been shown to retain their strength at these fluences at 1300 K. BeO surface-granule lifetime would thus be about 0.5 y. The effect of pulsed neutron bursts is unknown, but they may reduce the limiting fluence. If longer lifetimes are required, the granules could be coated with SiC.

There is no significant data on irradiation effects on Be₂C, except for some transient pulsing exposures²⁵ to very low fluences in

a Sandia reactor. The cubic crystal structure of Be_2C would indicate a good resistance to irradiation damage. However, the formation of the transmutation products of helium, tritium, and lithium by neutron absorption in the beryllium will no doubt have a significant effect on the irradiation behavior of the Be_2C . Swelling may result from the agglomeration of helium in the Be_2C granules, unless the helium can diffuse out at the high operating temperature. With coated Be_2C , the effects of irradiation on the Be_2C kernel could be accommodated with no significant damage to the particles.

Compatibility Temperature Limits. One of the strengths of Cascade is its ability to attain a tritium breeding ratio of 1.05 or greater with a variety of material combinations. Surface-layer materials considered are SiC, BeO, and Be_2C , coated with either BeO or SiC. Tritium-breeder materials considered are Li_2O and LiAlO_2 because these are the currently preferred breeders^{2b} for fusion application. The reaction-chamber and heat-exchanger material is SiC.

To establish the maximum temperature limits for the candidate materials, we performed a thermochemical assessment of all possible material interactions and examined the mass transport and overpressure that would result from vaporization, assuming thermodynamic equilibrium. The resultant temperature limits and limiting criteria are summarized in Table 7-8. We calculated the mass transport assuming that the reactions both proceed to completion and are very conservative; no credit is taken for possible kinetics effect. Postulated material interactions were evaluated for overpressure by calculating the heats of reaction ΔH and standard free energies of reaction ΔF at 1800 K using JANAF²⁹ data. The standard entropy ΔS was then calculated from the second law,

$$\Delta F = \Delta H - T\Delta S \quad (11)$$

Over a reasonable temperature range, it can be assumed that ΔH and ΔS are constant; ΔF can be calculated at other temperatures using the ΔH and ΔS values at 1800 K. Because the processes of interest include gaseous species, the equilibrium constant K given by

$$K = \exp\left(-\frac{\Delta F}{RT}\right) \quad (12)$$

Material	Temperature limit (K)	Basis
SiC, bulk temperature	2100	Thermal decomposition overpressure of 1 Pa ^a
SiC/T ₂ O interface	No limit	Acceptable material transport assuming complete reaction of T ₂ O
SiC/Li ₂ O interface	1100	Overpressure of 1 Pa
Front surface	1200	Overpressure of 10 Pa
Within the blanket	1500	Stability of SiO ₂ layer on SiC
SiC/LiOT	No limit	Acceptable material transport assuming complete reaction of LiOT
SiC/LiAlO ₂	1400	Overpressure of 1 Pa
	1500	Stability of SiO ₂ layer on SiC, overpressure of 4 Pa
SiC/BeO	1700	Overpressure of 1 Pa
BeO	2300	Material phase transition
BeO/LiAlO ₂	No limit	LiAlO ₂ has much lower (more negative) free energy of formation than BeO·Al ₂ O ₃
BeO/T ₂ O	No limit	Overpressures much less than 1 Pa at 2300 K; acceptable mass transport assuming complete reaction of T ₂ O
Be ₂ C, bulk temperature	1600	Thermal decomposition overpressure of 1 Pa
Be ₂ C/T ₂ O	No limit	Acceptable mass transport assuming complete reaction of T ₂ O
Li ₂ O, bulk temperature	1200	LiOT mass transport (Ref. 16)
LiAlO ₂ , bulk temperature	1500	0.8 × melting temperature (Refs. 16-18, 27, 28)

^aOne-tenth of the base pressure inside the reaction chamber.

is a measure of the overpressure. The background pressures¹⁶ of T₂ and T₂O used in the calculations are 13 and 1.3 Pa, respectively.

We considered first the direct thermal decomposition of SiC. The equilibrium constant in this case is equal to the overpressure P_{Si} due to silicon evaporation. We found that to maintain a silicon overpressure less than 1 Pa, i.e., one order of magnitude lower than the background pressure, the SiC temperature must be kept below 2100 K.

We next examined reactions of SiC with T₂O. For these reactions, the total partial pressure at equilibrium is given by 2K. We found that the total partial pressure would be unacceptable over the 1800 K temperature range. However, the T₂O source rate is only 1.25×10^{-6} kg·mole/s for 3000 MW of fusion power at a tritium breeding ratio (TBR) of 1.4. Assuming that the T₂O continuously and fully reacts only with the SiC, the maximum resultant SiC loss rate is 1.5×10^3 kg/y. These figures are to be compared with typical front-zone flow rates of $\sim 10^4$ kg/s and a total available SiC mass of 6×10^4 kg. It is thus concluded that the reactions of SiC with T₂O do not impose operational temperature limits.

Table 7-8. Material compatibility temperature limits.

In the reaction of SiC with Li₂O, the equilibrium constant has the form $K = P_{\text{CO}} \times (P_{\text{Li}})^6$. If the gas phase has the reaction stoichiometry of $P_{\text{Li}} = 6P_{\text{CO}}$, then the total overpressure for this reaction is given by

$$P_{\text{Li}_2\text{O}} = 1.51 K^{1/7} \quad (13)$$

We found that the temperature limit at a SiC-Li₂O interface is approximately 1100 K for a maximum overpressure of 1 Pa, and 1200 K for 10 Pa. The latter is assumed in the present application because these limits are particularly conservative because they apply to a solid-to-solid reaction in which kinetics must limit the overpressure due to point-to-point contact limitations. Examination of the products reflects the fact that SiC has kinetic stability in the presence of Li₂O due to the formation of a SiO₂ film. After formation of this film, further reaction is limited by diffusion of the reactants. The high-temperature capability of SiC in the presence of Li₂O is thus limited by the stability of the SiO₂ protective layer. We conservatively estimated that the SiO₂ protective layer is stable when the SiO₂ is amorphous (below 1500 K). For the surface layer, where the stability of the SiO₂ cannot be preserved due to the surface heat flux, the limits based on overpressure are applicable.

The reaction of SiC with LiOT can be analyzed on a mass transport basis similar to that of the reactions of SiC with T₂O. Assuming a bulk average Li₂O temperature of 1300 K, the LiOT production rate¹⁶ is 7×10^{-5} kg/s or 3×10^{-6} mole/s. The corresponding SiC loss rate assuming complete reaction is 600 kg/y, which is quite acceptable; thus, no SiC temperature limit arises from LiOT considerations.

The reaction stoichiometry is the same for LiAlO₂ as for Li₂O. A temperature limit of 1500 K can be established based on both overpressure and SiO₂ stability considerations.

BeO will contact the SiC reaction chamber and heat exchanger. Reaction stoichiometry dictates that the total product overpressure is given by $P_{\text{BeO}} = 3K^{1/3}$. A temperature limit of 1700 K is necessary to limit overpressure to 1 Pa, i.e., one-tenth of the base pressure inside the reaction chamber. In addition to the reaction of BeO with SiC considered above, the reactions of BeO with LiAlO₂ and T₂O also need to be considered. With respect

to interaction with LiAlO₂, Li₂O·Al₂O₃ has a much more negative free energy of formation than BeO·Al₂O₃ and there is the no likelihood of a reaction.

In the reaction with T₂O, the equilibrium constant has the form

$$K = \frac{P_{\text{Be(OT)}_2}}{P_{\text{T}_2\text{O}}} \quad (14)$$

The total overpressure is much lower than 1 Pa at temperatures as high as 2300 K and thus does not present a limitation. In addition, assuming complete reaction of the BeO with the available T₂O, the BeO loss rate is 3×10^{-5} kg/s, or 1000 kg/y. This is also acceptable compared to a typical front-zone flow rate of 5×10^3 kg/s and total available BeO mass of $\sim 10^5$ kg.

For the thermal decomposition of Be₂C, experimental results³⁰ were taken directly. The beryllium overpressure is given in this case by $P_{\text{Be}} = K$ and shows that a temperature limit of 1600 K would keep the overpressure due to decomposition below 1 Pa. This temperature is relatively low and suggests that a SiC coating that would contain the overpressure be applied to the Be₂C, allowing the limiting temperature to be that of the SiC, i.e., 2100 K. With a coating, reactions of Be₂C with T₂O are precluded. Nonetheless, as in the case of BeO, mass transport via reaction of uncoated Be₂C with T₂O would not present a temperature limitation. Finally, it is noted that though Be₂C hydrolyzes to methane at low temperatures, it does not do so at the temperatures in Cascade.

Thermodynamics of Blanket Options. To minimize the number of blanket designs investigated in detail, we performed a preliminary screening of the possible material combinations based on the following considerations:

- Li₂O is the only ceramic tritium-breeder capable of achieving an adequate TBR without a neutron multiplier. Thus, a blanket with SiC as the high-temperature surface layer material and Li₂O as the tritium-breeder is attractive.
- If the surface layer is a neutron multiplier, then LiAlO₂ is more attractive than Li₂O as the breeder because it has higher temperature capability and avoids LiOT compatibility concerns.

Table 7-9. Maximum compatibility temperature limits of prime-candidate blankets (K).

Materials reaction chamber/ breeder zone/ front zone	Reaction chamber/ breeder zone interface	Breeder zone	Breeder zone/ front zone interface	Front zone	Front zone/ coating interface
1. SiC/Li ₂ O/SiC	2100	1400	1400 (peak) 1200 (average)	1200	2100
2. SiC/LiAlO ₂ /BeO	2100	1500	1500	1500	2300
3. SiC/LiAlO ₂ /Be ₂ C-SiC	2100	1500	1500	1500	2100

- Among the ceramic beryllium neutron multipliers, BeO offers the highest temperature capability if uncoated. If a coating is required, Be₂C is the preferred kernel material.

The three prime material combinations selected are

- SiC surface layer, Li₂O tritium breeder.
- BeO surface layer, LiAlO₂ tritium breeder.
- SiC-coated Be₂C surface layer, LiAlO₂ tritium breeder.

These are shown in Table 7-9 along with the applicable temperature limits from Table 7-8.

The Cascade blanket consists of a high-temperature front zone followed by a tritium-breeding zone. From a thermodynamics perspective, the determination of the relative thicknesses of these zones is a neutronics-optimization problem constrained by the need to achieve an acceptable TBR while maximizing the energy delivered by the blanket and the allowable delivery temperature. The allowable delivery temperature depends on the nuclear heat deposition profile, the velocity flow, and the degree of granule mixing within the blanket.

Figure 7-9 is a schematic representation of the blanket parameters entering the calculation of the outlet temperature distribution. Note that a surface layer of high-temperature material is explicitly shown in Fig. 7-9 moving with a mass flow rate $\dot{m}_{sl}$. This surface layer accommodates the high surface-heat flux and is independently controlled to move at a much higher velocity than the rest of the blanket. The front zone and breeder zone move at rates of $\dot{m}_{fz}$ and $\dot{m}_{bz}$, with a parabolic velocity profile³⁰ given by

$$v(r) = v_0 \left[1 - 0.9(r - r_0)^2 \right] \quad (15)$$

where r is the distance from the center of the chamber, r_0 is the distance to the surface layer, and v_0 is the velocity of the front zone inside the surface layer. The

Fig. 7-9. Schematic representation of parameters in the calculation of blanket temperature distribution.

volumetric neutron heat generation rate profile in the high-temperature front and breeder zones is given by $\dot{q}_{fz}'''(r)$ and $\dot{q}_{bz}'''(r)$, respectively.

The temperature distribution within the blanket can be obtained by examining a differential volume element for conservation of energy. Thus, the energy deposited via neutron absorption is equated to the energy transported to the heat exchanger by

$$v(r)c_p \Delta T(r) \cdot 2\pi r dr = \dot{q}'''(r) \cdot 4\pi r^2 dr \quad (16)$$

where c_p is the specific heat of the material at radius r and $\Delta T(r)$ is the temperature rise of the differential volume element in traversing the reaction chamber. Although the Cascade reaction chamber has a conical geometry, little error is introduced by this spherical approximation. We solved this equation for $\Delta T(r)$ so that the outlet temperature distribution can be calculated for a uniform inlet temperature distribution. The optimum outlet temperature distribution is obtained by calculating v_0 so as to satisfy the temperature limits given in Table 7-9. The temperature rise of the surface layer is treated similarly, except that the velocity profile is constant through the cross section due to full mixing, and in addition, the

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surface layer handles the surface-heat flux. The Cascade blanket can thus deliver energy to the heat exchanger at the distinct temperatures of the surface layer, the front zone, and the breeder zone.

The most useful figure with which to evaluate blanket options at this stage in conceptual design is gross electrical output. With a single power-conversion-system working fluid, the gross electrical output P_3 is given by

$$P_1 = P_{\text{blanket}} \times \eta_1(T_1) \quad (17)$$

where P_{blanket} is the total thermal energy delivered by the blanket to the heat exchanger, and where $\eta_1(T_1)$ is the power conversion system efficiency for a working fluid peak temperature given by

$$T_1 = T_{\text{inlet,bz}} + \frac{\Delta T_{\text{bz}}}{f_{\text{bz}}} - \Delta T_{\text{hx}} \quad (18)$$

where $T_{\text{inlet,bz}}$ is the uniform temperature of breeder material entering the reaction chamber, ΔT_{bz} is the average temperature rise of the breeder zone, f_{bz} is the fraction of energy deposited in the breeder zone; ΔT_{hx} , the temperature drop across heat exchanger heat transfer areas, is assumed to be 300 K. If three power conversion system

working fluids are used, the gross electrical power output P_3 is given by

$$P_3 = P_{\text{blanket}} \times \eta_3 \quad (19)$$

where η_3 , the power-averaged efficiency, is given by

$$\eta_3 = \sum_{i=\text{sl,fz,bz}} \eta(T_{\text{outlet},i} - \Delta T_{\text{hx}}) \times f_i \quad (20)$$

where $T_{\text{outlet},i}$ is the average outlet temperature of zone i .

The thermodynamic performance was uniformly evaluated for three blankets of similar high-temperature zone neutronic thicknesses³¹ and total mass flow rates. The BeO/LiAlO₂ blanket had the highest temperature, efficiency, and gross electrical power production capability. It was thus selected for the limited optimization. It is followed closely by the SiC/Li₂O blanket. The efficiency can be increased by 3 to 6% by using three working fluids, each optimized to an individual zone or surface layer.

We then optimized the BeO/LiAlO₂ blanket by maximizing the energy delivered by the blanket and the blanket outlet temperature. The constraints on this optimization are the need to achieve adequate TBR and the compatibility temperature limits. The

Table 7-10. Thermodynamics of BeO/LiAlO₂ blankets (1500 MW fusion).

BeO zone thickness, m	0.10	0.10	0.20	0.30	0.75	0.75
Blanket power (MW _p)	1670	1670	1790	1790	1840	1840
Total m , m ³ /s	5.3	26	10.4	20	13	25
Surface layer						
Power fraction	0.29	0.29	0.27	0.27	0.26	0.26
$T_{\text{avg outlet}}$ (K)	2300	2300	2300	2300	2300	2300
η_1	0.65	0.65	0.65	0.65	0.65	0.65
T_{inlet} (K)	2100	2100	2100	2100	2100	2100
$\dot{m}$ (m ³ /s)	1.0	1.0	1.0	1.0	1.0	1.0
Front zone						
Power fraction	0.13	0.13	0.31	0.31	0.49	0.49
$T_{\text{avg outlet}}$ (K)	1510	1500	1510	1505	1540	1510
η_1	0.55	0.55	0.55	0.55	0.56	0.55
T_{inlet} (K)	1350	1475	1450	1450	1475	1475
$\dot{m}$ (m ³ /s)	0.6	3.3	3.7	7.4	10.6	21.2
Breeder zone						
Power fraction	0.58	0.58	0.42	0.42	0.24	0.24
$T_{\text{avg outlet}}$ (K)	1240	1450	1020	1240	820	1160
η_1	0.48	0.54	0.40	0.49	0.35	0.47
T_{inlet} (K)	1100	1425	950	1225	700	1100
$\dot{m}$ (m ³ /s)	3.7	21.8	5.8	11.5	1.5	3.0
T_1 (K)	1040	1170	820	1030	900	1050
η_1	0.50	0.54	0.46	0.50	0.48	0.50
P_1 (MW _e)	835	900	823	890	880	920
η_3	0.54	0.57	0.51	0.55	0.53	0.55
P_e (MW _e)	900	950	910	985	970	1010

^aAssumes $\Delta T_{\text{hx}} = 300$ K.

Fig. 7-10. Gross electrical power output vs blanket volume flow rate.

controlling variables in this optimization are the neutronic thicknesses of the two material zones and the mass flow rates. The thermodynamic performance of three BeO/LiAlO₂ blankets of selected BeO zone thicknesses and mass flow rates are shown in Table 7-10. The gross electrical power output figures, P_1 and P_3 are displayed as functions of blanket volume flow rate in Fig. 7-10. The neutronic data for these blankets were obtained from Ref. 32. All blankets achieve a TBR > 1. For a given mass flow rate, the inlet temperature to each zone was adjusted to maintain the peak temperatures below the limits given in Table 7-9.

Thicker BeO zones maximize the energy delivered by the blanket by increasing blanket energy multiplication through increased neutron multiplication. Thicker BeO zones, however, also thermalize neutrons to a greater degree, limiting the penetration of the neutrons into the LiAlO₂ zone and resulting in highly peaked volume heat-generation rates. The result is that the bulk of the energy in the LiAlO₂ breeder zone is deposited in a very narrow region near the BeO zone, resulting in more highly peaked temperature profiles of lower average temperature. This requires high mass flow rates to meet the compatibility temperature limits. Thinner BeO zone blankets have flatter volumetric heat-generation rates and temperature distributions, and thus have higher average outlet temperatures and lower mass flow rate requirements.

