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1982 Laser Program Annual Report

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Charles D. Hendricks

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National Laboratory**

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Preface

In this *1982 Laser Program Annual Report*, we present a summary of the unclassified activities and accomplishments of the Laser Fusion Program at LLNL for the calendar year 1982. This Program is supported by the United States Department of Energy.

Our purpose in publishing this report is to present our work to the inertial-confinement fusion community in sufficient depth to provide an understanding of the experimental and analytical bases of the LLNL Program. To keep the report complete in one volume, we have necessarily left out many details that would most certainly be interesting and useful to our readers. We can only suggest that readers desiring more information contact the authors of the various articles to obtain more in-depth information.

We have arranged the 1982 Annual Report in basically the same format used in previous years. Section 1 provides the reader with an overview of the entire Program and some highlights and accomplishments in the seven major functional areas. These functional areas—Laser Systems and Operations, Target Design, Target Fabrication, Fusion Experiments Program, Zeus Laser Project, Laser Research and Development, and Energy Applications—are described in more detail in Sections 2 through 8, respectively.

C. D. Hendricks
Scientific Editor

Acknowledgments

The production of an Annual Report for a rapidly moving program like the Laser Fusion Program is a difficult task. Potential authors for the various articles are all very busy and have other important things to do—some of which must take precedence over Annual Report writing. However, almost without exception, each of the staff members who volunteered, or were “volunteered,” to write an article did so with only a minimum of prodding and coercion. The total response of all the authors is therefore gratefully acknowledged. Special mention must be given to Bill Krue, the Section Scientific Editor for Section 3, Target Design, and the many authors and the clerical staff of the Target Design group, all of whom worked exceptionally hard and delivered most of the articles in Section 3 at a very early date in the preparation process. Indeed, the section on Target Design was the first to have all the articles written.

The clerical staff members in each program element deserve much more recognition than we can possibly provide. They are beset on the one hand by authors who want everything typed instantly once it is handwritten, and who want instant response to corrections, and on the other hand by editors who want the same things in addition to perfect text and instant delivery of multiple copies—and both authors and editors expect cheerful countenances! Our fervent thanks to these super people.

Without the staff of the Technical Information Department and their patience and expert help, this Report would neither have appeared so quickly nor would it be the quality document that it is. From the beginning, Jerry Grow, the Publication Editor, has provided the experience, coordination, background, and expert advice and assistance needed to bring this report together and into publication in a heretofore impossibly short time. Howard Lentzner provided overall guidance in the editorial effort, while the difficult task of editing complex technical articles was accomplished by Phil Anderson, Wally Clements, Thom Elkjer, Steve Greenberg, Bob Kirvel, Pete Murphy, Louise Rufer, and Bill Shissler. The indexes for this Report were compiled with great ingenuity by Thom Elkjer and Peter Link.

Don Cowden, Bill Fulmer, and Wilma Leon planned, guided and coordinated the artwork and layout. Bill Fulmer, ably assisted by Tim Pierce, exercised considerable imagination and creativity in designing the general format and in preparing the divider pages, section introductions, and covers. Wilma Leon, in particular, deserves special mention for the prodigious commitment of her time and talent, seeing the effort through from its inception to final printing. Gary Alexander and his coworkers turned in a magnificent performance in quickly but accurately preparing page layouts for camera-ready copy.

Elaine Price ensured that composition resources remained available despite heavy demands from other Laboratory programs and despite temporary personnel shortages. The flow of work through the Composition group was smoothly and effectively guided by Delores Mason, who also applied her exceptionally accurate and fast proofreading skills to a major part of this Report. Composition entries and corrections were quickly and accurately performed by Stacy Bookless, Cathy Johnson, and Kathy Marshall. Both Peter Link and Cathy Johnson were especially helpful in developing formats for on-line final page makeup, and this difficult, demanding task was ably performed by Stacy Bookless, with assistance from Cathy Johnson and Barbara Phillips. Yvonia Hatcher provided invaluable support in the tedious and exacting job of proofreading.

In the final production phase of this Report, Margaret Sylvester expertly handled development of the printing specifications and did a comprehensive review for format consistency. In turn, Mel Moura displayed his usual excellence in coordinating the printing and binding contracts and in maintaining our high standards of printing quality in the face of an extremely tight schedule.

Other Laboratory organizations provide a great deal of necessary help and assistance in ensuring that the Annual Report not only gets published, but that it is done correctly and carefully. We are indebted to Rick Pond and the Laser Fusion and Special Projects group in the Classification Office, who kept us aware of classification requirements and provided unusually rapid turnaround on their perusal of the articles in the Report. The task of ensuring that the sensitive issues of classification are addressed appropriately is a difficult one, and the help of the Classification Office is gratefully acknowledged.

Acknowledgments

The Patent Office, under the direction of Bert Weis, turned in an admirable performance in swiftly, but thoroughly, reviewing and clearing for public release the material contained in this Report.

The Section Scientific Editors—Ken Manes, Section 2; Bill Kruer, Section 3; Dale Miller, Section 4; Peter Lee, Section 5; Roger Haas, Section 6; Bill Hagen, Section 7; and Wayne Meier, Section 8—were invaluable. They provided section outlines, wheedled and cajoled the authors, coordinated the work with the clerical staff members, and provided very worthwhile advice and counsel during the preparation of the Report. Each of these Section Scientific Editors also did the initial, and very difficult, editing of the section for which he was responsible.

There were several individuals without whose help the Scientific Editor's task would have been much more difficult. Many thanks to Joanne Bell, Hallie Gibson, Patt Page, and Joan Seibel for their help in coordinating and preparing the reports. Sam Tindell's assistance in record keeping, in coordinating the distribution of information, and in following the flow of the report material was extremely valuable. Ollie Parker and the staff of the Multimedia group provided outstanding support to the authors and editors by obtaining prints and artwork, frequently on very short notice.

And, finally, let me thank John Holzrichter, without whose help I would not have been the Scientific Editor of this Report. That help has been a mixed blessing. However, when programmatic decisions were needed to expedite the preparation of this report, I have luxuriated in John's help. Without his hacking, the Report would not have been brought out in as timely a fashion as it has been.

C. D. Hendricks
Scientific Editor

The filter-fluorescer diagnostic is rapidly becoming the standard throughout the ICF community for measuring the high-energy x-ray spectrum from laser-irradiated targets. The thick lead shielding shown here is used to reduce the background from scattered high-energy x rays. This instrument can measure x-ray energies up to 300 keV. From the x-ray spectrum, we calculate the number and energy of suprathermal electrons that could preheat our laser fusion capsules, making them more difficult to compress. Our measurements show that, by using a shorter-wavelength laser, we suppress the production of hot electrons.

Section 1

Laser Program Overview

Laser Program Overview

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Section Editor: C. D. Hendricks

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Laser Program Overview

Introduction

J. F. Holzrichter

The Inertial Confinement Fusion (ICF) Program at the Lawrence Livermore National Laboratory (LLNL) experienced a time of transition during 1982. At the end of 1981, we had turned off the old lasers, Argus and Shiva, to free resources needed for construction of the new lasers, Novette and Nova. In 1982, we installed Novette, the 2-beam forerunner of the 10-beam Nova; we completed the Nova laser building and office building; and we began installation of the Nova laser itself.

The main emphasis of our Program remains: using ICF facilities and concepts to support military and weapons research in the near term and to conduct high-risk, high-payoff research for commercial applications in the long term. We will first use the Novette laser to conduct experiments relevant to the weapons program. Later, using ICF targets, we will obtain information useful for understanding high-gain physics.

The commercial ICF activities are minimal at present. However, during 1982, we continued to emphasize the long-range, high-potential aspects of the ICF process when applied to electrical power production. The entire fusion community (both magnetic and inertial fusion) is becoming increasingly concerned with the costs of fusion systems as compared to the competing technologies. Consequently, the objective of our ICF research for commercial applications is to conceptualize fusion reactor systems that are less costly and more flexible.

Future ICF activities received much attention during 1982. In particular, a new five-year plan was developed by the Department of Energy and the inertial fusion laboratories. This plan clearly shows both that the Novette and Nova lasers are the primary facilities for maintaining ICF momentum during the five-year time period and that a versatile large driver beyond Nova is needed. Addressing this need for a large new driver, we at LLNL are investigating an advanced laser suitable for a 5- to 10-MJ high-gain test facility. In addition, we have developed a program strategy for extending target performance from the 100-kJ Nova level to the 5- to 10-MJ high-gain level.

Highlights of our ICF activities during 1982 are described below by program element

• Laser Systems and Operations	• Zeus Laser Project
• Target Design	• Laser Research and Development
• Target Fabrication	• Energy Applications
• Fusion Experiments Program	

Laser Systems and Operations

The most important event of 1982 was the completion of both beamlines of the Novette laser (Fig. 1-1). This laser, designed as an interim system until the 10-beam Nova can be completed, is built up from borrowed Shiva parts, some special parts,

and large-aperture Nova components. In essence, Novette comprises 2 of the ultimate 10 beamlines of Nova.

In 1982, Novette achieved a performance of 11 TW with 1ω ($1.05\ \mu\text{m}$) light in one beam at a pulse width of 94 ps. In January 1983, the full two-beam Novette system provided 25 TW with 1ω light and 13 TW with 2ω ($0.53\ \mu\text{m}$) light, both at a 94-ps

Introduction

pulse width. This performance verifies the basic design goal of the Nova laser, which is 8 to 12 TW of 1ω light per beam. Furthermore, the beam has a fill factor of 0.85, which is even greater than our original estimate of 0.70.

At the end of 1982, diagnostics had been installed on Novette, and experiments were beginning. Preliminary experiments on the

conversion of 1ω to 2ω light had been carried out with a measured conversion efficiency of $\sim 60\%$.

The final configuration of the Nova laser was approved in April 1982. Nova is to provide approximately 100 kJ at $1.05\ \mu\text{m}$ and slightly less energy at 0.53 and $0.35\ \mu\text{m}$. The total estimated cost for Nova is \$176 million, with construction to be

Fig. 1-1. View of Novette laser.

Fig. 1-2. Nova office building provides space for 200 personnel.

completed by the end of FY-84. The Nova office and laboratory buildings (Figs. 1-2 and 1-3) have been finished and accepted, and the Nova laser system is presently being installed.

Target Design

The Target Design program is concentrating its resources on obtaining the greatest performance from targets designed to be irradiated by Novette and Nova. Test results from these targets will allow us to project, with added confidence, to the design of high-gain targets that are expected to be driven by multimegajoule lasers. In particular, we are preparing to test target designs for Nova that, hydrodynamically, will be very nearly equivalent to our high-gain multimegajoule targets.

In our direct-drive designs for Nova, we are concentrating on a single-shell design, with a double-shell design as the principal backup. Of these two designs, the single-shell design has a lower risk and may be more easily explored.

For these targets, we are emphasizing a target made up of cryogenic deuterium-tritium (DT) fuel and low-Z ablaters, as shown in Fig. 1-4. In addition, we are directing our diagnostic and target-fabrication programs to provide techniques for measuring the size of the final

compressed-fuel region and for verifying that central ignition in these targets will occur as required.

LASNEX, the principal target-design code, was significantly improved during 1982 by the inclusion of improved energy transport and charged-particle deposition. LASNEX is playing a key role in the design of various x-ray lasers, several of which will be tested in Novette. Additionally, our effort in plasma physics resulted in a self-consistent kinetic theory of classical light absorption, in improved LASNEX models for Raman backscatter, and in the development of improved plasma-simulation codes.

Target Fabrication

During 1982, we completed the fusion target development facility (FTDF).¹ The FTDF, a 47 500-ft² building, was constructed at a cost of \$72/ft² by using local construction techniques developed for light industry; this cost is about one-half that which would have been incurred by using traditional construction methods. With completion of this building, we are now able to concentrate our target fabrication activities in one place.

An important part of the effort of the Target Fabrication program is the development of special cryogenic-fuel/low-Z ablator targets required for the Nova program.

Fig. 1-3. Nova laboratory building provides 115 000 ft² of high-quality space for the laser, target system, and supporting functions.

Fig. 1-4. Typical Nova direct-drive target showing initial and compressed (~ 1000 times solid density) fuel conditions.

Introduction

During 1982, we directed major efforts to the formation and diagnostics of cryogenic layers. In the past, we used simple refreeze techniques to form layers of cryogenic DT in the targets; however, these techniques cannot reliably provide the thicker DT layers required by direct-drive targets designed for Nova (and larger drivers). We have also built a completely different series of targets for the special application of demonstrating an x-ray laser.

Fusion Experiments Program

During 1982, the major task of the Fusion Experiments program was to build the Novette target-diagnostic systems. This task required us to refurbish most of the Shiva target diagnostics and to bring up the Nova data-acquisition system to operate on Novette. In addition, we developed several specialized diagnostics, such as time-resolved x-ray spectrometers for low-energy x-ray laser studies and for subsequent studies of laser-plasma interactions on long-scale-length plasmas.

To enhance our rate of data analysis, we completed the fusion experiments analysis facility (FEAF), shown in Fig. 1-5. This facility is a data-acquisition system built around a dedicated VAX 11/760, which directly

accesses the Novette (and eventually Nova) instruments through a one-way fiber-optic link so that data analysis can begin immediately after a shot. We designed this facility to reduce the time required for data analysis so that experimenters can keep up with the laser shot rate.

Zeus Laser Project

We feel that attainment of an affordable 5- to 10-MJ single-pulse test facility is the most important next step in the effort to develop the ICF process. Consequently, we have concentrated our efforts on the design of a 10-MJ, 1000-TW laser, named Zeus, that we would like to construct for approximately \$250 million. Several new architectures have been invented as candidates for the Zeus system, including a ring architecture, a multipass architecture, and a harmonic-conversion architecture using a solid-state amplifier.

The Zeus laser system would provide high-gain output for weapons-effects simulation and weapons-physics studies and would provide invaluable information for commercial ICF, specifically, curves of target gain vs driver energy. A detailed understanding of the gain-vs-energy curves shown in Fig. 1-6 will be decisive in determining

Fig. 1-5. Fusion experiments analysis facility (FEAF) provides immediate target analysis via one-way fiber-optic link from diagnostics.

the future of commercial ICF, particularly in regard to driver development, reactor critical cost, and repetition rates.

Laser Research and Development

As part of our research and development effort in 1982, we have continued to firm up our data base for future laser systems. In particular, we have determined that the nonlinear index of refraction in fused silica is the same at 3ω as at 1ω and is low enough to permit the anticipated operation of Nova at 3ω ($0.35 \mu\text{m}$).

We have also made a series of damage measurements on materials used in the harmonic-conversion system. Particularly exciting has been the performance of the neutral-solution-processed optics that have been installed on the Novette laser system. Optics produced by this technology are wavelength insensitive and are at least twice as damage resistant as thin-film antireflection coatings.

Energy Applications

During 1982, we developed three new fusion reactor concepts, all more adventurous than HYLIFE, our previous baseline concept. In developing these new concepts,

Fig. 1-6. Target gain vs input energy curves relating to high-gain applications. Arrows show directions of target exploration.

which we call Pulse*Star, Sunburst, and Cascade, we attempted to test and define the asymptotic potential of fusion power plants, rather than limiting ourselves to using currently available technology. Although we are exploring new territory in fusion reactor concepts, our basic objectives remain the same: to produce fusion power plants that are economically and environmentally competitive with alternative types of power plants (e.g., coal-fired and fission).

In addition to our basic conceptual studies on advanced fusion power plants, we conducted both a trade-off study of the reactor implications of direct-illumination geometry and a commercialization study that produced a success-oriented plan for

Programmatic Resources

the development of fusion power. We also performed subsidiary studies on methods of calculating induced radioactivity in reactors containing many diverse materials and on the various interactions of liquid metals with reactor structures and materials.

Programmatic Resources

Resources employed in the ICF Program are categorized as operating funds, equipment funds, and work force, with the latter expressed in full-time equivalents (FTE). Figure 1-7 shows a 10-year history of these resources for the ICF Program.

Operating funds for FY-82 were \$52.0 million, down from \$57.5 million in FY-81. Although these funds nominally declined 10% in terms of dollars, they represented a decline of 20% in terms of purchasing power when inflation is taken into account. The FY-82 operating funds were distributed as shown in Table 1-1. Equipment funding

and the work-force level for FY-82 were also down about 10% each from the previous year: equipment funding was \$4.3 million, down from \$4.7 million; and the work force was 315 FTE, down from 354 FTE. In spite of this reduction in resources, the flexibility inherent in our philosophy of maintaining a high percentage of procurement funds relative to work-force expenditures enabled us to finish construction of the Novette laser on schedule.

We received an appropriation of \$37.5 million for the Nova project during FY-82, bringing the total Nova funding thus far to \$141.5 million. The project scope has been changed from a 20-beam, 1.05- μm system to a 10-beam, frequency-converted system, the total project cost of which will be \$176 million (rather than \$185 million if final funding had been spread over two years). During 1982, we completed construction and took beneficial occupancy of both the Nova office building and the Nova laboratory building.

We received the final \$1.1 million appropriation for the fusion target development facility (FTDF), the total project cost of which was \$7.6 million. We have substantially completed the FTDF and have begun working in this well-equipped facility.

Even though equipment funding declined 10% in FY-82, we were able to procure the necessary tools for diagnostics, target

Budget category	Operating funds (thousands of \$)	Work force (FTE)
Laser development	16 119	70
Fusion theory and experiments	35 856	245
	51 975	315

Table 1-1. Distribution of FY-82 operating funds and work force.

Fig. 1-7. Ten-year history of ICF Program resources.

fabrication, and laser research and development. We also acquired a computer-aided design system, which will significantly improve the productivity of our design and drafting efforts.

Author: L. P. Altbaum

Programmatic Facilities

Major construction work completed for the ICF Program in 1982 includes the new Nova office building (481) (59 000 ft²), the Nova laboratory addition (115 000 ft²), the fusion target development facility (FTDF) (47 500 ft²), and the fusion experiments analysis facility (FEAF) (2500 ft²). We installed new facilities for a computer-aided design/drafting system (CADD) in the Laser Fusion office building (381), including new rooms for the CADD computer, for associated work stations and drafting tables, and for a conference area. To improve internal communications, we installed LIDNET (Laser Information Distribution *NET*work)—a system of 18 strategically located color TV monitors linked to a central computer—to present information of Program-wide interest throughout the two office buildings (381 and 481).

Other major facilities work that we completed during 1982 includes

- Relocation of the Inner Loop Road about 30 ft south of its original location to accommodate the Nova office building.
- Renovation of the D-7 parking lot adjacent to the Nova office building.
- Addition of paving and grading around the Program's machine shop, and completion of the adjacent lunchroom/vending-machine area.
- Addition of a corporation yard on the Outer Loop Road to support the Nova laboratory.

Type of facility	Building space (ft ²)		Change (%)
	1981	1982	
Permanent offices	50 930	110 305	+116.6
Owned office trailers	74 007	74 007	0
Rented office trailers	28 365	23 565	-16.9
Permanent laboratories	207 806	371 591	+78.8
Technician support areas	49 949	52 949	+6.0
Total	411 057	632 417	+53.9

Table 1-2. Distribution of building space in the Laser Program.

Landscaping	1981	1982	New total	Increase (%)
Trees (No.)	4 651	3 299	7 950	71.0
Shrubs (No.)	796	1 717	2 513	215.0
Lawn area (ft ²)	50 031	160 441	210 472	320.0
Other cover (ft ²)	98 750	94 120	192 860	95.3

- Design and construction of laboratory and technician support rooms in the basement of the building 381 high-bay to support the Novette operating groups and in the basement of the Nova laboratory building to support the Nova/Novette Power Conditioning Group.
- Design and construction of facilities in the basement of the Nova laboratory building for neutral-solution processing of Novette and Nova optics.

Table 1-3. Landscaped area in the Laser Program.

Table 1-2 summarizes the relative distribution of building space in 1981 and 1982. The current total of 632 417 ft² represents an increase of 53.9% over 1981.

In addition to completing the several major building projects and site improvements during 1982, we more than doubled the area of ground cover and landscaping. We installed new landscaping in the areas around the electrooptical facility, the FTDF, and building 481, and we restored the landscaping in the area between buildings 381 and 481. Table 1-3 summarizes the areas of landscaping cover in 1981 and 1982. The present landscaped area of 413 805 ft² represents an increase of 168.3% over 1981.

Author: R. A. McCardle

1. 1981 Laser Program Annual Report, Lawrence Livermore National Laboratory, Livermore, Calif., UCRL-50021-81 (1982), p. 1-13.

Nova

● Phase-measurement equipment, designed and built by LLNL, that is compatible with the vendor's machines and tooling

nique uses absorption, and the second uses scattering, of the laser beam. In the former, we place an absorbing apodizer at the input

Section 2

Laser Systems and Operations

Laser Systems and Operations

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Section Editor: K. R. Manes

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Laser Systems and Operations

Introduction

R. O. Godwin

The year 1982 was a very exciting time for construction and activation. The Novette laser was completely installed and Nova's project uncertainties were resolved. The major events associated with these two laser systems are summarized in the following paragraphs.

- Novette—At the beginning of the year, the Novette laser existed only as a few components and as a construction schedule for the bare building that formerly held the Argus laser system. By the year's end, all of Novette's laser hardware had been installed (Fig. 2-1), and both beams had been operated at $1\text{ }\mu\text{m}$ with a 100-ps pulse width. (Novette incorporates two 46-cm amplifier heads, rather than three, as does Nova.) One of the laser beams demonstrated an output power of 11.1 TW, near the upper end of the projected performance for Nova. This same beam was also frequency-doubled by a 5×5 array of potassium dihydrogen phosphate (KDP) crystals in early harmonic-conversion tests. Novette's laser staging is identical to Nova, except for some slight physical differences arising from fitting the beams into the former Argus building. In addition, the Novette laser is only capable of frequency-doubling instead of both doubling and tripling.
- Nova—The uncertainties associated with the scope of the Nova project were resolved in April 1982, when the 10-beam, $1.05\text{-}\mu\text{m}$ Nova laser system (previously called Phase I) was authorized with the added capability for frequency-doubling (2ω) and tripling (3ω). Late in the year, the project's cost and

Fig. 2-1

schedule uncertainties were also resolved when Congress appropriated funds to complete the project. Although the full 20-beam system with the 2ω and 3ω capability was not approved, Nova became a much more versatile experimental facility.

The performance expected from the system is

- Powers of from 80 to 120 TW in 100-ps pulses at $1.05\ \mu\text{m}$.
- Energies of from 80 to 120 kJ in 3-ns pulses at $1.05\ \mu\text{m}$.
- Energies of from 50 to 80 kJ, with frequency conversion to the second harmonic ($2\omega = 0.53\ \mu\text{m}$).
- Energies of from 40 to 70 kJ, with frequency conversion to the third harmonic ($3\omega = 0.35\ \mu\text{m}$).

Because the project is now fully funded, Nova's completion has been accelerated to the fourth quarter of FY-84, at a total cost of \$176 million. By the end of 1982, approximately 60% of the laser hardware had been procured on schedule and within budget. Timely and successful activation of Novette gives us further confidence in our ability to meet the accelerated schedule.

Both the Nova laboratory building and the Nova office building were completed in midyear. The laser fusion staff has moved into the office building and laser installation has started in the laboratory building.

Following completion of the Novette construction, a Novette operations group was formed. Personnel not required to operate Novette will help to construct Nova. As a result of our demonstrated performance in the activation of Novette, and because the funding uncertainties associated with the Nova project were eliminated, we are confident that Nova will be completed on schedule, on budget, and will meet or exceed its projected performance specifications.

Nova

Laser Design and Performance

As far as the amplifier chains and their component configurations are concerned, the Nova laser design remains as it has been described in previous Annual Reports.^{1,2} Significant advances have been made in

- Evaluating laser performance at 1ω .
- Identifying and measuring the optical damage thresholds and temporal behavior of components at 1ω , 2ω , and 3ω for critical items in the Nova chain.
- Solidifying the baseline designs for frequency conversion and target focusing.

The following paragraphs discuss these items, as they pertain to Nova's performance. This article concludes with a summary of the construction status of Nova at

the end of 1982; details of the Nova construction effort are presented in subsequent articles.

In the current Nova chain design, the components most likely to suffer damage are

- Surfaces of spatial-filter entrance lenses.
- High-reflectivity coatings on turning mirrors.
- Potassium dihydrogen phosphate (KDP) crystals.
- Fused silica antireflection surfaces on the focusing lenses and crystal-array windows.

New ways for protecting these optics were developed in 1982. A proprietary graded-index treatment of silicate glass surfaces, known as neutral-solution (NS) processing, has been put into practice by the Schott Optical Co., at Duryea, Penn., and implemented for large-optic processing at LLNL. The NS process has been (and will continue to be) used as an antireflection surface treatment for spatial-filter lenses, since it is highly resistant to damage.^{3,4} Nova's turning mirrors are coated with dielectric oxide films (alternating between high- and low-refractive-index-film layers) to give a high reflectivity at 1ω . A process called "laser

Table 2-1. Seventy-percent values of the best damage thresholds at 1ω .

Pulse duration (ns)	0.1	1.0	3.0	5.0
NS-processed spatial-filter lenses	5.3	17	30	38
High-reflectivity turning mirrors	3.1	10	15	19
KDP (bulk)	3.1	10	17	22
Potator-polarizers	1.4	4.5	5.6	6.3

hardening" can give the mirror coatings and KDP crystals an improved resistance to fluence damage. Fused silica surfaces must be made antireflective at both 2ω and 3ω , and they must also possess a relatively high resistance to laser-induced damage. Unfortunately, NS processing is not applicable to fused silica. However, a graded-index surface treatment, which is currently being developed, shows great promise.

For all of these processes, the damage threshold's dependence on pulse duration is well approximated by a square-root law, as long as the surfaces are kept scrupulously free from absorbing contaminants. Table 2-1 lists values that are 70% of the best measured damage fluence, taken with the so-called "n-on-1" data technique (also called "laser hardening").³ These data are typical of similarly prepared surfaces for pulses ranging from 1 to 20 ns in duration.⁴⁻⁶ In general, the time-dependence data appear to fit a square-root power law; the entries for 0.1 ns were extrapolated using this rule.

The MALAPROP system-simulation code has again been used to assess the 1ω -performance of one of Nova's chains as a function of pulse duration. The assumptions made for the calculations were

- A temporally Gaussian pulse starts down the chain.
- There is an effective geometric amplifier filling factor of 0.85 (consistent with early Novette data).
- The amplifier gains are as shown in Table 2-2.
- The performance is limited by a damage threat to one or more of the components listed in Table 2-1.

The predicted maximum allowable laser energy per arm, for various pulse durations, is shown in Fig. 2-2.

We also used these analytical methods to evaluate Nova's performance with the additional 31.5- and 46-cm amplifiers, for which space is allowed in the baseline design. Our preliminary results indicate that, with pulses of 3 ns or longer, adding one 46-cm amplifier will give us 2 to 3 kJ more per arm.

Earlier,^{1,2} we had anticipated that the small-signal gain of the 46-cm amplifier would be about 10% higher than that shown in Table 2-2. We are currently conducting experiments to see if the performance can be improved. The most promising modification to be tested is NS processing for the clear-glass shields within

the optical cavity, which separate the disks from the flashlamps.

The design for the KDP frequency-conversion array is optimized for a fundamental drive intensity of 2.5 GW/cm^2 , spanning Nova's pulse-width range of from 1 to 3 ns. This operating range is consistent with other system constraints, such as those imposed by nonlinear propagation⁷ and by material fluence damage limits. A good dynamic range for conversion⁸ is achieved by optimizing all the crystal lengths at 1.0 cm for the input intensity range of interest. Further elaboration of the array design appears in the following article.

In Fig. 2-3, the 1ω , 2ω , and 3ω performance expectations for a single Nova chain are shown over the full temporal range of interest. It is possible to operate the system without component damage in regions below and to the left of each curve. The 1ω curve is representative of the laser drive levels that are available at the frequency-conversion array. The array's performance at 3ω is power-limited by the nonlinear growth of spatial irregularities as the beam propagates through the fused silica window and focusing lens.

The array and focusing lens are mounted on the target chamber, where the lens also serves as the vacuum barrier, as previously

Fig. 2-2. The energy (per arm) from the Nova baseline chain as a function of pulse duration.

Table 2-2. Component gain data.

Amplifier	MALAPROP small-signal gain	Novette nominal gain
Rod (5 cm)	25.0	25.9
Disk (9.2 cm)	6.6	6.58
Disk (15 cm)	4.2	4.2
Disk (20.8 cm)	2.3 ^a	2.0
Disk (31.5 cm)	1.85	1.82
Disk (46 cm)	1.75	1.75

^aNova's disks are 3.0 cm and Novette's disks are 2.5 cm in thickness.

described.⁴ The cross hair and retroreflector used as alignment aids are both retracted from the beamline during shots. The optical train for frequency conversion and beam focusing is very simple, and is shown schematically in Fig. 2-4. A dichroic beam dump transmits only the particular wavelength that is desired for each experiment. The return-beam filter absorbs 3ω light, protecting the optics and turning mirrors from solarization. Dispersion in the fused silica focus lens causes the three wavelengths to focus at different distances from the lens, as shown. Therefore, a target at or near the focal position for the desired wavelength lies in the shadow of the beam dump for the

remaining wavelengths. Care must also be exercised in positioning the crystal array and the focus lens relative to one another, to avoid "ghost" foci (back reflections from the lens surfaces) within the crystals. At present, the spacing between the lens and array is about 1.0 m. In fact, the primary ghost focus from the target lens may be useful for measuring the beam energy incident on the target (using a small calorimeter to collect the light). A shield is also needed to protect the focus lens from debris coming from the disintegrating target. (Otherwise, after only a few shots, the debris would seriously degrade transmission through the lens.)

Since Nova's target chamber was designed to accommodate 20 beams and we are only deploying 10 of them, we have seized the opportunity to mirror the east and west beam clusters so that the beams do not radiate into each other through the target chamber. All the beams in a cluster are equally spaced in angle upon the surface of a 100° cone, whose vertex is at the target. This configuration will allow us to use a full complement of experimental and diagnostic instruments. The five-beam overlap spot in the common focus will not exceed $250\text{ }\mu\text{m}$ in diameter, including allowances for alignment, positioning, and verification tolerances. This criterion applies for a lens with a focal length of 3.0 m.

The Nova target chamber was built and tested by the Chicago Bridge and Iron Company, of Chicago, Ill. In Fig. 2-5, it is shown ready for installation. This aluminum chamber, 2.3 m in radius, has 12.5-cm-thick walls to allow the mounting of components without misalignment resulting

Fig. 2-3. 1ω , 2ω , and 3ω performance limits for a single chain from Nova. (The 1ω curve is from Fig. 2-2.)

Fig. 2-4. Optical train for frequency conversion and beam focusing.

from deflections or strains. Aluminum was chosen because of its rapid recovery from radioactivity following a high-yield target shot.⁹ The target chamber has been fabricated, accepted as vacuum-tight, and delivered to LLNL.

The Nova laboratory and office buildings were completed in June, 1982. Installation of the tubular steel spaceframe supports for the laser components, the turning mirrors, output-sensor packages, and the target chamber is nearing completion. More than 60% of Nova's hardware procurements is currently under contract and in fabrication, and the assembly of components is well under way.

Power-circuit construction is approximately at the halfway point. The alignment, laser diagnostics, and control subsystems for Nova are extensively deployed in support of Novette.¹⁰ This two-beam Nova prototype employs Nova hardware wherever possible. Besides functioning as a test-bed laser system, Novette is also performing selected advanced target experiments. Novette's activation has been an extremely valuable learning experience for the coming activation phase of Nova. We confidently expect that Nova will meet its cost, schedule, and performance goals.

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Frequency Conversion and Focusing System

Nova Baseline. The Nova frequency-conversion and focusing-optics baseline design was further refined during 1982, and its present form is shown in Fig. 2-4. It now consists simply of a $2\omega/3\omega \times 3 \times 3$ KDP crystal array, described previously,¹¹ and a single fused-silica, coma-corrected aspheric lens (see the following article). This simplification was made possible by removing the requirement for a full-aperture diagnostic splitter, which was included in the earlier baseline (double lens) design.¹¹ Although implementing an accurate energy diagnostic will be more difficult, it still appears possible, and several options are being evaluated (see "Nova Alignment and Diagnostics," later in this section).

The beam dump has also been greatly simplified, in that a full-aperture absorber is

not required. Dispersion in the focus lens causes different wavelengths to be focused at different distances from it. Therefore, a target at or near the focal plane for the desired wavelength lies in the shadow of the small beam dump. (See Fig. 2-4.) The beam-dump material is chosen to transmit the desired harmonic wavelength.

The new baseline lens design provides a 30% reduction in the amount of glass through which the harmonic beam must propagate, thereby allowing a 30% increase in the maximum harmonic intensity on the targets. Optimizing the performance of Nova's frequency-conversion crystal arrays for the higher intensities is accomplished by selecting a crystal thickness that matches the maximum obtainable 3ω intensity, with this intensity being determined by nonlinear propagation effects and component damage thresholds. However, if the crystal thickness is much less than roughly 1 cm, the maximum achievable 3ω intensity occurs at a fundamental input intensity well outside the range of operational interest (1.5 to 3.5 GW/cm²). This range corresponds to

Fig. 2-5. Nova target chamber, ready for mounting on the spaceframe when delivered to LLNL.

Fig. 2-6. Harmonic intensity vs fundamental intensity for flat temporal and spatial profiles.

Fig. 2-7. Exploded view of the Novette second-harmonic KDP array.

(i.e., with greater than a 3-ns pulse width) would be significantly reduced.

We have chosen to optimize the array's performance for 3ω because of our present programmatic emphasis on the UV wavelength, and because of the higher sensitivity of Cascade 3ω generation^{12,13} to crystal thickness, L , as shown in Fig. 2-6. To take full advantage of the performance improvement that is available with the new baseline design at 3ω , we have chosen a crystal thickness of 1.0 cm, thereby achieving high performance over the anticipated operational range. Our analysis also indicates that, in this intensity range, the frequency-doubling and mixer crystals can have the same thickness with no performance penalty.

The performance of quadrature 2ω generation¹⁴ (also shown in Fig. 2-6) is significantly less sensitive to crystal thickness and remains efficient using the optimum 3ω values. Although better 2ω performance is possible with a crystal thickness optimized for the quadrature design, the resulting 3ω performance would be significantly degraded. The new baseline system performance at 2ω and 3ω is summarized in Fig. 2-3.

Novette KDP Crystal Arrays. The development and use of KDP crystal arrays have been prompted by the limited size of suitable frequency-conversion crystals. These arrays are precision optical assemblies, consisting of a matrix of KDP crystal segments placed edge to edge and mounted

the maximum Nova 1ω output intensity in the 1- to 3-ns pulse-width region. In addition, as the crystals are made thinner, the 2ω and 3ω performance at low intensity

between two windows in a housing filled with an index-matching fluid, as shown in Fig. 2-7. The array must, as nearly as possible, emulate the performance of an equivalent full-aperture single crystal. Previous reports^{15,16} have summarized the development and performance of small (15-cm aperture) crystal-array prototypes. We report here on the development and performance of two full-scale, 74-cm-aperture, KDP second-harmonic-generating crystal arrays, which were recently installed on the Novette laser system. These two arrays are identical, except for the size of the individual crystal segments. The first array is composed of a 5×5 matrix of 15-cm KDP segments, the second is a 3×3 matrix of 27-cm segments (similar to the Nova design). The limited availability of crystals during the accelerated Novette schedule necessitated our using the two different geometries.

The crystals for the two Novette arrays were selected from a group of thirty-four $15 \times 15 \times 1.8$ -cm and twelve $27 \times 27 \times 1.8$ -cm precision-oriented and diamond-machined KDP blanks. These crystal segments were oriented to within $\pm 30 \mu\text{rad}$ of the optimum 2ω -generating phase-match angle and were machined to accuracies of within $\pm 1 \mu\text{m}$ in thickness and $\pm 0.1 \text{ mm}$ in the other dimensions.

The design for the Novette arrays uses an internal support structure, within which the individual crystal segments were nested. The aluminum "egg crates" were diamond machined at the Oak Ridge Y-12 facility, are parallel to within $\pm 1 \mu\text{m}$, and are precisely $10 \mu\text{m}$ thicker than the (Novette 1.8-cm-thick) KDP segments.

When assembled, the egg crate/crystal matrix is sandwiched between two BK-7 windows, with a thin ($5 \mu\text{m}$) layer of halocarbon-56 index-matching fluid in the crystal/window gap. This fluid interface was found to be especially vulnerable to laser damage. Special clean-room handling and assembly techniques were developed to achieve acceptable damage threshold levels. A partial vacuum in the array forces the windows to register against the egg crate, thereby resulting in a stiffened structure capable of producing the required optical figure. The individual crystal tilt angles are

restricted by the proximity of the windows, yet the crystals are isolated from undesirable stresses by the egg crate.

A unique elastomer bonding technique was developed to mount the high-aspect-ratio (90-cm-diam by 4.3-cm-thick) array windows in a virtually stress-free condition, while simultaneously providing the required fluid and vacuum seal.¹⁵ The completed Novette 3×3 crystal array is shown in Fig. 2-8.

The two Novette arrays have been subjected to a series of optical and laser system performance tests, including full-aperture, transmitted wave-front interferometry and 2ω conversion-efficiency measurements. The interferometer measurements confirmed that the egg-crate structure employed in the arrays resulted in an adequate optical

Fig. 2-8. The assembled Novette 3×3 crystal array.

figure. Only one exception was noted. In the 5×5 array, a minor flaw resulted in an increased distortion near the window edge.

Near-field burn photographs and conversion-efficiency data taken on Novette (see "Test and Activation Results" at the end of this section) have confirmed that each of the crystal-array segments generates 2ω light at (or very near) the maximum predicted value.

Fig. 2-9. A Novette 2ω near-field burn pattern.

Fig. 2-10. Prototype 15-cm-diam discrete support structure for KDP crystal arrays.

Apodization. Apodization is required to reduce unwanted diffraction peaks from the hard-edge obscurations formed by the crystal interstices. A design for producing a high-damage threshold scattering apodizer was developed.¹⁷

To make the apodizers for Novette, we designed and built a special x-y stage for generating narrow diffuse stripes of adjustable width and controlled edge roll-off. The setup consisted of a rotary table for holding the 90-cm-diam apodizer substrates and a motor-driven track that traversed above the substrate. A high-pressure nozzle was mounted on the traversing stage and used to spray the surface with a fine abrasive. It can be adjusted in height, according to the roll-off pattern that is desired. By appropriately combining the transverse track position with table rotations, the apodizing pattern for either a 3×3 or a 5×5 array was produced on the substrate. Close examination of a Novette near-field 2ω burn pattern, such as that shown in Fig. 2-9, shows the effectiveness of these apodizers. The wide vertical obscuration comes from a 46-cm split-disk transmission-apodizer (now replaced by a scattering apodizer), located before the 46-cm amplifier. Diffraction peaks from the crystal interstices have been greatly suppressed and are no longer a problem.

Nova Arrays. The Novette array design demonstrates that good optical quality and conversion efficiency can be achieved with large-aperture crystal arrays. However, the diamond-machined aluminum egg crate is expensive, and it must be shielded from the incident laser light. This latter requirement necessitates our using a beam mask and a remotely located apodizer. Other crystal/window support structures that address these issues are being evaluated. One such design uses a "discrete" support structure, similar to that shown in Fig. 2-10. This small-aperture (15-cm) testing prototype produced a good reflection wave front, as shown in Fig. 2-11, and is in close agreement with the predictions from a finite-element analysis of the structure. Our calculations indicate that this configuration is a viable candidate for the 74-cm Nova array. The inexpensive, discrete glass pillars provide an internal support structure with a high damage threshold, which removes the beam mask requirement and allows the

apodizer pattern to be applied directly to the array input window.

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Optical Components for Nova

Laser Glass. Hoya Optics, Inc., Schott Optical Glass, Inc., and Kigre, Inc. are supplying the phosphate laser-glass disks for Nova. By the end of 1982, of the 545 Nova disks, 294 (about 54%) had been melted, edge-clad, and pre-finished, and 144 (about 49%) had been optically finished by Eastman Kodak, Inc. and Zygo Corp. The disks are meeting or exceeding specifications.

Fused Silica. Fused-silica lens and window blanks will be required on Nova for 3ω operation. Fused silica is the only material that has a high enough transmission

(>99% internal) for 3ω light. The windows in the 90-cm KDP array, and for the 80-cm, f/4 focusing lens in the target chamber, will also be made of fused silica.

Prior to 1982, the largest available fused-silica blanks were 50 cm in diameter. Late in 1981, working with Corning Glass (in Canton, N. Y.) and Heraeus Quarzschmelze (in Hanau, West Germany), LLNL embarked on a program to produce 90-cm-diam test blanks 4-cm thick, with good homogeneity and low inclusion densities. In the past, these two parameters have been the factors limiting the production of large fused-silica components. The goal of each vendor was to produce blanks having a homogeneity $\Delta n \leq \pm 2 \times 10^{-6}$ and no bubbles and inclusions exceeding a cross section of 0.03 mm^2 per 100 cm^3 of glass volume.

Both manufacturers had to improve their respective processes to produce test blanks that met all aspects of our specifications. In June, the first of the 90-cm-diam by 4-cm-thick blanks was inspected in Germany, and was found to have exceeded our specifications. (See Table 2-3.) By August, Corning had also met our requirements. This great progress, in a little over a year of intensive effort on the part of both vendors, is

Contour values (μm)

A = 0.40	F = 1.98
B = 0.717	G = 2.30
C = 1.03	H = 2.62
D = 1.35	I = 2.93
E = 1.67	J = 3.25

mks

Fig. 2-11. Test results for the discretely supported test array with tightened bolts and HC-56 index-matching fluid.

a major contribution towards achieving Nova's performance goals.

In December, contracts were placed with Corning to supply the sixteen 81-cm-diam by 8.9-cm-thick focus lens blanks, and eight of the sixteen 90-cm-diam by 4.3-cm-thick array window blanks. Heraeus' contract is for the remaining 16 window blanks.

Potassium Dihydrogen Phosphate. The KDP crystal production for Nova is well under way. We have ordered a total of 113 plates, each of which is 27-cm square, and 44 pairs of corner plates. In June, we placed contracts for 40% of the Nova KDP with Cleveland Crystals, Inc., Interactive Radiation, Inc., and Lasermetrics, Inc., all of

whom initiated crystal growth at that time. In December, we contracted for the remaining 60% with the first two of these manufacturers. Cleveland Crystals has approximately 20 crystal-growth runs started, with a potential yield of 20 crystal boules and 100 plates. In December, two of these boules were ready for cutting and finishing. Interactive Radiation has 11 runs in progress, with a potential yield of 35 boules and 80 KDP plates. Lasermetrics is growing 12 boules in 2 growth stations, and expects to supply us with 32 plates. The KDP vendors are on schedule for Nova.

The finishing of KDP by single-point diamond-turning, a technique developed at LLNL, has proved to be most effective from a cost standpoint and for optical performance. In 1982, we generated thirty-six 15-cm-square by 1.8-cm-thick KDP blanks, and twelve 27-cm-square by 1.8-cm-thick blanks for Novette. The following tolerances were achieved

- Surface flatness: $\leq 2 \lambda$ at 0.6328 μm .
- Phase matching angle: $\leq \phi$ 30 μrad (6 arc-sec).
- Thickness tolerance: $\leq 2 \mu\text{m}$ (80 $\mu\text{in.}$).
- Parallelism: $\leq 7.5 \mu\text{rad}$ (1.5 arc-sec).
- Surface roughness (rms): $\leq 125 \text{ \AA}$.

The average time needed to diamond-turn a Novette 27-cm KDP blank was 40 hours. This compares to an estimated time of from 250 to 300 hours each for finishing by conventional polishing methods. It is also extremely doubtful that conventional techniques would have produced the dimensional accuracy achieved by diamond turning. We completed the diamond finishing without breaking a single KDP blank. A 27-cm-square crystal is shown on the diamond-turning machine in Fig. 2-12. Figure 2-13 shows the crystal being measured for thickness variations with an LLNL-designed thickness-measurement gauge, which uses two opposing air-bearing linear variable-displacement transducers for measuring thickness variations to a tolerance of 0.1 μm .

The Nova KDP crystal plates will differ from the Novette plates in two ways

- Their thickness will be reduced from 1.8 to 1.0 cm, which will make them more difficult to finish.
- They will be phase matched with new LLNL-designed equipment, which uses

	Specification	Actual
Attenuation coefficient [†] in cm^{-1}	$\lambda_1 = 1.064 \mu\text{m} \leq 0.0015$	—
	$\lambda_2 = 1.053 \mu\text{m} \leq 0.0015$	—
	$\lambda_3 = 0.633 \mu\text{m} \leq 0.0015$	0.0004
	$\lambda_4 = 0.527 \mu\text{m} \leq 0.0015$	0.0009
	$\lambda_5 = 0.351 \mu\text{m} \leq 0.0015$	0.0012
Homogeneity	$\lambda = 0.6328 \mu\text{m}$	
	$\Delta n \leq 2 \times 10^{-6}$	$0.7 \text{ to } 2.2 \times 10^{-6}$
Stress birefringence	$\leq 6 \text{ nm/cm}$	$\leq 5 \text{ nm/cm}$
Bubbles and inclusions	$\leq 0.03 \text{ mm}^2/100 \text{ cm}^3$	$0.00038 \text{ mm}^2/100 \text{ cm}^3$

Table 2-3. Inspection data on the Heraeus 90-cm fused-silica blank.

Fig. 2-12. A 27-cm-square KDP crystal mounted in a vacuum chuck on the LLNL DL-1 oil-showered diamond-turning machine.

the 3ω wavelength rather than the 2ω wavelength used in phase-matching Novette's crystal plates.

In addition, the number of crystals needed will be 10 full squares and 8 corner pieces per array, for a total quantity of 157 blanks (including construction spares). It will take approximately 10 months from the first crystal delivery to complete the Nova assemblies.

Early in 1982, we decided that the Nova KDP finishing should be done by outside industrial sources, rather than at LLNL. The decision was in keeping with our continuing policy to transfer as much optical finishing as feasible to industry, thus spreading KDP finishing technology throughout the optical industry. Our goal was to have two vendors finishing crystals and to have contracts placed by January, 1983. This plan required an intensive evaluation and survey of the industry.

In March, 1982, representatives from 14 companies doing diamond turning attended a pre-procurement and technology-transfer meeting at LLNL. In April, informal letters of inquiry asked these companies for cost estimates for completing Nova's require-

ments. After reviewing the eight replies, we surveyed five of the vendor's facilities and invited these companies back to LLNL for in-depth briefings and demonstrations.

After these technical exchanges, three vendors were selected to participate in a qualification program and a formal fixed-price solicitation. The qualification program involved machining a 15-cm-square KDP crystal using LLNL-supplied tooling, and with LLNL-qualified personnel assisting in the machining sequence. Specifications for surface flatness and finish were established as goals, and all three vendors qualified in October. We reviewed the final bids of the three finalists and selected Pneumo Precision and Cleveland Crystals as suppliers.

Since neither one had diamond-turning equipment or test fixtures that could meet our requirements for the 27-cm-square crystals, LLNL is providing the following equipment to each vendor as part of the procurement

- A new single-point diamond-turning machine, produced to each vendor's own specifications.
- An 18-in. (46-cm) Fizeau interferometer, for use in the final optical evaluation.

Fig. 2-13. LLNL-designed thickness gauge used to measure variations in crystal thickness.

- Phase-measurement equipment, designed and built by LLNL, that is compatible with the vendor's machines and tooling.

Mirrors. Two manufacturers, Zygo Corp. in Middlefield, Conn., and Eastman-Kodak in Rochester, N. Y., are finishing the large optical mirrors for Nova. During 1982, both companies have been in full production, turning out finished mirrors at the combined rate of six per month. This is a significant accomplishment, considering the sizes of these mirrors, which range from 60.4 to 109 cm in diameter and weigh from 111 to 900 lb (55 to 446 kg). All are high-precision mirrors, with reflected wave front distortions $\leq \lambda/10$ over the clear apertures, and with transmitted wave-front distortions $\leq \lambda/8$. The transmitting mirrors also require a wave front gradient of $\sim \lambda/30$ /cm. Of Nova's mirrors that have been finished ($\sim 40\%$), we have not lost one in handling or because of poor quality.

Apodizers. Apodization is used to suppress diffraction ripples from the split at the 46-cm amplifier stage. We have taken two approaches for apodization: the first tech-

nique uses absorption, and the second uses scattering, of the laser beam. In the former, we place an absorbing apodizer at the input to the 31.5-cm spatial filter to symmetrically taper the amplitude of the incoming beam across the split. The obscuration at the 46-cm stage in the beam is approximately 6 mm. The apodizer is made by fusing a cylindrical section of index-matched, copper-doped BK-7 glass into a clear BK-7 substrate. The transmission of the apodizing stripe ranges from 50% at the center to 100% at the edge. The stripe's width is 50 mm, as projected onto the 46-cm aperture. This component was developed by Schott Optical Glass; they have achieved a nearly perfect fusion in two of their four prototypes. These two apodizers are currently in use in Novette.

In the second approach, we make an apodizing stripe by bead-blasting the glass surface so that it scatters the laser beam. This approach may provide adequate apodization, along with higher transmission and lower cost. It would be used on the output side of the 46-cm spatial filter. The width of the scattering stripe is about 15 mm. Prototypes for this alternative scheme were manufactured and tried late in 1982.

Coatings. We have demonstrated the ability to coat the 70- by 40-cm polarizers required for Nova with excellent (85%) production yields. Eight (40%) of Nova's 20 polarizers have already been produced. These coatings are probably the most difficult to produce for Nova because of the tight tolerance ($\leq 1\%$) required for the film thickness. Both coating vendors, Optical Coating Laboratories, Inc., and Spectra Physics, Inc., appear able to produce polarizer coatings on large substrates, to Nova specifications, with excellent production yields.

Both vendors have also produced a variety of mirror coatings on substrates ranging from 60.4 to 94 cm in diameter. A 94-cm-diam mirror, coated for $>99\%$ reflectivity at all three Nova wavelengths, is shown in Fig. 2-14. Five Nova coating designs have been produced, and three different sizes of mirror substrates (some of which are transmitting, for diagnostic purposes) have thus far been successfully coated. Table 2-4 outlines the mirror diameters and the reflectance specifications for the various mirror combinations. The production yield for 25

Fig. 2-14. A 94-cm-diam turning mirror with multilayer thin-film coating, used to reflect λ , 2λ , and 3λ wavelengths.

of these mirrors has exceeded 90%. The first four mirror designs listed in Table 2-4 will be employed on Nova; the last two will be replaced by a single design that will be used for all 94-cm and 109-cm turning mirrors. This design approach was necessitated by changes in the Nova diagnostic approach; we expect, however, that the production yield for the new design will be as high as for the other designs.

Our design requirements for spectral reflectance are $>99\%$ at 1.05 and $0.53 \mu\text{m}$, and $43.5 \pm 6.5\%$ at $0.35 \mu\text{m}$. A spectral (reflectance) curve for a preliminary Nova turning-mirror coating is shown in Fig. 2-15.

Antireflection Surfaces on Borosilicate Glass. About two years ago, Schott Glass Co. developed a chemical back-etch process for producing low-reflectance surfaces on certain glasses.^{18, 21} This neutral-solution (NS) process produced such promising results, with regard to reflectance and withstanding laser damage, that LLNL purchased the right to use it for the Laser Fusion Program.

NS processing, illustrated schematically in Fig. 2-16, has been scaled up from 24 cm and used successfully to treat components

as large as 90 cm in diameter. These components show excellent low-reflectance characteristics, good damage thresholds, and considerable promise for improvement in reflectance levels. Reflectance minima have been produced at 0.53 and $1.05 \mu\text{m}$, and the process can be controlled to yield minima anywhere from about 0.40 to at least $2.50 \mu\text{m}$. Figure 2-17 shows a two-surface transmission scan obtained from a 5-cm-diam BK-7 witness piece, processed with an 80-cm Novette focusing lens. This is a rather extreme result; more typical results show from 0.5 to 0.8% reflection losses through two surfaces in the peak transmission regions. However, without NS

Mirror diameter (cm)	Reflectance		
	$1.05 \mu\text{m}$	$0.53 \mu\text{m}$	$0.35 \mu\text{m}$
60.4	$>99\%$	—	—
60.4	$96.5 \pm 1.5\%$	—	—
82.0	$>99\%$	$>99\%$	$>99\%$
94.0	$96.5 \pm 1.5\%$	$50 \pm 20\%$	$50 \pm 20\%$
94.0	$>99\%$	$>99\%$	$>99\%$
94.0	$>99\%$	$10 \pm 2\%$	$>99\%$

Table 2-4. Coatings produced for the Nova and Novette laser systems in 1982.

Fig. 2-15. Reflectance for a preliminary coating design for Nova's turning mirrors.

Fig. 2-16. Block diagram of the operating sequence for the scaled-up NS process.

Fig. 2-17. Spectrophotometer transmission scan through both NS-processed surfaces of a 1-cm-thick BK-7 witness piece.

processing, these pieces should show transmission losses of about 8%.

The largest apertures processed in the developmental phase (24 cm) and in the scaled-up operational phase (90 cm) are shown in Fig. 2-18. The diameters and quantities of parts scaled up during 1982 are given in Fig. 2-19.

In Nd-doped solid-state pulsed lasers, for the 1-ns regime (at both the fundamental and frequency-doubled wavelengths), NS processing adds an important class of high laser-damage threshold ($>12.5 \text{ J/cm}^2$) to antireflection-treated BK-7 optical surfaces. The ease with which the technique can be scaled up to process large optics is very attractive.²¹

Sol-Gel Antireflection Surfaces for Fused Silica. For several years, we have carried out research with several vendors to obtain damage-resistant antireflection coatings for fused silica in the 3ω wavelength range. (For a report on the research done in 1982, see the preceding article.)

We decided that the Westinghouse process was the most advanced, and late in

1982 we negotiated with them to design and build full-scale production equipment for coating the Nova 80-cm-diam array windows and focusing lenses. This equipment consists of four major subsystems

- A solution mixer and storage facility, which was completed in 1982.
- A cleaning facility.
- An application and sintering facility.
- An etching facility.

The basic process consists of depositing a uniform layer of polymerized silicon alkoxyde onto the finished fused-silica surfaces, burning off the organic material, heat sintering the remaining pure (but porous) fused silica to a much smaller (10-nm) pore size (at about 500°C), and etching the resulting surface to produce a graded-index, broad-band antireflection coating.

Optical Design. During 1952, we made several major changes in the Nova focusing-optics design. The old doublet baseline design, with its KDP array and beam dump, has been previously described in detail.²²

The current baseline for the Nova focusing optics is shown in Fig. 2-4. A smaller axial thickness has resulted from the elimination of one lens element and from reducing the diameter and thickness of the beam dump. The dispersion of fused silica protects the target from unconverted 1ω and 2ω light passing around the 27-cm-square central beam dump. Effectively, the target sits in the "shadow" created by the beam dump.

These two design changes increase Nova's energy and power performance at 3ω , but some focusing and diagnostic capabilities are compromised. In particular

- Larger amounts of unconverted 1ω and 2ω light enter the target chamber.
- The focused spot sizes at 1ω , 2ω , and 3ω are larger than the diffraction-limited values.
- It is not possible to provide an adjustable diagnostic beam splitter.

The baseline singlet design is an aspheric lens, which is diffraction-limited throughout a 2-mrad field of view at a wavelength of $0.41 \mu\text{m}$. This choice was made after trading off the on-axis lens correction for all harmonics vs its correction for coma, so as to extend the field of view.

The focal length, F , of a thin lens is given by

$$\begin{aligned}\frac{1}{F} &= n_0 \left(\frac{1}{R_1} - \frac{1}{R_2} \right) \\ &= n_0 (c_1 - c_2) \\ &= n_0 c\end{aligned}\quad (1)$$

where:

n_0 = index of refraction,
 R_1, R_2 = radii of curvature of the lens surfaces,
 c_1, c_2 = inverses of the radii, with
 $c = c_1 - c_2$.

We can define a shape parameter β , so that $c_1 = \beta c$, and $c_2 = (\beta - 1)c$. Then, if we choose an aspheric surface, such that the lens is corrected on-axis at an index of refraction n_0 , the focal spot diameter, $D_{\Delta n}$, for a change in the index, Δn , is given by

$$D_{\Delta n} = \left| \frac{y^2 c^2}{2} \left(n_0 - \beta - \frac{\beta^2}{n_0} \right) \Delta n \right|. \quad (2)$$

The additional focal-spot diameter due to coma at a distance, h , from the optical axis is given by

$$D_c \approx \left| \frac{3h y^2}{2f^2 (n_0 - 1)} \left(\frac{n_0 + 1}{n_0} \beta - n_0 \right) \right|. \quad (3)$$

Note that, while both spot sizes are functions of the lens shape, only the on-axis spot varies as a function of the change in the index of refraction. Therefore, a singlet lens can provide good on-axis optical performance at all harmonics, or (as a trade-off) coma correction at all harmonics with

Fig. 2-18. Largest-diameter optics (24 and 90 cm) processed in the Nova development and operational phases, respectively.

Fig. 2-19. Diameters and quantities of parts treated in the NS-processing facility.

less well-corrected spherical aberration at the extreme wavelength.

These equations are graphically interpreted for an $f/4$ fused-silica lens in Fig. 2-20. The solid lines show the minimum spot sizes on-axis, as a function of lens shape. The curve marked $1\omega/3\omega$ is for a lens designed for equal spot sizes at 1ω and 3ω . This lens is diffraction limited on-axis, at the intermediate index of refraction. Since the index difference between 2ω and 3ω is less, the $2\omega/3\omega$ curve shows a smaller spot size. The dashed lines show the off-axis coma spot size for two fields of view.

The focus lens is used to view targets up to 2 mrad in extent. Furthermore, the lens should be relatively insensitive to at least 1 mrad of misalignment. Under these conditions, little is gained by departing from the

coma-corrected shape. A ray-trace analysis of the Nova f/4 singlet shows a focal-spot size of less than $40\text{ }\mu\text{m}$ at 2ω and 3ω , and about $75\text{ }\mu\text{m}$ at 1ω .

80-cm Interferometer. The large-aperture optical needs of the Nova and Novette laser systems present us with special requirements for detecting optical distortions caused in their manufacture or mounting. One-tenth wavelength distortions in the laser components are easily detected in the 80-cm diameter Fizeau interferometer built by the Zygo Corporation. This instrument was especially designed and built to operate at 1.06 and $0.63\text{ }\mu\text{m}$ (Fig. 2-21). It also incorporates a computerized automatic interferogram-reduction system. The optical cavity will accommodate assemblies up to 4 m long and weighing

up to 1364 kg . The entire instrument is mounted on an air supported vibration-isolated table. It has already proven invaluable for measuring the optical quality of the assembled large optical turn-mirrors and the assembled KDP crystal arrays.

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Nova Alignment and Diagnostics

Introduction. We have previously described in detail^{2,3} the functions and architecture of the Nova alignment and diagnostic systems, which will guide the 10 Nova beams to the target and characterize the laser beams. We have built and are operating the alignment and diagnostic systems (for 1ω and 2ω light) in Novette, the two-arm laser test bed for Nova. However, changes in the Nova target-illumination design have necessitated developing a new design for the output diagnostics. We have also developed a new beam-multiplexing output-alignment system. This design substantially reduces the cost of the $2\omega/3\omega$ alignment, while preserving the functionality of the alignment system.

Two changes have been made to the target-focusing geometry. First, instead of being a doublet, the target-focusing lens is now a singlet, which gives increased operational performance for the $\leq 1\text{-ns}$ pulse lengths used in 3ω operation. This results

Fig. 2-20. Focal-spot size vs lens shape for the Nova fused-silica f/4 singlet lens.

Fig. 2-21. The 80-cm Fizeau interferometer in the Shiva building. The computerized interferogram-reduction system is on the left.

from a decrease in glass thickness and the corresponding B-integral reduction, whose small-scale self-focusing effects limit the allowed average flux.²⁴ However, the singlet focus lens eliminates the slowly converging back reflection, obtained from the doublet, upon which the old diagnostic design was based. This has required us to develop a new plan for diagnosing the beam without affecting the improved performance that our new lens geometry offers.

Second, the illumination geometry has been changed. We no longer have two 5-beam groups, each group placed on a 100° open cone, with each beam in one cone directly opposite a beam in the other. Instead, we now have a nonopposed beam arrangement; the two cones are rotated with respect to each other, along their common axis of symmetry. In case a target is incorrectly positioned, this new arrangement will now prevent the beams from firing directly into each other's focusing optics, thereby protecting them from damage. The choice

of a nonopposing beam arrangement on the target chamber also affects the alignment procedures. It is no longer possible to view either the beam or a silhouetted target²⁵ from the opposing beam-alignment sensor. Other methods for examining the beam in the target plane, which were previously thought of as being optional alternatives, are now required.

In the following paragraphs, we describe the design concepts for these new alignment and diagnostics systems and the details of those portions that have evolved to a detailed design stage.

Output-Alignment Systems. The output-alignment functions include introducing and viewing full-aperture harmonic alignment beams, positioning the focused light (at all three wavelengths) onto the target, and tuning the KDP crystal arrays for maximum harmonic conversion. Since the direction of each pulsed harmonic beam is determined by the direction of the corresponding 1ω driver, the alignment beams

Fig. 2-22. Nova output-alignment systems. Three modes of operation are illustrated.

for all wavelengths must be coaligned. Previously, the three wavelengths were separated from each other in a central beamsplitter box,^{26,27} and a separate alignment and diagnostic module was attached for each wavelength. In the new alignment system, illustrated in Fig. 2-22, the 2ω and 3ω modules are no longer required on the output sensor. The remaining 2ω and 3ω diagnostics are performed in the redesigned splitter box and on the focus-lens structure, as described later in this article. The alignment viewing functions, previously done with the 2ω and 3ω modules, are now performed by two multiplexed harmonic-illuminator/viewer modules, one for each group of five beams. We kept the alignment

and viewing of the 1ω light as it had been originally designed, because we were too far along in our procurement and fabrication schedule to multiplex it through the spatial filter, as is being done with 2ω and 3ω .

The illumination and viewing of a beamline is selected by inserting two mirrors. The first selects the beamline to be multiplexed to the illuminator/viewer. The second is inserted near the focus of a final spatial filter, so that the $2\omega/3\omega$ local illuminator exits from the output lens as a full-aperture collimated beam, as shown in Fig. 2-22(c). Light reflected from the pointing reference on the centering mirror, or from a surrogate target, returns along the same

Fig. 2-23. The multiplexed harmonic illuminator/viewer, shown on the space frame.

path for viewing by the illuminator/viewer module. The capabilities of the illuminator/viewer are

- Operation with either 2ω or 3ω light.
- Local illumination at full aperture for output alignment.
- Three far-field imaging fields of view: 2 mrad, 1 mrad, 300 μ rad.
- Far-field focusing that is adjustable from 8 mm ahead of best focus in the target chamber to 4 mm past best focus.
- Two near-field imaging fields of view: 100 cm and 45 cm.
- Near-field focusing that is adjustable from the output lens of the final spatial filter to the target chamber focus lens.
- Two systems, each multiplexed for up to five beams.

Only two viewing systems are now needed with this arrangement, compared to the 20 separate modules (two per beam) that were previously needed. The relative location of the harmonic viewer and illuminator, in relation to the laser system's components, is shown in Fig. 2-23. The illuminator/viewer is located inside the spaceframe and is near the final spatial filter, through which beams are multiplexed to it.

A local 2ω or 3ω alignment beam is useful only if it follows the same path to the target as the pulsed harmonic beam. The insertable pointing reference, consisting of a corner cube and a gimbaled mirror (Fig. 2-22), is used for this task. First, the 1ω alignment beam is viewed by the output sensor in the far field, with both the mask and pointing reference inserted [Fig. 2-22(b)]. With the auxiliary mirror tilted slightly, the output-sensor display shows three spots. The brightest spot is the common focus of the light passing directly into the sensor through the mask's two holes. The other two spots are the foci of the light reflected from the pointing reference mirror and the corner cube. These are generally two separate spots, because the corner cube reflects directly back on itself, while the mirror's reflection is reported by an amount that depends on the mirror angle. By adjusting this angle, the two far-field spots can be made to coincide. When this is done, the mirror is normal to the 1ω beam.

When a harmonic alignment beam is introduced, as shown in Fig. 2-22(c), once again light from the corner cube and mirror combination is reflected back to the

Fig. 2-24. The target-plane imager positions the alignment beams on the target's surface by viewing the target-insertion region of the target chamber

harmonic viewer. However, now the two far-field spots on the viewer display can be superimposed by adjusting the direction of the harmonic beam, thus making it normal to the mirror. By this procedure, the harmonic alignment beams are made parallel to the main 1ω beam. A small residual error between the pointed harmonic alignment beams and the pulsed harmonic beams will remain because of dispersion and birefringence effects in the KDP assembly. This is a measurable fixed offset, which can be compensated for by operational procedures.

Target-Plane Imager. focused beams can be viewed in the target plane for accurate focus and position adjustments by numerous methods, but direct imaging from opposing beamlines is no longer possible. Therefore, the new baseline design includes a target-plane imager. One way to implement this capability is to position a frosted-glass reticle at the desired plane. This reticle

Fig. 2-25. (a) Incident and (b) reflected 3ω light; a slight curve in the KDP array's output window provides the converging 3ω (or 2ω) diagnostic signal.

can either be mechanically attached to the target-plane imager (as shown in Fig. 2-24), or to the target inserter (as shown in Fig. 2-22). It is aligned with the target position using the target viewers. Images of the incident beams on this reticle are relayed out of the target chamber to a charge coupled device (CCD) or to other TV-format cameras. Beams can be viewed and aligned at 1ω , 2ω , or 3ω . With the geometry shown in Fig. 2-24, the imager is inserted from the target chamber equator, and a mirror set at a 45° angle directs the light into the optical relay. A simple 180° rotation allows viewing first one, and then the other, set of five beams. An alternate approach would be to insert the instrument from a chamber pole (Fig. 2-22). In this position, the optical axis of the relay is normal to the reticle and images of all 10 beams can be viewed sequentially without moving either the reticle or the imager. Five beams are viewed by transmission, and five are viewed by reflection from the diffuse reticle. This arrangement allows an accurate relative positioning of all 10 beams in a selected plane. Other techniques, including using fiber-optic bundles, are under consideration.

The KDP crystal angle-tuning for generating 2ω light is done in a manner similar to that previously described.^{2b} A small-diameter 1ω beam is introduced ahead of the final spatial filter (Fig. 2-22). It propagates along the main beam path through the crystal array. The resulting 2ω signal is measured by a detector mounted on the edge of the centering mirror inserter, and the 2ω signal is maximized. The 1ω beam is inserted into one beamline at a time from two local illuminators: one for the five north beams, and one for the five south beams.

Tuning the angle of the frequency-conversion array to produce 3ω light will be done in the same manner as for 2ω light. Depending on our ability to detect low levels of 3ω , the 1ω source might require higher intensities than the local illuminator can generate. In this event, the output of the MOR will be inserted into the beamline by routing it through the local illuminator.

Alignment through the laser chain, and centering on the target-chamber cross hair with 1ω light, is done as before²⁷⁻³⁰ [Fig.

2-22(b)]. With the insertable mask and cross hair removed, the 1ω alignment beam from the chain has an unobstructed passage through the final spatial filter directly into the output sensor. When the insertable pointing reference is also removed, the full-aperture beam continues to the (closed) centering mirror and is reflected back to the output sensor from the auxiliary mirror for viewing.

Output Diagnostics. In our former baseline plan, the diagnostic signal for 2ω and 3ω was derived by a reflection from the final surface of the doublet focusing element (Fig. 2-79 of Ref. 26). The slowly converging sample of the frequency-converted 2ω and 3ω light, generated by the crystals, was tipped 3 mrad off-axis. This allowed the signal and the light reflected from the target to be separated in the far field and collected by the output sensor.

In analyzing the focusing design, it became apparent that the laser performance at 3ω with the doublet lens was limited in the ≤ 1 ns pulse-length regime by the doublet thickness, which added significantly to the B integral²⁴ (see "Frequency Conversion and Focusing System," earlier in this section). Using a singlet lens can increase the 3ω performance by approximately 30%. Therefore, the singlet has been chosen over the doublet design. However, a singlet focusing lens cannot provide the same back-reflection for diagnostics as the doublet. We

are evaluating two schemes for providing 2ω and 3ω diagnostics with the singlet. The first scheme is illustrated in Fig. 2-25. A slight curvature ($r = 120$ m) could be added to the output window of the frequency-conversion array, so that the reflection from this surface will provide a converging-beam sample for diagnostics.

In operation, the frequency-conversion array is tipped (to reach the phase-match angle of the desired frequency) and rotated (to achieve the proper polarization and mix ratio). The first crystal array in the assembly operates at a fixed angle and always generates 2ω light. The second crystal array of the assembly is tuned in an orthogonal plane to double the unconverted 1ω from the first array, or it is tuned to the 3ω conversion peak, where it mixes 1ω and 2ω to produce 3ω .^{25,26} For the $1.05\text{-}\mu\text{m}$ YLF oscillator wavelength, the 2ω and 3ω conversion peaks are 4.1 mrad (external angle in air) apart. By cutting the crystals so that the surface normal is between these two peaks, the array can be tuned by tilting it to ± 2.05 mrad to obtain 2ω or 3ω . These angles are doubled by reflection from the crystal-array window to yield a ± 4.1 mrad angle with respect to the incident 1ω beam. The convergence of the diagnostic signal must be equal to (or greater than) this value, in order to stay on the mirrors and within the beam tubes, to reach the output sensor. Additional convergence is added to

Fig. 2-26. The selected value for convergence of the diagnostic signal is below both the damage limits and solarization fluences at the turning mirrors and the input sensor's collecting telescope.

allow for changes in tuning angle as a function of temperature. The design allows for a $\pm 2^\circ\text{C}$ temperature change. Since the diagnostics signal intensifies with convergence, the maximum amount of convergence allowed is determined by the damage limit at 2ω and the solarization fluences for 3ω experienced by the turning mirrors and output sensor optics. The fluences, and the corresponding diameters of the diagnostic turning mirrors and output sensors intercepting the diagnostic signal, are shown in Fig. 2-26. In the Nova design, each beamline has a different intercept diameter because their turning mirrors and beam sensors are at different distances from the target chamber.

As just mentioned, the first crystal array always generates 2ω light and does not change its angle along the phase-match-sensitive axis. The crystals in the first array are cut identically to minimize the spares inventory and handling problems. Therefore, the first array operates at an angle of $+2.05$ mrad with respect to the normal. However, the vector sum of this angle and the orthogonal tuning angle of the second array exceeds the maximum allowed convergence. To compensate for this effect, a wedge could be added to the curvature on the array window.

At the output sensor, the frequency-converted drive signal (2ω or 3ω) is collected by the sensor's input telescope. The full-aperture, on-axis reflection from the target is also collected. These signals can be

separated from each other (as shown in Fig. 2-27) by passing the focus of the target-reflected light through a hole in the mirror directing the frequency-converted drive sample into the calorimeter doing the 2ω (or 3ω) incident and reflected-energy measurements. Signal samples are also split off and transported by fiber-optic cables to a streak camera for temporal pulse-shape measurement.

A second plan was also developed for providing energy and temporal diagnostics for 2ω or 3ω converted light and for light reflected from the target. In this scheme (shown in Fig. 2-28), the primary ghost focus, derived from the reflection at the back surface of the singlet, provides the diagnostic signal. When it is about 15 cm in diameter, the reflection begins forming caustics with about 14 times the beam's intensity at the lens surface. Collecting it here would obscure too much of the beam, so the signal is further reflected from the $1\omega/2\omega$ beam dump. With the dump separated from the array, it can be positioned to intercept the 2ω or 3ω ghost-foci signals where the beam is about 27 cm in diameter and there are no threatening caustics. The reflection from the dump (approximately 4%) decreases the intensity of the diagnostic reflection enough so that it can be collected (at about 6 cm diam) by a calorimeter. This 6-cm block in the beam obscures less than 0.6% of the aperture.

The temporal shape of the beam would be monitored by a mirror in the same

Fig. 2-27. Separating a 2ω (or 3ω) input sample from the target reflection is possible because the signals have different focal points.

6-cm-diam region that the calorimeter blocks from reaching the target. This sample would be directed by the mirror to a module mounted on the side of the focusing system assembly. Next, filtering and collecting optics focus the sample into a fiber-optic cable, which transports it to the streak camera (as shown in Fig. 2-27).

The 2ω or 3ω light reflected from the target can be collected by the output sensor's telescope. It would be separated from spurious reflections originating at the crystal array by the 4.50-mrad off-axis angle between the array and the on-axis target reflection (Fig. 2-29). In this plan, the array window is flat, so different fractions of the beam reflection will be clipped by the various beam tubes.

Neither of these schemes provides an equivalent target-plane image for focal-region characterization. In the first scheme, reflections from the array window would be optically poor because the window is a stressed and therefore distorted optical element. In the second case, the singlet lens must be optimized for focus in transmission, and its ghost foci have extensive spherical aberrations. However, the previ-

ously unavailable ports opposite each beamline offer a method for obtaining the equivalent target-plane characterization. We plan to design an imaging system, to be mounted on these opposing ports, that will characterize the equivalent plane. We will construct two modules, which will give us potentially better images of the target plane than previously possible. In the original baseline, the 2ω or 3ω diagnostic beam made three passes through a diagnostic mirror, which was not figured for transmission. It would also have passed back through the array to reach the output sensor. This would have added aberrations that were not experienced by light directly reaching the target plane. However, using the doublet signal allowed an equivalent plane characterization to be made for each beam on every shot and in a totally noninterfering manner. The proposed opposite-port target-plane image module would only be done for two beams and can only obtain an equivalent-plane image when shots are fired specifically for target-plane characterization. No target could be in place, since it would block the desired image plane. Images of the light reflected

Fig. 2-28. Second plan for $2\omega/3\omega$ diagnostics. The target lens' primary ghost focus supplies a full-aperture beam sample for energy diagnostics, and a partial-aperture sample for the temporal pulse shape.

by the target can still be obtained at the $2\omega/3\omega$ reflected diagnostic portion of the output sensor. As of December, 1982, we had not decided which of these two diagnostic plans to use.

Nova KDP Orientation System. To be used in the Nova frequency-conversion arrays,³³ the KDP crystal segments must be oriented and machined so that the angles between their peak conversion points and their surface normals³⁴ are all within $\pm 30 \mu\text{rad}$ of each other. A measurement system has been designed that will provide the necessary orientation data to the Nova crystal vendors, to enable them to meet this requirement. This crystal orientation measurement system (COMS) is illustrated schematically in Fig. 2-30. The COMS uses a laser source to angle-scan a sample and a standard crystal, to determine the location of their peak harmonic-conversion angles with high accuracy ($\leq 5 \mu\text{rad}$; 1 arc-sec).

The standard-crystal concept, as it is used in this design, is a significant innovation. It reduces a difficult absolute-alignment measurement to a much simpler relative comparison. A standard crystal, with the prescribed orientation, is cut under carefully controlled temperature and wavelength conditions. Its normal is located halfway between the 2ω and 3ω phase-matching an-

gles. This master standard is then cut into segments, one of which is put on each COMS. A COMS will then put parallel beams through its standard and the crystal sample being tested. Both crystals are scanned through an angular range, and the fraction of 3ω light generated as a function of angle for both crystals is simultaneously recorded. Since both the standard and the sample being tested are on the same vacuum chuck, their input faces maintain the same angular orientation with respect to one another. Therefore, any difference in angular separation between their conversion peaks is due to the angular error in the cut of the sample being tested.

Before the COMS method was developed, orienting a crystal required that the normal to the face of each piece be located with a pointing detector. This method relied on an operator's locating the normal with respect to a reference. Using the COMS alleviates problems arising from operator interpretation and offers other advantages. Variations in the phase-match angle (due to temperature and wavelength shifts) are cancelled out, since both crystals experience the same variations simultaneously.

The procedure for matching all the KDP crystals to the crystal standard consists of placing each 27-cm-square by 1-cm-thick

Fig. 2-29. Separating target reflection and unwanted reflection from the array window is accomplished by an angular separation in the far field.

crystal plate on the chuck, next to the standard. A measurement scan is done to determine what angular correction, if any, is required. The crystal plate is removed and placed in a diamond-turning machine, which makes a correcting cut. After this face is correctly oriented, the second is machined parallel to it (with the desired thickness).

Nova Alignment and Diagnostics

Status Summary. Most of the alignment and diagnostic elements designed for Nova are operating on Novette. The elements include input sensors, mid-chain sensors, 1ω output sensors, and 2ω sensors of the former design. Diagnostic elements include a network of diagnostic energy processors, photodiode and calorimeter energy sensors, temporal pulse-shape-recording streak cameras, timing monitors, and a system for determining the peak frequency-conversion angle for the KDP arrays. All components are operating reliably and have proven the design concepts for implementation in Nova.

Software controls for these systems have demonstrated closed-loop alignment, multiple set-point function control, and the use of cross-coupling matrices to provide orthogonal movement for mirror gimbal pairs. Diagnostic software has local control for all systems. The centralized control and display of energy, temporal pulse shape, and pulse timing have also been achieved. The refine-

ments developed in Novette's software controls will be incorporated into Nova, along with the expansion necessary for controlling 10 beams.

Novette's activation has proved valuable for verifying the alignment and diagnostics procedures to be used in Nova. Our experience on Novette has given us confidence that the design changes we are making for Nova will function properly, since they utilize concepts similar to those embodied in Novette's alignment and diagnostics systems.

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Nova Front-End Subsystem

The Nova front-end subsystem consists of all the components from the oscillators in the master oscillator room (MOR) to the input aperture at the start of each arm. The list of components includes all the preamplifiers, isolation, relays, splitter arrays, and the alignment and diagnostic systems.

In the original Nova front end, as designed for the 20-beam system,³⁵ we had to

Fig. 2-30. The COMS compares the 3ω conversion peak of a sample crystal to that of a standard and determines the required angular correction.

put most of the amplifiers in the MOR before the east/west split, so they could be shared by each of the two 10-beam systems. The original design also required a long path to both 10-beam systems, with many turns and beam relays.

During 1982, Nova was defined as a 10-beam system, with all the beams in the new Nova laser building. Because of this decision, we made major changes in the design and layout of Nova's front end. With the 10-beam system, now in one laser bay, we were able to optimize the design by shortening the path from the MOR to the laser, using only one turn in the path, and eliminating beam relays. We also were able to move all the 50-mm amplifiers and the other 50-mm components out of the MOR and place them on the laser spaceframe. The MOR now contains only the oscillator systems, 10-mm amplifiers, Pockels cell isolation, and the 10 to 50 mm beam expander (for relaying the beam into the laser bay).

Our redesign has several advantages

- With the 50-mm components out of the MOR, all the 20-kV electrical components are gone. Since the 10-mm ampifiers use 5 kV circuits, the MOR is now completely separated from the 20-kV capacitor banks for the laser amplifiers. The separation reduces electrical noise in the MOR and improves safety.

- There is now more space in the MOR, allowing flexibility for future oscillator systems. Future systems will probably include a pulse-shaping system and a system to provide 20-ps pulses for various diagnostics and target irradiation.
- Nova's redesigned layout is now almost identical to the Novette design, which worked well. Our experience on Novette can now be used to improve Nova's front end.

in redesigning the front end, we used a very conservative approach to obtain the best possible beam profile at the start of each arm, with minimum pre-pulse energy and amplified spontaneous emission (ASE) from the oscillator and preamplifier stages. The following sections describe the general staging considerations and the layout of the MOR, the preamplifier, and the splitter array. We also discuss some provisions for x-ray backlighting and for aligning the harmonic crystals by using the MOR output.

Basic Staging of the Nova Front End.

The basic staging of the Nova front end is shown in Fig. 2-31. The beam system is completely relayed, from the flat oscillator-output mirror to the aperture at the start of each arm. In the figure, the various relays are numbered from 1 to 5 and will be referred to as such in the text that follows. The oscillator is followed by one 10-mm

Fig. 2-31. Schematic layout for Nova's front end.

and two 50-mm amplifier sections, before the north/south (N/S) split. Up to this point, the system is almost identical to Novette, which has been described in detail.³⁶ Note that the function of spatial filter 2, at the end of the Gaussian beam section, is to clean up the beam and reduce ASE as much as possible (on Novette we have successfully used a 2X diffraction-limited pinhole in this spatial filter).

After the N/S split, each beam goes through another 50-mm amplifier, and then through another spatial filter to relay each beam to a five-way splitter array. These 50-mm amplifiers are placed after the N/S split to provide the maximum drive into each arm, particularly for long-pulse operation of several nanoseconds. Under these conditions, the 50-mm amplifiers will saturate and the maximum drive will be limited by damage. Therefore, it is extremely important that the beam profile through the 50-mm amplifiers be as smooth as possible.

We analyzed the beam profiles in the Nova front end (Fig. 2-31) using the MALAPROP code, particularly in the section from the hard aperture before spatial filter 3, to the apertures at the start of each arm. We found that the beam profiles in this section are almost completely controlled by the size of the pinhole in spatial filter 3. With no pinhole here, the 6.36-mm hard aperture can be relayed almost perfectly to the next relay plane (after the five-way beam splits). However, severe modulation will occur on the beam between spatial filter 3 and the same relay plane; the most severe modulation will occur close to spatial filter 3. With a pinhole in this spatial filter, we can control modulation. We could also reduce it near the spatial filter; however, modulation would then be introduced onto the beam in the next relay plane. We have selected a pinhole size that will give about the same beam modulation in the components as in the relay plane.

On Nova, front-end relaying is complicated by the different lengths of the arms, as measured from the hard aperture at the start of the arm to the target. In the nonopposing target-illumination configuration chosen for Nova, the maximum difference in length between two arms is 17.14 m. The difference gives various propagation distance from the 6.36-mm hard aperture

to the aperture at the start of each arm and yields different beam modulations at the start of each arm. The latter modulation is controlled for each arm by the size of the pinhole in spatial filter 3. We found that with a 10X diffraction-limited pinhole, we obtain beam profiles with a peak-to-average value of 1.31 at the aperture farthest from the relay plane (8.57 m) and a value of 1.17 at the relay plane itself. When we looked at the beam profiles at the start of each arm, as a function of pinhole size in spatial filter 3, the 10X diffraction-limited pinhole was about the best design possible. We also found that with this pinhole the beam profile 3.81 m from the relay plane was very smooth, with a peak-to-average value of 1.10. The beam profiles can be further improved only by using separate relays of different length at the start of each arm. Such an addition is expensive and complicated, and does not justify the small improvement that would be obtained.

We also studied several relaying systems that covered the entire distance from the 6.36-mm hard aperture to the start of each arm. Spatial filter 3 can be made long enough to relay the beam all the way to the arm inputs, without using spatial filters 4 and 5. For the same 10X diffraction-limited pinhole in spatial filter 3, we obtain identical profiles at each arm, with or without the splitter relays 4 and 5. The major differences, for a typical layout without these relays, are that the final 50-mm amplifiers after the N/S split are about 18.3 to 21.3 m from the relay plane, and we have a beam profile in these amplifiers with a peak-to-average value of 1.47. By using splitter relays 4 and 5, and by placing the amplifier 3.81 m from the relay plane, the peak-to-average value is reduced to 1.10. The result is a major improvement, particularly since we will be driving these amplifiers very hard. Note that although these relays are about 18.3 m long, they can easily be placed on the laser spaceframe.

Layout of the Nova MOR. In the baseline layout (Fig. 2-32) for the MOR, we use two oscillators: one to cover the range from 100 ps to 1 ns, and the other to cover the range from 1 ns to 10 ns (or longer). These oscillators have previously been described in detail.^{37,38} The rest of the MOR layout is almost identical to that used in Novette.³⁶

There is now sufficient space in the MOR for pulse-shaping equipment and a 20-ps pulse system for diagnostics or target irradiation. The beam out of the MOR is now in line with the beamline on the space-frame. The layout makes the beam path from the MOR to the laser bay very short and simple.

The MOR layout has been drawn up in detail,³⁹ and by the end of 1982, we were designing and ordering components to bring the MOR up in early 1983.

Layout of the Preamp Section and Splitter Array. As just mentioned, all of the 50-mm components were moved from the MOR to the spaceframe and were placed between the beamlines for the laser arms, at the system's input section. In the previous design,³⁵ this space was used by the trombone lines in the splitter section. We redesigned the splitter section to make space for the preamplifier section, as shown in Fig. 2-33. For all the target illumination schemes that we considered (including the final nonopposing beam design for Nova),

the longest arms were always 3 and 7 north, and 4 and 8 south. For these arms, as shown in Fig. 2-33, we placed the trombone delay lines in line with each arm, rather than folded back, as in beams 9 and 10. Thus, two slots were left open on each spaceframe for preamplifier components. In our new splitter array design, we still use the variable splitter, which is composed of adjustable half-wave plates and polarizers. However, all the beams are now in the same plane. The layout simplifies the splitter array and reduces the number of turning mirrors; however, we must add 90° quartz rotators to obtain the correct beam polarization for each arm. We have also placed an alignment sensor in front of the splitter array to ensure accurate beam alignment into the splitter array.

The layout of the preamplifier section, which follows the general staging just described, is shown in Fig. 2-34. The first section, where the beam is Gaussian, contains the first two 50-mm amplifiers (each with a gain of approximately 20). These are

Fig. 2-32. Layout for Nova's master oscillator room.

followed by two more 50-mm amplifiers in the second section, where the beam has a square spatial profile. Finally, after the N/S split, there is another 50-mm amplifier. With the 10-mm amplifiers in the MOR, we have much more gain than is needed, even for long-pulse operations. This gives us considerable flexibility in setting up the front end for the various pulse widths that will be used in the system. We generally want to have as much gain in the Gaussian beam section as possible, to reduce ASE with the spatial filter ($M = 1$) at the end of the section. For long-pulse operation, the gain in this section will be limited by pin-hole closure in this spatial filter. For short-pulse operation, we can set the gain in various sections to minimize $\Delta\delta$.

For preamplifier-section isolation, we use a mixture of Faraday rotators and Pockels cells, as shown in Fig. 2-34. This is a change from the previous design, which used the latter only. The Pockels cells are used

mainly for ASE control and are placed after every second 50-mm amplifier in each beamline. In each 50-mm Pockels cell, the contrast ratio of about 10^3 (as measured on Novette) exceeds the gain of the two amplifiers, so ASE cannot grow in the system. The remaining isolation between the amplifiers is obtained with Faraday rotators. The simplicity and reliability of their power conditioning makes them an attractive alternative to Pockels cells. The Faraday rotators also require less propagation distance between amplifiers for proper isolation. By the end of 1982 we had drawn up the final layout of the preamplifier and splitter array⁴⁰ and were designing the remaining components.

Component Development for the Nova Front End. We also made considerable progress during 1982 in the development of various front-end subsystems. The electronic controls for the oscillators and the associated fast-timing circuits were completed

Fig. 2-33. Layout for Nova's splitter array.

for Novette and worked very well. We made progress towards developing Pockels-cell drivers, using both planar triodes and hydrogen thyratrons, with new devices being installed and tested on Novette. The power conditioning for the 10-mm amplifiers is also being redesigned. Details of all these developments are given in "Nova Power Conditioning," later in this section.

X-Ray Backlighting and Crystal Alignment. We designed x-ray backlighting and frequency-conversion crystal alignment capabilities into the front end. In general, two types of x-ray backlighting experiments are planned. We use short-pulse backlighting (to freeze the target's motion), or a long pulse (to study its temporal evolution). Great flexibility is also required in selecting which (and how many) beams are to be used for x-ray backlighting. The possibilities include using one beam, two opposing beams, five beams from one side, or five beams from each side, chosen as symmetrically as possible. We cannot assign some beams to the task of x-ray backlighting, a priori, and then design a beamline from the MOR to these arms. On Nova, we will use the same scheme that has been demonstrated on Shiva and is planned for Novette.³⁶ Two pulses are time-multiplexed

into a single beamline in the MOR. The target-irradiation pulse is transmitted first, followed about 20 ns later by the backlighting pulse. In this scheme, we either take 20 ns out of the backlighting arm or add 20 ns to the target arms and use the Pockels cell at the entrance to each arm to select the appropriate pulse. On Nova, the trombone-line assemblies will be moved to obtain the appropriate propagation-time changes for the various beams. In the new splitter-array design, as shown in Fig. 2-33, a total of four of the trombone lines have been mounted in-line with the beams (two on the north and two on the south sides), and cannot be moved. However, after the N/S split, we have a trombone line in each beamline so that we can delay the five arms on each spaceframe ± 20 ns, relative to each other. In addition, the remaining three beams without the folded trombone lines can also be moved ± 20 ns on each side. Thus, we can satisfy all of our x-ray backlighting needs.

The Pockels cell at the start of each arm has a contrast ratio of about 10^3 so that the relative input energies for the pulses (and the gain in each arm) will have to be carefully selected, to keep the pre- and post-pulses on the target within acceptable

Fig. 2-34. Layout for Nova's preamplifiers.

limits. This constraint has not yet been analyzed in detail. It may even be necessary to add another Pockels cell to some beams.

For aligning the frequency-conversion crystal arrays, we are able to produce a 100-ps-long, 100-mJ pulse with a repetition rate of 10 s. As shown in Fig. 2-32, this pulse is diverted out of the main beam just before Pockels cell 1-2, to reduce the risk of damaging it. The pulse is then recombined with the main beam before the MOR output spatial filter. It remains orthogonally polarized and is taken out of the main beam at the start of the preamplifier section on the spaceframe, as shown in Fig. 2-33. The pulse is finally injected back into the system beamline at the output spatial filter for aligning the 3ω frequency-conversion crystals.

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Control System

Introduction. The Nova control system must provide an efficient and economical method of controlling the large number of diverse, controllable components in a laser of Nova's scale. Figure 2-35 is a block diagram of the control system concept, starting with the operators' consoles and proceeding through the organizational subsystems to the laser hardware itself. Nova has over 5000 separate control and data-acquisition points, most of which are either physically remote or difficult to operate without some form of computer control. The Nova control system, as previously reported,⁴¹⁻⁴⁴ consists of four major application-related subsystems (power conditioning, alignment, laser diagnostics, and target diagnostics) and a fifth, unifying, central control system.

Our major accomplishment in 1982 was the successful implementation on Novette of a similar, but smaller, version of the Nova control system. The long-range decision, made early in the planning of

Novette, to commit to the use of Nova controls is advantageous to both developers and users. Even a preliminary version of the Nova control system provides Novette with a substantial improvement in operational functionality and flexibility over previous control-system technologies (such as that on Shiva). The parallel development of Novette and Nova controls has meant better results for less effort. Moreover, Novette provides a test bed for Nova's controls, driving issues to their appropriate conclusions through practical implementation, testing, and use. In this article we describe this effort, which was almost entirely devoted to the implementation and activation of the four major control subsystems.

Alignment Control System. The alignment control system must steer each of 10 laser beams accurately over a path of 250 m, from the master oscillator to the fusion target.⁴⁵ The system does this by moving the optical components with over 1000 stepping motors controlled by 29 distributed LSI-11 computers. A central control system,⁴⁶ using three VAX-11/780 computers, provides an integrated central control for the alignment system. A fiber-optic local-area network, Novanet,⁴⁷ connects the distributed processors to the central system. Interactive color-graphics consoles, with a touch input system, interface menu-driven remote controls and status information. The digital processing of sensor images closes the feedback loop for automatic alignment. Strategically located multi-function sensors along the laser view both the beam and reference cross hairs. These sensors generate video information for pointing, centering, and focusing, which is transmitted to an image digitizer and array processor attached to one of the VAX computers. The alignment program is synchronized with manual controls, but runs without the operator's supervision.

The LSI-11 stepping-motor controllers (SMC)⁴⁸ serve as the primary alignment-control building blocks. The SMC units are customized (with patch panels and software tables) to match them to the layout and isolation requirements of each laser installation. Up to four motors are grouped together into generic "devices," such as for cross hairs, spatial-filter pinholes, and

gimbaled mirror pairs. The SMC units supply substantial local control and status information, near the alignment components, through a programmable alphanumeric control panel. Switch-function labels, displayed above the panel controls for the currently selected device, allow users to visualize the operation being performed.

Alignment Control Software. Figure 2-36 is a block diagram of the interrelated components of the alignment control software. Three major software programs, unique to the alignment system, are shown in the box labelled "VAX Application Programs." The first of these, the console alignment program (CAP) implements system configuration and manual controls. One copy of the CAP runs for each active console being used. The graphics system, a product common to the entire control system, interfaces the CAP programs to the operator through the console hardware. An extensive frame library contains definitions for particular graphic controls and status displays. Independent graphic status displays are maintained by the second program, the status display program (SDP). The automatic alignment of mirror gimbals and spatial-filter pinholes is accomplished by the third program, the automatic alignment program (AAP). These three programs communicate through a shared memory, implemented with global sections in disk-

based files; signaling for synchronizing memory access is done with event flags. Both these functions are supported by the manufacturer.

The user interface at the central control console brings these distributed components together in a manageable manner. Access to operator controls is layered (from general to specific functions) by traversing a tree of graphic touch-input menus. The four-level hierarchy is illustrated in Fig. 2-37. Level 1 selects the room or major area to be controlled. Level 2 controls the automatic configuration and device setup for various tasks. Level 3 allows operators to select devices from a schematic illustration (Fig. 2-38) for detailed control at level 4.

Two-sided CRT screens in each console are reserved for independent status reports about the alignment control system. Schematic representations of selected portions of the laser system are enhanced with color-coded symbols reflecting the current status of each alignment device.

Automatic Alignment. The AAP is a control process that is independent of the other alignment programs. The gimbal alignment loop is illustrated in Fig. 2-39. The CAP communicates with the AAP by placing a list of commands to be executed in a shared data area. The AAP determines from this list which sensor and field of view to use, and then performs image

Fig. 2-35. Block diagram for the Nova control system.

analysis to determine if alignment errors exist. To correct any detected errors, the necessary stepping-motor commands are calculated from the sensor information. The system is flexible with regard to changes, because only the methods are coded in software. The hardware-specific information is maintained in separate data tables. This allows the system to work on several laser configurations by changing only the data tables. The data tables contain information for the five phases of AAP operation (as illustrated here for a gimbal alignment):

Data table	Indexed by	Provides
1	Command	Devices needed for the given command
2	Command	Set points (stored positions) for devices
3	Sensor	Intensity control information
4	Gimbal	Specific information for pointing
5	Gimbal	Specific information for centering

The first phase ensures that all the devices needed for performing the com-

mands in the command list are available. The second phase sets the laser to the proper starting configuration for performing a command. In the third phase, an image is acquired, its intensity level is checked, and the sensor adjusted (if necessary). In this phase, the background noise from each sensor is also subtracted from the images, to give a more consistent image quality across sensors. In the fourth and fifth phases, the gimbals are actually moved to correct the measured pointing and centering errors. Alignment references are automatically inserted in the beam, or retracted, as needed.

The alignment controls described here were implemented on Novette during 1982. Work on Novette covered the MOR and preamplifier areas, as well as the main beamline. Reusable software tools were developed for the control of classes of devices, such as gimbals and sensor packages. Methods were designed to accomplish routine configuration setups and status information displays. The image analysis program has been demonstrated with an automatic alignment of the front-end gimbals.

Nova/Novette alignment control system

Fig. 2-36. Block diagram for Nova alignment-control software.

Image Analysis. Image analysis techniques extract information from video images acquired at the sensor packages along the laser chain, so that the automatic alignment program can correct pointing and centering errors. Several techniques are used to analyze the different kinds of images.

The first type of information consists of an image taken from a sensor set for viewing the errors in pointing. These images are characterized by a small, bright, well-defined spot. This lets us use a simple technique to determine its center: the image is thresholded and exponentiated, the background is removed, and the centroid is calculated.

To provide a reference point for centering corrections, the location of the beam's central point must be calculated with a second type of analysis. These images differ greatly from the pointing images just discussed. Typically, the beam has a large area with a somewhat circular shape and a bright center; the brightness tapers off towards the edge. Correlating a user-specified "binary disk template" with the image will determine the center of the brightest portion within the circle.

The third type of analysis is used to determine the center of a cross-hair shadow. A stored template of an idealized cross hair (noting its center point) is correlated with the current image acquired from the laser. The point of maximum overlap,

along with the template center, gives the center of the cross-hair shadow in the acquired image. The correlations are done by using frequency-domain techniques, which provide significant improvements in speed over similar spatial-domain techniques. To generate the stored template, the cross-hair shadow is enhanced (by subtracting an image that includes the cross hair from one without it). This method has been demonstrated to work, even with severely diffracted images.

Laser Diagnostic Control and Data-Acquisition Software. We demonstrated a subset of our Nova system in 1982, by bringing the laser-diagnostics system on-line in Novette. This system contains

- Transient digitizers.
- Charge-coupled device (CCD) streak cameras.
- Remote front-end processors (FEPs).
- VAX control software.

The first two devices measure the temporal characteristics of the laser pulses. The FEPs measure laser energies with photodiodes and calorimeters located at various points on the laser chain. By December, the Novette laser-diagnostics system included two streak cameras, five FEPs, and two transient digitizers. Provisions have been made for including up to five streak cameras and ten transient digitizers. The system can display the measured beam energies and plot calorimeter and transient-digitizer

Fig. 2-37. Menu-frame tree.

data on the console CRT screens, as well as plot streak-camera pulse contours for off-line analysis.

The transient-digitizer control software regulates the digitizers from the central operator consoles. Menus that mimic the local controls of the transient digitizer and its oscilloscope-type plug-ins make the system easy to use. A main menu provides for high-level functions, such as arming for a shot, storing data, and displaying data as a one-dimensional plot on the consoles. An independent program exists for providing hard copies of these plots. The latter feature will be added to the high-level control menu.

The software for controlling the transient digitizers translates the operator's touches at the console to control messages to the digitizers. These messages pass through Novanet to the interfaces, and finally over the IEEE-488 bus to the digitizers. Status information, indicating whether the requested operation was performed, is returned to the operator.

VAX programs control the streak cameras, which are connected to a Novanet multiple-device interface. These programs allow the operator to manually control the camera mode and to record data. The oper-

ator can also use the console to subtract the background, to display the data as a color-enhanced image, or to plot pulse contours from the data. These functions were designed to be consistent with Nova's more automated operations.

The Novette diode and energy processors (DEPs) can control either 20 diodes or 20 calorimeters (60 on Nova) from a console graphics frame, or from a membrane switch panel at the DEP. The DEPs report data and status information (upon command from the VAX programs) as was done in Shiva. On Novette, however, the error messages are presented in English text on the front panel and repeated at the VAX control console. The time needed for doing the sampling has been reduced, and the sensors report their identity to the DEP when they are polled. VAX console programs allow operators to load and configure the DEPs, to control operating modes, to retrieve data, and to display the energy data and DEP configuration.

Target-Data Acquisition and Control System. We finished installing the target-data acquisition and control system, as outlined previously.^{49,50} This system consists of three FEPs attached to the VAX (via Novanet) and connected to seven CAMAC

Fig. 2-38. MOR selection frame.

crates in the diagnostic room. A fourth FEP was installed in the target room to control the vacuum system.⁵¹ The trigger system was installed essentially as described.⁵² It includes a direct optical connection (via fiber optics) to the MOR, and a connection to the power-conditioning FEP for synchronization with the target FEP software.

The digitizing equipment, which consists mainly of CAMAC modules and transient digitizers, is controlled by the FEPs in the diagnostic room. A single-display Ramtek console (with a touch panel) was installed in the diagnostics room to allow convenient setup and control. This remote console is isolated from the VAX with fiber optics, but otherwise it performs the same functions as the main control room consoles.

The basic software architecture is shown in Fig. 2-40. Functionally, the system is segmented into two components: data acquisition and control, and quick-look processing. Both segments of the system have been implemented on Novette. It is a table driven system, with the network-shared memory (NSM) used for process-to-process communication and synchronization. The NSM stores the configuration definition,

setup parameters, diagnostic-to-hardware maps, command definition tables, and operational status information. It is also used to buffer the acquired data and any messages issued to the operator.

We completed the target-diagnostics FEP software in 1982. Each FEP can have up to four CAMAC crates attached to it, and can control 15 different types of CAMAC modules. For most of the preshot setup sequence, the FEP acts as a slave to the VAX. However, for the last 15 s before the target shot, the FEP independently handles the final interrupt triggers and shutter sequences. After the shot, it again becomes a slave to the VAX for collecting the shot data. This technique requires no diagnostic data-base information in the FEP.

We also implemented the operator-to-machine interface using color enhanced menus, with touch-selectable categories supported by the console graphics system. The menu set, currently consisting of 21 frames, permits the complete operation of the system's hardware.

Using the system at its uppermost level is quite straightforward. The execution of a single command causes the system to

Fig. 2-39. Gimbal alignment loop.

perform a comprehensive check to detect abnormal conditions, make the necessary setups, and put the system on line. All data acquisition and control commands are executed via the command processor. Only the lower-level commands are "hard coded." Most of the upper-level commands are defined (by table entries) in terms of other commands.

Once prepared for a shot, the operation of the data acquisition and control system is interlocked with the power conditioning system, and no further operator interaction is required. The system may still be accessed (via the menus) during a shot sequence, but only those commands that will not affect the operation of the system (other than an abort request) are permitted.

The operator is advised of the target system's activity (and its current state) by the displayed status information and by one-line text messages that are issued by the system and logged in data files. The operational state of each piece of hardware and a record of errors that might have occurred are also kept as an integral part of the shot data.

Keyed ASCII files are used to tailor and configure the system and to save the acquired data. The command definition, menu command page tables, and hardware configuration are loaded from such files, either at the system's start-up, or on demand. By editing the appropriate file, a command may be created, redefined, or deleted, the

command presentation changed, or a hardware configuration defined for a shot.

One other VAX process allows the operator to do a quick-look evaluation of transient and calorimeter digitizer equipment setups and post-shot data. The quick-look process is controlled by a pair of menus on the console (for touch response). It uses the commercial graphics package DISSPLA to plot the analog data directly on the Ramtek display and to generate files for use with a hard-copy Versatec printer/plotter.

The performance record of the target-diagnostics system on Novette is very impressive. The problems encountered have been few and relatively minor, and the system has been operational virtually 100% of the time. The system is extremely fast, with a typical data-acquisition time of less than 1 min. On the Shiva system, the data-acquisition time for a similar configuration would have been 10 to 15 min. The real advantage of this speed is that it gives us the ability to monitor the system's performance during a shot. The comprehensive error checking designed into the system has been of tremendous value. In conjunction with the status information displayed on the menus, this system enables the operators to quickly isolate and correct hardware problems that probably would not have been recognized on the Shiva system, or, if recognized, would have taken weeks to isolate.

Vacuum Controls. The integrated vacuum control system, which was installed on

Fig. 2-40. Target diagnostics software architecture.

Novette in 1982, controls all the chamber pumps and valves. It also implements all safety interlocks, ensuring that all vacuum requirements are fulfilled prior to valve openings. It continually monitors pressure and vacuum sensors and puts all valves in a safe condition if any leaks are detected.

The vacuum system is controlled with a system of individual panels, each one customized for a particular diagnostic. With our Novette experience, we noticed a trend of constantly changing the mechanical vacuum connections to track the diagnostic requirements. We designed easily changeable vacuum panels for Nova to handle these continual modifications. We simply used removable plastic overlays on standard panels. Each new panel that is required can be generated with an artwork change, without redesigning the basic panel.

Power-Conditioning Control. The major accomplishment for 1982 in power-conditioning control was completing and activating the software for operating this major subsystem from the operator's console on Novette. The overall power-conditioning control system architecture is shown in Fig. 2-41. The VAX resident console software accepts operator control requests through a touch panel mounted on the Ramtek displays and writes the perti-

nent status information on the same displays. Control and status information is communicated to the FEPs through a Unibus multiport memory. A backup auxiliary FEP and hard-wired control panels are provided for off-line diagnostics.

We designed and implemented the sequencer, a software tool that schedules the time-ordered control system events during a shot sequence. Software has been written that will summarize the power-conditioning system's performance for display or archiving. An interlock program was written to monitor and log all the entrances into the facility. Programs were also written to display and archive the lamp-circuit diagnostic waveforms.

As reported previously,⁵³ predefined configurations can be used to repeat common shot operations. Several enhancements to this configuration loader were completed. The loading time was reduced and the number of device types that can be supported was expanded.

The design and implementation of the power-conditioning event logger was completed. The log can now be examined at a terminal, or it can be archived for future examination at either a VAX or an FEP.

The final design and implementation of the power-conditioning front-end

Fig. 2-41. Power-conditioning control system.

system was also completed during 1982. Two FEPs are used to control and monitor the power-conditioning components. An auxiliary processor was added to allow operating the power-conditioning system without the VAX and to do other maintenance functions.

Central Controls. Installation of the entire central control system for Novette was completed during 1982, including a VAX-11/780 with an array processor, four Ramtek operator consoles, the Novanet Nodestar, and a terminal development area. We also moved from our prototype development area and began installing the Nova Control Center in the Nova laboratory building. This work included activating a large VAX-11/780 facility, which supports 64 development terminals in several areas (including the Nova office building).

The development of central hardware and software tools, as previously described,^{54,55} evolved (as planned) into the application efforts just reported. Integrating these tools into the four major control subsystems of Novette, and clarifying many specific laser control requirements, were essential elements of the feedback needed for producing a successful Nova control system.

We learned where our common tools were successful in actual use, and where little or no further effort was required. We also learned which enhancements could be postponed. Most importantly for Nova, we learned where changes in scope and implementation were required. The following paragraphs discuss one tool in each category.

Praxis, the high-level, block-structured computer language, designed and implemented for Nova, reached maturity in 1982. It supports separate module compilation, user-defined types, exact machine-accessing specifications, and extensive type checking—all of which are useful features for programming real-time control systems. Praxis generates code for both the VAX and LSI-11 computers. An important milestone was reached in the development of Praxis; it was considered good enough to be used as the language for all of Novette's control system software. A substantial amount of software has been written for the Nova control system and applied to Novette. During 1982, 22 authors developed over 250 000

lines of code (organized into approximately 2800 modules) for the Nova and Novette control systems.

The design of the operator-console graphics system remained as described in previous Annual Reports.^{56,57} During 1982, its implementation reached the point where the base-level operations were considered good enough for Novette. The implementation of some designed features and higher-speed operations (by internal buffering) were postponed so that we could concentrate on Novette control problems and graphics software reliability.

The system's capacity for manipulating graphics control pictures and their components was quadrupled. Up to 400 subpictures and 200 control frames can now be used. The system was further upgraded by adding interlocking mechanisms that allow other plotting packages to share the color displays. The consoles now display data and waveform plots, color-enhanced beam profiles, and streak-camera images. Control input from the operator and feedback from the computer are automatically suspended while these packages are using the displays.

When more than two Ramtek interfaces were plugged into the same Unibus, this particular implementation of the DEC Unibus specification suffered from bus timing failures. With Ramtek Corporation's assistance, we located and solved this complex bus-interface problem, which had plagued many other users of Ramtek displays (and certain other products) since the VAXs were first produced.

One tool that required a change of scope was the network shared memory (NSM), a distributed data base, containing command and status information tables called regions. As originally designed, the NSM was supposed to support communication between programs in the same computer, or between programs in different computers (connected by a multiported memory or by Novanet). During 1982, the Novanet feature was completed, the NSM feature having been completed earlier. However, the data rate was too low for use in these applications, so we changed to a direct connection with Novanet.

During 1982, the console graphics system was also modified. We are now using another interprocess communication scheme

(VMS mailboxes) instead of the NSM package. The changes made in both of these message-oriented applications point out the awkwardness of using shared tables for this type of task. The need to support message-oriented applications was overlooked in the initial design criteria for the NSM.

However, for table-driven commands and global status-information applications, shared regions are still appropriate. Therefore, we are currently using the NSM within the target-diagnostics subsystem on the VAX, for the power-conditioning subsystem (on the VAX and LSI-11), and for all cross-subsystem communication.

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Nova Power Conditioning

Introduction. A significant amount of physical construction was done on the Nova power-conditioning system during 1982. The Shiva capacitor bank was completely removed, and the racks were set up in the Nova basement. New Nova equipment was used to supply the full complement of racks, as needed. We set up the capacitor test stand and began assembling and testing Nova's capacitor modules. At year's end, 350 modules were mounted on the racks. A separate rod bank was constructed and outfitted with new circuits. All the MVA power supplies were tested and placed in the Nova substation. We also began installing conduit from the substation to bring power into the Nova building. The 100-kVA power supplies were installed in the building; the fanout racks were installed, and assembly was begun on them. Switch racks were placed in the bank, and their assembly was also begun. The flashlamp test and fast-pulse laboratories

were constructed. Many of the power-conditioning controls chassis were fabricated, and cable tray runs were installed. A switcher power supply was designed for the master oscillator room. The MOR's electronics were updated, based on our Novette experience. The Pockels-cell drivers were improved, both in performance and reliability, and procurement orders were placed for the Nova system. The plasma shutter effort concentrated on mechanical design, and system problems were addressed. In 1982, the Nova power-conditioning area changed dramatically. In the following sections, we discuss these accomplishments in more detail.

Component Testing. In 1982, we performed a substantial amount of component testing for quality control. As procurement arrived, all crucial and high-cost parts were subjected to an incoming inspection and test. A prime example of this process was the infant-mortality testing done on all high-density capacitors. The details of the testing are given elsewhere.⁵⁸ Capacitors rejected for leaking, or for failing the hipot test, were easily repaired at the factory; only capacitors failing the discharge test are worthless. Of the 2025 capacitors delivered, 64 were returned to the vendors for replacement or repair at no cost to LLNL.

Thirteen capacitors failed our discharge testing, a much lower number than we had anticipated. Part of the reason for this improvement over previous levels may have been due to our including the infant-mortality testing as part of the specifications; manufacturers knew in advance that any failures would be replaced at their expense.

Bank Staging and Layout. The 10-beam Nova laser system was authorized in 1982; the first two beams were built in Bldg. 381 and called Novette. The corresponding power-conditioning staging for Nova is given in Table 2-5.

Since the physical layout of the Novette bank was under severe space constraints, most of the circuits were built with 12.5-kJ capacitors. This accounts for the mixing of capacitor types on Nova's 20.8- and 31.5-cm amplifiers; two arms of circuit modules were built with 12.5-kJ capacitors for use in Novette. Because the Faraday rotators contain a large coil, no inductors are

included in their circuits. The pulse ionization lamp check (PILC) has been described elsewhere.⁵⁹

Power Supplies. Our power supply work was divided into two areas: installing the MVA units, and upgrading and installing the kVA units. The final two MVA units were delivered early in 1982 and passed their acceptance tests by February. In the fall, all seven units were installed in the Nova substation, as shown in Fig. 2-42.

The kVA units were upgraded. The silicon controlled rectifier (SCR) driver cards were replaced by more reliable ones. To eliminate a drift problem, the operational amplifiers in the reference and comparator modules were replaced with new, improved units. A new computer interface, designed to be compatible with Nova's controls, was also checked out and installed. In addition, the doubling capacitors were replaced with non-PCB-containing units. The nine power supplies were installed in the Nova basement. Two spare power supplies will be built into the Nova system, one of each type.

Capacitor Bank Design and Assembly. The Nova capacitor bank is installed on 16 rows of racks 19.25 m long, stacked seven double shelves high, for a total of 224 shelves, as shown in Fig. 2-43. The total

bank energy required for Nova's 10 arms will be 50.4 MJ.

The assembly process for each capacitor module consists of

- Installation of a capacitor and its associated bus work in a fabricated phenolic tray.
- A hipot test.
- A five-station assembly line for installing the pulse-forming network (PFN).
- A test consisting of 25 shots at 22 kV.
- Placement of a module in the capacitor bank for intercabling after a successful test.

The assembly line averaged 32 modules per day, including testing.

The high-voltage switches (at the end of each row in the bank) consist of two size-D ignitrons (connected in series) with the appropriate voltage dividers, circuit monitors, and heater and trigger transformers. All of the high-voltage cabling is installed in A.B.S. phenolic trays to ensure adequate isolation from ground.

Nova's high-voltage distribution system includes nine 100-kVA power supplies, with up to four channels of diode isolation in the high-voltage fanout racks. Each diode channel can be selected or rejected by computer control. The high voltage for all the disk amplifiers comes from the MVA power

Table 2-5. Nova power-conditioning staging.

Component	No	Circuit components	Stored energy	Total number of lamps	Capacitance per circuit (μF)	Inductor size per circuit (μH)	Power supplies	Ignition switches used by each component
Faraday rotator (9.4 cm)	11	11	231 — 3 kJ ^a	—	101	—	1 — 100 kVA	1
Faraday rotator (15.0 cm)	10	40	1 500 — 12.5 kJ	—	156	—	1 — 100 kVA	2
Faraday rotator (20.8 cm)	10	50	2 000 — 5 kJ	—	200	—	2 — 100 kVA	4
Faraday rotator (31.5 cm)	10	50	2 000 — 5 kJ	—	200	—	2 — 100 kVA	4
Amplifier (9.4 cm)	20	160	2 880 — 3 kJ	320 — 44 in. ^b	87	450	1 — MVA	10
Amplifier (15.0 cm)	10	120	2 160 — 3 kJ	240 — 44 in.	87	450	1 — MVA	10
Amplifier ^c (20.8 cm)	6	48	1 200 — 12.5 kJ	96 — 44 in.	104	450	1 — MVA	2
	24	192	4 032 — 3 kJ	384 — 44 in.	101	450		8
Amplifier ^c (31.5 cm)	8	80	3 000 — 12.5 kJ	160 — 44 in.	156	650	2 — MVA	4
	32	320	10 560 — 3 kJ	640 — 44 in.	160	650		16
Amplifier (46 cm)	30	480	18 000 — 12.5 kJ	2 400 — 19 in.	156	450	2 — MVA	30
Rod amplifier (3.8 cm)	14	14	700 — 12.5 kJ	84 — 19 in.	208	450	1 — 100 kVA	3
Pulse ionization lamp check (PILC)		7	42 — 3 kJ	—	29	—	1 — 100 kVA	7
Total		1572	19 905 — 3 kJ 4 000 — 5 kJ 23 700 — 12.5 kJ	1840 — 44 in. 2484 — 19 in.			6 — MVA 8 — 100 kVA 1 — spare MVA 1 — spare 100 kVA	101

^aSize of each capacitor.

^bLength of flashlamp.

^cBeing used on Novette.

supplies with 20 channels of diode isolation. There are seven 4-channel fanouts and six 20-channel fanouts for the MVA supplies in the Nova system.

The Nova safety-dump relay system, for discharging the capacitor bank, consists of 229 seven-circuit air-operated switches. When the switches are closed, energy absorbing resistors are placed across each capacitor circuit.

Nova Pulsed-Power Control System.

The design of the Nova pulsed-power control system has been previously described.⁽⁶⁾⁽⁷⁾ During 1982, our baseline design guided the fabrication of the hardware and software elements for the control system. The hardware devices and computer programs were integrated into an operating system and were used to control the Novette pulsed-power system. The Novette control system (described in the previous article) is a subset of the Nova system. The operational experience that we gained during 1982 with the Novette control system has proven that our hardware designs and fabrication methods will effectively support the Nova project.

Table 2-6 identifies the major elements of the control system being worked on, delineates our progress on each one during 1982, and gives their status as of the year's end.

The complex nature of our control system warrants using computerized tests for sev-

eral of the interfaces listed in Table 2-6. We constructed two test sets capable of testing both the bank interfaces and the lamp-circuit-diagnostic (LCD) data-acquisition interfaces. We also wrote interactive software for the test sets, to help technicians activate and repair these interfaces.

The devices shown in Table 2-6 are connected to the control computer via a communication network called Novabus.⁶² This network operates at 10 Mb/s, and uses fiber-optic technology to minimize electromagnetic interference problems and provide voltage isolation. Control and status information is exchanged between the control computer and the device being controlled in both a polling and a demand fashion. Devices are repetitively polled for status information at subsecond intervals. During timed operations, such as those occurring before a laser firing, commands are sent to specific devices at specific times. Precise timings (with an accuracy greater than 10 μ s) are not directly computer-controlled. Devices requiring such accuracies contain internal timers, which are synchronized by a universal trigger in the Novabus network. Devices that require timings more accurate than 1 μ s obtain their timing signals by a direct connection to the master-oscillator electronics.

During 1982, a set of computer programs was developed and proven on the Novette

Fig. 2-42. MVA power supplies, as installed in the Nova substation.

system; these programs will also be used for Nova. They allow the operator to control pulsed-power devices via touch panels and also display the status information for the equipment on colored graphic displays. This is done for Novette's

- Oscillator electronics.
- Fast-timing electronics.
- Personnel-safety system.
- Plasma-shutter electronics.

Master-Oscillator-Room Pulsed Power.

A new power supply was designed for the 10-mm rod amplifiers in the MOR. Each amplifier requires two power supplies; one to simmer the amplifier's lamps (maintaining voltage on them at all times), and one to flash them. On Nova, we will use a pair of switching power supplies. The supplies will use compact, lightweight, and efficient ferrite-core step-up transformers operating at a frequency of 20 kHz.

One of the power supplies must charge a large capacitance from zero to a high voltage, on a 10-s cycle. With conventional control methods, a trade-off must be made in the average power output, to keep the transistors from failing due to the high current flow at the start of the charging cycle. To

solve this problem, we developed a pulse-width control circuit that lets the power supply operate at a safe maximum current during the entire charging cycle.

Nova's Oscillator Controls. We designed an oscillator control system (see Novette "Electronic Systems" later in this section) and began assembling a small-scale prototype of the system in the fast pulse development laboratory. This prototype will test the functions of all the oscillator control printed-circuit boards, modules, and chassis prior to their being used on Nova. It will also be used as the diagnostic and repair center for malfunctioning controls components. A local LSI-11/23 executes the status and control functions for the prototype that would normally be performed by the Nova FEP.

Pockels Cell Driver. Development work was done on several Pockels cell driver designs during 1982. We have categorized the designs as

- Type 1: less than 2-ns rise time (1 ns preferred) to 5 kV into 50 Ω ; for pulse carving of the long-pulse oscillator.
- Type 2: Less than 4-ns rise time to 5 kV into 50 Ω (or into dual 93- Ω cables); for

Fig. 2-43. Nova capacitor-bank layout.

MOR switchout and preamplifier ASE suppression.

- Type 3: Less than 10-ns rise time to 10 kV into 10 Ω (or less): fanned out to laser bay ASE suppression Pockels cells.

The Type-1 requirements were met by a bulk silicon photoconductive (Auston) switch system. (Shiva had used a planar triode driver, which was too slow to carve pulses less than 3 ns in length. Since the short-pulse oscillator can only generate pulses out to 1 ns, a temporal performance gap existed.) Using 150 μ J from the short-pulse oscillator to activate the Auston switch, electrical rise times of 250 ps have been demonstrated (Fig. 2-44). The switch consists of two silicon chips (in series) in a transmission line structure. When simultaneously illuminated, they discharge a charge line with a round trip time equal to the pulse width. (A pulse generator charges the line by applying up to 10 kV for 1 μ s. A fast delay channel from the oscillator control triggers this generator, synchronizing it with the optical pulse.)

The Type-2 requirements had been met previously by planar triode drivers that were built in-house. These triodes are high-bandwidth vacuum tubes that were developed for microwave modulator applications. Based on the performance of three of these

units on Shiva, we carried out a fairly extensive redesign, breadboarded it, and completed the circuit development work. The goals of this effort are to increase the reliability, manufacturability, and flexibility of these triodes.

Three versions of the Type-3 drivers were tested during the year. An in-house design capable of simultaneously driving 20 parallel 50- Ω cables was assembled and tested. In this unit, 20 charge lines are switched out by a single hydrogen thyatron.

A similar, six-way driver was obtained from the Argus laser, and by using a magnetic sharpener, we were able to obtain the desired 9-ns rise time from its thyatron.

Late in the year, we replaced this six-way driver with a prototype five-way driver from Cardon Instrument Co. in England. This unit was built to LLNL specifications and uses a new, ultra-fast thyatron. Based on the success of this prototype, a production order has been issued, and Nova will use these drivers in the laser bay. We also ordered two similar prototype units, which will be capable of meeting the Type-2 requirements.

Plasma Shutter. Development work continued on the Nova plasma shutter, an optical switch based on the rapid generation of a dense plasma in the beam path. As reported previously,⁶¹ the plasma is created by discharging approximately 3 kJ of energy into a small aluminum filament, vaporizing it in less than 300 ns.

We concentrated on the detailed mechanical design of various components in the system, including the high-current pulse generator, the spatial-filter vacuum housing, the chip-changer mechanism, and the alignment mechanism (Fig. 2-45). By the end of 1982, these designs were nearly completed.

The pulser design is an evolutionary extension of the previous two prototypes. It has eight energy-storage capacitors arranged

Fig. 2-44. Output waveform of an Auston switch.

Table 2-6. Status of the Nova power-conditioning control system.

Device	Quantity	Progress	Status
LSI-11/23 front-end processor	4	Fabrication bid awarded	In fabrication
Capacitor-bank interface	22	Fabrication bid awarded	In fabrication
Lamp-circuit-diagnostics interface	22	Fabrication bid awarded	In fabrication
100-kVA power-supply interface	12	Fabrication completed	Being tested
MVA power-supply interface	6	Design drawings completed	In procurement
MOR power-supply interface	7	Prototype testing	In final design
Computer-system interface	18	Fabrication bid awarded	In fabrication
Control panels	13	Fabrication bid awarded	In fabrication
Safety-interlocks interface	2	Fabrication bid awarded	In fabrication
Master-oscillator interface	1	Prototype testing	In final design
Plasma-shutter interface	10	Prototype testing	In design

Fig. 2-45. Nova plasma-shutter installation.

Fig. 2-46. Chip design of the Nova plasma-shutter source.

in a hollow rectangle, feeding a central coaxial feed to the chip via a rectangular, four-section rail-gap switch. A silicone dielectric sheet is the primary high-voltage insulator. Some examples of changes that were made include

- Redesign of the switch to allow aligning the trigger blade with the cover off.
- Addition of a UV preilluminator and trigger-blade bias networks in molded switch covers.
- Design of a coaxial high-voltage interconnection system.
- Redesign of the chip/high-current interface.
- Design of a swing-away pulser housing door to allow in situ maintenance.
- Movement (where possible) towards easily replaceable, modular components.

The vacuum-spool assembly, which mates the plasma shutter to the optical beamline, is a complex unit with significant design challenges. A cryogenic vacuum pump with three gate valves is mounted on the spool, allowing for rapid vacuum isolation when the shutter chips need changing. The spatial-filter pinhole, with its associated manipulator, is mounted a few centimeters from the plasma-shutter aperture. An insertable lens, used for alignment purposes, is also in the vicinity.

The plasma-shutter alignment must be done in two axes. Manual alignment will be carried out on a support cradle, an adaptation of those used for the large $h \times x$ amplifi-

ers. The pulser housing can also be opened out into the aisle area for crane access.

The chip design⁶² has been cleaned up further and is shown schematically in Fig. 2-46. The chip body is made of polycarbonate, with a ceramic insert to provide a substrate for the vapor-deposited aluminum foil. Filament current enters the pulser through tapered, silver-plated copper pads. To minimize arcing, these pads are held in place with very high contact pressure at the time of firing.

The chip-changer mechanism is shown in Fig. 2-47. One chip magazine will hold approximately 16 chips. An air-driven ram picks a chip up, inserts it into the pulser's jaws, and holds it in place during firing. An extraction mechanism removes and discards the used chip before the next cycle starts.

The electronics control package for each plasma shutter will be distributed in rack mounts near that shutter. This concept was proven during the 1982 Shiva testing, when the electromagnetic compatibility of the plasma shutter with low-voltage digital electronics (near the pulser) was demonstrated. The Nova design will incorporate the upgrades and modifications that are necessary for handling the pneumatic chip changer. Each plasma shutter will communicate with the power-conditioning front-end processors via the standard Novabus fiber-optic link.¹ Triggering for the system

will be done with redundant fast, laser-diode fiber links from the MOR. Finally, a plasma-shutter prefire will initiate the transmission of a fiber-optic "abort" signal back to the MOR, to inhibit switchout of the laser pulse.

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Nova Mechanical Systems

During 1982, we began fabricating most of the mechanical engineering designs for Nova. It was also a year of intense manufacturing to meet the tight Novette schedule. Laser hardware was fabricated,

Fig. 2-47. Plasma-shutter chip interface, with the combination pinhole/cross hair in position.

assembled in a rebuilt clean room, and then installed and operated on the Novette test bed. The construction of Nova's spaceframe, with its associated hardware, proceeded concurrently with this effort. In this article, we discuss our progress with regard to the spaceframe, amplifiers, clean room, gas system, mirrors, spatial filters, Faraday rotators, cradles, and the interstage hardware.

Spaceframe. As previously described,⁶⁴ there are four rooms requiring spaceframes to support and stabilize the Nova laser: the MOR, laser bay, switchyard, and the target room.

The MOR frame was completed by the end of 1982, and the others were in various stages of construction, as shown in Fig. 2-48. Our shop at LLNL has constructed each frame in modules that are easily transportable, will fit through building doorways, and do not exceed the building's crane capacity. The laser-bay spaceframe is being fabricated in $12.2 \times 3 \times 3$ -m sections. When finished, the two separate spaceframes will each weigh 220 tons and will support five laser beamlines apiece. The switchyard frames consist of six separate towers (four for supporting turning mirrors

and two for supporting diagnostics equipment) with a combined weight of 50 tons.

The modules are stress-relieved and painted before delivery. In some cases, to ensure dimensional control, the entire structure was assembled in the shop and then cut apart for shipment. The modules were assembled in each room, precisely aligned, and welded together with nondistorting procedures. The required installation tolerance of ± 1.6 mm has been met. As each frame is completed, it is tested to ensure that its modal frequencies match the calculated values.

We have completely fabricated the three different types of cradles used to hang and align components on the frames. The cradles include units already used on Shiva to support components on a 30.5-cm offset centerline, cradles with a 45.7-cm offset centerline, and telescoping cradles for the 31.5- and 46-cm box amplifiers. Each cradle has been used successfully on the Novette test bed.

Amplifiers. Five different sizes of amplifiers have been built: rod amplifiers with 1- and 5-cm apertures, and disk box amplifiers with 20.8-, 31.5-, and 46-cm apertures. Each

Fig. 2-48. Laser-bay spaceframe during installation.

type has been clean-room assembled, and operated. By the end of 1982, approximately one-fifth (32) of the amplifiers necessary for Nova had been processed.

At that time, all of Nova's parts had not been delivered. Some of the reflective surfaces⁶⁴ were still being manufactured by electrodeposition—a slow, time-consuming process.

Fig. 2-49. A 46-cm disk amplifier being assembled in the clean room.

Fig. 2-50. Three 46-cm amplifiers are mounted close together and act as a single six-disk amplifier.

The largest amplifiers (46 cm) have two split-elliptical disks. Assembling them revealed slight disk distortions caused by the mount. Fortunately, the mount design contained sufficient latitude to accommodate the simple modifications necessary for bringing it up to our performance specifications. No assembly problems were encountered with the refurbished 10- and 15-cm Shiva amplifiers.

The few assembly problems that we did have were associated with new hardware and were mainly due to substandard manufacturing quality. In Fig. 2-49, a 46-cm amplifier is shown being assembled in the clean room, and Fig. 2-50 shows how the 46-cm amplifier is installed on a spaceframe.

Clean Room. Shiva's class-100 clean room was refurbished for Nova and was activated in March, 1982. The clean-room facility was previously described in detail.⁶⁵ By the end of 1982, only amplifiers had been assembled in this facility.

Components that required less stringent clean-room control were assembled in the Shiva laser bay, which is a class-10 000 clean room. Units that have been processed in these facilities have a higher quality than previously seen. The quality results from well designed tooling, processing

Fig. 2-51. The processing of 20.8-cm amplifiers. One is shown entering the high-pressure spray box for cleaning. The second (foreground) is ready for flashlamp installation on its assembly cart.

procedures, and validation techniques. Figure 2-51 shows 20.8-cm amplifiers being processed in the class-100 clean room. Spatial filters, mirrors, and Faraday rotators were assembled in the class-10 000 clean room.

Spatial Filters. The hardware for 70 spatial filters, in seven different sizes, has been completely fabricated. The spatial-filter units—vacuum tubes with lenses at each end and pinholes in the center—control beam propagation. Their unique features have been described previously.⁶⁶ Two units of each size have been assembled and operated on the Novette spaceframe. A 31.5- to 46-cm spatial filter is shown mounted on the Novette spaceframe in Fig. 2-52. These tapered vacuum tubes were processed in a specially constructed facility for electropolishing the inner surfaces. This process provides a clean (<50 particles/cm² with >5 μ m size), low-outgassing surface ($<3 \times 10^{-11}$ Torr L-cm⁻²/s). These characteristics are quite important; micrometer-sized particles can cause laser-induced damage to optics when they are struck by the laser beam, and low outgassing rates reduce the size of the vacuum pumps needed for evacuating the filters. No

Fig. 2-52. A 31.5- to 46-cm spatial filter installed on Novette. The pinhole and vacuum pump are at the waist; the lens positioners are at the ends.

industrial electropolishing facility existed that could process these tubes, with lengths of up to 7.9 m and with end openings ranging from 20 to 74 cm in diameter. In our electropolishing facility, a tube is mounted on a carriage. With a cathode and the electropolishing fluid inside, the tube is free to rotate slowly. With this arrangement, it

Fig. 2-53. The assembled linear actuators are attached to the mirror gimbal by a spherical bearing and supported at the mirror base by a hinge.

Fig. 2-54. An interferogram for a 94-cm-diam, 100% reflecting mirror in its bezel and mount, which shows only 1/4 wave of distortion.

takes about 30 min to remove 0.01 mm of the interior surface.

Mirrors. By the year's end, the fabrication of Nova's turning mirrors had neared completion. The hardware for 72 of Nova's 98 mirrors, in four sizes, had been delivered. Three of the four sizes had been assembled and operated. (A total of 13 mirrors had been assembled.) The accurate linear actuators for the mirrors were still in production. The only mirror type not assembled was the one with a 109-cm aperture. The detailed mechanics of these mirror

assemblies have been described previously.⁶⁷ The mirror gimbal's adjustable design, as shown in Fig. 2-53, allows the glass disks to be supported without distortion. A typical interferogram (Fig. 2-54) for a mounted mirror (Fig. 2-55) shows about 1/4 wave of distortion.

Interstage Hardware. Coupling hardware between each component in a chain seals the beamline from the atmosphere. Double-lead jack screws allow the quick removal of the bellows for maintenance. Sometimes the flanges accommodate nitrogen-gas inlet and outlet connections. There are over 40 different types of interstage hardware in each laser chain. All such hardware for Nova has been fabricated, and two full beams have been assembled and tested on Novette.

Nitrogen Gas System. Nova will be equipped with a cooling system that will circulate nitrogen gas, drawn as required from a 28 000-gal liquid nitrogen tank. Nitrogen gas will be recirculated, in quantities of up to 5000 scfm, with roots-type rotary blowers. It will be compressed by the blowers, recooled by a set of refrigerators, and trim heated to the operating temperature $\pm 1^\circ\text{C}$. In contrast, Shiva used a single-pass flow-through system and exhausted its

nitrogen to the atmosphere. It is expected that recirculation could save up to \$10 million in five years of Nova operations.

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Nova Target Systems

During 1982, the Nova target chamber (Fig. 2-5) was built by Chicago Bridge and Iron, Inc. (CBI) and delivered to LLNL in November. The chamber is made of 5083 aluminum alloy, and has an inside diameter of 4.6 m; it is 12.7 cm thick and weighs about 20 tons (see Fig. 2-5).⁶⁸ All openings were machined so that their axes intersect within 250 μm of the center of the chamber, to aid in the installation and alignment of the lens positioners, the target positioner, and the plasma diagnostic packages. During a test at the CBI Memphis plant, the vacuum tightness exceeded our specification of $1 \times 10^{-7} \text{ atm} \cdot \text{cm}^{-3}/\text{s}$, assuring us that we will achieve our design goal of a 1×10^{-7} Torr base pressure.

Two Nova lens positioners were completed and installed on Novette, as shown in Fig. 2-56 (Ref. 68). All of the design, fabrication, assembly, and installation of these units was done during 1982. Our Novette experience indicates to us that only the mounting details need changing to make the concept completely applicable to Nova's requirements (as currently defined).⁶⁸ Two sets of mounts for the KDP array, the 1 ω beam dumps, centering mirrors, and final alignment reticles were completed and installed on Novette. Some minor modifications made during the installation will be incorporated into the Nova design. The lens positioner and KDP mount are shown in Fig. 2-56, as installed on Novette. Extension spools are used on the target chamber to provide a mounting surface that is the same distance from the target as it would be in the Nova chamber.

We have done a finite-element stress and deflection analysis for the central ring of the target chamber, using the GEMINI code. The model is shown in Fig. 2-57, and Fig. 2-58 is a computer-generated printout of the deflections around various ports due to atmospheric loading when the chamber

Fig. 2-55. A turning mirror, shown ▲ mounted in its enclosure and gimbaled mount, on Novette. The two linear actuators are located 120° apart at the bottom.

Fig. 2-56. Nova lens positioner, as installed on Novette. ▼

is under vacuum. The maximum stress predicted is about 2500 psi ($1.7 \times 10^8 \text{ Pa}$), a small fraction of the yield strength of the chamber material. The port deflection is nominal, and is nearly symmetrical, ensuring that the pointing accuracy of target diagnostics sensors and alignment optics

Fig. 2-57. Model used for the finite-element study of the stresses in, and the deflections of, the central ring of Nova's target chamber.

Fig. 2-58. Predicted deflections for two ports in the central ring of Nova's target chamber.

Radical deflections of center ring, in inches

packages will not be seriously affected during the pumpdown cycle.

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Nova Project Management

During 1982, major uncertainties regarding the Nova project were resolved by Congress and the Department of Energy (DOE). Frequency conversion was approved, Phase II was indefinitely deferred, and the project was fully funded. These decisions clearly established the remaining technical criteria and schedule for the project. Authorized as a 10-beam system, with frequency conversion to the second (2ω) and third (3ω) harmonics, Nova is scheduled for completion in the fourth quarter of FY-84 at a total estimated cost of \$176 million.

Nova Financial System. The Nova financial system, based on performance-measurement techniques, is operated to conform to the basic requirements for DOE project management systems. The financial system tracks the budgeted cost of work performed (BCWP) against the budgeted cost of work scheduled (BCWS), for both commitments and costs. The fundamental soundness of the financial system designed for Nova was reaffirmed in 1982. A revised baseline was established on March 1, based on a projected project cost of \$185 million and completion in the second quarter of FY-85. This baseline was structured to allow an easy transition to a \$176 million plan with completion in the fourth quarter of FY-84, pending the full project funding that occurred in early FY-83.

The scheduling system for Nova⁶⁹ became fully operational in 1982. This system identifies and tracks progress towards key assembly milestones. Status, as measured against the milestones at the end of June, September, and December, showed that the Nova project was on schedule.

Master Assurance Program. The Nova master assurance program, which integrates the quality and safety assurance programs, continued to be an integral part of our activities throughout 1982. Its effectiveness was conclusively demonstrated by the safe and successful completion of Novette.

Conventional construction was completed in 1982, and a corresponding final audit report was issued. This report ensured verification of the completion of contractual requirements, and defined the status and review criteria for the final as-built documentation.

Design and procurement activity in the Nova special-facilities areas continued in 1982. Corresponding assurance-program activities included the design and procurement reviews shown in Table 2-7. These reviews were highly effective for assuring design quality and procurement-strategy viability, and for identifying safety issues through preliminary hazards analyses. Interface-related designs were controlled through the design change request (DCR) procedure, which was used nine times in 1982 to control and effect changes to special-facility components.

The updated *Master Assurance Program Plan* was released in the latter portion of 1982. At that time, noting that Nova was leaving the procurement phase and entering the activation phase, the Project Manager requested an internal audit of Nova's assurance program for special facilities. Representatives from the LLNL Quality Assurance and Hazards Control offices, and from the Laser Program's Activation Personnel, were assigned as a team to perform this audit. The audit committee was chartered to examine three areas:

- Implementation and compliance with the requirements of the master and subordinate assurance plans.
- Status of assurance program activity for optics and alignment, diagnostics, and controls.
- Effectiveness of the assurance program in achieving project objectives.

In addition, the committee was asked to report on the impact of the Novette test bed for Nova, and to be particularly thorough in its investigation of the safety aspects of the master assurance program.

By the end of 1982, the audit committee had completed its project-personnel interviews, as well as its documentary research and on-site inspections. The completed report, including findings and recommendations, was scheduled for release in early 1983.

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Milestones	Description
1	Order laser glass
2	Shiva laser hardware becomes available
3	Begin spaceframe fabrication
4	Begin laser glass delivery
5	Begin power conditioning/capacitor bank installation
6	Begin spaceframe installation
7	Begin control room installation
8	Order array windows, focus lenses, and KDP crystals
9	Complete installation of two laser chains in Novette
10	Complete installation of master oscillator room
11	Begin laser bay alignment systems installation
12	Begin delivery of array windows, focus lenses, and KDP crystals
13	Complete laser spaceframe installation
14	Begin laser installation on Nova spaceframe
15	Begin target chamber installation
16	Complete target chamber installation
17	Begin final turning mirror installation
18	Complete capacitor bank installation
19	Begin focus lens/KDP crystal array installation
20	Complete final turning mirror installation
21	Complete power conditioning installation
22	Complete laser installation
23	Complete delivery of array windows, focus lenses, and KDP crystals
24	Complete focus lens/KDP crystal array installation
25	Complete alignment systems installation
26	Nova operable: project completed

Table 2-7. Major milestones in Nova's activation schedule.

Novette

System Description

Construction of glass laser systems at LLNL over the past decade has been an evolutionary process.⁷⁰ The work in this area is well documented in the Annual Reports from 1973 to the present. Our system design has been gradually improved using successive generations of progressively more reliable, efficient, and cost-effective components.

Novette utilizes a master oscillator to drive a single-pass power amplifier; the master oscillator power amplifier design is known as a MOPA system. Our experience with high-power laser-pulse propagation in the dielectric media of Nd:glass amplifier chains has led to the use of spatial filters in optically relayed designs to control local spatial-intensity fluctuations in laser beams. By maximizing the efficient use of each amplifier's aperture, while minimizing the ripple of the laser field, the peak laser flux that each optic must tolerate to produce a given output power is minimized. Ultimately, it is always the damage threshold of the optical components that limits the single-shot performance for MOPA lasers. We have designed our system components

so that they can be easily serviced and will maintain their alignment during laser operations. Our design concept has led to the use of sealed modular components through which the oscillator pulse is progressively amplified and the beam expanded such that its flux remains below optical damage thresholds.

The conversion of laser light to its harmonics in nonlinear birefringent media is a well known process that was first demonstrated over 20 years ago.⁷¹ Because shorter wavelengths more effectively couple to fusion targets, considerable interest arose in using conversion to generate short-wavelength beams to irradiate fusion targets. In the last three years, efficient conversion of 1.05- μm (1ω) light to its second (2ω), third (3ω), and fourth (4ω) harmonics in KDP crystals has been accomplished in several laboratories.⁷²⁻⁷⁴ Material- and optical-coating limitations currently restrict the harmonic fluxes tolerable in large, high-power, fusion systems. Handling and focusing of 2ω (0.53- μm) light is possible with the same materials and coating technologies used for 1ω light. We must be concerned with solarization of most optical glasses, low damage thresholds, and high nonlinear index of refraction with 3ω (0.35- μm) light. While solutions exist to the problems associated with 3ω light, we have not been able to apply them in sufficient time to build and meet Novette's early experimental schedule. Novette was therefore assembled with 2ω -conversion capability only.

In designing the Nova system (see Nova "Laser Design and Performance," earlier in this section), we used a proven architecture from past systems, including repeatable high-quality master oscillators. Our master oscillators typically produce a 10^{-4}J output; thus, the remainder of the system must have a gain of 10^4 to produce a 100-kJ output. To achieve this required output from Nova, the master-oscillator pulse will be amplified by a glass rod preamplifier with beam-splitting elements to provide 10 parallel beams of $\sim 1\text{J}$ each. The preamplifier in this system will be followed by 10 single-pass parallel chains of rod and disk amplifiers, with each succeeding stage having an aperture larger than the previous stage. Our assemblies are designed so that the output flux of each stage approaches

the operating limit of the component. Spatial filters, capable of handling high-power beams, are inserted between stages to provide beam expansion and propagation control.

The Novette laser system is comprised of salvaged components from earlier lasers (such as Shiva), special frames, and two borrowed output-amplifier chains from Nova. The laser subsystems, including the master oscillators, preamplifiers, controls, and laser diagnostics are all of the current Nova baseline design. The target chamber and all parts through the 20.8-cm stage were removed from Shiva, with the chamber modified to accept two large-aperture horizontal beams. The target-diagnostics instruments in Novette include many modified components from Shiva and Argus. We are also utilizing several entirely new diagnostic systems.

The Novette power amplifier, optically very similar to two Nova arms, was completely assembled and activated in 1982, with the first target experiments completed late in the year.⁷⁵ Initial target experiments were conducted at 2ω , the monolithic KDP crystal arrays were located between the final turning mirrors and the target chamber.

Novette's oscillator and switch-out system has a shot-to-shot reproducibility of $\pm 2\%$ energy and $\pm 1\%$ pulse duration. The laser pulse generated in the MOR⁷⁶ emerges from an actively mode-locked Q-switched Nd:YLF oscillator among a burst of approximately 20 pulses spaced 8 ns apart. Each pulse in the burst exits the oscillator with subnanosecond accuracy relative to the Q-switch time; pulse timing is controlled through precise shaping of the oscillator flashlamp pulse. Timed by the Q-switch trigger, the same pulse in the burst is consistently selected by a Pockels-cell gate.

The oscillator output-mirror plane is optically imaged or relayed through the laser system and, since the oscillator is adjusted to lase only in its lowest-order transverse mode, the resulting output beam is spatially Gaussian with a 1.4 mm-FW at e^{-2} waist. Table 2-8 traces the temporal Gaussian pulse through the MOR components and preamplifier stage in the laser bay. The configuration chosen in the table is designed to deliver spatially uniform and temporally

Gaussian ~ 7 -mJ, 100-ps pulses to the inputs of Novette's two amplifier chains. The master oscillator and preamplifier system is

extremely versatile, since it provides outputs ranging from 20-ps-long, 1-mJ pulses to 40-ns rectangular pulses containing joules of

Table 2-8. Novette MOR and preamplifier components.

No.	Component	No. of components	Beam diam (cm)	Glass thickness (cm)	Refraction index	Nonlinear index of refraction	Bank kV	Net gain	B-incident total
1	Oscillator	1	0.1	0.0	1.000	0.00	0	1.00	0.00
2	Mirror	2	0.1	0.0	1.507	1.24	0	0.98	0.00
3	Lens	1	0.1	0.5	1.507	1.24	0	0.99	0.00
4	Polarizer	2	0.1	1.5	1.507	1.24	0	0.98	0.01
5	Pockels isolator	1	0.1	11.0	1.500	1.00	10	0.92	0.03
6	Polarizer	2	0.1	1.5	1.507	1.24	0	0.98	0.03
7	Pockels isolator	1	0.1	11.0	1.500	1.00	10	0.92	0.05
8	Polarizer	2	0.1	1.5	1.507	1.24	0	0.98	0.06
9	Mirror	2	0.1	0.0	1.507	1.24	0	0.99	0.06
10	Rod amplifier	1	0.1	15.0	1.523	1.05	5	0.92	0.08
11	Mirror	2	0.1	0.0	1.507	1.24	0	0.99	0.08
12	Polarizer	2	0.5	1.5	1.507	1.24	0	0.98	0.08
13	Pockels isolator	1	0.5	11.0	1.500	1.00	10	0.92	0.08
14	Polarizer	2	0.5	1.5	1.507	1.24	0	0.98	0.08
15	Lens	1	0.5	0.5	1.507	1.24	0	0.99	0.08
16	Rod amplifier	1	0.5	15.0	1.523	1.05	5	19.41	0.10
17	Polarizer	2	0.5	1.5	1.500	1.00	0	0.98	0.10
18	Pockels isolator	1	0.5	11.0	1.500	1.00	10	0.92	0.12
19	Polarizer	2	0.5	1.5	1.507	1.24	0	0.98	0.13
20	Mirror	2	0.5	0.0	1.507	1.24	0	0.98	0.13
21	Spatial filter lens	1	0.5	0.5	1.507	1.24	0	0.99	0.13
22	Pinhole		0.15						-0.13
23	Spatial filter lens	1	2.2	1.0	1.507	1.24	0	0.98	0.13
24	Mirror	5	2.2	0.0	1.507	1.24	0	0.98	0.13
25	Rod amplifier	1	2.2	48.0	1.523	1.05	20	18.48	0.16
26	Polarizer	2	2.2	0.3	1.507	1.24	0	0.98	0.16
27	Faraday rotator	1	2.2	1.5	1.670	2.10	20	0.94	0.16
28	Polarizer	2	2.2	0.3	1.507	1.24	0	0.96	0.17
29	Wave plate	1	2.2	1.0	1.500	1.00	0	0.05	0.17
30	Rod amplifier	1	2.2	48.0	1.523	1.05	20	18.48	0.19
31	Spatial filter lens	1	2.2	1.0	1.507	1.24	0	0.98	0.19
32	Pinhole		0.15						-0.06
33	Spatial filter lens	1	2.2	1.5	1.507	1.24	0	0.99	0.19
34	Hard aperture	1	0.6	0.0	1.000	0.00	0	0.08	0.19
35	Spatial filter lens	1	0.6	0.5	1.507	1.24	0	0.98	0.19
36	Pinhole		0.15						0.00
37	Spatial filter lens	1	3.4	1.0	1.507	1.24	0	0.99	0.19
38	Mirror	2	3.4	0.0	1.507	1.24	0	0.98	0.19
39	Rod amplifier	1	3.4	48.0	1.523	1.05	20	0.92	0.20
40	Polarizer	2	3.4	1.5	1.507	1.24	0	0.98	0.20
41	Pockels isolator	1	3.4	11.0	1.500	1.00	17	0.87	0.20
42	Polarizer	2	3.4	1.5	1.507	1.24	0	0.98	0.20
43	Rod amplifier	1	3.4	48.0	1.423	1.05	20	18.49	0.21
44	Polarizer	2	3.4	1.5	1.507	1.24	0	0.98	0.21
45	Faraday rotator	1	3.4	1.0	1.670	2.10	20	0.94	0.21
46	Polarizer	2	3.4	1.5	1.507	1.24	0	0.98	0.21
47	Pockels isolator	1	3.4	11.0	1.500	1.00	17	0.87	0.22
48	Mirror	4	3.4	0.0	1.000	0.00	0	0.99	0.22
49	Hard aperture	1	2.6	0.0	1.000	0.00	0	0.58	0.22
Nonoptimized chain total								70.65	0.22
Output energy =								7.07 mJ	

Input energy: 0.10 mJ
Pulse duration: 0.10 ns
Absorption loss per cm: 0.15%
Scatter loss per surface: 0.30%

energy; the system accomplishes such outputs by means of three synchronized oscillators and a regenerative amplifier.

Once triggered by Novette's computer control system, the master oscillator system provides the subnanosecond timing and trigger pulses required by the laser and target systems. The complexity of the master oscillator is evidenced by the beam diagnostics, automatic alignment features, and pulse shaping capability of the system. Figure 2-59 is a measured beam profile obtained at the output of the preamplifier. As shown in Table 2-8, this beam is apertured to remove all of the beam except the center at component 34 in the preamplifier and at the input to power-amplification component 49. Spatial filtering smooths the beam profile until it has less than 10% intensity ripple from the ring structure when presented to the two Novette power-amplifier chains.

At the output of the preamplifier, the beam is split in two. Relative pulse timing is adjusted by variable-path-length sections, and the beam amplitude is regulated by wave plate and polarizer pairs. The wings of the beams are truncated so that only the central 2.7-cm diameter of the beam is allowed to pass into the amplifier chain. Each beam passes through the components listed in Table 2-9. The energies and fluences shown are appropriate for the propagation of 100-ps-duration Gaussian pulses typically used during activation. Data are recorded at several sensor packages, including intermediate locations such as the input to the 46-cm amplifiers. Amplifiers, isolators, and spatial-filter beam-expanders comprise the three generic, modular components for each stage in the power-amplifier chain.

Figure 2-60 is an optical schematic of a Novette power-amplifier chain. The

Table 2-9. Novette south laser chain components.

No.	Component	No. of components	Beam diam (cm)	Glass thickness (cm)	Refraction index	Nonlinear index	Bank kV	Net gain	Energy (J)	B-incident	B-incident total	Flux (J/cm ²)	Experiment (J)
1	Mirror	2	3.4	0.0	1.000	0.00	0.0	0.98	0.02	0.00	0.00	0.00	
2	Polarizer	1	3.4	0.0	1.000	0.00	0.0	0.49	0.01	0.00	0.00	0.00	
3	Mirror	1	3.4	0.0	1.000	0.00	0.0	0.99	0.01	0.00	0.00	0.00	
4	Wave plate	1	3.4	1.0	1.450	1.00	0.0	0.99	0.01	0.00	0.00	0.00	
5	Mirror	2	3.4	0.0	1.000	0.00	0.0	0.98	0.01	0.00	0.00	0.00	
6	Polarizer	1	2.6	0.3	1.507	1.24	0.0	0.98	0.01	0.00	0.00	0.00	0.0083
7	Hard aperture	1	2.6	0.0	1.000	0.00	0.0	0.98	0.01	0.00	0.00	0.00	
8	Spatial filter lens	1	2.6	0.5	1.507	1.24	0.0	0.92	0.00	0.00	0.00	0.00	
9	Pinhole		0.15						0.00		0.00	0.00	
10	Spatial filter lens	1	3.6	0.5	1.507	1.24	0.0	0.99	0.00	0.00	0.00	0.00	
11	Polarizer	2	3.6	0.3	1.507	1.24	20.0	0.96	0.00	0.00	0.00	0.00	
12	Rod amplifier	1	3.6	40.0	1.523	1.05	20.0	19.18	0.09	0.04	0.04	0.01	
13	Pockels isolator	1	3.6	11.0	1.500	1.00	17.0	0.87	0.08	0.02	0.06	0.01	
14	Beam splitter	2	3.6	0.5	1.507	1.24	0.0	0.97	0.07	0.00	0.06	0.01	0.0715
15	Spatial filter lens	1	3.6	0.5	1.507	1.24	0.0	0.92	0.07	0.00	0.06	0.01	
16	Pinhole		0.15						0.07		-0.06		
17	Spatial filter lens	1	8.8	1.0	1.507	1.24	0.0	0.99	0.07	0.00	0.06	0.00	
18	Disk amplifier	6	8.8	2.4	1.523	1.05	20.0	6.24	0.42	0.01	0.07	0.01	
19	Polarizer	2	8.8	1.0	1.507	1.24	0.0	0.96	0.40	0.00	0.08	0.01	
20	Faraday rotator	1	8.8	1.5	1.670	2.10	20.0	0.94	0.38	0.00	0.08	0.01	
21	Polarizer	2	8.8	1.0	1.507	1.24	20.0	0.96	0.36	0.00	0.08	0.01	
22	Disk amplifier	6	8.8	2.4	1.523	1.05	20.0	6.32	2.28	0.06	0.14	0.04	
23	Spatial filter lens	1	8.8	1.2	1.507	1.24	0.0	0.92	2.10	0.01	0.16	0.05	
24	Pinhole		0.05						2.08		-0.10		
25	Spatial filter lens	1	14.4	1.2	1.507	1.24	0.0	0.99	2.06	0.00	0.16	0.02	
26	Mirror	2	14.4	0.0	1.000	0.00	0.0	0.98	2.02	0.00	0.16	0.02	
27	Faraday rotator	1	14.4	1.0	1.670	2.10	20.0	0.94	1.89	0.01	0.17	0.02	
28	Polarizer	2	14.4	1.0	1.507	1.24	20.0	0.96	1.82	0.01	0.17	0.01	
29	Disk amplifier	4	14.4	3.0	1.523	1.05	20.0	4.19	7.63	0.07	0.25	0.04	
30	Spatial filter lens	1	14.4	1.4	1.507	1.24	0.0	0.92	7.02	0.02	0.27	0.06	
31	Pinhole		0.075						6.95		-0.11		
32	Spatial filter lens	1	20.0	1.9	1.507	1.24	0.0	0.99	6.88	0.01	0.28	0.03	
33	Beam splitter	1	20.0	2.0	1.507	1.24	0.0	0.97	6.67	0.01	0.29	0.02	6.61
34	Polarizer	1	20.0	1.0	1.507	1.24	0.0	0.98	6.54	0.01	0.30	0.02	
35	Faraday rotator	1	20.0	1.0	1.670	2.10	20.0	0.94	6.15	0.01	0.31	0.03	

schematic also shows the evolutionary transition in beam size from the old Janus system, to Argus and Shiva, and finally to Nova and Novette. The beam exits the master oscillator and is transmitted to the 2.7-cm input aperture, where it passes through a spatial filter and is expanded to 3.75 cm before being amplified by a single 48-cm glass rod amplifier. After the rod amplifier, the beam passes through a series of beam-expanding spatial filters, Faraday isolators, and glass disk amplifiers. The 9.4-cm and 14.5-cm-clear-aperture amplifiers, as mentioned earlier, were removed from Shiva and fitted with phosphate glass disks. The 20.8-cm-aperture amplifiers use the more efficient Nova box design.

All Novette components through the first 20.8-cm amplifier are equivalent to those previously contained in a single Shiva system chain. By substituting Nova-style 20.8-

cm-diam amplifiers with spatial-filter optic surfaces coated with a high-damage-threshold neutral solution, the 20.8-cm-diam beam energy at the 20.8-cm diameter can be increased to more than 3.5 times the energy of a Shiva beam⁷⁷ (refer to "Optical Components for Nova," earlier in this section). Following the 20.8-cm stages, the beam is expanded to 31.5 cm before passing through the final Faraday-rotator isolator in the chain. The beam is then turned by two mirrors to direct it through the final leg in the chain, where it is further amplified by four 31.5-cm-aperture amplifiers and three 46-cm-aperture output amplifiers. Each of these amplifiers contains two disks. We have made provisions to add a 31.5- and 46.0-cm amplifier to both chains (dashed lines in Fig. 2-60) if we can obtain output optics capable of withstanding the additional flux generated by the increased

Table 2-9. Continued.

No.	Component	No. of components	Beam diam (cm)	Glass thickness (cm)	Refraction index	Nonlinear index	Bank kV	Net gain	Energy (J)	B-incident	B-incident total	Flux (J/cm ²)	Experiment (J)
36	Polarizer	1	20.0	1.0	1.507	1.24	0.0	0.98	6.02	0.00	0.31	0.02	
37	Disk amplifier	3	20.0	2.5	1.523	1.05	20.0	2.05	12.34	0.05	0.36	0.04	
38	Disk amplifier	3	20.0	2.5	1.523	1.05	20.0	2.04	25.23	0.11	0.47	0.08	
39	Disk amplifier	3	20.0	2.5	1.523	1.05	20.0	2.03	51.28	0.21	0.68	0.15	
40	Spatial filter lens	1	20.0	1.8	1.507	1.24	0.0	0.92	47.17	0.09	0.77	0.21	
41	Pinhole		0.10						46.70		-0.50		
42	Spatial filter lens	1	30.3	2.8	1.507	1.24	0.0	0.99	46.23	0.06	0.82	0.09	
43	Polarizer	1	30.3	2.0	1.507	1.24	0.0	0.98	45.31	0.03	0.86	0.06	
44	Faraday rotator	1	30.3	2.0	1.670	2.10	20.0	0.94	42.59	0.06	0.91	0.08	
45	Polarizer	1	30.3	2.0	1.507	1.24	0.0	0.98	41.74	0.03	0.94	0.05	
46	Mirror	2	30.3	0.0	1.000	0.00	0.0	0.97	40.49	0.00	0.94	0.08	43.7
47	Disk amplifier	2	30.3	4.3	1.523	1.05	20.0	1.84	74.39	0.16	1.10	0.10	
48	Disk amplifier	2	30.3	4.3	1.523	1.05	20.0	1.83	136.02	0.30	1.40	0.18	
49	Disk amplifier	2	30.3	4.3	1.523	1.05	20.0	1.81	246.54	0.54	1.95	0.32	
50	Disk amplifier	2	30.3	4.3	1.523	1.05	20.0	1.79	440.29	0.97	2.91	0.57	
51	Spatial filter lens	1	30.3	2.6	1.507	1.24	0.0	0.92	405.06	0.47	3.38	0.80	
52	Pinhole		0.15						401.01		-2.61		
53	Spatial filter lens	1	44.2	3.7	1.507	1.24	0.0	0.99	397.00	0.30	3.68	0.37	406.0
54	Disk amplifier	2	44.2	4.3	1.523	1.05	20.0	1.62	643.12	0.69	4.37	0.39	
55	Disk amplifier	2	44.2	4.3	1.523	1.05	20.0	1.60	1029.50	1.11	5.48	0.63	
56	Spatial filter lens	1	44.2	3.7	1.507	1.24	0.0	0.92	947.14	0.74	6.22	0.88	
57	Pinhole		0.23						937.66		-2.83		
58	Spatial filter lens	1	71.2	5.7	1.507	1.24	0.0	0.99	928.29	0.41	6.63	0.33	
59	Mirror	3	71.2	0.0	1.000	0.00	0.0	0.96	891.16	0.00	6.63	0.32	
60	Window	1	71.2	4.0	1.507	1.24	0.0	0.92	819.86	0.27	6.90	0.29	
61	Second harmonic generator	1	71.2	1.8	1.500	1.00	0.0	0.88	724.70	0.09	6.98	0.26	
62	Window	1	71.2	4.0	1.507	1.24	0.0	0.98	710.21	0.22	7.21	0.25	
63	Focus optics	1	71.2	12.7	1.550	2.48	0.0	0.91	646.29	1.30	8.51	0.23	
Nonoptimized chain totals								32	314.38	646.29	8.51		

Input energy: 20.00 mJ
Pulse duration: 0.09 ns
Absorption loss per cm: 0.15%
Scatter loss per surface: 0.30%

Novette

amplification. The final spatial filter in the chain expands the output beam to 74 cm to lower the flux on the turning mirrors, frequency-doubling array, and focusing optics. The frequency-doubling array, located between the final turning mirror and the target-focusing optics, is followed by a 1.05- μm beam dump to remove residual fundamental light.

Some changes from the baseline design were required in the Novette chains to fit the system into the existing building. Our modification included adding two 14.5-cm

folding mirrors, changing the location of the 31.5-cm rotator, and configuring the final turning mirror in each chain; otherwise, the chains are identical to the Nova chains.

The MALAPROP laser-design code has been used to predict the ultimate performance of Novette.⁷⁸ Figure 2-61 summarizes the calculated capabilities of a single Novette beam. The upper curve illustrates the maximum 1.05- μm energy projected by MALAPROP, while the lower band represents our best estimate of the 0.53- μm output energy vs pulse duration.

Figures 2-62(a) through (f) show the major components and their locations in the Novette laser chains. A typical spatial filter is shown in Fig. 2-62(a). Locations of all spatial filters in both chains are color-highlighted in the illustration accompanying the photograph. All other components in the remaining figures are similarly identified. A cutaway of a Novette Faraday-rotator isolator is shown in Fig. 2-62(b). The calibrated flange on the end of the isolator permits accurate adjustment of the unit when the system is initially aligned or when the isolator is replaced. Figure 2-62(c) is a cutaway of a typical Shiva-style cylindrical disk amplifier fitted with phosphate

Fig. 2-59 Beam profile at one of Novette's chain input sensors.

Fig. 2-60. Optical schematic of Novette power-amplifier chain.

(a)

(b)

(c)

(d)

(e)

Fig. 2-62. (a) Novette spatial filters; (b) Novette Faraday-rotator isolators; (c) Novette cylindrical disk amplifiers; (d) Novette 20.8-cm amplifiers; (e) Novette 31.5-cm amplifiers; (f) Novette 46.0-cm amplifiers.

glass for Novette. High-power beam requirements in Novette are accommodated by the 20.8-, 31.5-, and 46.0-cm aperture amplifiers shown in Figs. 2-62(d) through (f). The amplifiers store approximately 2% of the electrical energy in the capacitor bank as disk inversion at pulse-propagation time. The beams' single pass extracts up to 40% of this energy in the form of light energy from the phosphate glass disks. Average output flux during high-energy operation is typically from 3 to 3.5 times the saturation fluence of 4.3 J/cm^2 . Depumping from ASE within the disk imposes a size limitation on the glass disks in the amplifiers. With a good match of the index for the absorbing cladding on the disk edges, the ASE depumping limits the gain-length

product, gd , to between 3.0 and 4.0 along the major axis of the disk. To preserve high energy storage and efficiency, we have split the disks along their minor axis in the 46-cm amplifiers. At maximum gain, these disks operate with $gd = 3.2$. We have designed the Nova 46-cm output amplifiers so that they are equipped with flashlamps only on the sides facing the disks. Each amplifier is pumped with a nominal stored energy of 500 kJ in the capacitor bank at 20 kV, to produce nominal stored inversion energy of 8.2 kJ with a small-signal gain of 1.7.

Figure 2-63 illustrates the measured and calculated beam profile at the 20.8-cm amplifier diagnostic station in the laser chain. The measured and calculated profiles are similar in shape and amplitude. Novette's characteristic four concentric rings are formed by filtering the hard-apertured input beam at the first spatial filter in each chain. Passive phase errors, not included in the beam-profile calculation, begin to wash out much of the diffractive structure, leaving only the prominent outer ring by the time this beam exits the power amplifier. In practice, minute thermal changes in components, small misalignments, and minor density fluctuations in the nitrogen gas filling the beamline all modify the diffractive structure of the beam. When viewed as a statistical entity, the rms intensity fluctuation imposed upon the beam becomes a natural quantity to measure. Risk of local optical damage to the laser optics is presumed to depend only on the local fields

Fig. 2-61. Novette system performance at 1.05 and $0.53 \mu\text{m}$ projected by computer simulation.

Fig. 2-63. Beam profiles in the Novette 20.8-cm-amplifier stage.

experienced. Comparison of measured beams with simulations using "realistic" component imperfections has meaning only in a statistical sense.

Figure 2-64 compares calculated and measured fluences on a typical experimental trial. While the simulations may fail to exactly match actual measurements, they are adequate predictors that allow apparently safe system design and operation. It remains to be seen whether gain saturation, spatial-filter pinhole closure, and other laser system engineering issues are correctly modeled. The Novette test bed will be used in the coming months to study these questions. In "Test and Activation Results," later in this section, we describe high-power propagation tests performed on Novette during activation.

In Fig. 2-65, we show the spools that join the two laser chains to the target chamber and some of the diagnostic equipment interfacing the chamber. The spools are the large, circular assemblies connected on the left and right of the target chamber. Each contains an apodizer, frequency-doubling array, beam dump, and doublet focusing lens. The beams are reflected to the target chamber just beyond the output of the sensor packages. Approaching the target chamber from the east and west, the beams pass through the apodizers near the entrance to the spools. Bead-blasted bands softly

shadow the interstices between crystals in the frequency-doubling arrays. These low-damage-threshold regions of the array assembly are protected from the incident laser flux, and, by preventing diffraction from uncontrolled phase discontinuities between crystals, the focusing lenses are shielded from potentially damaging intensity ripples.

To produce 2ω light, Novette requires doubler optics with a diameter greater than 74 cm. To meet this requirement, we designed frequency-conversion arrays consisting of smaller crystals, which could be constructed in time for the initial activation of Novette. We built one 5×5 -element array with 15×15 -cm crystals and one 3×3 -element array with 27×27 -cm crystals (see "Frequency Conversion and Focusing System," earlier in this section). The crystals and egg-crate support structure are sandwiched between BK-7 windows that have neutral-solution, antireflection coatings on the outer surfaces. The window design ensures that all KDP array elements are oriented within $\pm 30 \mu\text{rad}$. Losses at the KDP and window interface are reduced by an index-matching fluid. The faces of KDP plates for each array have been precisely oriented to the phase-matching angle, and the surfaces were diamond-machined. Figure 2-66 shows the 5×5 -KDP crystal array used in Novette. The 3×3 -array has similar construction to that of the array

Fig. 2-64. Novette south-chain laser flux for a typical experiment.

shown in Fig. 2-66. Early harmonic-conversion data taken from the first Novette experiments are described in "Test and Activation Results," later in this section.

Beyond the frequency-conversion array, approximately 99% of the residual 1ω light is removed by a full-aperture filter (beam dump). The 2ω light is transmitted, almost unattenuated, and concentrated on the target by the $f/4$ doublet focusing lens. Measurements to survey the focal region will be carried out at high power in later experiments; in anticipation of future measurements, we have used MALAPROP to calculate instantaneous beam profiles near the center of the target chamber. The accumulated active and passive phase errors in the laser system and the crystal array are expected to result in the highly structured beam profiles shown in Fig. 2-67. Nevertheless, the first measurements (described in "Test and Activation Results" at the end of

this section) show an acceptable focal-energy distribution.

Summary. The rapid evolution of laser-system design has led to the construction of Novette. in only 13 months, we have constructed a relatively compact two-chain laser, with each chain output exceeding the output of the 20 chains in the old Shiva system. Novette has been built in less than one-third the time required to construct Shiva, and will require less than half the personnel to operate. The evolutionary progression that led to Novette has produced what could be characterized as a revolutionary step in glass-laser-system development.

The Novette laser system now provides a high-power target-experimental capability at the second harmonic of Nd:glass while delivering more than 10 TW to targets. The system has tested much of the Nova laser design and has provided the first test of

Fig. 2-65. Novette's final optical train and target area.

harmonic generation in large-aperture, multiple-element arrays of KDP. Novette now allows early tests of Nova target physics, 3ω generation, and focusing systems.

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Construction

At the beginning of 1982, the laser bay had been painted and the building facilities renovated in preparation for the Novette installation. The new target-chamber spaceframe had been delivered and assembled at one end of the laser bay. In the MOR, the optical tables salvaged from the Argus system had been sited and aligned before the computer flooring, air filters, and electronic-equipment racks were installed. Various optical components had been aligned on their tables, while other optical equipment awaited assembly. In the basement, flash-lamp circuits, capacitors, pulse forming net-

works, fuses, and other component assembly and testing had started. The VAX 11-780 computer system was delivered for installation, and equipment racks for the programmable, touch-sensitive control panels were sited in the control rooms.

During January and February, the laser spaceframe was set up and we started installing high-energy electrical cables, nitrogen plumbing, and the electrical cables and

Fig. 2-66. Novette 5×5 -array of KDP crystals.

Fig. 2-67. Calculated Novette focal-region energy distribution for a 5-TW beam.

optical fibers for the computer control system. In the MOR, some of the small optical components were assembled while power supplies, timing circuits, and control-system interfaces were installed in the equipment racks. Larger laser-bay optical components were prepared in the clean room before installation in the laser bay. The YAG crystal from a commercial cw laser was replaced by a YLF crystal and tested. We installed the capacitor-bank-charging power supplies salvaged from the Shiva system and connected them to minicomputers in the main control room. Control racks for laser and target-chamber equipment were installed beneath the north and south sides of the target spaceframe and at the east and west ends of the laser spaceframe.