For consistency, it is necessary to compare blankets at similar flow rates. Figure 7-10

Fusion power (MW)	1500
Blanket power (MW _t)	1670
Total $\dot{m}$ (m ³ /s)	10.6
Surface layer	
Power fraction	0.29
$T_{avg, outlet}$ (K)	2300
η_1^t	0.65
T_{inlet} (K)	2100
$\dot{m}$ (m ³ /s)	1.0
Peak exit velocity (m/s)	1.9
Front zone	
Power fraction	0.13
$T_{avg, outlet}$ (K)	1705
η_1^t	0.55
T_{inlet} (K)	1435
$\dot{m}$ (m ³ /s)	1.3
Peak exit velocity (m/s)	0.25
Back zone	
Power fraction	0.58
$T_{avg, outlet}$ (K)	1385
η_1^t	0.52
T_{inlet} (K)	1325
$\dot{m}$ (m ³ /s)	8.4
Peak exit velocity (m/s)	0.25
T_1 (K)	1130
η_1	0.53
P_1 (MW _e)	885
η_3	0.56
P_3 (MW _e)	935

^aAssumes $\Delta T_{hx} = 300$ K.

shows that at similar flow rates, the blankets have similar gross electric power output. The gross electric output is a mild function of blanket flow rate above ~ 10 m³/s for thinner BeO blankets. Because thinner BeO blankets have higher efficiencies, minimize beryllium inventory requirements, and allow lower mass flow rates, we selected a 10-cm BeO-zone blanket for the reference design.

The thermodynamic performance of the selected Cascade blanket is presented in Table 7-11. The resultant temperature and

Table 7-11. Thermodynamics of the selected BeO/LiAlO₂ blanket.

Fig. 7-11. (a) Temperature profile for the reference blanket. (b) Velocity profile for the reference blanket.

velocity profiles within the blanket are shown in Figs. 7-11(a) and 7-11(b), respectively. The blanket delivers BeO surface granules at an average temperature of 2300 K, the balance of the BeO front-zone granules at 1505 K, and the LiAlO_2 at 1385 K. This allows a power-conversion-system efficiency of 53% with a single-fluid peak temperature of 1130 K. The total blanket volume flow rate is $11 \text{ m}^3/\text{s}$, delivering 1670 MW of thermal energy to the heat exchanger. The gross power-conversion-system output is 885 MW_e. We believe this can be improved by careful design of the heat exchanger and power conversion system.

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Neutronics of Three Blanket Options for Cascade

A comparison has been made of the neutronic characteristics of three blanket options for the Cascade chamber. All three blankets feature a relatively thin inner region followed by a thicker outer region in which tritium is bred. The first option has an inner region of SiC, which is primarily used to protect the outer, Li_2O region from the x-rays and debris emitted by the target. The second option uses Be_2C in the inner region and LiAlO_2 in the outer region. In this case, the Be_2C not only protects the LiAlO_2 from the short-ranged radiation, but it also provides the neutron multiplication needed to obtain a tritium breeding ratio greater than unity with LiAlO_2 . The neutron multiplication is the result of $(n,2n)$ reactions

with Be. The third option uses BeO in the inner region and LiAlO_2 in the breeding region.

Neutronics Model. The model of the Cascade chamber used in this comparison is described in the article titled: "Blanket Optimization for the Cascade Reactor," later in this section. The inner zone is 10 cm thick, and the packing fraction of the material in this zone was varied to give different effective thicknesses for the three blanket options. The SiC in the first option has an effective thickness of 3.0 cm, which is more than enough to stop the short-ranged x rays and debris.

The Be_2C in the second option has an effective thickness of 4.2 cm. The Be_2C region must be thick enough to provide adequate neutron multiplication. The thickness of the BeO in the third option is also 4.2 cm. This thickness was selected based on the results of the optimization study reported in the article titled: "Blanket Optimization for the Cascade Reactor," later in this section. Note that while the density of Be_2C is less than BeO, 1.9 vs 3.0 g/cm^3 , the Be atom density in the two materials is essentially the same. Therefore, equal thicknesses of Be_2C and BeO should give about the same neutron multiplication.

In all three cases the breeding region is 90 cm thick and has a packing fraction of 50%, which gives an effective thickness of 45 cm. The Li_2O breeder uses natural lithium (7.42% ^6Li and 92.58% ^7Li), while the LiAlO_2 was enriched to 35% in ^6Li .

Results. The results of the neutronics calculations are given in Table 7-12. All three options achieve a tritium breeding ratio greater than unity. Previous studies³³ have

shown that without the inner SiC region, the Li_2O blanket gives a tritium breeding ratio of 1.35. With a 3-cm-thick protective layer of SiC, the breeding ratio is reduced to 1.14. The Be_2C and BeO options give slightly lower breeding ratios. With thicker multipliers, the breeding ratio could be increased if desired. The total energy deposited in the chamber per DT reaction is 18.90, 19.30, and 19.30 MeV for the SiC, Be_2C , and BeO cases, respectively. This total includes the 3.52-MeV alpha particle energy and the neutron energy deposited in the target. The fusion-energy multiplication factors, i.e., the ratio of the total energy deposited to the fusion energy of 17.6 MeV per DT reaction, varies from 1.074 to 1.097. The neutron leakage, which includes leakage through the beam ports, is 0.1 neutron per DT reaction in all three cases.

Conclusions. Based on the comparison of the neutronics characteristics of these three blanket options, there is no clear choice as to which is best. All are viable candidates from a neutronics perspective. The choice of a reference blanket for Cascade can thus be based on other considerations.

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Energy Deposition Profile in the Cascade Granular Blanket

Introduction. The reference granular-blanket concept for Cascade features an inner multiplier layer of BeO followed by a tritium breeding layer of LiAlO_2 . We performed neutronics calculations to determine the energy-deposition profile in the blanket. This information was used by GA Technologies to calculate the temperature profile in the blanket, the blanket flow requirements, and, based on the outlet temperatures and flow rates, the expected energy conversion efficiency for the plant. Three different effective thicknesses for the multiplier—4, 12, and 30 cm—were evaluated. (The effective thickness is the product of the geometric thickness and the 40% packing fraction of the BeO granules.)

Conical Neutronics Model. The basic geometry of the neutronics model used in these calculations is shown in Fig. 7-12. The vertical axis in this figure is a reflecting planar boundary, i.e., a plane of symmetry, and

Inner layer	SiC	Be_2C	BeO
Breeding material	Li_2O	LiAlO_2	LiAlO_2
Tritium breeding ratio	1.14	1.10	1.09
Total energy deposited (MeV) ^a	18.90	19.30	19.30
Energy multiplication factor	1.074	1.097	1.097
Neutron leakage ^a	0.10	0.10	0.10

^aPer DT reaction.

the horizontal axis is the axis of rotation. A source of 14.1-MeV neutrons is distributed throughout a region of compressed DT at the target zone. The ρr of this target zone is 3 g/cm^2 . The conically shaped chamber has a half-angle of 35° . The 0.75-m-radius hole along the axis of rotation represents the beam port. The inner radius of the BeO multiplier along the vertical axis is 3.78 m. At its closest point, the inner surface of the BeO multiplier is 3.10 m from the target. The entire two-layer blanket thickness is 1 m in all cases. The packing fraction in the LiAlO_2 layer is 50%; natural lithium, i.e., 7.42% ^6Li and 92.58% ^7Li , is used in all cases. Outside the LiAlO_2 breeding layer is a 2-cm-thick SiC chamber wall followed by a 2-cm-thick zone representing the SiC-fiber/Al-composite tendons.

Results. The energy-deposition profiles for the 4-, 12-, and 30-cm-thick multipliers are shown in Fig. 7-13. The heating rate is given in watts per gram based on a fusion power of 1500 MW. The peak heat rate in the BeO multiplier occurs at the inner edge of the blanket and is about 15 W/g in all cases. Note that this heating is due only to neutron interactions with the blanket; it does not include the heating due to x rays and debris emitted by the target.

A distinguishing feature of the energy-deposition profiles is the large peak in the LiAlO_2 breeder where it interfaces with the BeO multiplier. The magnitude of this peak increases with increasing BeO thickness. The peak heating rates in the LiAlO_2 are 21, 58, and 99 W/g for the 4-, 12-, and 30-cm-thick multipliers, respectively. This peak results from the absorption of neutrons in $^6\text{Li}(n,\alpha)$ reactions, which release 4.8 MeV. The BeO multiplier, also serves as an effective moderator so that the spectrum of neutrons entering the breeding layer is shifted toward thermal energies; as the thickness of the multiplier increases, the neutron spectrum becomes more thermalized. Since the cross section for the tritium breeding reaction with ^6Li increases with decreasing neutron energy, the thermal neutrons are

Table 7-12. Results of the neutronics calculations.

Cascade Reactor

Fig. 7-12. Conical neutronics model for the Cascade chamber.

Fig. 7-13. Energy deposition profiles in the Cascade blanket.

readily absorbed near the inner edge of the breeder.

Conclusions. Cascade uses a two-layer blanket, in which the inner layer multiplies

and moderates the neutron source, and the outer layer contains the tritium breeding material. With this type of blanket, a peak in the heating rate occurs in the LiAlO_2 breeder at the interface of the two layers. The heating-rate peak becomes a limiting constraint in determining the flow rate and average outlet temperature of the blanket.

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Heat Exchanger Design

Introduction. The transfer of heat from recirculating beds of solid granules has been used industrially for many years in chemical-process industries, heat-recovery systems, and catalytic reactors requiring continuous regeneration of the solid-granule catalysts. The use of recirculating solid granules as the primary heat transport medium in large-scale power-conversion systems presents new challenges in heat exchange to a working fluid.

The heat exchangers used to recover the energy from the stream of solid granules

Table 7-13. Cascade heat-exchanger design summary.

	Immersed-tube fluid bed	Direct-contact fluid bed	Gravity-flow cascade (helium)	Gravity-flow cascade (vacuum)
Helium pressure (MPa)	0.1	5.1	0.1	10 ⁻⁵
Pumping power (MW)	1.6	44	0	0
Fluidization velocity (m/s)	0.45	2.8	NA	NA
Number of tubes	6 200	NA	6 200	6 200
Tube diameter (cm)	2.5	NA	2.5	2.5
Total heat transfer surface area (m ²)	22 000	NA	22 000	84 000
Total frontal surface area (m ²)	510	200	500	500
Active height (m)	1.3	13	1.1	4.0
Helium transfer locks	yes	yes	yes	no
Intermediate heat exchanger	yes	yes with steam cycle, maybe with gas turbine	no	no
Steam cycle	yes	yes	yes	yes
Gas turbine	yes	yes	yes	yes

flowing in a vacuum are a key feature of the Cascade power-conversion system. An obvious approach is to use a simple shell-and-tube heat exchanger with the granules in vacuum flowing by gravity over the tubes with the secondary coolant inside. Because of concerns about low heat-transfer in a vacuum, we also investigated the use of helium gas in the granule interstices. Because Cascade operates in a vacuum (10 Pa), designs that include helium gas in the granule interstices require vacuum locks between the heat exchangers and the reaction chamber. (A possible vacuum-lock system is described in Ref. 34.)

We examined four heat-exchanger concepts:

- Fluid-bed immersed-tube.
- Fluid-bed direct contact.
- Gravity-flow cascade (with helium gas in the interstices).
- Gravity-flow cascade (with vacuum in the interstices).

Heat-transfer models and correlations³⁵ can be applied with reasonable confidence for the case of a moving granule bed flowing over heat exchanger surfaces that are under vacuum, or with interstitial helium gas. Fluid-bed heat-transfer correlations are also available³⁶ for immersed-tube heat-exchange surfaces.

The secondary fluid can be either helium or water, but for this study, we evaluated the heat exchangers for uniform secondary steam-side conditions of conventional, superheated steam turbine cycles of modern fossil-fired power plants.³⁷ All concepts use double-walled tubes to minimize tritium permeation into the steam and to reduce

the possibility of water leaks into the blanket. We chose granule inlet and outlet temperatures to the heat exchanger as 1200 K and 800 K. The resulting designs are summarized in Table 7-13 and discussed below.

Fluid-Bed Immersed-Tube Exchanger

Fluid-bed immersed-tube heat exchangers offer relatively high bed-to-surface heat transfer coefficients, very good bed mixing, and enhanced granule transport from fluidization. Some disadvantages of this type of heat exchanger are tube abrasion, vibration, solids dusting, and the need for a compressor system to circulate the fluidizing gas.

To minimize the granule-transport height differences for the Cascade power-conversion system, we evaluated a concept for a horizontally staged counterflow, fluid-bed immersed-tube heat exchanger, illustrated in Fig. 7-14. In this design, the hot granules enter the first stage, where they mix with the stage granule inventory that is fluidized around triangularly stacked tubes. The tubes contain either water or a mixture of water and steam. Triangular spacing is used to minimize bubble growth in the granule bed and granule inventory within each stage. Triangular spacing of the tubes is 3.25 cm for 2.5-cm-diam tubes. If we assume that the water or water/steam combination makes two passes per tube within each stage, and there are six stages, the required heat-transfer area is 22 000 m². With an active heat-exchange height of 1.3 m, the total frontal area is 510 m². The hot granules move from stage to stage through the granule cascade orifices. The helium gas enters through the bottom of each stage, blows through the granules, and exits

Fig. 7-14. Fluid-bed, immersed-tube, heat exchanger

Fig. 7-15. Fluid-bed, direct contact, counterflow heat exchanger.

horizontally. The small heat capacity and flow rate of the low-pressure helium fluidization flow does not significantly degrade thermal performance.

We estimate the fluidization pressure drop per stage as equal to the static head of the granules. The total pressure drop is 0.17 MPa. For fluidization pumping-power estimates, we doubled the bed fluidization pressure drop to account for ducting and gas-distributor losses.

Fluid-Bed Direct-Contact Heat Exchanger. The high effective surface area of a fluid-bed direct-contact heat exchanger is ideal for transferring heat from the granules to a fluid. However, the heat capacity of the gas requires high pressure and large frontal area to keep gas velocity low enough to maintain a stable fluidized bed.

A conceptual fluid-bed direct-contact staged-counterflow heat exchanger is shown in Fig. 7-15. Granules enter through the top of the apparatus and flow by gravity from stage to stage against the fluidization pressure gradient. A horizontal configuration may also be possible, but to achieve a countercurrent flow we must devise some means to move the particles horizontally against the pressure gradient between stages.

In a fluid-bed direct-contact heat exchanger, the high-pressure, high-flow rate helium must flow countercurrent to the granules to use the high outlet temperature. Hot helium leaving a fluid-bed contactor can be used for either a helium-heated steam boiler or a closed-cycle gas turbine.

The helium flow rate required to transfer heat is large (~1430 kg/s), and results

through the top. The helium recirculates to the next stage through tubes that run outside the heat exchanger.

The helium gas flows parallel with the granules, thereby using the stage-to-stage pressure difference to convey the particles

Fig. 7-16. Gravity-flow cascade heat exchanger.

in a total loop pumping power of 44 MW, including losses external to the direct contactor.

The fluid-bed direct-contact heat exchanger offers high effective surface area, provides gas near thermal equilibrium at the solids inlet temperature, and the secondary heat exchanger can be conventional single-wall tubes. The disadvantages of this design are: high pressure is needed to provide adequate heat capacity; gas flow rate limited by fluidization velocity; granule attrition; potential dust carryover; and high pumping power for gas recirculation.

Gravity-Flow Cascade Heat Exchanger. Gravity-induced flow of granules over the heat exchanger provides mixing similar to fluidization (Fig. 7-16). Helium in the interstices acts only to enhance heat transfer, and does not serve a fluidizing function. The increase of the surface-to-wall heat-transfer coefficient depends on the average contact time. Flow-splitter-type distributors can provide relatively short residence times for the hot granules while maintaining uniform flow over all surfaces. To improve mixing, we selected a staggered tube, horizontal array similar to that used for the immersed-tube fluid-bed concept. In this case, however, horizontal staging is not appropriate, and the once-through boiler tubes pass back and forth, ascending countercurrent to the granule flow. This heat-

exchanger arrangement has the advantage that it can operate with either interstitial gas or vacuum. (Characteristics calculated for both gas and vacuum cases are shown in Table 7-13.) Other advantages are good bed mixing and a relatively high surface heat-transfer coefficient.

Disadvantages for the gravity-flow cascade heat exchanger are difficult uniform granular distribution over a large frontal area and difficult removal of granule plugs.

Conclusions. The most significant conclusion of this study is that all of the heat-exchanger concepts evaluated are feasible and can result in practical designs. Fluid-bed concepts are self-distributing across their frontal area, whereas the gravity-flow concepts require mechanical distribution to minimize granule inventory. The direct-contact heat exchanger requires high helium pressure and has higher pumping-power losses. Fluid-bed heat transfer areas for the immersed tube and gravity flow with helium designs are not significantly different, because of the good agitation and short particle-contact time offered by the triangular pitch and flow baffles in the gravity-flow heat exchanger. Vacuum heat exchange surface area is approximately three times that of heat exchange area in a helium medium. However, the latter needs an active transfer lock system through which to pass granules in and out of the vacuum that is

required in the reaction chamber. Because of the simplicity of a system without vacuum locks, we selected the gravity-flow Cascade heat exchanger with granules flowing in vacuum as the reference design.

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Granule Fabrication Processes

A key feature of Cascade is the granular blanket. Of the many blanket material options open to Cascade, fabrication of Li_2O granules was felt to offer the greatest challenge. We explored available methods for initial Li_2O granule fabrication.

We identified three cost-effective processes for fabricating Li_2O granules: the VSM drop-melt furnace process,³⁸ which is based on melting and spheroidizing irregularly shaped Li_2O feed granules; the LiOH process,³⁹ which spheroidizes liquefied LiOH and uses GA Technologies' sphere-forming procedures; and the Li_2CO_3 sol-gel process,⁴⁰ used for making spherical fuel particles for the high-temperature gas-cooled reactor (HTGR). Each process is described below.