In March, the spaceframes were painted and surveyed and the laser bay cleaned. We determined the location of the laser beamline with respect to the spaceframe rails and installed laser component mounting hardware. Part of the MOR beamline was laid out and optical components were installed to direct the beam from the cw YLF alignment laser into the laser bay. The first alignment-sensor package was installed and made operational. Later in March, we

began final testing of the rod-amplifier flashlamp banks and exercised the power-conditioning control system. The touch-sensitive panels were installed in the control room, and software development proceeded on the computer programs that link the panels to the laser system.

By the second week in April, the laser bay was ready to receive laser optical components and the rebuilt Shiva target chamber. Figure 2-68 shows the renovated target chamber shortly after it was installed in the support cradle on the target spaceframe.

During the same period, we installed the first of the laser-bay alignment sensors, mirrors, rod amplifiers, and Faraday rotators on the beamline. This first alignment sensor was then used to maintain the beam alignment through the south side of the preamplifier. A week later, we started the first operational testing of the rod amplifiers and Faraday rotators. We measured amplifier small-signal gains and trigger timing, and determined the operating voltage required for 45° rotation of the optical polarization in the Faraday rotators. At the same time, we installed the second cw YLF alignment laser at the west end of the laser center spaceframe and steered the preamplifier

Fig. 2-68. Modified Shiva target chamber installed in its new home in Novette.

output to the insertion points in the two laser chains. This step allowed work on the amplifier sections to proceed in parallel with the preamplifier alignment. By the end of the month, we had installed all active components through the second 9.4-cm disk amplifier in the south chain and had started to align the spatial filters. Meanwhile, installation of the nitrogen plumbing was nearly complete in this section of the south chain, and preparations were completed in the power-conditioning system to fire a full shot to the 31.5-cm mid-chain sensor as soon as the 20.8-cm-aperture units were installed and made ready for testing.

Early in May, we began integrating the control system with the laser system. Initially, this affected only the laser-diagnostic and power-conditioning activities; however, much of the controls-group effort was still directed toward perfecting the system's basic structural features rather than direct control of the laser. By mid-May, the MOR installation had reached a point where pre-amplifier pulse-propagation tests could be interspersed with the small-signal gain measurements. In late May, the pace of activity increased when we started installing the target-chamber vacuum system.

In June, we demonstrated that the vacuum chamber and its pumping system were leak-tight. Off-line, the NS processing facility was activated and some of the processed output-spatial-filter lenses were delivered. Meanwhile, we started a full-production effort on the crystals for the frequency-doubling arrays.

One of our testing objectives was to propagate 100 ps pulses through a complete laser chain in preparation for the x-ray laser target experiments. To meet this objective as early as possible, we focused our attention on completing the south laser chain, thereby allowing us to propagate pulses to the south mid-chain sensor by the end of June. At the same time, the clean-room crew improved its component delivery schedule by assembling all of one type of component before proceeding to the next type. This approach allowed us to complete some phases of the north chain installation earlier than anticipated. By the end of June, we had installed most of the north-chain components through the 20.8-cm-aperture

disk amplifiers, tested some of the smaller disk amplifiers for small-signal gain, and prepared for later pulse-propagation tests.

On July 9th, we reached our first definitive operational milestone when we fired a 1 ω laser shot at 1 TW (96 J); pulse duration was 93 ps. The shot was measured at the south mid-chain sensor and had approximately three times the power required at the 31.5-cm aperture for operation at 100 ps with a 10-TW laser chain output. With this successful test, Novette became one of the most powerful lasers in the world.

Through the rest of July, we worked on installation of the large output amplifiers in the south-chain and completion of the north driver chain. We assembled several of the south-chain 31.5-cm disk amplifiers before installing all of the 31.5-cm-aperture turning mirrors, in anticipation of possible problems associated with the new design. This alteration to the schedule proved beneficial when we experienced problems with the laser-glass mounting techniques. To overcome these problems we developed new procedures for mounting the glass, to reduce stress and propagated wave-front distortion. During July, we installed the target local illuminator on the east end of the center laser spaceframe. We also started testing the software that controls the touch-sensitive panels in the main control room (shown in Fig. 2-69).

In August, the four target-alignment viewers were mounted on the target chamber. This operation required careful alignment of the viewers on target-chamber center, since proper location of the target at chamber center is critical to performance of target diagnostics. The rails supporting the final laser optical components, i.e., the frequency-doubling crystals, target-illumination lenses, and beam dumps, were installed east and west of the target chamber. The alignment of the rails along the intended beamline was another critical factor, since their misalignment could cause focus motion of the target-illumination lenses to generate a transverse offset of the laser-beam focus.

The electronic devices required to control the target-chamber equipment and output-alignment sensor were installed in the equipment racks under the target

spaceframe in August. To protect the optical surfaces from dust, we enclosed the beam path in beam tubes. During the month, we completed alignment of the 74-cm-diam output in the final spatial filter in the south chain. The alignment was completed without installation of the 46-cm-aperture disk amplifiers, since these units had not yet been delivered. At this point, we shifted our efforts toward completing the north-chain installation up to the mid-chain sensor.

By mid-September, the two laser chains were at the same stage of completion. Our final test consisted of propagating pulses to the output of the spatial filter between the 31.5- and 46-cm disk amplifiers in both chains. Calorimeters were installed at the spatial-filter output in both chains to measure the pulse power, recorded at 300 to 400 J (3 to 4 TW). During September, we performed a short series of tests, including pulse-energy measurements, with 825-ps pulses to check the performance of the completed portion of the system. Our measurements in these tests were made with calorimeters again mounted at the output of each 46-cm spatial filter. The north-chain shots produced 1.76 kJ, while the south chain yielded 2.07 kJ. In the meantime, several output turning mirrors were installed in the south chain, allowing us to start calibrating the south output sensor.

September and October were marked by rapid progress on the central control system as the touch-sensitive panels were integrated with laser-system operations. During this time, the laser and target-diagnostic software was successfully linked with the power-conditioning controls for the laser system. In Fig. 2-70, we show the extent of equipment installation in the target-diagnostics control room at the end of 1982. The target-diagnostic console, with its touch-sensitive control panel, is in the center of the figure, with the other diagnostic equipment installed in the racks shown in the background.

In October, we successfully fired the 46-cm disk amplifiers. We considered this a significant event, since these amplifiers represented our largest step beyond the laser technology with which we were familiar. At the end of October, all optical components intended for use in the first series of experiments had been installed in the south chain of the laser system. The laser configuration for this short-pulse series contained two 46-cm-aperture disk amplifiers in each chain. We tested the south chain capabilities during calibration of the output sensors by propagating a pulse to the target chamber entrance where we measured 1.097 kJ (11.1 TW). We then concentrated on completing installation of the north chain. In

Fig. 2-69. Main control room; control consoles are shown in left and right foreground.

late November, we successfully tested the north chain with a 612-J (6 TW) shot.

During the last two months of 1982, we achieved significant progress toward completing installation of the target chamber and target-diagnostics equipment, laser-beam focusing optics, frequency-doubling crystal arrays, and beam-path diagnostic hardware. We also installed a set of x-ray spectrographs for use with the target-diagnostics electronics as the primary diagnostics for early experiments. Secondary diagnostics, such as x-ray microscopes and pinhole cameras, were functionally tested for integration with the system. The filter-fluorescer spectrometer (FLEX) and Dante, major diagnostics required for the later experiments, were also installed.

The target-lens positioners and the $f/4$ lens pairs were the final major mechanical additions to be mounted on the target chamber in November. We then tested the vacuum integrity of the complete assembly. The lenses were aligned to the established target-chamber beamline prior to final adjustment of the output-alignment sensors, which required the reflection from the uncoated output surface of the second lens of the doublet. Later, we fired the laser oscil-

lator and preamplifier sections to perform a series of low-energy south-chain shots.

Metallic-foil targets were used in the tests to time the triggers for the time-resolved x-ray and laser-beam diagnostics. Comparison of these timing data with similar off-target north-chain shots enabled us to make a preliminary adjustment to equalize the optical path lengths for the two chains.

In early December, the 5×5 crystal array shown in Fig. 2-71 was installed in the south laser chain, east of the target chamber. After several days of developing the crystal alignment procedure, we fired several 2ω test shots through the array; all shots were fired at less than half the operational energy level (100 to 250 J). During these tests, two shots were made directly into a calorimeter to test conversion efficiency. The results of these two shots are described in "Test and Activation Results," at the end of this section. On the other shots, we used the burn-paper technique to obtain profiles of the 1ω and 2ω beams.

Figure 2-72 shows the Novette installation at the end of 1982. The system, as shown, is operational with nearly all of the installation complete. The connections between the main and target-diagnostics

Fig. 2-70. Target-diagnostics control room during installation.

Fig. 2-71. Installation of the 5×5 frequency-conversion crystal array.

Fig. 2-72. Novette laser bay at the end of December 1982.

control rooms and the time-resolved x-ray diagnostics are complete, and all scheduled laser components are in use. The control software allows us to align, charge, fire, and measure the results of firing the laser, via the touch-sensitive panels in the main con-

trol room. We have satisfactorily propagated short pulses through the system. Long-pulse propagation has been checked at approximately half energy up to the input of the 46-cm disk amplifiers. The laser alignment devices have performed satisfactorily and reliably indicated the adjustments needed for setting the laser mirror angles and pinhole sizes. The photodiodes and calorimeters for measuring laser beam energies have been calibrated. The frequency doubling-array alignment techniques have been tested and appear to agree with the theoretical conversion efficiency calculations for the 5×5 crystal array. Our progress on completing the Novette installation in 1982 was exemplified by three successful test shots that established the system as the most powerful pulsed laser in the world.

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Electronic Systems

The electronic systems for both the laser and target facilities were installed and

integrated with the Novette operating system during 1982. At the beginning of the year, installation of the control-room equipment had just started while the other areas in the facility were empty. By the end of the year, all electronic systems had been installed and we were nearly ready to start both laser and target operations. In the following sections, we highlight the hardware, software, and operational progress made on the Novette electronic subsystems.

Central Controls. Novette's control system is a smaller version of the system to be used in Nova (see Nova "Control System," earlier in this section). By the start of 1982, the central control electronics devised to support subsystem development in 1981 were installed. During 1982, we concentrated on activating control support for the four major subsystems shown in Fig. 2-73.

By early spring, the VAX 11-780 central control computer had been installed in the Novette computer room and was functioning with a standard V2.5 VMS operating system. The computer facilities used for controls development were duplicated so that problems would be minimized with the transfer to Novette operations. Two Nova style operator consoles were delivered and installed in the control room. During this time, we established a work area with editing terminals and printers in the control room. Throughout the remainder of the year, we steadily increased use of this area for software installation and activation.

For the first time, the central-control tools that had been developed were put to use in an operational environment.⁷⁰ The Novette control system was activated with the most up-to-date Nova central-control software, including the Novanet support package, console graphics, network shared memory (NSM), the TEXTIO input/output package for Praxis, and the Oracle data-base management system (see "Oracle Data-Base Management System at FEA" in Section 5). The most current conventions were used to control modification and release of these software packages.

Master Oscillator. The master-oscillator control system provides the drive signals to the A O modulator, the Q-switch, fast-tuning system, and oscillator lamps. The A O modulator and the Q-switch are part of the rf processing subsystem. Novette has

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two A-O modulators, two Q-switch drivers, and four lamp drivers installed; one set is presently in use. Both the lamp driver and the rf output-amplifier packages are designed to improve operation, serviceability, and system safety. Monitor outputs and a distortion filter have been added to the rf amplifier. The lamp driver was rebuilt to improve component mounting. A capacitor-bank dump relay and an interlock isolation relay were added to the lamp driver.

The master-oscillator fast-timing system is the source of timing signals distributed over optical fibers throughout the laser system. Timing resolution for these signals is 1 ns with less than 150-ps jitter. Delay times are entered from a control and status panel. A multicolor CRT format displays the current status of all timing channels. Currently, there are 28 channels installed in Novette, with 21 channels in use.

Alignment. The 1ω and 2ω portions of the alignment systems described in the 1981 Annual Report⁸⁰ have been installed on both Novette beams. Operation of these systems has been satisfactory, and only minor changes have been required. The video format for all alignment data was a significant improvement over previous centroid detection systems. This format allows operators to use the alignment sensors as

versatile viewers for positioning other components during installation and for identification of beamline irregularities.

During installation of the system, we identified a polarization problem that caused distortion of the KDP angle-tuning curve. To trace the source of the problem, we started with the 2ω tuning signal generated by the small-diameter 1ω local illuminator beam passing through the center KDP crystal. In our original system configuration, this signal was returned to the local illuminator by a corner cube inserted after the crystal array. The signal passed back through the array and was reflected from the main beam-turning mirrors to the local-illuminator detector. Because the corner cube considerably changed the polarization state of the beam, and because KDP is birefringent, the signal returning to the local illuminator had polarization properties that changed when the KDP angle was changed. This caused the polarization-sensitive turning mirrors to return a varying fraction of the signal to the local-illuminator detector and distort the tuning curve. To eliminate these adverse polarization effects and increase the signal level, we moved the 2ω tuning-signal detector and fastened it to the centering-mirror inserter immediately following the KDP array. The detector is now attached so that the 2ω detector is inserted when the split centering mirror is partially removed from the beam. In Fig. 2-74 we show an example of the quality of tuning curve we can now obtain for 2ω shots (see "Nova Alignment and Diagnostics," earlier in this section).

Laser Diagnostics. The laser diagnostics implemented on Novette have performed well within their design specifications. The operation of the laser diagnostics on Novette has verified our diagnostic configuration designs and given us confidence in the ability of the hardware and software in the acquisition network to provide real-time diagnostics data on a larger (Nova) scale. The software for the laser-diagnostic data-acquisition system allows the operator to configure the system for a shot and choose the data-readout display via the menus on the touch-sensitive panels in the control room. As we continue to develop the software, higher-level data processing and control functions continue to be added to the diagnostic system.

Fig. 2-74. Tuning curve for a 5×5 array of KDP crystals, showing relative KDP frequency-conversion efficiency vs array orientation angle.

The fundamental electronic elements of the Novette laser diagnostics are shown in Fig. 2-75. All energy diagnostics have been activated and are operational. The stability of the chain-diagnostics calibration has remained within 2.5% during the three-month period following the start of laser operations.

The temporal diagnostics for switchout and Pockels-cell timing have been operating routinely for setup of the laser and performance monitoring. We have also found this system to be valuable in diagnosing occasional malfunctions of the oscillator switchout and Pockels-cell drivers. The streak cameras for recording 1ω and 2ω diagnostic information have been installed and used to accumulate data.

Determination of the optimum linear angle of the frequency-doubling arrays repre-

sents a new measurement requirement for laser diagnostics. In the new arrangement, the output sensor aligns a local oscillator output with the main beam axis. The signal frequency is then converted by the center element of the array and synchronously detected by a photodiode and amplifier. The detected signal is compared with a fraction of the local oscillator frequency that has been split and converted by a stationary crystal in the local illuminator package. The comparison removes fluctuations caused by output-power variations of the local oscillator. The crystal array is then scanned through an angle around the expected peak conversion angle to produce a tuning curve like that shown in Fig. 2-74. In this position, the array is oriented at the peak conversion angle prior to propagation of the main laser pulse. We have used this technique in a

Fig. 2-75. Basic elements of the laser diagnostic system, showing interconnection with the control system.

series of Novette shots; we achieved close agreement with the calculated frequency-conversion efficiencies.

We developed a system for detecting the modulation that a flat-plate target with a pinhole puts on a probe beam passing through the pinhole when the surrounding cryogenic target shroud is ejected. The measurement system determines the ratio of the transmitted signal with respect to a reference split from the 1.05- μm cw laser source and, when the shroud is ejected, produces negligible target movement in relation to the alignment tolerance.

Target Diagnostics. In Table 2-10 we show the diagnostics on-line at the end of

1982. The following data-acquisition components were on-line and operational at the end of the year

- Three front-end processors (FEPs).
- Seven CAMAC crates.
- Seven R7912 transient digitizers.
- Three 7903 oscilloscopes.
- 80 integrator channels.

We have the capability to quickly add Dante, photodiode, calorimeter, neutron scintillator, and neutron counting channels to the system; we can also easily add other instrumentation for fast and slow diagnostics. The data-acquisition system has proved very reliable. Operators control the system via touch-sensitive panels that allow them to display menus for system or diagnostic operation. The target-diagnostics trigger system uses fiber-optic signal transmission from the MOR. This system has proved very stable, while producing signal levels that have been more than adequate for normal operations.

Power Systems. The high-voltage power supplies and the energy-storage capacitor bank were installed and made fully operational. Construction on the capacitor bank started early in the year when high-voltage and control-signal cable trays were installed. This was followed by installation of the capacitor racks and, early in March, placement of capacitor modules on the rack shelves. To speed installation, we adopted production-line methods for assembly of the capacitor modules. By May, all the modules were installed and cabled to the ignitron switches and junction boxes in the laser bay.

During May, we installed and initially tested the power supplies. Our tests were followed by more comprehensive testing as the capacitor banks were integrated with

Fig. 2-9. A section of the completed Novette energy-storage capacitor bank.

Table 2-10. On-line Novette target diagnostics, December 1982.

X-ray diagnostic	Light diagnostic	Type	Detectors	
			No.	Type measurement
Dante H		XRD-31 diode	7	Spectral and time resolution
Filter fluorescer (3 m)		fluor/PM	18	Spectral resolution, time integration
8 $\times$ microscopes (2)		film	8	Spatial resolution, time integration
X-ray crystal streaked spectrograph		film	1	Spectral and time resolution
Spatially discriminating streaked spectrograph		film	1	Spectral, spatial, and time resolution
	Energy balance (2 ω)	PIN diode	20	Time integration

the remainder of the system. The separate rod-amplifier bank was the first operational segment in the power system. Checkout began with 25-kV testing of all high-voltage cable shields and center conductors to ground. Each bank module was charged to its rated voltage and fired into dummy loads on the PFN boards.⁸¹ The rod bank was operational by the middle of June and the entire capacitor bank was operational by the end of August. Figure 2-76 shows a completed section of the capacitor bank with capacitor modules installed in the racks.

Between September and the end of the year, the pulse-ionization lamp check (PILC) diagnostic system was activated. This system provides a method that ensures flashlamp circuits are operational prior to high-energy operations. The PILC system proved itself when it prevented the loss of at least two amplifiers by detecting broken flashlamps in advance of a laser firing.

Facility Systems. The personnel-safety interlock, security-door monitor, audio, video, nitrogen-gas, and vacuum systems were installed and activated during the

year. These support systems provide many functions essential to the operation of the laser and target facility.

We completed installation and activation of the personnel-interlock system controlling access to the facility during hazardous operations. Safety procedures are activated either automatically or by an operator in the event a hazardous area is penetrated. Taped messages are automatically played to warn personnel of safety hazards. The interlock system has continued to operate reliably with all inputs and outputs monitored by the control computers.

The security-door monitoring system allows security officers to monitor access to the facility from a central point using a computer-generated display. Switches have been installed on all access doors that were not previously monitored by the safety interlock system. The outputs from these switches are connected to a display in the central computer room, where an operator can quickly locate an open interlock.

The Novette work areas are divided into four audio zones. The audio system provides paging functions for the laser facility.

Fig. 2-77. Power-supply status display on the touch-sensitive operator panel in the control room.

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Taped announcements are activated by the safety interlock system. Computer control of this system allows specific safety messages to be announced in any number of the Novette work areas. It can also simultaneously play separate messages in the various work areas.

The video system connects 40 remotely placed cameras with 20 video monitors. Monitors are located in the control room, diagnostic rooms, and laser bay. A channel selector is connected to each monitor to allow the operator to select any video camera in the system. One output of the video switcher is routed to a digitizer connected to the control computer. For automatic alignment functions, the image of the laser beam is digitized and processed by the computer to determine the adjustments necessary for alignment. When the computer determines the required adjustments, the information is sent to the appropriate motor controller where the correction is made.

The nitrogen-gas system is used to purge and cool the numerous components in the

laser chain. By the end of the year, this system was installed and working under manual control. The design of an electronic system for controlling the nitrogen-gas system is nearly completed.

The vacuum system associated with the target chamber and target diagnostics is locally controlled by a dedicated computer and control panels. As new diagnostics are installed, new control panels will be added to the vacuum system. To ensure proper operation of the vacuum system, on-line diagnostics have been placed under computer control. At the end of 1982, the three roughing and high-vacuum systems for the target chamber and auxiliary pumping systems and their safety interlocks were operating under control of the computer. In addition, we had started operational use of four of the six cryogenic compressor systems planned for the installation.

System Operation. The system configuration allows operators to control the laser with either the touch-sensitive panels on the consoles or to manually operate the

Fig. 2-78. Sequencer status display on the touch-sensitive panel in the control room.

system with the controls on the equipment racks. When operating the system with the touch-sensitive panels, the operator selects options from a menu displayed on the center screen of either of the two consoles in the control room. Options are selected by touching specified areas of the video display. System status is shown in color, by position, or with text displayed on the control console monitors. Control of the new system from a central console has increased operator efficiency, since operators are no longer required to attend several control stations in the control room. Another control feature in the system allows the operator to align laser components by issuing instructions from the control console.

To prepare to fire a target shot, a number of tasks must be performed, starting with laser alignment. After the laser is aligned, the power-conditioning, target-diagnostics and laser-diagnostics components must be preset for the shot. To perform this task, the console can be used to load a previously defined configuration. If the configuration is different from configurations previously committed to the computer memory, the operator can perform alignment from the control console. When pre-setting the system, the operator must check the status of the laser high-voltage power supplies. This step can be performed from the control console using the display shown in Fig. 2-77. To fire the laser, the operator starts the sequencer. The sequencer is a software entity that schedules time events. When sequencer time equals zero, the laser is fired. Figure 2-78 is a display of the sequencer status display on the control console. In this display the status frame shows the sequencer event being performed, the countdown time, and other information associated with the shot. Following a target shot, the data acquired from the control system is stored on disk and later forwarded to the fusion experiments analysis facility (FEAF) for additional processing.

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Test and Activation Results

We designed and built Novette to provide a high-power target-irradiation facility with the second harmonic of Nd:glass. To provide this capability, the fundamental 1ω laser beam must be sufficiently uniform in intensity and low in divergence to be efficiently frequency-doubled (to $0.53\ \mu\text{m}$) with KDP crystal arrays and focused to a diameter of less than $200\ \mu\text{m}$ at the target.

During the activation of Novette, we performed a detailed characterization of the laser-amplifier-chain performance to ensure that the beam met our requirements. Our tests included use of a cw probe beam to perform gain measurements for individual amplifier stages, use of pulsed beams from the MOR to perform gain measurements and determine beam quality of the laser chains, and measurement of focusability of cw alignment and pulsed beams in diagnostics sensors. Most of the pulse data was obtained with nominal 100-ps FWHM pulses. A limited data set was established with 800-ps pulses. Short-pulse data included the laser characteristics at 1ω , the conversion efficiency of the frequency-doubling arrays, and the 2ω beam characteristics at the target chamber. The test results indicate that Novette meets its performance goals of 2ω power, energy, and focusability.

As each set of amplifiers was activated, we measured their gain for a full-aperture 1ω cw probe beam for the duration of the

Table 2-11. Measured and calculated gain for a full-aperture 1ω laser beam.

Amplifier (size and type)	20-kV bank energy (kJ)	Measured gain	Design gain
3.75-cm rod	41.5	$19.4 \pm 6\%$ (unclad) $25.9 \pm 5\%$ (clad)	25.0
7.4-cm disk	145.0	$6.6 \pm 6\%$	6.6
15.0-cm disk	215.0	$4.15 \pm 10\%$	4.2
20.8-cm disk	165.0	$2.0 \pm 10\%$	2.1
31.5-cm disk	310.0	$1.82 \pm 4\%$	1.9
46.0-cm disk	500.0	$1.75 \pm 4\%$	2.0

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Fig. 2-79. (a) Novette cw gain measurement for full aperture beam. (b) Background signal correction for full-aperture beam.

Shot No. 92102704
31.5-cm amplifiers—chain 1
Flashlamp voltage—20 kV
cw baseline level—40 mV
Average gain/amplifier—1.8

Shot No. 92102705 (background)
31.5-cm amplifiers—chain 1
Flashlamp voltage—20 kV
Without cw

$$G = \frac{\text{Signal} - \text{background}}{\text{baseline}} = \frac{425 - 3}{40} = 10.55$$

Fig. 2-80. (a) Best focus of cw alignment beam from the Novette south arm. (b) Silhouette of a 500- μm ball at the center of the target chamber.

Fig. 2-81. Scan of horizontal diameter of a cw alignment beam in the south chain of Novette. ▼

fluorescence. Table 2-11 summarizes the measured gains and the design gains for each size amplifier⁸² in Novette. The measured gains are nearly the same as the flashlamp pump pulse. The signal was monitored with a photodiode that was placed in a focused beam in a diagnostic sensor following the pulsed amplifiers. The signal peak, corrected for the fluorescence and flashlamp background, gives the total gain of the amplifiers, excluding transmission losses; the time of the peak signal gives the correct firing time for the amplifier flashlamps relative to the oscillator pulse. Figure 2-79(a) is a typical data set showing the amplified cw signal. Figure 2-79(b) is the correction of the background signal caused by flashlamp light and glass

design gains; however, the 46-cm amplifier gain was approximately 12% low. This slightly lower gain does not prevent the system from meeting its performance goals, but we are investigating design changes to improve the amplifier (see "Laser Design and Performance," earlier in this section).

The far-field viewing capability of the output sensor⁸³ allows the focused beam to be directly observed at the laser output. Alternatively, the opposing target-focusing lens can be used as an imaging element to allow viewing of the focal spot at the center of the target chamber with the opposing output sensor. Figure 2-80(a) is a video display, viewed by the north output sensor, of a 1- ω cw beam transmitted through the south laser chain frequency-doubling crystal array and focused in the center of the target chamber. The scale for the display is determined by Fig. 2-80(b), in which a back-lighted 500- μ m ball is silhouetted in the target chamber center and imaged through the same beam optics. Since the size of the spot is dominated by aberrations in the beamline optics and frequency-conversion array, the display should be a good representation of the south 2 ω beam focused in the center of the target chamber.

Figure 2-81 is a horizontal line scan through the center of this image with the beam diameter at the 0.5- $I_{\max}$ and 0.1- $I_{\max}$ intensity levels shown. This scan was produced by the VAX computer from data taken by the charge-coupled diode array. These typical displays indicate that Novette 2 ω beams can be satisfactorily focused for target-irradiation experiments.

The initial pulse-propagation studies, conducted with 100-ps pulse durations, provided data on the pulsed gain of the laser chains. From the energies measured in the four diagnostics sensors in each laser chain, we obtained the overall chain gain and input and output energies for the sections of the chains between sensors. The input sensor,⁸⁴ located after the rod amplifier, is calibrated to read energy at the rod output. The mid-disk sensor,⁸⁵ at the output of the 15- to 20-cm spatial filter, is calibrated to give energy at the input to the 20-cm rotator. The mid-chain sensor,⁸⁶ behind the first 31.5-cm turning mirror, is calibrated to give energy at the input to the 31.5-cm amplifier.

Fig. 2-82. Output vs input energies of the disk amplifiers in the Novette laser chains.

Fig. 2-83. Near-field photograph taken at the output sensor of Novette south beam.

The output sensor, behind the first output beam turning mirror, is calibrated to give energy after the final beam turning mirror.

We compared the gains measured in each section of the chain with the lumped parameter calculations derived with SNOBAL (our system-simulation computer code) for the net amplification. The calculation was performed with the gains measured with the cw laser and nominal transmission losses⁸⁷ for the beamline components. The measured gains in the first two sections of the chain exceeded the calculated gains, and

we believe the disagreement results from pessimistic transmission losses assumed for components. We have adjusted the trans-

mission losses for the two sections to bring them into agreement. Figure 2-82 shows the measured data for chain output energy vs measured input-sensor, energy and the calculated values before the transmission losses were adjusted for the two Novette beams.

Following the low-power characterization and calibration shots, we increased the input to the chains to characterize high-power performance. The chain-energy data from these shots are included in the data plotted in Fig. 2-82. Near-field photographs of the beam provided a sensitive monitor of chain performance when we increased the power levels. In Figs. 2-83 and 2-84, we can see that the near-field beam profile remained smooth as the beam output power was increased to 11.0 TW. In Fig. 2-83, taken in the output sensor, we can see evidence of an underlying rectangular structure in some regions of the south beam. Beam energy for this shot was 1.1 kJ, with a pulse duration of 99 ps. A densitometer scan across the diameter of the beam, perpendicular to the shadow of the 46-cm disk split converted to intensity, indicates a peak-to-valley beam modulation of approximately $\pm 5\%$. Figure 2-84 shows data derived from a beam scan, to show the statistical distribution of intensities. The curve in the figure indicates that 50% of the power in the beam is below 3.05 GW/cm² and less than 2% of the power is greater than 4.0 GW/cm².

Assuming the beam fills the 74-cm aperture, we have derived an effective fill factor

Fig. 2-84. The beam intensity-power distribution obtained from the 11-TW beam in Fig. 2-83.

Fig. 2-85. Near-field photograph taken at 12.4-TW output power.

Fig. 2-86. (a) Component arrangement to determine dual-beam high-power conversion efficiency of the frequency-doubling arrays. (b) Component arrangement for 2ω beam output-sensor calibration.

from the data illustrated in Fig. 2-84 with the definition: fill factor $\times$ area $\times$ average J/cm^2 = total energy. When Novette data are reduced with this method, the results suggest that the fill factor is in a range from 0.80 to 0.85.

When the frequency-doubling system became operational, we accumulated additional high-power performance data on the laser system and initial data on the frequency-doubling efficiency of the large-aperture crystal arrays. The data base consists of 17 single-beam 1ω shots with output power ranging from 2.5 to 6 TW, and two dual-beam 1ω shots with approximately 12 TW output power per beam; nominal pulse duration for all shots was 100 ps. System beam quality at the lower powers was excellent, with approximately 75% internal conversion efficiency in the crystal arrays. (Internal conversion efficiency is derived by division of the 2ω energy at the output face by the 1ω energy at the input face.) Figure 2-85 is a typical near-field photograph taken at 12.4-TW output power.

When we operated the beam beyond 12 TW, the near-field photographs taken in the output sensors revealed severe beam break-up from self-focusing instabilities,⁸⁸ and the internal conversion efficiency of the doubling arrays dropped to approximately 60%. The severe intensity modulation is

Fig. 2-87. Initial 1ω -to- 2ω power conversion on the first set of frequency-doubling shots.

Shot No.	South chain			North chain		
	Array input ^a $I_0(GW/cm^2)$	External efficiency	Internal efficiency	Array input ^b $I_0(GW/cm^2)$	External efficiency	Internal efficiency
92120304 ^c	0.76	0.39	0.56			
92121302 ^c	1.06	0.50	0.71			
93010502 ^d	1.05	0.49	0.71			
93010504 ^d	1.20	0.53	0.76			
93010603 ^d	1.1	0.49	0.70			
93010606 ^d	1.6	0.54	0.77			
93010609 ^d	1.0	0.55	0.78			
93010702 ^d	1.69	0.55	0.78			
93010705 ^d	1.18	0.53	0.76			
93011004 ^e	1.17	0.52	0.75			
93011103 ^e	0.74	0.49	0.70			
93011204 ^e	0.92	0.54	0.76			
93011408 ^d				0.92	0.50	0.69
93011507 ^d				1.28	0.55	0.75
93011509 ^d				1.65	0.56	0.77
93011709 ^d				1.61	0.56	0.77
93011803 ^e				1.58	0.58	0.78
93012111 ^c	3.40	0.41	0.59	3.12	0.33 ^f	0.45 ^f
93012409 ^c	3.33	0.42	0.60	3.44	0.46	0.63

^a5 $\times$ 5 crystal array.

^b3 $\times$ 3 crystal array.

^c 2ω calorimeter between beam dump and target chamber.

^d 2ω calorimeter behind target chamber ($T_{tc} = 73\%$).

^e 2ω target-shot energy measured by 2ω output sensor.

^fLow efficiency attributed to misalignment of north beam KDP-crystal array.

$$\text{External conversion efficiency} = \frac{2\omega \text{ energy on target}}{1\omega \text{ to array apodizer}}$$

$$\text{Internal conversion efficiency} = \frac{2\omega \text{ energy at KDP output face}}{1\omega \text{ energy at KDP input face}}$$

Table 2-12. Novette 2ω conversion efficiency.

Fig. 2-86. Internal conversion efficiency for a Type II KDP-crystal frequency-doubling array obtained from initial Novette tests.

Table 2-13. Initial high-power shots performed at 1ω and 2ω with Novette.

Shot No.	Pulse length (ps)	South beam		North beam		Total	
		$1\omega^a$ (TW)	$2\omega^b$ (TW)	$1\omega^a$ (TW)	$2\omega^b$ (TW)	$1\omega^a$ (TW)	$2\omega^b$ (TW)
93012111	95	12.4	6.1	11.4	4.3	23.8	10.4
93012409	94	12.2	5.9	12.6	6.9	24.8	12.8

^aPower incident on crystal-array apodizer.

^bPower measured by calorimeter between beam dump and focusing lens.

near the maximum spatial frequency passed by the last two spatial filters in the amplifier chains.

Figure 2-86 shows the arrangement of the frequency-doubling and target-chamber optics and the location of the calorimeters for most of the 2ω shots. For the first two shots with the south beam, and for the high-power shots, we measured 2ω energy with large-aperture calorimeters installed behind the 1ω beam dumps in Fig. 2-86(a). Most of the low-power shots were used to calibrate the 2ω energy sensors with the configuration shown in Fig. 2-86(b). For these shots, the 2ω beam was transmitted through the target chamber to the calorimeter behind the target chamber. The reflection from the back surface of the input-lens second element was transmitted back to the 2ω sensor for energy measurement. We corrected the calorimeter energy to the center of the target chamber ($I = 0.73$) and to the crystal array [$I = (0.73)(0.965)$] using the measured transmission of the target-chamber optics (measured chamber optics transmission = $73\% \pm 2\%$). The remaining low-power shots were single-beam target-irradiation experiments in which 2ω energy was measured by the 2ω sensor. In all cases, 1ω energy was measured with the 1ω sensor, which was calibrated to give energy at the

input to the crystal-array apodizer. The input pulse width was measured by a streak camera in the MOR.

Figure 2-87 summarizes the measured power conversion from 1ω power incident on the apodizer to 2ω power incident on the target. Table 2-12 lists all of the shots in this data set, giving the 1ω intensity incident on the crystal array and the external and internal conversion efficiencies. Figure 2-88 is the measured internal conversion efficiencies superimposed on curves obtained from calculated conversion efficiencies. The conversion efficiencies were calculated for a Gaussian temporal pulse incident on the array with various internal angles of detuning from the phase-matching angle.⁸⁹ The low-power data provide a good match with previously recorded information obtained from Argus and other sources.^{90,91} The principal reason for loss of conversion efficiency at high power appears to be severe intensity and phase modulation, caused by the self-focusing instability.

Table 2-13 provides the 1ω and 2ω powers measured from the two high-power shots. The 2ω power was measured by the calorimeter shown in Fig. 2-86. To compensate for the measured transmission of the Novette target-chamber optics, these measurements must be multiplied by 0.86 to obtain the power on the target. The low conversion efficiency on the north chain (shot 93012111) resulted from an approximate $130\text{-}\mu\text{rad}$ misalignment of the crystal array.

After reviewing the near-field photograph from the first high-power shot (Fig. 2-85), we attempted to reduce beam modulation on the next shot by using smaller spatial-filter pinholes in the four largest spatial filters. Table 2-14 lists the pinhole sizes, spatial-filter cutoff wavelengths, and dominant wavelengths observed on the near-field photographs for the two shots. After the pinholes were changed, the beam was successfully propagated through the smaller pinholes. The output near-field photographs produced during the shot failed to reveal evidence of pinhole misalignment. Reducing the pinholes in the 15- to 20.8-cm and 20.8- to 31.5-cm spatial filters in the north beam did not produce noticeable improvement in the beam spatial

Spatial-filter size (cm)	No. 93012111 (both beams)		No. 93012409			
	Pinhole (mm)	λ_{co}^a (mm)	South beam		North beam	
			Pinhole (mm)	λ_{co}^a (mm)	Pinhole (mm)	λ_{co}^a (mm)
15.0 to 20.8	1.0 ^b	23.0	0.75 ^b	30.0	0.75 ^b	30.0
20.8 to 31.5	1.3 ^b	17.0	1.0 ^b	22.0	1.0 ^b	22.0
31.5 to 46.0	2.0 ^b	9.1	1.5 ^b	12.0	2.0 ^b	9.1
46.0 to 74.0	3.3 ^c	9.4	3.3 ^c	9.4	3.3 ^c	9.4
		10.5 ^d		12.0 ^d		10.5 ^d

^aSmallest spatial wavelength passed by spatial filter scaled to 74-cm output beam.

^bSide length for square pinhole.

^cSlit width for slit pinhole.

^dObserved near-field self-focusing cell.

modulation. Reducing the pinhole size in the 31.5- to 46-cm spatial filters in the south arm increased the wavelength of the dominant structure and reduced its modulation depth. The distinct, rectangular pattern that we had observed in earlier tests is still present, indicating that the beam remained on the threshold of catastrophic breakup. We now believe that beam breakup, at lower than expected power, is caused by higher than anticipated modulation imposed on the near-field beam by square spatial-filter pinholes. By returning to round pinholes, we should be able to sufficiently decrease the modulation to allow propagation of high-power beams through the system without danger of breakup.

The data collected during activation of Novette have confirmed that Nova will attain its high-power system performance milestones for 1ω and 2ω (Ref. 92). The system performance achieved in Novette with high-power short pulses has increased our confidence that high-energy, longer-pulse milestones will be attained as well. We are now looking forward to a year of productive target experiments.

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Table 2-14. Pinhole sizes used in largest spatial filters for shot numbers 93012111 and 93012409.

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A contour plot is shown of the plasma density in the corona of a laser-irradiated target. The ripples in density are due to filamentation and Brillouin sidescatter of the laser light, which is incident from the top of this picture.

Section 3

Target Design

Target Design

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Target Design

Introduction

J. H. Nuckolls

Our target-design program combines the efforts of the Plasma, Code Development, and Design groups in three 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.

During 1982, we designed targets for early Novette experiments, which are expected to achieve 200 times liquid density. Later Novette experiments may reach densities twice as high. We also designed targets for use in Novette coupling and weapons-physics experiments.

Significant effort was devoted to the design of targets for Nova. The principal objectives here are to test coupling at plasma scale lengths corresponding to megajoule-size targets, to achieve 1000 times liquid density, and to implode targets that are hydrodynamically equivalent to the high-gain, reactor-scale targets anticipated for the future. These experiments will define the output of the driver required to ignite high-gain targets and will provide confidence that high gains can be achieved. A limited effort was devoted to the physics of targets and beam transport for a reactor employing a heavy-ion beam as the driver.

LASNEX, our principal computer code for target design and experiment analysis, was substantially improved—particularly in areas of energy and charged-particle transport. We added more accurate physical models of the degree of ionization of compressed matter and of mixed materials. LASNEX has played a key role in the design of the laser-driven x-ray lasers to be tested on Novette in 1983.

Our Plasma Physics group gained a better understanding of the problems to be encountered in designing targets for Novette, Nova, and future reactors, as well as of the complex plasma processes that have already been observed. Improved plasma codes have been developed, which operate on time and space scales not accessible to LASNEX. Approximate descriptions of Raman and other plasma instabilities derived from parameter studies with these codes have been formulated and incorporated into LASNEX.

Additional effort is being devoted to the development of targets with higher performance. A preliminary analysis of implosion and burn physics suggests that the use of polarized fuel may be practical in inertial-confinement fusion (ICF) targets. If our analysis is correct, the required driver energy may be reduced by one-half or two-thirds.

Our estimates of fusion-reactor driver requirements and our calculated target gain curves are unchanged from previous years (see Fig. 3-1 in the 1981 Annual Report). We expect to achieve the "conservative" estimates of target gain in early experiments, and later to progress along a learning curve toward the "optimistic" performance levels with advanced target concepts. The experimental tests of our target designs, and our analyses of the results, are reported in Section 5, "Fusion Experiments Program."

Laser Targets

Introduction

Most of the effort of the laser-target design group has been in preparing for the experiments soon to begin on Novette and in developing advanced target designs for Nova, which will be hydrodynamically equivalent to the eventual high-gain, megajoule-scale capsules projected for future reactors. To experimentally demonstrate the hydrodynamic equivalence will require numerous advances in diagnostics and fabrication techniques. We have concentrated on designs that will allow us to use Novette as a test bed for developing several techniques. If these techniques are successfully implemented, Novette should produce implosions that reach densities from 2 to 4 $\times$ those obtained with Shiva.

We have extended our previous analytic work in examining the trade-offs involved in optimizing implosion efficiency, symmetry, and stability for direct-drive ICF targets. As suggested by S. Bodner of the Naval Research Laboratory, we find that, for a given degree of smoothing and constant implosion velocity, the hydrodynamic efficiency increases approximately as $\lambda^{0.25}$. However, we find that the pressure generated for a given degree of smoothing varies as $\lambda^{-1.7}$. Hence, the number of e -foldings for Rayleigh-Taylor instability would increase as $P^{1/2} \propto \lambda^{0.85}$ for a given implosion velocity.

If, for fixed wavelength, the implosion pressure must be increased, the implosion efficiency is approximately inversely proportional to pressure for $r_c/r_a > 1.2$, where r_c is the critical radius and r_a is the ablation radius.

In the past year, KMS Fusion, Inc. has used a short-duration probe pulse to map out experimentally the subcritical density profile for spherical targets. Our LASNEX simulations of the KMS profiles for low-Z materials are consistent with electron transport modeled with a flux limit of $f = 0.03$. Our analyses of the profiles for layered spheres coated with low-Z material are also consistent with flux-limited transport. In general, for the experiments to date, there appears to be no significant dif-

ference in transport between spherical and slab targets.

It is possible to reduce growth of the Rayleigh-Taylor instability at a density discontinuity by introducing a density gradient to bridge the discontinuity. This problem was first tackled by Lelevier in 1952. We have derived the shape of the optimal density profile for such a bridging layer, for the case of fixed mass or fixed distance between the high- and low-density layers.

It is also possible to approximate a continuous density profile with a series of density steps. We have worked out a formalism to treat the growth of the Rayleigh-Taylor instability for an arbitrary stepped profile. The results are valid for the linear regime, and, hence, would be valid for implosions in which all modes remain in the linear regime. However, the approach does not take into account important issues of mode coupling, other nonlinear effects, or the effects of the ablation process.

Author: J. D. Lindl

Trade-Offs Among Symmetry, Stability, and Implosion Efficiency for Directly Driven ICF Targets

One of the important considerations in symmetry is the standoff, i.e., the distance between the critical surface and the ablation surface ($r_c - r_a$), which usually appears in normalized form, $(r_c - r_a)/r_a = r_c/r_a - 1$, or simply as r_c/r_a . Last year, we reported that when both r_c/r_a and the ablation fraction m_i/m_0 (ratio of payload mass to initial mass), are fixed (see Fig. 3-1), then the implosion efficiency is essentially independent of laser wavelength. However, typical capsules require a fixed implosion velocity rather than a fixed ablation fraction. For fixed implosion velocity and fixed r_c/r_a , the implosion efficiency actually increases slowly with increasing wavelength.

The ablation fraction for optimal implosion efficiency is $m_i/m_0 \approx 0.2$. For a typical implosion velocity of 3×10^7 cm/s, a direct-drive rocket (which is the applicable concept here) typically ablates substantially less than the optimum amount, so that the rocket efficiency is less than optimal, going

in. Moreover, in order to keep r_c/r_a fixed as the wavelength is shortened, the temperature and hence the exhaust velocity of the rocket must increase, thus further decreasing the ablated mass and moving m_i/m_0 further from the optimum, resulting in still lower rocket efficiency.

Another consideration is that for fixed implosion velocity and r_c/r_a , the ablation pressure, P_a , decreases rapidly as wavelength is increased. In order to reach a fixed implosion velocity, we require that $P_a r_a / \Delta r \approx \text{constant}$, where Δr is the shell thickness. Since the Rayleigh-Taylor instability sets a maximum tolerable value of $r_a / \Delta r$, there is a minimum allowable pressure and hence a maximum useful wavelength.

Fig. 3-2 gives a global view of these results, based on calculations with a fixed velocity of 3×10^7 cm/s, a flux limit $f = 0.03$, and a capsule size appropriate to a high-gain driver (a few megajoules), giving $r_a = 0.3$ cm with a driving pressure of 20 Mbar. For pressures other than 20 Mbar, the radius was scaled as $P_a^{-1/3}$.

One set of curves in Fig. 3-2 plots rocket efficiency, ϵ , vs wavelength for fixed ablation pressure, showing $\epsilon \propto \lambda_L^{-1.7}$. The other set of curves plots efficiency vs wavelength

for fixed standoff, showing $\epsilon \propto \lambda_L^{0.25}$. As explained, fluid instability will set a minimum ablation pressure, while symmetry requirements will set a minimum value of r_c/r_a . If the required ablation pressure is known, then following the curve at that pressure will reveal the shortest wavelength that will yield the required smoothing. Any longer wavelength will also give the required smoothing, but the efficiency will drop as $\lambda_L^{-1.7}$.

As shown in Fig. 3-2, a relatively modest increase in standoff results in a substantial drop in efficiency. Hence there is a significant advantage to providing the smoothest possible beams and operating at the shortest possible wavelength. In addition, by choosing a wavelength to generate a

Fig. 3-1. Implosion efficiency vs normalized standoff distance is virtually independent of laser wavelength λ_L and only weakly dependent on f . (Ablation fraction fixed.)

Fig. 3-2. Rocket efficiency vs laser wavelength for six fixed ablation pressures and for five fixed normalized standoff distances. (Implosion velocity constant.)

Laser Targets

given standoff and pressure, an increase in the required pressure by a factor of 2 produces a corresponding drop in efficiency of almost a factor of 2, since $\epsilon \propto P^{-0.75}$.

Authors: J. D. Lindl, C. E. Max, and W. C. Mead

Analysis of KMS Electron-Transport and Density-Profile Data

In recent years, many experiments have pointed to the need for invoking anomalous (inhibited) thermal-electron transport in target simulations; when the thermal-electron mean free path becomes comparable to the scale length characterizing the temperature

gradient, the physical assumptions made in deriving the well known Spitzer conductivity break down. In this regime, we therefore invoke the constraint that the electron flux be no greater than the free-streaming limit, nvT . To match experimental data, however, the need has arisen to constrain the flux to $f\nu T$, where f (known as the flux limit) is a small number, as opposed to $f \approx 0.5$ in the classical, free-streaming case.

In this report we analyze new data, taken at KMS Fusion, Inc.¹ The new data, sensitive to the flux limit, are a series of snapshots of the electron density profile, taken holographically with a separate, frequency-quadrupled (4ω) probing laser. The density profiles decisively reinforce the need to invoke a more severe flux limit, at or around 0.03.

If electron transport is not inhibited, the heat is efficiently transported from the critical target surface, where the laser is absorbed, to the overdense plasma; the heated plasma then blows out into the underdense corona. If electron transport is inhibited, less overdense plasma is heated, and less is blown out into the corona. Thus, at a given time, at a given position not too far from the original target surface (200 μm in Fig. 3-3), there is more plasma for the noninhibited model than for the inhibited

Fig. 3-3. The qualitative behavior of the electron density profile changes depending on the value of the electron flux limit.

Fig. 3-4. Comparison of measured and simulated density profiles from a KMS gold-disk shot; the flux-limited models match the data more closely than does the free-streaming model.

model. On the other hand, the inhibited model bottles up all the energy in the underdense plasma and heats it to much higher temperatures. This very hot, tenuous plasma consequently moves out faster than the cooler, denser, and slower-moving plasma of the noninhibited model. Therefore, at a point very far from the target (350 μm in Fig. 3-3) there is actually more plasma according to the inhibited model than according to the noninhibited model.

Our analysis qualitatively explains Fig. 3-3, which shows quite graphically the difference in predictions for the underdense density gradient between our two LASNEX flux-limit models: $f = 0.5$, which inhibits electron transport to half the naive, classical free-streaming limit of mvT ; and $f = 0.03$, which restricts the electron flux to just a few percent of free streaming.

Gold-Disk Data. The first set of KMS data we examine was produced with 1.06- μm irradiation of gold disks. Previous work at LLNL showed the need to invoke a flux limit of 0.03 for gold-disk irradiations.² By setting $f = 0.03$, we were able to match such observable quantities as the fraction of laser light absorbed, the fraction converted to re-radiated x rays, and the fraction scattered by stimulated Brillouin scatter (SBS); we were also able to match estimates of the coronal electron temperature. Nonetheless, the validity of the classical flux limit remained open. This possibility required changing our assumptions about SBS (to match measured absorption and SBS) and our assumptions about material opacities (to match measured x-ray emissions). Thus, the density-profile data from KMS imposed an added constraint on the modeling with which to test energy transport.

The laser energy in the KMS experiments ranged from 100 to 200 J in a 750-ps flat-topped pulse; the incident laser intensity was $2 \pm 1 \times 10^{14} \text{ W/cm}^2$. The probing pulse at $t = 200$ ps was 80 ps long. We performed LASNEX simulations with these parameters, varying the flux limit. Since changing f changes the coronal conditions, and therefore the observable quantities mentioned above, we varied the amount of laser light legislated as lost due to SBS. We also varied the opacities such that, for each value of f , our simulation matched the measured absorption and x-ray emission. The

simulations were done in two dimensions with laser ray tracing.

In Fig. 3-4, we compare the LASNEX simulations with the KMS gold-disk data; the three curves for each model correspond to the profiles at the beginning, middle, and end of the probing time. Note that the profile data are matched either by a severe flux limit ($f = 0.01$) or by a slightly less severe limit ($f = 0.03$) when LASNEX is allowed to generate thermal-electric magnetic fields, which tend to contribute to flux inhibition. Note also that the profile corresponding to $f = 0.5$ lies well outside the error bars of the KMS data, showing clearly the need to invoke nonclassical thermal-electron transport.

As we noted earlier, the very underdense plasma moves out further in the inhibited case, since the plasma is hotter. In Fig. 3-5, we compare the KMS data with the predictions of our two flux-limit models for the two-dimensional position of the first holographic interference fringe. The fringe blowing out from the target at 200 ps corresponds to a rather low density of approximately $2 \times 10^{19} \text{ electrons/cm}^3$. Note that a hot, rapidly moving corona—as is created in a flux-inhibited model—is required to match the data.

Spherical-Target Data. The second series of experiments we analyze consists of

Fig. 3-5. Two-dimensional snapshot of the first holographic interference fringe from a gold-disk target; the inhibited transport model (with a hot corona) shows the best fit to the data.

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spherical targets, illuminated spherically by the triple-bounce illumination scheme at KMS. The spherical nature of these experiments greatly assists our computer simulations because we can use the much simpler one-dimensional modeling. The targets were roughly $100\text{ }\mu\text{m}$ in diameter; glass shells $6\text{ }\mu\text{m}$ thick were coated with (working outward) $0.5\text{ }\mu\text{m}$ of aluminum, $5\text{ }\mu\text{m}$ of plastic (CH), and finally $0.01\text{ }\mu\text{m}$ of titanium. The titanium flashes on briefly with x-ray emission at the onset of irradiance. The thermal-conduction front then propagates through the poorly emitting plastic, eventually reaches the aluminum, heats it, and lights up the target in x-ray emission once again. Time-resolved x-ray detectors recorded these emissions vs time, providing a clear record of the energy transport

through the CH. Density profiles were also taken on these shots.

Figure 3-6 compares actual density profiles from four shots with density-profile predictions for those shots. The laser intensities for these shots were 2 , 6 , and $9 \times 10^{14}\text{ W/cm}^2$. Once again the LASNEX model with $f = 0.03$ matches the data very well, whereas the $f = 0.5$ model does not.

Another way of viewing the conclusions from modeling this series is given in Fig. 3-7, in which we plot the position of the critical surface vs time. The data points come from the density-profile measurements; the number within each point represents the incident laser intensity in units of 10^{14} W/cm^2 . The red lines are experimental inferences of the critical radius trajectory based on measured Doppler shifts. The

Fig. 3-6. Density profiles from four different spherical, layered-transport targets; all require $f = 0.03$ to match the data.

shaded regions are LASNEX predictions for the two models, with the thickness of the regions representing a twofold increase in incident energy. Thus, the overwhelming sensitivity is not in the absorbed energy fraction but rather in the choice of flux limiters; once again, $f = 0.03$ shows the best agreement.

At this point, it is proper to mention those factors to which the computer simulations are most sensitive. In our spherical experiments, the amount of absorbed energy that wound up as suprathermal electrons was roughly 35% of the absorbed energy. To match the data, we need to inhibit the flux of the suprathermals as well as the thermals; if we inhibit one and not the other, the results resemble the $f = 0.5$ model. The profiles are rather insensitive to changes in the pulse shape, which nominally is a fast-rising (100 ps), flat-topped (1 ns) pulse; in reality, the pulse shape deviates by as much as a factor of 2 in rise time and flatness. The SBS model of Estabrook et al.¹ predicts no SBS for these experiments. As such, the $f = 0.5$ model predicts too much absorption by a factor of 2, even when the classical inverse bremsstrahlung

absorption opacity is multiplied by 0.5 to account for the Langdon effect. Thus, while we must throw away energy to make the $f = 0.5$ model match the absorptions, we really have no good justification for doing so. Finally, the implosion time depends sensitively on the flux limiter and, as such, could serve as yet another constraint on the modeling.

Since LASNEX models x-ray emission from heated materials, we can match the time history of x-ray emission from these layered spheres with the experimental observations. We present the data in terms of the mass ablation rate, dm/dt , operationally defined here as the thickness of CH multiplied by its density (1.1 g/cm^3), divided by the time interval between the titanium x-ray flash and the x-ray turn-on due to the heat front reaching the aluminum. We plot the mass ablation rate of CH in Fig. 3-8 vs the absorbed irradiance. Note that most of the data lie near the LASNEX $f = 0.03$ curve, and well below the $f = 0.5$ curve. There is possibly a single exception at the very lowest absorbed irradiance, corresponding to an incident irradiance of just under $1 \times 10^{14} \text{ W/cm}^2$.

Fig. 3-7. Time history of the critical surface radius. Data points are derived from holographic density-profile measurements; red lines are critical surface trajectories inferred from Doppler-shift measurements.

In Fig. 3-9, we plot data from other labs along with the KMS data and our LASNEX models. An intriguing picture seems to emerge. Below some threshold, corresponding to an incident intensity of 1×10^{14} W/cm², the data lie on the $f = 0.5$ curve implying no need for anomalous flux inhibition for those conditions. Above that threshold, however, the need for $f = 0.03$ emerges quite clearly. This threshold has been discussed earlier by McClellan et al., based on the analysis of temporal x-ray emission from disks of varying Z (Ref. 4). They concluded that the inhibition threshold is Z-dependent, and its value for CH is about 1×10^{14} W/cm². This value was also deduced some years earlier by Pearlman and Anthes⁵; it is also hinted at by our compilation of data in Fig. 3-9.

In summary, by analyzing the KMS density profiles both for high-Z disks and low-

Z coated spheres, we have seen the clear need to invoke a flux limit at or near 0.03. These new data confirm the modeling we have been using in the past few years to analyze the behavior of laser-driven targets.

Author: M. D. Rosen

Antimix-Layer Density Profiles Giving Least Rayleigh-Taylor Instability

The Rayleigh-Taylor instability causes small ripples on an interface between two fluids to grow if the interface is accelerating toward the denser fluid. The growth rate of such ripples can be reduced by sandwiching a layer of intermediate density between the two fluids. If this "antimix layer" is reasonably thick (thick enough to significantly reduce the growth rate), further reduction in the growth rate can be achieved by smoothly increasing the density of the layer from its interface with the low-density fluid to its interface with the high-density fluid. We can use variational calculus to find what the density profile for such an antimix layer must be to minimize the growth rate of the Rayleigh-Taylor instability.

Useful terminology for this problem is defined in Fig. 3-10. A gravitational field of acceleration g acts in a downward direction on a horizontally stratified fluid. The low-density fluid is gravitationally lowest in the stack, followed by the antimix layer, with the high-density fluid at the top of the stack. The lowest layer is called the supporting fluid, while the highest layer is called the supported fluid. The vertical density profiles in the supporting and the supported fluids, as well as the boundary conditions at the top and bottom of the stack, are assumed to be given.

The goal is to determine the vertical density profile in the antimix layer that minimizes the growth rate of the Rayleigh-Taylor instability. Some additional constraint must be placed on the profiles to be examined, otherwise an infinitely thick antimix layer with an infinitesimal density gradient will surely be the solution. Two constraints have been considered: either the thickness of the antiripple layer can be constrained to have a given value, or its mass per unit area can be constrained to have some given value.

Fig. 3-8. Mass ablation rate of CH vs absorbed intensity. Data points are derived from x-ray emission vs time, as the CH burns and allows the underlying aluminum to radiate.

Fig. 3-9. Compilation of data on mass ablation rate from various laboratories. Note the possible shift of data, from the $f = 0.5$ curve at low intensities to the $f = 0.03$ curve at higher intensities.

The problem is best solved by first constraining both the thickness and the areal mass of the antimix layer. Later, these solutions can be searched at constant areal mass or at constant thickness to determine which profile is best when only one of these quantities is constrained. The solution will be obtained by using variational techniques on an integral expression for the growth rate for an antimix layer of a given thickness and employing a Lagrange multiplier to constrain the antimix-layer mass.

The amplitude of the vertical (z) component of the displacement field is ξ . Since the displacement field is wavelike in the transverse direction, ξ is a function of z alone. The transverse wave number, k , is assumed to be given. The most convenient dimensionless measure of the growth rate γ of the instability is the stability parameter

$$\sigma = gk/\gamma^2 \quad (1)$$

which increases with decreasing growth rate. If the boundary conditions at surfaces 0 and 3 of Fig. 3-10 are those of a free surface or a rigid surface, then the following integral expression for the stability

parameter holds

$$\sigma = - \frac{\int_0^3 dz \rho (\xi'^2 + \xi^2)}{2 \int_0^3 dz \rho \xi'^2 \xi} \quad (2)$$

where primes denote $d/d(kz)$. Finally, the areal mass of the antimix layer is simply

$$m = \int_1^2 dz \rho \quad (3)$$

The task is now to extremize the right-hand side of Eq. (2), subject to the constraint in Eq. (3), for all admissible variations in the mode shape $\xi(z)$ and the density profile $\rho(z)$. Extremizing σ with respect to variations in ξ merely ensures that the mode shape $\xi(z)$ does, in fact, satisfy the equations of motion for the (unknown) density profile $\rho(z)$. Maximizing σ with respect to variations in ρ is, of course, the goal of the exercise. The two simultaneous differential equations that result from this procedure can be solved analytically for the mode shape and the optimal density profile.

The results of this calculation are

Fig. 3-10. Definition of terms.

$$\begin{aligned}
 & -k(z - z_1) \\
 & = \sigma \ln \left| \frac{-\sigma(\xi/\xi_1) \pm \sqrt{(\sigma^2 - 1)(\xi/\xi_1)^2 - (\ell/\xi_1^2)}}{-\sigma \pm \sqrt{\sigma^2 - 1 - (\ell/\xi_1^2)}} \right| \\
 & + \sqrt{\sigma^2 - 1} \\
 & \times \ln \left| \frac{\sqrt{\sigma^2 - 1}(\xi/\xi_1) \pm \sqrt{(\sigma^2 - 1)(\xi/\xi_1)^2 - (\ell/\xi_1^2)}}{\sqrt{\sigma^2 - 1} \pm \sqrt{\sigma^2 - 1 - (\ell/\xi_1^2)}} \right|
 \end{aligned} \quad (4)$$

and

$$\begin{aligned}
 \frac{\rho_1}{\rho} & = \left[\frac{-\sigma(\xi/\xi_1) \pm \sqrt{(\sigma^2 - 1)(\xi/\xi_1)^2 - (\ell/\xi_1^2)}}{-\sigma \pm \sqrt{\sigma^2 - 1 - (\ell/\xi_1^2)}} \right] \\
 & \times \frac{\sqrt{(\sigma^2 - 1)(\xi/\xi_1)^2 - (\ell/\xi_1^2)}}{\sqrt{\sigma^2 - 1 - (\ell/\xi_1^2)}}
 \end{aligned} \quad (5)$$

Fig. 3-11. Curves of constant growth rate for given transverse wave number k and supporting and supported fluid structure. The growth rate plotted is the minimum possible for antimix layers with the given values of km and $k\theta$.

The constant ℓ is related to the Lagrange multiplier associated with the constraint of given mass. The constant ξ_1 is arbitrary; it simply sets the amplitude of the perturbation. The constant ρ_1 and the sign of the square roots can be determined from the boundary condition at the lower boundary of the antimix layer. Equation (5) then de-

termines a value of ρ_2 , which must be consistent with the boundary condition at the upper boundary of the antimix layer.

From Eq. (4), the condition of consistency determines a unique value of the thickness of the antimix layer for which this solution can apply. Thus, although the thickness and areal mass of the antimix layer were originally regarded as given, the problem is more naturally expressed by the constants σ and ℓ/ξ_1^2 in Eqs. (4) and (5). These equations, together with the consistency of the boundary conditions, can be used to derive analytic expressions for the antimix-layer mass, m , and thickness, θ , associated with any particular values of the parameters σ and ℓ/ξ_1^2 .

Using the preceding expressions, we can construct a contour plot of the maximum stability parameter, or minimum growth rate, attainable by introducing an antimix layer of thickness θ and areal mass m . Such a plot is shown in Fig. 3-11; its details depend on the thicknesses of the supporting and supported fluids, the boundary conditions at their outer surfaces, the transverse wave number, and the density contrast ρ_1/ρ_0 . However, the qualitative features of the plot remain invariant.

The bold contour in Fig. 3-11 represents the growth rate in the absence of any antimix layer. Outside its two branches, the growth rate of the Rayleigh-Taylor instability with an optimal antimix layer is actually greater than it would be without any antimix layer at all. In the upper left part of the plot, the combination of low areal mass and large thickness forces a low density gap between supporting and supported fluids, which invariably causes a high growth rate. Conversely, in the lower right part of the plot, high areal mass and low thickness imply that the highest densities in the problem will actually be found within the antimix layer, again leading to high growth rates.

Indeed, to the right of the dashed line in Fig. 3-11, the method cannot find any optimal density profile. The reason can be seen by studying the sequence of optimal density profiles corresponding to the points on the horizontal line PQ . This sequence is shown in Fig. 3-12. At the dashed line in Fig. 3-11, the density profile becomes singular at the upper boundary of the antimix layer. To the

right of the dashed line, the optimal density profile would include a delta function singularity at the upper boundary of the antimix layer.

Thus, the only region of interest in Fig. 3-11 is the region between the bold curves, which represents a ridge of stability. In particular, the locus of points P (where P on a given curve is the point corresponding to minimum thickness) represents the density profiles with minimum growth rate for a given thickness, independent of mass. The locus of points R (where R on a given curve is the point corresponding to minimum mass) represents the optimal profiles for a given mass, independent of thickness. By setting appropriate derivatives of the analytic expressions for m and θ to zero, these points can be algebraically determined. The resulting equations can be solved in closed form in both cases.

For either the case where the thickness of the antimix layer, or the amount of areal mass, is constrained to take a given value, the optimal density profile is exponential, with density jumps at the upper and lower boundaries of the layer. The magnitudes of the density jumps depend on the given transverse wave number so that the instability growth at the jumps is comparable to the growth in the density gradient within the antimix layer. The optimal profile for given mass differs significantly from the optimal profile for given thickness only when the thickness of the antimix layer is less than about $1/k$, where k is the transverse wave number.

Since the optimal profile for a given areal mass is of greater practical interest, the formulas for it will be explicitly given. The mode shape and density profile in the antimix layer are

$$\zeta(z) = \zeta_1 = \zeta_2, \quad (6)$$

and

$$\left. \begin{aligned} \frac{\rho(z)}{\rho_1} &= \exp[k(z - z_1)/\sigma], \\ \frac{\rho_1}{\rho_a} &= \frac{A_a(\sigma)}{\sigma}, \quad \text{and} \\ \frac{\rho_2}{\rho_b} &= \frac{|A_b(\sigma)|}{\sigma}. \end{aligned} \right\} \quad (7)$$

The thickness and areal mass of the antimix layer are

$$k\theta = \sigma \ln[\rho_b A_b(\sigma)/\rho_a A_a(\sigma)] \quad (8)$$

and

$$km = \rho_b A_b(\sigma) - \rho_a A_a(\sigma). \quad (9)$$

Here, the boundary functions A_a and A_b are defined (see Fig. 3-10) by

$$A_a(\sigma) = \begin{cases} \sigma + \coth k\theta_a, & \text{if surface 0 is rigid,} \\ \frac{\sigma^2 - 1}{\sigma - \coth k\theta_a}, & \text{if surface 0 is free,} \end{cases} \quad (10)$$

and

$$A_b(\sigma) = \begin{cases} \sigma - \coth k\theta_b, & \text{if surface 3 is rigid,} \\ \frac{\sigma^2 - 1}{\sigma + \coth k\theta_b}, & \text{if surface 3 is free.} \end{cases} \quad (11)$$

The explicit dependence of optimal density profile on the transverse wave

Fig. 3-12. A series of optimal antimix-layer density profiles at fixed thickness with increasing mass.

number k proves that no single density profile can minimize the Rayleigh-Taylor growth rate for perturbations of all transverse wavelengths.

It is remarkable that the density profiles, which are optimum for a given mass, have the extremely simple mode shape represented by Eq. (6). Since the amplitude of the vertical displacement field is constant throughout the antimix layer, and since the (infinitesimal) displacement field is divergence-free (isochoric), there must be no transverse displacement of the fluid in this mode. Thus, the density profile that minimizes the instability growth rate for a given areal mass has the property that its fastest growing mode exhibits no transverse displacements within the antimix layer. Hence, the usual "bubbles and spikes" do not form.

Similarly, the density profile of given thickness, which minimizes the instability growth rate, is characterized by the fact that, for its fastest growing mode, the kinetic energy density is spatially constant within the antimix layer.

In most interesting situations, the supporting and supported fluids are of the order of one transverse wavelength, or more, in thickness. Then, the hyperbolic cotangents in the boundary functions in Eqs. (10) and (11) are essentially unity so that the expression for the areal mass in Eq. (9) takes on the especially simple form

$$km = \rho_b \left(\frac{gk}{\gamma^2} - 1 \right) - \rho_a \left(\frac{gk}{\gamma^2} + 1 \right). \quad (12)$$

This formula is a useful rule of thumb. If an antimix layer of areal mass no greater than m is to be placed between layers of density ρ_a and ρ_b , then the smallest growth rate γ for perturbations of wave number k is given by

$$\frac{gk}{\gamma^2} = \frac{\rho_b + \rho_a}{\rho_b - \rho_a} + k \frac{m}{\rho_b - \rho_a}. \quad (13)$$

The first term on the right-hand side of Eq. (13) represents the usual formula for the growth rate of the Rayleigh-Taylor instability in the absence of an antimix layer. The second term is simply kL , where L is the density-gradient scale length in the antimix layer. Again, the second term by itself is the usual estimate for the growth rate of the

Rayleigh-Taylor instability in a continuous density gradient. In a sense, then, while Eq. (13) has always been known, we have provided the specific question to which Eq. (13) is the answer.

Author: D. H. Munro

Graded-Density Antimix Layers for Suppression of Rayleigh-Taylor Instability

Fluid instabilities in ICF targets arise sometime during the history of the implosion, when a low-density material pushes and accelerates a shell of higher-density material. Typically, the low-density material is plastic or staging fluid of density about 1 g/cm^3 , while the heavy shell is a high- Z inner pusher, such as gold, of density about 20 g/cm^3 . Rayleigh-Taylor instabilities are caused by small density perturbations, at the interface between two fluids, which can grow and cause the heavy shell to break up and mix with the DT fuel, resulting in substantial degradation in yield.

The instabilities can be suppressed by making the transition smoother from the low- to the high-density regions. Smoothing can be achieved by introducing antimix layers at the interface. Our systematic study has yielded definite results as to the densities and thicknesses of antimix layers required to minimize the growth of Rayleigh-Taylor instabilities.

The hydrodynamic efficiency of implosion will change if, in introducing the antimix layers, we increase the mass to be accelerated. To avoid a change in overall mass, we can reduce the mass of the heavy shell, in which case our approach can be properly described as spreading the mass of the original heavy shell into some number n of subshells, keeping the total areal mass fixed (Fig. 3-13).

Our calculations are based on the linear theory of Rayleigh-Taylor instabilities.⁶ The linear theory is adequate because we are trying to suppress the initial growth, where the perturbations are small, and yet grow most rapidly (in fact, exponentially) with time. As the amplitude grows larger, nonlinear effects slow down the growth. We will make use of this fact later as a criterion for choosing the wavelength of the most dangerous perturbations.

Over the past year, we have developed⁷ new techniques for calculating the eigenvalues and eigenvectors for normal modes in a stratified fluid; we have applied our techniques to find how the Rayleigh-Taylor instability can be suppressed by graded-density layers.

In the original heavy shell, perturbations grow almost classically⁶ such that

$$\gamma_{\text{classical}} = [gk(\rho_h - \rho_l)/(\rho_h + \rho_l)]^{1/2} \quad (14)$$

where γ is the rate of exponential growth (perturbations $\propto e^{\gamma\tau}$, where τ is time), g is the acceleration due to gravity, $k = 2\pi/\lambda$, $\lambda =$ wavelength of the perturbation, and ρ_h (ρ_l) is the density of the heavy (light) fluid. In addition to Eq. (14), there is a stable mode, which instead of growing, oscillates in time with $\gamma = i\gamma_{\text{osc}}$.

Equation (14) leads to unacceptably large growth rates. When a certain fraction (to be determined) of the mass is spread into n subshells, we find $n + 1$ unstable modes in addition to the (harmless) stable mode. All the unstable modes are smaller than $\gamma_{\text{classical}}$, and we have developed codes to calculate and list them in decreasing order. We focus on the first (fastest growing) mode, which we call γ_{max} , and determine the fraction of mass spread to minimize this fastest-growing mode.

Normal Modes. Perturbations growing in the direction of acceleration $\mathbf{g} = g\hat{\mathbf{y}}$ can be analyzed in terms of normal modes with velocity $v_y = W(y)e^{ikx}e^{\gamma t}$. The growth rate γ is found by solving the differential equation

$$\frac{d}{dy} \left(\rho \frac{dW}{dy} \right) + \frac{gk^2}{\gamma^2} W \frac{d\rho}{dy} - k^2 \rho W = 0, \quad (15)$$

subject to proper boundary conditions. Here $\rho = \rho(y)$ is a specified density profile, and Eq. (15) must be viewed as an eigenvalue equation, where $W(y)$ is an eigenfunction corresponding to the eigenvalue γ , both of which remain to be found.

Stepwise Density Profile. Figure 3-14 shows the density profile for n antmix layers (or subshells) of arbitrary thicknesses t_i and densities ρ_i , $i = 1, 2, \dots, n-1$. The thickness of the low-density $\rho_0 = \rho_l$ shell is assumed to be large, and δt denotes the thickness of the remaining heavy shell whose density is ρ_h .

There are $n + 2$ interfaces, and we find⁷ that there are $n + 2$ modes whose growth rates are given by the eigenvalues of the matrix $\mathbf{M}$. [See Eq. (16) in Table 3-1.] In Eq. (16), $S_i = \sinh(kt_i)$, $i = 1, \dots, n + 1$ ($t_{n+1} = \delta t$), and the diagonal elements C_i are given by

$$C_i = \frac{1}{\rho_{i+1} - \rho_i} \times \left[\rho_{i+1} \left(T_{i+1} + \frac{1}{S_{i+1}} \right) + \rho_i \left(T_i + \frac{1}{S_i} \right) \right], \quad (17)$$

where $T_i = \tanh(kt_i/2)$, and $i = 1, \dots, n + 1$.

The matrix $\mathbf{M}$ depends only on the shell densities and their thicknesses and, of course, on $\lambda = 2\pi/k$. The eigenvalues of $\mathbf{M}$ give gk/γ^2 . The characteristic equation is

$$\det \left| \mathbf{M} - \frac{gk}{\gamma^2} \mathbf{I} \right| = 0, \quad (18)$$

which is a polynomial equation of degree $n + 2$, giving $n + 2$ values of γ^2 .

We developed codes⁷ to calculate all the eigenvalues and eigenvectors for a stratified density profile. In addition, we checked how these modes approach the modes of a

Fig. 3-13. (a) Capsule with a single-shell pusher of uniform density. (b) Antmix layers introduced at the pusher/ablator interface by spreading some of the mass of the original pusher into a number of subshells.

continuous exponential profile $\rho = \rho_0 e^{x/d}$ in the limit of large n , with various boundary conditions. We solved Eq. (15) exactly for the case of an exponential profile and compared our results for n discrete layers with the results of the continuous profile, which we denote by $n = \infty$.

Exponential Density Profile. The density profile shown in Fig. 3-15 has an infinite number of normal modes, only one of which is "trivial" and stable: $\gamma = \pm i\pi gk$. The remaining growth rates are given by

$$\frac{\gamma^2}{gk} = \frac{2cd}{x^2 + d^2 - c^2} \quad (19)$$

where $c = kt$, $d = \beta t/2$, t is the thickness of the antirix layer, and $\beta t = \ln(\rho_h/\rho_d) \simeq 3$. In Eq. (19), x is obtained by solving the transcendental equation

$$\tan x =$$

$$\frac{2cx}{x^2 + d^2 - c^2 - [x^2 + (c - d)^2] e^{-2\delta t}} \quad (20)$$

where, as before, δt is the thickness of the remaining heavy shell. The solutions are again listed in decreasing order. We focus on the fastest growing rate $\gamma_{\max}$ and deter-

mine the fraction of mass spread that minimizes $\gamma_{\max}$.

In our formalism, the thicknesses and wavelengths are arbitrary. However, physical considerations, to be discussed next, help us choose the relevant values for the two parameters.

Mass Constraint and Wavelength Selection. When no mass is spread, so that the heavy shell has its original thickness t_0 , the unstable mode is close to $\gamma_{\text{classical}}$. If f = the fraction of mass spread, then $\delta t = (1 - f)t_0$. We will often set $t_0 \equiv 1$ for scale. Requiring that the total areal mass remain fixed, i.e.,

$$m_{\text{antimax}} + m_{\text{remaining}} = m_{\text{original}} = \rho_h t_0$$

we obtain

$$(\rho_h - \rho_c)/\beta + \rho_h(1 - f)t_0 = \rho_h t_0 \quad (21)$$

for the exponential profile $\rho = \rho_0 e^{x/d}$. Equation (21) determines β for given ρ_c/ρ_h and f , and we find t from $\beta t = \ln(\rho_h/\rho_d)$.

For the stratified profile of finite n , the mass requirement is

$$\sum_{i=1}^n \rho_i t_i + \rho_h(1 - f)t_0 = \rho_h t_0 \quad (22)$$

We set t_i equal to a common value t_{subshell} and choose the densities ρ_i such that the Atwood number $(\rho_{i+1} - \rho_i)/(\rho_{i+1} + \rho_i)$ at each interface is the same, i.e., independent of i . It is easy to see that in the limit $n \rightarrow \infty$ such a profile goes over to an exponential. For given ρ_c , ρ_h , and n , requiring an equal Atwood number implies that $\rho_i \approx \rho_c(\rho_h/\rho_c)^{i/(n+1)}$, $i = 1, \dots, n$, which when substituted into Eq. (22) determines t_{subshell} so that

$$t_{\text{subshell}}/t_0 = \frac{f}{\sum_{i=1}^n (\rho_i/\rho_h)^{n+1-i}} \quad (23)$$

The last and perhaps most important parameter needed to determine the growth of Rayleigh-Taylor instabilities is the wavelength λ . A reasonable choice is $\lambda = \delta t$, motivated by the nonlinear behavior referred

Fig. 3-14. Density profile of the capsule with antirix layers introduced: ablator density $\rho_0 \approx 1 \text{ g/cm}^3$, remaining pusher density $\rho_{n+1} \approx 20 \text{ g/cm}^3$.

Table 3-1. The matrix M.

$$M = \begin{pmatrix} C_0 & \frac{-\rho_1}{S_1(\rho_1 - \rho_0)} & 0 \\ \frac{-\rho_1}{S_1(\rho_2 - \rho_1)} & C_1 & \frac{-\rho_2}{S_2(\rho_2 - \rho_1)} \\ & \frac{-\rho_2}{S_2(\rho_3 - \rho_2)} & \ddots \\ & & \ddots & C_n & \frac{-\rho_{n+1}}{S_{n+1}(\rho_{n+1} - \rho_n)} \\ 0 & & & \frac{1}{S_{n+1}} & C_{n+1} \end{pmatrix} \quad (16)$$

to earlier (rapid exponential growth continues until the amplitude of the perturbation is of the order of the wavelength, after which the growth slows down). To maintain the integrity of the last (Au) shell and prevent mixing with the DT fuel, we require saturation of amplitudes by the time they are of order δt , which implies that wavelengths $\approx \delta t$ are the most dangerous. Perturbations with $\lambda \ll \delta t$ saturate quickly, while perturbations with $\lambda \gg \delta t$ grow slowly. To explicitly show the dependence of our results on λ , we consider the range 0.5

smoother the transition), the more γ is reduced. Of course, experimental limitations determine how many layers we can coat between the high-density pusher and the low-density ablator.

Summary. The exponential growth rates of Rayleigh-Taylor instabilities can be substantially reduced by spreading a fraction of the mass of a heavy shell into a number of subshells. Our computer codes can calculate the growth rates for an arbitrary number of layers, assuming incompressible fluids, and

Fig. 3-15. Continuous density profile for antinix layer: $\rho = \rho_0 e^{\beta y}$ between $\rho_0 \approx 1 \text{ g/cm}^3$ and $\rho_t \approx 20 \text{ g/cm}^3$.

Fig. 3-16. Rate $\gamma_{\max}$ of fastest growing mode (normalized to γ_0) as a function of fraction of original mass of heavy shell spread into n subshells.

Results. In Fig. 3-16, we show $\gamma_{\max}/\gamma_0$, the normalized rate of the fastest growing mode, as a function of the fraction of the original mass spread into 4, 8, or ∞ (continuous) subshells. The optimum density profile is obtained when $\gamma_{\max}$ is minimum, which occurs at about 35%, 45%, and 100% for $n = 4, 8$, and ∞ , respectively, if we choose $\lambda = \delta t$. In Fig. 3-16, we also show the results for shorter ($\lambda = 0.5\delta t$) and longer ($\lambda = 2\delta t$) wavelengths. The rate γ_0 used for normalization corresponds to zero mass spread.

For $n = 4$ and $n = 8$, we have stopped the curves at 53% and 70%, respectively, because at these values $t_{\text{subshell}} = \delta t$, and spreading more mass would only decrease δt , making the remaining heavy shell thinner than the subshells and causing the perturbations to grow even faster. Note that the best profiles occur *before* this situation is reached.

In Fig. 3-17, we show the density profiles for $n = 8$. Figure 3-17(a) shows the original shell with no mass spread; Fig. 3-17(b) shows 45% of the original shell's mass spread into eight subshells; Fig. 3-17(c) shows 70% of the original mass spread into eight subshells. The best profile (for $\lambda = \delta t$) occurs at 45% mass spread, where the thickness of each subshell ($0.2t_0$) is less than the thickness of the remaining heavy shell ($0.5t_0$). We note in Fig. 3-16 that the 45% spread reduces $\gamma_{\max}$ by 1.6, whereas the 70% spread (where the thickness of each subshell is the same as that of the remaining heavy shell) reduces $\gamma_{\max}$ by only 1.3.

It is clear that the growth rates are reduced by making the transition from ρ_h to ρ_l smoother—the more subshells (i.e., the

predict the thickness of each layer and the fraction to be spread in order to suppress the fastest growing mode. Depending on the number of layers that can be coated, reductions in growth rate by a factor of from 1.4 to 3.5 are expected.

A potentially critical assumption made in our analytic work is that the fluids are incompressible. LASNEX simulations will be necessary to establish how the density profiles that we have found would change when compression and other effects, such as heat transfer and varying acceleration, are taken into account.

Authors: K. O. Mikaelian and J. D. Lindl

Fig. 3-17. Density profiles for $n = 8$.

Charged-Particle Fusion

Introduction

During 1982, the LLNL Charged-Particle Fusion program made advances in several areas

- We neared completion of symmetry and pulse-shape studies for heavy-ion reactor targets (at ~ 4 MJ of beam energy) and initiated similar studies for targets (at 1 to 2 MJ of beam energy) applicable to fusion-fission hybrid reactors.
- We modified our computed gain curves for single-shell, ion-beam targets to take into account effects that improved target gain.
- We further explored the relationship between heavy-ion target gain and accelerator phase-space to take into account more detailed modeling of both final-focus conditions and beam expansion before final focus.
- We performed the analyses that provide the bases for several alternate hydrodynamic-like numerical modelings of pinched-beam phenomena; such modeling will be useful in increasing focal-spot beam brightness in high-temperature heating experiments and for studies of pinched-beam propagation in fusion chambers.

In studies of heavy-ion reactor targets, we found that 3.8 MJ of ions with an equivalent range of 10-GeV uranium at 160 TW peak power can be used to drive a double-shell cryogenic target using the double-pulse shape of Fig. 3-18 instead of a single pulse with a long initial foot. Similar double-pulse shapes can be used for targets now being designed for fusion-fission hybrid reactor applications.

Nuckolls⁸ mentioned the possibility of imploding several targets nearly simultaneously in a single-target chamber to take full advantage of high-repetition-rate drivers, such as heavy-ion beams. To exploit the Nuckolls concept, we need not increase the number of beam-entrance ports and beamlines in proportion to the number of targets imploded. For heavy ions, we might use higher beam currents per beamlet,

which is possible if we neutralize the beams before entering the fusion chamber, or if we fully develop the possibility of propagating self-pinch beams. Issues involved in these two beam-propagation modes are discussed in "Influence of Target Physics on Ion Accelerators and Reactor-Beam Transport in Heavy-Ion Fusion" and "A Hydrodynamic Model for Simulation of Pinched Beams in Innovative Final Focusing and in Reactor-Beam Transport," later in this section.

Initial studies of the relationship between heavy-ion target gain and accelerator phase-space requirements were outlined in the 1981 Annual Report.⁹ Recently, we have refined the results to take into account the effects of

- Improved single-shell target gains.
- More detailed modeling of final focus and beam-transport conditions in the fusion chamber.
- Beam expansion to reduce possible emittance-growth effects in the fusion chamber.

The results are summarized in "Influence of Target Physics on Ion Accelerators and Reactor-Beam Transport in Heavy-Ion Fusion," later in this section.

In the 1981 Annual Report,¹⁰ we pointed out that reductions by factors of 2 or more in focal-spot radius for multigap accelerator test facilities are important in that they make disk-heating and other experiments more interesting, even at 1 to 10 kJ (such as experiments with Na^+ ions at ~ 100 MeV). In Ref. 10 we also suggested one possible collective final-focusing mechanism using the pinching B_z azimuthal magnetic field of an external-discharge-current channel. Moreover, self-pinch propagation of beams might provide a useful alternative to heavy-ion ballistic propagation for fusion-chamber designs in which smaller beam entrance ports are desired.¹¹ Self-pinch propagation might also be crucial for light-ion reactor scenarios.¹² It is therefore useful to pursue further investigation, including numerical simulations of three-dimensional pinched-beam-propagation effects, such as those that help us confirm the absence of instabilities as well as those we need to develop the process of beam merging. If applicable, hydrodynamic-like simulations would be highly desirable to reduce computational expense. In "A Hydrodynamic

Fig. 3-18. Suggested pulse shape for driving a double-shell, cryogenic, heavy-ion reactor target.

Model for Simulation of Pinched Beams in Innovative Final Focusing and in Reactor-Beam Transport," we outline analytical derivations of several plasma regimes in which hydrodynamic-like modeling is indeed applicable; we are currently pursuing numerical comparisons with particle codes and analytical results of simpler problems, prior to detailed application. Similar hydrodynamic-like models are available for the numerical simulation of particle beams in axial magnetic fields (found in some ion sources and collective plasma final-focus lenses) and also in accelerators with axial and azimuthal magnetic fields or other focusing fields.¹³

Author: J. W-K. Mark

Revised Gain Curves for Single-Shell Ion-Beam Targets

We previously reported an initial survey^{14,15} of target gains for ICF targets driven by ion beams. Recently we have revised the gain of cryogenic single-shell targets (see Fig. 3-19). The results reported here are suitable for drivers such as heavy-ion accelerators. These results are applicable to cases where the ion beam incident on the target has small divergences of energy and angle.

The important ion-beam parameters are total beam energy E , focal-spot radius r , and ion range R . Peak beam power P is also an important parameter, but for a given target design it is a function of beam energy—one of the parameters already mentioned.

Fig. 3-19. Single- and double-shell targets for ion-beam-driven inertial confinement fusion.

We made numerical target calculations using the LASNEX code¹⁶ for simulating inertial confinement fusion. One of the important considerations in determining the performance of ion-beam-driven ICF targets is the specific energy deposition, which varies as E/r^2R . Taking a cue from this relationship, and from the functional dependence of target gain obtained in our earlier survey,^{14,15} we again looked for some range of values of r and R within which the gain might depend on the parameter combination $r^{3/2}R$ rather than on r and R independently. Our calculations to date can be considered valid for $0.1 < r/E^{1/3} < 0.2$, and with somewhat less accuracy, up to $r/E^{1/3} \approx 0.33$, where r is in cm and E is in MJ. The input energy scales roughly as target mass, and therefore the expression $r/E^{1/3}$ simply reflects the fact that the mass of a target scales as r^3 .

Our new calculations of single-shell target gain are given in Fig. 3-20(a); results for double-shell targets, reported previously,¹⁵ are given for comparison in Fig. 3-20(b). The solid curves are best-estimate target gains for different values of $r^{3/2}R$, where r is the focal radius (cm) of the beam and R is the range (g/cm^2) of the ions. As a specific example of how these curves can be used, we note that a 10-GeV beam of lead ions focused to $r = 0.2$ cm gives $r^{3/2}R = 0.0112$ (cf. Fig. 3-14 of Ref. 14 or Fig. 4 of Ref. 15 for curves of R vs ion kinetic energy). Thus, the 0.01 gain curve would be roughly appropriate for this case.

In calculating the best-estimate curves, we have attempted to include the effects of fluid instabilities as well as reasonable imperfections in target fabrication and illumination symmetry, all of which affect target ignition and burn. Our assumptions are consistent with the conservative criteria used in Ref. 17. The dashed lines correspond to the "best-estimate" band calculated by Lindl for short-wavelength lasers ($\lambda \approx 0.25 \mu\text{m}$). We emphasize that our results are preliminary; target performance for heavy-ion fusion may improve as other designs and modifications are pursued. The dotted ideal-gain curves, which are based on one-dimensional simulations, represent reasonable upper bounds on gain for these types of targets. The best-estimate gain curves have been degraded from the ideal by including our estimates of the effects of instabilities, target imperfections, and illumination asymmetry.

The new gain curves for single-shell targets given in Fig. 3-20(a) are not as much different from those for double-shell targets as was the case in our initial survey.^{14,15} The new curves decrease less rapidly with decreasing energy; hence, single-shell targets are more interesting than double-shell targets at lower input energy. This is especially true for ion-beam fusion because the lower target gain is compensated for by the expected high efficiency of particle accelerators.

One of the important advantages that double-shell, ion-beam targets once seemed

Fig. 3-20. Best-estimate target gains as a function of energy of the heavy-ion beam.

to have over single-shell targets—a considerably higher gain—may not be as great as formerly thought. The estimated gains of double-shell targets with high- Z tampers are more affected than those of single-shell targets by serious uncertainties¹⁸ associated with unstable mixing of the high- Z shell into both the igniter and the main fuel. Although the use of low- Z tampers might reduce the undesirable effects of mixing in double-shell targets, there are other considerations that enter as well in comparing the two types of targets. For example, single-shell targets have the advantage that they are simpler and easier to fabricate. Double-shell targets have a further disadvantage, in that the threshold energy for ignition of the outer fuel occurs at an uncomfortably large driver energy (>3 MJ), which makes their near-term validation more difficult.¹⁸ Without near-term validation, we might have to move beyond this threshold by a factor of 2 to be conservative.

At present, the main advantage of double-shell targets for ion-beam drivers appears to be in reducing input-power requirements. For a given input energy, the double-shell target requires only about half as much power as the single-shell target. The presently estimated peak-power requirement is given as a function of energy

Fig. 3-21. Peak beam power required as a function of input beam energy for single- and double-shell targets.

in Fig. 3-21. The numerical labels on the curves refer to different values of $r^{3/2}R$. As in the case of target gain, these curves are somewhat uncertain, but they should be adequate for accelerator design. The pulse length is not equal to E/P because our targets require pulse shaping. Also, for a target with fixed yield, we are following a curve

of increasing pulse energy and power requirements when we increase the parameter $r^{3/2}R$.

Authors: J. D. Lindl and J. W-K. Mark

Influence of Target Physics on Ion Accelerators and Reactor-Beam Transport in Heavy-Ion Fusion

The interface between target physics and accelerator technology in heavy-ion fusion is defined by target requirements, which include the following simplified set of parameters: total energy E , pulse duration τ , spot radius r , ion mass number A , and kinetic energy T . From the accelerator side, it is useful to discuss the six-dimensional phase-space requirement, which represents an important constraint on all driver-system designs. The connection between the phase-space volume produced by the accelerator and that entering the target is made by properties of the final-focusing lens systems.¹⁹

At first, simple assumptions about the final beams were used, without considering space-charge effects in the transport lines leading to the lenses, within the lenses, and within the reactor. The justification for this approach was, in part, that for power-plant scenarios using singly charged heavy ions, the individual beamlet currents are small enough that perturbations up to and through the lenses due to space charge are not very large. In addition, various reactor regimes were considered, some of which appeared to allow for final beam propagation without unacceptable spreading. At present, it seems that significant corrections may be needed only for charge state Z greater than 1 and beamlet current greater than 1 kA.

From one point of view, there might now be less need to consider such corrections because all current driver designs use $Z = 1$. The original HIBALL design has been redesigned from $Z = 2$ to $Z = 1$ for several reasons, including the desire to minimize a dangerous longitudinal microwave resistive-wall instability in the ring system. On the other hand, Nuckolls²⁰ recently suggested exploiting the repetition-rate capability of drivers by igniting several targets in rapid succession in the same reactor vessel;

this is most attractive if the number of beamlines per target is minimized. Thus, there may be a new reason to consider currents larger than 1 kA. In any case, fusion considerations almost always favor a minimum number of beamlets and entrance ports. Among possible methods for handling larger currents are neutralization in the reactor and a beneficial nonlinear effect of beam expansion just upstream of the final lenses.

Accelerator Phase-Space and Final-Focus Considerations. In 1978, Judd²¹ gave an analysis of phase-space requirements for systems using rf linear accelerators, which he later²² extended to include induction linear accelerators. Judd dealt with the constraints arising from final-focusing lens systems only in the "ballistic limit," i.e., fully neutralized²³⁻²⁵ or low-current beams; corrections for some transverse-space-charge self-force effects are discussed below. A six-dimensional phase-space volume is defined by ranges in coordinates x, y, z and momenta p_x, p_y, p_z . For fusion targets, all six of these parameters have bounds, which, if extended, will lower the gain. The curves of Figs. 3-20 and 3-21 (see preceding article) display this effect for spot size, which limits x and y . The beam length z is limited by the peak-power requirement. Known properties of final-focusing lens systems place limits on p_x and p_y (related to transverse emittance), and on δp_z (through chromatic aberration).

A useful expression for the maximum permissible dilution factor $D_{\max}$ is derived in Ref. 22 (p. 21); $D_{\max}$ is defined as the ratio of the six-dimensional phase-space acceptance volume of the final lens system to the occupied six-dimensional emittance volume at some reference point in the accelerator. It represents the upper limit on the product of all actual phase-space dilution factors D_i incurred between that point and the final lenses

$$\prod D_i \leq D_{\max} \equiv \frac{V_{6f}(\text{acceptance})}{V_{6f}(\text{occupied})}.$$

Some of the transverse and longitudinal emittances and acceptances (four-dimensional and two-dimensional phase-space volumes) can be expressed in terms of spot radius, pulse duration, and other parameters (see, for example, Ref. 22, pp. 5

and 9). Total energy and ion kinetic energy are brought in through their relation to linac current I and on-time t_L ; $It_L = qeE/T$, with qe the ion charge. The result may be written as the product of three factors, each containing parameters of a part of the total system such that

$$D_{\max} = F_{\text{TI}} F_L F_{\text{FF}} \quad (24)$$

Only target and ion parameters enter the first factor

$$F_{\text{TI}} = \frac{\beta^3 \gamma^2 c \tau r^3 T}{4E} \quad (25)$$

with c the speed of light and β and γ the usual relativistic factors. (Although normalized differently, F_{TI} is equivalent to that in Ref. 26.) The second factor, F_L , is the "linac figure of merit" and depends only on properties of the linac driver; it is the electric charge per unit six-dimensional phase-space volume (i.e., six-dimensional brightness). The third factor, F_{FF} , depends only on final-focusing system parameters; its value is $n_f B_Q F_s / k_G$, in which n_f is number of final beamlets, B_Q is pole-tip field in the final lenses, and F_s and k_G are dimensionless parameters determined by lens design details.

For rf linac systems, we expect large dilution factors in the multistage linac tree and from multiturn injections, and we expect additional dilutions during acceleration, debunching, and the numerous other beam manipulations. For induction-line systems there are no manipulations; however, as is also the case for most of the rf linac-system manipulations, there is no practical experience to guide estimates of safety factors. For both types, it is desirable or necessary to strive for larger $D_{\max}$ because of the many uncertainties. The constraint is eased by increasing n_f , rather than decreasing it, as mentioned above. From the accelerator side, it is therefore desirable to increase F_{TI} to assist in meeting the phase-space constraint although increasing T would raise driver cost.

Calculated properties of both single- and double-shell targets have been used to estimate the dependence of target gain on E , T , and F_{TI} (Ref. 26); the dependence on T is weak for fixed E and F_{TI} . Gain decreases as F_{TI} increases, displaying the price paid to ease this driver design constraint; a balance must be sought. Because F_{TI} varies as the

cube of the spot radius, unwanted "effective emittance growth" in the reaction chamber could be serious; beam neutralization²³⁻²⁵ or other effects^{27,28} to combat this tendency might be needed. Preliminary estimates suggest that self-pinch beam propagation could possibly increase F_{TI} by a factor of from 2 to 3. From the accelerator side, more beamlets, brighter linacs,²⁹ and smaller dilutions would ease the constraint.

In our analysis, we have omitted space-charge forces in and beyond the final lenses. We now show that space-charge corrections may be small. It has been shown by Garren³⁰ that with space charge there is some focal-spot growth for a beam of uniform current density (flat profile), but for most parameter regimes the growth factor is rather small. Another source of growth, emittance growth due to nonflat beam profile, has been studied more recently²³⁻²⁵; for some parameters (larger ion charge and beam current, smaller ion kinetic energy, and less flat profiles) this effect may be more important. Such parameters will enter the relation between "intended" and actual spot radii r and r_s , and will therefore also appear in a revised expression for F_{TI} that would take spot growth into account. In particular, the functional dependence of F_{TI} on T would be modified.

There is a dimensionless parameter U , which characterizes the importance of transverse space-charge force for an unneutralized beam in vacuum; from Garren's work, it may be written as

$$\begin{aligned} U &= \frac{ZeI}{4\pi\epsilon_0 (\beta\gamma)^3 A m_p c^3} \\ &= 3.20 \times 10^{-8} \frac{Z I_b(A)}{A(\beta\gamma)^3} \\ &= \frac{(Ze)^2 (P/N_b)}{4\pi\epsilon_0 (\beta\gamma)^3 T A m_p c^3} \\ &= 1.34 \times 10^{-7} \frac{Z^2 P(\text{TW})}{N_b (\beta\gamma)^3 T(\text{GeV})} \frac{238}{A} \quad (26) \end{aligned}$$

with Ze the ion charge, $A m_p$ the ion mass, β and γ the usual relativistic factors, I_b the beamlet electric current, P the power, N_b the number of beamlets, and T the ion

kinetic energy. For a flat profile, Garren showed that space-charge force increases target spot size roughly by

$$\frac{r_s - r}{r} \approx \frac{2U}{\theta_0^2} \ln \frac{L\theta_0}{r} \quad (27)$$

with r_s and r spot radii with and without this growth, respectively, θ_0 the beam convergence angle at the lens exit, and L the standoff distance from lens to target; $L\theta_0$ is the beam radius r_0 at the lens exit. For example, consider a 1-kA beam of U^{+1} ions at 10 GeV, $\beta\gamma \approx 1/3$, $L = 5$ m, $r_0 = 0.2$ mm, $r = 2.5$ mm. Then $U \approx 3 \times 10^{-6}$, $\theta_0 = 4 \times 10^{-2}$ rad, and $(r_s - r)/r \approx 0.02$. As suggested earlier, we have found that space charge in final focus may not be a large effect.

Beam Emittance Growth During Transport Across Reactor Chamber. An analysis²³⁻²⁵ of the emittance growth in beams with a rounded radial profile is expressed in terms of rms beam radii and an "rms emittance," ϵ . The emittance growth (neglecting the small effect above and assuming no initial rate of growth at the lens exit) may be written as

$$\epsilon^2 = \epsilon_0^2 + (\lambda UL)^2 \quad (28)$$

with ϵ_0 the rms emittance at the lens exit and λ a dimensionless shape factor; λ is 0, 0.149, and 0.388 for flat, parabolic, and Gaussian profiles, respectively. Thus

$$R_s \approx R \left[1 + \left(\frac{\lambda UL}{\epsilon_0} \right)^2 \right]^{1/2} \quad (29)$$

where R_s , R , and ϵ_0 (unnormalized) are rms values, as distinguished from r_s , r , and ϵ_i (normalized) usually employed²² for flat beam profiles. We define the profile-dependent emittance growth parameter ϵ_g (or ϵ_{gv} for vacuum propagation) as

$$\epsilon_g \approx \epsilon_{gv} \equiv \lambda UL \quad (30)$$

For $\epsilon_0 \approx 0$, ϵ_g is the actual emittance growth.

For the example following Eq. (27), a beam with Gaussian profile has $\epsilon_{gv} \approx$

$0.07\epsilon_0$. This is still not a large effect; however, if the ion charge state is increased to $+2$ with the same particle current (twice the electrical current), the value of ϵ_{gv} is increased to about $0.4\epsilon_0$.

In a soft vacuum ($\sim 10^{-4}$ to 10^{-3} Torr), the ionization of the medium provides electrons that partially neutralize the beam²³⁻²⁵ and reduce emittance growth. This neutralization process requires the ion beam to have a potential drop sufficient to accelerate the plasma electrons to the beam velocity (~ 25 kV is needed for a 10-GeV U^{+1} beam). This process leads³¹ to the modified emittance growth parameter $\epsilon_{g'}$

$$\epsilon_{g'} \approx \epsilon_{gi} \equiv (1 - f_{\text{neut}}) \epsilon_{gv} \quad (31)$$

where f_{neut} is the effective fractional neutralization of the beam. For 1.2×10^{-3} Torr Li vapor pressure in the fusion chamber, $f_{\text{neut}} \approx 0.5$ (cf. Refs. 23, 25). As the vapor pressure in the fusion chamber is increased, the fractional ionization saturates roughly at the gas density of $n_g \approx (\sigma_{\text{ion}} L)^{-1}$, where σ_{ion} is the ionization cross section. For U ions propagating in Li vapor, this density is $n_g \approx 10^{13} \text{ cm}^{-3}$, for pressure $\sim 10^{-3}$ Torr. Increasing vapor pressure beyond this level no longer increases the favorable effect of f_{neut} , but only increases beam stripping.

Influence of Target Requirements, Final Focus, and Beam Transport on Accelerator Phase-Space. We now illustrate some of the restrictions that could occur if the considerations, above, are not carefully balanced. Target gain as a function of F_{TI} , E , and T can be calculated.^{26,31,32} The gains represented by Figs. 3-20 and 3-21 were used for such a calculation for $A = 238$, with results that were first shown in Ref. 26. The dependence on T is very weak. Our results, summarized in Figs. 3-22(a), 3-23(a), and 3-24(a) assume ballistic focusing of a fully neutralized beam in a hard vacuum ($< 10^{-4}$ Torr). Although calculations (cf. Refs. 23 and 25) suggest that the beam could be neutralized by conjection of electrons, more detailed simulations, and perhaps experiments, are needed to confirm the applicability of the process. On the other hand, electrical efficiency for liquid-metal-wall-protection schemes points to somewhat higher vapor pressures (cf. Refs.

33 and 34). A confluence of these requirements occurs at from 10^{-4} to 10^{-3} Torr, for which partial beam neutralization by ionization of the propagation medium (e.g., Li vapor) is modified somewhat by mild stripping of the beam ions.^{23,25}

Emittance growth results in a stronger dependence of F_{TI} on ion kinetic energy T . At fixed gain with a specified type of target, and for fixed beam-pulse energy, both the product $r_s^{3/2}R_i$ and the peak power P are determined (cf. Refs. 31 and 32; R_i is the ion range in g/cm^2). For a given heavy-ion species with energy above a critical value ($T > 8$ GeV for uranium), R_i is roughly proportional to $T^{3/2}$; hence, $r_s T$ is determined, or $r_s \propto T^{-2}$. On the other hand, the beam cur-

rent I_b is proportional to Z/TN_b at fixed power P (where N_b is the number of beamlets); and the ions are still nonrelativistic, so that $\beta\gamma \propto T^{1/2}$. For a fixed propagation distance from final focusing lenses to the target, the "intended" focal radius varies as

$$r \approx r_s \left[1 + \left(\frac{e_g}{e_0} \right)^2 \right]^{-1/2}$$

or

$$\frac{r}{r_s} = \left[1 - \left(\frac{r e_g}{r_s e_0} \right)^2 \right]^{1/2} \quad (32)$$

Fig. 3-22. Phase-space factor F_{TI} vs ion kinetic energy for single-shell targets, for several combinations of fixed target gain, beam pulse energy E , and peak power P : (a) For $e_{\text{gc}} = 0$. (b) For $e_{\text{gc}} = 1$.

Fig. 3-23. Phase-space factor F_{TI} vs ion kinetic energy for double-shell targets, for several combinations of fixed target gain, beam pulse energy E , and peak power P : (a) For $e_{\text{gc}} = 0$. (b) For $e_{\text{gc}} = 1$.

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If A , P , N_b , and $r_s T$ are held fixed (for fixed gain), the target-ion phase-space factor F_{TI} which is proportional to $r^3 T^{5/2} \propto T^{-1/2}$ without growth, has for the desired r_s (and therefore desired gain) with the dependence

$$F_{TI} \propto r_s^3 T^{5/2} \propto T^{-1/2} \left[1 - \left(\frac{r e_g}{r_s e_0} \right)^2 \right]^{3/2}, \quad (33)$$

displaying the tighter constraint on accelerator phase space. Many parameters enter into the correction term. It can be shown^{22,30} that for fixed A

$$\left(\frac{r}{e_0} \right)^2 = \left(\frac{\beta \gamma r}{\epsilon_1} \right)^2 \propto \frac{T^{1/2}}{X B_Q}, \quad (34)$$

with X the final quadrupole bore radius and B_Q its pole-tip field. From the equations above,

$$\frac{e_g}{r_s} \propto \frac{UL(1 - f_{neut})}{r_s} \propto \frac{\lambda Z^2 L(1 - f_{neut})}{T^{3/2} r_s T}. \quad (35)$$

Thus

$$F_{TI} \propto T^{-1/2} \left\{ 1 - C_1 T^{-5/2} Z^3 \right. \\ \left. \times [\lambda L(1 - f_{neut})]^2 \right\}^{3/2} \quad (36)$$

because $r_s T$ is held fixed; dependences on P , N_b , X , B_Q , and lens-design parameters are contained in C_1 .

For a neutralized beam, the correction vanishes, and the weak dependence on T is that mentioned in Ref. 26; but for partial or

little neutralization the dependence on T is altered for nonflat beams unless C_1 and Z are small. We neglect the dependence of X on T , because it varies only weakly due to an upper limit²² set by third-order geometric aberration

$$X = k_1 \left(\frac{\beta \gamma A m_p c r}{Z e B_Q} \right)^{1/2}, \quad (37)$$

where the constant $k_1 \approx 0.85$.

Typical values of emittance growth Δe are illustrated in Table 3-2; they were obtained from runs of the HIPPO code²⁵ for a 1-kA beam of 10-GeV U^{+1} ions (Gaussian profile and negligible initial emittance) propagating into a reaction chamber with lithium vapor at the indicated pressures and Li particle densities. For the same beam propagating different distances in vacuum,²³⁻²⁵ Δe scales roughly linearly, as suggested by Eq. (26). In the presence of Li vapor at fixed pressure, the linear dependence holds when the effective fractional neutralization f_{neut} has reached its equilibrium value (after a few meters of propagation). For a fixed f_{neut} (fixed pressure and, in the first few meters of propagation, also fixed L), our numerical data indicate that estimates of emittance growth (e_g in this case, where $e_{ti} \approx 0$) could be obtained by scaling the values in Table 3-2 according to Eqs. (26), (30), and (31). Beyond a pressure of 1.2×10^{-3} Torr, f_{neut} saturates.

Our sample numerical data are meant only to suggest qualitative conclusions

Fig. 3-24. Target gain vs F_{TI} for single- and double-shell targets, evaluated for U^{+1} ions: (a) For $e_{gc} = 0$. (b) For $e_{gc} = 1$.

regarding the validity of the scaling laws (26), (29), (30), and (31), which lead to Eqs. (32) to (34). For application to typical fusion-reactor designs, we incorporate initial emittance ϵ_0 according to Eq. (28). A first estimate of the parameter ϵ_g might now proceed according to Eqs. (33) and (34).

The behavior of F_{T1} vs ion kinetic energy T (for U^{+1} ions) is illustrated in Fig. 3-22 for single-shelled targets and in Fig. 3-23 for double-shelled targets. The fixed values of ϵ_{gc} from Eq. (34), target gain G , beam pulse energy E (MJ), and total beam power P (TW) are as shown in the figures.

Figure 3-24 illustrates the "typical" reactor-beam transport modifications to the relationship between target gain and F_{T1} for the ballistic focusing scenario. Effects such as those in Figs. 3-22 and 3-23 are averaged out. Figure 3-24 is, again, for U^{+1} ions, showing curves of fixed-beam-pulse energy E (MJ) for single- and double-shelled targets. Clearly the case of full neutralization, $\epsilon_g = 0$, is the same as that of Ref. 26.

Consider vacuum propagation ($f_{\text{neut}} = 0$) over $L = 10$ m of $N_b = 16$ total Gaussian ion beamlets ($T = 10$ GeV, $A = 238$, $Z = 2$) with beam shape factors $\lambda = 0.388$, $\lambda_i = 0.707$, and final focus parameters $XB_Q = 2.5$ (T·m), $k_G = 40$. In this case $\epsilon_{gc} = 8.3$ according to Eq. (34); but if the charge state of the beam ions is decreased to $Z = 1$, then $\epsilon_{gc} = 1.0$. On the other hand, propagation in 1.2×10^{-3} -Torr Li vapor would reduce $(1 - f_{\text{neut}})$ by a factor of 2 in the $Z = 1$ case, so that we have $\epsilon_{gc} = 0.25$, a small value. Active coinjection of an electron beam³⁵ with 1-keV transverse beam temperature into the $Z = 1$, but $N_b = 4$ (4-kA), beamlets reduces the $\epsilon_{gc} = 4$ 2 case to the actively neutralized beam value of $\epsilon_{gc} = 0.015$.

Finally, beam neutralization has a greater positive effect than the deleterious one of beam stripping (certainly for Li vapor at $< 10^{-3}$ Torr), because F_{T1} is only inversely proportional to the peak power P_{peak} , which is decreased somewhat by the process of stripping some beam ions (cf. Ref. 25).

Modifying Final Focus to Control Beam-Emittance Growth in the Fusion Chamber. Here we point out some limitations of the cold converging beam (CCB) model of reactor beam transport used in the foregoing discussion. Under some circum-

Pressure in Li (Torr)	Particle density in Li (cm ⁻³)	Emittance growth (cm-mrad rms, unnormalized)		
		$\Delta L = 2$ m	$\Delta L = 4$ m	$\Delta L = 7$ m
0	0	0.43	0.86	1.53
2.5×10^{-4}	3.1×10^{12}	0.40	0.76	1.22
5.0×10^{-4}	6.3×10^{12}	0.34	0.54	0.80
1.2×10^{-3}	1.5×10^{13}	0.23	0.42	0.67
2.4×10^{-3}	3.0×10^{13}	0.20	0.38	0.58
3.6×10^{-3}	4.5×10^{13}	0.17	0.26	0.51

stances, the CCB model could be pessimistic in its assumption of a cold beam with a rounded (nonflat) profile at the final lens. An alternative line of thought accounts for the beam expansion from its initially small size in the accelerator to its much larger size at the final lens by ballistic transport. In one scenario, a beam with initially rounded profile at the accelerator exit assumes a uniform (flat) profile as it expands to the final lens (Fig. 3-25). The final lens then focuses a flat-profile beam to the target pellet with little or no emittance growth. We call this model the expanded flat focus (EFF).

A semiquantitative analysis of the alternative EFF model illustrates the merits and limitations of the CCB model. First, we introduce a dimensionless shape factor, C , which describes the transverse potential energy of the beam, such that

$$C = -2 \int_0^b \frac{dr}{r} \left(\frac{I_r}{I_0} \right)^2 + \ln \frac{b^2}{R^2} \quad (38)$$

where b is an arbitrary, large integration limit and I_r is the current within a radius r . The quantity C assumes a maximum value of 0.192 for a uniform beam and has values for different beam profiles given in Table 3-3. We can show²⁴ that the emittance variation is directly related to changes in the beam shape by

$$\frac{de^2}{dz} = KR^2 \left(\frac{dC}{dz} \right) \quad (39)$$

In supporting phenomenological relations from particle-simulation studies with Eq. (39), Barletta²⁷ has estimated the beam emittance at the target to be

$$\epsilon_n \approx -0.1Z_r K + (0.1R_s R_r K)^{1/2} \quad (40)$$

for $0.1Z_r K \ll 0.1R_s R_r K$.

Table 3-2. Emittance growth calculated by the HIPPO code for a 1-kA beam of 10-GeV U^{+1} ions.

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In contrast, the CCB model predicts that the emittance at the focus, ϵ , is

$$(\epsilon_n^*)^2 \alpha^2 K^2 Z_f^2 \quad (41)$$

Hence, the question as to whether the CCB model is optimistic (and therefore not unduly restrictive) reduces to the inequality, $(\epsilon_n^*)^2 < \epsilon_n^2$; that is,

$$0.5KZ_f < (0.1R_a R_f)^{1/2} K^{1/2} \quad (42)$$

To proceed, we must evaluate the product $R_a R_f$ and find its functional relation with K .

The solution to the envelope equation for the expansion of a beam under its own space charge is the universal expansion curve shown in Fig. 3-26. The laminar flow assumed by Eq. (41) provides a way to estimate the relation among Z_{lens} , R_a , and R_f . Equating the emittance predicted by Eq. (39) with that from Eq. (28), we have

$$\alpha^2 K^2 (Z_{\text{lens}})^2 = \Delta C K R_a R_f \quad (43)$$

Plotting Eq. (43) with the universal space-charge curve (Fig. 3-26), we find that the product $R_a R_f$ is independent of U and that

Fig. 3-25. Definition of geometry for the EFF beam-transport model.

Table 3-3. Shape factors for various beam profiles.

Profile name	I/I	C	α
Flat	$\frac{1}{\pi a^2}$	0.193	0.000
Parabolic	$\frac{2}{a^2} \left(1 - \frac{r^2}{a^2}\right)$	0.182	0.149
Gaussian	$\frac{\exp(-r^2/R^2)}{\pi R^2}$	0.116	0.388
Annulus	$\frac{(r-a)}{2\pi a}$	0.000	—
Bennett, truncated at $r = 5a$	$\frac{26}{25a^2} \left(1 + \frac{r^2}{a^2}\right)^{-2}$	-0.136	0.885

$R_f \approx 3R_n$. Thus we can solve Eq. (42) for K to find the scaling relation,

$$K < 1.3 \times 10^{-6} \left[\frac{10^3}{Z_f(\text{cm})} \right]^2 \left[\frac{R_a(\text{cm})}{\sqrt{2}} \right]^2 \quad (44)$$

that must hold for the CCB model to be an optimistic estimate. For $Z_f = 10$ m and typical beam-in-accelerator sizes, the CCB model is unduly pessimistic for beams with charge state greater than +1. Nonetheless the EFF model is not a practical solution for beam transport. The beam size at the target is given by

$$R_n \approx \frac{\epsilon_n Z_f}{R_f} \quad (45)$$

For R_a in the range from 1 to 2 cm, Eq. (42) shows the beam size at the lens to be too small to allow a focal spot < 2.5 mm at 10 m for $K < 2 \times 10^{-6}$ (10 GeV, 1 kA, U^{+1}).

The analysis of the preceding section has been compared with a particle simulation of beam behavior made with the code DESTIN.^{24,36} The code pushes particles in a single disk using the single-particle equations of motion. The net force on the particles is evaluated self-consistently from the distribution of particles. The assumption of an axisymmetric distribution of particles in the field solver necessitates treating the final lens as a simple, thin-lens approximation.

Figure 3-27 compares the CCB model—both analytic and simulation—with the EFF model for a choice of K (given Z_f) just below the limit of Eq. (44). For $Z_f = 300$ cm and $R_a = 2$ cm, K should be less than 2.5×10^{-5} for CCB to be optimistic in its prediction of ϵ_n . As Fig. 3-27 shows, the analytic CCB prediction of ϵ_n is about 20% smaller than that of the EFF model for a choice of $K = 1.9 \times 10^{-5}$. Other simulations of beam behavior confirm the analysis based on Fig. 3-26 that $R_a R_f$ is independent of K , with $R_f \approx 3R_n$.

Because the beam radius at the time that the beam first assumes a uniform distribution is usually too small for efficient ballistic focusing over 10 m, we must next consider letting the beam continue to expand. Barletta²⁷ has shown that, as a beam with an initially Gaussian profile continues to expand, about 25% of the charge piles up at the now well-defined beam edge.

Eventually, space-charge forces in the region of accumulated charge will disperse

this feature as the beam tends asymptotically to a uniform distribution. Therefore, we expect C to decrease as the beam develops the accumulation. It will reach a minimum value when the concentration at the edge is maximum and then will increase. We denote the condition of maximum pileup by the subscript m . The decrease and rise of C will be accompanied by a decrease and subsequent increase in e^2 as prescribed by Eq. (39).

To estimate the beam emittance at $Z = Z_m$, we write e^2 as the sum of two contributions (i.e., that from $Z = 0$ to Z , and that from Z_t to Z_m):

$$e_m^2 = (0.1)R_a R_f K - (0.01)R_f R_m K \quad (46)$$

We can use the procedure of Fig. 3-26 to find that $R_m \approx 2R_f \approx 6R_a$. Hence, at Z the emittance is

$$e_m^2 \approx 0.02R_a^2 K \quad (47)$$

Comparison with simulation indicates that Eq. (47) is accurate to about 20%.

With a radius of $\sim 6R_a$ at Z_m , the beam is large enough for final ballistic transport. However, even after passage through the thin lens, the beam will tend to wash out the charge concentration (i.e., become flat) via space-charge forces. The shape factor C will increase, and therefore the emittance of the beam will initially increase. Nonetheless, preliminary studies³⁷ of beam focusing after expansion to R_m show transport efficiency to the target pellet far exceeding the equivalent-radius CCB predictions. Apparently beam expansion and passage through the lens creates a velocity distribution in the

beam, which resists the space-charge distortion of phase space as the beam tries to come to a tight focus. For example, for a 12-GeV U^{+2} beam with 1-kA particle current, the CCB model predicts a transport efficiency of 10% to a 2.5-mm pellet for $Z_{\text{focus}} = 10$ m, whereas the expansion to R_m leads to an efficiency of 53%.

The most important conclusion to draw from our analysis is that the entire transport from the accelerator to the target must be considered in choosing a final transport scenario.

Discussion: Is Active Beam Neutralization in the Reaction Chamber Necessary?

Although we have not found conclusive arguments to show that beam neutralization in the fusion chamber is essential, it is a helpful procedure deserving further study. If beamlet currents greater than 1 kA were used (cf. the initial discussion in this article, at ion charge state $Z > 1$), then the possible undesirable growth of effective beam emittance in the reactor chamber could be overcome by beam neutralization.

Fig. 3-26. Beam expansion as a function of propagation distance. Solid curve is universal space-charge expansion curve; dashed curve equates formulas for laminar flow and shape-change emittance growth.

Fig. 3-27. Comparison of the predicted limits of the CCB model with a simulation by the DESTIN code.^{24,36} (a) Change in beam shape during expansion and focal phases of transport. (b) Corresponding emittance variation. Target pellet is located at beam neck.

Alternatively, it could be compensated to some extent by redesigning the final focusing lenses (cf. "Modifying Final Focus to Control Beam-Emittance Growth in the Fusion Chamber" earlier in this article). But it is not clear that 4-kA beamlets can be propagated the required 8 to 10 m without neutralization. Such higher beam currents are of interest because they have the economic advantage of minimizing the number of beamlines and reactor entrance ports.

Neutralization also removes from further consideration the issue of target charge-up.³⁵ Typically, 1 MJ of heavy ions deposits about $10^{-4} Z_i(1 - f_{\text{neut}})$ coulombs of positive charge on target. Partial neutralization of the beam is accounted for by defining the fractional neutralization f_{neut} so that $Z_i(1 - f_{\text{neut}})$ is the effective average charge of the ion arriving at the target. The consequences of this charge accumulation for heavy-ion targets are complex. It is not even clear that there is a problem.

One plausible scenario is that a few percent of the heavy ions miss the main target, but before escaping are stripped to a high charge state by passing through a blowoff plasma. As a consequence, this small number of escaping ions leaves behind a substantial negative charge on the target, which might more than compensate for the positive charge brought in by the average ions with $Z_i(1 - f_{\text{neut}}) < 1$ to 2. By contrast, in the worst-case scenario the target is positively charged, and attracted electrons might receive a total of

$$\frac{Z_i e^2 f_{\text{neut}} (1 - f_{\text{neut}})}{8\pi\epsilon_0 r} = 4500 Z_i f_{\text{neut}} (1 - f_{\text{neut}}) \frac{1 \text{ cm}}{r} \text{ joules} \quad (48)$$

in kinetic energy. Most of these electrons do not penetrate deep enough into the target to cause preheat (especially in tapered targets). Preliminary simulations³⁵ using the HIPPO code suggest that preheat is not a problem. Preheat due to target charge-up remains a possibility that could be overcome by neutralization of the beams.

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A Hydrodynamic Model for Simulation of Pinched Beams in Innovative Final Focusing and in Reactor-Beam Transport

In last year's Annual Report,³⁸ we pointed out that for the suggested heavy-ion-fusion test facility, collective final focus might be useful to enhance the beam energy flux at the focal spot. The process we suggested last year relies on the use of the radial force on a beam due to an azimuthal-pinch magnetic-field, B_θ , of a discharge current channel. It has also been suggested, regarding reactor-beam transport, that self-pinched beams, or aided pinches (using a counter-injected electron beam), have the advantage of requiring only small entrance ports.³⁹ In addition, self-pinched beams present the possibility of transporting higher beam currents or requiring fewer beam ports; these factors may be useful if we consider the simultaneous implosion of more than one target in a single chamber.

Detailed simulations of such pinched-beam processes are still preliminary. Envelope codes are crude and do not allow freedom for instability studies. Particle codes are frequently too costly to run for three-dimensional studies, such as those required to investigate and prevent instabilities, as well as those needed to develop the process of the merging of multiple beams into an externally initiated propagation channel. In this article, we propose that several hydrodynamic-like descriptions could be used under certain conditions; in particular, the closure of hydrodynamical models for pinched beams is discussed within adiabatic assumptions mathematically analogous to those of Chew, Goldberger, and Low.⁴⁰ We first derive a simplified drift-kinetic equation that could already be used for closure. In addition, some plasma regimes allow more simple closure in terms of adiabatic equations of state.

Hydrodynamic models have been used in particle-beam research^{41,42} as approximations to the Vlasov theory, just as in other fields of plasma physics. Sometimes equations of state are assumed by dropping heat fluxes,⁴³ without a detailed attempt to model possible beam behavior, such as

anisotropic stresses, even in the transverse directions. Recently, we have had occasion to reexamine^{44,45} the foundations of such hydrodynamic models. Our reexamination was undertaken for the purpose of studying propagation of self-pinch and external-current-pinch beams for ion-beam inertial-fusion research.^{38,39,46,47} This study also applies to beams in axial magnetic fields (ion sources,⁴⁸ plasma lenses for final focus, and others) or beams in accelerators.⁴⁴ As a result of our reexamination, we now outline results showing that dynamically consistent equations of state can be derived from the Vlasov equation for some beam models as a mathematical analogue of the Chew, Goldberger, and Low adiabatic theory for magnetohydrodynamics⁴⁰ and of the theory of Berman and Mark for galactic dynamics.^{49,50} For our derivations, we of course assume similar generalized adiabatic conditions.

Generalized Adiabatic Conditions. One of the generalized adiabatic conditions relates to the character of the particle orbits, which we allow to be rather general (rosette orbits in central fields), except that the ratio of maximum to minimum radial excursion is limited by something less than a factor of 2. Thus, the particle orbits can deviate from laminar flow with circular orbits, but not by too much. We note that if the reactor beam has zero net angular momentum, then it is necessary to consider at least two counter-rotating beam components of this type. By remaining closely tied to the actual particle orbits, we retain the possibility of wave-particle resonances typical of resistive-mode instabilities. It is not clear that other recent attempts⁵¹ at fluid models have retained such resonances.

In terms of the usual cylindrical coordinate system (r, θ, z) with the beam axis in the z direction, the above and other adiabatic assumptions are summarized by

$$\begin{aligned} \epsilon &= 0 \left(\frac{\delta r}{r} \right) = 0 \left(\frac{r p_r}{p_\theta} \right) \\ &= 0 \left[\left[\frac{\partial}{\partial t} + (\dot{z} - v_{zp}) \frac{\partial}{\partial z} \right] \div \nu' \right] \\ &= 0 \left(\frac{\partial}{\partial \theta} \right) < 1 \quad , \quad (49) \end{aligned}$$

where δr is the deviation from average radius in the particle orbit mentioned above; p_r , p_θ , and p_z are the particle momenta; t is the time coordinate; and ν is the radial oscillation frequency of the rosette orbits [to lowest order; see Eq. (53) below]. Throughout this article, dots over symbols indicate time derivatives, such that $\dot{z}$ is the axial velocity of beam particles. We assume that there is some frame of velocity, v_{zp} , where $[\partial/\partial t + (\dot{z} - v_{zp})\partial/\partial z] = 0(\epsilon)$. This frame need not have the same velocity as the beam, a wave-packet disturbance, or the laboratory, but it is useful for showing the validity of the particular mode of derivation; the final results could be transformed to other frames. We also note that the assumptions in Eq. (49) will not require linearization of Vlasov's equation, so that nonlinear waves could be considered. For example, for nonaxisymmetric-nonlinear disturbances in the distribution function, a hypothetical $F[H, \epsilon b(r, \theta, z, p_r, p_\theta, p_z, t)]$, where H is the Hamiltonian of a moving-beam equilibrium (independent of θ), is consistent with Eq. (49).

For the present treatment, we also assume for simplicity that the dominant external force on a beam particle of charge q is the radial force $f_r = q(E_r - \dot{z} B_\theta/c)$, where E_r is the radial electric field, B_θ is the θ -component of the magnetic field, $\dot{z}$ is the z -velocity, and c is the speed of light. Generalizations to additional central forces are straightforward; nonradial force components are assumed to be of higher order in ϵ .

Transformation of Momentum Space. The mathematical framework follows that of Chew, Goldberger, and Low for magnetohydrodynamics, that of Berman and Mark⁴⁹ for galactic dynamics, and is initiated by transforming Vlasov's equation to a different coordinate system in momentum space (see Ref. 44 for details). In a cylindrical coordinate system about the beam axis, the particular choice of transformed variables is

$$\dot{r} = U(r, \theta, z, \dot{z}, t) + \alpha \sin \Psi \quad (50)$$

and

$$r^2 \dot{\theta} = r^2 \Omega(r, \theta, z, \dot{z}, t) + \beta \cos \Psi \quad (51)$$

This transformation identifies the general rosette orbit as consisting of a rapid circulation, Ω , and additional rapid oscillations represented by the phase, Ψ . Since U and $(r\Omega)$ will later be identified with low approximations of the hydrodynamic motions, the analogue of "gyro-motion" is relegated to the rapid quantities ($\alpha \sin \Psi$) and $(\beta \cos \Psi)$.

In the terminology of adiabatic orbit theories, U and $(r\Omega)$ are "drifts." As part of the transformation process, we have identified the drifts $(r\Omega)$ and U , and the fast gyration frequency $\nu \simeq \Psi$, with

$$\begin{aligned}\Omega^2(r, \theta, z, \dot{z}, t) &= -\frac{1}{m} f_r \\ &= \frac{q}{mr} \left(\dot{z} B_\theta - E_r \right), \quad (52)\end{aligned}$$

$$\nu_2(r, \theta, z, \dot{z}, t) = \frac{1}{r^3} \frac{\partial}{\partial r} (r^4 \Omega^2), \quad (53)$$

$$\begin{aligned}U(r, \theta, z, \dot{z}, t) &= \frac{2\Omega}{r\nu^2} \\ &\times \left[\frac{\partial(r^2\Omega)}{\partial t} + \Omega \frac{\partial r^2\Omega}{\partial \theta} + \dot{z} \frac{\partial r^2\Omega}{\partial z} - \frac{1}{m} f_\theta \right], \quad (54)\end{aligned}$$

and the ratio of amplitudes,

$$\frac{\beta}{\alpha} = \frac{r\nu}{2\Omega}. \quad (55)$$

In Eqs. (52) and (54), m is the mass of the beam particle and f is the body force.

The Drift-Kinetic Equation. Following Chew, Goldberger, and Low, we expand the distribution function,

$$F = \sum_{j=0}^{\infty} \epsilon^{j/2} F_j, \quad (56)$$

which begins with $0(\epsilon^{-2})$ because the spatial number density $n = 0(1)$ in ϵ . Under our adiabatic assumptions, it is straightforward to derive from the transformations in Eqs. (50) and (51) the "drift-kinetic" equation to first order in ϵ ,

$$\begin{aligned}\frac{\partial F_0}{\partial t} + U \frac{\partial F_0}{\partial r} + \Omega \frac{\partial F_0}{\partial \theta} + \dot{z} \frac{\partial F_0}{\partial z} + \dot{z} \frac{\partial F_0}{\partial \dot{z}} \\ = 0 \quad (57)\end{aligned}$$

Here we have utilized the fact that the distribution function, $F(r, \theta, z, \dot{z}, \mu, t)$, is independent of the rapid phase, Ψ , a fact that follows because the zeroth-order (in ϵ) form of Vlasov's equation is $\partial F_0 / \partial \Psi = 0$, as in the Chew, Goldberger, and Low theory. We have also simplified the drift-kinetic equation [Eq. (57)] by taking advantage of the existence of an adiabatic invariant, μ , expressed to lowest significant order in ϵ as

$$\mu = \frac{2}{\nu} \langle (\dot{r} - U)^2 \rangle, \quad \frac{d\mu}{dt} = 0. \quad (58)$$

The drift-kinetic equation [Eq. (57)] is of interest by itself for applications in which longitudinal Landau damping is important and a crude modeling of transverse dynamics is sufficient. We note that (p_r, p_θ) variables are absent from any differentials. Equation (57) provides full kinetic descriptions of longitudinal motions but retains only lowest-order drift motions in the transverse direction. In this process, the number of independent variables in Vlasov's equation is reduced by two.

Closure in the Adiabatic Regime. The relation of drifts to hydrodynamic motions follows from

$$\begin{aligned}v_r &\simeq \langle U + (\mu\nu)^{1/2} \sin \Psi \rangle_{F_0} \\ &= \langle U(r, \theta, z, \dot{z}, t) \rangle_{F_0}, \quad (59)\end{aligned}$$

$$\begin{aligned}v_\theta &\simeq \left\langle \Omega r + \frac{\mu^{1/2} \nu^{3/2}}{2\Omega} \cos \Psi \right\rangle_{F_0} \\ &= \langle r\Omega(r, \theta, z, \dot{z}, t) \rangle_{F_0}, \quad (60)\end{aligned}$$

and

$$v_z \equiv \langle \dot{z} \rangle_F = \langle \dot{z} \rangle_{F_0}, \quad (61)$$

where the averages are

$$\langle h \rangle_F \equiv \int h F \frac{\nu^2}{4\Omega n} d\mu d\Psi d\dot{z} \quad (62)$$

and the number density of particles in this beam coordinate is

$$\begin{aligned}n &= \int F \frac{1}{r} dp_r dp_\theta d\dot{z} \\ &= \int F_0 \frac{\nu^2}{4\Omega} d\mu d\Psi d\dot{z}. \quad (63)\end{aligned}$$

If we examine these velocities, we find that they correspond to determining v_θ by radial force balance, while v_z is determined from v_θ by $\dot{p}_\theta \approx m d(\Omega r^2)/dt$.

Solving Eq. (57) can result in determination of the time evolution of hydrodynamic stresses. Simultaneous evolution of the hydrodynamic equations with this closure for stresses can result in a more accurate description of the transverse motions. We may relate, to lowest significant order in ϵ , the components P_{rr} , $P_{r\theta}$, and so forth, of the stress tensor to F_0 by

$$P_{rr} = n \langle \mu v \sin^2 \Psi \rangle_{F_0} = \frac{1}{2} n \langle \mu v \rangle_{F_0}, \quad (64)$$

$$P_{\theta\theta} = n \left\langle \mu \frac{v^3}{4\Omega^2} \cos^2 \Psi \right\rangle_{F_0}, \quad (65)$$

$$P_{r\theta} = P_{\theta r} = n \left\langle \mu \frac{v^2}{4\Omega} \sin^2 \Psi \right\rangle_{F_0} = 0, \quad (66)$$

$$P_{rz} = P_{zr} \\ = n \langle (\mu v)^{1/2} \sin \Psi (z - v_z) \rangle_{F_0} = 0, \quad (67)$$

$$P_{\theta z} = P_{z\theta} \\ = n \left\langle \mu \frac{1/2 v^{3/2}}{2\Omega} \cos \Psi (z - v_z) \right\rangle_{F_0} = 0, \quad (68)$$

and

$$P_{zz} = n \langle (z - v_z)^2 \rangle_{F_0}, \quad (69)$$

which are determined by integrals over F_0 . The continuity and momentum hydrodynamic equations of relevance can be written for each rotating beam component (where the sum of such components may have zero net angular momentum), such that

$$\frac{dn}{dt} + n \left[\frac{1}{r} \frac{\partial(rv_r)}{\partial r} + \frac{1}{r^2} \frac{\partial j}{\partial \theta} + \frac{\partial v_z}{\partial z} \right] = 0, \quad (70)$$

$$\frac{dv_r}{dt} = \frac{j^2}{r^3} + \frac{1}{m} f_r \\ - \frac{1}{mn} \left[\frac{\partial P_{rr}}{\partial r} + \frac{1}{r} (P_{rr} - P_{\theta\theta}) \right], \quad (71)$$

$$\frac{dj}{dt} = \frac{r}{m} f_\theta - \frac{1}{mn} \frac{\partial P_{\theta\theta}}{\partial \theta}, \quad j \equiv r v_\theta, \quad (72)$$

and

$$\frac{dv_z}{dt} = \frac{1}{m} f_z - \frac{1}{mn} \frac{\partial P_{zz}}{\partial z}, \quad (73)$$

where

$$\frac{d}{dt} \equiv \frac{\partial}{\partial t} + v_r \frac{\partial}{\partial r} + \frac{j}{r^2} \frac{\partial}{\partial \theta} + v_z \frac{\partial}{\partial z}, \quad (74)$$

with the body forces, f_i including the dominant radial component $q(E_r - v_z B_\theta/c)$ assumed earlier. The transformations of hydrodynamic equations to cylindrical coordinates is well known.^{52,53}

In Eq. (58), the frequency, ν , plays the role of $(q B_z/mc)$ in the Chew, Goldberger, and Low theory. A physically suggestive "field equation" for ν is

$$\frac{d}{dt} \left(\frac{\nu^2 r^2}{n |j|} \right) = 0, \quad (75)$$

which can be derived from Eqs. (52) through (54), (59) through (61), and (70), provided z does not vary much over local beam elements. For example, in the axisymmetric case, a heuristic derivation follows from considering a ring of particles with inner and outer radii r_1 , r_2 and a constant particle number $\propto nr(r_2 - r_1)$. Then constancy of angular momentum, $j(r, t) = r v_\theta$, implies

$$\text{constant} = |j(r_2, t) - j(r_1, t)| \\ \simeq \frac{\partial |j|}{\partial r} (r_2 - r_1) \propto \frac{1}{nr} \frac{\partial |j|}{\partial r} = \frac{\nu^2 r^2}{2n |j|}, \quad (76)$$

which is simply Eq. (75) on this ring. For nonaxisymmetric problems, a similar heuristic derivation follows from Kelvin's Theorem for vorticity. Due to its physical basis, Eq. (75) is preferred over Eq. (53), especially when extrapolating beyond the adiabatic regime indicated by the assumptions of Eq. (49). The adiabatic-regime description in terms of Eqs. (52), (54), (57), (58), and (63) through (75) is partly hydrodynamic and partly kinetic.

Closure in the Adiabatic Hydrodynamic Regime. A fully hydrodynamic description with closure in terms of generalized adiabatic equations of state is possible if, in addition, $|\dot{z} - \langle z \rangle|/\langle z \rangle = 0(\epsilon^2)$ and $(\partial/\partial z) = 0(\epsilon^2)$ (for example, due to negligible

Landau damping). We find under these conditions that

$$P_{\theta\theta} = \frac{r^4 v^2}{4j^2} P_{rr} \quad (77)$$

and

$$\frac{d}{dt} \left(\frac{P_{rr}}{nv} \right) = 0, \quad \frac{d}{dt} \left(\frac{P_{zz}}{n} \right) = 0. \quad (78)$$

Alternatively, heat-flux terms might also be omitted when collisions with the background plasma are strong enough to inhibit longitudinal heat flux but not strong enough to cause isotropic stresses.

For the best fully hydrodynamic model, the closure equations, Eqs. (75), (77), and (78), should be used with the continuity and momentum equations, Eqs. (70) through (74), which are exact velocity moments of Vlasov's equation. Thus we follow long traditions (e.g., turbulence theory) in making approximations only on the closure of a system of moment equations.

Discussion. The usefulness of our hydrodynamic and partially hydrodynamic approximations in applications to rotating and nonrotating beams will require more accurate numerical delineation of the adiabatic regime, summarized in the preceding paragraphs. In the case of beams with no net angular momentum, in which we superpose several of such beam components with different senses of rotation, numerical comparisons with particle codes are particularly needed. Some of these issues await further study.⁵⁴ For galactic dynamics, Berman and Mark⁵⁰ found that similar approximations are valid even for WKB analyses in which $(\partial/\partial r) \propto 0(\epsilon^{-1})$, while the derivations of such hydrodynamic models assumed implicitly $(\partial/\partial r) = O(1)$.

For applications to particle beams, more intriguing questions arise in situations where B_θ (or B_z , in cases with axial field) might temporarily vanish, reminiscent of magnetic-field-tearing instabilities for magnetohydrodynamics. In magnetohydrodynamics, supplements to this description usually require resistive magnetohydrodynamics or even local kinetic treatment of tearing modes.⁵⁵ Apart from the possibility of tearing magnetic field lines, which may require additional resistivity terms for correct description, the model of Eqs. (70) through (78), for example, does not neces-

sarily break down even though the radial force suddenly becomes small. This is because v^2 is approximately carried forward in time properly by Eq. (75) or (76), due to conservation of angular momentum (or Kelvin's Theorem for vorticity). For similar accuracy, in the use of Eqs. (57) and (70) through (74) in a combined fluid-kinetic description, it is preferable to define $\Omega = v_\theta/r = j/r^2$ in Eq. (54) and use Eq. (75) to determine ν . Within the adiabatic regime, $\Omega = (j/r^2) + 0(\epsilon)$.

Summary. We have outlined three different regimes of adiabatic approximations for simulations of pinch-beam phenomena. The drift-kinetic equation, Eq. (57), could be used in conjunction with Eqs. (52) through (54) in a description that provides accurate longitudinal dynamics but simplifies the description of transverse motions involving only the adiabatic drift velocities. A more accurate account of the transverse dynamics is provided by the mixed kinetic-hydrodynamic adiabatic regime. The adiabatic, fully hydrodynamic description can be obtained in the regime in which longitudinal Landau damping is negligible or when collisions are strong enough to inhibit longitudinal heat flux but not strong enough to result in isotropic stresses.

The formalism given here can also be naturally generalized to include relativistic effects or strong magnetic fields, such as are found in some ion sources⁴⁸ or collective plasma final-focus lenses; we have also envisioned some accelerator applications.⁴⁴

Authors: J. W.-K. Mark and S. S. Yu

Laser-Plasma Theory and Simulation

Introduction

The ability to model the coupling of laser light with targets is essential both for designing fusion targets and for choosing drivers for those targets. Our theoretical program on laser-plasma coupling is a three-front effort consisting of

- Theory and simulation of coupling processes.
- Plasma code development.

- Design and interpretation of coupling experiments.

In this article, we describe our progress in these areas during 1982.

We have developed a self-consistent kinetic theory of classical light absorption and electron-energy transport. Our model accounts for the non-Maxwellian velocity distributions that are established when inverse bremsstrahlung absorption is balanced by thermal electron transport in high-Z targets. Numerical calculations with this improved classical theory show both lower absorption and a thermal electron transport that is significantly reduced from naive, free-streaming estimates or the Spitzer-Härm theory. The self-consistently determined electron distributions also allow improved calculations of the atomic physics of the plasma.

Stimulated Raman scattering (SRS), as a potent source of hot electrons, has continued to be a focus of our attention. We have incorporated a simplified prescription for Raman-backscattered light into our target design code and used it to help design and interpret a number of experiments. In some of these experiments, about 10% of the laser light was observed to be Raman-scattered. We have carried out many computer simulations of SRS to improve our understanding of the hot-electron generation and the stabilizing effects of density gradients. We have also investigated another coupling process: heating by the oscillating component of the ponderomotive force of a light wave. This process, known as $\mathbf{J} \times \mathbf{B}$ heating, is relevant to high-intensity CO_2 experiments. We have discovered a new regime of intrinsically nonsteady stimulated Brillouin scattering (SBS); our SBS calculations explain the complex spectra of the scattered light observed in a number of experiments. Finally, we are carrying out calculations of thermal self-focusing, with particular attention to reactor targets.

Important new advances have been made in the development of implicit particle codes and hybrid particle-fluid codes for long-time-scale plasma simulation. Our direct method for solution of the implicit equation has been recently adopted by other groups at the Institute for Fusion Studies and at l'École Polytechnique. This year, we developed a direct-implicit electromagnetic algorithm in two dimensions and

a prototype code that merges a particle treatment of weakly collisional plasma with a fluid description of strongly collisional plasma. We are pursuing applications in both ICF and other programs.

A major part of the Plasma Physics group's resources was again devoted to the design and interpretation of coupling experiments. This year, we completed analysis of the numerous wavelength-scaling experiments previously carried out on the Argus laser (now retired). Our experiments confirmed the predicted—and very favorable—wavelength scaling of laser-target coupling, x-ray conversion efficiency, and energy transport; the experiments also enabled us to test and improve our LASNEX models. Specific new results include detailed comparisons of experiments with 1.06 μm , 0.53 μm , and 0.26 μm light. The results of this collaborative effort with the experimental group are reported in Section 5, "Fusion Experiments Program."

Finally, we have designed a variety of coupling experiments for the Novette laser. These include experiments to test our current scaling models and to extend the models to the much larger plasmas appropriate to Nova and reactor targets. New information on the plasma conditions and the mix of coupling processes will significantly enhance our ability to extrapolate our models to future targets.

Author: W. L. Kruer

Laser Absorption and Heat Transport by Non-Maxwell-Boltzmann Electron Distributions

We have developed a new kinetic theory of Coulomb collisional electrons in laser fusion plasmas, which promises better agreement with experimental results than the classical theory for absorption and heat transport. The failure of the classical theory^{56,57} to explain laboratory experience with high-Z laser fusion plasmas is well known.⁵⁸ In general, the classical theory overestimates both laser-beam absorption and electron heat transport into the irradiated target.

Workers in the field have come to employ ad hoc normalization of the classical theory implemented in computer simulation codes in order to replicate experimental data. In particular, it is common practice in

such calculations to limit the electron heat transport by employing the flux-limited expression $Q_{\parallel} = Q/(1 + |Q|/F_{\parallel} n_e T/m T)$ in place of Q , the classical heat flux. When Q is unsuitably large, heat is effectively transported according to $Q_{\parallel} \simeq F_{\parallel} n_e T/m T$. The phenomenological constant $F_{\parallel}$ (the flux-limit parameter, whose value is often $\sim 1/20$) has come to assume considerable significance in the field because it represents the rate and, thereby, the efficiency with which the energy to do the work of driving the target may be transported.

Naturally, there has been reconsideration⁵⁹ of the classical theory and improvements⁶⁰⁻⁶² that relax its limitations. One can readily show that classical heat-transport theory should not be expected to describe laser fusion plasmas,⁶³ because of the failure of the theory to self-consistently govern the higher-energy electrons (that transport the heat) of the assumed Maxwell-Boltzmann distribution. The same failure of theory is true with regard to the lower-energy electrons that absorb the energy from the laser beam in the classical absorption theory.

The essential physics of our new theory is that the electron collision times are short compared to the time scales of interest in laser fusion plasmas, whereas the electron mean free paths are comparable to the scale lengths. We show that in typical laser fusion plasmas, electron-ion collisions drive the distribution to near isotropy over a short distance, while thermalization in energy due to electron-electron collisions occurs over a significant distance. We present a new *kinetic* (reduced Fokker-Planck) theory of Coulomb collisional electrons in laser fusion plasmas, which provides this first improvement to the classical theory. We have obtained self-consistent results giving lower-than-classical absorption and heat transport in numerical calculations appropriate to laboratory and target-design systems. Our results are determined by non-Maxwell-Boltzmann distributions that are established when laser absorption is balanced by heat transport into denser, colder plasma.

We shall describe in some detail the heat-transport problem for laser fusion plasmas, but we omit a parallel description of the absorption problem. We emphasize the ele-

mentary and significant dependence of the performance of laser fusion systems on absorption, and the fact that absorption and transport are closely coupled in laser fusion plasmas. We do not address the problem of anomalous behavior in laser fusion plasmas, because our new kinetic theory naturally provides a new equilibrium, which itself must be investigated for any tendency toward instability.

A review of classical heat-transport theory permits us to indicate its failure in laser fusion plasmas and to develop motivation for the new kinetic theory. In one dimension of spatial inhomogeneity, x , the electron Fokker-Planck equation is solved by ordering electron-ion collisions large so that the distribution is dominantly isotropic

$$f = f_0(\epsilon, x, t) + \mu f_1(\epsilon, x, t) + P_2(\mu) f_2(\epsilon, x, t) + \dots \quad (79)$$

Here, $\epsilon = mv^2/2 - e\phi$ is the electron total energy in the slowly varying electrostatic potential $\phi(x, t)$ (the electric field is $-\partial\phi/\partial x$), and $\mu = \hat{x} \cdot \hat{v}$ is the cosine of the pitch-angle between the velocity vector and the direction of inhomogeneity (note that it is also the first Legendre polynomial). We have

$$\begin{aligned} \frac{\partial f}{\partial t} - \frac{\partial e\phi}{\partial t} \frac{\partial f}{\partial \epsilon} + \mu v \frac{\partial f}{\partial x} - \frac{\partial}{\partial x} \left(\frac{1 - \mu^2}{mv} \right) \frac{\partial f}{\partial \mu} \\ = \frac{v}{\lambda_{ei}} \frac{\partial}{\partial \mu} \left[(1 - \mu^2) \frac{\partial f}{\partial \mu} \right] + \frac{\partial f}{\partial t_{ee}} \end{aligned} \quad (80)$$

Here, $v = [2(\epsilon + e\phi)/m]^{1/2}$, and $\lambda_{ei} = m^2 v^4 / 2\pi e^4 Z^2 N(\ln \Lambda_{ei})$ is the mean free path for electron-ion momentum scattering. The strength of anisotropy, f_1 , obeys

$$v \frac{\partial f_0}{\partial x} = - \frac{2v}{\lambda_{ei}} f_1 + \dots \quad (81)$$

Equation (81) is the coefficient of the first Legendre polynomial in the decomposition of the transport equation.⁵⁷ The dots indicate several neglected terms. The electron-electron collision terms in μ and in ϵ are small by order $1/Z$, the inverse of the plasma ionization state, compared with the electron-ion scatter retained. The term in f_2 , which represents higher-order pitch-angle

structure, is taken to be small in comparison with the term in f_0 . Finally, the rates of temporal relaxation of the distribution and potential are also assumed to be slow compared with the electron-ion collision rate.

Classical transport is obtained by employing the Maxwell-Boltzmann distribution,

$$\frac{n}{(2\pi T/m)^{3/2}} \exp\left[-\frac{(\epsilon + e\phi)}{T}\right], \quad (82)$$

for f_0 in Eq. (81) to evaluate the heat flux associated with the temperature gradient. This is appropriate whenever electron-electron collisions are dominant in determining f_0 (which obeys an equation to be displayed later). The electrostatic potential is determined by quasi-neutrality, $n = NZ$, through the condition that the flux in total electron number vanishes

$$\Gamma = \frac{4\pi}{3} \int_{-\infty}^{\infty} \frac{d\epsilon v}{m} v \left(f_1 = -\frac{\lambda_{ei}}{2} \frac{\partial f_0}{\partial x} \right) = 0; \quad (83)$$

$$\frac{\partial e\phi}{\partial x} = \frac{5}{2} \frac{\partial T}{\partial x} + T \frac{\partial (\ln n)}{\partial x}.$$

Employing this potential, we may write

$$f_1 = \frac{\sqrt{2\pi}}{32} \frac{Q}{n\sqrt{T/m} T} \left(\frac{mv^2}{2T} \right)^2 \left(\frac{mv^2}{2T} - 4 \right) f_0. \quad (84)$$

The heat flux associated with electrons of energy less than ϵ^* is

$$Q^* = \frac{4\pi}{3} \int_{-\infty}^{\epsilon^*} \frac{d\epsilon v}{m} v \frac{mv^2}{2} f_1 \quad (85)$$

and the classical heat flux is

$$Q = \lim_{\epsilon^* \rightarrow \infty} Q^* = -\frac{64}{\sqrt{2\pi}} n\sqrt{T/m} \lambda_{mfp} \frac{\partial T}{\partial x}. \quad (86)$$

Here $\lambda_{mfp} = (T/mv^2)^2 \lambda_{ei}$. We observe that the lower-energy electrons actually carry heat up the temperature gradient as they participate in the neutralizing return current; thus, we consider the higher-energy electrons, which contribute the flux of interest. We shall adopt $|f_1(\epsilon^*)/3| < f_0(\epsilon^*)$, where $Q^*(\epsilon^*) = 0.9Q$, as necessary for the self-consistent validity of the classical heat-

transport theory.⁶³ We find $\epsilon^* = 11.4T$ and $|f_1^*/3| = 25(|Q|n\sqrt{T/m} T)/f_0$. Thus, $|f_1^*/3| < f_0$ for $|Q|/n\sqrt{T/m} T < 0.04$, or $\lambda_{mfp}/|\partial(\ln T)/\partial x| < 0.002$. These inequalities are ubiquitously violated in laser fusion plasmas, so that the classical heat-transport theory fails to be self-consistently valid in its application. For example, for parameters of $Q = 10^{14}$ W/cm², $n = 10^{21}$ /cm³, and $T = 10^3$ eV, we find that $|f_1^*/3| \approx 13f_0$. For these parameters, the theory has failed for $\epsilon \ll \epsilon^*$, and we may expect large changes in the heat flux from an improved theory.

Previous efforts to improve electron heat transport theory for laser fusion plasmas essentially involve improving Eq. (81). (See Ref. 61.) It turns out that for modeling laser fusion plasmas, Eq. (81) is, in fact, adequate to describe f_1 ; our new kinetic theory involves the calculation of the isotropic but not necessarily Maxwell-Boltzmann distribution, f_0 , which obeys

$$\begin{aligned} \frac{\partial}{\partial t}(vf_0) - \frac{\partial}{\partial t}(e\phi) \frac{\partial}{\partial \epsilon}(vf_0) \\ + \frac{\partial}{\partial x} \left(\frac{v^2 f_1}{3} = -\frac{v^2 \lambda_{ei}}{6} \frac{\partial f_0}{\partial x} \right) = \gamma m (\ln \Lambda_{ee}) \\ \times \frac{\partial}{\partial \epsilon} \left(f_0 4\pi \int_{-\infty}^{\infty} \frac{d\epsilon' v'}{m} f_0 \right) \\ + \frac{\partial f_0}{\partial \epsilon} 4\pi \int_{-\infty}^{\infty} \frac{d\epsilon' v'}{m} \int_{-\infty}^{\infty} d\epsilon'' f_0 \\ + v_1^2 \frac{m^2}{6} \gamma Z^2 N (\ln \Lambda_{ei}) \\ \times \frac{\partial}{\partial \epsilon} \left[\frac{1}{1 + \left(\frac{2v}{\lambda_{ei} \omega_L} \right)^2} \frac{\partial f_0}{\partial \epsilon} \right]. \end{aligned} \quad (87)$$

Equation (87) is the coefficient of the zero-order Legendre polynomial in the decomposition of the transport equation.⁵⁷ First on the right-hand side of Eq. (87) is the electron-electron collision term, which fails to be dominant in laser fusion plasmas, so that f_0 is not necessarily Maxwell-Boltzmann. Here $\gamma = 4\pi e^4/m^2$. Second is the source of energy to the plasma, the laser-heating term recently investigated by Langdon.⁶² In the regime of interest, the coefficient of $\partial f_0/\partial \epsilon$ in this term is effectively

unity, so that it acts to heat the plasma by diffusing electrons away from the region of low kinetic energy of phase space. The equation for the laser beam, $(v_1)^2$, will be given directly.

It is also necessary that electron-electron collisions dominate spatial transport at ϵ^* . When the Maxwell-Boltzmann distribution is used in comparing the associated terms in Eq. (87), the condition that is actually to be applied to high-Z plasmas is obtained by

$$\frac{|Q|}{n \sqrt{T/m} T} < \frac{0.07}{[Z(\ln \Lambda_{ei})/(\ln \Lambda_{ee})]^{1/2}} \quad (88)$$

or

$$\lambda_{mf} \partial(\ln T)/\partial x < \frac{0.003}{[Z(\ln \Lambda_{ei})/(\ln \Lambda_{ee})]^{1/2}} \quad (89)$$

We have developed a computer code, which solves Eq. (87) numerically by fully implicit splitting techniques commonly used for two-dimensional, initial-value diffusion calculations.⁶⁴ The finite differencing plus iteration of the nonlinearities preserves the conservation laws associated with the transport equation. Equation (87) is coupled to WKB equations for the incident and returning (i.e., $\pm$) laser beams

$$\pm \frac{\partial}{\partial x} (v_x v_x^2) = -\frac{1}{3} Y \frac{Z^2 N}{n_c} \times (\ln \Lambda_{ei}) f_0 (\epsilon = -\epsilon\phi; v = 0) v_{\pm}^2 \quad (90)$$

Here, $v_{\pm} = eE_{\pm}/m\omega_L$, with $E_{\pm}$ the electric fields of the laser beams; in Eq. (87), $v_{\pm}^2 = v_{\pm}^2 + v_{\pm}^2$. The group velocity of light in plasma is $v_g = c(1 - n/n_c)^{1/2}$, where $4\pi e^2 n_c/m = \omega_L^2$ defines the laser critical density, n_c . The intensities of the laser beams are given by $I_{\pm} = v_g n_m v_{\pm}^2/2$. Note that the kinetically calculated laser opacity depends upon the value of the distribution function at zero velocity.

The electrostatic potential is determined by quasi-neutrality, as in the classical theory, such that

$$n = 4\pi \int_{-\infty}^{\infty} \frac{d\epsilon v}{v\phi} f_0 = ZN \quad (91)$$

In particular, the variations with $\epsilon\phi$ of Eqs. (87) and (91) (or the equivalent zero-current condition already given) are employed iteratively to find improvements in the potential and distribution that will cause the governing system of equations to be obeyed simultaneously.

Table 3-4 is an overview of three calculations of laser-heated plasmas according to Eqs. (87), (90), and (91). Common conditions for the calculations are shown in Fig. 3-28. In the expression for Λ_{ei} and in Eqs. (87) and (90), $Z^2 N$ is replaced by $Z(Z+1)N$. The Coulomb logarithms are $\ln \Lambda_{ei} = 2.5$ and $\ln \Lambda_{ee} = 5$ throughout. In each calculation, the ionization state, Z , is constant while a constant-intensity laser beam irradiates the plasma until the heat transport out of the absorbing coronal region into the colder, overdense region equals the

Table 3-4. Comparison of calculations of electron behavior in laser fusion plasma based on classical Coulomb collision theory, a mixed theory, and our new kinetic theory.

	Fraction absorbed	Corona temp (keV)	$\frac{Q_{\max}}{n\sqrt{T/m} T}$	f_1/f_0 at $Q_{\max}$	$\frac{Q_{\max, \text{kin}}}{Q_{\text{class}} T_{\text{kin}}}$	F_{η} at $Q_{\max}$
$Z = 10, I = 10^{14} \text{ W/cm}^2$						
$t_{\text{data}} = 125 \times 10^{-12} \text{ s}$:						
Classical	0.63	1.28	0.19	14	—	—
Mixed	—	1.16	0.16	12	—	—
Kinetic	0.46	1.33	0.14	1.4	0.34	0.21
$Z = 10, I = 10^{15} \text{ W/cm}^2$						
$t_{\text{data}} = 41 \times 10^{-12} \text{ s}$:						
Classical	0.36	2.20	0.48	35	—	—
Mixed	—	1.76	0.36	27	—	—
Kinetic	0.19	2.28	0.26	1.4	0.17	0.32
$Z = 40, I = 10^{14} \text{ W/cm}^2$						
$t_{\text{data}} = 168 \times 10^{-12} \text{ s}$:						
Classical	0.84	2.05	0.13	9.4	—	—
Mixed	—	1.85	0.11	8.3	—	—
Kinetic	0.63	2.20	0.09	0.70	0.31	0.13

(slowly varying) absorbed laser intensity. At this time, the coronal plasma has reached a quasi-steady state, which is taken to be the conclusion of our absorption and transport calculations. Beyond this state, the hydrodynamic ablation energy sink should be included. However, the coronal plasma condition may not be strongly influenced by this effect, except as it determines the density scale length. We suggest that the coronal plasma may reach the condition found here on a time scale as fast as the processes that determine the scale length. Thus, our present calculations may reveal the long-time-scale coronal-plasma condition.

We can most readily put the new results into perspective by comparing them with parallel classical calculations. Inspection of Table 3-4 shows that the problems cannot be reliably calculated classically. The classical results are seen to be significantly in error, finding excessive laser absorption and heat transport into the overdense plasma. Mixed calculations, in which the kinetically absorbed energy serves as the source in a classical transport calculation, also find

Fig. 3-28. Common ▲ conditions for the three calculations whose results are summarized in Table 3-4.

Fig. 3-29. Plasma fluid properties at the conclusion of the calculation for $Z = 10$, $I = 10^{14} \text{ W/cm}^2$.

excessive heat transport when compared with the kinetic calculations. The new kinetic calculations are seen to be essentially self-consistent. They give heat fluxes that are significantly reduced from naive, free-streaming estimates and from the classical heat flux associated with the kinetic temperature profile (our kinetically calculated temperature is two-thirds of the energy density divided by the density).

We now examine in detail the plasma conditions at the end of the calculation for $Z = 10$ and $I = 10^{14}$ W/cm² summarized in Table 3-4. Figure 3-29 shows the macroscopic plasma properties. The peak of the laser energy deposition is at critical density. The somewhat greater penetration of the classical and mixed temperature fronts is evident, resulting in comparable or lower coronal temperatures for greater or equal absorbed energy (the exponentially increasing density cause, some masking of the penetration of the energy density). Near the peak of the heat flow, just inside critical density, the classical heat flow associated with the kinetic temperature profile is con-

siderably greater than the actual kinetic heat flow. This is due to the fractionally higher temperature, which enters as $T^{7/2}$ in the classical heat flow, and also to the steeper temperature gradient there (compare the mixed calculation that carries the same heat flux by virtue of being driven by, and balancing, the same laser absorption). Note that the heat flow is an even smaller fraction of the characteristic quantity $n\sqrt{T}/m T$.

In Fig. 3-30, the kinetically calculated electron distributions at various positions in the plasma are shown together with Maxwell-Boltzmann distributions with the same temperatures. The subcritical plasma is characterized by the distribution (labeled M-B) shown as a function of velocity and energy at position B. The lack of dominance of electron-electron collisions over laser heating and spatial transport is clearly seen in the depletion of the distribution at low energy (which is responsible for the reduced laser absorption) and also at high energy (which is responsible for the reduced heat transport). The flow in phase space is at first upward in energy due to laser

Fig. 3-30. Electron distribution functions at positions B through F labeled on the temperature profile in Fig. 3-29.

heating, then into the colder plasma due to spatial transport. The electric field provides the return flow of lower-energy electrons from the transport-heated plasma to the laser-heated plasma. The quasi-steady distributions in the coronal plasma result from the balance between laser heating at lower energy and spatial transport at higher energy, with self-collisions mediating but not dominating. It is this balance between heating and transport that fixes their common value found here.

Previous work on absorption concerned similar but nonsteady distributions, which evolved under heating and self-collisions only⁶² and did not address this physics. We observe that laser heating diffuses electrons away from lower energy more rapidly than self-collisions can establish equilibrium. The reduced opacity to the laser beam results from the depopulation of lower-energy electrons relative to the Maxwell-Boltzmann distribution. In the laser-heated plasma we see that the thermalization mean free path of the higher-energy electrons is long enough that they transport spatially into the colder plasma before self-collisions can establish a Maxwell-Boltzmann tail. It is intuitively satisfying that the heat flow associated with the resulting truncated distributions is smaller than that associated with the Maxwell-Boltzmann distribution at the same temperature. Above critical density, the distribution relaxes to Maxwell-Boltzmann at low energy more rapidly than at high energy. The tail of the coronal distribution penetrates as preheat into the dense, cold plasma as evidenced by the strongly two-component distribution shown

at position *F*. Except in the preheat region, we have underfull tails on the distribution, and not the overfull ones of previous calculations in which hot Maxwell-Boltzmann plasma heats colder plasma.⁶¹

In Fig. 3-31 and the related Table 3-5, we display the self-consistency of the various calculations. At position *C*, 90% of the heat flux is carried by electrons with $v < v^* \approx 5.3 \times 10^9 \text{ cm/s}$, where $|f_1|/f_0 < |f_1^*|/f_0^* \approx 1.5$. Near the peak of the heat flux, especially, we see that the kinetic theory represents a vast improvement over the classical theory. It is not surprising that a considerable modification of the heat flux occurs near the peak. In the denser, colder plasma, contributions to the heat flux occur further and further out on the tail of the distribution, where the electrons are effectively streaming in from the corona.

In Fig. 3-32 and the related Table 3-6, we show the flux limitation implied by the kinetic temperature and heat-flux profiles. We have evaluated the spatially varying parameter, F_D , that would be required in a flux-limited calculation to simultaneously match the kinetic temperature and heat flux near the peak of the heat flux. A local flux limit of about 1/5 is implied. As indicated already, this is due to the higher kinetically

Position	Classical	Kinetic	Mixed
A	2.6	0.25	2.3
B	5.3	0.50	4.5
C	9.8	1.0	8.4
$Q_{\max}$	14.0	1.4	12.0
D	12.0	1.5	11.0
E	9.4	1.8	8.3
F	7.7	2.0	7.4
G	0	2.5	0

Table 3-5. Values of f_1/f_0 at $0.4Q$ for $Z = 10$, $I = 10^{14} \text{ W/cm}^2$.

Fig. 3-31. Self-consistency of the heat-transport calculation for $Z = 10$, $I = 10^{14} \text{ W/cm}^2$.

Position	$Q/\sqrt{T/m} \ T _{kin}$	$Q_{kin}/Q_{class} (T_{kin})$	F_{fl}
C	0.085	(0.60)	(0.21)
CD	0.122	0.20	0.15
$Q_{max,kin}$	0.138	0.34	0.21
D	0.135	0.50	0.27
DE	0.124	0.74	0.48
E	0.114	0.90	1.14

Table 3-6. Calculation of F_{fl} for $Z = 10$, $I = 10^{14}$ W/cm².

Fig. 3-32. Determination of the flux-limit parameter required to match flux-limited heat transport evaluated with the kinetic-temperature and heat-flux profiles at the conclusion of the calculation for $Z = 10$, $I = 10^{14}$ W/cm². (See Table 3-6.)

calculated temperature and its steeper gradient. It is intriguing that this is the same region where the flux limit acts in typically normalized simulations, and that the value we calculate accounts for a substantial portion of the parameters that have been employed in the past to simulate difficult experimental data. In the coronal plasma, the kinetic heat flux is small and is seen to be up the temperature gradient, while in the overdense plasma the preheat flux exceeds that associated with classical conduction in the weak temperature gradient there.

The other two calculations summarized in Fig. 3-28 and Table 3-4 exhibit data that are qualitatively similar to that of Figs. 3-29 through 3-32 and Tables 3-5 and 3-6. It is remarkable that the high- Z , low-intensity run is, in general, the most well behaved calculation in all respects and exhibits the greatest flux limitation. The low- Z , high-intensity calculation is the least well behaved and exhibits the least flux limitation. Both these calculations also exhibit the reduced absorption associated with the kinetic laser heating and opacity modeling.

From our calculations we find that increasing the laser intensity or the plasma Z

enhances the non-Maxwell-Boltzmann nature of laser-heated plasmas. Clearly, these non-Maxwell-Boltzmann effects will persist at somewhat longer density scale lengths (laser pulse lengths) and somewhat shorter laser wavelengths.

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Stimulated Raman Scattering in Inhomogeneous Plasmas and by Large-Bandwidth Lasers

As larger targets are irradiated with longer and more energetic pulses of laser light, the size of the underdense plasma increases. In large regions of underdense plasma, stimulated Raman scattering (SRS) can become an important coupling process. The SRS instability corresponds to the decay of an incident light wave into a scattered light wave plus an electron plasma wave (epw); as can be seen from the frequency-matching conditions, this process takes place in the underdense plasma having density less than or equal to quarter-critical density ($1/4n_c$). Since the unstable plasma wave can have a high phase velocity (on the order of the velocity of light), very energetic electrons can be generated. SRS is thus a significant pre-heat concern and has been a focus of our attention. Because gradients in plasma density influence SRS, we report some recent results on hot-electron generation by Raman scattering in inhomogeneous plasmas. A more extensive discussion of our work can be found in Ref. 65.

To examine the nonlinear behavior of Raman scatter, we used a 1.5-dimensional particle code (OREMP) to follow the propagation of light through an inhomogeneous slab of underdense plasma. The code follows the plasma evolution until a quasi-steady state is attained, and monitors the scattered light, absorption, and hot-electron temperatures.

Figure 3-33 presents the code's simulation of Raman absorption in an inhomogeneous plasma. The plasma density initially was a ramp from $n/n_c = 0.01$ to $n/n_c = 0.32$ in 87 wavelengths; $T_e = 5.5$, $T_e/T_i = 30$, and the incident intensity $I_{\mu m}^2 = 3 \times 10^{14}$ to

$2 \times 10^{16} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$. In some runs, 60% of the incident light reaching the high-density boundary condition was reflected, emulating a critical-density surface. The plasma-boundary condition on the high-density side was a numerical surface that allowed escaping particles to be replaced by cold particles at T_e or T_i . Some room was allowed for expansion on the laser side. As expected, the absorption due to SRS drops near the minimum gradient threshold intensity [$I_T \sim 10^{17}/(L_{\mu\text{m}}^{4/3} \lambda_{\mu\text{m}}^{2/3}) \sim 1.4 \times 10^{14} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$ for these parameters, where T refers to the threshold] but rapidly proceeds to a sizable level well above that intensity.

In a simulation using the plasma conditions, above, except with fixed ions, SRS still absorbs 20% of the light for $\lambda_{\mu\text{m}}^2 = 8 \times 10^{14} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$, which is about a factor of 2 over the density-gradient threshold. In general, ion dynamics reduce the level of SRS either due to competition with other processes such as light loss from stimulated Brillouin scattering (SBS) or laser absorption on the short-wavelength Brillouin waves (with low density and/or high T_e increasing $k\lambda_D$),⁶⁶⁻⁶⁸ or due to the steepening of the density profile. A simulation with the same initial density profile as above ($\lambda_{\mu\text{m}}^2 = 10^{15} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$ and $T_e = 1 \text{ keV}$) showed essentially no Raman scattering due to quite vigorous Brillouin scattering.

An example of the profile steepening in a simulation with moderate SRS is shown in Fig. 3-34. The density profile is steepened due to the ponderomotive force of the epw electrostatic field, the transverse field of the scattered light, and the heated electrons escaping and leaving a positive charge that accelerates some of the ions away. At high intensities, the pump can form holes in the plasma; this effect has also been seen in two dimensions⁶⁹ and is analogous to bubbles at the critical density.⁷⁰ Since it depends on the polarization, bubble formation is inherently a three-dimensional phenomenon and is only partially resolved in one-dimensional simulations. The Raman absorption, heated-electron temperatures,⁶⁶ and ω_s are characteristic of the roughly constant density shelf and the steep gradient near $1/4n_c$. The actual density profile also depends on the $2\omega_{pe}$ (Ref. 68) and Raman sidescatter instabilities as well as the spheri-

cal divergence and flow of the plasma. There is indirect experimental evidence from spectra that the shelf is at a much lower density. Consequently, Fig. 3-34 can only be considered schematic, showing that such steepening can occur from Raman backscatter in isolation. In detail, of course, the Raman absorption and heated-electron temperatures, T_{hot} will depend on the actual profile.

Fig. 3-33. Fraction of light absorbed from 1.5-dimensional kinetic simulations of Raman backscatter in a density gradient.

Fig. 3-34. Density profile from a 1.5-dimensional simulation at time = 1500 laser cycles.