VSM Drop-Melt Furnace Process. This process requires high-purity Li_2O kernels as feed material to the melt-spheroidizing furnace. Because Li_2O readily absorbs H_2O and CO_2 , we will need a process step to remove H_2O and CO_2 impurities before spheroidizing in the VSM furnace. The major steps of the VSM process are:

1. H_2O and CO_2 are removed from the Li_2O feed kernels by heating in a vacuum furnace to 570 K at a vacuum of 0.3 Pa.
2. The irregularly shaped feed kernels are converted to spherical granules by melting as they drop through the hot zone of a VSM furnace operating at 1900 to 2000 K. The spheres solidify by cooling below the Li_2O melting point (1700 K) in the cold zone at the bottom of the furnace.
3. The spherical Li_2O granules are screened to specific sizes by a sonic sifter.
4. Oversized granules are recycled by crushing and converted to smaller granules in a pelletizer.
5. Undersize granules are recycled directly to the pelletizer.

6. After crushing, the recycled material is combined with the feed stream to the vacuum furnaces.

The VSM drop-melt process is also applicable to other materials, such as LiAlO_2 .

LiOH Process. This process requires high-purity LiOH as the feed material to the granule forming equipment. H_2O and CO_2 impurities are removed, but nonvolatile and metallic impurities remain in the product granules.

The major steps for the LiOH process are as follows:

1. The LiOH is melted by heating to about 870 K (melting point ~ 730 K).
2. The liquid LiOH is transferred to the spheroidizing feed tank, also at a temperature of 870 K.
3. The liquid LiOH is pumped at a constant rate to nozzles at the top of a cooling tower.
4. The spherical LiOH droplets are solidified by a countercurrent flow of argon gas in the cooling tower.
5. The solidified granules are collected at the bottom of the cooling tower and transferred to batch vacuum furnaces.
6. The spherical LiOH granules are converted to Li_2O in the vacuum furnace (at 370 K and 0.3 Pa).
7. The Li_2O granules are sintered to high density by heating in a sintering furnace to ~ 1300 K in an argon atmosphere.

The LiOH process appears to be useful only for Li_2O , because it requires a relatively low-melting-point precursor (such as LiOH).

Li_2CO_3 Process. The Li_2CO_3 process can use commercial-grade LiOH as the raw material because water-soluble impurities in the LiOH would be removed during the gelled-sphere washing step of the process. Volatile impurities would also be removed during the conversion step to Li_2O as in the LiOH process.

The Li_2CO_3 process consists of the following major steps:

1. LiOH dissolved in water reacts with CO_2 to form a Li_2CO_3 sol.
2. The sol is pumped at constant rate to spheroidizing nozzles at the top of the gelling column.
3. The Li_2CO_3 granules from the spheroidizing nozzles are gelled in a CO_2 gas and collected in an H_2/CO_2 solution.

4. The granules are washed in a countercurrent wash column with H_2O or aqueous carbonate solution.
5. Washed granules are dried in a continuous-drying oven at up to 420 K.
6. Dried Li_2CO_3 granules are converted to Li_2O in a batch vacuum furnace at about 0.3 Pa and 870 K.
7. The Li_2O granules are then sintered to about 90 to 95% of theoretical density in a vertical tube furnace at 1270 K in an argon atmosphere.

The sol-gel process is highly dependent on the chemistry of the gel production and gellation steps. A sol-gel process has been successfully developed at GA Technologies for production of $LiAlO_2$ and $LiAl_2O_3$.

Fabrication Cost Estimate. A cost-estimate summary for granule fabrication is shown in Table 7-14. The high cost of raw materials (1.4- to 2.3-mm, high-purity Li_2O feed kernels) for the VSM process makes it noncompetitive with the other two processes for fabrication of the granules for the reaction chamber. Although the LiOH process also requires high-purity feed material, it offers the opportunity for some impurity removal. The LiOH is the major cost item for granule manufacture, but it costs much less than the Li_2O for the VSM process. The Li_2CO_3 process uses commercial-grade LiOH, which means that the raw-material costs for the Li_2CO_3 process should be the lowest of the three.

The labor, energy, and equipment (depreciation) costs presented in Table 7-14 are based on actual GA Technologies operating experience with the VSM furnace in the manufacture of HTGR granules, and sol-gel processes used to manufacture HTGR alumina and $LiAlO_2$ granules. The raw-material cost of the Li_2O is based on estimates of commercial-scale purification and vacuum-dewatering of LiOH obtained from Ref. 41 and other materials costs from Ref. 42. We obtained the chemical plant equipment costs by following standard industry practice,^{4,5} and assuming a 1-MT/day plant. The fabrication costs shown in Table 7-14 should be used only as a comparison; these would change if the cost of high-purity lithium compounds were reduced by large production demands. Labor costs for all three processes would drop, based on normal learning curves for new production facilities and cost-sharing of

	Costs in \$/kg Li_2O		
	VSM process	LiOH process	Li_2CO_3 process
Labor	13	15	16
Raw materials	150	33	17
Energy at \$0.07/kWh	0.32	0.69	0.93
Depreciation	1.6	1.6	2.3
Total	170	50	37

personnel with other fusion-related tasks at the power plant. The processes are nearly automated to current U.S. practice and no additional major cost reductions are foreseen through further automation.

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Chamber Clearing Time

With each fusion pulse, some material is vaporized from the surface of the granular bed. This vaporized material must be removed from the chamber before the next pulse can be fired. In general, because of the large chamber size and short time between pulses, it is impractical to use mechanical pumps to remove the vaporized material from the chamber. Our studies indicate that the recondensation of the vaporized material on the surfaces inside the chamber may determine the maximum permitted pulse rate.

There have been many studies⁴⁴⁻⁵² of the phenomena important to recondensation in reactor cavities containing liquid-metal sprays, jets, or curtains, but the phenomena are still not well understood and relevant data are sparse. The Electric Power Research Institute (EPRI)^{51,52} has recently undertaken an experimental program to obtain some of the necessary data on fragmentation of isochorically heated liquids. Among the computer codes that attempt to model some of the recondensation phenomena is CONRAD, being developed as a successor to the widely available MFFIRE⁵³ code. We used the CONRAD code to estimate the recondensation time in the Cascade chamber.

In this article, we separate the problem of determining the chamber clearing time into four parts. First, we estimate the amount of mass initially vaporized from the granular blanket and the amount of energy in that mass. Second, we estimate the

Table 7-14. Li_2O granule fabrication cost summary.

temperature profile in the remaining material. Third, we estimate how rapidly the vapor flows into the granular blanket. And fourth, we use the CONRAD code to estimate the energy and mass transfer between the vapor and the wall. We discuss these parts, in order, below.

Vaporization of Material. When the laser pulse travels from the final focusing lens to the target, the chamber pressure should be no more than about 0.1 Torr. At this time, we expect that all gases will have recondensed, except for the residual tritium, deuterium, and helium from the fuel pellet, and the tritium and helium born in the blanket. The implosion and burn of the target, and the emission of x-rays and neutrons and their subsequent absorption in surrounding material, all take place in less than 1 μ s. Approximately one-third of the 300-MJ pellet yield is in the form of x-ray and debris energy, while the remaining two-thirds is in neutron energy.

We used the BUCKL code⁵⁴ to estimate the deposition of x-ray energy in the first layer of BeO granules. The energy-deposition profile is

$$q'''(x) = \phi [A \exp(-\mu_1 x) + B \exp(-\mu_2 x)] \quad (\text{kJ/kg}) \quad (21)$$

for the first 10 μ m, and thereafter,

$$q'''(x) = \phi C x^{\mu_3} \quad (\text{kJ/kg}) \quad (22)$$

where x is the distance (in m) from the granule surface, ϕ is the surface energy flux (in kJ/m^2), and A , B , C , μ_1 , μ_2 , and μ_3 are constants calculated for each material. The cohesive energy for BeO is calculated as the difference between the enthalpy of BeO vapor at the vaporization temperature (4200 K) and the enthalpy of the solid BeO at the

initial temperature (2300 K). This is about 23×10^6 J/kg. We assume that all BeO with a specific energy greater than 23×10^6 J/kg is vaporized. Our calculations showed that about 3 μ m of material on each surface granule is vaporized. The total granular material vaporized is about 1 kg, which is much more than the few milligrams of DT fuel injected into the chamber or the 0.3 g of background gas at 0.1 Torr. The vaporized BeO expands into the chamber and meets the debris. We estimated that the BeO has much more momentum than the debris, which leads us to assume that the kinetic energy and mass of the debris are absorbed in the BeO vapor. Thus the 1 kg of BeO vapor initially contains about 100 MJ of energy.

Temperature Profile. We derived the ambient temperature of the granules from neutronics and granular-flow calculations discussed in the articles "High-Temperature Blanket Design Options" and "Energy Deposition Profile in the Cascade Granular Blanket," earlier in this section. We conclude that the inner 2 cm of the granular blanket will be at 2300 K and that the temperature will fall to 1300 K at the outside edge.

Vapor Flow Through the Granular Blanket. The vaporized granular material will expand toward the center of the chamber, rebound, and eventually reach the blanket again. During this time, the energy exchange is between the vapor and the inside surface of the granular blanket. However, when the vapor reaches the blanket again and starts to penetrate it, the available surface area for energy and mass exchange increases considerably. Because of this increase, it is important to estimate the rate at which the expanding vapor penetrates the blanket. We used a simple, one-dimensional geometry (shown in Fig. 7-17) and the following assumptions to estimate the penetration rate:

- The chamber inside the blanket is filled with hot vapor at uniform pressure and temperature, the blanket porosity is ϵ , and the initial pore pressure p_0 is zero.
- The vapor penetrates the blanket without condensing and acts as a perfect gas.
- The blanket is a one-dimensional column with a sealed end.
- The driving pressure and temperature in the volume at the edge of the blanket are constant.

Fig. 7-17. Simplified geometry for calculating vapor flow into granular blanket. The granular blanket has a thickness L , a porosity ϵ , and an initial pore pressure p_0 of zero.

- All material properties are constant, and there is no heat transfer (isothermal flow).

We calculated the driving pressure and temperature using the MIXERG⁵⁵ code. The vapor pressure is about 10^5 Pa at a temperature of about 11 600 K (1 eV).

Permeability k is determined from the equation⁵⁶

$$k = \frac{\beta \epsilon^2}{\lambda (1 - \epsilon)} d^2, \quad (23)$$

where $\beta = \pi/6$ for spheres, λ is an empirical constant, ϵ is the porosity, and d is the sphere diameter. For a nominal value of λ , a porosity of 0.5, and 1-mm-diam spheres,

$$k = 1.5 \times 10^{-9} \text{ m}^2 = 1500 \text{ darcy}. \quad (24)$$

From Ref. 57 we can calculate the viscosity of BeO vapor at 11 600 K as

$$\mu = 1.41 \times 10^{-4} \text{ Pa} \cdot \text{s}. \quad (25)$$

In the method described in Ref. 58, the governing equation, which combines momentum, energy, continuity, and the perfect-gas equation of state, is

$$\frac{\partial}{\partial x} (\rho u) + \epsilon \frac{\partial \rho}{\partial t} = 0, \quad (26)$$

where x is the axial dimension of the porous column, ρ is gas density, u is velocity (equal to the volume flow rate per unit area including voids and solids), ϵ is porosity, and t is time. Using Darcy's law and the perfect gas law with isothermal flow,

$$u = -\frac{k}{\mu} \frac{\partial p}{\partial x} \quad (27)$$

and

$$p = \rho RT, \text{ with } T \text{ constant}, \quad (28)$$

then

$$\frac{\partial}{\partial x} \left(p \frac{\partial p}{\partial x} \right) = \frac{\epsilon \mu}{k} \frac{\partial p}{\partial t}. \quad (29)$$

With the dimensionless variables

$$\bar{p} = \frac{p - p_0}{p_1 - p_0}, \quad (30)$$

where p_1 is the driving pressure and p_0 is the initial pore pressure in the blanket,

$$\bar{X} = \frac{x}{L}, \quad (31)$$

where L is the blanket thickness,

$$N = \frac{p_1}{p_0}, \quad (32)$$

and

$$\tau = \frac{k(p_1 - p_0)t}{\epsilon \mu L^2}, \quad (33)$$

we obtain

$$\frac{\partial \bar{p}}{\partial \tau} = \frac{1}{2} \frac{\partial^2}{\partial \bar{X}^2} \left(\bar{p}^2 + \frac{2\bar{p}}{N-1} \right). \quad (34)$$

The pressure at the end of a sealed column and the pressure distribution in a sealed column are shown in Figs. 7-18 and 7-19,

Fig. 7-18. Pressure at the end of a sealed column.

Fig. 7-19. Pressure distribution in a sealed column.

respectively. Values for $N = \infty$ are close to those for $N = 45$. Figure 7-18 shows that the vapor first reaches the outside edge of a 1-m-thick blanket at $\tau = 0.38$. Using Eq. (33) thus results in time $t = 160$ ms.

From Fig. 7-19 we see that the time for the pressure in the entire blanket to reach 80% of the chamber pressure is $\tau = \sim 0.8$ or $t = 330$ ms. Thus the available surface area will be much larger than the surface of the first layer of granules. Note that if the vapor filled the blanket, more than half the vapor would contact the granules.

Energy and Mass Exchange Between the Vapor and the BeO Blanket. The CONRAD code 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 contained in any zone. The maximum permitted mass change in any zone is an input variable. Non-ideal equations of state for use in the CONRAD code are obtained in tabular form from the MIXERG code, which calculates ionization level, opacity, etc., to create such tables. The CONRAD code uses a one-dimensional, finite-difference method

to determine the time-dependent temperature profiles in the solid or liquid layer.

Heat transfer across the interface between the vapor and liquid is due to mass transfer (either condensation or vaporization), thermal radiation from hot vapor, and x rays. All of these deposit energy either on the interface, or within the solid, both of which have a deposition profile of a multigroup sum of exponentials. Vapor within the 1-mfp-thick boundary region exchanges heat and mass with the layer of solid or liquid material. 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 only valid within the hydrodynamic approximation, the CONRAD code stops when the thickness of the boundary layer equals the thickness of the entire vapor region.

No one-dimensional code can yet model in detail the granular blanket or the vapor flow into it. Soon after the initial vaporization of BeO by the x rays, we use the simple geometry shown in Fig. 7-20. We assume the chamber to be spherical and to be divided into three regions: a solid-BeO spherical shell (between R_1 and R_2), a high-density BeO-vapor region (between R_2 and R_1), and low-density BeO-vapor region (between 0 and R_1). The outer boundary is made immovable; there should be no heat transfer across this immovable surface, because it will represent the center of each granule. The heat-transfer condition is not yet available in the CONRAD code. Instead, we held the temperature at the outer boundary constant. Containing the 100 MJ of energy would have raised the temperature of the solid wall while decreasing the temperature gradient across it. However, if all the energy is deposited in the 170 kg of BeO used in our calculations, the wall temperature will be raised only 280 K, or about 10%. From the x-ray deposition calculations above, we chose the mass of BeO vapor to be 1 kg, and the energy in the BeO vapor to be 100 MJ (the sum of the x-ray and debris energy). Initially, the high-density BeO vapor was spread over a few centimeters to make equal-mass zoning feasible. The properties of BeO used in the calculation are shown in Table 7-15.

Fig. 7-20. Initial problem geometry for the CONRAD calculations.

Table 7-15. Material properties of BeO.

	Reference No.	BeO
Melt temperature, T_m (K)	59	2800
Vaporization temperature, T_v (K)	59	4200
Heat of fusion, H_f (J/kg)	60	2.6×10^6
Heat of vaporization, H_v (J/kg)	61	19.8×10^6
Density, ρ (kg/m ³)	59	3000
Specific heat at 1270 K, c_p (J/kg·K)	62	2100
Thermal conductivity at 1270 K, k (W/m·K)	62	28
Young's modulus, E (GPa)	62	325
Tensile strength, σ_t (MPa)	62	330

We have run a series of calculations in which we vary the surface area in the solid BeO to represent various portions of the granule blanket accessible to the hot gas. The ideal area multiplier X_A is

$$X_A = 3n \left(\frac{R_3 - R}{r} \right)^2 \text{ for } R_2 < R < R_3$$

$$\text{and } r = R_3 - R_2 = 0.5 \text{ mm} \quad (35)$$

where n = number of granule layers included, R is an arbitrary location between R_2 and R_3 , r = granule radius, R_3 = 3.1005 m, and where we assume the bed porosity to be 0.5. For a correct thermal gradient, the thickness of the BeO shell must equal the granule radius. However, we must decrease the area as a function of R to represent correctly the conductivity into each granule. This area should be used to calculate heat transfer from the gas to the solid BeO surface and from the surface at R_2 into the bulk. Up to 50% of the gas will be in the bed at any given time, and the multiplier at the surface R_2 is $3n$.

If there is no heat transfer across the outer boundary at R_3 , the solid BeO will heat up. To represent the total heat capacity, we must change the mass in the BeO each time we change the area in the BeO. Thus, the mass multiplier is

$$X_M = n \left(\frac{R_3 - R}{r} \right)^3 \text{ for } R_2 < R < R_3 \quad (36)$$

The mass multiplier should be used to calculate the temperature changes in the solid BeO. We must have $R_2 = 3.1$ m to make the total energy density correct. The initial temperatures chosen for R_2 and R_3 depended on the number of granule layers accessible to the gas. Table 7-16 contains the values that would apply if the gas were in close contact with 1, 10, 100, or 1000 layers of 1-mm granules (1000 layers represents the entire blanket). We calculated the temperatures by averaging the granular blanket temperatures weighted by area. For one layer of granules, we assume that the surface of each granule is at 4200 K. In all other cases, we assume that most of the bed is at a temperature determined by the time-averaged power density for the

n	Initial T at R_2 (K)	Initial T at R_3 (K)
1	4200	2300
10	2300	2300
100	1750	1750
1000	1300	1300

location, and each granule has no internal thermal gradient. For one layer, the ideal initial temperature gradient would be an exponential between R_2 and R_3 .

Making the area and mass multipliers functions of radius proved difficult with the CONRAD code. Thus, for the first runs, we chose constants that bound the problem. We ran two cases initially, both representing only a single layer of granules ($n = 1$). In the first case we chose a constant multiplier of $M = 3$ and a thickness $R_3 - R_2 = 0.17$ mm. In these CONRAD calculations, the multiplier is used to multiply the calculated surface heat flux (composed of condensation/vaporization and radiation flow) and the thermal diffusivity in the solid BeO. This is the correct multiplier for the total surface of one layer of granules and the correct thickness for the total heat capacity of a single layer of granules. The multiplier and thickness represent an upper limit to the heat transfer rate into the granules because for the real granules the thermal gradient will be only one-third as large, and the surface area through which the heat flows decreases as the granule centers are approached.