Figure 2 of Ref. 65 presented T_{hot} vs density for $T_e = 1$ and 10 keV in a long, homogeneous plasma. Our Fig. 3-35 presents T_{hot} from our simulations of Raman backscatter in a density gradient. Except for the run at $\lambda_{\mu\text{m}}^2 = 2 \times 10^{16}$ W/cm², 60% of the incident light was reflected from an artificial critical surface. This reflected light can also drive Raman in the backward direction and may account for the difference in the temperatures measured by averaging over space and through the high-density boundary. At these low intensities and absorptions, T_{hot} is inherently difficult to measure; the high and low temperatures in Fig. 3-35 may be considered as error bars.

Fig. 3-35. The heated electron temperatures in a density gradient.

Fig. 3-36. Heated electron temperatures, T_{hot} , and fractional absorption from Raman backscatter for two gradient lengths.

Figure 3-36 presents more simulation data for other gradient lengths as noted. The simulations for these points did not include the artificial critical surface described above, having instead a transmitting electromagnetic boundary condition. The predicted gradient threshold intensities⁷¹ are $\lambda_{\mu\text{m}}^2 \sim 1.5 \times 10^{15}$ W·μm²/cm² for $L/\lambda_{\mu\text{m}} = 23$ and $\lambda_{\mu\text{m}}^2 \sim 1.5 \times 10^{14}$ W·μm²/cm² for $L/\lambda_{\mu\text{m}} = 139$, in reasonable agreement with the data. Note that T_{hot} again rapidly rises to a value characteristic of SRS at densities of about $1/5n_c$. It is interesting to note that Raman scatter near $1/4n_c$ has a lower gradient-threshold intensity than the $2\omega_{\text{pe}}$ instability, provided T_e (keV) $\geq 20 (\lambda_{\mu\text{m}}/L)^{1/3}$.

Effects of Bandwidth. Finally, let us note that the gain of SRS can be reduced by the use of laser light with a large bandwidth; experiment⁷² and simulations⁷³ have shown that bandwidth can also reduce SBS, an instability similar to SRS. A bandwidth of $\Delta\omega \simeq 10$ to 20% can be made up with about 10 different lines, as suggested by Eimerl et al.⁷⁴ Reference 66 showed that use of a large bandwidth can drive SRS below its gradient threshold in an inhomogeneous plasma.

To further explore the effect of a large bandwidth on SRS, we carried out simulations using a plasma with uniform density. As Fig. 3-37 shows, the fraction of the light absorbed by SRS is significantly reduced when the bandwidth is quite large. If we plot the inferred SRS gain vs $\gamma/\Delta\omega$, we find that the gain is proportional to $\gamma/\Delta\omega$ for $\Delta\omega/\gamma > 10$; here $\Delta\omega$ is the total bandwidth, and γ is the instability growth rate. This result is consistent with the predictions of Thomson and Karush.⁷⁵

Since the Raman instability is resonant for densities less than or near $1/4n_c$ lines for which $n > 1/4n_c$ for that particular line interact only weakly. The problem with discrete lines is that if ω_p equals the "difference frequency" between any of the lines, the beam frequency can directly drive SRS (Ref. 76) or enhanced filamentation⁷⁷ at that density. For a bandwidth in which the maximum difference between the lines is 20%, for example, the resonant density would be $n/n_c = 0.04$. Since the density here is fairly low, the backscatter epw phase velocity is low, and rather benign heating is expected. Forward scatter may be

considerably enhanced with associated hot electrons, however.

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J×B Heating by Very Intense Laser Light

The absorption of very intense light is a problem of great interest to anyone studying ICF applications of lasers, particularly CO₂ lasers. There is, of course, the well-known mechanism of resonance absorption,⁷⁸ which can be strongly enhanced by both cratering⁷⁹ and rippling⁸⁰ of the critical density surface. We have recently explored another little-known mechanism, which can play a significant role at the very high intensities commonly used in CO₂ experiments. This so-called J×B heating⁸¹ is due to the electrostatic field driven by the oscillating component of the ponderomotive force of the laser light. Here we present a model calculation to show that the driven electrostatic field is sizable for high-intensity light. We then present computer simulations of the J×B heating, which demonstrate significant absorption of very intense light into electrons with a modest temperature.

Model Calculation. We begin with a model calculation to illustrate the physics of this coupling mechanism. Consider a linearly polarized light wave with electric amplitude $\mathbf{E} = E_L(x)(\sin \omega_0 t)\hat{\mathbf{y}}$. The ponderomotive force, $\mathbf{f}_p$, exerted by this field is readily derived in the usual way by

$$\mathbf{f}_p = -\frac{1}{4} \frac{\partial}{\partial x} v_L^2(x) (1 + \cos 2\omega_0 t) \hat{\mathbf{x}}, \quad (92)$$

where $v_L = eE_L/m\omega_0$, and, for simplicity, we have neglected relativistic effects. The time-averaged part produces steepening of the density profile, and the oscillating component can lead to heating. In effect, the oscillating ponderomotive force corresponds to driving the plasma with an electrostatic field E_d , where

$$\frac{eE_d}{m} = \frac{1}{4} \frac{\partial}{\partial x} v_L^2(x) (\cos 2\omega_0 t). \quad (93)$$

Since the magnitude of E_d is proportional to $v_L \partial v_L / \partial x$, significant absorption is expected only for very intense light.

To estimate the size of the driver field, we consider an intense, planar light wave normally incident on an inhomogeneous slab of plasma and use the model of Lee et al.⁸² to compute $E_L(x)$ in the strongly steepened light-plasma interface. For intense fields, the density variation is large, and a numerical solution of the model equations is needed. Figure 3-38 shows a typical result for the spatial variation of the amplitude of the transverse electric field and the plasma density in the interface. For this example, $(v_{FS}/v_e)^2 = 10$, where $v_{FS} = eE_{FS}/m\omega_0$, E_{FS} denotes the free-space value of E_L , and v_e is the electron thermal

Fig. 3-37. SRS absorption vs bandwidth.

Fig. 3-38. Spatial variation of the transverse electric field ($v_L = eE_L/m\omega_0$) and the plasma density in the light-plasma interface.

velocity. As expected from the pressure balance, the plasma is steepened by the intense, reflecting light wave to a density of $n_c(v_{FS}/v_e)^2$, where n_c is the critical density. Both the transverse field and the density vary quite rapidly in the interface.

Given $v_i(x)$, we can evaluate the magnitude of the driver field due to the oscillating ponderomotive force. For example, at the location where the driver field would be most enhanced by the plasma response ($n \approx 4n_c$), v_i and that

$$\frac{eE_d}{m} \approx 0.6 \left(\frac{v_{FS}}{v_e} \right)^{1.8} \frac{v_e^2 \omega_0}{c} \quad (94)$$

Although relativistic effects should be included for more quantitative estimates, the magnitude of this field is clearly significant when the incident light is very intense. The oscillating electrostatic field has a strong spatial variation in the light-plasma interface and will heat electrons that interact nonadiabatically with it.

Simulations. To compute the absorption and electron heating, we have carried out computer simulations using a 1.5-dimensional particle code with electromagnetic fields and relativistic particle dynamics. In our simulations, we consider the propagation of an intense, planar light

wave from vacuum onto a plasma slab. The plasma begins as an overdense slab with density equal to the upper density expected in the final steepened state; i.e., $n/n_c = (v_{FS}/v_e)^2$. The reflected light waves are allowed to escape into vacuum at the left boundary, and heated electrons are replaced with incoming thermal electrons at the right boundary. The evolution is followed until a quasi-steady state is attained.

Our simulations demonstrate significant heating when the light is very intense. As illustrated by the plot of electron phase space in Fig. 3-39, electrons are accelerated and then beamed into the plasma by the oscillating ponderomotive force. The effective temperature of the heated electrons is relatively modest, as shown by the plot of the electron distribution function in Fig. 3-40. In this example, the product of incident intensity and wavelength squared is $10^{18} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$ (i.e., $I = 10^{16} \text{ W}/\text{cm}^2$ for CO_2 light), and the initial electron temperature of the dense plasma slab is 4 keV. The time-averaged absorption is about 11%. Note that the heated temperature is about 50 keV, whereas the heated temperature due to resonance absorption would be of order 500 keV for these parameters.

Table 3-7 shows the averaged absorption measured in a number of simulations with different incident intensities and background electron temperatures. As expected, the absorption is strongly dependent on intensity, ranging from 1 or 2% at $I\lambda_0^2 = 10^{17} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$ to between 10 and 15% at $I\lambda_0^2 = 10^{18} \text{ W} \cdot \mu\text{m}^2/\text{cm}^2$. A weak dependence of the absorption on background electron temperature is also evident.

In summary, we have demonstrated a novel mechanism for the absorption of very intense laser light. The oscillating component of the ponderomotive force generates an electrostatic field, which leads to electron heating. A model calculation shows that the electrostatic field is sizable for very intense light. Computer simulations with a relativistic particle code demonstrate significant absorption into electrons with a relatively modest temperature. Of course, this process is supplemented by the well-known process of resonance absorption on a deeply cratered and rippled critical-density surface.

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Fig. 3-39. Electron phase space from a computer simulation of $\mathbf{J} \times \mathbf{B}$ heating.

Nonsteady Stimulated Brillouin Scatter Induced by a Reflective Critical Surface

Stimulated Brillouin scatter⁸³ (SBS) is of considerable interest in laser fusion because it is a mechanism by which incident light may be scattered from the target before significant absorption can occur.⁸⁴ In conventional analyses of SBS, outgoing wave conditions are applied to the scattered beam at the material boundaries or at infinity.⁸³ In actuality, however, physical systems often possess a reflective internal boundary, such as the critical-density surface in a laser fusion plasma.

In this article, we describe the results of studies showing that an internal boundary that is reflective to light waves can significantly alter the usual SBS physics. In the linear regime, we find large changes in the absolute instability threshold, growth rate, and scattered-light frequency shift. In the nonlinear regime, where the scattered light wave is comparable in magnitude to the incident pump wave (although the acoustic wave may remain linear), numerical calculations show that the scattered light exhibits a rich frequency spectrum, including shifts corresponding to both multiple and fractional harmonics of the fundamental acoustic-wave frequency.

As the pump intensity, acoustic-wave dissipation, and system length are varied, the time-asymptotic behavior of the system (as evidenced by the intensity of the scattered light, for example) displays transitions among simple oscillatory steady states, complex limit cycles, and chaos of no discernible periodicity. For parameters typical of high-intensity laser-plasma-interaction experiments, SBS often falls in the chaotic regime. As we will show, calculations of highly structured scattered-light spectra in this regime compare favorably with experimental data.⁸⁵

Our new results may be understood by considering that a reflective internal boundary alters both the equilibrium in which SBS grows and the wave-coupling process of the instability itself. The reflected incident wave acts as a second, counter-propagating pump wave; the incident and reflected waves beat together, supporting an

"equilibrium" perturbation of the medium that couples them. If the medium and reflective boundary have no relative velocity, this coupling results in a simple phase shift. When the relative motion is nearly sonic, large energy transfers from the incident to the reflected pump wave can occur, significantly reducing transmission to the internal boundary.⁸⁶

In the usual analytic case of SBS without a reflective boundary, there is a single pump wave; two pairs of acoustic and backscattered light waves satisfy the frequency-matching and wave-number-matching conditions required for three-wave coupling with the pump wave. In one pair, an acoustic wave propagating parallel to the pump wave couples with a frequency-downshifted (Stokes) backscattered light wave and is enhanced. In the other pair, an acoustic wave propagating antiparallel to the pump couples with upshifted (anti-Stokes) backscattered light and is evanescent. These acoustic waves have a wavelength one-half that of the light waves (the coupling of very-long-wavelength acoustic waves and forward-scattered light waves is neglected here).

In the presence of a reflective internal boundary, on the other hand, there are two pump waves, so that each acoustic wave

T_e (keV)	I_p^2 ($W \cdot \mu m^2/cm^2$)	f_{abs}
4	10^{17}	0.015
4	3×10^{17}	0.06
4	10^{18}	0.11
12	3×10^{17}	0.05
12	10^{18}	0.17
40	10^{18}	0.13

Table 3-7. Time-averaged absorption (f_{abs}) measured in simulations with different incident intensities and background electron temperatures.

Fig. 3-40. Plot of the electron distribution in a computer simulation of $J \times B$ heating, with same parameters as Fig. 3-39.

couples to both the Stokes and the anti-Stokes light waves. For example, the acoustic wave propagating parallel to the incident pump wave couples to the incident pump wave and its associated Stokes wave, while simultaneously coupling to the reflected pump wave and its associated anti-Stokes wave. The former coupling is the usual destabilizing influence and represents an energy source for the acoustic wave; the latter coupling is stabilizing and represents an energy sink. In addition, the reflective surface directly couples the scattered light waves of the same frequency (the Stokes waves associated with the incident and reflected pump waves, for example) and indirectly couples them as they propagate through the perturbation supported by the incident and reflected pump waves. For these reasons, the system including the reflective boundary may be expected to display new, stable and unstable behavior.

Linear Theory. To demonstrate these effects, we first consider SBS in a homogeneous medium at rest with respect to the reflective boundary. We write the left- and right-propagating electromagnetic vector potentials and the (linear) acoustic response as products of a slowly varying envelope and a rapidly varying phase such that

$$A(x, t) = \text{Re} [A_+ \exp i(kx - \omega_0 t) + A_- \exp i(-kx - \omega_0 t)] \quad (95)$$

and

$$\tilde{m}(x, t) = \text{Re} [\tilde{m} \exp (-2ikx)] \quad (96)$$

where ω_0 is the laser pump frequency and k is the wave number. In the dimensionless units of previous workers,⁸³ the envelopes obey the equations

$$\epsilon \frac{\partial A_+}{\partial t} + \frac{\partial A_+}{\partial x} = -i A_+ \tilde{m}^{(*)} \quad (97)$$

and

$$\frac{\partial^2 \tilde{m}}{\partial t^2} + 2\beta \frac{\partial \tilde{m}}{\partial t} + 2\Omega_s \frac{\partial \tilde{m}}{\partial x} + \Omega_s^2 \tilde{m} = -2\Omega_s A_+ A_-^* \quad (98)$$

Here x is measured in units of the growth length,

$$L_1 = \sqrt{c_s c_l} / \gamma_0,$$

t is measured in units of

$$t_1 = (\gamma_0 \nu \epsilon)^{-1},$$

$\Omega_s = \omega_s t_1$, $\beta = \nu t_1/2$, and $*$ means conjugate for the plus (+) equation only. In these units, c_s is the ion sound speed; c_l is the group velocity of light; ω_s and ν are the real frequency and damping rate, respectively, of undriven acoustic waves; γ_0 is the infinite-medium SBS growth rate; and $\epsilon = c_s/c_l$ is a small parameter. Note that Ω_s , as defined here, is inversely proportional to the pump strength.

In analyzing the linear regime of the SBS instability, we split $A_\pm$ into the incident and reflected pump waves ($A_{\pm 0}$) and the

Fig. 3-41. Contour plots of (a) the normalized SBS growth rate, γ , in the presence of a reflective internal boundary and (b) the normalized SBS frequency shift, $\Delta\Omega$, relative to the usual value.

Stokes and anti-Stokes backscattered waves ($C_{\pm}$ and $D_{\pm}$, respectively). Similarly, we split the acoustic response, $\tilde{m}$, into the perturbation supported by the beat of the pumps, $\tilde{m}_0$, and left-hand right-going waves, $\tilde{m}_{\pm}$ so that

$$A_{\pm}(x, t) = A_{\pm 0}(x) + C_{\pm}(x, t)e^{i\Omega t} + D_{\pm}(x, t)e^{-i\Omega t} \quad (99)$$

and

$$\tilde{m}(x, t) = \tilde{m}_0(x) + \tilde{m}_+(x, t)e^{i\Omega t} + \tilde{m}_-(x, t)e^{-i\Omega t} \quad (100)$$

Substitution of Eqs. (99) and (100) into Eqs. (97) and (98) results in a system of nine equations for the three quantities, $A_{\pm 0}$ and $\tilde{m}_0$, and the six decay waves, $C_{\pm}$, $D_{\pm}$, and $\tilde{m}_{\pm}$; for the sake of brevity, we do not record them here. If the reflective boundary (of reflectivity f) is located at $x = 0$ and the

laser is incident from the right at $x = L$ then the appropriate boundary conditions are $A_{+0}(0) = fA_{-0}(0)$, $C_{+0}(0) = fC_{-0}(0)$, $D_{+0}(0) = fD_{-0}(0)$, $\tilde{m}_+(0) = 0$, and $\tilde{m}_-(L) = 0$. Imposing these boundary conditions and setting $\partial/\partial t$ of the envelopes equal to zero defines a sixth-order eigenvalue problem for the frequency, $\Omega = \Omega(\epsilon, \Omega_0, L, f, \beta)$. For the usual case of small ϵ , the absolute instability threshold condition is the usual one that the growth rate in the absence of damping exceeds the damping rate: $-\text{Im} \Omega(\dots, 0) > \beta$.

As an alternative to this difficult eigenvalue problem, we have set $\Omega = \Omega_s$, neglected $\partial^2/\partial t^2$, and solved the resulting initial-value problem for the temporal evolution of the envelopes. Figure 3-41 is a contour plot of the growth rate, $\gamma(\epsilon, \Omega_s)$, for $L = 10$, $|f|^2 = 0.9$, and $\beta = 0$. The usual result obtained in the absence of a reflecting surface is $\gamma = 2$, (nearly) independent of ϵ and Ω_s ; it is evident in Fig. 3-41, however, that the reflecting surface has resulted in

Fig. 3-42. Amplitudes of the electromagnetic eigenfunctions, $|C_{\pm}|^2$ and $|D_{\pm}|^2$, and corresponding acoustic eigenfunctions, $|\tilde{m}_{\pm}|^2$, for parameters corresponding to point A in Fig. 3-41.

substantial ϵ - and Ω_s -dependence of the growth rate. For example, at point A ($\Omega_s^{-1/2} = 1/\sqrt{55} = 0.135$, $\epsilon^{1/2} = \sqrt{0.001} = 0.031$), the value of the growth rate has been increased to about 7.0, several times the usual value! The absolute instability threshold is correspondingly reduced.

Figure 3-41(b) displays the departure of the Stokes-wave frequency from its usual value, i.e., the pump frequency shifted by the acoustic frequency. The reflective boundary induces large frequency shifts, as may be seen, for example, at point A, where $\Delta\Omega = -10$. In physical units, this corresponds to a Stokes-frequency shift of about $-1.2 \omega_s$.

In Figs. 3-42(a) and 3-42(b) we display the amplitudes of the electromagnetic eigenfunctions, $|C_{\pm}|^2$ and $|D_{\pm}|^2$, and the acoustic eigenfunctions, $|\tilde{m}_{\pm}|^2$, for the parameters corresponding to point A in

Fig. 3-41. For comparison, in Figs. 3-42(c) and 3-42(d) we display $|C_{+}|^2$ and $|\tilde{m}_{-}|^2$ for the same parameters as in Figs. 3-42(a) and 3-42(b), except that $f = 0$. In the presence of the reflective boundary, the acoustic waves nearly fill the region; hence for a fixed level of scatter, they have a much smaller amplitude than in the $f = 0$ case, for which the waves are appreciable only over a few growth lengths. Consequently, for $f \neq 0$ the growth of the scattered light waves is also considerably extended. The anti-Stokes light wave, $|D_{+}|^2$, exits the region with nearly 1/10 the intensity of the Stokes $|C_{+}|^2$ wave.

Substantial modification of the current understanding of SBS in inhomogeneous or drifting media and convective SBS may also be expected to emerge from linear analyses that include a reflective internal boundary.

SBS in the Nonlinear Regime. To pursue the foregoing analysis into the nonlinear

Fig. 3-43. (a) Time history of the scattered light emerging at $x = L$, showing linear growth, saturation, and asymptotic limit-cycle behavior; (b) frequency spectrum just prior to saturation and (c) in the asymptotic state.

regime, a numerical solution of Eqs. (97) and (98) would be sufficient. Instead, we take up more general equations that are appropriate for drifting, inhomogeneous media and that allow comparison with data from laser-plasma interaction experiments⁸⁶

$$\frac{\partial A_{\pm}}{\partial t} \pm c_s \frac{\partial A_{\pm}}{\partial x} = -i \gamma_0 A_{\pm} \tilde{m}^{(*)} \quad (101)$$

and

$$\begin{aligned} \frac{d^2 \tilde{m}}{dt^2} - 2i\omega_s \left[\left(\eta + i \frac{\nu}{2\omega_s} \right) \frac{d\tilde{m}}{dt} - c_s \frac{\partial \tilde{m}}{\partial x} \right] \\ - \omega_s^2 \left(\eta^2 - 1 + i \frac{\nu}{\omega_s} \right) \tilde{m} \\ = -2\omega_s \gamma_0 A_{-} A_{+}^{*} \end{aligned} \quad (102)$$

In Eq. (102), $d/dt = \partial/\partial t + v_0 \partial/\partial x$ is the convective derivative in the drifting medium (having drift velocity v_0 in the laboratory frame where the scattered light spectra are observed); and $\eta = \Delta\omega/\omega_s + v_0/c_s$ where $\Delta\omega = \omega_- - \omega_+$ is the frequency difference of the left- and right-going light waves due to reflecting-surface Doppler shifts and the changing optical depth of the medium.⁸⁷ Frequency shifts due to SBS are contained in the phases of the envelopes, $A_{\pm}$.

In Fig. 3-43 we display numerical solutions of Eqs. (101) and (102) for a homogeneous, stationary plasma slab initialized to include the perturbation in near equilibrium with the incident and reflected pump waves. In the units of our previous analysis, the parameters are $L = 10$, $|f|^2 = 0.9$, $\Omega_s = 28$, $\epsilon = 0.001$, and $\beta = 3.0$. Figure

3-43(a) is a time history of the scattered light, $A_s(t) = [A_+(t) - A_{+0}]/A_{+0}$, emerging at $x = L$, the vacuum interface. For these parameters the data of Fig. 3-41 predict a linear growth rate $\gamma \approx 7.3 - \beta = 4.3$, and a frequency shift of $\Omega = -\Omega_s - \Delta\Omega(\epsilon, \Omega_s) = -42$. Converting to the units of Fig. 3-43 [$t \sim (k_0 c_s)^{-1}$], we obtain $\gamma = 0.26$ and $\omega - \omega_0 = 2.5 = -1.5 \omega_s$, which agree well with the observed linear growth and frequency shift.

After saturation at $t = 30 (k_0 c_s)^{-1}$, the instability evolves slowly toward its asymptotic state, a complex limit-cycle behavior characterized by multiple and half harmonics of the linear frequency shift Ω . Figure 3-43(b) shows the frequency spectrum just before saturation; note that the Stokes and anti-Stokes lines at $\Delta\omega = \pm 2.5 k_0 c_s$ are prominent and the nonlinearly generated multiple harmonics have emerged. The frequency spectrum in the asymptotic state is shown in Fig. 3-43(c); half harmonics are clearly visible at $\Delta\omega \pm 1.25 k_0 c_s$.

As the system parameters are varied (particularly as L is increased), the asymptotic limit-cycle behavior can become more and more complex, with the generation of higher-order half harmonics ($1/2^n$, $n > 1$) and their multiples. Finally, the asymptotic state can become chaotic, i.e., not discernibly periodic. These chaotic states are characterized by highly structured frequency spectra of light returning from the plasma. Such chaotic behavior is shown in Fig. 3-44(a), which is a calculated time history of reflected power $|A_+(t)|^2$. The corresponding calculated reflected-light spectrum is shown in Fig. 3-44(b).

The systems that exhibit chaotic SBS are not limited to the simple homogeneous

Fig. 3-44. Calculated (a) and (b) and measured (c) time histories of the scattered and reflected intensity and associated spectra, showing chaotic behavior.

ones just described. In fact, the background plasma conditions in the simulation that produced the curves in Fig. 3-44 correspond to those characterizing recent high-intensity, short-pulse laser-plasma-interaction experiments.⁸⁵ The plasma density, drift velocity, and temperature profiles taken from hydrodynamic simulations of these experiments exhibited the very strong gradients typical of this regime.

The qualitative similarity between the light spectra from the simulation [Fig. 3-44(b)] and the experimental data [Fig. 3-44(c)] is striking. Moreover, the agreement can be made semiquantitative. At peak laser power ($t = 60$ ps) the hydrodynamic simulation yields an electron temperature $T_e = 15$ keV at the peak laser intensity of 5×10^{15} W/cm². The normalized sound speed at this temperature is $\epsilon = c_s/c = 2.8 \times 10^{-3}$. The calculated frequency shifts in Fig. 3-44(b) are converted to wavelength shifts in A by multiplying by $-(\lambda_0)^{-1} = -30(\text{\AA})$ for Nd:glass laser light. Thus, both the relative spacing and placement of the peaks in the simulation spectrum are in agreement with the data.

Note that if we appeal to multiple-harmonic generation (rather than fractional-harmonic generation) to explain the experimental spectra, the sound speed derived from the spacing of the spectrum peaks, $c_s/c = \Delta\lambda/(2\lambda)$, yields an electron temperature of ~ 0.4 keV; the value is much too low for plasmas illuminated at intensities of 5×10^{15} W/cm² and is not in agreement with measurements of T_e in similar experiments. In our simulation we set the parameters $|f|^2 = 0.1$ and $v_d = 0.4 \omega_p$; similar spectra may be obtained with somewhat different values of these parameters.

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One-Dimensional Time-Implicit Hybrid Simulations

Recent time-implicit particle-simulation methods, using moments^{88,89} or a direct approach,⁹⁰ allow large time steps $\Delta t > \omega_p^{-1}$, where $\omega_p = (4\pi e^2 n_e / m_e)^{1/2}$ is the plasma frequency, n_e the electron density, $-e$ the electron charge, and m_e the electron mass. However, these methods remain limited by the condition $\Delta t \ll (kv_{th})^{-1}$, which specifies

that an electron near the thermal velocity v_{th} must not move across a large fraction of a characteristic wavelength $\lambda = 2\pi/k$ during a time step. Taken together, these conditions imply $k\lambda_D \ll 1$, where $\lambda_D = v_{th}/\omega_p$ is the Debye length. Thus, implicit methods yield improved efficiency only in the case of long wavelengths or in dense, cold plasmas. However, dense, cold plasmas also have large collision frequencies ν_c , and implicit codes operating in this regime must therefore allow time steps with $\nu_c \Delta t > 1$.

In inertial confinement fusion, an area of application of implicit methods concerns interactions between sparse, hot regions, such as the corona, and dense, cold regions, such as the target, which are typically several orders of magnitude denser and colder than the corona. Such problems yield naturally to a hybrid formulation, in which the sparse, hot component is represented as particles and satisfies the conditions $\omega_p \Delta t \ll 1$, $kv_{th} \Delta t \ll 1$, and $\nu_c \Delta t \ll 1$, while the dense, cold component is given a fluid representation and is characterized by $\omega_p \Delta t > 1$, $kv_{th} \Delta t \ll 1$, and $\nu_c \Delta t > 1$.

One-dimensional hybrid codes of this type are still in the development stage, but they have interesting possibilities. In our code,⁹¹ we model ions and thermal electrons in terms of Braginskii's two-fluid equations,⁹² including ion and electron inertia, friction and thermal forces, and electron and ion temperatures. We represent fluid quantities on a uniform Eulerian grid, which we also use for particles. We use flux-corrected transport to minimize numerical diffusion; vacuum regions are allowed, and the boundaries are absorbing and re-emitting, with either fixed external charges or a fixed potential difference across the system.

We represent energetic electrons as weighted particles with position x , parallel velocity v_x , and transverse velocity $v_\perp$. These particles interact with the fluid via the electric field and by collisions involving energy exchange and Monte Carlo scattering. Particles are emitted and reabsorbed within the fluid according to the local collision-frequency, with emission when $\nu_{fluid} < \nu_1$ and reabsorption when $\nu_{particle} < \nu_2$, where ν_1 and ν_2 are problem-dependent collision-frequency thresholds. The threshold ν_2 and time step Δt are assumed to satisfy the condition $\nu_c \Delta t \ll 1$, allowing collisions to be applied separately, as shown in Fig. 3-45.

The implicit field computation is done by a direct method.⁹⁰ Let

$$G_j \equiv E_{j+1} - E_j - \frac{a^2 \Delta x}{2} [Z(n_{ij} + n_{ej}) - n_{ej} - n_{ej+1} - n_{pj} - n_{pj+1}] \quad (103)$$

where E_j , n_{ij} , n_{ej} and n_{pj} are respectively the electric field, ion density, fluid electron density, and particle density at cell j , Δx is the mesh size, $a = L/\lambda_D$, L is the characteristic length, and λ_D is the Debye length corresponding to the characteristic density n_0 and temperature T_0 . Here, lengths are measured in units of L and velocities in units of $v_{th} = (T_0/m_e)^{1/2}$. The Poisson equation, $G_j = 0$, is solved using Newton iterations by solving

$$(\partial G_j / \partial E_j)^q \delta E_j^{q+1} = -G_j^q \quad (104)$$

where q is the iteration level. Both fluids and particle quantities are computed implicitly within the iteration loop, and the matrix $\partial G_j / \partial E_j$ is computed from the linear dependence of the n_{ej} 's and n_{pj} 's on the E_j 's in the fluid- and particle-pushing algorithms. This matrix has four diagonals in the present case.

Several examples involving expansion of a plasma slab into vacuum and electron current across a diode gap with plasma electrodes are presented in Figs. 3-46 through 3-49 to illustrate the capabilities of the code. The expansion results shown in Fig. 3-46 correspond to a plasma slab initially occupying a region of width $2 \times 10^3 \lambda_D$ located at the center of the system.

Fig. 3-45. Implicit hybrid simulation of interaction between hot (corona) and cold (target) regions.

Fig. 3-46. Simulation of a plasma slab expanding into a vacuum.

A mass ratio $m_i/m_e = 900$ is assumed with $T_i/T_e = 10^{-2}$ and a collision time $\tau_c = 0.528 \omega_p^{-1}$ (corresponding approximately to a 1- μm -thick foil with initial electron density $\sim 10^{23} \text{ cm}^{-3}$ and temperature $\sim 500 \text{ eV}$). The slab is initially represented as fluids only, but energetic electrons become particles as they expand into vacuum, as shown in the phase plot in Fig. 3-46 at $t = 400 \omega_p^{-1}$. The energetic electrons are trapped by strong electric fields near the plasma-vacuum boundary and oscillate across the slab, as shown in the phase plot at $t = 4000 \omega_p^{-1}$. During their transit across the slab, the electrons lose energy by collisions with thermal electrons so that some of them are reabsorbed into the fluid.

The diode results, shown in Fig. 3-47, correspond to a system of length $L = 1.96 \times 10^5 \lambda_{D0}$ and a potential $\phi = 2.3 \times 10^3 T_{e0}/e$ (where T_{e0} and λ_{D0} are electrode temperature and Debye length). The simulation is initialized with fluid electrons only, and particles are emitted as a pure electron gas is formed on the cathode surface (right). These electrons are accelerated across the

gap and are absorbed as they reach the anode. Ohmic potential drops in the electrodes reduce the potential across the gap to $1.65 \times 10^3 T_{e0}/e$, which corresponds to a Child-Langmuir limit $J_{CL} = 5.0 \times 10^{-6} en_{e0} v_{th0}$. In this example, the diode current, $J_p = 0.8 \times 10^{-6} en_{e0} v_{th0}$, is limited by the cathode resistivity, but the Child-Langmuir limit is approached in simulations with low electrode-collision frequency.

The results, shown in Fig. 3-47, correspond to fixed ions and to constant temperatures. Simulations in which these constraints are removed, and in which the complete set of fluid equations is used, display interesting features. For example, the phase plot of Fig. 3-48 shows a case in which energetic electrons have been emitted suddenly within the cathode. The sudden emission occurs when the electrons in a slice of the cathode, near the surface, are sufficiently heated by ohmic losses to become collisionless. The emission is accompanied by large electric fields generated by expansion of the hot electrons into the gap. When the sudden emission is inhibited by

Fig. 3-47. Simulation of an electron current across a diode gap with plasma electrodes.

lowering the emission threshold ν_i , the diode operates as in the case of Fig. 3-47, but ions are now pulled from the anode surface and begin to close the gap, as shown in Fig. 3-49.

The time step in our simulations was adjusted dynamically to satisfy the conditions $u_e \Delta t < 0.2\Delta x$, $v_{the} \Delta t < 0.2\Delta x$, $(eE/m_e \Delta x)^{1/2} \Delta t < 0.2$ (where u_e and v_{the} are local fluid electron velocity and particle thermal velocity). For example, the time step for the case shown in Fig. 3-47 varied from $\omega_p \Delta t = 2$ to 10 ($\Delta x = L/100$), but for some simulations these conditions can yield time steps $\omega_p \Delta t < 1$.

In addition to having the desirable feature of a large time step, our simulation also has the ability to follow steep gradients with density ratios of several orders of magnitude across interfaces. Note, for example, the density ratio n_e/n_p of $\sim 10^7$ across the cathode surface in Fig. 3-47 and the ion density ratio of $\sim 10^6$ across the anode surface in Fig. 3-49. We expect that such codes can be generalized to two dimensions and to formulations that include electromagnetic and relativistic effects.

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Electromagnetic Direct-Implicit Algorithm

During 1982, we extended the "direct" implicit particle-in-cell (PIC) simulation approach⁹³⁻⁹⁵ to apply to simulations that include full electromagnetic fields. The new algorithm offers advantages relative to "moment" implicit electromagnetic algorithms⁹⁶ and may help in our quest for simpler, more robust implicit codes.

Implicit fields reproduce electromagnetic waves at long wavelengths ($\gg c\Delta t$). At short wavelengths, the electrostatic, magnetostatic, and inductive electric fields are retained, as in a "Darwin" code.⁹⁷ At all wavelengths, Langmuir waves are stabilized, as in a direct-implicit electrostatic code. The electrostatic fields are accurate for wavelengths longer than the electron transit distance ($v_{the} \Delta t$). These properties make an implicit electromagnetic code attractive for modeling intense electron flow, which is subject to pinching, the Weibel instability, and other processes that generate mag-

netic fields, which, in turn, alter the electron flow.⁹⁸

Implicit Time Differencing of the Particle and Field Equations. To begin our outline of an implicit algorithm, we select the D_1 time-differencing scheme⁹⁹ for the particles

$$\frac{x_{n+1} - x_n}{\Delta t} = v_{n+1/2}, \quad (105a)$$

$$\begin{aligned} \frac{v_{n+1/2} - v_{n-1/2}}{\Delta t} \\ = \bar{a}_n + \frac{v_{n+1/2} + v_{n-1/2}}{2} \times \frac{qB_n}{mc}, \end{aligned} \quad (105b)$$

where

$$\bar{a}_n = 1/2 \left[\bar{a}_{n-1} + \frac{q}{m} E_{n+1}(x_{n+1}) \right]. \quad (105c)$$

Fig. 3-48. Simulation of sudden emission of particles caused by heating of a slice of the cathode ($L = 1.96 \times 10^3 \lambda_{D0}$).

Fig. 3-49. Simulation of ion expansion into the diode gap ($L = 1.96 \times 10^3 \lambda_{D0}$).

Desired features of the implicit differencing of the Maxwell equations include

- Accuracy in dispersion, $\text{Re } \omega(k)$, and weak damping (for example, $\text{Im } \omega(k)/ck = O(ck\Delta t)^3$, where k is the wave vector) at long wavelengths.
- At short wavelengths ($\approx 2\Delta x$), stability (preferably damping) despite $c\Delta t \gtrsim \Delta x$ (violation of the Courant condition for explicit differencing), and dissipation of inaccurately calculated short wavelengths.
- Compatibility with implicit particles.
- Adaptability to general boundary conditions.
- Simplicity.
- Economy in storage.
- Optional recoverability of the centered second-order scheme now commonly used for the fields.
- Optional recoverability of the centered Darwin scheme.⁹⁷

For the fields, we adapt implicit schemes developed for the particle equations of motion. For example, in the particle equations in Eq. (105) we drop the $v \times B$ term, replace x by E , v by $c\nabla \times B$, and $\bar{a}$ by $-c^2 \nabla \times \nabla \times \nabla \bar{E}$ to obtain the Maxwell equations (in rationalized cgs units)

$$c\nabla \times B_{n+1/2} = I_{n+1/2} + \frac{E_{n+1} - E_n}{\Delta t} \quad (106a)$$

and

$$-c\nabla \times \bar{E}_n = \frac{B_{n+1/2} - B_{n-1/2}}{\Delta t}, \quad (106b)$$

where

$$\bar{E}_n = 1/2[\bar{E}_{n-1} + E_{n+1}] \quad (106c)$$

is the result of a recursive low-pass filter with phase error $O(\Delta t^3)$. The phase error is an advance instead of the lag that we get if E_{n+1} is not used; it provides stability at $ck\Delta t \gg 1$.

The code must solve the coupled set of Eqs. (105) and (106). A price we pay in employing implicit differencing is that time-cycle splitting of the particle and field-time advances is more complicated. To advance the field values implicitly, we eliminate E_{n+1} or $B_{n+1/2}$ from the coupled Eqs. (106a), (106b), and (106c) to obtain a single elliptic equation. Eliminating $B_{n+1/2}$ to form an

equation for E_{n+1} yields

$$\begin{aligned} E_{n+1} + \frac{1}{2} c^2 \Delta t^2 \nabla \times \nabla \times E_{n+1} \\ = E_n - I_{n-1/2} \Delta t \\ + c \Delta t \nabla \times \left(B_{n-1/2} - \frac{1}{2} c \Delta t \nabla \times \bar{E}_{n-1} \right); \quad (107) \end{aligned}$$

alternatively, we can eliminate E_{n+1} to form an equation for $B_{n+1/2}$ such that

$$\begin{aligned} B_{n+1/2} - \frac{1}{2} c^2 \Delta t^2 \nabla^2 B_{n+1/2} \\ = B_{n-1/2} + \frac{1}{2} c \Delta t \nabla \times (I_{n-1/2} - \bar{E}_{n-1} - E_n). \quad (108) \end{aligned}$$

In either case, the right-hand side is composed of known fields, and the left-hand sides have well-behaved elliptic operators.

To form a B_n for use in the particle mover, we use, for example,

$$B_n = B_{n-1/2} - \frac{1}{2} c \Delta t \nabla \times E_n. \quad (109)$$

The Direct Method for Implicit Particles and Fields. The essence of the direct method is that we work directly with the particle equations of motion and the particle/field coupling equations. The equations are linearized about an extrapolation estimate for their values at the new time level, $n+1$. The future values of $\{x, v\}$ are divided into two parts

- Increments $\{\delta x, \delta v\}$, which depend on the (unknown) fields at the future time level $n+1$.
- Extrapolations $\{x_{n+1}^{(0)}, v_{n+1/2}^{(0)}\}$, which incorporate all other contributions to the equation of motion.

The increments $\{\delta x, \delta v\}$ are evaluated by linearization of each equation of motion⁹³⁻⁹⁵; here, we have

$$\begin{aligned} \frac{\delta x_{n+1}}{\Delta t} &= \delta v_{n+1/2} \\ &= \frac{1}{2} (I + R) \cdot \left(\frac{q \Delta t}{2m} \right) E_{n+1}(x_{n+1}), \quad (110) \end{aligned}$$

where the operator R effects a rotation through angle $-qB_n \Delta t / mc$. The corresponding

densities $\{\rho_{n+1}^{(0)}, j_{n+1/2}^{(0)}\}$ and $\{\delta\rho, \delta j\}$ are inserted into Maxwell's equations.

Evaluation of the Extrapolated Densities. The extrapolated densities, $\{\rho_{n+1}^{(0)}, j_{n+1/2}^{(0)}\}$, are evaluated as in explicit codes, such as ZOHAR (Ref. 100) and WAVE, from x_{n+1} , $v_{n+1/2}$, and x_n . At the grid point located at X_i ,

$$\rho_{n+1}^{(0)} = \sum qS(X_i - x_{n+1}^{(0)}) \quad (111)$$

and

$$j_{n+1/2}^{(0)} = \sum qv_{n+1/2}^{(0)} \times S\left[X_i - \frac{1}{2}(x_n + x_{n+1}^{(0)})\right] \quad (112a)$$

or

$$= \sum qv_{n+1/2}^{(0)} \frac{1}{2} [S(X_i - x_n) + S(X_i - x_{n+1}^{(0)})] \quad (112b)$$

To correct the small error in $\nabla \cdot j_{n+1/2}^{(0)}$ (due to the slightly nonconservative but otherwise beneficial method of collection of j [Refs. 100 and 101]), we replace $j_{n+1/2}^{(0)}$ with

$$j_{n+1/2}' = j_{n+1/2}^{(0)} - (\nabla\psi/\Delta t) \quad (113a)$$

where

$$-\nabla^2\psi = \rho_{n+1}^{(0)} + \nabla \cdot (\Delta t j_{n+1/2}^{(0)} - E_n) \quad (113b)$$

Evaluation of the Increments Due to Future Fields. The care with which we express the increments due to future fields, $\{\delta\rho, \delta j\}$, is a compromise between complexity and strong convergence.^{93,95} If necessary, we may evaluate increments rigorously as derivatives of Eqs. (111) and (112) ("strict differencing"⁹³) or as simplified difference representations^{93,95} of

$$\delta\rho = -\nabla \cdot (\rho\delta x) \quad (114)$$

and

$$\delta j = \rho\delta v - \frac{1}{2} \nabla \times (j \times \delta x) \quad (115)$$

for each species. This form for δj trivially conserves charge: $\delta\rho + \Delta t \nabla \cdot \delta j = 0$. This property can easily be preserved in the spatial differencing of δj .

The terms in Eq. (115) have both analytic and pictorial justifications (see Fig. 3-50). A heuristic derivation of δj uses an analogy to magnetization current. The magnetic moment of the current loop in the last diagram in Fig. 3-50 is

$$\left(\frac{1}{2c}\right) \int x \times dx = \frac{q}{2c\Delta t} \delta x \times (v_{n+1/2}^{(0)} \Delta t) \quad (116)$$

The current due to a density n of these terms is

$$\begin{aligned} \delta j &= c \nabla \times M \\ &= c \nabla \times \left[n \frac{q}{2c\Delta t} \delta x \times v_{n+1/2}^{(0)} \Delta t \right] \\ &= \frac{1}{2} \nabla \times (\delta x \times \rho v_{n+1/2}^{(0)}) \quad (117) \end{aligned}$$

which leads to the last term in Eq. (115).

We now have everything we need to write an equation for E_{n+1} . On substituting our expressions for $\rho_{n+1}^{(0)}$, $\delta\rho$, $j_{n+1/2}^{(0)}$, and δj into the field equations [Eqs. (106) and (107)], we have

$$\begin{aligned} c \nabla \times \left[B_{n+1/2} + \frac{j_{n+1/2}^{(0)} \times \delta x}{2c} \right] \\ = j_{n+1/2}' + \rho_{n+1}^{(0)} \delta v + \frac{E_{n+1} - E_n}{\Delta t} \quad (118a) \end{aligned}$$

and

$$\begin{aligned} -\frac{1}{2} c \nabla \times (\bar{E}_{n+1} + E_{n+1}) \\ = \frac{B_{n+1/2} - B_{n-1/2}}{\Delta t} \quad (118b) \end{aligned}$$

The implicit terms $\rho \delta v$ and $J \times \delta x$, with Eq. (110), are seen to be

$$\rho_{n+1}^{(0)} \delta v = \left[\sum_s \frac{\rho_s q_s \Delta t}{2m_s} \frac{1}{2} (I + R_s) \right] \cdot E_{n+1} \quad (119a)$$

and

$$J_{n+1}^{(0)} \delta x = \left[\sum_s \frac{J_s q_s \Delta t^2}{2m_s} \frac{1}{2} (I + R_s) \right] \cdot E_{n+1} \quad (119b)$$

The square brackets indicate what must be collected from the particles of each species, s . If only the electrons are implicit, only they appear in Eq. (119). In this case, the terms in square brackets require only a knowledge of the electrons' ρ and J [in addition to the net ρ and J used on the right side of Eq. (118a)].

The divergence of the Ampere-Maxwell equation recovers exactly our electrostatic

implicit field equation,^{93,94}

$$\rho_{n+1}^{(0)} = \nabla \cdot (I + \chi) \cdot E_{n+1} \quad (120)$$

where the implicit susceptibility represented as $\chi = (\rho_{n+1}^{(0)} q \Delta t / 2m) (I + R)/2$ is a tensor due to the rotation R induced by B . Equations (118) and (119) are the simplest yet proposed for implicit field prediction, both in themselves and in what we must accumulate from the particles.

Generalizations. To include relativity, we would linearize the relativistic particle equation of motion.^{100,101} Electron-ion collisions ($\nu \lesssim \Delta t^{-1}$) may be added as an additional rotation in the equation of motion.

If a component of the plasma is modeled by fluid equations, then those equations can be linearized to find $\{\delta \rho, \delta J\}$ (see the previous article). Combining fluid and particle descriptions is difficult, but no more so in the direct method than in the moment method.

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Fig. 3-50. Geometric interpretation of the terms in $J = J^{(0)} + \delta J$.

While $J_n + v_2$ corresponds to moving the particle directly from x_n to x_{n+1} , it can be regarded as the sum of three motions:

1. motion from x_n to $x_n^{(0)} + 1$, giving $J_n^{(0)}$;

2. motion from $x_n^{(0)} + 1$ to x_{n+1} , giving $\rho \delta v$;
3. plus a circulation term, $-\frac{1}{2} \nabla \times (J \times \delta x)$, to cancel the "detour" to $x_n^{(0)} + 1$. This is not needed to get ρ_{n+1} , but it does affect B and E_T .

Implementation Issues in Direct-Implicit Electrostatic Simulation

The PIC computer codes for plasma simulation, which employ implicit finite differences as approximations to the time derivatives in the particle equations of motion, can be far more efficient than traditional explicitly differenced codes. During the past year, we have further developed our direct method for implicit PIC simulations. Although our approach is very simple in concept, there are many important implementation issues that must be resolved in the construction of a practical computer code.

A desirable property of any simulation algorithm is that it algebraically preserve conservation of momentum. Here, we demonstrate the feasibility of momentum-conserving direct-implicit schemes. In addition, we describe a simple but efficient one-dimensional electrostatic simulation code that implements the momentum-conserving algorithm.

The basic concepts of the direct method have been described elsewhere¹⁰²⁻¹⁰⁶; we

summarize them briefly here. In the direct method, we write the future position of a particle, x^{n+1} , as the sum of an explicitly computable position $\tilde{x}$ and a displacement δx due to the (as yet unknown) future field such that

$$\delta x = x^{n+1} - \tilde{x} = \beta \Delta t^2 a^{n+1} \quad (121)$$

where $1/2 \leq \beta \leq 1$ is a constant of the time-differencing chosen. The charge density is the sum of $\tilde{\rho}$ from the $\tilde{x}$'s and a correction

$$\begin{aligned} \rho^{n+1} &\simeq \tilde{\rho}^{n+1} - \nabla \cdot [\tilde{\rho}^{n+1}(x) \delta x(x)] \\ &\simeq \tilde{\rho}^{n+1} - \nabla \cdot [\chi(x) E^{n+1}(x)] \quad (122) \end{aligned}$$

$$\chi(x) \equiv \beta \Delta t^2 \frac{q}{m} \tilde{\rho}^{n+1}(x) = \beta \omega_p^2(x) \Delta t^2 \quad (123)$$

is an effective susceptibility. The field equation is therefore

$$\nabla \cdot \mathbf{E} + \chi(x) E^{n+1}(x) = \tilde{\rho}^{n+1}(x) \quad (124)$$

Thus, by computation of $\tilde{x}$ and accumulation of $\tilde{\rho}^{n+1}$ at the end of the particle step from x^{n+1} to x^n , we can compute the field E^{n+1} one step before we know the x^{n+1} .

The above direct-implicit equations could be differenced in a straightforward manner, yielding a tridiagonal field equation. The simplicity of such an approach recommends it for application to a two-dimensional or electromagnetic code. However, because of the finite size of the simulation particles, such a differencing introduces errors at short wavelengths. Here, we describe a strict scheme with more desirable properties. At each cell j ,

$$\begin{aligned} \rho_j^{n+1} &= \sum_k F_{jk} (x_k^{n+1}) \\ &\simeq \sum_k \left[\tilde{\rho}_{jk} + (x_k^{n+1} - \tilde{x}_k) \frac{\partial \tilde{\rho}_{jk}}{\partial \tilde{x}_k} \right] \quad (125) \end{aligned}$$

where $\tilde{\rho}_{jk} = \rho_{jk}(\tilde{x}_k)$ is the charge density in cell j due to particle k at $\tilde{x}_k$. From Eq. (121), we approximate

$$x_k^{n+1} \simeq x_k^{(1)} \equiv \tilde{x}_k + \beta \Delta t^2 a_k^{n+1}(\tilde{x}_k) \quad (126)$$

Use of $a_k^{n+1}(\tilde{x}_k)$ not $a_k^{n+1}(x_k^{n+1})$ leads to errors of order $\omega_{tr}^2 \Delta t^2$, where ω_{tr} is the electrostatic trapping frequency. Using Eq. (126) in Eq. (125), we see that the strict field equation is

$$\begin{aligned} (-\nabla^2 \phi_j^{n+1}) &= \sum_k \tilde{\rho}_{jk} \\ &+ \beta \Delta t^2 \sum_k \frac{\partial \tilde{\rho}_{jk}}{\partial \tilde{x}_k} a_k^{n+1}(\tilde{x}_k) \quad (127) \end{aligned}$$

Introducing the finite-size-particle shape function S , we write

$$\tilde{\rho}_{jk} = q_k \frac{S(\tilde{x}_k - x_j)}{\Delta x} \quad (128)$$

$$a_k^{n+1}(\tilde{x}_k) = \frac{q_k}{m_k} \sum_i S(\tilde{x}_k - x_i) E_i^{n+1}$$

and thus the second term of the right member of Eq. (127) is just a linear operation

$$\sum_i \mathbf{W}_{ji} \phi_i$$

For the usual tent-function choice of S , $\mathbf{W}$ is pentadiagonal and nonsymmetric.

The effective charge-deposition locations one time-step ahead are well-defined as $x_k^{(1)}$, where S associated with (linearized about) $\tilde{x}_k$ is used. If S is a one-dimensional tent function, linearization is exact, provided a_k^{n+1} does not induce a cell crossing. Thus, if we subsequently advance particles using $x_k^{n+1} = x_k^{(1)}$, we obtain exact charge prediction [equality in Eq. (125)] in the weak-field regime where cell-crossing errors are absent. More significantly, we can obtain algebraic momentum conservation (with arbitrary field amplitude) by arranging that the field used to advance a particle be interpolated at the same place its charge is deposited (to avoid a self-force) such that

$$x_k^{n+1} = \tilde{x}_k + \beta \Delta t^2 a_k^{n+1}(x_k^{(1)}) \quad (129)$$

i.e., a reinterpolation but still using cells nearest $\tilde{x}_k$. The result does not depend upon S being a tent function and is

unaffected by cell crossings. Furthermore, the S in the second equation of Eq. (128) need not even be the same S as all the others; use of nearest-grid-point weighting here can lead to a symmetric tridiagonal $\mathbf{W}$, while retention of a tent function in Eq. (129) preserves smoothing. A spatial filter can be applied to the field used in Eq. (129) after the field solution is complete. While inconsistent (in the sense of Ref. 105), such a filter need not destroy momentum conservation. A truly consistent filter must be incorporated into the field solution itself; this may be implemented for an arbitrary filter with conjugate-gradient iterative techniques, which are far easier to implement for symmetric systems.

From the above, we see that the momentum-conserving scheme (and in fact any one-pass scheme, since cell crossings will in general occur) does not lead to truly exact field prediction. Furthermore, the scheme is inexact in the sense that the equation of motion should ideally be

$$\dot{x}_k^{n+1} = \dot{x}_k + \beta \Delta t^2 a_k^{n+1}(\dot{x}_k^{n+1}) \quad (130)$$

Equation (130) can only be made to conserve momentum by modifying the field equation so that the system becomes nonlinear and intractable without multiple passes through the particles. We have implemented a simple corrector iteration,¹⁰³ which converges at least linearly to a solution of Eq. (130), while conserving momentum at each step. However, we have found that on most problems, corrector iterations are not necessary.

We have developed a simple unmagnetized simulation code using the methods outlined above and in the references. The

code DIMPLES (for Direct IMPLICIT Electrostatic Simulation) implements the optimized C_1 scheme of Ref. 106 in a manner whereby all quantities are stored at integral time levels, and code start-up requires little special computation. The time differencing scheme in its original leapfrog form is

$$\begin{aligned} \dot{x}^n &= \dot{x}^{n-1} + v^{n-1/2} \Delta t + c_0(a^n - a^{n-1})\Delta t^2 \\ &\quad + c_1(a^{n-1} - a^{n-2})\Delta t^2, \\ v^{n+1/2} &= v^{n-1/2} + a^n \Delta t. \end{aligned} \quad (131)$$

The algorithmic steps, assuming E^n , x_k^{n-1} , v_k^{n-1} , a_k^{n-2} , and a_k^{n-1} are known, are:

1. $\dot{x}^n \leftarrow \dot{x}_k^{n-1} + v_k^{n-1} + (1/2 + c_1 - c_0)a_k^{n-1} - c_1 a_k^{n-2}$.
2. $\ddot{a}^n \leftarrow q/m E^n(\dot{x}^n)$.
3. $a^n \leftarrow q/m E^n(\dot{x}^n + c_0 \ddot{a}^n)$ (S associated with $\dot{x}^n$).
4. $x_k^n \leftarrow \dot{x}^n + c_0 a^n$.
5. $v_k^n \leftarrow v_k^{n-1} + (1/2)(a_k^{n-1} + a^n)$.
6. Store a_k^{n-1} in a_k^{n-2} array; store a^n in a_k^{n-1} array.
7. $\dot{x}^{n+1} \leftarrow \dot{x}_k^n + v_k^n + (1/2 + c_1 - c_0)a_k^n - c_1 a_k^{n-1}$.
8. Deposit ρ at $\dot{x}^{n+1}$.
9. Solve field equation.

In the above, particle array quantities have the subscript k , and we have set $\Delta t = 1$. For the optimized scheme, $c_0 = 0.302$, $c_1 = 0.04$. At time step $n = 0$, we set a^n into the a_k^{n-1} array and perform steps 7 through 9 with $c_0 = 0.25$, $c_1 = 0$.

Our coding is comparable in complexity and length to that for an explicit one-dimensional magnetized code; DIMPLES runs essentially at the same speed (<2.5 μ s per particle per time step) as an explicit unmagnetized code on a vector computer such as the Cray-1, since the same gather/scatter work is involved.

DIMPLES uses either consistent local digital filters for spatial smoothing of its field or a Fourier filter implemented by a functional iteration (not yet conjugate gradients), which converges for reasonably uniform systems. The effect of consistent spatial filtering is altered for $\chi \gg 1$; the change is such that local digital filters yield an adequately smooth fall-off as a function of wavelength. Both strict and simplified

Fig. 3-51. Ion phase space (w_x vs x) at $t = 20,000 \omega_{pe}^{-1}$.

(symmetric tridiagonal) field solutions have been implemented, and the results agree very closely with each other.

As a practical example of code performance, we consider the free expansion of a plasma slab into vacuum.¹⁰⁷ A rarefaction moves into the slab at the ion sound speed, while electrons and ions leave the slab in an ambipolar process. This problem was run on DIMPLES with moderately strong spatial filtering and the following parameters: electron plasma frequency $\omega_{pe} = 1.0$, time step $\Delta t = 8.0$, number of cells $n_g = 512$, number of particles of each species $N_e = N_i = 8192$, mass ratio $m_i/m_e = 900$, temperature ratio $T_e/T_i = 10$, cell size $\Delta x = 20\lambda_D$ (where λ_D is the Debye length), system length $L = 1.0 = 10\,240\lambda_D$, initial slab width $L_s = 2560\lambda_D$, electron thermal velocity $v_e = 9.8 \times 10^{-5} L\omega_{pe}$, ion thermal velocity $v_i = 1.03 \times 10^{-6} L\omega_{pe}$, and ion sound speed $c_s = 3.3 \times 10^{-6} L\omega_{pe}$. The problem was run until time $t_{\max} = 20\,000\omega_{pe}^{-1}$ was reached. In Fig. 3-51, the ion phase space (v_i vs x) at the end of the problem is shown. In Fig. 3-52, the ion charge density as a function of position x is shown at various problem times. At each side, the rarefaction can be seen to move inward to $c_s t_{\max} = 0.066 L = 675\lambda_D$ by the end of the problem. During the run, the electrons give up about half their thermal energy to the ions, while the field energy is much smaller than electron or ion energies. Total energy was conserved within 2.5%; as expected, total momentum was conserved within roundoff tolerances. The run took slightly over two minutes of Cray-1 time.

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Smoothing and Spatial Grid Effects in Direct-Implicit Particle Simulation

The recently developed direct-implicit particle simulation method^{108,109} relaxes the severe stability constraints on size of the time step that exist in explicit formulations (e.g., $\omega_p \Delta t < \mathcal{O}(1)$, where ω_p is the plasma frequency). This permits more efficient computer simulation of long-time-scale phenomena at long wavelengths. The time-step constraints that do arise are much less stringent, having to do with accuracy of the

particle orbits and of the plasma dielectric response.¹¹⁰ The direct-implicit method readily extends to a two-dimensional magnetized algorithm¹¹¹ and to an electromagnetic version (see "Electromagnetic Direct Implicit Algorithm," earlier in this section).

Here, we investigate convergence and linear-dispersion properties in direct-implicit particle codes. We show that inconsistent spatial smoothing and simplified differencing are closely related and that care must be exercised to avoid degradation of both convergence and stability. Our analysis has already been used to correct instability in practical applications.¹¹¹

The essence of the direct-implicit method is a linear expansion of the charge density around a known particle position plus an additional linear increment due to the acceleration at the next time step.^{108,109} Moving the resulting charge-density increment to the left side of Poisson's equation gives an elliptic equation for the potential ϕ , e.g.,

$$-\nabla \cdot [(1 + \tilde{\chi}) \nabla \phi_{n+1}] = \tilde{\rho}_{n+1} \quad (132)$$

where $\tilde{\chi} \propto \beta \omega_p^2 \Delta t^2$ is the effective linear susceptibility, β is a parameter controlling implicitness, and $\tilde{\rho}_{n+1}$ is the charge density at time level $n+1$ computed from known particle data. The relative error in the expansion is $\beta \omega_p^2 \Delta t^2 = \beta q E \Delta t^2 / m L_i$, where ω_{tr} is the electrostatic trapping frequency and L_i is the length scale over which E varies.

Dispersion characteristics of time-integration schemes are investigated in

Fig. 3-52. Ion charge density vs x at various problem times (t in units of ω_{pe}^{-1}).

Ref. 110 for $\Delta x = 0$, and spatial differencing issues are examined in Ref. 109. Rigorous adherence to systematic and conservative differencing leads to good convergence and stability characteristics.¹⁰⁹ Convergence and stability can be degraded by ad hoc spatial differencing and inconsistent spatial smoothing, which are related effects.

In Ref. 109, Poisson's equation is cast in an iterative form,

$$-\nabla \cdot [(1 + \chi^{(r)}) \nabla \delta \phi^{(r+1)}] = \rho^{(r)} + \nabla^2 \phi^{(r)}, \quad (133)$$

where $\chi^{(r)}$ is the effective susceptibility constructed from known data,¹⁰⁹ and r and $r + 1$ are iteration levels at time level $n + 1$. Convergence can be measured by calculating the ratio of successive residuals $|\rho^{(r+1)} - \nabla \cdot \mathbf{E}^{(r+1)}|$. When strict spatial differencing and a consistent formula for $\chi^{(r)}$ are employed, the calculated residual vanishes through $O(\beta \omega_p^2 \Delta t^2)$ except for errors due to particles that cross cell boundaries.¹⁰⁹ All code applications so far have used a simplified susceptibility that omits a Newton-Raphson correction¹⁰⁹; this weakens convergence, and residuals scale as $\beta \omega_p^2 \Delta t^2$.

Consistent spatial filtering in the direct method¹⁰⁹ takes the form

$$-\nabla \cdot [(1 + S_2 S_1 \chi^{(r)}) \nabla \delta \phi^{(r+1)}] = \rho^{(r)} + \nabla^2 \phi^{(r)}, \quad (134)$$

where S_1 and S_2 are smoothing operators, $\rho_s \equiv S_2 S_1 \rho$, and $(\phi_s, \delta \phi_s) \equiv S_2(\phi, \delta \phi)$ are smoothed quantities. Convergence is weaker for inconsistent smoothing, and, when simplified, ad hoc differencing schemes are employed to reduce the bandwidth of the matrix field equations.^{109,111} A representative calculation gives

$$\frac{\rho^{(r+1)} - \nabla \cdot \mathbf{E}^{(r+1)}}{|\rho^{(r)} - \nabla \cdot \mathbf{E}^{(r)}|} = \frac{[(\nabla \cdot \chi^{(r)} \delta \mathbf{E}^{(r+1)})' - \nabla \cdot \chi^{(r)} \delta \mathbf{E}^{(r+1)}]}{|\nabla \cdot (1 + \chi^{(r)}) \delta \mathbf{E}^{(r+1)}|}, \quad (135)$$

where $\nabla \cdot \chi \delta \mathbf{E}$ uses strict differencing and consistent spatial filtering, and $(\quad)'$ indicates a simplified differencing version with the option of inconsistent spatial filtering. Rapid variations in the plasma density and

the electric field degrade convergence via differences in χ and field.

Analysis of the difference equations for the particle equations of motion, as applied to the simple harmonic oscillator,¹¹⁰ gives information on stability at large $\omega \Delta t$ and accuracy at small $\omega \Delta t$. Consider implicit schemes with $\mathbf{E}_{n+1}^{(0)} = 0$ used in obtaining the right side of Eq. (132) or (133). For the two classes of time integration schemes in Ref. 110,

$$\left. \begin{aligned} \frac{X}{A \Delta t^2} &= c_0 + \frac{c_1}{z} + \frac{c_2}{z^2} + \dots + \frac{z}{(z-1)^2} \\ \text{or} \\ z^2(z-1)^{-2} &\left(d_0 + \frac{d_1}{z} + \frac{d_2}{z^2} + \dots \right)^{-1}, \end{aligned} \right\} \quad (136)$$

where $(x_n, a_n) = (X, A)z^n$, $z \equiv \exp(-i\omega \Delta t)$, and $d_0 + d_1 + d_2 + \dots = 1$. Using the methods of Ref. 112, we obtain a dispersion relation for simulation of cold plasma,

$$\epsilon = 1 + (X/A \Delta t^2) \omega_p^2 \Delta t^2 \times \sum_p \frac{\kappa(k_p) S_m^2(k_p)/K^2}{p} = 0, \quad (137)$$

where $k_p = k - p k_{\text{gr}}$, $k_{\text{gr}} = 2\pi/\Delta x$, and $-K^2, \kappa(k_p)$, and $S_m(k_p)$ are the Fourier transforms of d^2/dx^2 , d/dx , and the m th order particle spline—e.g.,

$$S_m(k) = [\sin(k \Delta x/2)/(k \Delta x/2)]^{m+1}. \quad (138)$$

For $\omega_p^2 \Delta t^2 \sum_p \kappa(k_p) S_m^2(k_p)/K^2 \gg 1$, solutions of Eq. (137) are identical to those described in Ref. 110 for $\omega_p^2 \Delta t^2 \rightarrow \infty$ and $k \Delta x \rightarrow 0$. For finite $\omega_p^2 \Delta t^2$, grid effects become increasingly important as k approaches $\pi/\Delta x$ (Ref. 112). When inconsistent spatial filtering or ad hoc spatial differencing is introduced and $-\nabla \cdot [\chi^{(r)} \nabla \delta \phi^{(r+1)}]$ in Eq. (133) is altered, the modified dispersion relation becomes

$$\epsilon = \beta \omega_p^2 \Delta t^2 \sum_p (\tilde{S} - \tilde{S}') k_p \kappa(k_p) S_m^2(k_p)/K^2, \quad (139)$$

where ϵ is as given in Eq. (137), $\beta = c_0$ or d_0^{-1} , and $(S - S')$ is the Fourier transform

of the consistent smoothing operator minus the inconsistently smoothed operator, possibly with simplified differencing. The right side of Eq. (139), which would ideally be zero, is similar in form to the numerator of the right side of Eq. (135).

AN example of the resulting change in linear dispersion is given as follows. An effective c_0 can be calculated by subtracting the right side of Eq. (139) from the left side, and is reduced for positive $(\tilde{S} - \tilde{S}')$. Reducing c_0 changes $\text{Re}(\omega/\omega_p)$, but not $\text{Im}(\omega/\omega_p)$, for $\omega_p^2 \Delta t^2 \ll 1$ and lowers the $\omega_p^2 \Delta t^2$ instability threshold.¹¹⁰ Quantitative predictions of Eqs. (137) and (139) have been confirmed in one-dimensional simulations. We have extended our analysis of linear dispersion to the effects of using a finite initial guess for $E_{n+1}^{(0)}$.

Our analysis of inconsistent smoothing and ad hoc spatial differencing, including code experimentation to study their effects, indicates that they can produce substantial changes in convergence and linear dispersion. Code experience in one- and two-dimensional applications at very large $\omega_p^2 \Delta t^2$ has been encouraging (see Ref. 111 and the preceding article). Analysis, as described here, continues to guide both algorithm design and code performance.

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LASNEX and Atomic Physics

Introduction

The complexity of laser fusion target physics requires that most target design and analysis of experiments be carried out using sophisticated computer simulations. Our continued development of the LASNEX computer code consists of analyzing physical processes, developing numerical models, and packaging the resulting algorithms so that they work together consistently and can be used easily. In 1982, we made important improvements in several LASNEX models, including those for atomic physics, thermonuclear burn, and radiation transport. We summarize our improvements here

and present greater detail in the following four articles

LASNEX Upgrades. The degree of ionization of the target material is the single most important atomic-physics parameter. We have incorporated a Thomas-Fermi ionization model into LASNEX, resulting in a much improved degree of ionization for compressed matter. In addition, we have used the Thomas-Fermi atomic model to scale our equation-of-state tables so that they can be used for continually changing material compositions.

The degree of ionization of laser fusion target material also strongly affects the energy-loss rate of suprathermal electrons. We have modified our electron-transport routines to account for this, and have found an increase in the fast-electron range in high-Z matter. We have also incorporated suprathermal-electron bremsstrahlung formulas that account for partial ionization; thus, LASNEX now produces the correct amount of bremsstrahlung when energetic electrons are injected into cold matter.

We are investigating the possibility of using spin-polarized DT fuel to enhance thermonuclear reaction rates and lower the ignition threshold. It appears that if the fuel can be polarized, it will remain polarized throughout the implosion and burn. With complete polarization, DT reaction rates would be increased by about 50%; this change would mean a reduction in reactor-driver energy requirements by a factor of 2 to 3. We have not yet modified LASNEX to account for spin polarization; we have, however, added several packages that aid in diagnosing thermonuclear burn. Our packages include multigroup gamma diffusion, neutron pinex, and spectral output.

Radiative energy flow is one of the principal means of coupling laser energy to an imploding target. Because of its importance, we have concentrated on improving the LASNEX radiation-transport models. The multigroup-diffusion radiation model now uses an iterative scheme to couple energy to matter; the result is a consistent matter temperature for all photon frequencies and an elimination of the error that used to dominate certain important radiation-flow problems. We have also modified the diffusion and Monte-Carlo radiation routines so that they can be used simultaneously in different parts of the same problem; the primary

impact of this change is saving the enormous amount of computer time that would otherwise be wasted doing Monte-Carlo simulations in optically thick regions. Finally, we now have a new, operational package that solves the two-dimensional radiation-transport equation using the discrete-ordinate method. This package will not suffer from statistical noise problems, which can affect Monte-Carlo results.

User Support. The LASNEX computer code is used by many people who are not involved with developing or maintaining the code. To allow them to work most efficiently, it is essential that the code be well documented and designed to minimize human effort and chances of error. In 1982, we completely rewrote and published the LASNEX user's guide. We also designed and implemented an interactive graphics language that will provide a concise and consistent medium for users to set up problems and analyze results.

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Nuclear Spin-Polarized Fuel in Inertial Fusion

The use of spin-polarized DT fuel has been proposed for magnetic confinement fusion.¹¹³ We are now examining the possibility of using spin-polarized DT fuel for ICF. The increased thermonuclear reaction rate of spin-polarized fuel would improve target performance and reduce the temperature required for target ignition. A simple energy-scaling model proposed by Brueckner¹¹⁴ implies a driver energy proportional to $(\sigma v)^{-3}$ for constant target gain (assumed to be determined by the reactor configuration). Even an incomplete spin-polarization giving a modest 25% increase in the fusion cross section would yield a 50% reduction in the size of the required laser driver. The specific numbers may be debated, but they indicate that spin-polarized ICF fuel offers significant advantages.

One crucial practical question is how to achieve nuclear spin-polarization in a target suitable for laser irradiation. It is also important to ask whether the target modifica-

tions required for polarizing the spins would themselves reduce target performance. We will not address these essentially technological questions, but will instead consider a more fundamental issue: whether an (assumed) initial spin-polarized state would survive the ICF conditions of target irradiation and implosion.

The spin moment, S_i , of an individual nucleus precesses under the influence of the local magnetic field according to

$$dS_i/dt = \gamma S_i \times B, \quad (140)$$

where $\gamma \simeq g 4800/(G\cdot s)$, with $g = 0.86$ for deuterons and 5.9 for tritons (Ref. 115), and $B = B[R_i(t), t]$ is the magnetic field at the position $R_i(t)$ of the nucleus. For order-of-magnitude estimates, we replace the vector equation by a scalar equation for the rotational phase-accumulation $\Delta\phi$, such that

$$\Delta\phi = \gamma \int B dt. \quad (141)$$

If the magnetic field occurs as a pulse that obeys $\Delta\phi \ll 1$, it is unable to depolarize the spins.

In the following discussion, we consider various magnetic fields that arise in ICF. In general, there are fields on three size-scales: atomic fields arising in collisions, microscopic fields associated with surface instabilities, and macroscopic fields that extend throughout the target.

Collisional depolarization during thermonuclear burn is discussed by Kulsrud et al.^{113,116} A plasma particle (an ion or electron) produces a large magnetic field either by its magnetic dipole moment, $B \simeq \mu/r^2$, or by Lorentz transform of its electric field, $B \simeq ve/cr^2$. This field acts over the collision duration, $\Delta t \simeq r/v$, where r is the classical or quantum distance of closest approach for ions or electrons, respectively. The depolarization cross sections are readily estimated using Eq. (141) and $\sigma \simeq (\Delta\phi)^2 \sigma_c$, where $\sigma_c \simeq \pi r^2 \log \Lambda$ is the Coulomb cross section. (This method of obtaining spin-flip cross sections is a classical equivalent of the Coulomb-Born approximation; in dense plasmas, $\ln \Lambda \simeq 1$, and it does not matter whether $\ln \Lambda$ is included in the cross section.)

Our simple classical estimates are plotted for tritium in Fig. 3-53, which also shows a similar estimate of the cross section for quadrupolar depolarization of deuterons (tritium has no quadrupole moment). The spin-orbit cross section is the same for incident electrons or ions. Both the spin-orbit and spin-spin cross sections are evaluated for tritium nuclei; the corresponding cross sections for deuterium are smaller. Because high-density ICF plasmas have small Coulomb logarithms,¹¹⁷ the spin-orbit cross section is about an order of magnitude smaller than those given by Ref. 113. For reasonable particle fluences, the cross sections are so small that collisional depolarization may be neglected.

We next consider micrometer-scale inhomogeneities in $B(r, t)$. One question is whether two reacting nuclei might experience sufficiently different $B[(R(i), t)]$ histories to acquire different spin directions. Using the ion-ion Coulomb cross section, we estimate the distance of ion diffusion in a time, t such that

$$d = 1.5 \mu\text{m} \cdot kT^{5/4} (t/\rho)^{1/2}, \quad (142)$$

where kT is in keV, t is in ns, ρ is in g/cm^3 , and $d < 0.1 \mu\text{m}$ for most plausible target-plasma conditions.

A field inhomogeneity would have to be extremely strong to produce significant dephasing by this mechanism. The damping time $\tau = 4\pi\sigma (L/c)^2$ for a $1\text{-}\mu\text{m}$ field gradient in a kilovolt plasma is ~ 20 ps. Therefore, it would be difficult to sustain field inhomogeneities large enough to depolarize reacting nuclei, except, perhaps, close to an unstable interface. Estimates of $\text{field } J_s$ generated by the Rayleigh-Taylor instability range up to several MG (Ref. 118), which would be strong enough to depolarize spins of nuclei close to an unstable fuel-pusher interface. This fuel region is not likely to burn in any case, however, because of contamination by pusher material.

We next consider nuclear spin-resonance as a depolarization mechanism. To address this issue, we consider a target subject to an initially uniform field, B_0 , and assume that the spins are initially polarized along B_0 . During implosion, the field is compressed; to estimate $B(t)$ we employ Kidder's ideal-

ized model of homogeneous, isentropic compression,¹¹⁹ which predicts a time-dependent uniform field $B(t) = A(t) B_0$. The area compression factor, $A(t) = 1/h(t)^2$, is also the fractional increase in fuel areal density, ρR . For the Kidder model,

$$h(t) = (1 - t^2/t_0^2)^{1/2}, \quad (143)$$

where t_0 is the total time until complete spherical collapse. This form applies only up to a final time, t_f , determined by the symmetry of the implosion. The final field and its time integral are

$$B(t_f) = \frac{B_0 t_0^2}{t_0^2 - t_f^2} \quad (144a)$$

and

$$\int_0^{t_f} B(t) dt = \frac{1}{2} B_0 t_0 \log \frac{t_0 + t_f}{t_0 - t_f}, \quad (144b)$$

respectively.

For an implosion that raises ρR by a factor of 100 with $t_0 = 1$ ns and $t_f = 0.995(t_0)$, we find that an initial field of 10 kG gives a large final field (~ 1 MG). Despite this large field, the phase integral, $\Delta\varphi$, remains small enough that the spins do not precess even one cycle. Assuming that this field is not exceeded by some other spontaneously generated field, there is no spin-resonance.

Finally, we must ask whether larger fields would be generated during the implosion

Fig. 3-53. Approximate cross sections for spin-depolarization compared with the DT thermonuclear cross section.

itself. There is no doubt that the energy density in ICF plasmas is large enough to support enormous magnetic fields, but a target implosion with good spherical symmetry will have no mechanism for generating a large magnetic field. For a slightly unsymmetric implosion, on the other hand, we estimate the magnetic-field generation by integrating the rate of field generation¹¹⁴

$$\frac{\partial \mathbf{B}}{\partial t} = \frac{c}{en} \nabla(kT) \times \nabla n, \quad (145)$$

where n is the electron density. The density-temperature gradients of the idealized isentropic-implosion model are purely radial so that the vector product is zero. An unsymmetric implosion may be modeled by using gradients from the isentropic flow solution and taking the vector product to be a constant fraction, $a \sim 1/10$, of the product of the absolute values ($|\nabla n| |\nabla kT|$). The field generated is then time-integrated along a particle track (the Lagrangian trajectory); the result is

$$B(r, t) \cong a F \left(\frac{ct}{R_0} \right) \left(\frac{kT}{c} \right) c, \quad (146)$$

where the shape factor, $F(r, t)$, is

$$F(r, t) = 3\beta \left(\frac{\beta r^2}{\beta r^2 + h^2 R_0^2} \right) \leq 3\beta. \quad (147)$$

Here R_0 is the initial target radius, r is the current particle radius, c_s is the initial sound speed, T_c is the current temperature at the target center, $h = h(t)$, and $\beta = 1/3(R_0/c_s t_0)^2$ describes the nature of the implosion. $\beta \ll 1$ is a slow implosion, appropriate to a gas-filled target, while $\beta \geq 1$ is a fast (sonic) or hollow-shell implosion. The formula predicts fields ~ 50 kG, too small to have any significant effect.

Our estimates indicate that an initial state of spin-polarization would have a reasonable likelihood of surviving implosion and fuel burn. Note that the analysis assumes that two essential conditions can be met

- The spins must be initially polarized without adversely perturbing the target design.
- Large magnetic fields produced directly by the laser must be kept out of the fuel either by screening the fuel with a conducting pusher or by adopting a favorable target-irradiation geometry.

Summary. Under ICF conditions, nuclear collisional-depolarization cross sections are reduced because of the smaller Coulomb logarithm. We do not encounter possible resonance-depolarization mechanisms (such as Doppler-shifted static-field inhomogeneities) for two reasons: the ions move diffusively, and, in any case, there is not enough time during ICF implosions for spin-resonance to occur. In these two respects, the ICF situation appears favorable for employing spin-polarization of DT fuel.

Author: R. M. More

Suprathermal-Electron Bremsstrahlung

The laser fusion concept is based on irradiating solid or liquid targets with a laser to compress them to very high densities. Very hot, suprathermal electrons are created by the laser-plasma interactions occurring at and below the critical density (e.g., by the Raman instability, $2\omega_{pe}$ instability, and resonance absorption). Due to their long range, suprathermal electrons can preheat the fuel and prevent the achievement of the very large compressions necessary for thermonuclear ignition and burn. Correctly simulating the suprathermal-electron physics has always been an important part of our work. This year, we have improved our modeling of the suprathermal-electron bremsstrahlung radiation and the resultant electron-energy loss.

For thermal electrons, the emission of bremsstrahlung radiation is normally calculated by application of the classical Kramers formula, which gives a good approximation for slow-moving electrons in hyperbolic orbits in the Coulomb field Z^*e/r of the ion effective charge, where Z^* is the ionization state. A modest quantum correction is required, and it is available in the form of the Karzas-Latter Gaunt factor.

For suprathermal electrons, however, a more accurate treatment is needed. Such electrons may have relativistic energies, so that they often penetrate deeply into the bound electron distribution, where they see an effective charge larger than Z^* . Hard x rays emitted by suprathermal electrons are an important diagnostic indicator for determining the number and energy of such

electrons—especially for plasmas produced by long-wavelength lasers, in which suprathermal electrons are more prominent.

For laser-beam interaction with low- Z exploding-pusher targets, producing a fully ionized plasma, suprathermal bremsstrahlung can be adequately calculated from the Bethe-Heitler formula,¹²⁰ which is derived by application of the relativistic Born approximation for collisions with a pure Coulomb field. In this approximation, the bremsstrahlung-emission cross section depends on plasma composition only through the factor Z^2 , where Z is the atomic number. In practice, one set of energy-dependent cross sections for photon production (and electron downscatter) is calculated at the generation of a problem. The Bethe-Heitler cross section is then adjusted to apply to the material in any zone by multiplying by the appropriate value of Z^2 .

These cross sections have recently been improved by incorporating a parameterized fit to the bremsstrahlung spectrum developed by Feng and Pratt.¹²¹ Their algorithm interpolates between cross sections for fully ionized and neutral atoms as follows

$$\sigma(Z, Z^*) = \left(\frac{Z^*}{Z}\right)^3 \sigma(Z, Z) + \left[1 - \left(\frac{Z^*}{Z}\right)^3\right] \sigma(Z, 0), \quad (148)$$

where $\sigma(Z, Z^*)$ is the cross section for material with atomic number Z at ionization state Z^* . Feng and Pratt compared their results from Eq. (148) with their ionic data obtained by detailed quantum-mechanical calculations and found the error to be no more than 15%, except for Mo and W at 1 keV.

In Feng and Pratt's formulation, the Coulomb or fully ionized cross section $\sigma(Z, Z)$ is chosen from among four limiting forms according to the values of three parameters

- Z = nuclear charge of the fully stripped ion,
- T_1 = incident electron's kinetic energy (units of mc^2),
- k/T_1 = fraction of initial energy transferred to photon of energy k .

The energy conservation equation for the interaction is

$$T_1 = k + T_2, \quad (149)$$

where T_2 is the kinetic energy of the scattered electron (again in units of mc^2). The cross sections are split into the classical approximation, which applies when

$$T_1^{1/2} < 0.00737 \left(1 - \frac{k}{T_1}\right) Z \quad (150)$$

and the Born-Elwert approximation applies for higher incident kinetic energies. The classical approximation uses the expansions of Florescu and Costescu.¹²² The break point between their hard- or soft-photon approximation is

$$\frac{k}{T_1} = 116.28 \frac{T_1^{1/2}}{Z}. \quad (151)$$

The high-energy approximation recommended for fully stripped cross sections is the simplified nonrelativistic Born-Elwert equation,¹²³ applied when

$$T_1 \leq 4.54 \times 10^8 Z^4 = 2\alpha^4 Z^4, \quad (152)$$

where α is the fine-structure constant.

The Bethe-Heitler-Elwert (BHE) equation is used in the opposite (high energy) limit. It is the Bethe-Heitler formula corrected by the Elwert factor.¹²³

Screening due to atomic electrons causes the potential energy of an atom to drop off faster than $1/r$. In the Thomas-Fermi model, the potential energy is approximated by¹²⁴ $(Zc^2/r) \exp(-Q_s r)$, where Q_s is the inverse of the atomic radius, given by $Z^{1/3}/192$ in units of inverse Compton wavelengths.

Feng and Pratt recommend replacing Q_s by an effective value three times as large, giving more screening. The enhanced screening is then used with the screened Elwert-Born form-factor approximation to obtain the neutral cross section, which is a function of only T_1 , k , and Z —just like the fully ionized cross section.

The new bremsstrahlung cross sections have a more complex dependence on the nuclear charge than the simple Z^2 dependence of the Bethe-Heitler approximation. Therefore we evaluate and store, at problem generation time, the cross sections for all possible bremsstrahlung transitions determined by the structure of the

suprathermal-electron bins and the photon groups in both the neutral and fully ionized limits for each material. At each time step, the photon-production and electron-downscatter cross sections for a LASNEX zone are obtained by combining the corresponding neutral and Coulomb cross sections for each material in the zone, according to Eq. (148) and the current Z^* , and then averaging over all the materials present, weighted by the ion number density.

A simple calculation illustrates the effect of these new cross sections. We modeled a thin, gold disk with a source of 150-keV electrons present with LASNEX using both new and old bremsstrahlung cross sections. Figure 3-54 plots the photon-output spectra from our calculations. The old model calculates too few photons, especially in the range of high photon energy.

Authors: J. A. Harte and Y. T. Lee

Suprathermal-Electron Energy Loss in Partially Ionized Matter

Issues in atomic physics arise in every aspect of the dynamics of high-density inertial-fusion plasmas, especially for targets that contain materials of moderate ($Z = 10$ to 15) to large ($Z \approx 30$) atomic number. As we continue to extend our understanding of atomic processes, we have made significant improvements in the ability of our LASNEX hydrodynamics code to describe the behavior of real materials, especially at plasma temperatures for which the atomic ionization remains incomplete. For heavy atoms, partial ionization occurs at all temperatures of practical interest.

Fig. 3-54. Photon-output spectrum from a thin, gold disk with a source of 150-keV electrons present, calculated using Bethe-Heitler cross sections and Feng and Pratt's parameterized cross sections.

Various workers at LLNL have commented on the discrepancy between LASNEX calculations of hot-electron x-ray production and well established experimental data taken on cold, solid targets.¹²⁵ During 1982, we developed a practical solution to this difficulty.

Analysis. Suprathermal electrons (with energies of 10 to 500 keV) are produced in the laser-plasma interaction and lose energy by Coulomb interactions with bound and free electrons of the target plasma. As the suprathermals slow down, they produce energetic x rays by scattering from ion cores.¹²⁶ This suprathermal bremsstrahlung is an important practical diagnostic of the hot-electron number and energy.