In the second case, we used a multiplier of $M = 1$ and a thickness of 0.5 mm. $M = 1$ is the average value of the expression in Eq. (35) for one layer of granules, and because the thickness is equal to the granule radius, the heat flow is calculated correctly. However, in this case, the mass and, thus, the total heat capacity is a factor of 3 too large. Thus in the first case, the heat flux in the solid BeO is a factor of 9 larger than in the second case, but the total heat capacity is correct. In the second case, the heat flux is accurate, but the heat capacity is a factor of 3 too high.

The results of our calculations are shown in Figs. 7-21 and 7-22. To maintain the BeO pressure below the 0.1-Torr pressure of the tritium, deuterium, and helium, the BeO density must be below 10^{15} atoms/cm³. In

Table 7-16. Initial temperatures as a function of the number n of 1-mm granule layers exposed to the gas.

Fig. 7-21. Average density for the first case (see text).

Fig. 7-22. Average density for the second case (see text).

neither case did the CONRAD code allow the calculation to proceed to 10^{15} atoms/cm 3 , and we must extrapolate to this value. Within the numerical fluctuations in the calculations, the times for recondensing to this value in both cases appear to be about 25 ms.

Conclusion. The combination of the CONRAD code results and the estimate of gas flow into the granular blanket shows that the recondensation time is likely to be considerably less than the desired 200-ms interpulse time. However, there are a number of calculational uncertainties that could alter the results. The CONRAD code contains a very simple condensation model: if a vapor molecule is within one mean free path of the wall and is moving toward the wall, it is assumed to be condensed. No account is taken of the effect of the presence of noncondensable gases, sticking coefficients not equal to unity, dissociation and recombination of the BeO molecules, and any effects of surface chemistry. All these effects could make the condensation time

much longer than calculated here. On the other hand, the real geometry and possible contact with the very large surface area of the granular blanket would accelerate condensation. We must further investigate all potential effects to provide a definitive answer.

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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. The 1-GW $_t$ design⁶³ uses a liquid lead/lithium eutectic (Li $_{17}$ Pb $_{83}$) to transfer heat from the fusion chamber to the steam generators, to breed tritium for fuel, and to protect the bell jar, which constitutes the structural element of the fusion chamber.

Work on the Pulse*Star concept this year has concentrated on three areas: polonium issues, tritium containment, and the use of lithium in place of LiPb. We discuss these areas in the following three articles.

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Polonium Issues

The large quantities of lead in the Li $_{17}$ Pb $_{83}$ (LiPb) coolant of Pulse*Star can lead to production of significant quantities of ^{210}Po —an alpha emitter with a 138-day half-life—by neutron activation. This half-life is long enough that the ^{210}Po would not decay away before human contact, in the hypothetical case of a polonium release from a reactor, and short enough that the alpha-emission rate from a given mass of ^{210}Po is much higher than from an equivalent mass of longer half-life plutonium. The allowable concentration of ^{210}Po in air is extremely low: 7×10^{-12} Ci/m 3 = 1.6×10^{-15} g/m 3 , compared to 10^{-12} Ci/m 3 = 1.6×10^{-11} g/m 3 for ^{239}Pu .⁶⁴

Fig. 7-23. Production of polonium from lead and bismuth.

Polonium can be produced from two pathways, as shown in Fig. 7-23. Successive absorption of two neutrons by ^{208}Pb followed by two beta decays can produce ^{210}Po , but the 3.3-h half-life of the intermediate ^{209}Pb minimizes this contribution. The fast decay of ^{209}Pb , however, produces ^{209}Bi , the only stable isotope of bismuth. This activation-produced bismuth and the bismuth impurity normally found in lead can absorb a neutron and then decay to ^{210}Po : this is the primary source of polonium in fusion reactors by perhaps five orders of magnitude.

Comparison of Pulse*Star and MARS.

Both the Pulse*Star and the mirror-advanced-reactor-study (MARS) reactors⁶⁵ use LiPb coolant. The ^{210}Po inventory in Pulse*Star is shown in Fig. 7-24 as a function of initial bismuth impurity level. The 41.7-wppm bismuth-in-lead data point corresponds to the initial bismuth level in MARS. According to Ref. 65, the ^{210}Po production rate after 30 y of operation would be 0.027 Ci/s, with a total production of 16.8 MCi. The production rate at the end of the reactor's 30-y life would equal the decay rate for an inventory of 0.47 MCi, which is about 20% of the calculated Pulse*Star inventory. The possible reasons for the different polonium levels include:

1. Possible error in either the activation code or cross section libraries at LLNL or the University of Wisconsin (which performed the MARS calculation).
2. In MARS, the LiPb is in the balance of plant 86.4% of the time and in the high-flux region for 13.6% of the time. In Pulse*Star, the LiPb is always in either

Fig. 7-24. Generation of ^{210}Po in Pulse*Star as a function of initial bismuth content. Shutdown follows 30-y operation at 70% capacity factor and 1 CW₀.

the high-flux region or the low-flux pool region.

3. In MARS, there are several competitors for thermal neutrons. The ^{208}Pb reaction rate is reduced because of the increased reaction rates of ^6Li (90% enriched in MARS, natural 7% enriched in Pulse*Star) and steel (in the MARS reflector).

The magnetic fusion and ICF projects will run a simple, one-dimensional test problem for the Sixth Topical Meeting on

Fusion, March, 1985, to evaluate the first and second possibilities. The third possibility is design dependent; however, we note that enrichment to 90% ^6Li is expensive, and that absorption of neutrons in steel creates additional nuclear energy at the expense of increased activation of the steel's constituents.

Reduction of Polonium Production by Tailoring the Lead. Different lead ore bodies have different isotopic abundances, which result from the production of ^{206}Pb and ^{207}Pb from the decay of uranium, and ^{208}Pb from the decay of thorium. Inspection of 11 ore bodies from around the world⁶⁶ showed that the lowest ^{208}Pb abundance was 51.25% (in a galena sample from Joplin, Mo.) compared to 52.30% for radio-genically uncontaminated lead. Thus, the ^{208}Pb abundance variation is too low to significantly reduce ^{210}Po production by selection of an ore body with low ^{208}Pb . However, isotopic tailoring of the lead to reduce the bismuth impurity level could significantly limit polonium production.

The bismuth impurity in lead ranges from a low of 41.7 wppm for the purest-available lead to 1982 wppm for standard lead. The impact of the initial bismuth content is shown in Fig. 7-24. Using the purest-available lead reduces the final ^{210}Po inventory by about a factor of 3, but the ^{210}Po inventory is still very large. Removing all the remaining bismuth before startup provides almost no additional improvement; however, continuous removal of bismuth to very low levels could provide a great improvement.

Polonium Vapor Pressure and Chemistry. The vapor pressure P_0 associated with pure polonium liquid was given in Ref. 67 as

$$\log P_0 (\text{Torr}) = 7.2345 \pm 0.0068 - \frac{5377.8 \pm 6.7}{T} \quad (37)$$

in the 438 to 745°C range. Although Ref. 67 states that the temperature units are in degrees celsius, the correct units are kelvin, as stated in Ref. 68, where the equation was originally published. At 500°C (227 K), Eq. (37) gives $P_0 = 2$ Torr, which is in agreement with the MARS report.⁶⁵ Such a high vapor pressure would allow removal of polonium in the vacuum system.

The vapor pressure above a polonium-in-lead solution is

$$P = P_0 \gamma n \quad (38)$$

where γ is the activity coefficient and n is the atom fraction of polonium in the liquid lead. The activity coefficient for polonium in lead is unknown, but the MARS report gives the relation⁶⁹

$$\log \gamma_{\text{Po-Bi}} = 1.1176 - \frac{2728.3}{T} \quad (39)$$

for polonium in bismuth, where T is in kelvin. At 500°C, the activity coefficient of polonium in bismuth $\gamma_{\text{Po-Bi}}$ equals 0.0039, which reduces the polonium vapor pressure by a factor of 250 from ideal. A similar situation may exist for polonium in lead.

Polonium vaporization will probably be in the form of H_2Po , where the hydrogen is provided by the unburned deuterium and the tritium. H_2Po , a chemical analog to H_2S , is a gas at room temperature, suggesting a potential containment difficulty. Additional information on polonium chemistry is contained in the German⁶⁷ and Soviet⁷⁰⁻⁷³ literature.

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Tritium Containment

Because tritium is not very soluble in the $\text{Li}_{17}\text{Pb}_{83}$ eutectic we chose as the coolant for the Pulse*Star reactor, concern was expressed that tritium will diffuse into the Pulse*Star steam supply and escape at unacceptable rates into the environment. While acceptable tritium leakage rates for a fusion reactor are still subject to much debate, values of 10 to 500 Ci/d (corresponding to 1 to 50 mg/d tritium) have been widely used for conservative design goals. Because ~1.4 kg of tritium are injected into the chamber each day, very stringent measures might be required to reduce tritium leakage to an acceptable level.

We assume the tritium vapor in the Pulse*Star chamber to be in equilibrium with the tritium dissolved in the $\text{Li}_{17}\text{Pb}_{83}$. Some $\text{Li}_{17}\text{Pb}_{83}$ vaporizes with each fusion pulse and mixes well with the tritium, deuterium, and helium gases from the target.

This mixture is in close contact with the fine $\text{Li}_{17}\text{Pb}_{83}$ droplets in the spray zone as the $\text{Li}_{17}\text{Pb}_{83}$ recondenses. Tritium and helium exist both in the vapor (from the target) and in the $\text{Li}_{17}\text{Pb}_{83}$ (from the neutron reactions). Each $\text{Li}_{17}\text{Pb}_{83}$ droplet experiences many fusion pulses as it falls to the bottom of the chamber. In this case, the equilibrium atomic concentration of tritium in the $\text{Li}_{17}\text{Pb}_{83}$ is $C_T = P_T^{1/2}/K_5$, where P_T is the tritium pressure in the chamber and K_5 varies between 1.0 and 1.3 Torr^{1/2}/appm for 500 and 400°C $\text{Li}_{17}\text{Pb}_{83}$ (Ref. 74), respectively. Last year, we established an upper tritium pressure limit in the chamber of 1.7×10^{-3} Torr to minimize the tritium release in the event of a $\text{Li}_{17}\text{Pb}_{83}$ spill. If we take $P_T = 1.7 \times 10^{-3}$, we find $C_T = 0.034$ appm in the $\text{Li}_{17}\text{Pb}_{83}$, or a tritium inventory of ~ 20 g for the whole $\text{Li}_{17}\text{Pb}_{83}$ inventory of 3.2×10^{10} g. The $\text{Li}_{17}\text{Pb}_{83}$ inventory circulates through the heat exchangers over 500 times per day.

Without a tritium-extraction process in the $\text{Li}_{17}\text{Pb}_{83}$ circulation loop, the tritium driving pressure in the $\text{Li}_{17}\text{Pb}_{83}$ is 1.7×10^{-3} Torr. The permeation rate of tritium through a 4.8-mm-thick iron heat-exchanger tube wall has been calculated on the DIFFUSE code at Sandia National Laboratory, Livermore.⁷⁵ The results are shown in Fig. 7-25 as a function of wall temperature T and the temperature drop across the wall ΔT . The temperature of the tube wall varies from $\sim 500^\circ\text{C}$ at the top to $\sim 400^\circ\text{C}$ at the bottom. The eight Pulse*Star heat exchangers have a total area of 14 700 m². Using a weighted-average value of 3.1 mg/m²·d from Fig. 7-25, we find a total transfer of ~ 46 g/d (neglecting back diffusion). If we assume that all tritium entering the steam cycle escapes to the environment, then the leakage would be three or four orders of magnitude greater than the presently used goals.

Several methods of reducing the tritium leakage to the environment can be used:

- Dilution of the chamber gases with hydrogen or helium.
- Adding a liquid getter to the $\text{Li}_{17}\text{Pb}_{83}$ to trap tritium.
- Removing tritium from the $\text{Li}_{17}\text{Pb}_{83}$ with counterflowing helium.
- Using double-walled heat exchanger tubes with an oxide layer and/or flowing-gas sweep system.
- Introducing a liquid sodium intermediate loop, and cold trapping tritium from the sodium.

Fig. 7-25. Permeation rate of tritium through 4.8 mm iron with no barriers for a drive pressure of 1.7×10^{-3} Torr.

- Removing tritium from the water in the steam cycle.
- Substituting liquid lithium for $\text{Li}_{17}\text{Pb}_{83}$ as the coolant.

We are examining all methods quantitatively, and we have performed preliminary calculations for hydrogen dilution, cold trapping in sodium, and the use of pure liquid lithium instead of $\text{Li}_{17}\text{Pb}_{83}$.

If Sievert's law applies to the diffusion of tritium through a thin metal wall, then the rate of diffusion through the heat-exchanger wall should be proportional to the square root of the tritium partial pressure (assuming zero tritium pressure on the steam side). Thus, to reduce the diffusion rate by more than four orders of magnitude, we must reduce the partial pressure by more than eight orders of magnitude.

If hydrogen or helium gas is added to the chamber and the chamber pressure is simultaneously kept below 5×10^{-3} Torr, then the partial pressure of tritium is reduced by the resulting dilution. If hydrogen gas is used for the dilution, then cold trapping in an intermediate sodium loop can be used to further reduce the tritium driving pressure in the steam generators. However, using hydrogen in the chamber complicates the process of separating the tritium from the other gases in the target factory because of the isotopic separation techniques that must be used. Although using helium as the diluent removes these complications, cold trapping in sodium is not possible with

Fig. 7-26. Rate of tritium transfer from Pulse*Star reaction chamber into steam for three methods of control.

helium as the diluent because the total hydrogen isotope content in the sodium would not be great enough (unless we added hydrogen only in the sodium loop).

Some hydrogen will enter the chamber by diffusing from the steam side of the heat exchanger into the $\text{Li}_{17}\text{Pb}_{83}$. Hydrogen (^1H) is formed by the corrosion reaction between hot H_2O and the steam-generator tubes. Some of this nascent hydrogen combines with other hydrogen atoms to form molecular hydrogen that goes back into the water. The majority of the nascent hydrogen atoms, however, diffuse through the metal wall. Measurements on 2-1/4 Cr-1 Mo steam generators used in the liquid-metal fast-breeder-reactor (LMFBR) program show the ^1H flux from this source to be about $0.2 \times 10^{-6} \text{ g/m}^2 \cdot \text{s}$ (Ref. 76). Because each of the eight Pulse*Star steam generators has an area of 1840 m^2 , $\sim 32 \text{ g/d}$ of hydrogen will diffuse through the walls of each steam generator. Thus, the total input per day of ^1H diffusing towards the $\text{Li}_{17}\text{Pb}_{83}$ pool is $\sim 250 \text{ g/d}$. For each individual steam generator, the 32 g-atoms of ^1H (32 g) diffusing toward the $\text{Li}_{17}\text{Pb}_{83}$ liquid is 17 times greater than the 1.9 g-atoms (6 g) of ^3H diffusing out of the $\text{Li}_{17}\text{Pb}_{83}$. However, 250 g of ^1H is small compared to the 1.4 kg of tritium that is injected in the form of targets each day. Much greater dilutions will be necessary, and the hydrogen will have to be

injected directly into the chamber. Our calculation shows that even if we use helium as the diluent, some hydrogen will enter the chamber anyway, making isotopic separation in the target factory necessary.

If we add an intermediate sodium loop between the $\text{Li}_{17}\text{Pb}_{83}$ and the steam circuits, we can cold trap tritium from the sodium. We estimate⁷⁷ that cold trapping reduces the total tritium diffusion rate into the steam by a factor of 10.

Finally, we could replace $\text{Li}_{17}\text{Pb}_{83}$ with liquid lithium as the coolant. In the lithium loop, hydrogen isotopes could be kept below 1 wppm by electrochemical removal.⁷⁷ We could reduce the tritium diffusion rate more than three orders of magnitude using this method.

Figure 7-26 shows the maximum total tritium transfer rate for three combinations of limiting measures: hydrogen dilution, hydrogen dilution plus sodium cold trapping, and pure lithium instead of $\text{Li}_{17}\text{Pb}_{83}$ with hydrogen dilution. (We assume a driving tritium pressure of $1.7 \times 10^{-3} \text{ Torr}$ and a 4.8-mm-thick heat-exchanger tube wall.) Only the use of pure lithium can limit the total tritium transfer rate to 1 to 50 mg/d. In future work, we will examine methods of reducing tritium leakage in quantitative detail, and account for the effects of back diffusion. For now, we can only conclude that control of tritium leakage in the Pulse*Star reactor requires considerable attention.

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A Version of Pulse*Star Using a Liquid-Lithium Coolant

Our concern about tritium leakage and polonium formation in the Pulse*Star design that uses $\text{Li}_{17}\text{Pb}_{83}$ as a coolant has prompted us to design a version of Pulse*Star that uses pure liquid lithium as the primary coolant and liquid sodium as the secondary coolant. In our original design, all the components were contained in the liquid-metal pool.⁷⁸ In our new version, the primary circuit components are contained in a 17-m-diam lithium pool; the secondary circuit components are just outside

the pool wall. We selected specific configurations by minimizing capital costs and pumping power. We estimated the overall power balance of the plant.

The heat-transfer circuits are designed to transfer 3300 MW_t to the steam generators. The hot leg and cold leg temperatures are set at 500°C and 400°C, respectively, for the lithium circuit, at 482°C and 316°C for the sodium circuit, and at 460°C and 243°C for the steam circuit (pressure $P = 15.3$ MPa) based on the HYLIFE design.²⁰ The thermal conductivity, density, viscosity, and heat capacity of lithium⁸¹ at 450°C are 47.8 W/m·K, 490 kg/m³, 3.73×10^{-4} kg/m·s and 4.19 kJ/kg·K, respectively. The values of the same parameters for liquid sodium⁸¹ at 400°C are 72.3 W/m·K, 856 kg/m³, 2.85×10^{-4} kg/m·s, and 1.28 kJ/kg·K, respectively.