The fraction, f , of the initial energy E_0 of a hot electron converted into hard x rays is known to be roughly

$$f \approx ZE_0/(900 \text{ MeV}) \quad (153)$$

for electrons stopping in cold matter.¹²⁷ This fraction is sensitive to two atomic quantities: the bremsstrahlung cross section, and the atomic stopping power (dE/dx) that determines the path length, or range, that an electron travels before stopping. The fraction f is an average of the ratio of these quantities: a longer electron range would allow the electron to produce more bremsstrahlung before stopping.

The bremsstrahlung-emission cross sections used by LASNEX now incorporate parameterized fits to quantum-mechanical calculations (see the previous article). This change alters the spectrum of hard x rays while having little effect on the total emission fraction.

The electron energy loss is approximately given by

$$-\frac{dE}{dx} = \frac{4\pi e^4}{mv^2} (n_b L_f + n_f L_b) \quad (154)$$

where e , m , and v are the electron charge, mass, and velocity, respectively; n_b and n_f are the number densities of bound and free electrons; and L_b and L_f are the stopping numbers for bound and free electrons.

In recent years, we used LASNEX to calculate dE/dx by ignoring electron binding, i.e., by taking $n_b = 0$ and n_f equal to the total electron density. This approximation is reasonable in a hot plasma but not very accurate for cold matter. To improve on this

approximation, however, we require a practical expression for L_b for arbitrary partially-ionized target atoms.

Adopting a result due to Bethe,¹²⁸ we get

$$L_b = \log \left(\frac{mv_0^2}{2\bar{I}(Z, Z^*)} \sqrt{\frac{e}{2}} \right), \quad (155)$$

where e is the natural base of logarithms. In Eq. (155), the essential atomic property is the mean excitation-ionization potential, $\bar{I}(Z, Z^*)$, for ions of nuclear charge Z and ionization state Z^* . This quantity is important for small ion charges ($Z^* \ll Z$), where the bound-electron stopping is dominant.

For $\bar{I}(Z, Z^*)$ we have a few quantum-mechanical calculations performed by G. McGuire of Sandia National Laboratory,¹²⁹ and there are experimental data for many neutral atoms ($Z^* = 0$). In addition, Thomas-Fermi theory gives approximate values for $\bar{I}(Z, Z^*)$ for all atoms and ions.¹³⁰ We have fit the results of the Thomas-Fermi theory with an approximate formula,

$$\bar{I} = \frac{a Z \exp \left[1.294 \left(\frac{Z^*}{Z} \right)^{(0.72 - 0.18 Z^*/Z)} \right]}{\sqrt{1 - \frac{Z^*}{Z}}}. \quad (156)$$

The numerical coefficient, a , is predicted to be 6.54 eV in the Thomas-Fermi theory, but experiments (with $Z^* = 0$ and $Z > 10$) and McGuire's calculations are both best fit by taking $a = 10$ eV (see Fig. 3-55).

- Electrons are confined
- Radiation escapes
- f = fraction of incident electron energy escaping as hard x rays ($E > 10$ keV)

Fig. 3-55. Quantum-mechanical calculations of the mean excitation-ionization potential, $\bar{I}(Z, Z^*)$, show good agreement with Thomas-Fermi theory.

Fig. 3-56. LASNEX now includes atomic correction of suprathermal-electron energy loss, which improves our calculation of hard x-ray signals.

With this improved stopping power, we find a significant increase in the calculated range of hot electrons and as much as a two-fold increase in the bremsstrahlung-emission fraction, f (see Fig. 3-56).

Conclusions. Although it is true that hot electrons are better able to penetrate pushers or other shields with the corrected dI/dx formula, it would be an oversimplification to suggest that this change will always reduce the predicted performance of ICF targets. In many cases, the number of hot electrons has been inferred from the bremsstrahlung signal; the change in f described here will allow us to replicate observed x-ray signals using target calculations that assume fewer hot electrons and, hence, reduced preheat levels.

The new LASNEX stopping formula gives greatly improved agreement with both experiments on cold matter and with other (SANDYL) calculations of hot-electron bremsstrahlung. Further improvements may

include incorporating the effects of high-density plasma conditions on the coupling of energetic electrons to partially ionized target material.

Authors: J. A. Harte and R. M. More

Fully Implicit Solution of Coupled Multigroup Radiation-Diffusion and Electron-Temperature Equations

We have implemented a new algorithm for the solution of the coupled multigroup radiation-diffusion, electron-temperature equations in LASNEX. In our previous "partial temperature" scheme, large swings of the intermediate values of T_e during a single time step often led to nonphysical results, necessitating an additional control on time-step size to limit the swings. Furthermore, errors were often worst in steady-state situations. Since the new scheme iterates to an implicit radiation distribution

and T_e , it requires no limit on the time step, and it relaxes to the correct answer in steady state.

The new implicit multigroup algorithm is related to a scheme of Axelrod and Dubois.¹³¹ It may be thought of as a form of alternating-direction implicit iteration, in which one direction corresponds to photon energy (group number) and the other to spatial location (zone number). Each step of the iteration consists of two substeps. In the first, an explicit guess for the T_e field determines the source term for the implicit diffusion equation solved for each group. In the second, the off-diagonal diffusion terms are treated explicitly, and the electron temperature is solved for algebraically. A scheme similar to this second splitting has been employed in studies of LTE model atmospheres.¹³²

More precisely, the iteration is composed of two regular splittings¹³³ performed alternately. The vector of unknowns is

$$\mathbf{x} = (f_{11}, f_{21}, f_{31}, \dots, T_1, f_{12}, f_{22}, f_{32}, \dots, T_2, \dots)^T, \quad (157)$$

where f_{ij} is the radiation energy density per unit frequency interval in group i and zone j , and T_j is the electron temperature in zone j . The first splitting advances the unknowns from iteration level (k) to $(k + 1/2)$ such that

$$\mathbf{M}_1 \mathbf{x}^{(k+1/2)} = \mathbf{b} - \mathbf{N}_1 \mathbf{x}^{(k)}. \quad (158)$$

The sparsity patterns of $\mathbf{M}_1$ and $\mathbf{N}_1$ are indicated in Fig. 3-57 for a one-dimensional problem with three photon groups and three zones. In this splitting, couplings between spatial zones are treated implicitly (they appear as nonzero elements of $\mathbf{M}_1$), while the intergroup couplings via T_e are explicitly contained in $\mathbf{N}_1$. We solve only for the f_{ij} and not the T_j , since the latter are not needed for the second splitting. Note that no guess for f_{ij} is needed in this step if the diffusion equations are solved by a direct method; since in two dimensions we solve them using an ICCG iteration,¹³⁴ we find it useful to retain an updated guess. The user who wishes to do so can save storage by inhibiting retention of this latest-iterate f_{ij} at the expense of some run time.

Fig. 3-57. First splitting for a one-dimensional problem with three photon groups and three zones: sparsity patterns of $\mathbf{M}_1$ and $\mathbf{N}_1$.

The second splitting advances the unknown to $(k + 1)$,

$$\mathbf{M}_{II} \mathbf{x}^{(k+1)} = \mathbf{b} - \mathbf{N}_{II} \mathbf{x}^{(k+1/2)}, \quad (159)$$

implicitly treating couplings between photon groups (see Fig. 3-58). Here, the off-diagonal part of the diffusion operator acting on f_{ij} from the first substep is treated explicitly, and only the T_i are solved for, since the f_{ij} are not needed as input in substep (I). Compton scattering and electron and ion heat conduction are handled by operator splitting, as before.

We introduce the notation $f_i \equiv (f_{i1}, f_{i2}, \dots)$, $T \equiv (T_1, T_2, \dots)$, etc. The equation solved in splitting (I) for each photon group i is then

$$\frac{f'_i - f_i}{\Delta t} = \mathcal{D}_i f'_i + \gamma_i \left[B_i(\tilde{T}) + (T' - \tilde{T}) \frac{\partial B_i(\tilde{T})}{\partial \tilde{T}} - f'_i \right], \quad (160)$$

where

$$\gamma_i \equiv c \rho \kappa_i, \quad (161)$$

f_i is the radiation-energy-density field for group i at the beginning of the time step, f'_i is the new iterate at level $(k + 1/2)$ for which we solve, Δt is the time step size, $\mathcal{D}$ is the diffusion operator for frequency group i (a function of f_i or f'_i to incorporate flux limitation), c is the speed of light, ρ is the mass density, κ_i is the opacity to photons in group i , B_i is the Planck function, T' is the most recent iterate of the electron temperature,

$$\tilde{T} = \alpha_e T' + (1 - \alpha_e) T \quad (162)$$

is the temperature about which the Planck function is linearized, α_e is a user-specified implicitness parameter for $\tilde{T}$, and T is the electron temperature at the beginning of the step. Equation (160) is solved in much the same manner as the corresponding equation in the partial-temperature method.

For splitting (II), we solve for an improved level $(k + 1)$ iterate T'' ; we also introduce a corresponding f'' , for which we

need never solve explicitly. The diffusion operator $\mathcal{D}_i$ is written as the sum of a diagonal operator $\mathcal{D}_{Di}$ and an off-diagonal one $\mathcal{D}_{ODi}$, and the diffusion equation becomes

$$\frac{f''_i - f'_i}{\Delta t} = \mathcal{D}_{Di} f''_i + \mathcal{D}_{ODi} f'_i + \gamma_i \left[B_i + (T'' - \tilde{T}) \frac{\partial B_i}{\partial \tilde{T}} - f''_i \right]. \quad (163)$$

Using Eq. (160) to express $\mathcal{D}_{ODi} f'_i$ in terms of $\mathcal{D}_{Di} f'_i$ and then solving Eq. (163) algebraically for f''_i yields

$$f''_i = f'_i + (T'' - T') G_i, \quad (164)$$

where

$$G_i = \frac{\gamma_i \frac{\partial B_i}{\partial \tilde{T}}}{\frac{1}{\Delta t} - \mathcal{D}_{Di} + \gamma_i} \quad (165)$$

Fig. 3-58. Second splitting: sparsity patterns of $\mathbf{M}_{II}$ and $\mathbf{N}_{II}$.

The denominator of G_i is just the diagonal coefficient used in the ICCG solution of splitting (I).

The equation for the time evolution of the electron temperature is

$$C_v \frac{T'' - T^*}{\Delta t} = - \sum_i \Delta \nu_i \gamma_i \times \left[B_i + (T'' - \tilde{T}) \frac{\partial B_i}{\partial T} - f_i'' \right], \quad (166)$$

where C_v is the electron heat capacity, T^* is the modified T field resulting from the action of other (nonradiation) sources, and $\Delta \nu_i$ is the frequency interval associated with group i . Using Eq. (164) for f_i'' and solving Eq. (166) for T'' yields

$$T'' = \frac{S_1}{S_2}, \quad (167)$$

where

$$S_1 = C_v \frac{T^*}{\Delta t} + \sum_i \Delta \nu_i \gamma_i \times \left[f_i' - G_i T' - B_i + \tilde{T} \frac{\partial B_i}{\partial T} \right] \quad (168)$$

and

$$S_2 = \frac{C_v}{\Delta t} + \sum_i \Delta \nu_i \gamma_i \left(\frac{\partial B_i}{\partial T} - G_i \right). \quad (169)$$

The sums S_1 and S_2 are accumulated as we pass through the groups in splitting (I). The improved values T'' are used to replace the T' values for the next iterate, using the same storage.

To achieve algebraic energy conservation, we could arrange that any change in the radiation energy be accompanied by a cor-

responding opposite change in the energy associated with the electron temperature field actually produced as output. To this end, we can explicitly compute a temperature T_{expl} from the implicit f_i' at each iteration, using

$$C_v \frac{T_{\text{expl}} - T^*}{\Delta t} = - \sum_i \Delta \nu_i \gamma_i \times \left[B_i + (T' - \tilde{T}) \frac{\partial B_i}{\partial T} - f_i' \right]. \quad (170)$$

However, since T_{expl} is generally a much poorer estimate of the true temperature than is T'' , we choose T'' to be the final result and return the discrepancy

$$\sum C_{vi} (T_{\text{expl}} - T'')_i \quad (171)$$

in a new row of the energy-balance diagnostic. This discrepancy can be made arbitrarily small by choice of a stringent relative-convergence test (at present, applied to the L_2 -norm of the difference between iterates).

Our new method is more efficient than the partial-temperature algorithm on many problems and is always more reliable and robust. Convergence is rapid when the characteristic time scales for transport (τ_i) and coupling (τ_c) differ greatly, or when Δt is less than the maximum of τ_i , τ_c , or $C_v \tau_c / C_R$, where C_R is the radiation heat capacity.

We are considering various techniques to speed the rate of convergence of the iteration, a refinement that is expected to be useful only on unusually stiff problems. On most problems, our present procedure achieves convergence to one part in 10^4 within 10 iterations.

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The automated sphere mapping (ASM) system is used to provide graphic displays of the exterior of spherical objects with resolution capabilities of $0.2\ \mu\text{m}$ in height and $2.0\ \mu\text{m}$ in lateral dimensions. The ASM uses a phase-reference system with a computer-controlled sphere manipulator to provide the automated data collection and display.

Section 4

Target Fabrication

Target Fabrication

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Target Fabrication

Introduction

I. F. Stowers

The Fusion Target Fabrication program's primary responsibility is the construction of targets for current research. Our effort, of necessity, entails constant development of fabrication technology, which includes synthesis and handling of target materials, assembly, inspection, and characterization. The required precision and microscopic scale of inertial confinement fusion (ICF) target fabrication has led to a significant investment in manpower and facilities. Program personnel now include materials scientists, physicists, chemists, mathematicians, engineers, and a large staff of skilled mechanical, electronics, and chemical technicians. The new fusion target development facility (FTDF) was completed during 1982, for the first time housing all technical personnel and equipment under one roof.

Historically, glass fuel balls have contained only the DT fuel for implosion experiments. More recently, we have been requested to deliver DT fuel balls also containing small amounts of tracer gases such as Ar, Br, or Kr. We have developed an ion-drilling technology that will allow us to place virtually any tracer material in glass balls.

During 1982, we developed an analytical model to better understand the microsphere-formation process occurring in a drop tower. The model predicts the final diameter and wall thickness of spheres when given the temperature profile of the furnace and the initial liquid content of the glass-forming materials.

We completed construction of the automated sphere mapping (ASM) system that has been under development for several years, and we have successfully measured the topography of a nominally smooth sphere. The ASM instrument maps the full surface area of a sphere with high accuracy. In the final articles of this section, we discuss target assembly and specifically address soft x-ray laser targets constructed for testing on Novette and an improved technique for hemishell inspection.

Fuel Container

Introduction

The major objective of our fuel-container effort is to develop and maintain the technology necessary for the fabrication and filling of fuel capsules. In general, the processes we develop must be flexible enough to handle the various material and size requirements of the ICF program. Our daily production rate must be quite high (several thousand target-quality spheres per day) to meet the demands of both target production and other R&D efforts within the Target Fabrication program.

During 1982, our unclassified activities focused on three areas:

- Routine production of target-quality batches of glass microspheres to support upcoming Novette activities.
- Development of a computer model to simulate drying and microsphere formation in the droplet column.
- Design and start of fabrication of an ion-beam drill for drilling submicrometer holes in target-sphere materials.

The work leading to the current method of routine glass-sphere production has been discussed in depth elsewhere.¹ In the articles that follow, we discuss our work on computer modeling and development of the ion-beam drill.

Author: J. H. Campbell

Modeling Microsphere Formation in a Droplet Drying Column

The method we commonly use for the preparation of glass microspheres involves drying and fusing droplets of glass solution in a vertical oven.¹⁻³ This method has also been used by Kool et al.⁴ at KMS Fusion, Inc. (KMSF), for making polymeric shells. The initial step involves the drying of liquid droplets containing dissolved solids of either inorganic glass formers or, in the case of the KMSF work, organic polymers. If proper drying conditions are used and if the material properties meet certain criteria, a hollow sphere is produced at the end of the drying step. In the case of glass microspheres, the dried, hollow spheres are subsequently fused at high temperatures (1500°C) into the final glass product.

We must optimize several experimental parameters to produce target-quality shells from any given material. In the past, we used a largely empirical approach to determine the proper tower operating condi-

tions and droplet composition. Although successful, the empirical method was time-consuming and manpower-intensive.

To help guide and interpret our current experimental work on glass microspheres and other materials, we developed a simple one-dimensional numerical model to simulate the droplet drying and initial sphere-formation process. In brief, the model treats the simultaneous heat and mass transport to and from a liquid droplet containing some level of dissolved solids. The model is based largely on earlier investigations of the fundamental aspects of spray drying (for example, see Refs. 5 and 6).

Model Description. A schematic representation of the drying process is given in Fig. 4-1. The first step treated by the model is evaporation during the droplet deceleration period, i.e., between leaving the droplet generator and reaching terminal velocity. Significant evaporation can occur during the deceleration period if the gas temperature is above 150°C.

Having reached terminal velocity, the droplet enters the period of drying during which the moisture flux from the droplet surface is approximately constant [Fig. 4-1(a)]. During this first stage of drying, the droplet temperature remains at the adiabatic saturation temperature, and a dynamic equilibrium is established between the rates of heat ($\dot{Q}$) and mass transport ($\dot{M}_{H_2O}$). As evaporation continues, the droplet concentration reaches a critical value. The evaporation rate from the particle surface is equivalent to the sensible heat supplied by the surrounding gas. Eventually the moisture content of the particle diminishes to the so-called critical concentration⁶ and a surface film forms [Fig. 4-1(b)]; this initiates the next drying step treated by the model.

At this point the evaporation rate begins to drop due to the added resistance of the film to mass transfer. As a consequence, the temperature of the droplet begins to increase, and ultimately reaches the solution boiling point. The temperature of the droplet remains at the boiling point as the balance of the free moisture evaporates, forming a hollow shell [Fig. 4-1(c)]. This period of drying is perhaps the most critical. Whether an unruptured, hollow sphere forms depends strongly on the type and strength of the surface film formed and its permeability to water vapor and the furnace purge gas. The hollow sphere, containing

Fig. 4-1. Schematic representation of the different stages in the drying process for a droplet of alkali silicate glass solution in a droplet column.

gas at pressure P_b , begins to grow and ultimately produces a large, thin-walled, hollow sphere of unfused silica glass containing a gas at pressure P_b^1 . Finally, once all the free water has been removed from the droplet, the shell heats rapidly to the temperature of the furnace [Fig. 4-1(d)].

The basic equations used to calculate the droplet velocity and rates of heat and mass transfer are summarized in Table 4-1. Also given are brief physical descriptions of the corresponding drying conditions. The variables and dimensionless numbers used in our calculations are defined in Table 4-2.

The evaporation rate for a droplet in a moving gas stream is simply a product of the mass-transfer coefficient, h_m , the droplet surface area, and the difference in vapor pressure across the gas boundary layer

$$\frac{dm}{dt} = h_m \pi D_p^2 \frac{\Delta p_v}{RT} M_w \quad (1)$$

A similar equation also describes the rate of heat transport

$$\frac{d\dot{Q}}{dt} = h_q \pi D_p^2 (\Delta T) \quad (2)$$

where the driving force is the temperature drop, ΔT , across the boundary layer, and h_q is the heat-transfer coefficient.

Both heat- and mass-transport coefficients can be represented in terms of the dimensionless Nusselt (Nu) and Sherwood (Sh) numbers

$$Nu = h_q \frac{D_p}{k_f} \quad (3)$$

and

$$Sh = h_m \frac{D_p}{D_v} \quad (4)$$

Table 4-1. Summary of drying stages and governing equations used in model calculations.

Drying stage	Physical description	Governing equations
1a	Unsaturated solution; droplet deceleration to terminal velocity; droplet at adiabatic saturation temperature.	<p>Deceleration: $-\frac{dv}{dt} = \frac{3 C_d \rho_g v^2}{4 D_p \rho_p}$</p> <p>Mass transfer: $\frac{dm}{dt} = \pi D_p D_v Nu \left(\frac{\Delta p_v M_w}{RT} \right)$</p> <p>Heat transfer: $\frac{dQ}{dt} = \pi D_p k_f Sh (\Delta T)$</p>
1b	Unsaturated solution; droplet at terminal velocity and adiabatic saturation temperature.	<p>Velocity: $v_T = \left[\frac{4 (\rho_p - \rho_g) g^2}{225 \rho_g \mu} \right]^{1/3} D_p ; Re > 0.4$</p> <p>$= \frac{g(\rho_p - \rho_g) D_p^2}{18\mu} ; Re < 0.4$</p> <p>Mass transfer: } Same as step 1a</p> <p>Heat transfer: }</p>
2	Saturated solution; skin-formation occurs; mass-transfer rate decreases.	<p>Velocity: same as in step 1b</p> <p>Mass transfer: $\frac{dm}{dt} = h_o \pi D_p^2 \frac{(\Delta p_v M_w)}{RT}$</p> <p>Heat transfer: same as in step 1a</p>
3	Membrane encloses vapor and liquid phases. Temperature $\approx$ boiling point of solution.	<p>Velocity: same as in step 1b</p> <p>Mass transfer: same as in step 2</p> <p>Heat transfer: $\frac{dQ}{dt} = h_q A_{eff} (\Delta T)$</p>
4	Liquid evaporation complete; particle heats to furnace temperature.	<p>Velocity: same as in step 1b</p> <p>Mass transfer: $dm/dt = 0$</p> <p>Heat transfer: same as in step 1a</p>

Fuel Container

Substitution of the dimensionless groups into Eqs. (1) and (2) leads to the expressions given in Table 4-1.

Numerous experimental measurements of heat- and mass-transfer rates under different stream conditions have established the following well known empirical correlations for the Nusselt and Sherwood numbers in terms of other dimensionless groups^{6,7}

$$\text{Nu} = 2 + 0.6 (\text{Re})^{1/2} (\text{Pr})^{1/3} \quad (5)$$

and

$$\text{Sh} = 2 + 0.6 (\text{Re})^{1/2} (\text{Sc})^{1/3} \quad (6)$$

The Reynolds, Prandtl, and Schmidt numbers in Eqs. (5) and (6) are defined,

Table 4-2. Nomenclature.

A_{eff}	Effective surface area for heat transfer (cm^2)
C_d	Drag coefficient (dimensionless)
C_p	Heat capacity ($\text{cal/g} \cdot \text{K}$)
D_p	Droplet diameter (cm)
D_v	Vapor diffusivity (cm^2/s)
g	Gravitational constant (cm/s^2)
h_b	Mass-transfer coefficient through barrier (cm/s)
h_m	Mass-transfer coefficient (cm/s)
h_o	Overall mass-transfer coefficient (cm/s)
h_q	Heat-transfer coefficient ($\text{cal/cm}^2 \cdot \text{s} \cdot \text{K}$)
K	Membrane permeability [$\text{cm}^3(\text{STP}) \cdot \text{cm}/\text{cm}^2 \cdot \text{s} \cdot \text{atm}$]
k_f	Thermal conductivity of gas boundary layer ($\text{cal}/\text{cm} \cdot \text{s} \cdot \text{K}$)
l	Membrane thickness (cm)
m	Mass (g)
$\dot{M}_{\text{H}_2\text{O}}$	Rate of water vapor loss for unsaturated solution (g/s)
$\dot{M}_{\text{H}_2\text{O}}$	Rate of water vapor loss through film (g/s)
M_w	Gram molecular weight (g/g-mole)
p_a, p_b	Gas pressure in hollow sphere (atm)
Q	Energy (cal)
$\dot{Q}$	Rate of energy transport (cal/s)
R	Gas constant ($\text{cm}^3 \cdot \text{atm}/\text{g} \cdot \text{mole} \cdot \text{K}$)
T	Temperature (K)
t	Time (s)
v	Droplet velocity (cm/s)
v_T	Droplet terminal velocity (cm/s)
ΔH_p	Latent heat of vaporization (cal/g)
Δp_f	Vapor pressure driving force across film membrane (atm)
Δp_v	Vapor pressure driving force across boundary layer (atm)
μ	Viscosity of drying medium (P)
ρ_g	Density of drying medium (g/cm^3)
ρ_p	Particle density (g/cm^3)
	Dimensionless numbers:
Nu	Nusselt number
Pr	Prandtl number
Re	Reynolds number
Sc	Schmidt number
Sh	Sherwood number

respectively, as

$$\text{Re} = \frac{D_p v \rho_g}{\mu} \quad (7)$$

$$\text{Pr} = C_p \frac{\mu}{k_f} \quad (8)$$

and

$$\text{Sc} = \frac{\mu}{D_v \rho_g} \quad (9)$$

Using Eqs. (5) through (9), it is now possible to calculate the heat- and mass-transfer rates from measured physical properties of the system.

During the initial drying stages (before a film forms around the droplet), the rates of heat and mass transport are in dynamic equilibrium. Under these conditions, the particle temperature remains at the adiabatic saturation value, and the rate of evaporation is simply

$$\frac{dm}{dt} = \frac{1}{\Delta H_p} \left(\frac{dQ}{dt} \right) \quad (10)$$

where ΔH_p is the latent heat of vaporization.

Once the droplet solution reaches the critical concentration and the film begins to form, the mass-transfer rate begins to drop dramatically due to the vapor barrier formed by the film. During this drying stage, the rate of mass transport is calculated using an overall mass-transfer coefficient, h_o , given by

$$h_o = \frac{h_m h_b}{h_m + h_b} \quad (11)$$

where h_b is the mass-transfer coefficient across the barrier membrane of thickness, l , given by

$$h_b = \frac{K \Delta p_i}{l} \quad (12)$$

As the liquid continues to evaporate, a gas bubble begins to form inside the film membrane (provided the membrane permeability is low enough). As a consequence, the effective area for heat transfer, A_{eff} , is equivalent to the area of the spherical segment containing the remaining liquid (see

Table 4-1 and Fig. 4-1). Mass transfer, however, continues over the entire droplet surface.

Model Results. Input to the model includes the column length and temperature profile; purge-gas composition and flow rate; and initial temperature, size, composition, and velocity of the droplet. Basic physical and transport properties of the individual components are also input as functions of temperature. These properties are used with well-established empirical correlations to estimate the effective heat-, mass-, and momentum-transfer rates in the system; the critical concentration for film formation was determined from independent laboratory experiments on bulk glass solutions.

Output from the calculations includes the droplet temperature, size, velocity, and composition as functions of time and distance down the column. Typical examples of output for droplet drying in Ar are plotted in Fig. 4-2 for a 230- μm -diam droplet of our standard alkali silicate glass solution containing 18 wt% solids. The furnace temperature over the first 2 ft of the column is $\sim 260^\circ\text{C}$, decreasing slightly for the remaining length of the drying section.

The initial droplet velocity is about 500 cm/s as it leaves the generator. The velocity rapidly decelerates to terminal velocity over about the first half meter, as shown in Fig. 4-2. During this time, the droplet decreases from 230 to about 150 μm in diameter and remains at an adiabatic saturation

temperature of from 40 to 45°C. The temperature slowly increases as the solution concentration increases.

After the droplet has fallen roughly 1.2 m, the solution concentration reaches saturation, and a film begins to form. Figure 4-2 clearly shows the increase in droplet temperature and diameter. As the droplet begins to grow, the terminal velocity falls off dramatically.

As the surface film reduces the mass-transfer rate, the temperature of the droplet increases from the adiabatic saturation temperature to the solution boiling point. Once all the water has been evaporated, the sphere diameter stops changing and the temperature increases to that of the surrounding gas.

We have found that the results of our model calculations agree well with the experimental observations reported by Rosencwaig et al.¹ Both the column length needed to dry the spheres and the measured residence time of the droplets in the drying section agree with our calculations. We have measured the diameter of the hollow spheres produced at the end of the drying section of the column and found they range from about 1300 to 1450 μm with an average of 1385 μm (for an initial 230- μm droplet with 18 wt% solids). The model predicts a hollow-sphere size of 1300 μm .

We are continuing to use the model for designing and analyzing our droplet column experiments. It has been used also to

Fig. 4-2. Results from droplet drying calculations showing the diameter, velocity, and temperature of the droplet as a function of time.

Target and Material Measurement

evaluate the effects of droplet size, solvent systems, purge gas, and thermal profiles on the sphere-formation process. Finally, work is under way to extend the model to simulate the glass-fusion process that occurs after the drying section.

Authors: J. H. Campbell and J. Z. Grens

Major Contributor: J. F. Poco

Target and Material Measurement

Introduction

Both the quality of targets produced for inertial-confinement fusion experiments and the rate of progress in development of advanced targets depend directly upon our ability to perform measurements on the targets and the materials used in construction.

During 1982, we added new capabilities to the wide range of physical measurements and material analyses we routinely apply.

We completed construction of our automated sphere mapping device, and successfully measured sphere surfaces with a precision never before possible. In the following two articles, we describe these new developments in target characterization.

Author: W. G. Halsey

Automated Sphere Mapping

During 1982, we completed construction of our automated sphere mapping (ASM) system,⁸ the design of which has been thoroughly described elsewhere.⁹⁻¹¹ The ASM performed very well in start-up testing and preliminary applications; its implementation opens a new era in target-characterization accuracy and resolution.

The completed ASM, including the optics table, interferometer, and precision sphere rotator, is shown in Fig. 4-3. During final construction, we solved problems with the sphere rotator having to do with long-term drift; the mechanical system is now stable to $\sim 10\text{ }\mu\text{m}$ for 1 h of continuous operation. With the optical interferometer, we have

Fig. 4-3. Completed ASM system.

achieved excellent stability, 200-Å vertical (height) sensitivity, and a spatial (horizontal) resolution of $\sim 3 \mu\text{m}$. Timing errors in phase-shift detection currently limit data acquisition to 10^4 samples/s. Figure 4-4 shows sample data for a test sphere having a few defects in the $10\text{-}\mu\text{m}$ range.

Application of the ASM is illustrated in Fig. 4-5. In the two graphs, the solid lines represent the range of typical limits for sphere-surface defect heights; the shaded regions show the ranges for which various characterization techniques can detect the indicated defects. The contribution of the ASM to sphere characterization is apparent.

Author: W. G. Halsey

Major Contributors: A. Hernandez, G. F. Jacobson, A. L. Richmond, M. D. Saculla, R. L. Thombley, B. W. Weinstein, and K. R. Zeringue

radiography, microanalysis, and surface analysis.

With our newly installed x-ray diffraction machine, we can examine the crystal structure of target materials to assist in process development and optimization. X-ray diffraction provides information on material grain size, composition, and crystal orientation. To complement structure-of-material measurements, we have designed a system to measure thin-film strength. A new ellipsometer has also been installed to measure film thickness and refractive index.

By the end of the year, the FTDF computer system, a VAX-750, was installed and functioning as a control simulator for Novette experiments. Because the VAX allows for extraction of more information

Fig. 4-4. Sample ASM data.

Improved Measurement Capabilities

Our ability to provide timely measurements of target materials and dimensions is a key factor in the development and delivery of laser targets. Our analytical capabilities must constantly be expanded and upgraded to keep pace with evolving target fabrication requirements. During 1982, we added several new measurement systems to complement our existing capabilities^{12,13} in optical and electron microscopy, interferometry,

Fig. 4-5. Application of ASM to sphere characterization: solid curve shows critical defect dimension.

from radiographs, we also installed a code to analyze information for the x-ray micro-radiography system.¹⁴ Other applications are under development.

Author: W. G. Halsey

Target Assembly and Microfabrication

Introduction

The requirements for producing targets were greatly reduced during 1982 since both the Argus and Shiva lasers had been dismantled in 1981. However, a number of targets were required for use in developing target-fabrication technology. We prepared several flat-disk targets of glass, parylene, and gold to obtain data for the ion-drilling project, and we constructed ball-on-film and ball-on-stalk targets to evaluate the properties of DT fuel in glass and coated glass shells.

In October, we initiated a major development effort to build a short-wavelength, low-energy x-ray laser cavity. As a test bed for determining the flashlamp to be used with the lasing cavity, we coated several candidate materials such as manganese, nickel, chromium, and iron onto thin-film polymeric substrates. The films were supported on plastic washers and mounted for irradiation by the Novette laser. Our efforts are described in the following article.

An improved measurement device for evaluating target hemishells was developed at Rocky Flats. This device, described in the final article in this section, is based on an air-spindle movement and white-light interference techniques; for target hemishells, it measures the surface contour, the surface finish, the true dimensions, and the deviation of contour.

Author: C. W. Watcher

X-ray Laser Targets

The construction of targets for x-ray laser experiments required development of new fixture and assembly techniques. The target—actually a lasing cavity—is a long, narrow slit measuring 14 by 0.05 mm. The

slit is bounded on two facing sides by metallic knife edges and on the perpendicular sides by thin, composite films of Formvar and a metallic element; the ends are closed with thin films of Formvar. The lasing cavity is supported by an aluminum housing (Fig. 4-6), which also provides a connection to the target positioner and the lasing gas supply.

The knife edges for the laser cavity are made from 60- μ m-thick gold foil, 1 cm wide by 1.4 cm long. Each knife blade has a 60° precision-lapped angle on one long edge; all lapping is done on a cast-iron lap using 0.1- μ m-thick diamond lapping paste. Typical finished blades have less than a 5- μ m radius on the finished edge and are straight within 3 μ m over the entire 1.4-cm length. One blade of each set has a slot, 200 μ m wide and 30 μ m deep, milled on the center line and perpendicular to the lapped edge. The slot provides the connecting passage between the lasing chamber and the external gas supply.

The thin substrate foils and output windows (Fig. 4-6) are prepared by techniques used routinely by manufacturers of integrated circuits. The Formvar film is coated onto a glass slide and removed by flotation on distilled water. The film is removed from the water using a wire-loop support; the wire loop supporting the film is then placed in a vacuum vapor-deposition coater where the appropriate flashlamp or window material is applied by RF sputtering or thermal vaporization.

The gold collimating pinholes are produced by machining a copper mandrel to the required dimensions (Fig. 4-7). The mandrel is then electroplated with gold to the desired thickness. The cylinders are scribed at the desired length and submerged in a nitric acid bath. When the supporting copper mandrel is removed, a cylindrical tube with a closed end having a small (100- μ m-diam) pinhole is produced (Fig. 4-6).

The vacuum holding fixture (Fig. 4-8) used to locate the knife blades is fabricated from a precision lapped brass block. Machined into the lapped surface are two rectangular grooves, each 1 mm deep, 1 cm long, and 0.7 cm wide. A 2-mm-diam hole intersects each groove from the bottom to provide a vacuum pumping port. The two rectangular grooves are located parallel to each other and are separated by a 1-mm

Fig. 4-6. Short-wavelength, high-energy x-ray laser cavity and components.

slot through the brass block; this slot allows light to pass through from below to aid in the blade-alignment procedure. Perpendicular to the ends of each of the rectangles is a line-bored 3-mm-diam round slot used to precisely locate the gap alignment spacer (Fig. 4-9) during the assembly procedure.

Prior to assembly, all pieces are deburred and cleaned in the ultrasonic cleaner using chloroform and alcohol. After cleaning, all parts are inspected using a binocular microscope. The first step in assembly is placing the finished knife blades on the vacuum holding fixture (Fig. 4-8). The gap between the blades is determined by sliding each blade lightly against a precision-machined spacer of the proper thickness (Fig. 4-9). As each piece is brought into nominal

Fig. 4-7. Pinhole mandrel.

Fig. 4-8. Vacuum holding fixture with one blade in place.

Target Assembly and Microfabrication

contact with the spacer, the vacuum valve is opened and the blade is firmly held in place. The tolerance on gap spacing and parallelism is $\pm 2 \mu\text{m}$. A bridge consisting of a gold foil measuring $300 \times 100 \times 5 \mu\text{m}$ is bonded across the gap at each end of the blade to maintain the gap spacing. A $0.15\text{-}\mu\text{m}$ -thick Formvar film is bonded over the blades to form one side of the pressure chamber; this film must be wrinkle and pinhole free. One half of the aluminum body is now attached to the assembly. After the assembly is turned over, the remaining gold bridges and Formvar film are bonded over the blades. The gas chamber is now complete, with the exception of the end caps and output windows.

Output windows for the high-energy target consist of $0.15\text{-}\mu\text{m}$ -thick Formvar films

mounted on an aluminum framework having a $100\text{-}\mu\text{m}$ -diam hole. Because the hole will serve as the laser output window, it must be accurately located on the axis of the knife blades so that any output from the lasing cavity remains unobscured. The cylindrical 3-mm -diam body containing the $100\text{-}\mu\text{m}$ -diam pinholes is now located on the same axis and attached to the aluminum body. The second half of the aluminum body is then placed over the assembly.

Each target is tested prior to any subsequent operations to ensure that it meets pressure and vacuum specifications. After testing, the flashlamp material is vacuum-deposited on the Formvar films. After deposition, the optical alignment aids, consisting of mirrors and gold washers, are attached, and the target is mounted to its bayonet base. The base serves as a mechanical mount in the target chamber and also serves as a gas feedline through which the proper mixture of lasing gases is fed into the cavity.

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Fig. 4-9. Detail of gap-alignment spacer.

Fig. 4-10. Schematic of the improved noncontact measuring device.

Improved Noncontact Measuring Device

We designed, fabricated, and used an improved noncontact measuring device to measure the physical characteristics of inertial-fusion target hemishell components. The noncontact nature of the device is important due to the fragile nature of target components. Our improvements permit concurrent inside and outside contour measurements and protect against outside thermal effects.

The improved noncontact measuring device is shown in Fig. 4-10. It uses two white-light interference microscopes as null positioners against master-reference standards. We use digital measuring probes to

measure the deviation from the standards, with a resolution of $0.1\text{ }\mu\text{m}$. The measuring device is operated by manipulation of the part to be inspected on a vertical-axis air-bearing spindle. Interferometric microscopes permit us to examine the surface finish and to measure surface defects concurrently with the contour. Any component that has a surface sufficiently reflective to provide interferometric fringes in the microscope can be measured along the axis of the interferometers. Components from $100\text{ }\mu\text{m}$ to 4 mm in diameter can be characterized. The accuracy of the device is $\pm 0.25\text{ }\mu\text{m}$ in contour.

Author: C. W. Hatcher

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Shown is a color-enhanced x-ray micrograph of the imploding core of a glass microballoon irradiated by the Shiva laser. The line profile shows the intensity at the location along the horizontal line.

Section 5

Fusion Experiments Program

Fusion Experiments Program

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Section Editor: P. H. Y. Lee

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Fusion Experiments Program

Introduction

L. W. Coleman

Our experimental program relies on state-of-the-art diagnostics and utilizes modern methods for timely and accurate reduction and analysis of data to interpret experiments. We report here on progress and accomplishments in these areas of our program in 1982.

Diagnostics

Fusion Experiments Diagnostics

With no experimental laser facilities active during 1982, we concentrated on the final reduction and analysis of Argus and Shiva data. In addition, we have been upgrading our diagnostics capabilities in preparation for Novette, and assessing and beginning work toward future diagnostics needs. Our work was comprised of the following major efforts

- Evaluating low-energy, x-ray spectroscopy by grouping instruments according to their prime applications, to obtain information on initial laser-plasma conditions.
- Focusing more resources on the advanced stages of the laser fusion process (implosion dynamics and final fuel conditions) where diagnostics are sparse.
- Developing new instruments to diagnose the x-ray laser experiments to be conducted on Novette.

Toward these ends, we improved existing high-priority x-ray spectrometers so that they can be used more accurately and routinely. In addition, we set out in new directions with regard to developing diagnostics that will shed light on more advanced problem areas.

We have developed a sequence of sophisticated x-ray spectrometers for measuring subkilovolt x rays. Our x-ray streak camera has been used with various x-ray energy-dispersing elements as both a detection and recording medium for each spectrometer. The advantages of high temporal resolution and multichannel recording provided by this instrument became obvious during the development program. This year, we made final improvements in the x-ray streak camera systems to achieve easier and more reli-

able use and better spatial resolution. With this capability, combined with the simplicity and absolute calibration of the improved Dante spectrometers, we will routinely produce accurate x ray spectra as a function of time.

LLNL has pioneered x-ray microscopy for ultra-high-resolution x-ray imaging of targets irradiated by lasers. Similarly, we are at the forefront of one-dimensional, time-resolved imaging with our streak-camera technology. Now we are focusing attention on the more difficult problems of two-dimensional dynamic imaging. We have generated several promising designs for a gated ultrafast-framing camera. In the next two years, we anticipate a physics and engineering effort that will culminate in high-spatial-resolution, subnanosecond, two-dimensional snapshots. This year, we were successful in developing thick zone plates for time-integrated imaging of high-energy x rays and charged particles. Our cameras will play an important role in diagnosing final fuel conditions for targets of intermediate density on Novette.

The initial Novette laser shots will be devoted to x-ray laser experiments. During 1982, we developed three new instruments for diagnosing the experiments. A high-resolution crystal streak camera was built to time-resolve the x-ray pump lines. Two other low-energy instruments were also developed to measure the actual laser output energy.

We have met the challenges of the past by providing ample diagnostics for the initial stage of the laser fusion process. Now, we are proceeding with confidence toward narrowing our options in the developed diagnostics areas to meet the challenges of new and difficult diagnostics regimes.

Authors: L. W. Coleman and V. W. Slivinsky

Development of Soft X-ray Streak Cameras

The first soft x-ray streak camera developed at LLNL was described in previous Annual Reports and in Ref. 1. This camera ("model 1") was used in Shiva and Argus experiments as the detector for spectrometers (diffraction gratings, crystal, and filter-mirror combinations) and axisymmetric reflection and pinhole microscopes.

A program to improve the camera's performance and reliability and to simplify its construction and operation resulted in two new soft x-ray streak cameras: model 2, for general use, and model 3, best adapted to crystal spectrometers (see Figs. 5-1 and 5-2). The new cameras use the same streak tube and similar transmission photocathodes as model 1.

The Model 2 Camera. Figure 5-3 shows a cross section of the model 2 camera and the two valves used to pump down the

camera and to vent it to atmospheric pressure with dry nitrogen. A mechanism positions the fragile transmission photocathode so that it is no longer a pressure barrier between the streak tube and the front of the camera. The photocathode is positioned in the streak tube only during data collection, when both streak tube and camera front are at high vacuum. At other times the photocathode is retracted to a position behind a window, where it can be inspected without affecting the camera vacuum. This window also permits observation of the photocathode when it is inserted in the streak tube, simplifying optical alignment of the camera.

The photocathode can be extracted through the window port after the camera is vented to atmospheric pressure. The camera is not moved during this process, so it need not be realigned. The front valve that connects the streak camera with the associated instrument has a glass window, and another transparent valve connects the instrument to the target chamber; these

Fig. 5-1. Cross section of model 2 soft x-ray streak camera.

valves permit optical alignment of the entire instrument independent of the vacuum condition of the target chamber, instrument, and camera.

The photocathode positioner, front valve, vacuum valve, and vent valve are operated with electrical motors, rather than with pneumatic actuators. This simplifies installation and removal and also simplifies the transportation and storage of the camera when it is under vacuum.

An electronic controller permits manual or remote camera operation.

The spatial resolution of the model 2 camera was optimized by modifying the geometry of the electron optics and the voltages of the focusing electrodes. Measurements of spatial resolution give 10% modulation for ten line-pairs per millimeter.²

The Model 3 Camera. The model 3 camera was designed to operate in conjunction with a crystal spectrometer. Its main characteristics are a cathode-to-front-flange distance of 75 mm, a 60° acceptance angle, and a 25-mm-long cathode slit; corresponding values for the model 2 are 300 mm, 12°, and 13 mm.

Alignment of X-ray-Imaging Pinholes.

Closely related to the work on soft x-ray streak cameras was the development of a technique to align pinholes for x-ray imaging. A typical pinhole microscope used to study laser-irradiated targets has a target-to-camera distance of several meters and an imaging pinhole a few tens of micrometers in diameter; the target size is a few hundred micrometers.

A satisfactory alignment procedure must place the pinhole on the straight line connecting the target and the photocathode slit to within a fraction of the pinhole diameter. The difficulty of this task is compounded by the need to perform frequent adjustments to overcome mechanical and vacuum drifts, by the restricted environment of the laser facility, and by the use of grazing incidence mirrors along the x-ray beam for spectral discrimination.

The alignment technique, which permits location of the pinhole to within a few micrometers of the lens optical axis, consists of mounting the imaging pinhole at the optical center of an optical lens whose focal length projects an image of the target on the photocathode slit. Figure 5-4 illustrates

Fig. 5-2. View of model 2 camera.

Diagnostics

the concept, which amounts to two coaxial cameras, one for x rays (using the pinhole) and the other for visible light (using the lens). Both images of the target are coin-

Fig. 5-3. View of model 3 camera.

cident in location and orientation, even if x-ray grazing-incidence mirrors are used. An orifice is drilled on the lens in correspondence to the pinhole to allow propagation of the soft x-ray beam defined by the pinhole.

With this implementation, a backlit target provides a dark profile on a bright field in correspondence with the photocathode. The pinhole-lens combination is moved to achieve the desired position of the target image on the photocathode slit.

Author: H. Medeck

Time-Resolved Spectroscopic Measurements Using X-ray Transmission Gratings

We have developed a time-resolved x-ray spectrometer in which a free-standing x-ray transmission grating is coupled to a soft x-ray streak camera. The instrument yields continuous x-ray spectra with 20-ps time resolution and moderate spectral resolution ($\Delta\lambda \geq 1 \text{ \AA}$) over a broad spectral range (0.1 to 5 keV) with high sensitivity and large information-recording capacity. The spectrometer's capabilities are well suited to the investigation of laser-generated plasmas, and they nicely complement other techniques for time-resolved spectroscopy. The

Fig. 5-4. Alignment of pinhole for x-ray imaging.

spectrometer has been used on a variety of laser-plasma experiments, many of which have been reported.³⁻⁵

In the spectrometer^b an entrance aperture (a narrow slit parallel to the grating lines) limits the divergence of incoming x rays. An x-ray transmission grating immediately behind the slit disperses the x-ray spectrum along the streak-camera entrance slit, which is aligned perpendicular to the grating lines. The dispersed spectrum is streaked to produce a time-resolved x-ray spectrum.

To assess the impact of the transmission-grating spectrometer on the diagnosis of laser-produced plasmas, it is instructive to compare its characteristics with those of other instruments currently being used for the time-resolved spectral analysis of laser plasmas. This comparison illustrates the complementary nature of the different measurement techniques and indicates the most appropriate operating regime for each. A rigorous, quantitative comparison of instrumental capabilities is not easy, because many parameters of interest vary with experimental configuration. Nevertheless, Table 5-1 lists reasonable representations of the capabilities of the transmission-grating streak camera, a crystal streak camera, and two fast multichannel analyzers (a multilayer-mirror streak camera and the Dante diagnostic).

A number of features in Table 5-1 are noteworthy. The crystal streak camera, because of its high linear dispersion, has a

spectral resolution about 2 orders of magnitude greater, and a spectral range 2 orders of magnitude less, than the other instruments listed. (It also operates at higher x-ray energies than the others, for the same reason.)

The Dante multichannel analyzer has the poorest time resolution, because of the limitations of x-ray diodes and of oscilloscope readout.

Because they provide continuous spectral readouts, the transmission-grating and crystal spectrometers have a much higher information-recording capacity than either of the multichannel analyzers. [The information-recording capacity, which we define as the number of distinct resolution elements in a data record, is the product of two quotients: the duration of the x-ray signal (which we arbitrarily set at 2.5 ns) divided by the time resolution, and the full spectral range divided by the spectral resolution.]

The transmission-grating and crystal spectrometers, the only instruments that combine high time resolution with continuous spectral coverage, complement each other nicely. The transmission-grating spectrometer operates with moderate resolution over a broad spectral range at low x-ray energies. The crystal spectrometer operates with high spectral resolution over a narrow spectral range at higher energies. The relatively high sensitivities of the transmission-grating spectrometer and the

Table 5-1. Characteristics of instruments for time-resolved spectroscopic investigation of laser-generated plasmas.

Instrument description	Spectral resolution	Time resolution (ps)	Spectral range (Å)	Approximate information capacity ^a	Relative sensitivity
Transmission-grating streak camera	1-2 Å	20	$\Delta\lambda_R \approx 120$ (2 - 124 Å)	10^4	High
Crystal streak camera	0.02 Å	70	$\Delta\lambda_R \approx 1-2$	3000	Low
Multilayer-mirror streak camera ^b	$\Delta\lambda/\lambda \approx 0.1$	20	$\Delta\lambda_R \approx 120$ (2 - 124 Å) (5 channels)	600	High
Mirror-filter pairs with x-ray diodes (Dante) ^b	$\Delta\lambda/\lambda \approx 0.2-1$	200-500 ^c	$\Delta\lambda_R \approx 50$ (8 - 60 Å) (10 channels)	50	Low

^aThe information capacity of the instrument is defined as the total number of distinct resolution elements in a data record. For this calculation a data record 2.5 ns long is assumed.

^bFor these instruments $\Delta\lambda$ typically varies with the channel reflector chosen. Rough values for spectral resolving power are provided.

^cIn a typical Dante, only one channel has a fast scope recorder ($\Delta t \approx 200$ ps); other channels typically have ~ 500 ps time resolution.

multilayer analyzer arise from their moderate spectral resolution, their good detector efficiency (because they use streak cameras), and the high efficiency of the dispersive element (grating) or filter (multilayer mirror). The lower sensitivity of the crystal spectrometer results from its very high spectral resolution (i.e., the high dispersion of its crystal). The lower sensitivity of the Dante system, on the other hand, results from the lower response function (in C/keV) of its detectors.

We have used the time resolved transmission grating spectrometer in the

Fusion Experiments program to investigate rapidly varying laser-generated plasmas. The laser pulses were short (100 to 700 ps), of high intensity (10^{11} to $10^{16} \text{ W}/\text{cm}^2$), and at three wavelengths: 1.06, 0.53, and $0.35 \mu\text{m}$. We irradiated various targets in studies of x-ray conversion efficiency, line production, and plasma transport effects. To demonstrate the capabilities of the spectrometer, we discuss data from an experimental series on the Argus laser facility.

Figure 5-5 shows data from a shot in which a gold disk target was illuminated at $3 \times 10^{11} \text{ W}/\text{cm}^2$. The laser parameters for

Fig. 5-5 Data from a gold-disk target in radiation, assembled to demonstrate the capabilities of the time-resolved transmission-grating spectrometer.

this shot were as follows: wavelength, $0.53\text{ }\mu\text{m}$; pulse duration (FWHM), $\sim 700\text{ ps}$; energy on target, 147 J ; nominal laser-spot diameter, $350\text{ }\mu\text{m}$. In Fig. 5-5(a), the photographic record of the time-resolved spectrum, we see the bright, zeroth-order (undiffracted) component, with its long tail of low-energy emission, and a discontinuity at the $44\text{-}\text{\AA}$ C_K absorption edge of the parylene cathode window.

Figure 5-5(b) shows the time histories of the emissions in zeroth order and in spectral bands centered at 150 eV ($\sim 83\text{ \AA}$) and 500 eV ($\sim 25\text{ \AA}$). Both the zeroth-order emission and the 150-eV continuum have FWHMs of about 830 ps ; long residual emission tails extend for almost 2 ns beyond the peaks. The similar shapes of these two spectra illustrate the extent to which the zeroth-order component is dominated by the low-energy continuum. The 500 eV emission more closely follows the shape of the incident laser pulse in time. It has a FWHM of about 700 ps , and no discernible residual emission. The emission at higher energy responds more rapidly to cooling processes in the plasma.

The principal spectral features of Fig. 5-5(a) are easily identified in Fig. 5-5(c), which shows the raw spectral data at $t = 0$ (pulse peak), $t = 300\text{ ps}$, and $t = 400\text{ ps}$. More subtle features, one centered at $83\text{ \AA}$ (150 eV) and another in a band at around $20\text{ \AA}$ ($500\text{ to }600\text{ eV}$), are also visible. The spectral shape of the emission does not change appreciably over the 700-ps span shown, although (as might be expected) the band emission at around $20\text{ \AA}$ is stronger at the peak of the emission, when the plasma is hottest.

Low-energy features of the emission are more easily seen in the unfolded x-ray spectrum in Fig. 5-5(d), which shows relative x-ray intensity at $t = 0$ and $t = 400\text{ ps}$. The dominant feature is the continuum peak at 150 eV . Other notable features are the gold band structures at around 250 and 500 eV . As noted above, the higher-energy band structure becomes less prominent as the plasma cools.

An x-ray transmission grating, coupled with a fast, spatially discriminating x-ray detector, can provide a powerful, versatile spectroscopic capability. With this detector we have obtained the first time-resolved, continuous soft x-ray spectra from laser-

irradiated targets and have made the first time resolved measurement of the low energy band structure of a multiply ionized high- Z plasma. Instrument calibration details such as grating efficiency, multiple diffraction orders, streak-camera window absorption, and photocathode response were appropriately handled. These measurements presage a new generation of experiments that will exploit this powerful spectroscopic capability for investigating the properties of rapidly varying laser generated plasmas.

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Fabrication of Very-Thick-Zone-Plate Coated Apertures

Free standing gold Fresnel zone plates for coded-aperture imaging of x-rays and particles have been produced in the past with thickness up to $38\text{ }\mu\text{m}$ (Ref. 8) using conventional UV photoresist technology. Gold zone plates up to $75\text{ }\mu\text{m}$ thick will be needed to image the high energy x-rays and charged particles from future laser fusion targets. This article describes the fabrication of $75\text{ }\mu\text{m}$ thick zone plates; the next article describes the results of imaging tests on these zone plates.

The fabrication of thick gold zone plates involves a number of steps, which are shown in Fig. 5-6. Briefly, the process involves creating a zone plate mold in a layer of polymer on a flex glass substrate, microplicating gold into the mold, attaching a support ring, and chemically dissolving the mold and the substrate.

As shown in step 1 in Fig. 5-6, the glass substrate is coated with a chrome adhesion layer, a gold plating base, and a chrome "etch-stop" layer (required for compatibility with subsequent steps). A polymer layer of Spurr low-viscosity embedding material (SIVEM),¹⁰ somewhat thicker than the desired zone plate, is added. A metal layer, typically aluminum, is thermally evaporated onto the polymer and coated with a UV-sensitive photoresist.

Fig. 5-6. Procedures used in making thick, free-standing gold zone plates using RH and gold-microplating technologies.

Fig. 5-7. A 75- μm -thick, 12-mm-diam gold zone plate made as described in Fig. 5-6. It has 100 zones and a 30- μm minimum zone width.

Fig. 5-8. Progressive taper of polymer side-walls during RIE can lead to collapse of thin zone walls.

In step 2, the zone-plate pattern is exposed into the resist. In step 3, the resist is developed and the uncovered portions of the aluminum layer are removed, leaving an aluminum/resist zone-plate pattern on top of the polymer. In step 4, the pattern is etched through the exposed polymer layer to the chrome etch-stop layer by reactive-ion etching (RIE) in an oxygen-discharge plasma.

In step 5, the aluminum mask is removed, as is the etch-stop layer (to uncover the gold plating base). The polymer pattern is then plated to the desired thickness in a gold electroplating bath.

In step 6, the polymer pattern and the plating base are removed. A ring holder is attached to facilitate handling of the zone plate.

Finally, in step 7, the glass support substrate is dissolved in hydrofluoric acid.

Figure 5-7 shows a 75- μm -thick, 12-mm-diam gold zone plate made in this way. It has 100 zones and a minimum zone width of 30 μm .

The production of these thick apertures required deeper reactive-ion etching and thicker electroplating than in previous zone-plate work, and involved the use of a novel material (SLVEM) for this application. As a result, we encountered a variety of new phenomena that made the work quite challenging. The phenomena fall into five categories:

1. Sidewalls of thick, high-aspect-ratio polymer structures can become tapered during RIE, leading to wall collapse (see Fig. 5-8).
2. RIE etching of the polymer is faster in wider zones; the etch-stop layer prevents sputtering of the gold plating base from wide zones.
3. Spikes of polymer ("grass") can form on the polymer surface during RIE; embedded in the gold zone plate, "grass" reduces the plated density and thus

increases the signal-to-noise ratio in coded images.

4. Zones in the polymer mold may be undercut and lifted off during removal of the chromium etch-stop layer.
5. Stresses in thick plated gold are high enough to pull up portions of the glass substrate.

Techniques developed to meet these difficulties are discussed in detail in Refs. 11 and 12.

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Tests of High-Energy X-ray and Proton Imaging

As target performance continues to improve, charged-particle imaging schemes that rely on relatively low-energy particles (such as 3.5-MeV DT alphas or 3-MeV DD protons) lose their usefulness because such particles do not escape higher-density targets. Hence it becomes necessary to look toward particles and photons of higher energy on which to do compression diagnostics. Zone-plate-coded imaging (ZPCI)¹¹ of particles and photons at these higher energies can increase the collection solid angle by as much as 10^4 over pinhole cameras of equivalent resolution. We have extended ZPCI to energies of 100 keV for x rays and to 6 MeV for protons.

We conducted ZPCI experiments for simulated point sources of 100-keV x rays and

6- and 10-MeV protons. (We used nonpoint x-ray and proton sources to make contact radiographs of the zone plate; the resulting images are the same as would be produced by a point source.) For the experiments reported here we made zone plates whose thickness and zone width meet the energy and resolution requirements of future laser fusion experiments (see the preceding article). The resolution of the ZPCI technique is approximately equal to the minimum zone width (i.e., that of the outermost zone). We also developed techniques for recording multi-MeV protons, using solid state particle-track detectors.¹¹

Figure 5-9 shows the zone-plate shadowgraphs. Optical reconstruction of these coded images gives the point response of the zone-plate camera to the various sources. Figure 5-10 shows pseudo-three-dimensional plots and lineouts through image maxima for the first-order reconstructed data. The FWHM values of the point-response distributions are as follows

- 100-keV x rays, 37 μ m.
- 6-MeV protons, 37 μ m.
- 10-MeV protons, 47 μ m.

The high quality of the coded images for 100-keV x rays and 6-MeV protons made possible the reconstruction of the data in higher order with improved point response. Figure 5-11 shows the results; the FWHM values are

- 100-keV x rays (7th order), 4 μ m.
- 6-MeV protons (3rd order), 9 μ m.
- 10-MeV protons (3rd order), 28 μ m.

The poorer resolution at 10 MeV than at 6 MeV, in both the first- and third-order reconstructed images, is due to the use of an

Fig. 5-9. Zone-plate shadow patterns recorded using (a) 100-keV x rays, (b) 6-MeV protons, and (c) 10-MeV protons.

Fig. 5-10 First-order reconstructed point-response images.

absorber foil to match the absorber-plus-zone-plate thickness to the proton range at 10 MeV. Protons scattered in the foil blurred the shadows of the narrower (outer) zones, thereby reducing the shadowgraph resolution and degrading the point response of the system.

These experiments give us initial estimates of the target yield required for proton imaging of the thermonuclear burn region in laser fusion targets. The proton data shown in Figs. 5-9 through 5-11 simulate exposures from an isotropic point source of 5×10^7 protons at a source-to-detector distance of 3 cm. Data taken at an equivalent target yield of 5×10^7 protons produced images of insufficient signal-to-noise ratio for high-resolution imaging. Analysis of image data at an equivalent yield of about 5×10^8 protons is continuing.

These results indicate that zone-plate-coded imaging is possible, with the resolutions indicated above, at x-ray energies up to 100 keV and proton energies up to 6 MeV. Thicker apertures (250 to 300 μm) will be needed to image protons of higher energy (e.g., 14.7-MeV DHe⁺ protons).

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Soft X-ray Spectrograph/ Radiograph with Large Angular Acceptance and Dynamic Range

On the basis of schemes proposed by Hagelstein,¹⁷ we are searching for x-ray lasers at wavelengths of 53.6, 75.0, 119.0 and 153.0 eV at the Novette facility. In connection with this search, we have designed and built a versatile diagnostic system (called the "Tail Safe" system) of large angular acceptance and large dynamic range. The system can be used as a medium- or high-resolution grating spectrograph or as a broadband x-ray radiograph.

Large angular acceptance is important in this application because of the uncertainty in the x-ray laser direction arising from possible errors in target alignment and from possible nonuniformity in the lasing medium. The ability of this system to cope with a relatively large directional uncertainty in the x-ray laser led to the name "Tail Safe."

Large dynamic range, important because of uncertainty in the total output energy of the x-ray laser, is obtained by employing a microchannel plate for x-ray detection. For uniform x-ray flux, the dynamic range obtainable with a microchannel plate operated

Fig. 5-11. Higher-order reconstructed point-response images.

in single-photon detection mode is given by the number of available channels multiplied by the dynamic range of individual channels, which is about 100. Thus for an x-ray laser uniformly illuminating an area containing 10^5 microchannels, and assuming negligible background, the microchannel plate can be operated linearly for $1/q$ to $10^7/q$ incident photons, where q is the quantum efficiency. The total output of the microchannel plate is of course the sum of the outputs of individual channels, regardless of whether the output is obtained by integrating electronically or by integrating the density on the film of individual channels. A further advantage of the microchannel plate over x-ray film is that it is very sensitive to x rays and that its gain is adjustable. In this report we present an overview of the diagnostic system and its most important features.

Figure 5-12 shows the layout of the diagnostic system. An adjustable slit determines the angular acceptance. Appropriate absorption filters are used to transmit energy regions of interest (and at the same time to reduce contributions from higher-order diffraction), and to attenuate the x-ray intensity to desirable levels. Depending on the choice of component A or B (see Fig. 5-12), any of three options can be realized.

Option 1 – Medium-Resolution Spectrograph. In this option, component A is a plane grating and component B is the x-ray detector, a microchannel plate coupled to a phosphor screen that is viewed by an optical camera. A plane grating having 1200 to 3600 lines/mm will give an energy resolution of 1 to 3 eV for 50- to 150-eV x rays. This simple option permits measurement of the two-dimensional beam profile and the total laser output energy. With a 100-eV blackbody background, the instrument can detect a laser gain-length product, $g \cdot l$, of about 7. Concave gratings reduce the geometrical divergence of the beam and improve the energy resolution by a factor of 4 to 5. (Note that this configuration is *not* a conventional grazing-incidence spectrograph with components on the Rowland circle.)

Option 2 – High-Resolution Spectrograph. In this option, component A is a spherical mirror that astigmatically focuses the x rays onto the entrance slit of component B, a Spectral Precision model 1000S 1-m grazing-incidence spectrograph. This spectrograph is upgraded with a gated, curved microchannel plate that is coupled to a phosphor screen and imaged onto an optical film by fiber optics. With a 3600-line/mm grating, this system can give a resolution of 0.03 eV, thus lowering the

Fig. 5-12. The "Fail Safe" diagnostic system.

detection threshold to $\alpha/\sim 4$ to 5 for a 100-eV blackbody background.

Option 3—Broadband X-ray Radiograph. In this option, component A is a plane x-ray mirror and component B is identical to that of option 1. With proper combination of the filter absorption and the critical reflection of the mirror, it is possible to obtain a 30-eV bandpass. This option is useful only when the output of the x-ray laser is high. The detection threshold is $\alpha/\sim 10$ for a 100-eV blackbody background.

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Summary of Novette Diagnostics

We have assembled a large complement of diagnostics for Novette, using equipment from Shiva and Argus along with several

new instruments. Details of the Argus and Shiva diagnostics were described in the 1979, 1980, and 1981 Annual Reports; the new diagnostics are discussed in the previous articles, above. Here we briefly list the diagnostics planned for the first six to nine months of Novette operations. (Figure 5-13 shows a Novette target chamber map, with the port locations of each diagnostic.) These instruments will allow us to test LASNEX predictions and theoretical predictions of laser-plasma coupling, x-ray production, and other key issues for laser fusion.

X-ray Diagnostics. X-ray diagnostics serve two purposes. First, we use them to make measurements related to radiation-driven targets, which we prefer for inertial confinement fusion. Second, we use them to make detailed measurements of x-ray spectra in order to determine the electron density and temperature in the plasma produced at the target.

Our x-ray diagnostics include the following instruments:

1. Dante. Two absolutely calibrated 10-channel soft x-ray spectrometers¹⁰ measure the soft x-ray spectrum from 0.1 to 1.5 keV. Time resolution is ~ 400 ps (~ 150 ps if Thomson oscilloscopes are used) and energy resolution is $E/\Delta E \sim 3$.

8x	8x microscope	SC	SCS
⊗	Blocked	SA	SCS alignment
CR	CR-39 knockon detector		
CR	CR-39 track detector	DS	SDSS
Cu	Cu activation		SDSS alignment
CL	Cryo laser	LC	SLS light collector
CV	Cryo viewer	ST	Streaked transmission grating
		TI	Target illuminator
RS	DCRS	Tag pos	Target positioner
DA	Dante A	TV	Target viewer
DH	Dante H	Top port	Top port
E	Energy-balance measurement		
FF	FFLEX	UV	UV probe in
H1	Henway #1		UV probe out
H2	Henway #2		
N2	N ₂ fill	ZP	Zone plate alignment
NA	Near-axis imaging		Zone-plate camera
N	Neutron time of flight		
OX	OX-2		
Cass tel	Optical-imaging Cassegrainian telescope		
Pb	Pb counter		
PH	Pinhole camera		
RC	Radiation-chemistry diagnostic		

2. FFLEX. The absolutely calibrated, 18-channel filter-fluorescer spectrometer^{17,18} measures the x-ray spectrum from 3 keV to over 250 keV.
3. 22 $\times$ microscope. This diagnostic^{19,20} uses Wölter axisymmetric or multilayer-coated Kirkpatrick-Baez x-ray mirrors to image the x rays from a target with a magnification of 22 $\times$. It has one energy channel, whose location (within the range 0.2 to 10 keV) is selected by the mirror and a filter.
4. 8 $\times$ microscopes. These two microscopes²¹ (magnification 8 $\times$) use

Kirkpatrick-Baez mirrors to produce four x-ray images of the target with energy channels from 1 to 5 keV.

5. SDSS. The spatially discriminating "streaked spectrometer (SDSS)"²² has two configurations. It can use several filters to measure time-resolved soft x-ray spectra from 0.1 to 1.2 keV, or it can use a pinhole and one filter (or two pinholes and two filters) to obtain a time-resolved one-dimensional image of the target in a chosen energy range (see "Development of Soft X-ray Streak Cameras," earlier in this section).

Fig. 5-13. Novette target chamber map with port locations of all diagnostics.

6. Transmission-grating spectrometer. This diagnostic²³ uses state-of-the-art transmission gratings coupled to an LLNL soft x-ray streak camera to obtain a continuous time-resolved soft x-ray spectrum. It should produce a more complete spectral shape and better time resolution than Dante, but it is not yet absolutely calibrated (see "Time-Resolved Spectroscopic Measurements Using X-ray Transmission Gratings," earlier in this section).
7. Crystal streak spectrometer. An x-ray streak camera is attached to a Bragg crystal spectrograph to give high-resolution, time-resolved x-ray spectra for energies above 800 eV. This diagnostic is particularly useful for measuring electron temperature, mass ablation rates, and radiative preheat, all of which may be time-dependent.
8. "Tail Safe" diagnostic. This diagnostic (see the previous article) uses a large-acceptance-angle reflection grating and a large (7.5-cm-diam) microchannel plate to measure the time-integrated x-ray spectrum in the vicinity of proposed laser transitions (50 to 150 eV) during x-ray laser experiments. For future experiments it can be converted into a high-resolution, time-integrating spectrograph or a high-sensitivity pinhole camera.
9. Zone-plate cameras. These cameras²⁴ produce time-integrated, two-dimensional images of target x-ray emissions in several energy bands between 3 and 100 keV. The images have a spatial resolution as good as $1\text{ }\mu\text{m}$, and a field of view of up to a few millimeters (these performance limits cannot be obtained simultaneously).
10. Crystal spectrographs. Three absolutely calibrated crystal spectrographs^{25,26} give time-integrated x-ray spectra. They can measure x-ray line emissions for plasma diagnostics and can measure x rays related to radiation drive and preheat.
11. Pinhole cameras. We use x-ray pinhole cameras to obtain time-integrated, two-dimensional images of x rays emitted from the target. K-edge filters are used to select the x-ray energies.

In addition to the diagnostics described above, we are installing an x-ray backlighting source, to be driven by a laser pulse

from an independent oscillator copropagated down one of the Novette chains. We expect to illuminate a backlighter disk with $\sim 300\text{ GW}$ of $0.53\text{-}\mu\text{m}$ (2ω) light. The backlighting x rays will be used to obtain one-dimensional and two-dimensional images of imploding capsules and planar foils.

Energy-Balance Diagnostics. To compress targets efficiently or to produce enough x rays for radiation drive, laser energy must be efficiently absorbed by the target. We measure the incident, reflected, and scattered energies to determine the absorption and to understand the physics of laser-target coupling.

Our energy-balance diagnostics include:

1. Incident and reflected energy. These diagnostics¹⁶ are part of the Novette laser system. They use diodes and imaging optics to provide absolute measurements of the 2ω energy incident on and reflected from the target.
2. Incident laser-light time history. A streak camera,²⁷ located in the output sensor package, is used for this measurement.
3. Reflected- 2ω time history. A streak camera,²⁷ located in the diagnostic sensor, is used for this measurement.
4. Imaged incident and reflected laser light. A time-integrated, two-dimensional image (somewhat degraded by the KDP crystal array) is recorded on film in the beam diagnostic system.²⁷
5. Scattered- 2ω photodiodes. An array of 46 2ω photodiodes (using bandpass filters and pin diodes) measures the energy fluence and angular distribution of the 2ω light scattered from the target. (Reference 28 shows a sketch of similar photodiodes used on Shiva. The other photodiodes mentioned below differ from these only in choice of filters.)
6. Plasma calorimeters. To complement the photodiodes, 22 type LC-27 plasma calorimeters²⁹ measure the energy of the plasma and x rays produced by the target.
7. Raman-light photodiodes. Raman scattering from some targets is significant. An array of 31 photodiodes measures the energy fluence and angular distribution of these emissions.
8. $3\omega/2$ photodiodes. Seventeen photodiodes measure the $3\omega/2$ emissions, which indicate the presence of plasmons near quarter-critical density ($1/4n_c$).

Laser-Plasma Interaction Diagnostics.

Additional optical diagnostics for measuring the plasma profile and coupling physics include the following:

1. Holographic UV probe. We are producing a high-quality 296-nm probe beam to measure the plasma-density profile by holographic interferometry.^{29,30} The beam may also be used for Faraday rotation and other optical measurements.
2. Optical/x-ray streak camera. This diagnostic (the OX-2)³¹ measures the relative timing and magnitude of incident laser light, 25- to 60-keV x rays, scattered 2ω light, Raman-scattered light, and $3\omega/2$ light. It will help determine which instabilities produce hot electrons.
3. Scattered-light streak and integrating spectrometers. To detect the optical signatures of plasma instabilities, we use standard optical spectrometers, with or without optical streak cameras, to measure visible-light spectra produced by the target.
4. Optical-imaging Cassegrainian telescope. This diagnostic measures light with wavelengths between 170 and 3000 nm emitted in the horizontal plane and at right angles to the laser beam. It uses $f/5$ optics and has a magnification of $8\times$. It can obtain time-integrated, two-dimensional images or time-resolved, one-dimensional images. It can measure emissions arising from plasma instabilities or from heating of a target by shock waves or preheat.
5. Raman discrete-channel spectrometers. These two diagnostics each use eight photodiodes (with narrow-bandpass filters) to measure the energy and spectrum of the Raman light emitted from the plasma, both in and out of the plane of polarization of the laser light.
6. Near-axis imaging system. This diagnostic, with a magnification of $10\times$, will use $f/10$ optics to obtain a time-integrated, two-dimensional image of light scattered from the target in the range 200 to 7000 nm. It will be located on or near the laser-beam axis. The optical system is achromatic and has a near-diffraction-limited design.

Neutron Diagnostics. We diagnose implosions using x-ray diagnostics and using the neutrons produced by thermonuclear fusion of the compressed DT. We determine

the yield and other important characteristics of the compressed target by measuring the number and the effects of the neutrons.

Our neutron diagnostics include:

1. Copper neutron-activation counter. This diagnostic³² determines total target neutron yield by measuring the radioactive decay of neutron-induced activations in copper. Its threshold yield is 5×10^8 neutrons per shot.
2. Three-tube scintillation detectors. The two diagnostics³³ determine neutron yields between 10^5 and 10^{12} by measuring the scintillations produced by neutrons in plastic. The diagnostic uses three photomultiplier tubes per detector.
3. Lead neutron-activation counter. This diagnostic³⁴ determines neutron yields above 7×10^4 by the same principle as the copper neutron-activation counter.
4. Silicon activation system. This diagnostic³⁵ determines the areal density, of the glass pusher that surrounds the fuel in an implosion target, by measuring the activation of the silicon in the glass.
5. Neutron-track detector. Using neutron-track detectors, we are developing diagnostic techniques that will enable us to measure the spectra of neutrons emitted by compressed targets. For sufficiently dense targets and sufficiently high yield, this spectrum contains information on the fuel ρr . (The areal density, or ρr , of high-performance Novette targets will be roughly 1% of that required to stop a 14-MeV neutron.)

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Data Management and Analysis

Introduction

Our experiments on target irradiation are designed to develop an understanding of ICF target physics. The design of future experiments is strongly based on the outcome of previous efforts. Thus, data from each target experiment must be analyzed promptly, accurately, and completely so that correlation with other experiments, simulations, and theories can proceed in a

timely fashion. The procedures used to treat raw data, calibrations, experimental parameters, and results determine the quality of the data available for shot-to-shot correlations and comparisons. The procedures also affect our ability to reprocess data in light of future discoveries.

Our goal in the management and analysis of data is to streamline the processes through which we develop an understanding of ICF target behavior. We accomplish the goal through design, development, utilization, and maintenance of resources that enhance our access to target results, ensure information quality, and reduce tedious involvement in the mechanics of data management and analysis. Recently, we have been automating many of our repetitive and mechanical tasks to make interactive functions friendlier to the experimenter.

In 1982, we made substantial progress in the development and implementation of hardware and software for data management and analysis. We completed the fusion experiments analysis facility (FEAF), a 2500-ft² addition to Bldg. 381. The VAX 11/780 computer system was moved from its temporary site to the permanent and centralized FEAF addition. The system was certified for classified operations, remote graphics terminals were added in staff offices to improve accessibility, and an image digitizer was received and installed.

We developed and implemented a one-way data-communications link. This link transfers raw data from the Novette control computer to the FEAF system, where complete analysis is performed and results are archived. For the first time, we have a hardware and operating environment that permits collection and analysis of data from all of our target experiments without the time-consuming transfer of data between laboratory computer systems via magnetic tape. We can, within minutes after a shot, collect experimental data, transfer it to FEAF for analysis, and then make results available to experimenters.

Our ability to manage data depends upon the integration of software and procedures with the hardware and physical environment. We have acquired commercial software packages that were selected to aid us with system development, applications,

analysis, and user interaction. This year, we installed major new releases of our operating system, a data-base management system, graphics packages, and programming languages. Our new software includes a command line editor, two additional programming languages, and two mathematics and statistics subroutine libraries.

We developed procedures and software for storage, resource management, and security. The procedures permit us to operate the system in an efficient, reliable, consistent, and secure manner. Novette-to-FEAF data-transfer procedures and software, used with the one-way link, were operational for the first Novette target shots. We continue to add to our base of data-analysis software, and in 1982, we added several processing codes for x-ray diagnostics.

The future poses a continuing challenge. We have excellent resources; a modern facility with extensive hardware capability; and powerful system, application, and analysis software. Now we must continue to develop user-interface software that permits easy access to manipulation of target data.

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Fusion Experiments Analysis Facility

The fusion experiments analysis facility (FEAF)¹⁶ is an integrated computer system for collecting, reducing, analyzing, and archiving inertial-confinement fusion (ICF) data obtained from unclassified and classified experiments conducted at the Novette and Nova laser facilities.

FEAF's primary purpose is to increase the productivity of the fusion experiments program staff by

- Reducing time spent on routine tasks such data collection and shot-book preparation.
- Improving shot-data turnaround time, accuracy, and readability.
- Supporting automatic shot archiving, shot-book preparation, data correlation, and shot-to-shot comparisons.

The facility also provides general-purpose computing capability for tasks related to

Fig. 5-14. The fusion experiments analysis facility (FEAF).

Fig. 5-15. View inside main computer room of FEAF.

laser fusion experiments and diagnostic design.

FEAF was designed to be user-oriented, rather than development-oriented. FEAF is intended to provide continuous timeshare operation 24 hours a day, 7 days a week except for scheduled maintenance, installations, and back-ups. FEAF permits an experimenter to examine target results in his office minutes after a shot.

Hardware and the New Facility. Construction of the 2500-ft² FEAF addition to

our office building (Bldg. 381) was completed in September, 1982. The facility has an independent air conditioning system, a centralized computer power center with clean power, fire and thermal protection, raised computer floors, and wireways for cabling. FEAF supports terminals in users' offices, and line printers nearby.

As shown in Fig. 5-14, FEAF comprises a computer room, a workroom, and a document vault. The computer room is a vault-type room housing the main computer

Fig. 5-16. Work stations at FEAF.

hardware and all system storage devices. The workroom is equipped with user work stations and output and reception areas; CRT terminals, color monitors, a digitizing tablet, an image digitizer, and hard copy output are available. Noisy computer hardware and I/O equipment are conveniently located, but acoustically separated from the work stations.

The FEAF computer system consists of three major subsystems: the host computer, an image digitizer, and a one-way data link from the Novette control-system computer. Provisions have been made to include a "born classified" data link in the future.

The host computer (Fig. 5-15), a DEC VAX 11/780, supports timeshare, batch, and real-time operations. Timeshare users interact with the system from CRT terminals at 10 work stations (Fig. 5-16) in the workroom and in 16 staff offices. Office terminals are connected to the computer over a multiplexed fiber-optic link. A second multiplexer, to support 16 more office terminals, is being installed.

User terminals provide graphics for viewing and analyzing target shot data. The

standard monochrome user terminal is the DEC VT100, upgraded with 480- $\times$ 640-element Retro-Graphics 640 graphics hardware. Color graphics are displayed on three Ramtek GM859 color monitors (with 512- $\times$ 640-element capability) attached to two Ramtek 9400 display processors. Other output devices provide flexibility in output presentation. They include

- Three Versatec model 1200 high-speed line-printer/plotters.
- A letter-quality printer.
- A color hard copy unit equipped for 35-mm film and 4- $\times$ 5-in. and 8- $\times$ 10-in. Polaroids, color negatives, and transparencies.
- A Versatec V80 printer/plotter with transparency capability.

To permit easy transfer of information between FEAF and other computer systems, we maintain a variety of I/O devices, including the following

- An 8-in. double-density floppy disk drive.
- A 10-Mbyte RL02 disk drive.
- Two 9-track magnetic tape drives.

Devices for direct data transfer between the FEAF VAX and other machines include

the one-way link (discussed below) and a connection to a WANG Office Information System (being installed).

Images from diagnostic instruments that record information on photographic film must be digitized and analyzed quickly. This year we installed an Eikonix 785 Image Digitizer, which can digitize an image of up to 2048×2048 elements in 1 min. The digitizer is controlled by a DEC PDP 11/34 computer, which has the following storage devices

- Two 10-Mbyte RL02 disk drives.
- An 8-in. floppy disk drive.
- A 9-track magnetic tape drive.

The system is undergoing testing and program development. We will use DECNET to transfer digitized data between the PDP 11/34 and the VAX, where final image analysis is carried out.

One-Way Link. A communications link speeds the transfer of shot data from the Novette VAX to the FEAF VAX for analysis. A one-way data channel is needed because the FEAF VAX produces classified information, while the Novette VAX can handle only unclassified data. The link ensures unidirectional data flow no matter what applications software resides at either end of the link.

In our design, an LSI-11 in the FEAF computer vault, running a ROM-based program, acts as a third-party message arbitrator between the two VAX computers. The hardware link between the LSI-11 and the VAX computers is achieved with fiber-optic cables and a communications protocol called NOVALINK. The software used to implement the one-way link consists of hardware device drivers to interface hardware for reading and writing data from the NOVALINK and a "one-way" policy algorithm that controls data flow through the LSI-11.

The one-way link allows transfer of unstructured data objects from the Novette VAX to the FEAF VAX. Software on the two VAX computers permits the transfer of structured objects such as VAX files. A transmitter program on the Novette VAX disassembles a file-structured object into a bit stream, which it sends across the link. A receiver program in the FEAF VAX rebuilds the bit stream into a file-structured object

of the same form as the original file. We can transfer files of all types and formats including sequential, direct-access, and indexed. Although developed specifically for the one-way link, our file-transport software can be generalized to transfer files between any two VAX computers connected with NOVALINK hardware. The one-way link was fully operational for the first Novette target shot.

Security. During 1982 we developed and implemented a security plan for the FEAF computer system. We established special hardware and procedures for our operations to permit remote terminals to have 24 hour access to classified data.

Unattended terminals are disconnected from the computer after normal work hours. Specially designed hardware-disconnect units permit the system operator to connect or disconnect all terminals by entering a preset number into a local keypad. Momentary-contact switches permit users who access the computer vault after work hours to connect and disconnect individual terminals.

The FEAF operating system normally spools computer files to system printers when instructed to do so by users. We have developed software, hardware, and procedures to prevent spooling of classified files to unattended printers. When such output is requested, a random number is returned to the terminal screen. To initiate output, the user goes to an output area and keys this number into a keypad attached to a printer.

Before obtaining a system account, a user must attend a one-hour security briefing, which gives an overview of FEAF, discusses the nature of the classified data on the machine, and details procedures to be followed when working with the system.

System Software and Procedures. We have acquired software and have developed operating procedures that maintain effective utilization of our hardware system while providing improved capabilities for our users. The FEAF system includes commercial software packages and hardware devices that provide excellent data-management and data-analysis capability. To maintain a robust, reliable, user-friendly software environment, we must have well-defined system management and operating

procedures and good software tools for data management and analysis.

In 1982 we developed operating procedures that control system resources efficiently and with little impact on users. Improved periodic disk-file backups improve system reliability and reduce down time. Backups are done often and generally do not interrupt system operation. System accounting procedures now include well defined user accounts and detailed audit trails of computer utilization. With these tools we can better understand user requirements and can make adjustments to accommodate changing demands on the system. New policies define computer access, user privileges, storage allocation, and process priorities.

Our policies and procedures are documented in a manual, *FEAF Operating Procedures*,¹⁷ which describes the responsibilities of FEAF operating personnel and the policies and procedures relating to system maintenance, supplies, start-up, shutdown, backups, accounting, security, hardware and software installation and documentation, and general operation. The manual is useful to persons involved with the day-to-day operation of the FEAF computer system.

We have also increased the availability and extent of system documentation. A library of software and hardware manuals is available in the FEAF work area, and extra copies of documentation can be checked out. Further documentation exists in on-line help packages, which include FEAF-specific information. An on-line inventory of FEAF-related hardware and software includes locations, dates, and pointers to more detailed documentation. A *FEAF User's Guide*¹⁸ gives an overview of the FEAF computer system and references to more detailed documentation.

Software maintenance contracts provide revisions, upgrades, and telephonic consultation service. Upgrades are installed promptly. The benefits of maintaining currency in our software system more than outweigh the expense of software maintenance contracts and the effort required to perform the upgrades. Below we give brief functional descriptions of commercial software available at FEAF.

VAX/VMS Operating System. Interaction between computer and user is controlled by an operating system. Utility programs provide users with general service functions.

The FEAF operating system, called VAX/VMS (Virtual Address Extension/Virtual Memory System), supports multi-user interactive and batch operations. The VAX/VMS operating system includes utility programs to aid the user with program development and file management. We have used the powerful VAX/VMS command language in the development of nearly all of our operating procedures. We upgraded VAX/VMS during the year to version 3.0.

Command Line Editor. We purchased a command line editor for interactively editing the Digital Command Language (DCL). The line editor allows correction of typographical errors without retyping an entire command string. It also stores up to 20 of the most recent commands in a buffer for easy recall and submittal. Editing characteristics are based on our standard text editor, EDT, making use of the line editor easy.