To produce 3300 MW_t with a 100°C temperature difference, the lithium flow rate through the intermediate heat exchangers (IHXs) must be 16 m³/s. We obtain a reasonable flow rate of 4 m³/s per pump by dividing the total flow into four parallel circuits. A low flow velocity through each circuit is necessary to minimize the piping pressure loss component of the pumping power. We can obtain a 3.5-m/s flow rate by using 1.2-m-diam piping. At 3.5 m/s, we estimate a minimal pressure loss of about 3×10^{-1} m (Li) of head for each meter of pipe length.

To remove 3300 MW_t from the IHXs with a sodium temperature difference of 166°C, the sodium system must have an 18-m³/s flow rate. Dividing this flow rate into four heat-transfer circuits results in a flow rate of 4.5 m³/s per heat-transfer circuit. We can obtain the reasonably low flow velocity of 3.9 m/s by using 1.2-m diam piping (at this flow velocity, the pressure loss through the pipe is minimal).

Liquid Metal Pump Selection. A pump is required in each of the four primary and four secondary heat-transport circuits. In each primary circuit, the liquid lithium is forced through two parallel IHXs. In each of the secondary heat-transport circuits, the liquid sodium is pumped through the two parallel IHXs, and then into two parallel steam generators. The pump characteristics and dimensions selected⁸² are summarized in Table 7-17.

	Primary loop	Secondary loop
Number of pumps	4	4
Flow per pump (m ³ /s)	4.0	4.5
Developed head (m)	17.4	13.2
Developed pressure (kPa)	86.2	111
NPSH ^a available (m)	2.5	3.0
NPSH ^a required (m)	1.9	2.5
Impeller diameter (m)	1.2	1.2
Pump speed (rpm)	445	420
Net fluid power per pump (MW)	0.34	0.50
Electric power per pump (MW _e)	0.43	0.63
Pump heat absorbed by fluid (MW _f)	0.40	0.59

Table 7-17. Pump characteristics and dimensions for lithium Pulse*Star.

Fig. 7-27. Heat exchanger capital cost and pump operating cost vs P/D .

Heat Exchanger Design. The thermal energy extracted from the primary circuit is transferred to the secondary circuit through the IHXs. The energy from the secondary circuit is used to generate steam by circulating sodium through the shell side of the steam generators (SGs). Based on an earlier study,⁸³ we selected countercurrent heat exchangers for both the IHXs and the SGs. A heat exchanger height of 12 m was chosen to correspond to the reaction chamber height. To complete the layout of the reactor building, we must find the necessary heat exchanger diameters, which will

Pulse*Star Reactor

depend on the required heat exchanger surface area.

The total required heat-transfer surface area A for any heat exchanger is given by

$$A = Q/U \Delta T_{lm} \quad (40)$$

where Q is the total power requirement (3300 MW_e), U is the overall heat-transfer coefficient (W/m²·K), and ΔT_{lm} is the log mean temperature difference. The values of ΔT_{lm} for the IHXs and the SGs are 43°C and 45°C, respectively.

The overall heat-transfer coefficient U is given by

$$\frac{1}{U} = \frac{A_o}{A_i} \frac{1}{h_i} + \frac{r_o \ln(r_o/r_i)}{k} + \frac{1}{h_s} \quad (41)$$

where

A_o is the outside heat-transfer surface area,
 A_i is the inside heat-transfer surface area,
 r_o is the outside tube radius,
 r_i is the inside tube radius,
 k is the thermal conductivity,

Fig. 7-28. Conceptual layout of the pumps and circuits in the lithium Pulse*Star.

h_t is the average heat-transfer coefficient for the tube side,
 h_s is the average heat-transfer coefficients for the shell side.

Both h_t and h_s are complex expressions of the material properties and geometry. We developed these expressions from turbulent flow analysis. Both the pressure drop through the shell side of heat exchangers and the average heat-transfer coefficient for the shell side of the heat exchangers are dependent on tube spacing. Because of this, we can optimize the heat-exchanger design with tube spacing that will give the most cost-effective combination of heat-exchanger area and pumping power. As the tube spacing increases, flow resistance and pumping costs decrease. However, the heat-transfer coefficient is also reduced, resulting in increased tube surface area and therefore increased heat exchanger cost. Our objective is to select a tube spacing that will minimize the total cost (pumping plus heat exchanger area) for the 30-y plant lifetime.

Economic parameters used for the present optimization: are plant electrical costs at 3¢/kW_eh, 80% pump efficiency, 80% plant capacity factor, 30-y plant life, and a 12% per year economic rate of return. We neglect the slight dependence of the pump capital cost on pump head in this analysis. The capital cost of a heat exchanger is proportional to its heat-exchange surface area: we used \$4500/m² to calculate the heat-exchanger cost.

Figure 7-27, taken from Ref. 82, shows the cost of various items as a function of P/D , where P is the spacing between tube centerlines and D is the tube diameter. To find an optimum for the IHX tube spacing, we must combine the cost of lithium pumping with the cost of the IHX. This curve produces a minimum at $P/D = 1.3$. At this minimum, we calculated a necessary heat-transfer area of 7840 m², which indicates an IHX diameter of 1.2 m for each of eight IHXs and a corresponding shell-side pressure loss of 5.4 m of head.

For the SGs, Fig. 7-27 shows a minimum cost at $P/D = 1.5$. At this minimum, the total required heat-transfer area is 16 160 m². Dividing this into eight steam generators produces a diameter of each of 1.5 m with a corresponding shell-side pressure loss of 2.2 m of head.

Overall Layout and Power Flow.

Figure 7-28 shows a conceptual layout of the pumps and circuits for the lithium pool. The geometry is fairly compact, even with the sodium secondary circuits. The rest of the plant layout is very similar to the original Pulse*Star layout.⁸³

The lithium Pulse*Star power balance is shown in Fig. 7-29. The main differences between the lithium and Li₁₇Pb₈₃ Pulse*Star energy flow designs are the smaller system energy multiplication factor and auxiliary pumping requirements in the lithium design.

Table 7-18 compares the major design features of the lithium Pulse*Star design to the previously reported Li₁₇Pb₈₃ Pulse*Star

Fig. 7-29. Lithium Pulse*Star power balance.

	Lithium	Li ₁₇ Pb ₈₃
Primary circuit coolant	Lithium	Lead-lithium
Secondary circuit coolant	Sodium	None
Number of pumps	8	4
Number of intermediate heat exchangers	8	None
Number of steam generators	8	8
Total heat-exchanger area (m ²)	24 000	14 700
Pool height (m)	16	17.5
Total required pumping power (MW)	4.2	35.5
Pressure drop through secondary circuit (m)	13.2	—
Pressure drop through primary circuit (m)	17.9	14
Total heat exchanger capital cost and pump operating cost (\$M)	163	173

Table 7-18. Comparison of the Li₁₇Pb₈₃ and lithium Pulse*Star configurations.

design.⁸³ The significant differences are the heat exchanger area (almost twice that of the $\text{Li}_{17}\text{Pb}_{83}$ Pulse*Star) and the pumping power (only one-eighth that required for the $\text{Li}_{17}\text{Pb}_{83}$ version). The pool height is smaller for the lithium version because the net-positive suction head (NPSH) required by the pump is smaller.

Summary. We have proposed a compact configuration for a Pulse*Star reactor using lithium as the primary coolant and sodium as the secondary coolant. Conceptual designs of the pumps and heat exchangers are within the range of existing technology. The recirculating power and plant net efficiency are reasonable. The addition of the secondary circuit into this design does not increase the combined pump-power and heat-exchanger costs compared to the $\text{Li}_{17}\text{Pb}_{83}$ design. The layout of the proposed design compares the $\text{Li}_{17}\text{Pb}_{83}$ plant to one of pure lithium with a sodium secondary circuit. Further work is

needed to compare the lithium design with the $\text{Li}_{17}\text{Pb}_{83}$ design in greater depth.

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New Concepts

Starfish Reactor

The Starfish Reactor concept⁸⁴ (Fig. 7-30) maximizes the conversion of fusion energy into kinetic energy, which is then directly converted to electricity using turbine generators. The concept involves placing a liquid-lithium (or lithium-lead) falling blanket as close as possible to the central axis of the chamber. If looked at in cross section the blanket forms an annulus, unlike the cross section in the HYLIFE design,⁸⁵ which forms a series of discrete jets.

Fusion energy in the form of x rays, fuel pellet debris, and neutrons is absorbed into the lithium blanket. High pressure developed in the blanket's central core accelerates the lithium radially. The lithium hits the chamber wall, and a portion of it is channeled through converging passages to turbine generators. As the lithium interacts with the cone-shaped flow enhancers that make up the surface of the inside chamber wall (Fig. 7-30), kinetic energy from the majority of the lithium blanket is transferred to those portions of the blanket directed to the turbine generators. The phenomenon is similar to that of a large amount of sea water transferring kinetic energy to a smaller portion of sea water as it splashes against rocks at the shoreline.

After kinetic energy is extracted from the lithium that passes through the turbine generators, that lithium is collected with the lithium that has fallen out the bottom of the chamber. Once collected, it is directed to heat exchanger loops where thermal energy is extracted and converted into electricity via a standard steam cycle.

In 1983, Lasche⁸⁶ showed that a maximum of one-half the fusion energy can be converted to kinetic energy in his spherically symmetric Sunburst reactor concept. Our geometry is cylindrical, and while we expect less conversion to kinetic energy

Fig. 7-30. Starfish reactor concept.

(a) Top view

(b) Cross section view

than with the Sunburst reactor, we expect a total conversion efficiency greater than that yielded by the Sunburst or HYLIFE reactor. The Starfish will allow us to directly convert four-tenths of the fusion energy to kinetic energy, with 90% efficiency. Then, with 40% efficiency, we can convert the remaining six-tenths of the fusion energy to electricity using the standard steam cycle. Together, the two conversion processes can give a total efficiency of 60% (Table 7-19).

Fuel pellets, used to produce the fusion energy, are injected into the central core of the chamber through the top opening. As each pellet reaches the chamber center it is illuminated with laser or ion beams from both the top and the bottom of the chamber. The chamber is kept at vacuum to permit beam propagation. The lithium blanket shields structural components from radiation damage and produces tritium to replace that burned in the fusion fuel pellets.

Starfish is a novel concept because it retains the advantages of the HYLIFE and Sunburst designs while improving on their limitations. Starfish is expected to have a 40% higher conversion efficiency than HYLIFE, and it does not require magnetic fields around the chamber, as the Sunburst reactor does.

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10-MJ Laser Target Chamber Options

The next step beyond Nova will be to produce high gain in a fusion fuel pellet. We estimate that 1 to 10 MJ of driver energy will be required to produce the requisite high gain. Such a facility, if upscaled from Nova, could be very expensive. Our efforts, then, are to minimize costs of each component.

We present three options for a low-cost target chamber. Each option is based on a target with high gain (~ 100), which would yield a total energy of 1 GJ. Environmental effects can be substantial when fusion yields of this magnitude are involved, and many options (including some presented here) may be rejected for this reason.

Option 1—No Target Chamber. In the first option [Fig. 7-31(a)], the fusion target is placed at the center of a depression in the

Method of conversion	Fraction of energy converted (%)	Unit conversion efficiency (%)	Net conversion efficiency (fraction $\times$ unit) (%)
Direct turbine generator	0.4	90	36
Standard steam cycle	0.6	40	24
Total net conversion efficiency = 60%			

ground. A small plastic-covered vacuum barrier, approximately 100 mm in radius, is placed around the target. Evacuated beam ports extend horizontally to the laser final turning mirrors and necessary diagnostics. The vacuum barrier and evacuated beam ports, made of inexpensive wood frames, must be replaced following each shot.

A special nonactivating granular fill (perhaps boron carbide or processed borax), with a thickness of 3 m and a theoretical density of 1/2, protects diagnostics and laser optics permanently installed in the surrounding earth formation. This fill absorbs energy from x rays, fusion fuel pellet debris, and neutrons without transmitting shock waves. However, because neutrons activate the surrounding air, access to the site would have to be restricted.

The advantage of option 1 is its low cost, but environmental effects are a concern. First, tritium would be released into the atmosphere. The yield for completely burned deuterium-tritium is 340 MJ/mg. The mass of tritium contained in a 1000-MJ shot is 5.5 mg, with an associated activity of about 55 Ci. While the normal burn fraction of tritium is 0.32, from 66 to 100% of the tritium amount could be released, depending on the success of the burn. Although this amount is small, it may not be environmentally acceptable.

A second concern with this option is the prompt dose from the atmosphere caused by high-energy neutrons and gamma rays. Singh⁸⁷ used the TART⁸⁸ code, calculating that the dose would reach 1 mSv (100 millirem) per shot at a radius of 2 km. The proposed DOE standard⁸⁹ of 1 mSv per year suggests that no personnel would be permitted within this radius. Furthermore, site selection may have to be restricted because flora and fauna present at closer radii would receive significant doses.

We also calculated activation of ^{13}N , ^{41}Ar , and ^{16}N in the air, as well as the resulting dose at the site boundary (with a 5 m/s wind speed and average turbulence).^{90,91} For

Table 7-19. Calculated efficiency for the Starfish reactor.

Fig. 7-31. Three options for a 10-MJ laser target chamber: (a) no target chamber, (b) disposable target chamber, and (c) reusable target chamber.

a site boundary at 2 km, the dose is only $6 \mu\text{Sv}$. This value, which is insignificant compared to the prompt dose, poses no additional hazard.

Beyond the technical issues addressed here lie political issues dealing with the relative benefits of fusion, the release of tritium, and the activation of the atmosphere. These issues would have to be resolved before this target chamber option could be used.

Option 2—Disposable Target Chamber.

Here we suggest a 3-m inner radius dome, also made of a nonactivating granular material (boron carbide or processed borax) 3 m thick, at 1/2 density [Fig. 7-31(b)]. The B_4C sand in the upper hemisphere is held in place by a wooden geodesic dome structure over which plastic is placed to form a vacuum barrier. Laser beam ports and diagnostic tubes extend from the chamber, each surrounded by granular material. Following every shot, x-ray and debris energy vaporize the plastic, but the gases formed in the process permeate the porous granular

material, dissipating energy and recondensing before reaching the outer radius.

Neutrons are absorbed in the granular material. Tritium and other toxic materials are contained in the target chamber room (or in a similar large containment structure), closed to the atmosphere. These toxic materials, including the same 55 Ci of tritium released with option 1, are then removed and either recycled or placed in an appropriate disposal. The wooden dome structure is destroyed during each shot; however, the granular material that collapses on the floor can be scooped up by a front loader and reused on the next shot.

The cost of this option is low, and no toxic materials are released into the atmosphere. Singh⁸⁷ also calculated the neutron and gamma dose for a 3-m inside radius chamber with B_4C walls. If the walls are 3 m thick, the neutron dose is reduced eight orders of magnitude, to a total of 1 mSv at the outside surface.

The plastic liner that holds the B_4C sand in place is exposed to x rays and fuel pellet debris during each shot, and is thus destroyed. We calculated that the pressure inside the chamber that results from vaporization of the inside surface of the plastic is about $4 \times 10^5 \text{ Pa}$ (4 bar). The vaporized material flows into the B_4C sand, where it condenses. Increased pressure in the target chamber room is insignificant and presents no problem. Before condensation is complete, the gas pressure accelerates the granular material. This acceleration is small, however, and the maximum vertical motion of a granule is only 6 cm above its original position.

Option 3—Reusable Target Chamber.

This option involves a permanent aluminum target chamber lined with B_4C sand that is held in place by a plastic liner [Fig. 7-31(c)]. The liner absorbs x-ray and debris energy and must be replaced following each shot. Like the disposable chamber, the liner is supported by a wooden frame. The sand dissipates any shock waves caused by the fusion reactions, and the aluminum target chamber contains all gases so that no toxic materials are released into the target chamber room or the atmosphere. While beam and diagnostic ports would be attached in the same fashion as the Nova target chamber, all ports would be shielded with granular materials to reduce neutro

leakage. The advantage of this option is that the chamber can be used to structurally support diagnostics and beam tubes; however, there will be some activation of the chamber itself. This option is the most expensive of the three, but it is also the most reliable and flexible.

Stresses in the chamber wall come from two major sources: pressure that results from vaporization of material on the inner liner surface, and momentum generated by x-ray and debris energy deposition. For a 4-cm-thick aluminum wall, the stresses due to pressure and to x-ray and debris are 55 and 25 MPa respectively. The total of 80 MPa is slightly less than the value allowed by the American Society of Mechanical Engineers (1/3 the yield strength of the aluminum). Hence a 4-cm-thick wall could be used to contain toxic materials and could be reused from shot to shot for the 10-MJ laser target chamber.

We have identified three options for the 10-MJ laser target chamber. While we recognize advantages and disadvantages of each, final selection requires further examination.

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General Studies

Identifying Heavy-Ion-Beam Fusion Design and System Features with High Economic Leverage

Introduction. One of the primary objectives of the ICF Applications Group is to develop power plant concepts that will be economically competitive with other long-term electric power producers (e.g., with fission and coal). We are developing systems models for ICF power plants that allow us (1) to identify design features that have the highest leverage for improving ICF economics and (2) to do first-order economic comparisons with other power sources.

In this article we consider a heavy-ion-beam (HIB) fusion power plant that consists of a driver, a target factory, and one or more power units. A power unit is

defined as all the buildings and equipment needed to generate electric power; provided the target and beams are delivered to the reaction chamber. Because the maximum achievable pulse rate in a single chamber is limited, more than one reaction chamber may be required to achieve the desired output of a single power unit. We distinguish between multiple power units and multiple reaction chambers so that we can examine separately the effects of increasing the number of reaction chambers at a constant net power and of increasing the power level by driving more power units with a single driver.

We conducted studies to investigate the effects on the cost of electricity (COE) of variations in several design parameters. In particular, we examined the effects of maximum achievable chamber pulse rate, driver cost, target gain, electric conversion efficiency, and net electric power. We found that with a combination of improvements over our base case, HIB fusion can be economically competitive with present and future power sources.