Programming Languages. We upgraded VAX-11 Fortran during the year to version 3.0. We added VAX-11 Pascal V1.2, VAX-11 C V1.1, and VAX-11 Basic V1.4 to the system.

Graphics Software Packages. Two ISSCO software packages for graphics-related data analysis are available on the FEAF. The first, called DISSPLA, comprises a variety of graphics subroutines callable from Fortran, Pascal, or C. DISSPLA presents a high-level language interface to the FEAF graphics hardware. We upgraded DISSPLA to version 9.0 during the year. The second package, called TELL-A-GRAF, is an easy-to-use conversational computer-graphics system that accepts English-like commands to create graphs using automatic scaling, layout, and annotation. TELL-A-GRAF interfaces with all the FEAF graphics hardware. We upgraded TELL-A-GRAF to version 4.0 during the year. An extension to TELL-A-GRAF, called CUECHART, allows generation of stencils and templates for standardization of graphical output, and speeds up the generation process.

Mathematics Packages. FEAF maintains several libraries of subroutines for scientific programming. The DEC RT-11 Scientific Subroutine Package provides users with more than 100 mathematical and statistical routines. During 1982 we added the International Mathematical and Statistical Library (IMSL), a large, widely used set of Fortran mathematical and statistical

subroutines for single- and double-precision arithmetic, and upgraded this package to version 9.0.

Data-Base Management System. Our data-management strategies are based on the powerful ORACLE data-base management system. ORACLE is based on the relational data model and has excellent user-interface features. We upgraded ORACLE to version 3.0 in 1982.

Summary. This past year we merged our computer hardware, system software, and operating environment into a computer system that has great potential. The system has already significantly enhanced our staff's capabilities. Raw data from each target-detonation experiment is now transferred directly to FEAF for analysis. Experimenters can display and analyze the results on CRT terminals in their offices. High-quality output of physics results can be generated quickly and easily.

With FEAF we can logically centralize data analysis, archive raw and processed data, and easily access, manipulate, and display shot results and summaries. To make the system more effective, we continue to design our applications software packages so that the boundaries between our database, graphics, scientific-calculation, and system-utility programs are transparent to users. Our goal is to combine hardware, system and applications software, operating environment, and user requirements into a system that relieves the staff of the need to know the details of the components of the system. To do this in an efficient and orderly way while meeting user requirements represents a significant challenge for the future.

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Processing of Novette Shot-Data Files

Structure and Contents of Novette Raw Shot-Data Files. We have completed and implemented the format for raw shot-data files. The raw data is put in two files, one

holding raw data from CAMAC-based instrumentation on the Novette target data acquisition system, the other holding raw data from charge-coupled devices.

The design for the CAMAC data file was based on three requirements:

1. The file contents must be written in formatted ASCII. This allows direct reading of the file and permits the use of text editors to correct incorrect character patterns caused by a processing malfunction. (With text editors we can recover data that would be lost if the file format was unformatted binary.)
2. The file contents must contain parameters that specify the setup of the electronic module in the data acquisition system for a particular experiment. These parameters can be used to reinitialize the data-acquisition system to the conditions for that experiment at any later time.
3. Each item of data must be preceded by a label. This makes it possible to make random searches of data and to make additions to the file without changing routines that read the data.

Table 5-2 shows how the setup parameters and data for a typical module appear in the shot file. The Novette target-data-acquisition system³⁰ collects and organizes data according to electronic module (module-label and id), front-end processor (node), and CAMAC crate and station. The setup consists of the position data (node, crate, and station) and the instrument

```
// Description of module
$$module-label
xxxx
**id
xxxx
**node
xxxx
**crate
xxxx
**station
xxxx
**operation-status
xxxx
**error-status
xxxx
**setting-wt
xxxx
**settings-rd
xxxx
**data-rd
.
.
**end
```

Table 5-2. Generic format for a set of electronic-module data in the Novette shot-data file for CAMAC-based instrumentation.

settings. Organizing the data according to a diagnostic system would be slower, because diagnostic detectors are connected to various modules in various crates. We relate data for this file to the corresponding diagnostics by a mapping (described in the next article), using the module id as a primary index.

As Table 5-2 shows, each data item or parameter is preceded by a label record. This allows the file to be modified without requiring modification of programs that read the file. That is, the software looks for a particular label when reading the file, and does not depend on the order of the data.

All data values in the file are written in a format appropriate for the data type. In several instances the 16-bit pattern, and not the number, is significant for data processing. For example, for the *caltd* module (which represents a 16-channel calorimeter waveform digitizer), each 16-bit data word contains two important numbers. Bits 1 to 12 contain the digitized waveform value (0-4095) and bits 13 to 16 contain the digitizer channel number (0-15).

A typical shot file for a Novette experiment is 200 blocks long (512 bytes/block). A file this size can be transmitted over the one-way link from Novette to FEAF in about 20 s.

The second type of file contains readouts from arrays of charge-coupled devices

(CCDs). CCD arrays are used to record the outputs of optical and x-ray streak cameras; we plan to use them for digital imaging of laser-beam profiles as well. CCD arrays are of two sizes: 320×512 -element arrays, used for streak cameras, and 320×256 -element arrays, used for imaging cameras. The data are written in an unformatted file. Each CCD-element data value is contained in 1 byte. CCD data files are transported across the one-way link in 20 to 30 s. Various codes at FEAF display and manipulate the data. Figure 5-17 shows the result of a FEAF program that produced a color-enhanced image of the CCD array on the Novette input streak camera.

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ORACLE Data-Base Management System at FEAF. We have begun to realize the potential of the new ORACLE data-base management system¹⁰ for processing and analyzing experimental data.

ORACLE already plays a major role in converting raw data into the forms in which it is used to interpret experiments. We have begun work on an ORACLE-based "electronic shot book," which will be implemented in 1983. The electronic shot book will allow an experimenter to display and correlate data from selected experiments on a computer terminal, instead of requiring searches through hard-copy shot books.

Our first application of ORACLE has been to implement a data base containing a diagnostic-instrument-to-electronic-module mapping. This mapping specifies the connection between a cable on a diagnostic detector and the corresponding channel of an electronic module (analog-to-digital converter, amplifier, or power supply). The module is part of a CAMAC crate, which in turn is connected to a front-end processor.

The mapping data base is used extensively by data-processing software to extract data from shot files (described in the preceding article). Data in a shot file are labeled by CAMAC parameters and electronic-module ID. Such a labeling is efficient for a data-acquisition system, but it

Fig. 5-17. Color-enhanced image for the contents of a CCD data file for the Novette input streak camera.

is inconvenient for the experimenter, who needs to label data by diagnostic system and detector. This conflict gives rise to the need for a mapping.

Figure 5-18 shows the mapping. A single piece of data is identified as passing through the connection between a detector and a single input/output (I/O) channel of an electronic module, which is connected to a station in a CAMAC crate. When a module has only a single output, the mapping locates the output by the connection number in the station called a UNIT. When a module has multiple outputs, the I/O channel of the module is used to label the corresponding data.

The mapping is easily implemented into an ORACLE data base. As explained in Ref. 40, ORACLE is a relational data base, in which the relations between quantities can be visualized as tables. Each column of the table holds quantities of the same type; each row holds quantities related in some unique way.

We have created two tables for the mapping. In the MODULES table [Fig. 5-19(a)], an electronic module is identified by type (MODULE-TYPE) and by an identification number (ID). The rest of the entry gives the node, CAMAC crate, station, and unit to which the module is connected. In effect, the relationship expressed in the MODULES table answers the question, "to

what node, crate, station and unit is a specified electronic module connected?"

In the DIAGNOSTIC-LINK table [Fig. 5-19(b)] a detector I/O connection is identified by a diagnostic system label, or name (DIAG-SYM), a detector number, and a detector I/O number (DET-OUTPUT-NO) that identifies the connections to a detector ADC-CHANNEL, gives the number of the module I/O channel to which the detector output is connected. TOTAL-ATTEN and TOTAL-DELAY give the total attenuation and time delay along the signal path. In this table, the relationship expressed answers the question, "to which electronic module I/O channel is the I/O of a specified detector of a specified diagnostic connected, and what is the total attenuation and delay along the signal path?"

The entries MODULE-TYPE and ID are common to both tables. ORACLE allows the two tables to be "joined" with these common entries, allowing the user to obtain CAMAC location parameters (node, crate, station, and unit), which are entries in the MODULE table, by specifying diagnostic name (DIAG-SYM) and detector number which are entries in the DIAGNOSTIC LINK table.

Processing software, using the MODULES and DIAGNOSTIC LINK tables, performs the following functions to map shot data:

Fig. 5-18. The diagnostic-to-electronic module mapping for the Novette target-data-acquisition system.

(a) Table MODULES with sample entries

MODULE-TYPE	ID	NODE	CRATE	STATION	UNIT	ON-LINE	- Additional columns for supplementary information
r7912	5	201	3	15	0	Yes	
1s2249	3	203	5	12		Yes	

(b) Table DIAGNOSTIC-LINK with sample entries

DIAG-SYM	DET-NO	DET-OUTPUT-NO	MODULE-TYPE	ID	ADC-CHANNEL	TOTAL-ATTEN	TOTAL-DELAY
FFLEX1	12	2	1s2249	3	7	100	11.5
DANTE	5	1	r7912	5		10	2.5

Fig. 5-19. ORACLE data-base tables for diagnostic-to-electronic-module mapping.

- When the experimenter specifies a diagnostic system, the software, using ORACLE interface routines, retrieves from the mapping data base the identities of all modules that collect data for that diagnostic, and the corresponding CAMAC parameters.
- Using the CAMAC parameters and module names as keys, the software extracts the data from the shot file and stores it in memory in the same form as it is stored in the module (by ADC-CHANNEL number).
- Using the mapping relationships in the DIAGNOSTIC-LINK table, the software gets the module name, ID, and ADC-CHANNEL for the data corresponding to a specified diagnostic detector output, uses the module parameters as indexes to get the data, and stores the data in memory or in a file indexed by diagnostic detector.

Mapping operations with data in memory are extremely fast. The advantage of putting the mapping information in ORACLE is that the data may be entered in any order, but can be rapidly and easily extracted in the desired order using the query language in ORACLE. Otherwise, FEAF programmers would have to set up mapping data files and write extensive software for such tasks as sorting. The query language also makes it easy to display the mapping data in the many forms in which it is needed by experimenters, engineers, and technicians.

After raw shot data are processed to the level where they become a set of analog signals (e.g., voltage or charge) organized

by diagnostic detector output, they are processed further to become true "physics data." This processing requires a large set of detector calibration and response values.

Some of these values are measured; others must be calculated from a large set of detector-component parameters. Storage and manipulation of these parameters are greatly simplified when the parameters are implemented into an ORACLE data base. As an example of this implementation, we describe tables in the ORACLE data base that are used to process data from scattered-light photodiodes. A more extensive application, involving x-ray detectors, is described in "Calculation of Response Functions for Dante X-ray Detectors" later in this section.

Figure 5-20 shows the arrangement of a typical scattered-light photodiode, and defines some important parameters. Photodiodes around the target chamber measure the light energy scattered and emitted by the target within wavelength ranges determined by filters. The scattered-light energy distribution (energy/solid angle) depends on the diode's sensitivity, S_p , its position (r , θ , ϕ) with respect to the target, and the transmission, T , of the filter window.

Each photodiode assembly (diode and filter) is assigned an identification number. A given array of photodiodes is used for many experiments; it is therefore efficient to associate each photodiode in an array with the corresponding configuration ID number, and to associate a shot with a configuration ID rather than with a list of photodiodes.

Fig. 5-20. Parameters relevant to analysis of data from scattered-light photodiodes.

(a) Table PD-CONFIG with sample entries

PD-CONFIG-ID	PD-ID	THETA	PHI	R-DIST
PDAA01	100	120.	250.	51.0
PDAA0J	105	135.	90.	51.0

(b) Table PD-CALIB with sample entries

CALIBRATION DATA						
PD-ID	SENS	FILTER-TRANS	FILTER-WAVELENGTH	MO	DY	YR
100	1.2E + 04	.92	.53	06	23	82
105	1.6E + 04	.95	.53	04	18	82

Fig. 5-21. ORACLE data-base tables for photodiode configuration and calibration data.

A photodiode can be moved to many positions without changing its sensitivity, and different filters can be used with a given photodiode. It is therefore logical to store diode position data (uniquely related to array configuration) and sensitivity and transmission data in separate data-base tables.

The PD-CONFIG table [Fig. 5-21(a)] contains the position data, indexed by PD-CONFIG-ID (the photodiode configuration ID just described) and PD-ID, the photodiode identification number. THETA, PHI, and R-DIST denote the quantities θ , ϕ , and r defined in Fig. 5-20. Values in the PD-ID column correspond to values in the DET-NO column of the DIAGNOSTIC-

LINK table; this correspondence allows the processing software to correlate shot data with calibration parameters.

The PD-CALIB table [Fig. 5-21(b)] contains the diode sensitivity and filter characteristics, indexed by photodiode ID (PD-ID).

The ORACLE query language makes it possible to arrange photodiode data from tables PD-CONFIG and PD-CALIB in many ways, such as by PD-ID, θ , or ϕ .

The scattered-light energy distribution that is output by the photodiode processing software is stored in another ORACLE data-base table, which contains the shot number, photodiode ID, and energy value for each photodiode. These data are used

for further analysis, such as determining energy balance and correlating scattered-light distributions for several shots.

The ORACLE data-base management system is invaluable in the processing of experimental data. Calibration and diagnostic data can be accessed easily by different programmers, all using the same software. ORACLE speeds up data storage and access. This experience leads us to expect that an ORACLE-based "electronic shot book" will give experimenters a versatile tool for organizing and analyzing their data.

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Software on the FFAF VAX

Calculation of X-ray Reflectivity. Data on x-ray reflectivity for a wide variety of materials are necessary in both the design of many x-ray diagnostics and the subsequent analysis of data. This year, to support data processing of x-ray diagnostics on the FFAF VAX, we have set up a library of the atomic scattering factors, f_1 and f_2 , for 94 elements and for the x-ray energy range from 0.1 to 2.0 keV. Reflectivities are easily calculated from f_1 and f_2 , which are a function of the material and photon energy. They are related to the photoelectric cross section, $\mu(E)$, where E is the photon energy calculated by

$$f_1(E) = Z + \frac{1}{\pi r_0 h c} \int_0^E \frac{E' \mu(E') dE'}{E^2 - E'^2} \quad (1)$$

and

$$f_2(E) = \frac{E \mu(E)}{2 r_0 h c} \quad (2)$$

where r_0 is the classical electron radius, h is Planck's constant, c is the speed of light, and Z is the atomic number.

Whereas f_2 is easily calculated from the cross section, calculations of f_1 for hundreds of energy values involve many thousands of integration steps. Calculation of the integral every time a value of reflectivity is needed would result in long program-

execution times (many minutes) and a large burden for a moderate computer system like VAX. While the scheme for reflectivity calculations is presently implemented on the Octopus system, calculating the reflectivity for hundreds of energy values takes minutes even on that large system. To optimize the situation, we have set up a library on the VAX consisting of f_1 and f_2 values for 94 elements and a wide energy range. Thus, the integral for f_1 need only be calculated once. For energy values not in the library, f_1 and f_2 can be calculated by either linear or logarithmic interpolation. Once f_1 and f_2 are extracted from the library, reflectivity can be calculated in a few microseconds of computer time. The result is that each time a reflectivity calculation program is run, execution time is a few seconds rather than a few minutes.

The FFAF VAX library on the atomic scattering factor is contained in a keyed-access file. Data are grouped by f_1 and f_2 values for the energy range from 0.1 to 2.0 keV. The key for each data set is the atomic symbol of the element corresponding to the data. Only data for elements are contained in the library. The f_1 and f_2 data for compounds are calculated by

$$f_1 = \frac{1}{N} \sum N_i f_{1i} \quad (3)$$

and

$$f_2 = \frac{1}{N} \sum N_i f_{2i} \quad (4)$$

where N is the total atoms per molecule and N_i is the atoms per molecule of element i .

Because our present range of data for f_1 and f_2 , with respect to energy coverage, is insufficient for some applications, we will extend the library's upper limit of photon energy coverage. The present data were obtained from a compilation of atomic scattering factors,⁴¹ which contained energy coverage to only 2 keV. In expanding our coverage to 10 keV, we will use the same method of computation and the same photoelectric cross sections.

A Fortran subroutine has been written, which allows programmers to extract all I_1 and I_2 data on a specified element for use by diagnostic data-processing programs. Our new software has been incorporated into software that allows the user to calculate response functions of x-ray detector channels with mirrors. Applications of our software include channel design of both the x-ray streak camera and the mirror Dante low-energy spectrometer. With the response-function software, the user need only specify the mirror material, the grazing angle of incidence, and an energy range for which the response function is to be calculated. The software automatically extracts the I_1 and I_2 data, calculates the reflectivity, and displays the results. Figure 5-22 shows the reflectivity of a nickel mirror with a 4° grazing angle of incidence. Figure 5-23 shows the response of an x-ray streak-camera channel consisting of a carbon mirror with a 4° grazing angle and a $3\text{-}\mu\text{m}$ -thick carbon filter, both of which are folded with the streak-camera photocathode response.

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Calculation of Response Functions for Dante X-ray Detectors.

The Dante x-ray detector system is a spectrometer designed to measure sub-keV x-rays emitted by laser-irradiated targets. The system comprises several channels, each consisting of an x-ray diode (XRD) and appropriate filters. In response to a source spectrum, $I(E)$, the output charge, Q_n , from the n th channel is

$$Q_n = \int_0^E I(E) R_n(E) dE, \quad (5)$$

where $R_n(E)$ is the channel response function (the charge produced per unit energy interval centered at E for a unit input spectrum level).

The Dante data-analysis problem is to obtain $I(E)$ by unfolding Eq. (5) for each active channel. Here we describe an automated method of calculating the response functions $R_n(E)$; the method requires only that the user enter a channel name.

Fig. 5-22. Reflectivity of a nickel mirror to x-rays in the range of 0.1 to 2.0 keV.

Fig. 5-23. Response of an x-ray streak camera channel with a carbon mirror and carbon filter.

Individual channel components (XRDs and filters) are given identification numbers and corresponding calibration data are stored using the ORACLE data base management system (see "ORACLE Data Base Management System at LLNL" earlier in this section) installed on the ULA VAX. When a channel is designed, the appropriate XRD and filter identification numbers, and a name and other data identifying the channel, are entered in another table. An example is shown in Fig. 5-24.

For each channel, R_n is calculated at each of 250 equally spaced energies between 0.1 and 2 keV. Let E_i denote one of these energies. $T_i(E)$ be the net x-ray transmission of all of the filters in the channel, and $S_i(E)$ be the sensitivity of the channel XRD (for clarity, we have dropped the channel identifier n). Then the response function at E_i is

$$R_i(E) = S_i(E) T_i(E), \quad (6)$$

Fig. 5-24. Page from the Dante data-base forms.

```

***** DANTE CHANNEL CONFIGURATION FORM *****

CHN_ID
DT060601

THETA  PHI  AREA  DIST
90      90    1.13  221

CATHODE  EDGE  FILTER1  FILTER2  FILTER3  FILTER4  MIR
ALJ 25   23     PY9999951  FV1 51   FV1 52

```

Fig. 5-25. Calculated final response function for Dante XRD channel DT060601.

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The net transmission, $T(E_i)$, is given by

$$T(E_i) = \prod_j T_j(E_i) \quad (7)$$

where $T_j(E_i)$ is the transmission of the j th filter in the channel. The transmission of an individual filter depends on the (energy-independent) transmission, T_i^{mask} , of its

supporting mask (if any), on its total x-ray cross section, $(\mu/\rho)_i$, and on its areal density, $(\rho\Delta)_i$:

$$T_i(E_i) = T_i^{\text{mask}} \exp \left[-(\mu/\rho)_i (\rho\Delta)_i \right] \quad (8)$$

If the filter foil is a chemical compound, its total x-ray cross section is given by

$$(\mu/\rho)_i = \sum_k W_{ik} (\mu/\rho)_k \quad (9)$$

where W_{ik} is the weight fraction of the k th elemental constituent of the filter, and $(\mu/\rho)_k$ is the total x-ray cross section of that element at E_i .

Some of the quantities needed for the transmission calculations (the identities and areal densities of the filters, and the mask transmission factors) are stored in the data base. If the filter is a chemical compound, the software extracts the identities of the constituent elements and their corresponding weight fractions from a COMPOUND table. Cross sections are calculated directly from data in the FEAF version of the x-ray

cross-section library CSLOW, using log-log interpolation to find the cross section at E_c from values defined on the CSLOW energy grid.

The XRD sensitivity data comprise a set of calibration values (in units of pC/keV) at selected energies between 0.1 and 9 keV. These data are interpolated to give sensitivities, $S(E_c)$, at the desired energies. An example of a raw calculated response function is shown in Fig. 5-25.

One more step may be performed. An XRD channel may be used many times to take data between calibrations. The detector response may vary with time, as indicated by the calibrations. If both pre- and post-calibrations exist for an XRD, with respect to the time of the data collection, the pre and post-calibration response functions are both generated and a linear interpolation is performed to get as accurate a response function as possible. If only one calibration is available, the response function closest to the time of the experiment is used.

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Processing of 22× Microscope Data. We made extensive use of the 22× streaked x-ray microscope¹² in the first x-ray back-lighting experiments at Shiva. Data were obtained for backlighting accelerating-slab targets and for imploding fuel-capsule targets. The data consisted of spatially resolved streak records recorded on film. We wrote software to obtain position-vs-time data from the streak records. This article describes the digitized data and the functions of the processing software.

Developed film records from the 22× microscope are digitized with a photometric densitometer system (PDS). Besides the streak image, the film also contains an image of a step wedge. This portion of the film is also digitized to relate film density to relative exposure, which in turn can be related to x-ray source strength.

The digitized data from the PDS are transported on magnetic tape to the FEAR VAX, where the data are read, converted to

VAX-compatible format, and stored on the FEAR data disks.

The software for analyzing the 22× microscope data allows the user to carry out the following functions

- Process the film-calibration data to obtain a density-log exposure (D-log E_c) table.
- Obtain a color-enhanced image of the streak record.
- Take vertical and horizontal lineouts of the streak image and make lineouts from averages of selected regions of the image to reduce film noise.
- Combine lineout data with camera streak speed, magnification, and the D-log E_c table, to obtain spatial and temporal profiles of the source.

Program output consists of the following items

- A color-enhanced image of the digitized streak record on a Ramtek color monitor.
- A hard-copy contour plot of the streak image with time and distance scales based on calibration of the x-ray microscope and camera. Locations of all lineouts are indicated on the contour plot.
- A plot of the film-calibration step-wedge data (film density vs relative exposure).
- For each vertical lineout, a plot of relative x-ray intensity vs time.
- For each selected horizontal lineout, a plot of relative x-ray intensity vs position.

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Processing of X-ray Filter Transmission Data. We have developed a routine to expedite the analysis of filter-transmission calibration data using the FEAR VAX computer and its ORACLE data-base management system. The routine presently handles monoenergetic x-ray transmission data from the LLNL IONAC calibration facility, but data from other monoenergetic sources can be easily accommodated.

Tables have been developed using ORACLE to facilitate entering the transmission data and composition of each filter. A table has been established to store specifications of frequently used compounds. Thus, both pure elements and complex materials

can be easily used as components in a specific filter. Each data set for a given compound is keyed by a three-character identification. For example, PYN indexes the nominal density of parylene-N together with the atomic symbols and weight fractions of each of its constituent elements.

After receiving a filter identification number, the processing routine first interrogates the data-base tables for the symbols of all constituent elements, and then extracts the necessary element cross section values from the CSLOW library. (Based on data obtained from LLNL, I Division, CSLOW has been installed on the VAX as a keyed-access library.) Total cross sections for compounds are computed from weight fractions obtained from the standard compound table. The process code then uses the compound or element cross section to gen-

erate a least-squares theoretical fit to the appropriate transmission data, with the areal density of each constituent as a fitted parameter. The calculated areal densities are automatically stored in the corresponding filter-specification table.

Figure 5-26 shows theoretical fits for a few filters used in a Dante x-ray spectrometer. Areal density ($\mu\text{g}/\text{cm}^2$) and filter thickness (μm) are given for the specific filter component, as determined by the code. Results shown in Fig. 5-26(a), consisting of parylene-N mounted on nickel mesh, are typical of the blast shields employed for nearly every Dante channel. When the mesh optical-transmission value of 86% is used along with the standard parylene material specification, an excellent fit is obtained for all IONAC transmission points. The thickness value of $0.20 \mu\text{m}$ returned by

Fig. 5-26. Theoretical fits to x-ray filter-transmission data with the HAF VAX processing routine.

*If the statistical weight = 0, that data point is not used in the filter-thickness calculation.

the code is in good agreement with the value quoted by the vendor.

The results shown in Fig. 5-26(b) suggest that our adopted standard specification for Formvar is reasonable, since a good fit is obtained for all data points except those on either side of the oxygen K-edge. (The disagreement of two data points is most probably due to the slight spread of IONAC source L-lines extending across the edge.)

Figures 5-26(c) and 5-26(d) show the results for cobalt and copper, respectively. We have found that the best results are obtained by specifying the element and its oxide, while excluding the unreliable data points falling on the filter absorption edge. The best fit for oxide-thickness values determined by the code are quite reasonable. A modification is planned that will provide a weighted least-squares fit to the transmission data; this modification will more rigorously take into account experimental errors and permit the use of marginal data with significantly reduced weights.

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Experiments and Analysis

Introduction

During 1982, we completed the reduction and analysis of data from Argus and Shiva experiments. In addition, we have devoted efforts to planning and preparing for Novette experiments.

With respect to Argus, the wavelength scaling of laser-plasma coupling was a major area of experimental study. We conducted experiments at irradiating wavelengths of 1 μ , 2 μ , and 3 μ . Our absorption measurements have shown the dominant role of inverse bremsstrahlung as the laser wavelength is decreased and the target Z is increased. The suprathermal-electron pro-

duction, as inferred from high energy x-ray measurements, also decreases with decreasing wavelength for fixed irradiation conditions. Optical measurements show that a small fraction of the incident energy is coupled into plasma instabilities, consistent with the limited laser energy and plasma size in the Argus experiments. While the efficiency of conversion of incident laser light to thermal x-rays generally increases at shorter wavelengths, we have not yet explained the unexpected decrease in conversion at intensities on the target below 10^{12} W/cm². In addition, layered target electron transport experiments at 1 μ and 2 μ have shown that the results cannot be matched with modeling in terms of a single flux limiter.

One of our final experimental series with Shiva was to study x-ray line emission from foil targets as candidate flashlamps for an x-ray laser. We characterized the intensity and angular distribution of the candidate x-ray spectrum and verified that uniform irradiation of the flashlamp foils produces uniform x-ray emission. We also measured the x-ray line energies, the duration of the line emission, and demonstrated that the line brightness achieved in these experiments was nearly that required for the proposed photo-pumped x-ray laser schemes.

During the last experiments on Shiva, we collaborated with the Naval Research Laboratory to examine target hydrodynamics and the interaction of long laser pulses with long scale length plasmas. In addition to the customary optical and x-ray diagnostics, we made time-resolved measurements with x-ray backlighting and double foil targets to measure target velocities.

X-ray spectroscopy will play an important role in understanding the processes involved in the large-corona plasmas of future targets. We participated in preliminary experiments at the Rutherford Appleton Laboratory, which were designed to evaluate the diagnostic promise of x-ray spectroscopy. Streaked x-ray spectroscopy, together with multicomponent-disk targets, provided us with measurements of line widths and intensity ratios for analysis.

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Wavelength-Scaling on Argus

Shorter laser wavelengths should improve inertial-confinement fusion performance by reducing the number of suprathermal electrons and thus reducing target preheat. Other benefits should include slightly higher absorption and slightly higher x-ray conversion efficiency. We completed the analysis of the four Argus wavelength-scaling experiments during 1982. The experiments were conducted at the following laser wavelengths and maximum energies: $0.53\text{ }\mu\text{m}$ (2ω , 40 J); $0.35\text{ }\mu\text{m}$ (3ω , 40 J); $1.06\text{ }\mu\text{m}$ (1ω , 90 J); and 2ω , 200 J. The first two segments were reported in the 1980 and 1981 Annual Reports.^{41,45} Last year's Annual Report⁴⁵ also contained preliminary results from the 1ω and high-energy 2ω experiments. In this report we discuss the final results of these two experiments. For details see Refs. 43 through 46.

X-ray Conversion Efficiency. The x-ray conversion efficiency, which we define as the ratio of the energy emitted in x rays to the absorbed laser energy, is a sensitive test of our ability to model correctly the complicated radiation-physics and energy-transport processes taking place during the irradiation of high-Z disk targets. The majority of the x-ray energy is emitted in the spectral region $0.1 < h\nu < 1.5\text{ keV}$. We used two Dante systems, the eight-channel "T" and the five-channel "L," to measure the x-ray energy over this spectral range at four different angular locations. Rather than move the diagnostics, we rotated the targets 12° or 30° with respect to the laser beam.⁴⁷ We assumed that this would not significantly affect the physics; various measurements support this assumption. For example, the absorption was measured for both orientations⁴⁵ and found not to change for these small tilts. The angular distribution of x-ray energy was built up over a number of shots; the integral overall solid angle gives the total x-ray energy. These measurements were performed for nominal incident intensities of 3×10^{13} , 3×10^{14} , and 10^{15} W/cm^2 . The results are summarized in Fig. 5-27. This figure includes data from the first three experimental series, which used laser energies of up to 40 J at 2ω and 3ω and up to 90 J at 1ω . LASNEX predictions are also shown.

In the 40-J (low-energy) 2ω experiments,^{41,44} which necessarily used small spots (down to $80\text{ }\mu\text{m}$), the x-ray conversion efficiency from gold-disk targets stayed constant, or increased slightly, as the laser intensity was increased from 3×10^{13} to $3 \times 10^{14}\text{ W/cm}^2$. This result is contrary to LASNEX predictions. These conversion efficiencies were checked at larger spot sizes during the Argus high-energy 2ω experiments. This check consisted of two shots at each of the two intensities, with the target at 30° to the laser, using the Dante T only. No laser absorption measurements were made, but the fractional backscatter through the f/2 lens was similar to the low-energy results. The focal spot size was $350\text{ }\mu\text{m}$ for all shots. Laser energy was 22 J for the $3 \times 10^{13}\text{ W/cm}^2$ shots (similar measurements were taken during the low-energy experiments), and 170 J for the $3 \times 10^{14}\text{ W/cm}^2$ shots. Assuming that the x-ray angular distribution and the laser-light absorption were the same as in the low-energy experiments, we find the conversion efficiency at $3 \times 10^{13}\text{ W/cm}^2$ to be $48 \pm 3\%$ (the low-energy value was 45%—see Fig. 5-27); at $3 \times 10^{14}\text{ W/cm}^2$ the efficiency was again $48 \pm 2\%$ (the low-energy value was 53%). The uncertainties represent only the spread between the two shots; no other errors are included. Because of the assumptions just mentioned, especially those concerning the angular distributions, we estimate the relative accuracy (compared to previous experiments) to be no better than 20% of the stated number.

The large-spot data support the conclusions from the small-spot data: for 2ω light, at incident intensities of $3 \times 10^{13}\text{ W/cm}^2$, there is a substantial difference between the LASNEX-calculated and measured x-ray conversion efficiencies. The reasons for this discrepancy are being investigated. Areas under consideration include

- Emission from material at density near n_c , the critical electron density for laser light.
- Laser light absorption in the corona at densities below n_c .
- Non Maxwellian electron distributions.
- Improved modeling of the opacity.
- Improved modeling of electron transport.

Hard X-ray Measurements. Measurements of hard x rays ($h\nu > 10\text{ keV}$) reveal the typical energy and level (total energy) of suprathermal electrons, which are

Fig. 5-27. The data points show measured x-ray conversion efficiencies at various incident intensities for the three laser wavelengths. Solid lines show LASNEX calculations.

Fig. 5-28. Hard x-ray spectra from gold-disk targets for 1ω and 2ω ; average intensities are 10^{15} W/cm 2 , $I_{\lambda} = 90.1$. The 1ω data bars, showing the range of values obtained in five shots, indicate excellent reproducibility.

responsible for target preheat. Figure 5-28 shows the hard x-ray spectra recorded in high-intensity (10^{15} W/cm 2) irradiations of gold-disk targets at 1ω and 2ω . Except for differences in the beam profile, the irradiation conditions were the same for both wavelengths. Since the high-energy slopes of the x-ray distribution functions (which correspond to "hot temperatures" of about 18 keV) are nearly the same, the fluence of suprathermal electrons is immediately seen to be reduced by approximately a factor of 10 at the shorter wavelength, despite the fact that the 2ω beam contained more "hot spots" (regions of very high intensity) than did the 1ω beam. The decreased level of suprathermal electrons is attributed to the reduction of parametric instabilities with the shorter-wavelength light.

The hot-electron temperatures indicated by the measured hard x-ray spectra do not scale with the most straightforward expectations for resonance absorption.^{48,49} Theory and collisionless-plasma simulations have suggested that

$$T_{\text{hot}} \sim (I\lambda^2)^{0.4} \quad (10)$$

for the hot electrons produced by resonance absorption. Previous short-pulse experiments at $1.06 \mu\text{m}$ have supported this intensity scaling.⁵⁰ However, in the Argus

wavelength-scaling experiments, we do not observe this wavelength-dependence for T_{hot} . This indicates either that the resonance absorption scaling is complicated by the beam nonuniformity, or by processes such as filamentation of the laser light, or that other processes, such as parametric instabilities, are responsible for the small hot-electron populations observed here.

Stimulated Raman Scattering. Stimulated Raman scattering (SRS), a potential source of preheat, can occur at wavelengths between λ and 2λ , where λ is the laser wavelength. This region is difficult to cover

in 1ω experiments because of the relative insensitivity of the available diagnostics (see Ref. 51, however). During the high-energy 2ω experiments, we were able to obtain, for the first time, spectra of Raman-scattered light with good time resolution (~ 30 ps). We were able to detect SRS only at the highest irradiation intensities (average intensities of $\sim 5 \times 10^{15}$ W/cm²). Details of these measurements are reported in Ref. 52. Figures 5-29 and 5-30 show examples of the data obtained from these shots on disk targets. The time integrated spectra are similar to those obtained at other wave lengths,^{47,53,54} and show the presence of the

convective Raman instability. If we assume that the corona temperature is about 2 keV, the observed spectral range corresponds to densities from 0.04 to 0.16 n_e .

The time resolution in these spectra reveals that the emission occurs in rapid pulsations over a wide range of wavelengths. This is most obvious in Fig. 5-29, but it also can be seen by close examination of the original data for the other shots. The origin of these pulsations is unknown; they could be produced by plasma processes leading to unstable wave structures or by nonsteady hydrodynamic flows. Examples are filament formation and collapse, or chaotic Brillouin

Fig. 5-29. Backscattered SRS from gold-disk target irradiated with 132 J. (a) Streak-camera data; fiducial is at lower right. (b) Spectrally integrated time history. The time $t = 0$ is the peak of the incident pulse.

Fig. 5-30. Backscattered SRS from beryllium disk target irradiated with 128 J. (a) Streak-camera data. (b) Time-integrated Raman spectrum.

scattering (see "Nonsteady Stimulated Brillouin Scatter Induced by a Reflective Critical Surface" in Section 3), or nonsteady density-gradient steepening near the critical⁵⁶ or quarter-critical surfaces.⁵⁷ Some plasma profiles may be intrinsically nonsteady as well.⁵⁷ These or similar processes could modulate the Raman behavior by modifying the plasma conditions on a short time scale.

However, the idea of large-scale profile modification, at densities of $1/4n_c$ or greater, has several problems. First, any profile modification must occur at speeds well above the ion-acoustic speed, or else the unmodified profile must be steeper ($l \sim 10 \mu\text{m}$) than believed. We infer this by noting the apparently simultaneous (within about 30 ps) pulsation over a wide spectral (and, by inference, density) range. Second, time-resolved measurements of the $3\omega/2$ harmonic, believed to be emitted from near $1/4n_c$, do not show pulsations on these shots.

The other interesting feature of these data is the absolute timing of the signal. The zero time on the figures corresponds (to within about 100 ps) to the peak of the incident laser pulse. The Raman light does not begin until near the peak of the laser pulse. This is in reasonable agreement with hydrodynamic calculations, which show that it takes about this much time for the density profile to evolve to a relatively long scale length (~ 150 wavelengths). The fact that the scattering continues, even while the incident intensity drops, may indicate that filamentation is occurring in the plasma. Filamentation could provide locally higher intensities (this instability is well seeded by the already highly nonuniform beam). While we have no direct evidence for this instability, x-ray photographs and backscattered-light pictures showing "hot spots" (which are different from the incident beam) suggest its presence.⁴³⁻⁴⁵

The Raman spectra show only the convective instability, with no evidence of scattered light coming from densities between 0.2 and $0.24 n_c$. The usual assumption⁵⁵ is that the density profile in this region has been steepened by the two-plasmon decay instability. Scale lengths of $10 \mu\text{m}$ or less should inhibit SRS, even in the most intense hot spots measured for our incident laser beam. We have no direct experimental

measurements of such steepening; however, measurements showing the presence of the two-plasmon decay instability are discussed below.

The energy in Raman-scattered light (spectrally integrated from 1.2 to 2 times the laser wavelength) was measured at only a few angles, making estimates of total SRS energy accurate only to an order of magnitude. Furthermore, these levels fluctuated strongly with small variations in incident energy. For the 2ω experiments, approximately 10^{-5} of the laser energy was detected as Raman-scattered light when the median incident intensity^{1,2,4,16} was 10^{14} W/cm^2 . At five times higher median intensity, the SRS fractional energy was typically 2×10^{-4} . Because of the strong intensity dependence, comparison with data from other experiments is difficult. However, the most comparable 1ω experiments, those at 10^{15} W/cm^2 , seem to show about an order of magnitude more energy in the emitted Raman light.

Two-Plasmon Decay. In the two plasmon decay instability (another source of suprathermal electrons) the incident laser light decays into two plasma waves in the vicinity of $1/4n_c$. Evidence for the two-plasmon decay is indirect. We must therefore rely on auxiliary processes to generate scattered light at $3/2$ and $1/2$ the laser frequency, as a result of nonlinear mechanisms involving the plasma waves. Quantitative interpretation is difficult; in fact, our data show that certain models⁵⁸⁻⁶⁰ do not correctly predict the observed $3\omega/2$ spectra at large observation angles. In this section we summarize our results, which are reported in more detail elsewhere.¹⁶

The $\omega/2$ measurements were made during the 2ω experiments. For these experiments, a light collector was located 30° in the plane of polarization from the incident laser beam. A fiber-optic cable transferred the light from the collector to a 0.5-m grating spectrometer. The disk targets were oriented either normal to the laser beam or normal to the spectrometer light collector. We were able to record $\omega/2$ spectra only at the highest intensities used in these experiments ($I_{0,0} = 5 \times 10^{15} \text{ W/cm}^2$, focal spot size $75 \mu\text{m}$, $E_{\text{laser}} = 150 \text{ J}$).

Figure 5-31 shows spectra from gold-disk targets recorded under the irradiation conditions just stated. The upper curve is for a

target oriented toward the spectrometer collector; the lower curve is from a similar shot with the target oriented toward the laser beam. In these experiments, the small amount of 1ω light still present in the laser beam furnished an accurate wavelength marker. For clarity, this signal has been removed from the data; its location is indicated by the line labeled " $\omega/2$."

Figure 5-32 shows the spectra for a beryllium-disk target oriented toward the collector. The laser energy for this shot was 35% higher than in the shots in Fig. 5-31; the scattered-light signal near $\omega/2$ increased by a factor of 5. We believe that this in-

Fig. 5-31. Spectra near $\omega/2$ from gold-disk target. (a) Disk oriented toward light collector. (b) Disk oriented toward laser. 30 from light collector.

Fig. 5-32. Spectra near $\omega/2$ from beryllium-disk target, with disk oriented toward light collector.

crease is due both to a higher level of two-plasmon decay waves and to the decreased signal absorption in the low-Z plasma. Beryllium targets irradiated at normal incidence show only the blue-shifted peak, similar to the gold-disk data in Fig. 5-31(b).

None of the data have been corrected for Doppler shifts, which should be small. For example, for an expansion velocity of 2×10^7 cm/s, the spectra shown should be moved ~ 15 A toward a longer wavelength to correct for Doppler effects.

We observed $3\omega/2$ spectra during 1ω irradiations of flat disk targets. The data presented here were obtained with irradiations at average intensities of 10^{15} W/cm² (90 J, 125- μ m spot, 0.7-ns FWHM). Similar data have been obtained by others.⁶⁰ The spectrometer light collector was placed either 30° or 85° from the laser beam, out of the plane of polarization.

Figure 5-33 shows $3\omega/2$ spectra obtained from gold-disk targets irradiated at normal incidence. Curve (a) shows data obtained with the spectrometer light collector at 30° ; curve (b) shows data obtained at 85° . The widths of the spectral splittings observed at 30° are summarized in Fig. 5-34 for the various target materials. We believe that the increased splitting observed with targets of higher Z is due to higher electron temperatures in these plasmas.

We obtained $3\omega/2$ spectra on a few 2ω experiments. Unfortunately, the diagnostic had relatively poor sensitivity and spectral resolution. These data will not be discussed here; we note that, as with the 1ω results, an increasing frequency deviation from $3\omega/2$ with target Z was observed.

We now discuss possible interpretations of these data, beginning with the $\omega/2$ results. In these experiments we should be above the two-plasmon decay threshold^{55,61,62} even if the density profile is moderately steep (~ 10 - μ m scale length); at least 10% of the laser energy is incident at intensities greater than 10^{16} W/cm². Therefore, we presume that a wide spectrum of plasma waves is excited, with wave numbers up to $k_1 = 0.3/\lambda_1$. (The subscript 1 refers to the higher-frequency plasma wave, and λ_1 is the Debye length.) This maximum plasmon wave number, determined by Landau damping, is typically (2 or 3) k_0 for our plasma conditions, where k_0 is the

laser-light wave number in the plasma. We consider only the case of normal incidence. The frequency shift of the plasma waves from $\omega/2$ is given by

$$\Delta\omega = \pm \frac{3}{2} \omega_p \lambda_D^2 k_0 (k_{\parallel} - k_0/2) \quad (11)$$

where $+x$ is the direction of the density gradient. The transverse (y) component of $k_{\parallel}$ must be determined from theory; it is predicted^(6,2,3) to be small, for our plasmas—on the order of k_0 or less. The exact value has only a small influence on the frequency of the plasma wave (ω_p depends weakly on the wave number), since phase matching requires that both plasma waves have the same $|k_{\parallel}|$. Therefore, we can estimate the frequency spectrum of the parametrically excited plasma waves with the only free parameter being the electron temperature.

We now state the origin of the dotted lines in Figs. 5-27 and 5-28. The lines labeled "(1)" show the predicted wavelength at which $k_{\parallel}\lambda_D = 0.3$, i.e., the maximum frequency splitting allowed by Landau damping. To fit the data, we chose a temperature of 3 keV for the gold case (Fig. 5-31) and 2 keV for the beryllium case (Fig. 5-32). The dotted lines labeled "(2)" show the location of the wavelengths at which $k_{\parallel} = k_0$ for these temperatures. This is the minimum wave number allowed for the higher-frequency plasma wave. These two representative wave numbers lead to shifts that bracket the experimental data.

The peculiar spectral asymmetry evident in the limited extent of the red-shifted peak in Fig. 5-31 is probably due to the mechanism that generates the scattered light, rather than to the two-plasmon decay itself. Several possibilities are discussed in detail in Ref. 46. We believe that the most likely mechanism is the nonlinear interaction between a two-plasmon decay wave and light at the laser frequency. For this process, we find that the red-shifted scattered light is limited to the small spectral shifts observed; the waves that would generate light with larger red shifts cannot satisfy the required phase-matching conditions.

We now turn to the mechanism for $3\omega/2$ generation. If we assume, as previously, that the plasma wave spectrum is bounded by $k_{\parallel} = k_0$ for small splittings, and by

$k_{\parallel}\lambda_D = 0.3$ for large splittings, then the data in Fig. 5-33 are fitted reasonably well by $T = 2$ keV. The data recorded at large scattering angles cannot, of course, be addressed by simple one-dimensional models. Further work is needed to quantify the theory, including realistic (at least two-dimensional) geometry, the effects of non-zero k_y , and possible interactions with back- and side-scattered laser light.

Thus our scattered-light data show the presence of the two-plasmon decay instability in laser-produced plasmas. The data are within the bounds predicted by simple models, using reasonable temperatures. For fixed laser conditions, these temperatures change as expected with target Z . The models do not attempt detailed explanations of the mechanism(s) of generation of scattered light. They do not predict how the level of

Fig. 5-33. $3\omega/2$ spectra (7095 Å) from gold-disk targets oriented toward the laser. (a) Light collector at 30° from the target normal. (b) Light collector at 85° from the target normal.

Fig. 5-34. $3\omega/2$ peak-to-peak spectral splittings for constant irradiation conditions, at ω_0 , for various target materials. The horizontal axis roughly represents increasing Z and thus increasing temperature.

scattered light is related to the level of the plasma waves, nor do they address the question of how the wave numbers of the (large-amplitude) plasma waves are related to the wave numbers predicted, by linear theory, to have the greatest growth rates. Many of the details of the spectra await further theoretical work on the mechanisms of scattered-light generation. Nevertheless, it is gratifying that the observed spectra can be bounded by simple models using reasonable wave numbers.

Mechanisms for Hot-Electron Production. The energy involved in parametric instabilities that generate plasma waves can be measured easily only in stimulated Raman scattering, since in this case one of the scattered waves is a light wave that can be detected directly (unless absorption is a problem). We examined the level of SRS in our highest-intensity (5×10^{15} W/cm²) disk target experiments, conducted at 2ω . Since this light is observed to be from the convective instability, which originates from low densities, we neglect absorption of the scattered light. The Raman-light energy was typically 0.02% of the laser energy.

We assumed that all of the plasma wave energy was converted by wavebreaking into electron kinetic energy, and further assumed that these electrons lost all their en-

ergy in the dense target material; then we calculated the resulting bremsstrahlung. Figure 5-35 shows the results of these calculations, along with the measured hard x-ray flux. In spite of our optimistic assumptions, which should lead to an overestimate, the level of SRS is not enough to account for the observed hard x-ray level. Therefore, we must conclude that in these experiments other instabilities, such as the two-plasmon decay or filamentation in conjunction with resonance absorption, are responsible for most of the suprathermal electrons. While we have no experimental means of measuring the level of the two-plasmon decay instability, we can use $3\omega/2$ and $\omega/2$ data to obtain estimates of the plasma wave numbers as indicated above. We can then estimate the wavebreaking "temperature" due to these plasma waves. When we do this for the data shown in Fig. 5-31, we obtain 33 keV, which roughly agrees with the 25-keV value we obtained from the x-ray spectrum for this shot.

We have already mentioned the possibility of filamentation. We can estimate the intensification required to bring resonance absorption temperatures up to the ~ 20 -keV level observed in the experiments. Using the $(I\lambda^2)^{0.4}$ scaling for 5×10^{15} W/cm² and 2ω , we find that the required intensification is a factor of from 1 to 4, depending on the background temperature.⁶ Thus filamentation is certainly a candidate for producing the observed hot electrons.

Summary of Wavelength-Scaling Experiments. The wavelength-scaling experiments have shown the expected increase in absorption at shorter wavelengths.⁶⁴ The x-ray conversion efficiency has also increased with the shorter wavelength, although the efficiency at low intensities is not completely understood. The level of suprathermal electrons has decreased dramatically, as evidenced by the hard x-ray spectrum. As expected, various parametric instabilities, including Brillouin⁶⁵ and Raman scattering, have decreased with shorter wavelength.

All the measurements with shorter-wavelength light show results favorable for ICF. A word of caution is in order, however. With the limited amount of energy available for these experiments, the plasmas produced were small (no larger than 100 μ m at the higher intensities). In spite of this,

Fig. 5-35. Hard x-ray spectrum and upper limit of x-ray spectrum predicted (as described in text) from SRS measurements.

evidence for parametric instabilities could still be seen, albeit at a low level. Estimates indicate that these instabilities could be a problem with Nova- or reactor-size plasmas.⁶⁶ Experiments on the 2 ω Novette facility will help to answer this question.

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X-ray Laser Flashlamp Experiments on Shiva

In 1983 we will begin a series of experiments on Novette to study the feasibility of lasing at x-ray wavelengths. The proposed scheme⁶⁷ is to use 2 ω laser light to irradiate two thin metal foils, which will then emit x rays that ionize neon gas between the foils and preferentially excite upper laser states in the ionized neon, which then lases at soft x-ray wavelengths. Appropriate stripping and pumping intensities should be obtainable with foils of Z between about 24 and 28.

Quantitative evaluation of this scheme will require that we know the characteristics of the x-rays emitted by the flashlamp foils in response to laser irradiation. The last series of experiments on the Shiva laser was performed primarily to characterize the x-ray emission spectra of some candidate flash-lamp materials.

The flashlamp targets consisted of iron, iron-copper, chromium, and nickel foils about $0.04\text{ }\mu\text{m}$ thick and 2 mm in diameter (Fig. 5-36). Each foil was deposited on an $0.5\text{-}\mu\text{m}$ -thick parylene disk that provided mechanical strength; the disk was then suspended across a thin ($25\text{-}\mu\text{m}$ -thick) Mylar hoop 3 mm in outside diameter.

Shiva's 10 lower beams were overlapped to give a nominal spot diameter of 1 mm. The laser wavelength was $1.06\text{ }\mu\text{m}$. Incident energies ranged from 142 to $37\frac{1}{2}$; typical peak intensities were $1\text{ to }2 \times 10^{14}\text{ W/cm}^2$. Pulse widths varied from 100 to 120 ps.

Target Absorption. Efficient absorption of laser light by the flashlamp foil is essential to the production of strong stripping and pumping radiation. We measured target

absorption with an array of scattered-light pin diodes (see Fig. 5-37), filtered to record wavelengths of $1.06\text{ }\mu\text{m}$, and with calorimeters for incident and backscattered light. For these measurements we oriented the iron and iron-copper targets at normal incidence, to facilitate angular integration over the pin-diode array.

Two experiments gave laser-light absorption values of $(50 \pm 5)\%$ for an iron foil and $(39 \pm 4)\%$ for an iron-copper foil. The average absorption value for these two experiments is $(45 \pm 9)\%$; the uncertainty reflects errors due to calibration uncertainties and angular integration.

The dependence of target absorption, and thus of bulk-radiation efficiency, on the angle of incidence of the laser light should be negligible (at least for angles less than 30°), because of the large effective aperture ($f/1$) of Shiva's beam clusters and the laser duration used (100 to 120 ps). This assumption is supported by experiments with the Argus laser in which we used a box calorimeter to measure reflected and scattered laser light. Absorption should not depend significantly on the atomic number of the target material, at least over the modest range of Z covered in these experiments.

Fig. 5-36. Flashlamp-foil target (metal foil deposited on a parylene disk suspended across thin Mylar hoop).

Fig. 5-37. Diagnostic instruments used in the Shiva x-ray experiments (angles are with respect to disk normal).

Fig. 5-38. (a) Densitometer tracing across spatial image of incident laser beam, showing acceptable uniformity. (b) Image of 0.8- to 1.2-keV x rays emitted from target. The spatial resolution of the x-ray microscope is about 10 μm .

Fig. 5-39. Low-resolution time-integrated x-ray spectra of laser-irradiated chromium and nickel foils.

Source Uniformity. Uniform illumination of the flashlamp foil is important in determining the uniformity with which the laser medium is pumped. A film record of the laser-beam profile [Fig. 5-38(a)] shows acceptable uniformity. We obtained time-integrated images of the foil's x-ray emission with a polar 8 $\times$ Kirkpatrick-Baez microscope, which viewed the rear of the foil target along the laser axis. An aluminum-Mylar filter in front of the microscope transmitted x rays with energies between about 0.8 and 1.2 keV.

Figure 5-38(b) shows a processed image from a nickel-foil experiment. The emission profile is satisfactorily uniform. The rela-

tively soft edge of the image may result from time integration and from overlapping of Shiva's 10 beams. The image shows no small-scale structure (time integration may again be responsible), although there is a relatively long-scale variation (1.5:1) in the film density over a scale length of about 500 to 700 μm . We also used this image to determine the area of the source in estimating the brightness of candidate pump lines from the foil, as discussed later.

X-ray Spectra. In our measurements of the flashlamp x-ray emission, we were particularly interested in the x-ray power and total yield, the angular dependence and relative strength of the x-ray lines, and the energy and brightness of candidate pumping lines. Total x-ray power and yield determine the rate at which the flashlamp x rays strip the lasing gas, and the extent of stripping to the desired helium- or hydrogen-like ionization state. The energy (wavelength) of individual flashlamp emission lines determines whether the overlap required for pumping will occur or not; the brightness of these lines determines both the possibility of creating a population inversion and the resultant laser gain.

To study these quantities, we combined measurements from several different instruments (Fig. 5-37). We measured the time history of the x-ray emission from the target (resolved to about 20 ps) using a gold transmission grating coupled to a soft x-ray streak camera.^{68,69} The spectral resolution ($E/\Delta E \approx 10$) was insufficient to resolve individual emission lines.

Figure 5-39 shows time-integrated spectra, obtained by integrating the streak-camera records, of chromium and nickel foils. The spectra display strong L-shell emission (the result of an M-shell electron filling an L-shell hole). Although the instrument does not resolve individual lines, it does show the expected increase with Z of the energy centroid of emission.

Figure 5-40 shows the time history of the L-shell emission from a chromium foil. The 100-ps rise time is consistent with a driving-laser pulse of 100 ps, but the emission decays more slowly, with a characteristic time of about 220 ps. The FWHM of the x-ray emission is about 190 ps, 1.5 to 2 times greater than that of the laser pulse. This pulse shape should provide marginally adequate x-ray power for stripping.

We recorded x-ray emission spectra from the front and rear of the foil targets with high-resolution crystal spectrographs.⁷⁰ Figure 5-41 shows a chromium spectrum from the rear of the target. We used a 932.6-eV 3d-2p transition from lithium-like chromium for energy calibration. The spectrum emitted by the laser-irradiated chromium foil was one of the most important measured in these experiments, because a Be-like Cr line, predicted to lie at 899.58 eV, was a prime candidate for pumping helium-like fluorine to its $n = 4$ level (transition energy 899.78 eV), leading to laser action at 42.3 eV.

The four crystal spectrographs (see Fig. 5-37) indicated that x-ray emission is strongest in the forward direction, which is favorable for pumping the laser cavity with a foil flashlamp.

Figure 5-41 shows the candidate pumping line, which lies among five distinctive groups of 3d-2p lines of Be-like Cr lying between about 0.88 and 0.92 eV. Despite the presence of the strong calibration line at 932.6 eV, the uncertainty in our energy calibration was ± 4 eV. Because we require that the pumping line and the transition to be pumped agree to better than 0.2 eV, we were unable to confirm that the candidate line was indeed suitable for pumping. (Subsequent, high-precision measurements at KMS Fusion, Inc., indicate that the candidate line has an energy of 900.1 eV, which makes it unsatisfactory; the 42.3-eV laser scheme has therefore been dropped.)

Line Brightness. To facilitate comparison with the kinetics calculated with the

Fig. 5-40. Time-resolved x-ray emission at the peak intensity point in the chromium spectrum of Fig. 5-38.

Fig. 5-41. High-resolution chromium x-ray spectrum from the rear of the foil, showing the $n = 3$ to $n = 2$ calibration line and a candidate laser pump line.

XRASER code, we express the brightness of proposed pump lines in terms of modal photon density (photons per mode). Most laser-photoexcitation schemes require flashlamp photon densities of at least 0.01 photons per mode.

The modal photon density is a measure of a line's brightness at a given frequency. It is convenient to compare line brightnesses with that of an optically thick source at a temperature, T ; for such a source, the modal photon density is

$$N_p = \frac{1}{e^{h\nu/kT} - 1}, \quad (12)$$

where h is Planck's constant, ν is the emission frequency, T is the temperature, and k is Boltzmann's constant. For the optically thick source, N_p increases rapidly with kT for a given photon energy; this emphasizes the importance of heating the flashlamp to a relatively high electron temperature.

The modal photon density of an emission line can be expressed in terms of measurable parameters as

$$N_p = \frac{2.56 \times 10^{-24} I (\Delta E_{\text{exp}} / \Delta E_{\text{nat}})}{E_{\text{line}}^3 \Delta \tau} \quad (13)$$

where I is the x-ray intensity in units of keV/keV over 4π , $\Delta E_{\text{exp}} / \Delta E_{\text{nat}}$ is the ratio of the experimental (FWHM) and natural (Doppler-broadened) line widths, E_{line} is the transition energy in electron volts, A is the emission area of the foil plasma in square centimeters, and $\Delta \tau$ is the duration of the transition (FWHM) in seconds.

We estimate the modal photon density of the strongest line shown in Fig. 5-40 (the calibration line at 932.6 eV) at 0.004 to 0.005. [On the basis of Kirkpatrick-Baez microscope data we assumed an emission area of $A = 6.4 \times 10^{-3} \text{ cm}^2$. In one experiment, we obtained the value $\Delta \tau = 190 \text{ ps}$ from the transmission grating and streak camera (see Fig. 5-41); in others, we assumed $\Delta \tau = 200 \text{ ps}$. The experimental line width, ΔE_{exp} , was about 1.8 eV FWHM for the rearward spectrograph; we assumed a natural line width, ΔE_{nat} , of about 0.3 eV.]

In our experiments, I is uncertain by $\pm 50\%$. We did not measure the uncertainty in $\Delta \tau$ because we did not generally determine $\Delta \tau$ for individual lines. The ratio $\Delta E_{\text{exp}} / \Delta E_{\text{nat}}$ is uncertain by $\pm 50\%$. These uncertainties lead to a large uncertainty in N_p . Future experiments at KMS and with Novette will focus on reducing the uncertainty in each of these parameters.

We estimate the brightness of the candidate pumping lines to be about half that of the calibration line, and thus well below the 0.01 photons per mode necessary for pumping. Although the data are quite uncertain, we will attempt to increase the value of N_p in the early phases of the Novette x-ray laser experiments, using a 2ω laser wavelength and the higher intensities that Novette will provide.

These experiments demonstrated that flashlamp foils can produce bright, uniform stripping radiation; the experiments were at least partially successful in characterizing the foil's line (pumping) radiation. Several problems remain; we must improve the accuracy with which we determine the flashlamp line energies and widths, and we must measure the time duration of individ-

ual lines (e.g., with a high-resolution spectrograph). Greater line brightness is required, and uncertainties in estimates of line brightness must be reduced.

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LLNL/NRL Long-Scale-Length-Plasma Foil Acceleration Experiment

An area of vital interest to the laser fusion effort is the study of plasmas and target hydrodynamics for conditions close to those required for the implosion and ignition of directly illuminated laser fusion targets. To simulate these conditions requires plasma scale lengths, $L/\lambda > \sim 1000$ (with $L \sim 1 \text{ mm}$); electron temperatures, $T_e < 1 \text{ keV}$; laser pulse lengths, $\tau_l > 3 \text{ ns}$; and laser intensities, $I_l > 10^{14} \text{ W/cm}^2$. Large scale length is desirable because it allows significant lateral energy transport and thus smooths the ablation pressure and material acceleration, thereby helping to minimize the growth of hydrodynamic instabilities. Ignition requires ablation pressures $> 10 \text{ Mbar}$, shell velocities $> 150 \text{ km/s}$ (Refs. 71 and 72), target-velocity uniformities of 0.5 to 1%, and low levels of preheat.

Large lasers are needed even to approximate these conditions. Previous attempts to achieve these conditions elsewhere have been hampered by limited energy and by the small spot size and short pulse length required to reach the intensities of interest. Experiments at the Naval Research Laboratory (NRL)^{73, 74} have attempted to model the early stages of pellet implosion through studies of planar targets ablatively accelerated by direct illumination with 500-J, 3-ns pulses from a Nd:glass laser.

In December, 1981 we undertook with NRL a series of experiments⁷⁷⁻⁸⁰ to explore a region of parameter space much closer to the region of direct-illumination ignition than had been previously explored. We studied a wide range of physics phenomena associated with the laser-target interaction and the ablative acceleration of the target material. Areas of prime importance were

- Laser-plasma coupling; absorption, stimulated-scattering mechanisms, and suprathermal x-ray production.

- Energy transport: classical or inhibited transport, lateral energy transport, suprathermal-electron transport, and x-ray- and shock-induced preheat.
- Hydrodynamics: target velocity and acceleration, scaling of mass ablation rate and of ablation pressure, foil velocity profile and uniformity, and foil decompression.

The experiments were carried out at the 20-beam LLNL Shiva laser facility. The 10 lower beams were defocused and overlaid to form an ~ 1 -mm-diam spot (at $\sim 90\%$ energy content) on a carbon foil $\sim 10\ \mu\text{m}$ thick and 2.5 mm wide. The foil was made much wider than the focal spot, to isolate the cooler rear side from the laser-irradiated (front) side. Overlapping the focal spots of the 10 laser beams should statistically smooth the intensity profile at the target by the factor $10^{1/2}$. Time-integrated x-ray microscope images indicate $\sim 50\%$ spatial-intensity uniformity on 10-beam shots in which the pulse length was reduced to 100 ps to limit lateral energy transport. The 3.0- to 3.2-ns (FWHM), 2.8- to 3.4-kJ laser pulse resulted in an average peak incident intensity of 1.0 to $1.5 \times 10^{14}\ \text{W}/\text{cm}^2$.

Figure 5-42 shows the generic target used in these experiments. Among several target variations, the simplest was a single carbon foil; others had a second (impact) foil about $200\ \mu\text{m}$ behind the first. A third variation omitted the backlighter, which sometimes interfered with other diagnostics. Arrays of diagnostics monitored the laser-target interaction, including the absorption fraction, the scattered-light distribution, the yield and spectrum of the $3\omega/2$ and stimulated Raman scatter, suprathermal x-ray production, and optical emission from the rear of the target. The motion of the accelerated foil was measured using streaked x-ray shadowgraphy^{81,82}; the double-foil technique²⁵ was used to measure the uniformity of the velocity profile.

Laser Coupling and Transport. The scattered light from the target was measured with an array of pin diodes filtered for 1ω light. The target absorption was inferred from the difference between the incident energy and the total scattered light.

Figure 5-43 shows the angular distribution of the scattered light, plotted against the polar angle, θ (measured from the laser axis). Because of the radial orientation of

Fig. 5-42. X-ray backlighting target used at Shiva for LLNL/NRL long-scale-length-plasma foil acceleration experiments.

Fig. 5-43. Angular distribution of scattered light from three experiments (curves a, b, and c) and from 2-D LASNEX predictions (curve d).

the polarization vectors of the 10 Shiva beams,⁸³ no azimuthal (ϕ) variation is expected, and none was observed. Curve a shows data from a carbon-target experiment at $I_1 = 10^{14}\ \text{W}/\text{cm}^2$, with $\tau = 3$ ns and spot diam = 1 mm. The distribution is very broad; substantial scattered-light intensity is observed up to $\theta = 70^\circ$. Curve b shows the

scattered-light distribution obtained with iron targets ($I_l = 10^{14}$ W/cm², $\tau = 100$ ps, spot diam = 1 mm). Curve c shows the results of earlier CH-target experiments carried out at NRL with higher peak intensities ($I_l \sim 2.5 \times 10^{14}$ W/cm²), similar pulse widths ($\tau \sim 3$ ns), but smaller focal spots (diam < 100 μ m).⁵⁴ The distributions shown in curves a through c are all normalized to 3-keV incident energy. Curve d shows the scattered-light distribution predicted by two-dimensional calculations with the LASNEX hydrodynamic code⁸⁵ for the experimental conditions of curve a with no stimulated-scattering model included.

The angular breadth of the scattering shown in curve a of Fig. 5-43 strongly suggests that stimulated Brillouin sidescattering (SBS) is occurring in these large (axial and transverse) scale length plasmas.^{86,87} For the coronal conditions in these experiments (as estimated using either LASNEX or simple scaling arguments, with the amplification length equal to the spot size), the number of growth lengths⁸⁸ for SBS is greater than 10. Strong SBS would not be expected to occur in either the short-pulse or small-radial-scale-length experiments shown in Fig. 5-43 (curves b and c).

Evidence for SBS was found in earlier 1-ns single-beam experiments, in which the scattering out of the plane of polarization depended strongly on laser spot size.⁸⁷ The lack of a unique plane of polarization at the Shiva facility prevented us from obtaining such evidence for stimulated sidescatter in these experiments, so large-angle refraction of the incident light cannot be ruled out. However, such large-angle scattering is not consistent with large-angle refraction, in this largely one-dimensional corona, according to simple estimates of refraction or to two-dimensional LASNEX simulations (Fig. 5-43, curve d) of light-ray refraction using geometric optics alone, and with no model for SBS included in the simulations.

The total scattering, obtained by integrating the scattered-light intensity (Fig. 5-43, curve a) overall solid angle, gives a target absorption of $60\% \pm 10\%$. One-dimensional LASNEX simulations incorporating inverse bremsstrahlung, a phenomenological resonance-absorption model (described below), and a model for stimu-

lated Brillouin scattering (backscatter)⁸⁹ predict an absorption of 65% with a thermal flux limiter⁹⁰ of 0.5 (which corresponds to classical heat flow). If the flux limiter is reduced to 0.03, the predicted absorption drops to 53%. Both one-dimensional and two-dimensional calculations without SBS give absorption fractions greater than 95%, regardless of the value of the flux limiter. Thus the agreement of absorption calculations (including SBS) with the observed absorption implies the presence of significant SBS.

We measured the level of suprathermal x rays produced by electrons heated by various laser-plasma coupling processes, with a multichannel filter-fluorescer x-ray spectrometer.⁹¹ Figure 5-44 shows the experimental results and those of a LASNEX simulation. If we assume that the x rays are produced by Maxwellian electrons stopped in the overdense carbon target, the data indicate that about 80 to 100 J (I_{11}) of laser energy is absorbed into an 8- to 10-keV (T_{11}) suprathermal electron distribution. Suprathermal electron losses by other mechanisms (such as fast ions) are not accounted for, so that the E_{11}/I_{11} fraction is a minimum value.

The LASNEX simulation includes a suprathermal-electron model based on resonance absorption only, with a peak intensity of 10^{14} W/cm²; 30% of the light that reaches n_c goes into the production of suprathermal electrons⁹² whose temperature, T_j , scales as $(\lambda \lambda_D)^{1/4}$, where l is the intensity at n_c . Figure 5-44 shows a considerable discrepancy between the measurements and the LASNEX predictions. Much higher intensities (by a factor of ~ 15) are required if resonance absorption is to explain the data. However, no bright spots are seen in near-field photos of the Shiva beams. These photos, which should be representative of an individual Shiva beam in the target plane, show only a 75% peak-to-valley intensity modulation $[(r_{\max}/r_{\min}) - 1]$. (Interference effects due to the overlapping of the 10 Shiva beams are not considered.) Furthermore, x-ray microscope photographs taken during 100-ps target irradiations show only a 2:1 modulation in the x-ray emission, indicating a similar smoothness in the beam profile. Thus, unless

the beams have filamented in the long, underdense plasma, resonance-absorption modeling does not explain the results.

A more plausible explanation for the suprathermal electrons observed in these experiments is the occurrence of processes such as two-plasmon decay ($2\omega_{pe}$), stimulated Raman scattering (SRS), and filamentation.⁸⁸ Gradient-stabilization threshold arguments indicate that the $2\omega_{pe}$ instability should be present in these experiments, but not SRS. (However, if substantial beam filamentation occurs for $n < 1/4n_c$, the SRS threshold may be exceeded as well.) Simple theories and particle simulations extrapolated down to these moderate intensity levels in fact predict that the $2\omega_{pe}$ instability would produce T_{11} values in the observed range.⁷⁸

We looked for the signature of the $2\omega_{pe}$ instability using arrays of pin diodes, filtered at $3\omega/2$ with 300- to 350-Å bandpass interference filters. We observed substantial levels of $3\omega/2$ light in the 3-ns experiments, representing a total yield of about 0.03 to 0.04% of the incident laser energy. The distribution of the $3\omega/2$ light was approximately isotropic, as was that of the ω scattered light. This distribution is not unreasonable, because the sidescattered SBS light could be involved in the generation of the $3\omega/2$ light. In the 100-ps experiments, the $3\omega/2$ yield was about 2 to 6×10^{-6} of the corresponding laser light, corresponding to a generation efficiency about 1% of that seen in the 3-ns experiments.

Although the details of the production of $3\omega/2$ light are not well understood, it is generally assumed⁸⁸ that the light arises from the presence of electron plasma waves near $1/4n_c$. One source of these waves is the $2\omega_{pe}$ instability, whose intensity threshold should be exceeded in these experiments. The $3\omega/2$ light arises from second- and third-order processes involving the coalescence of an electron plasma wave and an electromagnetic wave or of three electron plasma waves. The level of this light depends on such parameters as the localized density fluctuations and the size of the interaction region near $1/4n_c$. Since these quantities are not known, it is not possible to determine the efficiency of the $2\omega_{pe}$ instability from the energy of the $3\omega/2$

Fig. 5-44. High-energy x-ray spectra (LASNEX spectrum calculated with nominal resonance absorption model).

light. However, the large enhancement in the relative $3\omega/2$ emission observed as the laser pulse width is increased from 0.1 to 3 ns suggests that processes such as $2\omega_{pe}$ become more important as the corona dimensions increase.

The absolutely calibrated optical (pyrometer) streak camera (filtered for ~ 4200 Å), which images the rear surface of the single-foil target, can also provide clues to the source of the hot-electron production. It does this by time-resolving the optical signature of the preheat, which results from the creation and rapid transport of the hot electrons. Figure 5-45 shows the time of preheat onset, whose significance is discussed below. The figure also shows the time-dependence of several coronal quantities, as obtained from a two-dimensional LASNEX simulation: electron temperature, T_e ; scale length of $1/4n_c$ (scaled to a linear density ramp from n_i to vacuum); laser intensity, I_l , at $1/4n_c$; and intensity threshold, $I_{2\omega_{pe}}$, for the two-plasmon decay instability.

In this LASNEX simulation, 40% of the energy was thrown away to simulate SBS losses, and a flux limiter of 0.5 was used. At these moderate intensities, decreases in the flux limiter change the profiles only slightly. The temperature remains near 1 keV during the few nanoseconds surrounding the laser

peak. The laser intensity at $1/4n_c$ follows the broad Gaussian shape of the incident pulse. The plasma scale length L increases strongly with time, however, so that the intensity threshold for the $2\omega_c$ instability due to gradient stabilization, which scales as $1/L$, drops rapidly.²⁶ The laser intensity crosses that threshold at about 1 ns ($t = 0$) is the peak of the laser pulse! Within the parameters of the modeling (for example, the SRS depleted intensity at $1/4n_c$ is not considered), the preheat onset time scales almost with the calculated $2\omega_c$ threshold crossing.

For the nominal values of the laser intensity, the absolute Raman threshold is still too high in fact (it is at the top of Fig. 5-45) to be $\sim 2 \times 10^{14}$ W/cm² at $t = 1$ ns. The number of plasma gain lengths G for parametric amplification also grows with L , however, $G \sim 5$ for $t \sim 1/4n_c$, so that filamentation may be contributing to the electron generation.²⁷

We looked for SRS using an array of eight cameras, filtered to respond to wavelengths between 1.5 and 2.6 μm . At 95% scattered light was detected (the minimum detectable yield was about 0.00125%). Given the 80 to 90% of supra-thermal electrons inferred from the x-ray and optical data, the Manley-Rowe rela-

tions²⁸ indicate that SRS (if it is occurring) should give a similar amount of scattered-light energy. If the SRS occurs near $1/4n_c$ however, substantial collisional absorption of the 2.13- μm light can occur. The inverse bremsstrahlung absorption length for this scattered light (assuming the coronal conditions calculated by LASNEX) is 250 to 300 μm . This is to be compared with the calculated scale length of 1000 to 2000 μm near the peak of the laser pulse. (The light could also be resonantly absorbed in a nipped, magnetized $1/4n_c$ surface.)

Optical measurements with the absolutely calibrated streak camera mentioned earlier provide another way of estimating the temperature and energy content of the hot electrons.²⁹ The streak-camera images show that the heating is initially confined to the center of the laser illuminated region, indicating that little energy is being transported around the target foils.

Time-resolved rear-surface temperatures for single- and double-foil targets are given in Fig. 5-46. (The absolute calibration of the emissivity from which the temperatures were calculated is accurate to within a factor of 2 for these measurements. This corresponds to accuracies of a factor of 2 for $T_e = 2$ eV, with better accuracy below 2 eV.) The carbon target was 11 μm thick in both cases; an 8 μm thick carbon impact foil was used in the double-foil experiment. The rear surface of the single foil exhibits a sudden temperature rise soon after initiation of the laser pulse. LASNEX simulations indicate that this early heating is due to shocks produced by the leading edge of the laser pulse. As discussed in Ref. 77, shock heating and radiation preheat of the laser-irradiated foil produce temperatures of up to 5 eV. The continued, slower rise of the rear surface temperature of the single foil to about 10 to 15 eV is consistent with heating by transport through the target of electrons with the above mentioned T_e and T_{10} , where $T_{10} = 80$ to 100 eV and $T_{10} = 8$ to 10 eV (Ref. 76). Hot electrons can also account for the observed preheat of the impact foil before collision. The rear surface of the impact foil is further heated after collision by high pressure shocks. Temperatures of up to 50 eV were observed for a 17- μm -thick carbon foil irradiated at 2×10^{13} W/cm² and impacting at ~ 60 km/s.

Fig. 5-44. LASNEX calculations of the quantities in the caption, including the time to threshold for the two-plasma-decay instability. The peak of the laser pulse is at $t = 0$.

Laser-Driven Foil-Acceleration Experiments. During these experiments we used a streaked axisymmetric x-ray microscope (ASXRMS)^{80,96,97} to observe the laser-driven ablative acceleration of the carbon-foil targets to velocities greater than 10^7 cm/s. Figure 5-40 in last year's Annual Report⁸⁰ shows the target geometry. The target had a density of about 2 g/cm^3 , a thickness of $10 \mu\text{m}$, and a width of 2.5 to 3 mm. Measured areal densities were used in all calculations. Impact foils, typically about $8 \mu\text{m}$ thick and 2.5 to 3 mm wide, were mounted about $200 \mu\text{m}$ behind the target foil. In some cases the impact foil was omitted, and only the target foil was observed.

The carbon foils were backlit with $\sim 2.9\text{-keV}$ line radiation from a laser-irradiated palladium backlighter foil mounted on the same structure as the target and impact foils and about 3 to 4 mm away from them. A $75\text{-}\mu\text{m}$ -thick beryllium foil transmitted the 2.9-keV radiation but shielded the carbon foil from preheat by soft x rays ($<1 \text{ keV}$) and suprathermal elec-

trons ($<50 \text{ keV}$) from the backlighter. The backlighter was irradiated at $2 \times 10^{14} \text{ W/cm}^2$ by eight to ten beams from Shiva's top beam cluster, overlapped to provide statistical smoothing. The backlighter pulse was as long as the drive pulse, but was delayed by 1 ns to permit more of the target's trajectory to be seen (the target moves only a short distance during the first 1 ns).

The pressure ($\sim 6 \text{ Mbar}$) of the ablated carbon drove the foil away from the laser. The foil was much wider than the laser spot, and the part of the foil not irradiated by the laser did not move. The figure cited above shows a streaked image of a single-foil target accelerated to $1.2 \times 10^7 \text{ cm/s}$. The light region in the photo is the backlighter, which filled the camera's field of view. The dark triangular area that gets wider from left to right is the shadow of the target. The lower edge of the triangle represents the part of the foil that did not move.

Preliminary analysis of the backlighting data from the single-foil experiments involved fitting second- and third-order polynomials to the upper edge of the shadow of

Fig. 5-46. Brightness temperature of the rear surface of an accelerated single foil and the impact foil of a double-foil target.

the moving carbon foil.⁸¹ First a smooth curve was fitted to the photograph by eye. This was done because we did not yet have image-processing software to extract the edge and correct for backlighter intensity variations in the presence of substantial statistical noise (maximum signal levels were about 100 detected photons per pixel). Curves obtained in this way appear to be quite reproducible and accurate to about one pixel (~ 6 to $7 \mu\text{m}$). The smooth curve was digitized, and the polynomial fits were performed with an interactive program on the VAX computer in the fusion experiments analysis facility. The position and velocity histories looked reasonable and agreed well with the LASNEX calculation. However, small errors in the original position data lead to increasingly large errors for higher derivatives of the acceleration data. Because of these errors, we extracted only average accelerations (which appeared good) and only a limited amount of $a(t)$ information.

One approach to obtaining accurate $v(t)$ and $a(t)$ curves is to establish theoretically the functional form of the trajectory and then adjust the free (but empirically constrained) parameters to fit the data. We have tried this approach with reasonable success. The procedure is carried out by formulating the equations of motion (the rocket equations, which arise from Newton's second law) using theoretical relations⁹⁸⁻¹⁰⁷ for the dependence of ablation pressure, blowoff velocity, and mass ablation rate on laser intensity. These relations, which are supported empirically¹⁰⁸ and by LASNEX calculations, are

mass ablation rate

$$\dot{m}(t) = \dot{m}_0 [I_a(t)]^\alpha, \quad (14)$$

blowoff velocity

$$v(t) = v_0 [I_a(t)]^\beta, \quad (15)$$

ablation pressure

$$P(t) = \dot{m}_0 v_0 [I_a(t)]^{\alpha+\beta}, \quad (16)$$

where v_0 and $\dot{m}_0$ are the blowoff velocity (in cm/s) and mass ablation rate (in

g/cm²·s) at 10^{14} W/cm² effective absorbed intensity. (Note that v_0 and $\dot{m}_0$ are *not* initial values.)

In the theoretical literature,⁹⁸⁻¹⁰⁷ the mass-ablation scaling parameter, α , ranges from about 0.50 to 0.56, and the blowoff-velocity scaling exponent, β , ranges from about 0.22 to 0.25. The values are empirically consistent¹⁰⁸ with $\alpha = 0.6$ and $\beta = 0.2$; we assumed these values in the analysis described below. The effective absorbed intensity, $I_a(t)$, is defined as follows:

$$I_a(t) = \frac{I_i I_0}{c}, \quad (17)$$

where

$$\int_{-\infty}^{\infty} I_a(t) dt = \frac{E_i}{A} \times 10^{-14} \quad (18)$$

and where

- I_a = the effective absorbed laser intensity (in units of 10^{14} W/cm²),
- I_0 = fraction of laser light absorbed,
- c = ratio of the effective spot area to the nominal beam-spot area,
- I_i = the incident laser intensity (in units of 10^{14} W/cm²),
- E_i = incident laser energy (J),
- A = nominal laser-spot area (cm²).

The shape of $I_a(t)$ is obtained from optical streak-camera records of the laser pulse, and the magnitude is obtained from calorimetry. We assume a constant absorbed fraction f_a .

The effective spot area is determined from the area of the shock breakout at the back of the impact foil (on double-foil shots), as indicated by the streaked optical pyrometer. Light is emitted when the shock emerges from the back of the impact foil. The fastest part of the target foil reaches the impact foil sooner and therefore causes an early shock breakout; slower portions arrive later. We define the effective spot area as the target area that is accelerated to at least 60% of the peak velocity. The spot area measured in this way was 2.2 times larger ($c = 2.2$) than the nominal laser spot area (determined from laser-alignment photos); this corresponds to a spot diameter 1.5 times larger. This enlargement is

reasonable: the plasma blowing into the $f/1$ cone of the incident laser light at 10^7 cm/s for 3 ns would increase the spot diameter from 1 to 1.3 mm, neglecting any effects of refraction or lateral transport. The larger effective spot area and lower effective laser intensity is presumably due to any or all of several energy-transport mechanisms, including light refraction in the blowoff plasma, lateral thermal conduction, Brillouin sidescatter, and suprathermal-electron transport.

The rocket equations and the scaling laws [Eqs. (14) to (16)] yield the scaling law for the foil acceleration, $a(t)$:

$$a(t) = \frac{P(t)}{M(t)} = \frac{\dot{m}_0 v_0 [I_0(t)]^{1/\epsilon}}{M_0 - \dot{m}_0 \int_0^t [I_0(t')]^{1/\epsilon} dt'} \quad (19)$$

$M(t)$, the denominator on the right-hand side of Eq. (19), is just the initial areal mass, M_0 , reduced by the time-integrated mass ablation. In our experiment the ablated mass is small compared to the initial mass, so the numerator of Eq. (19) dominates the time dependence of $a(t)$. Thus a measurement of $a(t)$ is nearly a direct measurement of the pressure, with only a small correction necessary. The velocity and position are obtained by integrating Eq. (19) with zero initial velocity and acceleration:

$$v(t) = \int_0^t a(t') dt' \quad (20)$$

$$x(t) = \int_0^t v(t') dt' \quad (21)$$

Equations (19) to (21) can be integrated in closed form for some simple approximations to $I_0(t)$. In this experiment, I_0 was measured (as were f_s , ϵ , and M_0), and it was advantageous to integrate the equations numerically. We adjusted $\dot{m}_0$ and v_0 to give a best fit to the data.

Figures 5-47(a) through 5-47(d) show the time histories of the observed laser intensity, foil acceleration, foil velocity, and foil position, respectively, positions from the LASNEX calculations are also shown. The center of the laser pulse occurs at 3.8 ns. The timing, relative to the analytic model and to the LASNEX calculations, is absolute to within about 100 ps.

The positions given by Eq. (21), with the best-fit values $\dot{m}_0 = 3.17 \times 10^5$ g/cm²·s and $v_0 = 5.25 \times 10^7$ cm/s, and with $f_s = 0.6$ and $\epsilon = 2.2$, matched the polynomial fit to within about $2 \mu\text{m}$ over the entire region measured by the streaked x-ray microscope (2.5 to 6 ns in Fig. 5-47). The peak pressure was 6.3 Mbar ($\pm 15\%$).

The measurement of the pressure is sensitive to the product $\dot{m}_0 v_0$, which is directly related to the acceleration and the pressure but not to either factor separately. This is because as long as the product remains constant, either factor can vary as the reciprocal of the other without significantly changing the result. This is of course only true if the ablated mass is small, as is the case in these experiments. If the ablated mass were significant, the experiment would be much more sensitive to the parameters in the denominator of Eq. (19), $\dot{m}_0$ and v_0 . The best-fit values for $\dot{m}_0$ and v_0 agree quite well with the values given in Ref. 108 when corrections are made to account for experimental differences in spot-size definition.

This measurement provides accurate values for the ablation pressure and the efficiency with which laser light is coupled into kinetic energy. At a terminal velocity of 1.2×10^7 cm/s [see Fig. 5-47(c)], the central 1-mm-diam portion of the target foil has a kinetic energy of 90 J, or about 7% of the laser light energy incident on that portion of the foil. With a preheat level of about 15 eV (measured by streaked optical pyrometry), the foil has about 10 times as much directed kinetic energy as internal thermal energy.

X-ray backlighting experiments were also performed on double-foil targets, as shown in Fig. 5-48. In these experiments the irradiated foil reached a velocity of 7×10^6 cm/s before impact with the second foil a little more than halfway through the laser pulse. At the time of the collision, both foils had nearly equal areal densities of 1.5×10^5 g/cm². The collision gave the second foil a velocity of about 7×10^6 cm/s, as shown in Fig. 5-48. The decompression of the impact foil due to preheat, at about 10^7 cm/s, is also shown in Fig. 5-48. This velocity is consistent with the preheat temperatures measured with the optical pyrometer.