While economic factors are not the only ones that we consider when making decisions about future power sources, such factors as safety, environmental impact, and public acceptance may not weigh heavily unless ICF appears to be economically competitive. Although the definition of economic competitiveness is subject to debate, a common figure of merit in any such analysis is the COE. Because of economies of scale, the COE decreases monotonically with increasing power plant size. In a more extensive economic analysis, other figures of merit that reflect total capital investment would necessarily be considered. However, in this initial step we limit our discussion to the COE as the figure of merit. The cost models we used are reported in detail in a separate report.⁹² The essence of the models is summarized here.

Power Unit Cost. The model we have assumed for the HIB fusion power plant uses a conventional steam cycle to produce electricity. Each power unit looks like a fission power plant, except that the fusion chamber replaces the fission reactor vessel. Because of the structural similarities between the two, we have derived our cost-scaling models from recent estimates for future fission plants. Fission power plants

recently completed and currently under construction are significantly more expensive to build than coal-fired power plants. However, researchers in the Engineering Technology Division at Oak Ridge National Laboratory and United Engineers have estimated the cost of nuclear power plants for the 1990s based on the presumed effects of licensing reform.^{73,94} The 1990's nuclear plants have direct capital costs that are essentially equal to those of coal plants. We assume that the fusion plant can achieve similar direct capital costs for those elements not unique to fusion. On this basis, the direct capital cost of a 1-GW_e power unit is \$0.7 billion.

As previously stated, each power unit may have more than one reaction chamber. To account for this in our cost algorithms, we break down the direct capital cost into three parts: the containment building cost, the steam-supply-system (SSS) cost, and the remainder of the power unit cost. The report by United Engineers gives a detailed breakdown of the cost of the pressurized water reactor for the 1990s.⁹⁴ It shows that the reactor containment building cost is 7.5% of the total direct cost and that the SSS accounts for 18.2% of the total direct cost. We allocate additional containment building space and SSS equipment for each additional fusion chamber using these same fractions. The direct capital cost of the power unit (C_{ud}) can be written as

$$C_{ud} = 0.705 (0.743 P_g^{0.6} + 0.075 P_c^{0.3} N_c + 0.182 P_c^{0.6} N_c) \$B \quad (42)$$

Fig. 7-32. The gross electric power required to maintain a constant net electric power increases with increasing driver pulse rate.

where P_g is the gross electric power for one power unit, P_c is the gross electric power associated with each chamber, and N_c is the number of chambers per power unit. Hence, $P_c N_c$ is equal to P_g for each power unit. Experience has shown that there are savings in both the direct and indirect capital costs for multi-unit power plants.⁹³⁻⁹⁵ Typically, each additional unit has a direct capital cost equal to 80% of the direct capital cost of the first unit on the site and an indirect capital cost equal to 60% of the indirect cost for the first unit. These factors are included in our analyses.

Driver Cost. The cost of an HIB driver has been estimated based on radio-frequency accelerator⁹⁶ and induction-linac⁹⁷ technology. The driver direct capital cost (C_d) is given by⁹⁷

$$C_d = 0.64 E_d^{0.4} + 0.1 (N_c N_u - 1) \$B \quad (43)$$

where E_d is the driver beam energy in MJ. The second term in Eq. (43) accounts for the additional cost of transporting beams to the individual reaction chambers.

Other Factors. In addition to the driver and power unit costs, we include \$100 million for the direct capital cost of the target factory and 1¢/target for consumable materials. The total capital cost is 1.9 times the direct capital cost. This factor accounts for contingency and indirect construction costs (10% and 50%, respectively), as well as interest during construction (15%). Annual operation and maintenance costs, equal to 2% of the total capital cost, are also included. The fixed charge rate used in our constant-dollar analysis is 10%, and the plant capacity factor is 70%.

Results. The base-case parameters used in our calculations are a net power per power unit (P_n) of 1 GW_e, one power unit per plant ($N_u = 1$), maximum chamber pulse rate (ω_{cm}) of 10 Hz, and an electric conversion efficiency (η_e) of 40%.

Figure 7-32 shows the constant net power, driver and auxiliary power requirements, and gross electric power as a function of the driver pulse rate for the base case. The driver power requirement increases with increasing pulse rate. This follows since a higher pulse rate corresponds to a lower target gain, and, therefore, more power must be recirculated to run the driver. The gross electric power required to

maintain a net output of 1.0 GW_e increases from 1.15 to 1.45 GW_e as the driver pulse rate increases from 5 to 50 Hz.

Maximum Chamber Pulse Rate as a Parameter. The COE as a function of the driver pulse rate is shown in Fig. 7-33 for four different maximum chamber pulse rates (ω_{cm}): 5, 10, 20, and >50 Hz (unlimited). The discontinuities, which occur every ω_{cm} Hz, reflect the cost of additional SSS equipment-containment building space and beam-transport lines incurred each time an additional chamber is added.

Focusing on the base case with $\omega_{cm} = 10$ Hz, we see that the COE initially decreases with increasing driver pulse rate. This is because driver cost decreases with increasing driver pulse rate, since higher pulse rates correspond to lower driver energies. The minimum COE is $7.0 \text{ ¢/kW}_e\text{h}$ for the base case and occurs at 10 Hz. At this point, the driver energy is 3 MJ and the gross electric power is 1.18 GW_e .

Note the flatness of the dashed line in Fig. 7-33. If a chamber could be designed to operate at an unlimited pulse rate (or if no cost were associated with increasing the number of chambers), the minimum COE would be $6.75 \text{ ¢/kW}_e\text{h}$, which is realized between 25 and 30 Hz. Hence the 10-Hz limit on chamber pulse rate increases the COE by less than 4%. If the chamber pulse rate is limited to 5 Hz, the COE is $7.4 \text{ ¢/kW}_e\text{h}$, or less than 10% higher than the $6.75 \text{ ¢/kW}_e\text{h}$ minimum. The curves also rise at high pulse rates since the cost of the power unit is increasing with increasing gross power.

Driver Cost as a Parameter. Figure 7-34 illustrates the sensitivity of the COE to the cost of the driver. The three curves, from top to bottom, are the base case, the driver cost reduced by 25%, and the driver cost reduced by 50%, respectively. The cost reductions are applied to the total driver cost, including the additional beam-transport lines required for power units with more than one chamber. In all three cases, the minimum COE occurs at 10 Hz. Because the driver cost is about half of the total capital cost, reducing the driver cost leads to a significant reduction in the COE. Each 25% reduction in the driver cost gives a 13% reduction in the COE.

By looking at the dashed curves in Fig. 7-34, one can see that if the chamber pulse rate is unlimited, the optimum driver

Fig. 7-33. The minimum COE decreases as the maximum achievable chamber pulse rate increases.

Fig. 7-34. Reducing the driver cost reduces the COE.

pulse rate decreases with decreasing driver cost. In addition, the COE becomes less sensitive to driver pulse rate (i.e., the curves tend to flatten out).

Target Gain as a Parameter. We also investigated the effect of improving target performance. The reference target gain corresponds to the $\rho^{3/2}R = 0.01$ case given in Ref. 98. It is fit by the expression

$$G = 20 + 113 \ln (E_d/1.7) \quad (44)$$

As a variation on this reference case, we assumed that the driver energy required to achieve a given target gain is reduced by a factor of 3. This is consistent with estimates for polarized fuel. (See the article "Spin-Polarized DT" in Sec. 4.) At 10 Hz the COE is reduced by 12%. With no limit on chamber pulse rate, the COE minimum shifts toward higher pulse rates as gain improves at lower driver energy. Moreover, the COE also becomes more sensitive to pulse rate.

Electric Conversion Efficiency as a Parameter. We examined the effect of increasing the electric conversion efficiency by assuming that the conversion efficiency can be increased without increasing the cost per kW_e . We found that increasing the conversion efficiency from 40 to 60% reduces the COE by only 12%. Remember that the net electric power is held constant; therefore, a more efficient power unit is smaller and suffers an economy-of-scale penalty. That is, the cost per kW_e is only the same for plants of equal thermal power.

Plant Power as a Parameter. As noted previously, there are two ways to consider net plant power as a variable. The number of units can be set equal to one, and we can vary the power of the unit. On the other hand, we can keep the power of each unit

constant and vary the number of power units in a plant.

First, we considered a single-unit plant. Figure 7-35 shows the COE for net power outputs of 0.5, 1.0 and 2.0 GW_e . The COE from the 0.5- GW_e plant is 66% greater than the base case at 10 Hz (i.e., 11.6 vs 7.0¢/kW_eh). The 2.0- GW_e plant produces electricity at 4.3¢/kW_eh, or 39% less than the base case. It is clearly advantageous, in terms of the COE, to use the driver to produce as much power as possible. Other considerations, such as following load growth and minimizing total capital investment, may favor smaller net powers.

We also considered the advantages of using a single driver and target factory to operate more than one power unit. At a driver pulse rate of 20 Hz, a two-unit, 2- GW_e power plant produces electricity at 4.8¢/kW_eh, or 31% less than the single-unit, 1- GW_e plant. A four-unit, 4- GW_e power plant reduces the COE to 3.6¢/kW_eh, or 49% less than the single-unit plant. However, the total capital investment increases as the number of units increases. The total capital costs for the 1-, 2-, and 4-unit plants are \$3.55, \$4.81, and \$7.35 billion, respectively.

We found that as the number of units increases, the COE becomes less sensitive to the driver pulse rate. In fact, the four-unit plant could operate at a driver pulse rate of 12 Hz and still be within 5% of its minimum COE. The required chamber pulse rate in this case is only 3 Hz.

Figure 7-36 shows the COE as a function of the total net power for plants made of different size units: 0.5, 1.0, and 2.0 GW_e net. The COE from a 1.0- GW_e power plant consisting of two 0.5- GW_e units and a single driver is 7.7¢/kW_eh. This is 10% higher than the COE from a single-unit power plant generating 1 GW_e . This higher COE is the result of the economies of scale in the power unit cost and the additional cost of beam-transport lines for the two-unit plant. The COE from a 2- GW_e power plant with a single unit is 4.3¢/kW_eh, compared to 4.8¢/kW_eh for a two-unit plant, and 5.8¢/kW_eh for a four-unit plant.

At a fixed net power, the COE is lower for a single-unit plant than for a multi-unit plant. The multi-unit plant does, however, achieve most of the economy-of-scale benefits of a single large unit. For example, going from 0.5 to 2 GW_e in a single-unit plant

Fig. 7-35. The COE decreases as the net power of the single-unit plant increases.

Fig. 7-36. For the same net power, plants with fewer, larger power units have lower COE.

reduces the COE by 63%, compared to the 50% reduction achieved by a four-unit plant. An advantage of the multi-unit plant is its operational flexibility. The plant can begin to operate and produce revenues from one unit while decisions are pending on the need for additional units. Capital costs for the future units are, of course, also deferred.

Several Parameters Simultaneously.

Most of the parameters considered have a modest impact on the COE. No single improvement considered is able to reduce the COE to a point where it would be competitive with coal (4.6¢/kWh) or nuclear (3.5¢/kWh) power.⁹³ Hence it is important to make advances in as many areas as possible. In Fig. 7-37 we consider several simultaneous improvements: a driver cost reduction of 25%, improved target gain, an electric conversion efficiency of 50%, and two-unit power plants. We found that a 2-GW_e plant produced power for 3.6¢/kWh, which is economically competitive with coal and nuclear power.

Conclusions. We have examined the COE of HIB fusion power plants to determine the economic impact of various design and system improvements. Our conclusions fall into four major areas.

Pulse Rate. For the single-unit, 1-GW_e power plant, the optimum pulse rate (for both driver and chamber) is 25 to 30 Hz, resulting in a COE of 6.75¢/kWh. However, the COE is not very sensitive to pulse rate. For example, if the chamber pulse rate is limited to 5 Hz, the COE increases by less than 10%. This weak dependence of COE on pulse rate is important because there are major uncertainties associated with predicting the maximum achievable chamber pulse rate. Decreasing the driver cost or increasing the plant size results in a lower optimum pulse rate, and the cost of electricity becomes even less sensitive to pulse rate. Improving the target gain (which corresponds to reducing the required driver energy for a given gain) results in a higher optimum pulse rate and a greater sensitivity of the COE to pulse rate. Increasing the electric conversion efficiency has little effect on the location of the optimum pulse rate.

Multiple Reaction Chambers. For the conditions examined here, using multiple reaction chambers at constant net power never improved the COE. The optimum driver

Fig. 7-37. With several simultaneous improvements, the COE from HIB fusion power plants will be competitive with nuclear and coal power plants.

pulse rate in every case was at the maximum chamber pulse rate of a single reaction chamber. This was due to the large cost (\$100M direct) of adding beam-transport lines for each new chamber and to the insensitivity of the COE to pulse rate.

Economy of Scale. As expected, the COE decreases with increasing net electric power. The reduction is significant, whether the power is increased by increasing the power of a single unit, or by using a single driver to operate several power units. These savings result from the driver being such a large fraction of the plant cost. Only if the cost of the driver is significantly reduced will the penalty for small power units be significantly affected.

For a given net power, a multi-unit plant achieves most of the economy-of-scale advantage of a single large plant. Because the difference is so small, starting with a small plant (0.5 GW_e or less) and adding power units might be particularly attractive. In this way revenues can be collected from the first units, and the later units can be constructed as needed. Other advantages of smaller units, such as reduced on-site construction costs due to increased factory fabrication and shortened construction periods, have not yet been considered in our studies. These will tend to reduce the COE penalty of plants comprising several smaller power units.

Various Design Improvements. Reducing the HIB driver cost is important: a 25% reduction in the driver cost leads to a 13% reduction in the COE. Improving target performance is also important. If the same target gain can be achieved with a factor of 3 less driver energy, the COE can be reduced by 12%. Improving the electric conversion efficiency is not as advantageous as

might be expected. Increasing the conversion efficiency from 40 to 60% reduces the COE by 12%, assuming the power-unit cost per kW_e does not increase with increasing efficiency. If it does increase, the savings will be even less.

Only a combination of improvements reduces the COE to values in the range of those forecast for future coal and fission plants. Such a combination of improvements is not unreasonable, however. If the driver cost is reduced 25%, the target gain is improved at low E_d , the electric conversion efficiency is increased to 50%, and a two-unit plant is built (2 GW_e total), then the COE is competitive.

Summary. We find that the COE from HIB fusion can be competitive with other future energy sources. Staged construction of multi-unit plants may make smaller initial plant sizes more attractive. Chamber pulse rates of 5 to 10 Hz appear adequate to achieve these economic goals.

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Blanket Optimization for the HYLIFE Reactor

Introduction. Optimal blanket design is a key element in effective fusion reactor design. We have developed a methodology for optimizing the blanket design as a function of several variables.^{99,100} In this article, we apply this methodology to a modified version of the HYLIFE chamber as a function of two variables: x_1 (the ${}^6\text{Li}$ fraction in lithium) and x_2 (the effective lithium blanket thickness).

The HYLIFE concept^{101,102} is based on an array of liquid-lithium jets, which forms a 1-m-thick blanket that protects the first structural wall from direct exposure to the x rays, debris, and high-energy neutrons emitted by the fusion target. The lithium blanket reduces the 14-MeV neutron flux by a factor of 200, and neutron damage levels are low enough to allow the 5-m-radius chamber wall to operate for more than 30 y without replacement.¹⁰³ With a 1-m-thick blanket of natural lithium (7.42% ${}^6\text{Li}$ and 92.58% ${}^7\text{Li}$), HYLIFE achieves a tritium breeding ratio of 1.75 (Ref. 104).

Modified HYLIFE Design. We have modified the design of the HYLIFE chamber to increase the energy deposited in the

chamber per fusion reaction. Our approach is to reduce the tritium breeding ratio and capture the excess neutrons in ${}^{55}\text{Mn}$. Each neutron capture in ${}^{55}\text{Mn}$ adds 9.8 MeV to the energy balance. The ${}^{55}\text{Mn}(n,\gamma)$ reaction releases 7.3 MeV, while the decay of the product, ${}^{56}\text{Mn}$, releases an additional 2.5 MeV of recoverable energy. The half-life for the β^- decay is 2.6 h.

There are two simple ways of reducing the tritium breeding ratio in HYLIFE: reduce the fraction of ${}^6\text{Li}$ in lithium, or reduce the effective thickness of the lithium blanket.¹⁰⁵ In the modified HYLIFE design, we also replace the lithium-cooled graphite reflector, which surrounds the first structural wall, by a sodium-cooled, manganese-steel blanket. All the tritium breeding must occur in the lithium blanket.

The ratio of total energy deposited to fusion energy released per DT reaction is called the fusion energy multiplication factor M_f . For our calculations,

$$M_f = \frac{E_n + 3.5}{17.6} \quad (45)$$

where E_n is the total energy (in MeV) deposited as a result of neutron reactions within the chamber, including the neutron energy deposited in the compressed fuel region of an ICF target, 3.5 is the fusion alpha-particle energy in MeV, and 17.6 is the total fusion energy released per DT reaction in MeV. In general, M_f will be slightly greater than unity because of exoergic neutron reactions such as the ${}^6\text{Li}(n,\alpha)$ reaction, which releases 4.8 MeV. For HYLIFE, M_f is 1.16.

Increasing M_f by capturing neutrons in structural material has been proposed.¹⁰⁶⁻¹⁰⁸ For the calculations in this article, we use a 20 wt% manganese steel, known as Nippon Steel Alloy NM-1.¹⁰⁹ The composition of NM-1 is 77.5 wt% iron, 20.0 wt% manganese, 2.0 wt% chromium, and 0.5 wt% carbon.

An increased M_f offers the potential advantage of a lower cost of electricity from a fusion electric power plant by making more power available for a given investment in either an ICF or a magnetic fusion reactor. For an ICF reactor, the size of the laser and the characteristics of the target fix the fusion energy released per pulse. If this energy can be increased as a result of neutron reactions in the blanket, more thermal energy will be available for conversion to

electricity. In our blanket optimization problem, the characteristics of the driver, target, and pulse repetition rate are fixed: x_1 and x_2 relate only to the blanket design.

Design Constraints. Reducing the blanket thickness is advantageous in that it reduces the lithium flow rate, which saves pumping power and reduces capital costs for liquid metal pumps and piping. Reduction of the blanket thickness is limited, however, because of the maximum allowable radiation damage rate in the first structural wall. For HYLIFE this rate was such that the wall was expected to maintain its structural integrity for the 30-y life of the power plant. The life-limiting radiation damage mechanism for HYLIFE was found to be the repeated displacement of atoms from their lattice positions.¹⁰¹

Currently there are insufficient data to set absolute damage limits for structures in fusion reactors. We know that ferritic steels, such as the 2-1/4 Cr-1 Mo steel used in HYLIFE, are less susceptible to the effects of displacement damage than austenitic steels.¹¹⁰ A damage limit of ~ 200 displacement per atom (dpa) appears to be a reasonable estimate for ferritic steels.¹¹¹

In this article, the total displacement damage in the first structural wall of the modified HYLIFE chamber is limited to 200 dpa over an operating period of 30 years at a 70% plant capacity factor. This gives a constraint on the displacement damage rate of 9.5 dpa per full-power year. The dpa rate is expected to depend primarily on x_2 , but it may also depend somewhat on x_1 .

Another constraint in our optimization problem is the requirement for a sufficiently high tritium breeding ratio to ensure a self-sufficient fuel cycle. A tritium breeding ratio ≥ 1.05 has been chosen to satisfy this constraint. The tritium breeding ratio will be a function of both design variables.

Figure of Merit for the Modified HYLIFE Reactor. The figure of merit we used for this optimization problem is

$$F = C_t/P_n \quad (46)$$

where C_t is the total plant capital cost, including the costs of the reactor, laser, and target factory, and P_n is the net electric power (in kW_e) produced and available for sale.

The cost of electricity in cents per kW_eh from a fusion electric power plant is

expected to be dominated by the carrying charges on the capital investment and, therefore, proportional to F . Hence, minimizing the plant capital cost per unit of net electric power is equivalent to minimizing the cost of electricity.

The total capital cost C_T is expressed as

$$C_T = C_R + C_P + C_L + C_{TF} \quad (47)$$

where

C_R = cost of the reactor including the balance of plant required for heat transfer and conversion to electricity,

C_P = cost of lithium recirculating pumps and piping required to maintain the flowing lithium blanket,

C_L = cost of the laser driver,

C_{TF} = cost of the target factory.

The cost of the reactor scales as the plant thermal power P_t raised to an exponent less than unity. For our calculations, we conservatively use 0.8 as the exponent,¹¹² so

$$C_R \propto P_t^{0.8} \quad (48)$$

The thermal power is simply the fusion power P_f times the fusion energy multiplication factor M_f . Hence,

$$C_R \propto (M_f P_f)^{0.8} \quad (49)$$

As previously stated, P_f is fixed but M_f will be a function of both the ^6Li fraction, x_1 , and the lithium blanket thickness, x_2 .

The power and cost of the pumps will be proportional to the total lithium flow rate through the chamber. The flow rate is proportional to x_2 , hence

$$C_P \propto x_2 \quad (50)$$

The power and cost of the laser and the cost of target factory are fixed and are independent of x_1 and x_2 .

The costs in this article are based on those reported for HYLIFE in Ref. 113. The normalized values for the original HYLIFE chamber are given in Table 7-20. Using these values, the normalized total capital cost of the modified HYLIFE chamber is

$$C_1 = 0.62(M_f/1.16)^{0.8} + 0.10x_2 + 0.28 \quad (51)$$

Table 7-20. Original HYLIFE plant characteristics.

General Studies

Tritium breeding ratio (T)	1.75
Neutron energy deposition (E_n), MeV	16.9
Fusion power (P_f), MW	2700
Fusion energy multiplication factor (M_f)	1.16
Thermal power (P_t), MW _t	3130
Thermal conversion efficiency (η_t)	39%
Gross electrical power (P_g), MW _e	1220
Laser power consumption (P_l), MW _e	135
Auxiliary power requirements (P_a), MW _e	75
Lithium pumping power (P_p), ^a MW _e	30
Net electrical power (P_n), MW _e	980
Normalized capital costs	
Reactor (C_R) ^b	0.62
Lithium pumps (C_p)	0.10
Laser (C_L) ^c	0.21
Target factory (C_{TF})	0.07
Total (C_T)	1.00

^aBased on a lithium blanket thickness of 1.0 m and ~50% efficient electromagnetic pumps.

^bReactor includes the chamber and balance of plant required for heat transfer and conversion to electricity; it excludes the lithium recirculating pumps.

^cHalf the estimated cost of a 4.5 MJ laser. We assume that the laser drives two full-sized reactors by switching beams.

Fig. 7-38. The optimal design point lies at the intersection of the tritium breeding and displacement damage constraints.

The net electric power is given by

$$P_n = P_g - P_l - P_a - P_p \quad (52)$$

where

$$P_g = \text{gross electric power (MW}_e\text{)},$$

$$P_l = \text{laser power requirement (MW}_e\text{)},$$

$$P_a = \text{plant auxiliary power requirements (MW}_e\text{)},$$

$$P_p = \text{lithium pumping power (MW}_e\text{)}.$$

The gross electric power is equal to the thermal conversion efficiency η_t times the thermal power P_t and, therefore, is proportional to M_f . The laser and auxiliary power requirements are independent of the blanket design. Hence,

$$P_n = 1220(M_f/1.16) - 30x_2 - 210 \quad (53)$$

If we combine Eqs. (46), (51), and (53), and normalize to the HYLIFE net power of 980 MW_e, the result is

$$F = \frac{0.62(M_f/1.16)^{0.8} + 0.10x_2 + 0.28}{1.24(A_f/1.16) - 0.03x_2 - 0.21} \quad (54)$$

Our goal is to minimize F subject to the constraints on the displacement damage rate (≤ 9.5 dpa/y), and the tritium breeding ratio (≥ 1.05).

We used the methods presented in Refs. 99 and 100 to write analytical expressions for M_f , T , and the displacement damage rate (dpa/y) as functions of x_1 and x_2 . With this approach, we could calculate F at any point (x_1, x_2). We then used the direct-search optimization algorithm described in Refs. 99 and 100 to optimize F subject to the two constraints.

Optimal Design Point. Figure 7-38 shows that F decreases with decreasing ^{6}Li fraction and lithium blanket thickness. The constraints on the tritium breeding ratio and displacement damage rate limit the degree to which the ^{6}Li fraction and blanket thickness can be decreased to minimize F . The optimal design point lies at the intersection of the two constraints: a ^{6}Li concentration of 0.09% and a lithium blanket thickness of 0.91 m.

We performed a transport calculation for a 0.91-m-thick lithium blanket and a 0.09% ^{6}Li fraction to check the accuracy of our estimated results at the optimal design point. These calculations were performed with the TART Monte Carlo transport code.¹¹⁴ The results are compared in Table 7-21. All the results agree quite closely. The tritium

breeding ratio is off by ~0.5%, the total energy deposition is within 1% of the transport calculation result, and the dpa rate is high by ~0.5%.

Summary. The modified-HYLIFE plant parameters at the final optimal design point are compared to the reference HYLIFE parameters in Table 7-22. With the modified design, the fusion energy multiplication factor is increased to 1.34. As a result, the modified design produces 19% more electric power. The plant capital cost, however, is only 6% higher for an 11% reduction in the cost per kW_e.

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Blanket Optimization for the Cascade Reactor

Introduction. In 1984 we performed blanket optimization studies for the Cascade reactor concept.¹¹⁵ The primary feature of Cascade is a rotating chamber in which a blanket of solid lithium ceramic granules breeds tritium, acts as the heat transfer medium, and protects the chamber wall from the damaging effects of neutrons, x rays, and target debris.^{116,117} Granules are injected at each end of the chamber and are held against the wall by centrifugal action. The granules cascade toward larger radii and are then removed and sent to heat exchangers. Heat and tritium are extracted and the granules are recirculated to the chamber.

In the original design, we used Li₂O as the solid breeding material.¹¹⁷ While Li₂O is a good tritium breeding material, there are some concerns about the corrosive effects of LiOH (Refs. 118 and 119). From a compatibility standpoint, LiAlO₂ is a more attractive ceramic tritium breeding material because it does not form LiOH (Ref. 120). Unfortunately, LiAlO₂ will not give a tritium breeding ratio greater than unity unless a neutron multiplier is placed between the fusion neutron source and the LiAlO₂ breeding blanket.

The Cascade Optimization Problem.

We investigated a Cascade chamber that uses a LiAlO₂ breeding blanket and a beryllium oxide (BeO) neutron multiplier.

General Studies

	TART	Estimated
Reactions per DT reaction		
⁶ Li(n,T)α	0.357	0.350
⁷ Li(n,n'T)α	0.698	0.700
T	1.055	1.050
⁷ Li(n,γ)	0.016	0.016
Mn(n,γ)	0.353	0.353
Fe(n,γ)	0.390	0.382
Energy deposition ^a		
Target	1.82	1.84
⁶ Li	1.67	1.70
⁷ Li	9.45	9.32
Structures	6.17	6.03
Mn decay	0.88	0.88
Alpha particle	3.52	3.52
Total	23.51	23.29
Displacement damage rate ^b	9.45	9.50

^aMeV per DT reaction.

^bdpa per full-power year based on 2700 MW of fusion power.

Table 7-21. Comparison of TART and estimated results at $x_1 = 0.03\%$ and $x_2 = 0.91$ m.

Table 7-22. Comparison of HYLIFE and modified HYLIFE parameters at $x_1 = 0.09\%$ and $x_2 = 0.91$ m.

	HYLIFE	Modified HYLIFE
Tritium breeding ratio	1.75	1.05
Neutron energy deposition (MeV)	16.9	20.0 ^a
Fusion power (MW)	2700	2700
Fusion energy multiplication factor	1.16	1.34
Thermal power (MW _t)	3130	3607
Thermal conversion efficiency (%)	39	39
Gross electrical power (MW _e)	1220	1407
Laser power consumption (MW _e)	135	135
Auxiliary power requirements (MW _e)	75	75
Lithium pumping power (MW _e)	30	27
Net electrical power (MW _e)	980	1170
Normalized capital costs		
Reactor	0.62	0.69
Lithium pumps	0.10	0.09
Laser	0.21	0.21
Target factory	0.07	0.07
Total	1.00	1.06
Normalized cost per MW _e	1.00	0.89

^aBased on results of transport calculation.

Beryllium is an excellent neutron multiplier and has a threshold of only 1.85 MeV. Because of its high melting point, BeO will allow high-temperature operation, which is one of the goals of the Cascade concept.¹²¹ We can maintain a flowing layer of BeO granules on the surface of the LiAlO₂ blanket by making the BeO granules larger and less dense than the LiAlO₂ granules.¹²²

Figure of Merit for Cascade. We considered three blanket design variables in the optimization problem:

x_1 = ⁶Li fraction in lithium,

x_2 = LiAlO₂ blanket thickness (m),

x_3 = BeO multiplier thickness (m).

The figure of merit F chosen for Cascade is the sum of x_2 and x_3 , or,

$$F = x_2 + x_3 \quad (55)$$

By minimizing F , we would minimize the size and cost of the rotating chamber.

At the present stage in the development of the Cascade concept, it is not possible to optimize a more general system parameter, such as the cost of electricity, because cost estimates for the reactor plant have not yet been made. The minimum inner radius of the blanket is assumed to be fixed by the damaging effects (ablation and vaporization) of the x rays and target debris.

Design Constraints. Three constraints are imposed on the Cascade design. The first is a requirement for a tritium breeding ratio T greater than 1.05, or

$$T \geq 1.05 \quad (56)$$

The second constraint relates to the mechanical design of the Cascade chamber. GA Technologies has proposed a concept¹²³ in which the rotating chamber is constructed of individual SiC panels held together by tendons composed of aluminum in which SiC fibers are embedded for greater tensile strength. Based on a temperature limit of 400°C for the tendons, the heat-generation rate G in the tendons due to neutron and gamma heating must be less than 0.85 W/cm³ (Ref. 123), or

$$G \leq 0.85 \text{ W/cm}^3 \quad (57)$$

The heat-generation rate in the tendons is calculated from

$$G = P_f(E_f/17.6)/V_1 \quad (58)$$

where

P_f = the fusion power (W),

E_f = energy deposited in the zone representing the tendons per DT reaction (MeV),

V_1 = volume of the zone representing the tendons (cm³).

The total energy released per DT reaction is 17.6 MeV. At the time of this study, P_f in Cascade was 3000 MW.

A third constraint is placed on the total neutron leakage rate L from the Cascade chamber. This constraint indicates the effectiveness of the blanket design in performing one of its primary functions, that of capturing the fusion neutrons.

The beam ports at the ends of the Cascade chamber provide a direct leakage path from the chamber. The two ports subtend 1.25% of the total solid angle. The neutron leakage through the ports, however, will be greater than 1.25% of the fusion neutron source for two reasons.¹²⁴ First, the neutrons entering the blanket can be scattered out through the ports. Second, neutron multiplication in the BeO region will tend to increase the neutron leakage through the ports.

There will also be some neutron leakage through the LiAlO₂ blanket itself. The constraint on the total neutron leakage is set at 0.1 neutrons per DT reaction, or

$$L \leq 0.1 \quad (59)$$

We evaluated each of the three constraints as a function of the three design variables using successive two-point variational interpolation.¹¹⁵

Neutronics Model for the Cascade Chamber. The neutronics model of the Cascade chamber is shown in Fig. 7-39. The double-cone-shaped chamber is approximated by a sphere. The 14.1-MeV neutron source is uniformly distributed throughout a region of DT compressed to a density-radius product of 3.0 g/cm². The target (zone 1) is located at the center of the chamber. The region between the target and the blanket is void.

The innermost blanket region (zone 3) contains the BeO neutron multiplier. Zone 3 is 0.1 m thick and has an inner radius of 3.4 m. The density of BeO within zone 3 is varied to represent variations in x_3 .

The LiAlO₂ breeding blanket (zone 4) extends from a radius of 3.5 to 4.5 m. The material density is varied to represent variations in x_2 . The ⁶Li fraction of lithium in this zone is x_1 .

Outside the breeding blanket is zone 5, which is a 2-cm-thick SiC shell that represents the chamber wall. Zone 6 is a 2-cm-thick region that represents the SiC-fiber/Al-composite tendons. Zone 7 represents

the two beam ports in the blanket. The radius of the beam ports is 0.72 m so that the solid-angle fraction subtended at a radius of 4.5 m is 1.25%.

The volume of zone 6 is $5.09 \times 10^6 \text{ cm}^3$. Thus from Eq. (58), the heat-generation rate in the SiC-fiber/Al-composite tendons for 3000 MW of fusion power is

$$G = 33.5 E_t \quad (60)$$

Table 7-23 lists the composition of the materials used in the neutronics calculations. The tendons are 65 vol% Al and 35 vol% SiC fibers. These fibers contain CO and SiO₂ impurities, which give rise to the indicated oxygen content. The fiber composition is 57 wt% silicon, 31 wt% carbon, and 12 wt% oxygen.

Results of the Initial Transport Calculations. Eight reference-point transport calculations are required for the Cascade chamber optimization problem. The eight points are defined by the combinations of two values for each of the three design variables. The reference values and results are given in Table 7-24, which lists the reaction rates, neutron balance, and energy deposition per DT reaction. The neutron balance indicates the net neutron gain or loss in the various regions of the chamber. The small remainder is the neutron capture in the wall and tendons.

Estimated Neutronic Performance. The neutronic performance is estimated as a function of the three design variables by successive, two-point variation interpolation between values obtained in the eight reference-point transport calculations. We must determine the constraints on the tritium breeding ratio T , the tendon heat-generation rate G , and the total neutron leakage L .

The tritium breeding ratio increases with increasing ⁶Li fraction x_1 and with increasing blanket thickness.¹¹⁵ To meet the constraint of $T \geq 1.05$ with a 0.05-m-thick BeO multiplier, we need a LiAlO₂ blanket thickness greater than ~0.37 m if the lithium is enriched to ~40% ⁶Li. With a 0.15-m-thick BeO multiplier, T exceeds 1.05 over the entire range of ⁶Li fractions and LiAlO₂ blanket thicknesses, except for thin blankets (~0.3 m) with denatured lithium (less than ~3% ⁶Li).

Fig. 7-39. Neutronics model of the Cascade chamber.

Material in zone	Density (g/cm ³)	Isotopic fractions (%)
1	100	D 50.00
		T 50.00
3	3.01	Be 50.00
		O 50.00
4 ^a	2.55	⁶ Li 1.86
		⁷ Li 23.14
		Al 25.00
4 ^b	2.55	O 50.00
		⁷ Li 12.50
		Al 25.00
5	3.20	O 50.00
		Si 50.00
		C 50.00
6	2.60	C 19.71
		O 5.72
		Si 15.49
		Al 59.08

^aNatural lithium.

^bLithium enriched to 50% ⁶Li.

Table 7-23. Composition of materials used in the Cascade neutronics calculations.

Table 7-24 indicates that the energy deposition in the tendons is independent of the ⁶Li fraction. The heat-generation rate G decreases as the thickness of either region increases. To keep $G \geq 0.85 \text{ W/cm}^3$, the required LiAlO₂ thickness is ~0.38 m with

General Studies

Reference point	1	2	3	4	5	6	7	8
BeO thickness (m)	0.05	0.05	0.05	0.05	0.15	0.15	0.15	0.15
LiAlO ₂ thickness (m)	0.30	0.30	0.50	0.50	0.30	0.30	0.50	0.50
⁶ Li Fraction (%)	7.42	50.00	7.42	50.00	7.42	50.00	7.42	50.00
Reactions ^a								
⁶ Li(n,T) α	0.797	0.948	0.999	1.100	1.158	1.217	1.263	1.276
⁷ Li(n,n'T) α	0.070	0.038	0.075	0.040	0.023	0.012	0.023	0.012
Be(n,2n)	0.269	0.266	0.264	0.267	0.562	0.569	0.561	0.556
Neutron balance ^a								
Target	1.057	1.056	1.055	1.057	1.052	1.051	1.061	1.057
BeO	0.189	0.188	0.187	0.184	0.321	0.318	0.325	0.315
LiAlO ₂	-0.893	-1.046	-1.104	-1.179	-1.179	-1.252	-1.290	-1.305
Port leakage	-0.046	-0.033	-0.043	-0.030	-0.057	-0.051	-0.056	-0.050
Blanket leakage	-0.303	-0.161	-0.094	-0.031	-0.136	-0.065	-0.039	-0.016
Remainder	0.004	0.004	0.001	0.001	0.001	0.001	0.001	0.001
Energy deposited ^a (MeV)								
Target	1.80	1.85	1.85	1.83	1.81	1.83	1.84	1.83
BeO	3.05	3.02	3.06	3.03	6.89	6.82	6.94	6.82
LiAlO ₂	9.19	9.98	10.77	11.28	7.77	7.91	8.36	8.38
SiC wall	0.11	0.11	0.03	0.02	0.05	0.04	0.01	0.01
Tendons	0.059	0.059	0.013	0.014	0.026	0.024	0.007	0.006
Port leakage	0.20	0.20	0.19	0.18	0.21	0.20	0.18	0.19
Blanket leakage	0.66	0.65	0.14	0.13	0.28	0.26	0.06	0.06

^aPer DT reaction.

Table 7-24. Results of reference point neutron calculations.

Fig. 7-40. Three constraints as a function of ⁶Li fraction and LiAlO₂ thickness with 0.05 m of BeO.

0.05 m of BeO. The required thickness decreases with increasing BeO thickness to ~0.30 m with 0.15 m of BeO.

The neutron leakage L is calculated from a neutron balance, with the net gain in the BeO region and the net absorption in the LiAlO₂ blanket being estimated by successive two-point interpolation:

$$L = 1.056 + N_g - N_a \quad (61)$$

where N_g is the net neutron gain in BeO and N_a is the net neutron absorption in LiAlO₂. The quantity 1.056 is the number of neutrons emitted by the target per DT reaction; it exceeds unity because of (n,2n) reactions with deuterium and tritium in the compressed fuel zone.

The leakage decreases with increasing ⁶Li fraction and with increasing blanket thickness because in both cases more neutrons

are captured by ${}^6\text{Li}$ as these variables increase. With a 0.05-m-thick BeO multiplier and with the lithium enriched to 80% ${}^6\text{Li}$, a 0.41-m-thick LiAlO_2 blanket is required to keep the neutron leakage below 0.1 per DT reaction. With a 0.15-m-thick BeO multiplier, the neutron leakage is lower. For the 0.15-m-thick BeO multiplier, the minimum required breeding blanket thickness decreases from ~ 0.49 m with natural lithium to ~ 0.33 m with lithium enriched to 80% ${}^6\text{Li}$.

Figure 7-40 shows the three constraints as a function of the ${}^6\text{Li}$ fraction and the LiAlO_2 blanket thickness for 0.05 m of BeO. The hash marks are on the unacceptable side of the constraint lines. Neutron leakage is clearly the limiting constraint over the entire range of ${}^6\text{Li}$ fractions.

Optimal Design Point. The direct search algorithm gives the optimal design point as 63.6% ${}^6\text{Li}$, 0.397 m of LiAlO_2 , and 0.040 m of BeO.¹¹⁵ This result is shown in Figs. 7-41 and 7-42. Figure 7-41 shows the minimum blanket thickness as set by the neutron leakage constraint and the heat generation constraint as a function of the BeO thickness. Note that the neutron leakage constraint goes through a maximum at ~ 9 cm of BeO. The rise and fall in the required blanket thickness is related to the two

modes of neutron leakage. As indicated in Table 7-24, increasing the BeO thickness increases the neutron leakage through the ports and decreases the neutron leakage through the blanket. In some cases, the increase in port leakage can exceed the decrease in blanket leakage. As a result a thicker breeding blanket is required to maintain a constant total leakage rate of 0.1 per DT reaction.

Fig. 7-41. Minimum blanket thickness as a function of BeO thickness. The lower curves show the minimum LiAlO_2 thickness as set by the neutron leakage and heat-generation rate constraints. The top curve is the total blanket thickness including the BeO multiplier.

Fig. 7-42. First estimate of the location of the optimal design point with 0.04 m of BeO. The point lies at the intersection of the neutron leakage and heat-generation rate constraints.

For less than 0.04 m of BeO, the heat-generation rate is the limiting constraint. Between 0.04 and 0.16 m of BeO, neutron leakage is the limiting constraint. Beyond 0.16 m, the LiAlO_2 blanket thickness is again limited by the heat-generation rate. The tritium breeding constraint, which is not shown, only comes into play as the BeO thickness approaches zero.

Figure 7-41 also shows the total blanket thickness. There are two local minima, one at 0.04 m of BeO and the other at 0.16 m. The minimum at 0.04 m, however, is the lower of the two.

The constraints as a function of the ^6Li fraction and the LiAlO_2 thickness are shown in Fig. 7-42 for the optimal 0.04 m of BeO. The optimal design point lies at the intersection of the neutron leakage and heat-generation rate constraints. Higher ^6Li enrichments give the same minimum thickness because the heat rate is independent of this variable. The minimum acceptable ^6Li fraction of 63.6% was chosen from a resource consideration.

Comparison with Transport Calculation at the Optimal Point. We performed an additional neutron transport calculation at 63.6% ^6Li , 0.40 m of LiAlO_2 , and 0.04 m of BeO. We found that the estimated tritium breeding ratio was $\sim 3\%$ higher than the TART result and that the estimated neutron leakage was lower than the TART result by $\sim 6\%$. While these are acceptable differences, the heat-generation rate in the tendons is actually 22% higher than the estimated result.

Fig. 7-43. Final estimate of location of optimal design point.

Improved Estimate of Location of Optimum. There was a significant difference between the estimated heat-generation rate in the tendons and the results of the neutron transport calculation. We postulated that the difference was largely due to the fact that a significant fraction of the energy deposited in the tendons is photon energy that originates in the LiAlO_2 blanket.¹¹⁵ The variational interpolation formula used to predict the energy deposition in the tendons only accounts for variations in the neutron flux in that region; it does not account for variations in the photon source in adjacent regions and the transport of those photons into the tendons.

We used an alternative approach to estimate the heat-generation rate in the tendons as a function of the BeO and LiAlO_2 blanket thicknesses. We assumed that the heat-generation rate in the tendons decreases exponentially with BeO and LiAlO_2 thickness. That is,

$$G = G_0 \exp(-\mu_2 x_2) \exp(-\mu_3 x_3) \quad (62)$$

where

G_0 = heat-generation rate in the tendons with no blanket (W/cm^3),

μ_2 = attenuation coefficient for the LiAlO_2 thickness (m^{-1}),

μ_3 = attenuation coefficient for BeO thickness (m^{-1}).

Using the energy deposition results given in Table 7-24, Eq. (62) becomes

$$G = 27.5 \exp(-7.37x_2) \exp(-8.59x_3) \quad (63)$$

From Eq. (63), we find that the predicted energy deposition for the optimal design point of $x_2 = 0.40$ m and $x_3 = 0.04$ m is 0.031 MeV or 1.03 W/cm^3 . This is exactly the same as the TART result at that point.

We ran the optimization problem again, using the exponential estimate for the heat-generation rate in the tendons. In the rerun, the optimal design point was 34.2% ^6Li , 0.424 m of LiAlO_2 , and 0.042 m of BeO. The constraints and location of the optimal design point for this case are shown in Fig. 7-43.

We performed a final transport calculation at the new optimal point and compared the calculation results to the estimated results in

Table 7-25. The agreement is close for all the relevant parameters.

Summary. The Cascade chamber can be designed with a LiAlO_2 breeding blanket if a BeO neutron multiplier is used. The configuration that minimized the total blanket thickness is 0.042 m of BeO followed by 0.424 m of LiAlO_2 . The lithium must be enriched to at least 34.2% in ^6Li . Because the blanket is a granular bed, the actual thickness is equal to the effective thickness divided by the granule packing fraction. If we assume a 50% packing fraction, the result is a total blanket thickness of 0.93 m.

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Fragmentation of Suddenly Heated Liquids in ICF Reactors

Introduction. In 1984, we proposed four hydrodynamic models to predict the fragment size of liquid metals in ICF reactors.¹²⁵ In this article, we review the predictions from Model 3, the energy-minimization model, and compare them with results from molecular-dynamics experiments. Models 1, 2, and 4 did not predict fragment size accurately, and we refer to these models only to gain physical insight. Detailed information on all four models is available in Ref. 125.

In an ICF reactor, fusion neutrons are suddenly deposited into liquid metals within the vacuum chamber. The neutron energy is deposited so quickly that the liquid has virtually no opportunity to expand during the deposition time; the deposition is effectively isochoric. We consider energy deposition to be isochoric when the deposition time is less than the sound wave transit time of the liquid region. The suddenly increased temperature and energy content of the liquid produces a sudden increase in pressure. Because the free surfaces are exposed to effectively zero pressure, tensile relief waves propagate inward. Tension eventually causes the liquid to cavitate, leading to fragmentation in liquids with available expansion volume. The fragments have velocities due to acceleration by the pressure difference across the tensile wave.

Liquid-metal fragmentation has two effects on ICF reactor design. First, the

Parameter ^a	TART	Estimated
$^6\text{Li}(n,T)\alpha$	1.019 (1.0) ^b	1.029
$^7\text{Li}(n,T)\alpha$	0.040 (1.1)	0.057
Tritium breeding ratio	1.059	1.086
Total neutron leakage	0.097 (3.7)	0.100
Energy deposited in tendons (MeV)	0.024 (3.0)	0.025

^aPer DT reaction.

^bPercent standard deviation.

fragments eventually collide with the reactor structure, causing a pressure load (which can be estimated from the average velocities of the fragments). Second, fragment size determines the liquid surface area available to interact thermally with the superheated gas or plasma produced by the fusion pulse. Because cooling and condensation times depend on the cooling surface area, it is possible that the degree of fragmentation can limit the chamber pulse rate.

Energy-Minimization Model. An energy-minimization model (Model 3) was used to predict mean fragment size. This model minimizes the sum of the surface energy density and the total fragment dilational energy density as a function of fragment size. (Models 1 and 2 predict the fragment size by balancing the fragment surface tension with the isochoric heating, or the fragment surface tension with the dynamic expansion pressure, respectively; Model 4 minimizes the sum of the fragment surface energy and the dilational energy near the fragment surface.)

The energy-minimization process, which is similar to the Griffith theory of fracture,¹²⁶ was proposed by Grady¹²⁷ and extended by Glenn.¹²⁸ Fragment size a is used as an independent variable. The kinetic energy of each fragment is divided into dilation and the motion of its center of mass. The dilational kinetic energy is added to the surface energy, and the sum is minimized with respect to the fragment size. The velocity divergence, $\nabla \cdot \mathbf{v} = -\dot{\rho}/\rho$, is the independent variable in Grady's result,

$$a = \left[\frac{(D+2)D^2\gamma}{\rho(\dot{\rho}/\rho)^2} \right]^{1/3}, \quad (64)$$

where D is the number of dimensions (2 or 3) and γ is the surface tension. (This result is based on a homogeneous expansion, i.e., constant velocity divergence.)

Glenn performed a series of hydrodynamic calculations of isochorically heated spheres in which tensile fragmentation was

Table 7-25. Comparison of TART calculation and estimated neutronic performance at $x_1 = 34.2\%$, $x_2 = 0.424$ m, and $x_3 = 0.042$ m.

Table 7-26. Predicted fragmentation in several ICF reactors calculated from Model 3.

General Studies

Reactor	Fluid	Thickness (cm)	Fragment diameter (cm)	Suddenly added specific energy (J/kg)	Translational specific kinetic energy (J/kg)	Dilational specific kinetic energy (J/kg)
HYLIFE (Ref. 129)						
Inner jet	Li	20	0.21	800 000	15 800	1
Outer jet	Li	20	1.3	50 000	61.6	0.16
HIBALL (Ref. 130), (inner jet, ignoring SiC stock)	LiPb	3	1.8	500	0.14	0.013
Modified Wetted Wall ^a (Ref. 131)	Li	4	2.7	3 500	0.22	0.081
EAGLE (Ref. 132)	Li	0.05	0.14 ^b	3 700	0	0.34
Pulse*Star (Ref. 133)	LiPb	1	0.53	1 000	0.58	0.042

^aLiquid motion is not in diverging geometry.

^bNo breakup predicted by the energy-minimization model.

not allowed. The resulting velocity profile was relatively flat for $r/R > 0.1$, and the velocity divergence was even flatter. Hence, for most of the liquid mass, the velocity divergence was roughly constant (Dv_s/R), where R is the initial liquid radius. Then

$$a = \left[\frac{(D + 2)}{\rho_m} \gamma \left(\frac{Rc}{\Gamma e} \right)^2 \right]^{1/3}, \quad (65)$$

where e is the isochoric heating per unit mass, c is the sound speed, ρ_m is the mass density, and Γ is the Grüneisen constant. Equation (65) represents Model 3.

If the sphere size is chosen to avoid fragmentation ($a = R$), Model 3 agrees well with Models 1 and 2 because all the kinetic energy is dilational. Thus, for practical purposes, the jump-off velocity gives a much better estimate of the available energy for surface formation than does the total added energy. Further, if fragmentation is extensive, only a small fraction of the jump-off kinetic energy is balanced against the increased surface energy; the larger fraction becomes translational energy of the fragments.

We extended Eq. (65) to include the velocity divergence as a function of radial position within the original drop rather than as a single value.¹²⁵ The fragments near the drop center are up to one order of magnitude smaller, and the outermost fragments are up to 31% larger than the prediction based on constant velocity divergence. Even after this correction, Model 3 still assumes local homogeneity within the fragment. Hence, for large fragments (a approaching R), Model 3 continues to break down

because velocity divergence is then not constant within the fragment.

Table 7-26 lists the Model 3 fragment sizes for several ICF reactors and shows that the fragment dilational kinetic energy is small compared to the translational energy. In the low-energy-density reactor concepts, the liquids are predicted to have little or no fragmentation.

Because Model 3 does not consider the dynamic process of fragmentation, we used numerical experiments to study the fragmentation dynamics and test the model results. We used molecular dynamics to track the evolution of up to several thousand liquid atoms following sudden heating.

Molecular-Dynamics Validation of Model 3. Model 3 gave intuitively reasonable estimates of fragment size. However, Model 3 determines only the lowest final energy state of spherical or circular fragments. We tested the ability of Model 3 to predict fragment size as a function of initial pressure (Fig. 7-44) and system size (Fig. 7-45). We compared our Model 3 predictions with molecular-dynamics predictions of fragment size. Our calculations are described in detail in Ref. 125. Here, we present only our overall results.

We performed molecular-dynamics calculations on 2-D circular regions of liquid containing 169, 721, 2611, or 14 491 atoms. (The four system sizes correspond to the number of atoms in regular hexagonal crystals with 8, 16, 30, and 70 rows of particles on a side.) The crystals were melted and brought to equilibrium at a constant density (within a confining wall) at various temperatures between saturation and an energy

level above that of HYLIFE. The confining wall was then removed. During the subsequent fragmentation, we monitored the translational, rotational, thermal, and potential energies, and the population for each fragment.

Figure 7-44 shows the fragment sizes N_c (in atoms) for two 2611-atom computer runs as a function of initial pressure (or heating level). The two groups of symbols (A and B) correspond to separate computer runs at different initial pressures. Each symbol represents a different method of defining fragment size: the largest fragment, the average of the largest fragments comprising 30% of the system population, and the average of the 10 largest fragments. In each case, the raw data have been slightly adjusted to uniform initial conditions (2611 non-monomer atoms, and specified density, sound speed, and surface tension) using Model 3. The only remaining variable is the initial pressure. The molecular-dynamics data fit shows the average slope of the three methods of defining fragment size. Only Model 3 has the correct dependence on initial pressure (slope of N_c vs $P = -2D/3$). The slope of the curve indicates the power-law dependence.

Figure 7-45 shows the fragment sizes for each system size. The data were slightly adjusted to uniform initial pressure, density, sound speed, and surface tension. The only remaining variable is the system size N (number of atoms). The slopes of the molecular dynamics data fit (between the 721-atom and 2611-atom data points) for the three methods of determining fragment size are 0.75, 0.69, and 0.64, respectively.

(The molecular-dynamics data fit shown in Fig. 7-45 is the average of the three slopes.) These values bracket the predicted slope of $2/3$ for Model 3. We did not use the 169-atom data to compute the slope because of minimal fragmentation. The sizes of the fragments in the 14 491-atom system were reduced because of a sloshing motion present in the system at the time the confining boundary was removed. However, when we determined the slopes of the line segments between the 721- and 14 491-atom data, the results were 0.54, 0.50, and 0.41, respectively, for the three methods of determining fragment size. Therefore, the

Fig. 7-44. Fragment sizes (atoms) vs initial pressure for two 2611-atom runs.

Fig. 7-45. Fragment sizes (atoms) vs system size for similar heating levels. (Symbols identical to those in Fig. 7-44.)

molecular-dynamics data show a possible system size power-law dependence of 0.67 with an uncertainty of ± 0.15 . Again, this result is consistent with Model 3 and inconsistent with the other models.

Although the data in Figs. 7-44 and 7-45 support the Model 3 power-law dependence of fragment size on initial pressure and system size, the sizes of the fragments (in atoms) from the molecular-dynamics calculations are about one order of magnitude lower than the prediction

from Model 3. We will need to perform laboratory-scale experiments to determine if this is an effect of the small molecular-dynamics system size or if it extends to experimental and reactor sizes.

Summary. The molecular-dynamics experiments produced reasonable agreement with Model 3. The qualitative dependences of fragment size, system size, and heating level were correct, and the larger fragment sizes were within one order of magnitude

of the predictions. Thus, we conclude that Model 3 is a useful tool for ICF reactor design. Model 3 predicts significant fragmentation only in high-energy-density reactors, such as HYLIFE. Low-energy-density designs, such as Pulse*Star and EAGLE, must provide adequate droplet surface area by means other than the hydrodynamic response to fusion-neutron pulses.

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