Streaked Optical Pyrometry. We determined the velocity profiles across the accelerated targets using the double foil technique with a second (impact) foil

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200 μm behind the first. Nonuniformities in the impact time of the accelerated foil are monitored by observing the light emitted from the rear of the impact foil.

Figure 5-49 shows a streak record of light emitted from the rear of the impact foil after being struck by a carbon foil ablatively accelerated (by an incident intensity of 10^{14} W/cm^2) to about 10^7 cm/s . X-ray shadowgraphy showed that the collision occurred near the end of the laser pulse. We

determined target-velocity nonuniformities from the measured target velocity, the foil spacing, and the differences in impact times of different sections of the target. The velocity nonuniformities can be resolved into an 8% tilt across the central 1 mm of the target and modulations with 7% peak-to-valley amplitude $[(v_{\text{max}}/v_{\text{min}}) - 1]$. X-ray microscope pictures show that the laser beam was brighter on one side than the other; we speculate that this could account

Fig. 5-47. (a) incident laser pulse, (b) acceleration history, (c) velocity history, and (d) position history for carbon foil, with corresponding LASNEX results.

Fig. 5-48. Streaked x-ray microscope record of a colliding-foil experiment.

for the tilt in the velocity profile. When sampled over short scale lengths (50 to 200 μm) in the central 800- μm region irradiated by the laser, the peak-to-valley velocity nonuniformity was only 5%. The resolution limit for velocity nonuniformities was about 3%. (This resolution limit, which was due to limited time resolution of the streak camera, was manifested in granularity of the streak photographs.) The actual target-velocity uniformity over short scale lengths may thus have been better than 5%.

Factors that could cause velocity nonuniformities include hydrodynamic instabilities and nonuniformities in laser irradiation and target mass. Mass thickness (areal density) nonuniformities are probably not significant, because they are less than 2% across the focal diameter in carbon foils typical of those used in the experiment. The laser beams were defocused and overlapped to minimize irradiation nonuniformities. Each beam had intensity nonuniformities with peak-to-valley amplitude ratios of about 2:1 (for scale lengths from about 0.1 to 0.5 of the beam diameter) in the near field. Ignoring contributions from phase nonuniformity, this would produce modulations of similar amplitude at the 1-mm focus. Averaging due to beam overlap should decrease the irradiation nonuniformity on the target by a factor of about $10^{1/2}$. No actual measurement was made at the target plane, however. We estimate that the on-target illumination nonuniformities had peak-to-

valley amplitude ratios of about 30%. The observed short scale-length velocity nonuniformities are at least 4 to 6 times smaller than the estimated intensity nonuniformity indicating significant lateral smoothing.

One dimensional hydrodynamic calculations and theories^{17,18} predict that the steady state absorption to ablation separation should be about 1 mm at 10^{14} W/cm^2 . The divergence of the blowoff plasma and the limited laser pulse length prevented us from attaining a separation this large in our experiment. A 9 ns pulse will be required to reach steady state. Two dimensional LASNEX calculations for the conditions of the experiment predict separations of 200 to 300 μm . The observed lateral smoothing appeared to be consistent with this prediction. Lateral heat conduction, which helped smooth out beam variations, should have also been manifested by the acceleration of a larger region of the target foil than the laser focal spot. The double-foil collision data, in fact, show that an area 2.2 times that of the laser focal spot was accelerated to 60% of the peak target velocity, in agreement with LASNEX calculations.

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Fig. 5-49. Streak record of light emission from the rear of an impact foil and corresponding isodensity trace (at 40% of peak).

Measurement of Electron Distribution Function Using a Free-Bound Continuum

We have investigated laser energy deposition effects important to inertial confinement fusion by studying the instantaneous electron distribution in the corona. Langdon¹¹⁰ has pointed out that non-Maxwellian distributions can arise and cause changes in laser-target coupling. We study the electron distribution by measuring the time-resolved shape of the K x-ray radiation recombination continuum.¹¹¹

Figure 5-50 shows the experimental arrangement. We used the single-beam laser at the Rutherford Appleton Laboratory's Central Laser Facility to irradiate an Al disk target. The laser pulse length was 1 ns (FWHM); wavelength was $0.53 \mu\text{m}$. The

laser was focused to a minimum spot diameter of $\sim 50 \mu\text{m}$, which produced intensities $I_l \sim 0.5$ to $1.0 \times 10^{16} \text{ W/cm}^2$. Intensity variations were due to fluctuations in laser energy output (15 to 30 J).

The laser beam struck the target at normal incidence. A streaked x-ray spectrograph,¹¹² looking at right angles to the incident beam, recorded the x-ray emission. The spectrograph uses a rubidium acid phthalate (RAP) crystal to diffract photons into an x-ray streak camera. The range of x-ray energies measured is defined by the geometry of the spectrograph and by the width (2 cm) of the streak-camera cathode. The range chosen (2.2 to 3.0 keV) covers the Al K α recombination edge and higher energies. Intensity calibration of the streak spectrograph was accomplished by normalizing to a spectrum obtained simultaneously using a conventional time-integrating spectrograph.

As shown in Fig. 5-50, the region viewed by the spectrograph was restricted by a mask of 25- μm -thick Ti, which is opaque to any x-rays of interest. Density profiles obtained by interferometry¹¹³ for laser-produced plasmas generated under similar conditions indicate that this mask ensured that we looked only at the underdense corona ($n_e \sim 4 \times 10^{21} \text{ cm}^{-3}$).

The signal-to-noise ratio was high enough to permit observation of the spectrum for about 2 ns, considerably longer than the laser pulse length. The streak timing cannot be absolutely correlated with the laser pulse, but the time history of the x-ray emission gives an approximate correlation.

Figure 5-51 shows an x-ray spectrum, observed during the laser pulse. The spectrum is flat over the whole irradiation time. (The rolloff above 2.9 keV is due to the limited width of the streak-camera cathode slit.) For comparison, we include the slope of the spectrum of a thermal plasma at 1 keV, the value obtained from the time-integrating spectrograph that viewed the entire plasma. The observed spectrum is consistent with a free-bound spectrum for $kT > 6 \text{ keV}$, or with a nonthermal electron distribution.

Langdon¹¹⁰ demonstrates that nonthermal electron distributions can be produced in inverse bremsstrahlung absorption. Simulations show that an initially monoenergetic electron distribution, v_{te} , under the perturbation of electron-ion collisions (assuming

Fig. 5-50. Experimental arrangement and example of data from x-ray crystal spectrograph.

Fig. 5-51. Al K x-ray spectrum during the laser pulse.

Fig. 5-52. Al K x-ray spectra after the laser pulse.

$Z \sim 1$ and neglecting electron-electron collisions), progress toward a flat (or constant) distribution for $0 \leq v \leq v_0$. Langdon shows that this occurs if the electrons gain energy from the laser faster than they can thermalize (through electron-electron collisions) or if

$$Z \left(\frac{e E_1}{m_e \omega_1 v_1} \right)^2 \gg 1, \quad (22)$$

where E_1 is the laser energy, ω_1 is the laser frequency, and v_1 is the average ("thermal") electron velocity.

The photon distribution corresponding to a Langdon-type electron distribution is

$$\log n_\gamma = \text{const}, \quad h\nu \leq I_1 + \frac{1}{2} m v_0^2 \quad (23)$$

and

$$\log n_\gamma \approx 0, \quad h\nu > I_1 + \frac{1}{2} m v_0^2, \quad (24)$$

where I_1 is the ionization energy—that is, a flat distribution with a rapid decrease above some cutoff photon energy.¹¹⁰ By contrast, the Maxwell-Boltzmann electron distribution gives the familiar photon distribution

$$\log n_\gamma \propto -h\nu/kT. \quad (25)$$

The energy range covered in our experiment did not permit us to observe the rapid cutoff, but the data are clearly consistent with the Langdon-type electron distribution.

After the laser pulse, both the total intensity and the slope of the free-bound spectrum decrease rapidly. Figure 5-52 shows two typical spectra; the corresponding electron distributions are Maxwellian, with tem-

Fig. 5-53. Observed and predicted temperature history of the cooling plasma after the laser pulse.

peratures of 210 eV at 1700 ps and 90 eV at 2300 ps. (The hump at about 2.9 keV is caused by cathode inhomogeneity; it is accounted for in the analysis.)

The cooling rate of the coronal system is a sensitive monitor of the conditions in the underdense plasma. Indeed, the energy flow of the ablating plasma can be monitored by these observations. Figure 5-53 shows the observed and predicted temperature history of the cooling plasma; the calculation was made using the MEDUSA one-dimensional code,¹¹¹ which assumes an adiabatic expansion. The agreement is excellent.

In summary, we have demonstrated, for the first time, a technique for measuring the instantaneous electron distribution in the underdense region of a laser-produced plasma. During the laser pulse, we observe a very flat x-ray spectrum, which is consistent with nonthermal heating of the corona (or with a Maxwellian electron distribution at a temperature much higher than

simulation predicts). After the pulse, the spectrum is that of a Maxwellian electron distribution, and the coronal cooling rate agrees with the results of computer simulation.

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Measurement of Electron Temperature and Electron Density Using X-ray Line-Intensity Ratios

X-ray spectroscopy has been applied to a wide range of different plasma sources.¹¹⁵⁻¹¹⁷ We have attempted to widen the application of x-ray spectroscopy by using seed materials to diagnose laser-produced plasmas. We use low seed concentrations: this leaves the laser-plasma interaction and energy transport unchanged, and it minimizes opacity effects, thus reducing interpretation difficulties. Using a variety of seed materials allows us to probe plasma parameters over wide ranges.

These experiments were performed in collaboration with Imperial College, London, and the Central Laser Facility¹¹⁸ at the Rutherford-Appleton Laboratory. We used the single-beam 0.53- μm laser with 1-ns pulses of about 20 J. The beam was focused to a 50- μm spot size.

In our initial studies, reported here, we introduced silicon into laser-produced CH

and PbO plasmas to investigate the effects of these plasmas on the silicon x-ray spectrum. Two principal observations can be drawn from these studies

- Silicon lines from higher ionization states are more intense in the PbO plasma, indicating a higher electron temperature than in the CH plasma. This is to be expected,¹¹⁹ because the laser energy is more efficiently absorbed in the PbO (because of its higher Z) than in the CH.
- The silicon spectrum cannot be interpreted in terms of a unique electron temperature. We believe that this is due to the strong spatial gradients that exist in most laser-produced plasmas. These gradients are a persistent difficulty in the use of x-ray line spectra for laser-plasma diagnostics.

Three principal diagnostics measure the x-ray emission from targets in the Central Laser Facility. These are a time-integrating, spatially resolving spectrograph, a time-resolving x-ray crystal spectrograph,¹²⁰ and an x-ray pinhole camera. Both spectrographs view the planar disk targets nearly edge-on and along an axis perpendicular to the incident laser beam. In the time-integrating spectrograph a 50- μm slit, 5 mm from the target, spatially resolves the x-ray spectrum along the direction of the incident laser beam with a magnification of from 5 to 7. The time-resolving spectrograph has no spatial resolution, but gives 80-ps resolution using an x-ray streak camera. The pinhole camera, 20 cm from the incident laser direction, images the plasma emission at around 1.5 keV using a 10- μm pinhole. The

Fig. 5-54. X-ray spectrum of H- and He-like silicon lines emitted from laser plasma. (a) silicon in CH; (b) silicon in PbO.

size of the irradiation area can be determined from the size of the pinhole image, and some information on the size of the coronal blow-off can be obtained.

The targets consisted of disks of CH and PbO. The CH was impregnated with SiO_2 ; the resulting silicon concentrations were less than 5 at.%. The seeding of PbO with silicon was achieved using an ordinary lead-silicate glass, in which the silicon concentration was 23 at.%. This concentration of silicon in the PbO targets was needed to obtain measurable silicon signals.

Examples of the silicon K spectra from CH and PbO plasmas are shown in Fig. 5-54. The spectra are from the interaction region, defined by the region of most intense emission. The density in this region is expected to be around n_e (about $4 \times 10^{21} \text{ cm}^{-3}$ for 0.53- μm light). The PbO spectrum has a stronger background because of greater bremsstrahlung and M-band emission from the Pb around 2.2 keV. The background in the CH spectrum is near the baseline set by film fogging. Spectra from the coronal regions of both plasmas were similar to those shown in Fig. 5-54, except that intensities were lower; the continuum level in the PbO spectrum is near background, indicating an absence of bremsstrahlung and M-band emission.

Relative intensities of various lines shown in Fig. 5-54 have been analyzed to estimate electron temperatures and densities. Ratios of line intensities from the two spectra are listed in Tables 5-3 and 5-4. Line identifications are those indicated in Fig. 5-54. The Li-like Si^{11} dielectronic satellite lines are not resolved, and only the sum of their intensities is used in the analysis. Some of the Si^{11} satellites are unresolved from the He-like $^3\text{P}_1$ peak. We have included the effects of this overlap in the analysis below.

We have compared the line ratios with predictions obtained from a steady-state level-population model.¹²¹ For ionization stages through Be-like ions, only ground states were included; for Li-, He-, and H-like ions, both excited and ground states were included. For Li-like and He-like ions, autoionizing states that form satellites of the resonance transition of the next ionization state were also included.

Results of the calculations for the ratio of resonance lines from successive ionization

stages are plotted in Figs. 5-55 and 5-56. Figure 5-55 shows the intensity ratios $I_{\gamma \beta/2} / I_{\gamma \alpha/3} {}^1\text{P}_1$ vs electron temperature for two assumed electron densities. The intensity ratios depend on both temperature and density in this region because the plasma state is intermediate between coronal equilibrium and LTE, so that dielectronic recombination may be important.

The line-intensity ratios from Table 5-3 are indicated in Fig. 5-55 by the horizontal arrows. Both sets of CH plasma line intensity ratios give $T_e = 0.84 \text{ keV}$ at $n_e = 10^{20} \text{ cm}^{-3}$ and $T_e = 0.63 \text{ keV}$ at $n_e = 10^{21} \text{ cm}^{-3}$. That is, the electron temperature implied by two ratios is the same for a given electron density, but the implied temperature changes with density. For PbO the two ratios imply different temperatures. The ratio for the $n = 3$ states ($I_{\gamma \beta/3} / I_{\gamma \alpha/3} {}^1\text{P}_1$) implies $T_e = 1.4 \text{ keV}$ at $n_e = 10^{20} \text{ cm}^{-3}$ and $T_e = 0.73 \text{ keV}$ at $n_e = 10^{21} \text{ cm}^{-3}$. The ratio for the $n = 2$ states gives a slightly lower temperature; this is probably due to trapping of the $\text{Ly}\alpha$ line in the PbO plasma. Trapping occurs in the PbO plasma because its silicon content is higher than that of the CH plasma and because the PbO is apparently hotter, producing a greater H-like population.

Figure 5-56 shows the calculated ratio of the intensity of the He-like resonance line to that of the Li-like dielectronic satellites.

Lines	CH plasma		PbO plasma	
	Intensity ratio	T_e (keV)	Intensity ratio	T_e (keV)
$\text{Ly}\alpha/2 {}^1\text{P}_1$	0.79	0.83	0.73	0.83
$\text{Ly}\beta/3 {}^1\text{P}_1$	0.25	0.85	0.45	1.05
$2 {}^1\text{P}_1/\text{Si}^{11} \text{ a}$	8.5	0.40	130	0.53

^aComponents ABCD + IKL + QR (notation of Ref. 115).

Table 5-3. Line- $\blacktriangle$ intensity ratios and corresponding electron temperatures, T_e , in CH and PbO plasmas.

Table 5-4. Line-intensity ratios and corresponding electron densities n_e in CH and PbO plasmas. $\blacktriangledown$

Lines	CH plasma		PbO plasma	
	Intensity ratio	n_e (10^{20} cm^{-3})	Intensity ratio	n_e (10^{20} cm^{-3})
$2 {}^1\text{P}_1/2 {}^3\text{P}_1$	2.7	1.0	3.3	1.1
$3 {}^1\text{P}_1/2 {}^3\text{P}_1$	0.71	1.0	0.68	1.0

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$2^1P_1/\text{Si}^{+11}$, as a function of electron temperature. Only those Si^{+11} lines whose energies are in this cluster are included in

Fig. 5-55. Calculated ratio of H- and He-like silicon resonance-line intensities vs electron temperature. The ratio is primarily a temperature diagnostic.

Fig. 5-56. Calculated intensity ratio of the He-like silicon resonance line and the dielectronic satellites of the Li-like silicon ion vs electron temperature.

calculating the intensity ratio. (The Si^{+11} lines are identified by Gabriel's letter designations.¹¹⁵) The L shell is near LTE, so electron-density effects are not nearly as important for this ratio as for those plotted in Fig. 5-55. The observed line-intensity ratios, from Table 5-3 (indicated by the arrows in Fig. 5-56), give electron temperatures of 0.4 and 0.53 keV for the CH and PbO plasmas, respectively. These values are lower than those implied by the H-like-to-He-like ratios, but they are similar in that the PbO plasma indicates a higher electron temperature than the CH plasma. As mentioned above, this is to be expected from laser-plasma scaling.¹¹⁶

The electron density was estimated using the ratio of the intensities of the resonance and intercombination lines ($2^1P_1/2^3P_1$) of the He-like species. This ratio has been used in a number of other applications.¹¹⁶ The calculated ratios are plotted in Fig. 5-57 vs electron density. Also plotted in Fig. 5-57 are curves for the intensity ratio $3^1P_1/2^3P_1$, which we observed for reasons explained in the next paragraph. The solid curves are the ratios calculated using only the intercombination-line intensity, while the dashed curves include the contribution from the Li-like satellites that are unresolved from the intercombination line. The correction for these unresolved lines is substantial at high electron densities, at which collisional quenching of the 2^3P_1 population reduces the intercombination-line intensity.

The observed line-intensity ratios, from Table 5-4, are indicated by arrows in Fig. 5-57. For each plasma, both ratios indicate the same electron density, confirming the assumption that the resonance lines are optically thin. The implied density is about 10^{20} cm^{-3} , which is a factor of 40 less than the critical density for 0.53- μm light. The steady-state calculations show that the Li-like intensities are not significant in this region, so contributions from Li-like silicon ions cannot explain such a low inferred electron density.

The line-intensity ratios thus imply electron temperatures between 0.4 and 1 keV and an electron density of $\sim 10^{20} \text{ cm}^{-3}$, so that the observed region of the plasma can not be characterized by a single electron temperature and density. Intensity ratios of lines from consecutive ionization stages indicate different electron temperatures. The

implied electron density is much less than critical density, where most of the emission is expected to originate.

The trends of the data are consistent with the axial and radial electron-temperature and density gradients calculated for planar-disk irradiations. The line-intensity ratios indicate that emission from lower ionization states will arise in colder, denser regions while that from higher ionization states will arise in hotter, less dense regions. The steep gradients may account for the low value of the inferred electron density: the intercombination line is more intense at lower densities, thereby weighting the average to the lower densities.

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Suprathernal X-ray Spectra from Shiva Disk Targets

It is important to understand suprathernal x-ray production in the laser-plasma interaction, because high-energy x rays and electrons preheat targets and thus affect target implosion.

Suprathernal electrons stopping in the target produce suprathernal x rays as bremsstrahlung. The spectra of those x rays that escape the target not only indicate the presence of suprathernal electrons but give their temperature and number density as well. During Shiva's lifetime we accumulated fragmentary data on suprathernal x rays from a number of disk targets using the filter-fluorescer¹²² and hyper-filter-fluorescer spectrometers¹²³ and filtered NaI(Tl) crystals.¹²² By grouping related data, we have studied the dependence of suprathernal x-ray spectra on the atomic number of the target and on laser intensity, energy, and pulse length. Because the disk targets have a simple geometry, we have also used the data to verify predictions of the LASNEX simulation code.

Experimental Data. Figures 5-58 through 5-62 show spectra of suprathernal x rays from various disk targets irradiated with 1.06- μm light at various laser energies, intensities, and pulse lengths. The solid lines

Fig. 5-57. Calculated intensity ratio of the He-like silicon resonance line to the intercombination line ($1s-2^3p$) vs. density.

are freehand fits to the data. Figure 5-58 shows that the yield of suprathernal x rays rises with laser intensity and target atomic number. Figures 5-59 and 5-60 show that short pulses yield far more suprathernal x rays than do long pulses, even though the temperatures are quite similar. Figure 5-61 shows that higher laser intensity almost offsets lower atomic number, so that the yields of suprathernal x rays from Zn and Au targets are nearly equal. Figure 5-62 shows that a combination of higher laser intensity and shorter pulse length more than offsets lower atomic number, so that a Zn target yields more suprathernals than an Au target.

Comparison of LASNEX Calculations with Data. We replaced the Bethe-Heitler bremsstrahlung cross sections in the LASNEX simulation code with a new, parameterized fit to the more detailed theoretical calculations of Feng and Pratt.¹²⁴ The calculated spectra are indicated by dashed lines in Figs. 5-58 through 5-62; T_{exp} and

Fig. 5-58. Supra-thermal x-ray spectra from Ni, Zn, and Ti disk targets: high intensity and atomic number give higher yields of suprathermals.

Fig. 5-59. Supra-thermal x-ray spectra from Au disk targets: shorter pulses yield more suprathermals.

T_{calc} are experimental and calculated supra-thermal x-ray temperatures, respectively.

The absolute intensities of the calculated spectra are arbitrary, because the intensity is controlled by the fraction of the energy that reaches the turning point that is resonantly absorbed. This is a free parameter in LASNEX, called the "percent of dumpall." We did not have absorption data for these shots from which we could determine this quantity, so we set it to 5% for all the calculations to give reasonable agreement with the observed spectra.

For short pulses, the temperatures obtained from the LASNEX calculations agree fairly well with the experimental temperatures. For long pulses, however (see Figs. 5-59 and 5-60), the calculated temperatures are much lower than the experimental values. This is because the only absorption mechanism (before the ray turning point) in these LASNEX calculations is inverse bremsstrahlung (collisional absorption), which does not produce many supra-thermal electrons. The collisionless absorption processes that absorb energy to produce a hot

Fig. 5-60. Suprathermal x-ray spectra from Au disk targets: larger pulses yield fewer suprathermals.

Fig. 5-61. Suprathermal x-ray spectra from Zn and Au disk targets: high intensity almost offsets effect of lower atomic number on suprathermal x-ray yield.

spectrum in the underdense plasma (such as Raman, $2\omega_{pe}$ and filamentation) are not included. Only resonance absorption at the ray turning point creates hot electrons in these LASNEX calculations.

In a long-pulse calculation, which is characterized by the formation of a long, underdense plasma and thus by a large path length for absorption by inverse bremsstrahlung, most of the light is absorbed by this mechanism at the thermal-electron temperature. The amount of light that reaches the ray turning point is quite small, and the intensity-dependent model for the hot temperature of resonance absorption gives a hot-electron spectrum that is not nearly hot enough. In reality, the collisionless processes that occur in the underdense plasma compete with classical absorption, and create much hotter electrons, as is clear from the observed spectra.

Figures 5-59 and 5-60 show that the LASNEX calculation was greatly improved by the incorporation of Brillouin¹²⁵ and Raman backscattering mechanisms. The Raman backscattering model essentially cal-

culates the fraction of light absorbed and scattered by the Raman process, using the analytic theory of Kruer¹²⁶ and OREMP (kinetic, electromagnetic, relativistic) simulations; then it puts the heated electron

Fig. 5-62. Supra-thermal x-ray spectra from Zn and Au disk targets: high intensity and short pulse more than offset lower atomic number.

Fig. 5-63. Steady-state thermal-electron temperature in the underdense plasma vs. microwave power, as measured in Prometheus III.

Fig. 5-64. Electron current to a planar probe in the underdense plasma vs. probe voltage. The non-Maxwellian line is a best fit to the data.

spectrum into the electron bins, assuming a T_{hve} corresponding to $(1/2)m_e v_p^2$, where v_p is the phase velocity of the electron plasma wave of the Raman process.¹²⁶ The $2\omega_{\text{pe}}$ absorption process¹²⁷ also heats electrons to high temperatures, but is not modeled here. The short- and long-pulse shots absorb

about 0.9 and 1.3% of the laser light, respectively, into Raman-heated electrons that radiate into the spectra shown in Figs. 5-59 and 5-60.

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Microwave Simulation of Laser-Plasma Interactions

We have extended our measurements¹²⁸ of microwave-plasma interactions to model laser-plasma interactions over the wide range of conditions that occur in laser driven pellets. Our work in 1982 was primarily concerned with modeling laser-plasma interactions in the long, underdense plasmas that occur in pellet designs driven by moderate laser power.

Hot-Electron Generation by the IAD and OTS instabilities. The ion-acoustic decay (IAD) and oscillating two-stream (OTS) instabilities can be excited on the steepened density profile near the critical surface or on the density shelf below the critical surface.¹²⁹ In the Prometheus I device at the University of California at Davis (UCD) we found that resonance absorption is the dominant absorption mechanism if the shelf is short enough. The Prometheus III device was designed so that the unperturbed density-gradient scale length could be much longer than in Prometheus I ($l \lesssim 5 \times 10^3 \lambda_{\text{De}}$), so that the underdense shelf could be much longer. The waveguide radius was increased to 30 cm so that the microwaves are cut off at a plasma density $n_{\text{cut}} \sim 0.9 n_c$. The underdense shelf in Prometheus III is now quite long ($l \sim 2 \times 10^3 \lambda_{\text{De}}$), and probe measurements indicate that strong high-frequency and ion-acoustic waves ($\delta n/n \sim 0.2$) are excited on the shelf. Supra-thermal electrons are heated as they travel through the shelf region. In addition, thermal electrons are strongly heated in the shelf region.

Thermal-Electron Heating and Heat-Flux Inhibition by Ion Turbulence.

Figure 5-63 shows measurements¹²⁹ of the thermal-electron temperature, T_e , from Prometheus III; T_e increases as $T_e \sim P_0^{1/2}$, where P_0 is the incident microwave power. The figure also shows the scaling, $I\lambda^2/T_{e0}$

to show how these experiments could be extrapolated to some laser conditions.

We have made extensive measurements¹³⁰ of the velocity-distribution function, $f(v)$, of the heated electrons. The thermal electrons are heated by scattering with the ion turbulence, and electron-electron collisions are weak, so $f(v)$ should not be Maxwellian: instead, theory¹³¹ gives $f \sim \exp[-(v/v_{th})^2]$. Figure 5-64 shows the current collected by a disk probe in the underdense plasma as a function of the probe voltage (the I - V characteristic). The data match the I - V prediction based on the non-Maxwellian $f(v)$, and not that for a Maxwellian distribution.

Hot-Electron Generation by the $2\omega_{pe}$ Instability. We have measured¹³² the properties of the two-plasmon decay ($2\omega_{pe}$) instability in Prometheus III. The measured threshold is very low: $(v_{os}/v_e)_1 \approx 0.1$. Our theoretical prediction¹³³ for p-polarized standing electromagnetic waves is $(v_{os}/v_e)_1 \sim 0.7/(k_0 l)^{1/2}$, where l is the plasma-density scale length. S-polarized, traveling waves result in a value of v_{os}/v_e about 5 times greater. Using the measured value ($l \approx 80$ cm), we obtain a predicted threshold of $(v_{os}/v_e)_1 \approx 0.14$, close to the measured value.

We also have measured¹³² the growth of the instability, the emission of microwaves with a frequency near $3\omega/2$, modification of the plasma density profile near the quarter-critical surface, microwave absorption, and electron heating. The ponderomotive force of the high-frequency waves strongly modifies the plasma-density profile near the quarter-critical surface. The sheath step density profile is formed in steady state. Hot electrons are observed for $P_{in} < 10P_1$, where P_1 is the threshold microwave power. The hot-electron temperature and density increase slowly with power, i.e., $T_{H1} \propto P_1^{1/2}$ and $n_{H1} \propto P_1^{1/4}$. The thermal electrons are only weakly heated. Figure 5-65 shows the measured T_{H1}/T_e , which fits the theory for resonance absorption¹²⁹ (solid line) very well: the figure also shows that $T_{H1} \sim I^{1/2}$, which agrees with particle-simulation calculations for laser parameters.¹³⁴

Stimulated Brillouin Scattering. We have made extensive measurements¹³⁵ of strong SBS in the UCD Cerberus device. In this experiment, standing microwaves are

Fig. 5-65. Ratio of hot-electron and thermal-electron temperatures T_{H1}/T_e vs $I\lambda^2/T_e^{3/2}$, which is predicted¹³³ to yield a straight line for resonance absorption.

Fig. 5-66. Ion-wave growth rate γ_i vs v_{os}/v_e from the Cerberus device. Solid line is the theoretical prediction.

set up in the region of microwave-plasma interaction. A semiparallel beam of microwaves ($\omega/2\pi = 2.9$ GHz) is produced by an $\epsilon = 1.5$ dielectric lens and reflected by a screen reflector (the plasma ($v \lesssim 1/2u$), which has a gentle axial density gradient, is created by electron bombardment of argon gas at 2.5×10^{-4} Torr. The electron temperature, T_e , is typically ~ 2 eV during the discharge and ~ 0.5 eV during the afterglow. The ion temperature, T_i , measured by an energy analyzer, is ~ 0.5 eV. The maximum microwave power incident on the plasma is 0.3 MW, and the minimum- ϵ (FWHM) beam diameter is 20 cm, so $(v_{os}/v_e)_1 \approx 0.1$.

The ponderomotive force of the standing microwaves drives a resonant ion wave ($k \sim 2k_0$) that acts as a large amplitude noise source for SBS. The ion-wave grows for microwave powers above a clearly defined threshold that agrees with the theoretical¹³⁶ threshold $(v_{os}/v_e)_1 \approx 0.15$.

As shown in Fig. 5-66, the ion-wave growth rate agrees with recent theoretical predictions.¹³⁷ The ion-wave grows to large amplitude ($\delta n/n \gtrsim 1$), and the incident microwaves are scattered within a few wavelengths ($l_s \gtrsim 2\lambda_0$). However, the process does not reach steady state. The scattering region moves toward the source at high

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velocity ($V \lesssim 5C_s$). Ion waves grow in the plasma in front of the scattering region (closer to the microwave source). The microwaves are excluded from the plasma behind the scattering region, resulting in ion-wave collapse and strong ion heating.

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A model is shown of a large-aperture segmented-disk amplifier for the Zeus laser. In the real amplifier, there would be 25 to 100 segments, each of which would be 30×60 cm.

Section 6

Zeus Laser Project

Zeus Laser Project

Contents

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Zeus Laser Project

Introduction

J. T. Hunt

It has become apparent over the past several years that advanced technologies that require igniters or drivers with high energy density may be explored in the laboratory using high-power, focused, short-wavelength lasers. The advanced technologies that can be explored include

- Testing and development of high-gain inertial-confinement fusion (ICF) targets, for simulation of nuclear-weapons effects and commercial power generation.
- Testing of nuclear-weapons physics.
- Pumping of x-ray lasers.

The Zeus Laser Project was organized during 1982 to address the development of a laser for these purposes. The ultimate goal of the Zeus Project is to design a laser-driven test facility that will meet the following technical objectives¹

- 10 MJ of energy, rated at 10 ns.
- 1000 TW of power, rated at 10 ns.
- Wavelength $\approx 0.5 \mu\text{m}$.
- Pulse width from 30 ps to 100 ns.
- Flexible focusing.
- Thermonuclear yield of 1 ton.
- Cost of \$250 million in 1983 dollars, excluding conventional project costs such as S & E, contingency costs, and inflation factors.

Several known laser media are capable of generating high-power, short-wavelength laser pulses; such laser media include Nd:glass, excimers (such as KrF), iodine, and several advanced solid-state media. Based on demonstrated technology, however, existing laser systems cannot be simply scaled to Zeus performance levels for the specified cost. Thus, to achieve our technical objectives using known laser media will require substantial innovation.

In Phase I of the Zeus Project, we will self-consistently evaluate known laser technologies, new approaches that may evolve from them, or alternative solutions that appear independently. For each reasonable architecture, Phase I evaluation will consist of identifying and resolving

- The physics risks and method of risk reduction.
- The component costs and performance goals necessary to meet the Zeus target cost of \$250 million.
- The fraction of the total system cost for each subsystem.
- The relative overall system costs for each of the various architectures.
- The relevant test-bed requirements.

During 1982, we carried out the first three steps of Phase I for Nd:glass systems. Following a completed assessment of the Nd:glass systems, we will turn to studies of the remaining candidate systems.

In Phase II of the Zeus Project, we will make more detailed studies of the likely costs and performance parameters for the various system components of candidate architectures. These studies will include analytical assessments, laboratory experiments, and vendor surveys. In Phase III, we will complete a detailed system design and cost study for the most attractive approach, as identified in Phases I and II.

In this section, we first present our cost-scaling scheme and compare several candidate laser architectures. We then present the results of our amplifier analysis carried out during the year. The balance of the section is devoted to discussion of various innovative architectures, frequency-conversion schemes, and optics technology.

Zeus Cost-Scaling and Candidate Amplifier Architectures

Our joint goals for Zeus are to meet the technical objectives stated in the Introduction to this section, while limiting the project cost to \$250 million. In studying how to meet both goals, the Zeus project team has addressed the following cost-related issues

- The magnitude of component-cost reductions necessary to complete Zeus at the constrained cost.
- The most reasonable and fair distribution of cost reductions across the various subsystems.

In this article we first present our preliminary answers to the cost-related issues. We next present the Zeus baseline energy requirements and then examine three candidate amplifier systems in light of those requirements.

We began our cost-scaling effort by determining Nova component costs, to obtain a point of departure for assessing required cost reductions for Zeus. Nova hardware costs were derived from Nova budget

Table 6-1. Comparison of costs for Nova, 100× Nova output power, and Zeus.

	System costs in 1983 \$ millions		
	Nova	100× Nova	Zeus
Laser-chain hardware			
Amplifiers	20	2025	152
Extraction optics, beam transport, and frequency conversion	21	693	51
Subtotal	41	2718	203
Other system components			
Output-sensor package	5		0
Chain-invariant hardware	29		17
Laser laboratory building	26		10
Laser office building	10		0
Target systems	12		10
Design labor	28		0
Fabrication labor	10		10
S & E	5		0
Contingencies	10		0
Subtotal	135		47
Total system costs	\$176 million		\$250 million

data and component-cost charts. From these sources, we determined that the cost of amplifiers scales approximately as stored energy and that the remainder of the laser system (beam transport, energy extraction and frequency conversion) scales approximately as aperture area.

As an intermediate step, we used the Nova cost data and the approximate scaling factors to estimate the cost of a similar system but with an output energy 100× that of Nova (as we require for Zeus). Assuming constant Nova component costs and performance levels, our intermediate step resulted in a chain hardware cost of \$2718 million. Of this amount, \$2025 million would be apportioned for amplifiers; \$693 million would be apportioned for the extraction, beam transport, and frequency conversion systems. Using this rough breakdown and our scaling factors, we then made a preliminary breakdown of required Zeus subsystem costs. The cost goals—and at this stage they are only goals—are shown in Table 6-1. As an example of how to meet these goals, Fig. 6-1 shows the cost reduction scheme we envision for the Zeus laser chain hardware.

The Baseline Zeus Laser System

The “reasonable” component performances required for the baseline Zeus laser system are based on physics limitations and our intuition as to what is achievable. We now describe the baseline system and then examine three candidate amplifier architectures for use in the system.

Figure 6-2 shows the energy-flow diagram for the baseline Zeus system. The energy on target is given: 10 MJ of blue (0.35- μ m) light. If we assume that the frequency-conversion optics are 70% efficient (the same as Nova) and that we have 95% transmission from the converter to the target (roughly the same as Nova), we find that the laser driver must deliver 15 MJ to the frequency converter.

Next we relate the energy extracted from the amplifier to the energy it stores. The energy relationship is

$$E_{\text{ext}} = \eta_{\text{ex}} E_{\text{st}} \quad (6-1)$$

where $\eta_{\text{ex}} = \eta_{\text{ext}}/\eta_{\text{st}}$. Here, the overall extraction efficiency is 0.5, the extraction efficiency for a beam having a filling factor of 1, η_{ex} , is 0.7; and the beam filling factor, η_{st} , is 0.7. We have chosen the value of 0.7 for η_{ex} because it is equal to the result of computer simulations for the basic amplifier gains and transmission of a 0.7 fill factor on the one hand, and is based on what we have actually achieved with past laser systems. The last parameter required to complete the calculation is the optical storage efficiency. We have chosen a value of 3.75%, which represents an improvement over Nova by a factor of about 3. This optical storage efficiency results in an 800 MJ capacitor bank and an overall system efficiency of 1.25%.

expanding-ring amplifiers, which are shown schematically in Figs. 6-3 and 6-4, respectively. For comparison, we have also calculated the performance of a third architecture, a single pass through a linear chain (such as on Nova). In all cases, the amplifier module consists of an amplifier and a spatial filter. The spatial-filter pinhole is also the location of an exploding foil Q switch; at the present time, this component does not affect the beam in any way. The amplifier itself consists of flashlamps arranged along two walls facing an array of rectangular glass amplifier segments. The size of the amplifier is nominally 3 m × 3 m × 50 m; it contains 10 disks. The spatial filter consists of a single pinhole and two segmented lenses that also serve

Fig. 6-1. Cost-scaling scheme for the Zeus laser-chain hardware, based on Table 6-1.

Laser-chain hardware cost for 100× Nova		\$2718 million
Amplifiers	Extraction, beam export, and frequency conversion	
	\$693 million	
\$2025 million		
1.5× reduction; damage limits from 0.5 to 1 μm are already available		
1.4× reduction; architecture		\$462 million
\$1446 million		
1.4× reduction; architecture		\$330 million
10× reduction; forced		\$152 million
6× reduction; forced		\$51 million
Zeus laser-chain hardware cost = \$203 million		

Candidate Amplifier Performance

We are exploring various laser architectures to meet our baseline Zeus performance goals. We have compared laser-chain performance for three potential architectures under similar rudimentary conditions using the SNOBAL design code (no attempt has been made to optimize these designs as yet). Each system consists of the same amplifier, with a different optical configuration; the stored energy is the same throughout. In addition, each system is driven to the point where the incremental B integral between pinholes on the last pass is 3.0. This enables us to isolate the differences in output as due to differences in extraction efficiency and the effect of losses external to the cavity.

Two new candidate architectures are based on the regenerative-multipass and

Fig. 6-2. Energy flow from electrical banks to target for Zeus baseline system.

as vacuum barriers (design of this component may be a significant technological problem). The pulsed power is placed in close proximity to the amplifier.

For the regenerative-multipass system, the amplified pulse is turned back along its original path several times. An exploding-foil Q-switch holds off the gain during the pumping cycle. The input pulse is applied through a leaky (3% transmission) mirror at one end of the cavity and is reflected from a dichroic mirror at the other end. While circulating in the cavity, the beam passes through a nonlinear crystal array held at a minimum conversion efficiency by an electric field applied to the crystal; on the last pass, the field is removed from the crystal array to allow optimal frequency conver-

sion. The frequency-converted beam passes through the dichroic mirror, through a spatial filter, and proceeds towards the target.

In the expanding-ring architecture, the external optical system consists of four mirrors that circulate an annular beam sequentially outward into new areas of an active medium on successive passes; the spatial filter provides the necessary magnification. If the magnification ratio is properly chosen to match the loaded gain, the system can be operated at constant fluence across the output aperture. The output is an annular beam that passes through an additional spatial filter and then through the crystal array for conversion to 0.35 μm .

In all three configurations, the target is protected from prepulses by an exploding foil opening switch placed at the focus of a spatial filter. Present calculations assume that the switch component does not modify laser performance in any way.

Table 6.2 gives the nominal specifications for the glass and external components of potential Zeus architectures. The glass parameters are those of a phosphate glass with a cross section somewhat lower than for the glasses used on Nova; the choice of these parameters is not sensitive to changes in the optical system. The storage density is typical of the smaller Nova amplifiers.

Table 6-2. Nominal specifications for amplifier modules for Zeus candidate laser architectures.

Glass	
Cross section	$3 \times 10^{-20} \text{ cm}^2$
Saturation fluence	5.1 J/cm^2
Saturation characteristics	Conservative phosphate $[\text{Es}^{II} - 0.81(\text{Es})]$
Nonlinear refractive index	1.25 ESU
Storage density	0.35 J/cm^3
Gain	5.6%/cm
Loss	$0.1\%/ \text{cm} + 0.3\%/\text{surface}$
Disk thickness	4.2 cm
Number of disks	10
Small-signal gain	16.9
External components	
Loss—transmitting mirrors	1%/surface
Loss—KD*P crystals	5%/pass
Frequency-conversion efficiency	70% for 3 ω (blue) light

Fig. 6-3. Schematic of the regenerative-multipass architecture considered for the Zeus baseline system.

while the disk thickness is the same as in the larger Nova amplifiers; 10 such disks give a total small signal gain of 16.9. The transmitting optics are assumed to be antireflective coated. The values for loss in the glass are typical for presently available glasses, as are the losses in the external components. Losses in the wavelength-conversion material in the non-converting mode are assumed to be 5% per pass. The overall conversion efficiency is assumed to be 70%.

Table 6-3 details the performance of the three candidate architectures, given the parameters of 1.1 to 6.2. For those parameters, none of the three systems reaches the limiting extraction efficiency or nominal damage limit for a 1 μ m laser beam, all are either all limited or damage limited at 0.35 μ m. Note that the output fluence for the expanding ring and linear chains is higher than for the regenerative system and that the linear amplifier has the highest output and the highest extraction efficiency of all three (the driver energy required to achieve this result is, of course, also high).

The efficiency for the five pass expanding ring system is lower than for the linear chain because of the losses in the turning

mirrors and a nonoptimal distribution of flux across the aperture. We chose a 3 $\times$ area magnification, which is optimum for extraction-limited performance; for performance limited by the B integral (long pulse operation), however, the flux on the earlier passes is less than optimum.

The difference in efficiency and output between the regenerative-multipass and expanding-ring systems is not significant. In fact, an optimized expanding ring design should have slightly higher efficiency than the regenerative multipass system because of the larger single cavity loss attributed to the multipass switch.

Table 6-3. Performance characteristics of the regenerative-multipass, expanding-ring, and single-pass linear laser architectures.

Controlled parameters			
Aperture	3 m $\times$ 3 m		
Fill factor	0.7		
Stored energy	2520 kJ		
Pulse length	10 ns		
	Regenerative multipass	Five-pass expanding ring	Single-pass linear
Energy out (0.35 μ m)	741 J	687 J	1002 J
Energy in (1 μ m)	2 J	0.1 J	400 J
Gain factor	1.0	0.67	1.0
Efficiency	0.29	0.27	0.32
Fluence (1 μ m)	18	24	24
Fluence (0.35 μ m)	13	17	17

3x3 optical component array

Fig. 6-4. Schematic of the five-pass expanding-ring architecture considered for the Zeus baseline system.

Summary

During 1982, we developed the baseline technological requirements for the Zeus laser system and began to assess the cost reductions and performance improvements in present technology that will be necessary to achieve those requirements at a constrained cost of \$250 million. We have apportioned such cost reductions and performance improvements over the various Zeus cost centers, with an emphasis on laser-chain hardware, which commands 80% of the total Zeus budget.

We have also undertaken computational comparisons of various amplifier architectures with an eye toward maximizing cost-effectiveness and total output fluence. We will now turn to more detailed assessment, and perhaps consolidation, of our preliminary cost scaling. We will likewise continue to examine various laser architectures for their applicability to the Zeus baseline requirements.

Authors: J. T. Hunt and W. L. Gagnon

Amplifier Analysis: Fluorescence Amplification in Different Disk Geometries

When laser amplifier disks are pumped, the natural rate of inversion loss due to fluorescence is increased by amplification in the disk. This fluorescence amplification reduces the available gain and, thus, reduces the system performance. Prior calculation of this effect¹ handled only the case of a simple elliptical disk; today's larger amplifiers use split elliptical disks, however, and tomorrow's amplifiers will probably use rectangular slabs. During 1982, we calculated the effect of fluorescence amplification in elliptical disks, half-elliptical disks, and rectangular slabs of various aspect ratios. We also devised a simple method of using the results of these calculations in existing computer programs for comparison with experimental data.

Fluorescence amplification starts when spontaneously emitted photons are emitted isotropically from each volume element in the disk, in proportion to the elements' inversion density, and having the spectral shape of the fluorescence line. Each photon travels through the disk and is exponentially amplified with a gain coefficient that depends on its wavelength. Photons striking the flat faces of the disk will experience one of two results. If the photon strikes at an angle below the angle of total internal reflection (TIR), it will pass out of the disk and be lost; if it strikes a flat disk face above the TIR angle, however, the photon will bounce back and forth between the faces until it strikes the edge of the disk, where it will be absorbed by an edge coating.

The processes of emission, amplification, reflection, and absorption were simulated with a computer program that followed photons with a single gain coefficient as they traveled through a disk. Photons at initially random positions, traveled in random directions. Runs were made with 10,000 to 20,000 photons to reduce errors in the results to a few percent. Results were calculated for across-disk gains ranging from ~ 1 to the maximum values attained at line center in the highest-gain disks actually used.

Table 6-4 shows sample results from a series of runs for increasing gains. Fluorescence either exits from the faces or bounces back and forth between the faces until it reaches the absorbing edge. Note that the face power barely increases with increasing gain across the disk since the path to the face is short on the average. In contrast, the bouncing path to the edge may be quite long, and the power reaching the edge can be thousands of times the emitted power. It is this large amplification that increases the loss rate, reduces the gain, and drops system performance.

Because of the great range of gain that results from these calculations, it is convenient to present the results in terms of the "effective length" of a photon path. The effective length is the path length that would give the same results as the actual calculation if all photons had that path. Figure 6-5 presents the effective length as a function of the gain coefficient times the short axis of

the laser disk, for four disk geometries of interest for fusion laser amplifiers. As expected, split elliptical disks show the least fluorescence amplification, and rectangular slabs show the most. For all disk shapes, the effective length increases as the gain coefficient increases because long paths contribute more heavily at high gain.

The results in Fig. 6-5 allow us to approximate the effect of geometry changes using programs (such as GAINPK) that have the results for an unsplit elliptical disk already built in. We use a modified disk long axis in the code. By matching the equivalent lengths, at a given gain coefficient, times the long axis of ~ 1.5 , we derive the following:

Shape	Long axis
Ellipse	1.00
Half ellipse	0.71
Square	0.80
2:1 Rectangle	1.12

To test our method, we used GAINPK to simulate the full-disk and half-masked gain measurements on a phosphate disk; Fig. 6-6 shows the data points for the two cases. GAINPK was fit to the full-disk (lower) points, and then the long axis was set to 0.71 times its actual value, and a prediction of the half-masked results was made. As can be seen, the prediction is in excellent agreement with the actual data.

Author: J. B. Trenholme

Reflex-Ring and New Multipass System Architectures

High-power Nd:glass laser systems built at LLNL over the past two decades have utilized the master-oscillator, power-amplifier (MOPA) architecture. This evolution has produced increasingly powerful lasers, from Janus through Cyclops, Argus, Shiva, Novette, and Nova. Nova will have 10 amplifier chains, each over 100 m long and reaching a clear-aperture size of 74 cm.

Laser systems of the next generation will be required to produce 10 to 100 times the power and energy of Nova: a simple

scale-up of the linear MOPA architecture to $100\times$ Nova power may not be cost-effective, however. So far, our search for

Gain	ASE power exiting the glass		
	Edges	Faces	Total
0.2	0.8810	0.2449	1.1259
0.4	1.0324	0.2516	1.2840
0.6	1.2347	0.2586	1.4933
0.8	1.4916	0.2627	1.7543
1.0	1.8582	0.2638	2.1220
1.2	2.2968	0.2676	2.5644
1.4	2.9114	0.2705	3.1819
1.6	3.8426	0.2711	4.1130
1.8	5.1892	0.2714	5.4606
2.0	6.7640	0.2826	7.0466
2.2	9.3418	0.2774	9.6192
2.4	13.5950	0.2819	13.8769
2.6	19.1881	0.2855	19.4786
2.8	26.3711	0.2960	26.6671
3.0	38.4508	0.2913	38.7421
3.2	58.1801	0.3026	58.4827
3.4	93.6017	0.3058	93.9075
3.6	142.2252	0.3085	142.5337
3.8	231.9242	0.3141	232.2383
4.0	347.5272	0.3056	347.8328

Table 6-4. Comparison of the power of amplified spontaneous emission (ASE) exiting the edges and faces of a rectangular slab of laser glass.

Fig. 6-5. Fluorescence amplification in four disk geometries.

Fig. 6-6. Comparison of simulation with full-disk and split-disk gain measurements for a phosphate disk.

Fig. 6-7. (a) Reflex-ring and (b) MOPA laser configurations.

(a) Reflex-ring architecture

(b) MOPA architecture

alternatives has focused mostly on regenerative or multipass systems, each with its unique set of technical issues. We have recently considered a new architecture called the "reflex ring," based on the unstable ring resonator.

The conventional isofluence linear MOPA is constructed of a number of stages, each consisting of amplifiers with gain and spatial filter having a magnification $M > 1$. A small diameter laser beam enters this long laser chain with an input fluence of a few J cm^{-2} ; the beam is then amplified, filtered, and expanded so that the output fluence equals the input fluence, i.e., $M = G$. This periodic structure suggests that we might employ a ring configuration in which the output is fed back to the input of one large aperture stage.

Beam propagation through the MOPA and reflex ring configurations is compared in Fig. 6-7. The reflex ring is driven through a hole in the input mirror. The small diameter beam passes through the amplifier and then through a spatial filter, which magnifies it and redirects the beam to the input via a series of mirrors. As the beam reflects off the input mirror for the next pass, the central part of the beam is lost through the hole; the beam is now annular (in the case of a round beam) as it propagates around the ring. On the last pass, the annular beam increases in size to the point that it passes around the output mirror residing in the magnified shadow of the input hole.

The beam spatial profiles for three passes through a reflex ring are illustrated in Fig. 6-8. In a typical isofluence reflex ring the central 25% of the aperture acts as a preamplifier for driving the remaining 75% in a single pass. Note that the reflex ring shares many features of existing laser systems, such as spatial filtering and beam expansion, allowing the extraction efficiency and the quality of reflex ring lasers to be the same as for the linear MOPA.

Development of the reflex ring as a viable fusion driver requires addressing several issues:

- **Backward propagation.** As the beam propagates, the first pass gain in a high-gain reflex ring.
- **Isolating the target from the beam.** This has to be achieved through frequency conversion and beam dumping.

- Target back-reflection isolation won't have to be achieved through frequency conversion.
- Lasing off of the spatial filter pinhole must be avoided (there are no Faraday rotators).
- Pinhole closure must be avoided.

Despite these issues, the reflex-ring architecture illustrates the simplicity and cost-effectiveness we must achieve in future laser drive.

A serious technical issue for multipass architecture is the requirement that the amplifier must be rapidly dumped between cycles. During 1982, we conceived a fast dump aperture switch for this purpose; the design employs noncollinear phase-matched second harmonic generation.

The most common harmonic generation technique is a three-wave process in which two photons, one from each of two pump waves, are added together to produce one photon at the second harmonic. Since both pump waves are required to generate the third, we envisioned using one pump wave to "switch out" the other. These three waves can be physically distinguished by polarization (Type II collinear phase-matching), propagation direction (non-collinear phase-matching), or wavelength (mixing).

A schematic of such a fast switch is shown in Fig. 6-9. Two 1ω pulses are injected into separate multipass ring lasers having unequal cavity lengths. By adjusting the injection timing and difference between the two cavity lengths, we can arrange for the two 1ω pump pulses to arrive at the frequency-conversion crystal simultaneously; after a small number of passes through the gain medium, the pulses will be switched out of the two cavities by the frequency-conversion process, with the 2ω pulse propagating along a direction determined by conservation of momentum, i.e.,

$$\vec{k}_1 + \vec{k}_2 = \vec{k}_3$$

The noncollinear phase matching switch features:

- Fast, passive large-aperture switching.
- Target ASE suppression.
- Isolation from target back-reflection.
- No requirement for a 1ω beam dump.

The technical issues to be solved before we can employ this large-aperture switch include:

- 1ω amplitude mismatch ($\vec{k}_1, \vec{k}_2$ mix ratio).

▲ Fig. 6-8. Spatial profiles of a triple-pass reflex ring.

Fig. 6-9. Non-collinear phase-matching switch for fast dumping of two multipass amplifier cavities.

- Vignetting in the frequency-conversion crystal.
- Crystal absorption if we use KDP for the frequency converter.
- Cavity loss if we use "leaky-mirror" injection.
- Effects of unconverted 1ω light.
- Oscillation suppression.

Despite technical issues, a large-aperture switch design appears worthy of experimental characterization. If the design becomes a viable option, it will, in turn, support the development of a new generation of multipass architectures.

Author: M. A. Summers

Cassegrainian Disk Amplifier

Multipass architectures are appealing for large laser systems since they provide high stage gain (E_{out}/E_{in}) in a single amplifier

Fig. 6-10. Schematic of the triple-pass Cassegrainian amplifier architecture (the Brewster's-angle disks in the gain volume are not shown).

Cassegrainian Disk Amplifier

stage with a minimum number of components. A simple multipass architecture widely used in gas lasers may be of some interest for glass lasers as well, because it requires little space and can be implemented without developing new components (such as a large-aperture optical switch). The architecture used in gas lasers is a triple-pass configuration designed around a Cassegrainian reflective beam-expanding telescope. In this article we examine the major advantages and disadvantages of the Cassegrainian amplifier.

As Fig. 6-10 shows, a Cassegrainian amplifier is not a true triple-pass amplifier but a double-pass power amplifier with a small preamplifier pass sharing the same gain volume. Model calculations show that the greatest change in performance and cost between single-pass and multipass amplifiers is for an increase of from one to two passes, though the theoretically optimum extraction may require more passes. In the latter situation, the cost of additional components required for an elaborate multipass or regenerative amplifier design makes the low component count of the Cassegrainian architecture very attractive.

The Cassegrainian geometry, with its built-in preamplifier, compares favorably with a full-aperture double-pass amplifier having a separate preamplifier. First, a double-pass amplifier requires that the beam direction change at each amplifier. It is, of course, possible to lay out laser chains that do this, as in the active-mirror amplifier geometries.³ These geometries, however, have complex beam paths, while the Cassegrainian geometry keeps the beam going in the same direction. Second, it is much easier to stack Cassegrainian optical trains side by side to form a large amplifier array, and the small convex mirror in the

Cassegrainian beam picks up very little of the output aperture. It should be much cheaper to increase the output aperture of a final Cassegrainian amplifier by a few percent (to provide for an integral preamplifier) than it would be to build a separate, smaller preamplifier and its additional components and controls.

The main drawback of the Cassegrainian amplifier is that its spherically diverging beam picks up aberrations when it passes through a disk at Brewster's angle. This difficulty has prevented any serious consideration of the Cassegrainian geometry for disk lasers, despite its popularity in other systems. Previous analysis¹ shows that the most important aberration is astigmatism. Consider a beam with a half angle α passing through a Brewster's-angle disk having a thickness t and a nonlinear refractive index of 1.5 to 1.6. This beam picks up N waves of astigmatism from the center to the edge of the beam, where

$$N = 0.45 \frac{t\alpha^2}{4\lambda} \quad (1)$$

As a specific example, suppose we have 60 cm of glass in a cavity (fourteen 4.3-cm disks, for example) and can accept five waves of astigmatism in the output beam. Then α must be less than 9 mrad, or, in other words, the length of the amplifier must be greater than $60D$, where D is the diameter of the output beam. The axial space occupied by fourteen 4.3-cm disks at Brewster's angle is about $224t$ —so that a bit less than two-thirds of the Cassegrainian cavity will be open space. We prefer to make more economical use of space, but even with this limitation the Cassegrainian amplifier uses space about as well as other designs. If it is possible to correct a larger

aberration, then the cavity can become spherical, though for short cavities the deviation from Brewster's angle must be considered. A rapidly expanding beam also has effects on extraction and the B integral.

The astigmatism in the Cassegrainian disk amplifier is correctable in several simple ways. The easiest method is to fabricate one or several of the disks in the shape of a thin wedge; it can be shown that the aberrations introduced by an expanding beam are easily corrected in this way.⁵ We could also figure the small convex mirror in the telescope to compensate for the distortion introduced by the disks. Finally, if a full-aperture half-wave plate were available to rotate the beam polarization, rotating half of the disks 90° around the optical axis would also correct the aberration.

The maximum small-signal gain in a Cassegrainian geometry will be limited by ASE or parasitic oscillations, and by nonlinear effects (the B integral) at short pulse lengths. The maximum tolerable small-signal gain between isolators in a practical design is probably between about e^9 and e^7 , which implies a single-pass small-signal gain in the amplifier of e^3 to e^4 (20 to 50). For 4.3-cm phosphate-glass disks pumped to the same energy storage per unit volume as Nova 31-cm amplifiers, the amplifier will have between 10 and 14 disks.

A simple Frantz-Nodvik analysis shows that a 14-disk Cassegrainian amplifier should have a telescope magnification of about 5 and give a stage gain of about 2000 at 75% extraction. The B-integral increment for a 5 ns pulse in this amplifier would be about 3.2, which is not outrageous, but 5 ns is clearly a short pulse for this design. A 10-disk Cassegrainian amplifier requires a magnification of 4.5, and would have a stage gain of roughly 200 at 75% extraction, with a somewhat lower B-integral increment. The fraction extracted from both of these amplifiers could be increased by from 5 to 10% at small cost by driving the amplifier somewhat harder, but, as with all amplifiers, B rises much faster than extraction, making this strategy useful only for very long output pulses.

In summary, the Cassegrainian amplifier has the following positive features

- High stage gain and extraction.
- Very simple geometry.
- Very low component count.

- No new technology or components.

- Easy stacking into a large array.

The disadvantages of the Cassegrainian architecture are as follows

- Brewster-angle components introduce aberrations in the laser beam, though these effects can be compensated for in straightforward fashion.

- The optics of the simple Cassegrainian are not image-forming and, thus, cannot be used to form part of an image relay, though they will transmit an image formed by optics earlier in the laser chain. Propagation of the laser beam through Cassegrainian amplifier stages will therefore be less controlled than in fully relayed chains.

- The expanding Cassegrainian beam complicates the analysis of nonlinear propagation and extraction; we will have to modify our SNOBAL and MALAPROP codes to allow multidimensional propagation and extraction before we can have full confidence in the accuracy of present amplifier models.

- The small convex mirrors at the end of the preamplifier pass in the Cassegrainian amplifier block about 4 to 5% of the output-aperture of the amplifier array.

- If the output beam is apodized, the edges of the beam around the convex mirror must be shaped as well as the outer periphery of the beam.

The simplicity and low component count of the Cassegrainian architecture, and its likely low cost, suggest that we should analyze carefully how serious these disadvantages will be in practice.

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Electrooptic Harmonic-Conversion-Switched Multipass System

We have developed a new approach to generating short-wavelength light with a multipass Nd:glass amplifier; the system is shown schematically in Fig. 6-11. A laser pulse at the fundamental Nd wavelength, 1.06 μm (1 ω), is injected into a multipass

cavity through a partially transmitting mirror (as in Fig. 6-11) or reflected from a destructible foil near the focus of the spatial filter. After the injected 1ω pulse has been amplified by repetitive passes through the Nd:glass amplifier, it is switched out of the cavity by electrooptic harmonic conversion. A pulsed electric field is used to control the phase-matching conditions for harmonic conversion in an electrooptic crystal (such as KDP, KD*P, or CD*A). The amplified 1ω light is converted to second-harmonic (2ω) light (in this example) and leaves the cavity by transmission through a dichroic end mirror (as in Fig. 6-11) or reflection from a polarizer (for Type I frequency-doubling, the 1ω and 2ω beams are polarized orthogonally). The dichroic mirror reflects 1ω light and transmits 2ω light.

This optical-pulse switchout technique simultaneously combines the functions of switchout and harmonic conversion. Hon,⁶ for example, has used this effect to compensate for detuning due to self-heating of CD*A during average-power harmonic conversion with a 1ω laser; an applied electric field was used to control phase-matching and provide efficient (57%) harmonic conversion.

To understand the operation of the switching technique described above, consider electrooptic harmonic conversion with CD*A. An isomorph of KDP, CD*A has a $\bar{4}2m$ point-group symmetry; the optimal harmonic-conversion configuration is shown in Fig. 6-12. The normally degenerate crystal axes, x_1 and x_2 , are oriented at 45° with respect to the propagation wave vector, $k_{1\omega}$, of the 1ω beam; as shown in

Fig. 6-12, this beam is polarized in the x_1x_2 plane. The 2ω beam produced by the non-linear interaction of the 1ω beam with the crystal is polarized along the x_3 axis.

If we apply a short-pulsed (10 to 100 ns), electric field, E , along the x_3 direction, x_1 and x_2 are rotated in their plane by 45° to x_1' and x_2' . The x_2' axis is then parallel to the electric field, $E_{1\omega}$, of the 1ω beam. The index of refraction seen by the 1ω beam as it passes through the CD*A crystal is

$$n_{x_1,1\omega} = n_{o,1\omega} + \left(\frac{n_{e,1\omega}^2}{2} \right) E_{1\omega} / E_0 \quad (2)$$

where $n_{o,1\omega}$ is the ordinary refractive index at 1ω , and $E_{1\omega}$ is the appropriate electrooptic coefficient. Equation (2) is valid provided the duration of the applied electric pulse is short compared to the typical mechanical response time (~ 10 μ s) of the bulk crystal. Otherwise we must account for the indirect contribution by the crystal strain, induced by either the E field or the reverse of the piezoelectric elastooptic effect. In the configuration of Fig. 6-12, the applied electric field, E , has no effect on the 2ω beam, which is polarized with the electric field, $E_{2\omega}$, along the x_3 axis. The index of refraction seen by the 2ω beam as it passes through the CD*A crystal is $n_{x_3,2\omega}$ —the extraordinary refractive index at 2ω .

Efficient frequency doubling requires that the phase of the 1ω and 2ω electromagnetic waves be matched. In the harmonic conversion region where the 1ω beam is not depleted significantly, the 2ω power generated is⁷

Fig. 6-11. Electrooptic harmonic-conversion-switched multipass system.

$$P_{2\omega, \text{max}} = \left| \frac{\sin^2 \Delta k \ell / 2}{\Delta k \ell / 2} \right|, \quad (3)$$

where ℓ is the crystal length, and

$$\Delta k = k_2 - 2k_1 \quad (4)$$

Using $k_1 = 2\pi/\lambda_1$ and $k_2 = 2\pi/\lambda_2$, then

$$\Delta k = \frac{2\pi}{\lambda_1} \left(n_{\text{ord},1} - n_{\text{ord},2} \right) \quad (5)$$

From Eq. (5), maximum 2ω power is generated when the 1ω and 2ω beams are phase-matched, i.e., $\Delta k = 0$. According to Eq. (3), as Δk is increased from zero, the 2ω power decreases, reaching zero when $\Delta k \ell / 2 = \pi$.

Using these results, we may now describe the operation of an electrooptic harmonic-conversion switch using CD*A. Suppose that no electric field is applied ($E = 0$), and the CD*A crystal is temperature- or angle-tuned to $\Delta k \ell / 2 = \pi$. According to Eq. (5), this requires that

$$n_{\text{ord},1} = n_{\text{ord},2} = \frac{\lambda_{2\omega}}{\ell} \quad (6)$$

Given this condition (and the 1ω beam polarization and CD*A crystal orientation shown in Fig. 6-12), each time the 1ω beam passes through the CD*A crystal, no 2ω light is generated. After the 1ω beam has

been amplified to the desired level, however, a pulsed electric field, E , is applied to the crystal to produce phase-matching ($\Delta k = 0$). Now when the 1ω beam passes through the crystal, it is efficiently converted to 2ω ; the 2ω pulse is polarized perpendicular to the 1ω beam and may exit the multipass cavity through a dichroic mirror or by reflection from a polarization splitter.

The electric field required to turn on the harmonic conversion may be estimated using Eqs. (5) and (6), and is

$$E \sim \frac{\lambda_{1\omega}}{n_{\text{ord},1}^2 r_{63} \ell} \quad (7)$$

Suppose, for example, that the CD*A crystal is 2 cm long, $r_{63} = 14.1 \times 10^{-12}$ cm/V, and $n_{\text{ord},1} \approx 1.5$; then $E \approx 11.1$ kV/cm.

At the high intensities that yield efficient harmonic conversion, depletion of the 1ω beam takes place and Eq. (3) is no longer valid; the value of Δk at which no harmonic conversion occurs becomes intensity-dependent and shifts to smaller values than required by Eq. (3). This shift in Δk may be troublesome, since it leads to prepulses and increases the optical loss of the 1ω beam as it propagates inside the multipass cavity. This problem may be overcome by applying a field initially to prevent phase-matching, and then removing the field after the multipass extraction has occurred. In addition, the switch configuration with a transverse applied field (Fig. 6-12) is useful for small apertures but not readily scalable

Fig. 6-12. Crystal-axis orientation and beam polarization for transverse-field-excited electrooptic harmonic conversion using CD*A.

to the large apertures required for Zeus and reactor-driver laser systems. The problem of scaling for large apertures may be overcome by using other electrooptic crystals and by applying the electric field along the direction of propagation. For example, for KDP with Type I phase-matching, the longitudinal (along the propagation direction) voltage requirement is ~ 45 kV.

Author: R. A. Haas

Rimfire Power Amplifier

In 1982 we proposed a design for a rimfire power amplifier, dubbed "Rimfire." The design shown in Fig. 6-13 consists of a glass disk (2 to 3 cm diameter) with the input beam profile as indicated, and the output beam profile as indicated. The 360° pump beam is directed at the

Fig. 6-13. (a) Single-taper and (b) double-taper Rimfire amplifier.

refocused by a convex output/turning mirror that rims the glass disk. The disk is assembled from pie-shaped glass segments ~ 10 cm thick. The segments of the single-taper Rimfire amplifier are of uniform thickness, while double-taper segments expand in thickness with increasing radius (see Fig. 6-13(b)). The glass segments are joined through the two largest faces; the beam of the pump does not propagate through these large pump faces; they do not require optical finishing. Each segment requires little preparation on its two interior faces, but the output disk segments would be set to a precise alignment. The excitation considerations are that the pulse is injected into all the segments through a single cone reflector located in the center of the glass disk; the resulting design is, in essence, an axicon with a conical beam. The transverse area of each segment increases linearly with radius in the single-taper geometry of Fig. 6-13(a) so that quasi-isofluence is maintained naturally if the glass is operated in the heavily saturated mode of a power amplifier. An additional control on the fluence as a function of radius is obtained by tapering the glass segments in the other cross sectional direction (see Fig. 6-13(b)) and diverging the beam in that dimension by figuring the cone reflector in the double-taper geometry; the beam transverse area increases quadratically with radius.

The expanding-area feature of the large Rimfire disk is an advantage, compared to conventional rod operation. Quasi-isofluent propagation means the nonlinear refraction is proportional to γL , instead of $\gamma L(e^{dI} - 1)/\alpha$. For both geometries, the rim mirror is naturally off-loaded (without lenses) with respect to damage fluence, by the radially expanding area. B-integral considerations would affect the allowed radial path length.

We could achieve uniform gain over the cross sectional area of the Rimfire amplifier by smoothly increasing the Nd doping concentration with distance from the pumped faces. Alternatively, we could implement a zig-zag^{2,11} propagation path within the glass slabs, as shown in Fig. 6-14. A zig-zag beam path averages the gain over the glass volume and has the additional advantage of double passing the disk, thereby increasing extraction efficiency and cancelling primary pump induced optical distortions. If we employ the latter option, we incur the cost of optically finishing the large pump faces. In addition, a saturable absorber may be required to prevent runaway ASE in the thick disks.

A stacked Rimfire amplifier is depicted in cross section in Fig. 6-15 (for comparison, both the single-taper and double taper options are shown simultaneously). Surrounding the lamps with laser glass should

Fig. 6-14. Rimfire single-taper disk segments figured for (a) simple and (b) zig-zag propagation; zig-zag beam path would achieve uniform gain.

Fig. 6-15. Cross-sectional view of a stacked Rimfire amplifier, showing both single-taper and double-taper configurations.

contribute to high overall efficiency; large pump faces over which a coolant could be flowed would be useful for operation at a high repetition rate.

Author: W. L. Smith

Frequency Conversion of Zeus

There are several approaches to frequency conversion of Zeus. The first is simply the use of present technology—the same crystal design, array fabrication, and construction techniques—without any changes; as frequency conversion is in its infancy, we expect considerable cost reductions in available technology in the next few years. The second option is an upgrade of present technology, with performance and cost scaled from our present experience; this approach would include new nonlinear materials, for example, and streamlined array-fabrication and construction techniques. The third approach is a completely new frequency-conversion technology, one whose performance and cost cannot be reliably predicted from present experience.

In this article, we look at the general requirements of performance and cost for

Zeus frequency conversion, the applicability of present technology to those requirements, and the known possibilities for technology upgrades and new technologies. At present, however, we lack the technical information and other data with which to make a complete assessment of the costs and performance of these latter approaches.

General Requirements

As discussed in Zeus Cost Scaling and Candidate Amplifier Architectures, earlier in this section, the target cost for Zeus frequency conversion is \$24 million. At this cost, the system must deliver a fourable output at 3ω (0.35 μm) or 4ω (0.26 μm) with a high (~ 70 to 80%) conversion efficiency, over a fairly large (0.03 to $3 \text{ GW}/\text{cm}^2$) range of intensities. The system must be compatible with the laser beam profile (including variations in intensity, phase, and polarization), operate efficiently in a wide dynamic range, and occupy an appropriate location in the laser system. Finally, the frequency conversion system must not only match the properties of the laser system itself, but its effect on laser performance must be minimal.

If the frequency conversion system is external to the final amplifier cavity, for example, it will have essentially no impact on the laser performance; an external architecture also allows some flexibility in the 1ω parameters. If the frequency conversion system is located inside the final amplifier cavity, however, it may have a sizeable impact on the amplifier's performance, by introducing birefringence into the cavity, for example. If the cavity is operating in a multipass mode, the conversion efficiency on all but the 1ω pass must be small; any light converted to 3ω or 4ω in prior passes appears as a loss in the cavity and detracts from the laser performance.

Nova Frequency-Conversion Technology

For present technology, high conversion efficiency to 3ω requires that the detuning budget in the laser beam be less than 150 grad, the polarization error be less than 3° , and the intensity range for the KDP conversion

crystals be between 1 and 3 GW/cm². Regarding these parameters, the frequency-conversion system used for Nova and Novette is acceptable. Nova technology also meets the requirement of 70% conversion efficiency in the range of 1 to 3 GW/cm².

In the low-intensity range (from 10 to 100 MW/cm²), however, conversion efficiency is below 10%. In addition, interesting pulse shaping is required: approximately one third of the energy is in the tail of the pulse.

For conversion efficiency from 10% to 55% at 3 μ m, at 20–40 GW/cm², the gain is from 80% to around 55 to 60% depending on the details of the spectral and pulse shape. Therefore, Nova technology is able to meet the nominal Zeus specifications over dynamic range.

We are also concerned about damage and absorption issues with present frequency-conversion technology. Depending on the peak-to-average ratio in the more intense Zeus laser beam, the damage fluence of KDP may be from 1 to 2 times too low. Furthermore, the absorption fraction of KDP is probably 3 times too high for use inside the Zeus cavity and 2 times too large for use outside the cavity. Though the damage and absorption issues are not as serious as the issue of low conversion efficiency due to pulse shaping, they do indicate the necessity of seeking materials other than KDP and KD*P.

Nova Technology Upgrades and New Technologies

Frequency-conversion efficiency at low intensity is proportional to the product of the nonlinear coupling coefficient and the crystal length. Increasing either parameter by a factor of 5 to 10 enables us to meet the design goal of 70% conversion efficiency. Since material costs tend to scale directly with crystal volume, we prefer the option of increasing the nonlinear coupling coefficient, rather than using a material having the same nonlinear coefficient as KDP and a much lower angular sensitivity. In searching for new materials, however, we must consider both options for meeting the requirement of 70% conversion efficiency.

In the area of totally new technology, we have explored frequency conversion in atomic vapors, such as lithium and rubid-

ium, using xenon gas for the phase matching; these schemes have typically achieved conversion efficiencies of less than 10%, however. Another technique involves doping a nonlinear crystal with an absorbing ion or some other very narrow line absorber, using anomalous dispersion for the phase matching. Preliminary work suggests that for dopings of 10 to 1000 ppm, a strongly allowed transition (given an oscillator strength between 0.1 and 1) might meet our technical requirements.

Cost

For Nova technology, the cost of KDP and finishing is 66% of the total cost of the array. This cost is loosely proportional to the volume of KDP, is inversely proportional to the growth rate, and is also inversely proportional to the size of the boule. The major cost is for fabrication labor. This suggests that the 6 $\times$ cost reduction required for Zeus frequency conversion (see "Zeus Cost-Scaling and Candidate Amplifier Architectures" earlier in this section) could be achieved by

- Making the size of the KDP boule much larger than the size of the crystal segment.
- Artificially reducing transverse losses to suppress Raman scattering.
- Increasing the number of parallel operations during crystal growth.
- Automating the growth process.

Along with these options, we expect to derive a considerable cost improvement factor from the learning curve and from volume scaling.

Summary

It is clear that neither KDP nor KD*P will meet the Zeus frequency-conversion specification of 70% conversion efficiency over a wide dynamic range. We do believe, however, that presently available technology holds the promise of a 6 $\times$ cost-reduction factor, which qualifies it for consideration for Zeus on the basis of cost alone. Other technologies, involving KDP isomorphs or other nonlinear materials, cannot yet be assessed technically because we lack the necessary information. With continuing research in the areas of both current and

entirely new technologies, we expect to perform more complete cost-benefit analyses of the options for frequency conversion of Zeus.

Author: D. Eimerl

Large-Aperture Optics Technology

Nova's optical components, excluding the laser glass, contributed about 10% of the cost of the total project. For Zeus, we intend to maintain or reduce this fractional cost while improving performance. Given the total Zeus target cost of \$250 million (see "Zeus Cost-Scaling and Candidate Amplifier Architectures" earlier in this section), Zeus' optical components excluding frequency conversion must cost no more than about \$25 million. The total area of finished optical surfaces for Zeus is expected to be about 1000 m²; this gives us a target cost of \$2.50/cm² of optical surface, compared to the Nova procurement cost of \$10/cm² (see Table 6-5). Meeting this challenge will require substantial innovation.

Some improvement in energy per unit aperture will be obtained through an increase in damage fluence and a reduction in the effects of nonlinear self-focusing at longer pulse lengths; both improvements will allow operation at higher extraction efficiency. The major difference between the optical systems of Zeus and Nova, however,

will be Zeus' larger total optical aperture facing the target. Our experience with amplifier construction suggests that the cost per unit of output energy decreases substantially as amplifier size increases. Construction of large optical surfaces using arrays of smaller segments in passive prealigned structures has been suggested. The cost of this approach will have to be compared to the cost of approaches that rely on individual, sophisticated control systems, which increase the cost of the laser system. At some size, the cost of optical components is expected to increase with aperture. The particular aperture where the balance between various approaches occurs, is dependent on specifications and technology, and is as yet unknown.

Zeus Optical-Surface Specifications

The specifications for specific optical surfaces are directly related to the aperture area and the desired beam quality or spot size. A nominal spot size for a single beam of Zeus is set by the requirement that the intensity in the spot be between 100 and 1000 TW/cm²; the requirement gives a nominal spot diameter of 5 mm for a single beam—significantly larger than for Nova. The f/No. of Zeus' focusing optics will have to be kept moderately large, to minimize aberrations and the cost of constructing aspheric surfaces; larger f/Nos., in turn, require higher pointing accuracy and a better optical figure to produce a specific spot size.

We expect that spatial filters will be required to remove small-scale amplitude modulation in the laser beam; the angle for maximum growth is ~ 1 mrad, and the pin holes in the spatial filter must intercept radiation at angles of ~ 100 μ rad to be effective. The slope error for the surface of the optical elements for the spatial filters must be sufficiently low to maintain the major portion of the beam well within this angular tolerance. For multipass systems that pass the beam through the same pin hole in the spatial filter several times, the slope errors must be even lower, since the angular deviations on each pass tend to add. In the same sense, the cumulative aberrations from all other optical surfaces in the system must be consistent with the spot sizes and angular pointing accuracy at the

Table 6-5. Comparison of Zeus and Nova optical-component goals and specifications.

Parameter	Zeus		Nova
	Range	Nominal	Nominal
System goals			
Energy (MJ)	10	10	0.1
Wavelength (μ m)	0.20–1.0	0.35	0.35
Pulse length (ns)	0.1–100	10	3
Cost (\$ millions)	25	25	15
Optical specifications			
Beam aperture (cm)	100–500	300	71
Number of beams	2–40	10	10
Fluence (J/cm ²)	10–20	15	3
Optical surfaces per beam	5–15	10	15
Total surface area (m ²)	500–2000	1000	150
Spot size (mm)	1–10	5	0.1
f/No.—focusing	5–10	10	4
f/No.—spatial filter	10–40	20	20
Angular tolerance (μ rad)	5–100	15	10
Cost (\$/cm ²)	1–8	2.50	10

target and spatial filters. Thus, the figure and pointing tolerances for the optical elements on Zeus will not be radically different from those for Nova.

Cost-Reduction Options

To achieve a significant saving in the cost of the Zeus optical system, the total number of optical components in each beamline must be kept as low as possible. A minimum of six large-area surfaces, exclusive of the laser disks, appears to be required: one surface for focusing onto the target, two for frequency conversion element, two for a spatial filter, and one for coupling the driver to the final amplifier. Additional components may be required for vacuum interfaces, debris shields, diagnostics, interstage isolation, multipass options, and

for the driver itself. A total of 10 surfaces, which would represent a significant improvement over prior designs, seems to be a realistic goal for the Zeus system.

Various technologies may reduce the cost of generating large optical surfaces.¹² Our approaches include

- Extension of conventional optical finishing techniques, using computer control.
- Use of thinner substrates with integrated support structures.
- Stressed mirror polishing.
- Segmentation and active electronic control.
- Lightweight cast-pyrex substrates.
- Replication from masters.
- Diamond machining.

We are actively pursuing such cost-reducing approaches in the optical-telescope and laser communities; none represents a radical departure from current technology.

Fig. 6-16. Large-aperture membrane mirrors.

One possibility that does break with tradition is the concept of membrane, or "pillow," mirrors and windows.¹³

Membrane Windows and Mirrors

The function of an optical component is to maintain a surface that is fixed with respect to some reference point such as a target. For conventional optical components, the stability of the surface is maintained largely by the substrate material itself. As the aperture of the component increases, the thickness of material must increase to resist deformation by its own weight. Thus, the weight and cost of large optical components tend to increase faster than the surface area. Thinner materials can be used, but the surface figure must then be maintained by external forces.

If a differential pressure is applied uniformly across a thin plate, it deforms into a curved surface. Such a surface might be used as an optical element, as shown in

Fig. 6-16; as a transmitting element, the surface could be used as a vacuum window. A pair of such windows could be used to generate a weak, variable positive lens. If one surface is a reflector, then the two configurations shown in Fig. 6-16 could be used to provide converging or diverging mirrors. In Fig. 6-16(a), the center space is above ambient pressure; a second transmitting surface might be required to contain the pressure in this configuration. In Fig. 6-16(b), the center space is at reduced pressure.

The primary issue is whether a surface with the desired focal length and optical figure can be generated with realistic materials. To explore this issue, we have analyzed the geometry shown in Fig. 6-17. A thin circular sheet of uniform thickness, h , is supported at radius, a . The deflection of the surface is referenced to a cylindrical coordinate system (r, θ) at the apex of the curve. The components of a vector describing the deformation of a point initially at radius, r , are $u(r, \theta)$, which correspond to local deformations in the directions θ . Cylindrical symmetry is assumed. If pressure, p , is applied normal to the inner surface of the plate. External forces, N , and bending moments, M , per unit length may be applied at the edge of the plate.

The surface shape is maintained by a balance of radial and tangential forces and bending moments in each surface element. For sufficiently thin plates, the contribution of the bending moments can be neglected. In this case, the pressure applied over the surface is transmitted to the edge by radial forces, N , acting at angle, α , where $\tan \alpha = dr/dh$ is the local slope of the surface. Thus,

$$p \cdot \pi \cdot r^2 = N \cdot 2 \cdot \pi \cdot a \cdot \sin \alpha = 0 \quad (8)$$

For small deflections, $\sin \alpha = \tan \alpha$, $\alpha = dr/dh$, and

$$\frac{dr}{dh} = \frac{p \cdot r}{2 \cdot \sigma \cdot h} \quad (9)$$

where $\sigma = N/h$ is the stress at radius, r . If the stress, σ , in the foil is independent of radius, then Eq. (8) can be integrated to give a parabolic surface.

$$\frac{dr}{a} = \left(\frac{dr}{r} \right) \cdot \left(\frac{r}{a} \right)^2 \quad (10)$$

Fig. 6-17. Geometry used for the calculation of membrane deformations.

where $w_0/a = (p \cdot a)/(4 \cdot \sigma \cdot h)$ is the relative displacement at the edge of the plate.

Unfortunately, the stress in the plate is not uniform. The stress arises from two sources: external forces applied prior to deflection, and reaction forces from the deformation of the plate. The approximation that the stress is constant is good only if the stress from external forces, or prestress, is much larger than the stress from the reaction forces. The radial stress, σ_r , in the membrane satisfies the equation¹⁴

$$\frac{d}{dr} \left(\frac{1}{r} \frac{dw}{dr} \right) + \left(\frac{1}{2} \right) \cdot \left(\frac{dw}{dr} \right)^2 = 0. \quad (11)$$

The stress and deflection are coupled by the nonlinear terms $N \cdot dw/dr$ in Eq. (8) and $(dw/dr)^2$ in Eq. (11).

A solution of Eqs. (8) and (11), accurate to second order in the shape of the surface, gives

$$\frac{w}{a} = \frac{w_0}{a} \left(\frac{r}{a} \right)^2 \left[1 - \frac{3}{2} \left(\frac{r}{a} \right)^2 \right]. \quad (12)$$

where

$$\beta = \frac{1}{\left[\frac{3}{1-\nu} \cdot \frac{p}{E} - 4 \left(\frac{\sigma_0}{E} \right) \left(\frac{w_0}{a} \right)^2 \right]}$$

ν is Poisson's ratio, and E is Young's modulus. The total stress in the membrane is $\sigma = \sigma_r + \sigma_\theta$, where σ_0 is the prestress from the external forces. The relative deflection, w_0/a , at the edge of the foil is given by

$$\left(\frac{\sigma_0}{E} \right) \cdot \left(\frac{w_0}{a} \right) + (0.96) \cdot \left(\frac{w_0}{a} \right)^3 = \frac{p \cdot a}{4 \cdot E \cdot h}, \quad (13)$$

where ν has been taken to be 0.3. The first term is the contribution from the external forces, and the second term is the contribution from the reaction forces. There are two asymptotic limits. The first, which corresponds to large prestress and negligible reaction forces, is given by Eq. (11). The second, which corresponds to large deflections with no prestress, gives

$$\frac{w_0}{a} = \left[\left(\frac{1}{3.8} \right) \cdot p \cdot \left(\frac{a}{E \cdot h} \right) \right]^{1/3}. \quad (14)$$

In this case, the deflection varies as the cube root of the pressure, rather than linearly as for the case of constant stress.

Figure 6-18 shows the $f/\text{No.} = 1/8(a/w_0)$ of the surface plotted against the parameter $X = (p \cdot a)/(E \cdot h)$. The lower curve is the $f/\text{No.}$ in the absence of prestress; the upper curve is the $f/\text{No.}$ in the absence of reaction forces prestressed to the value,

Fig. 6-18. Variation of $f/\text{No.}$ as a function of normalized pressure.

$\sigma_v/T = 3 \times 10^{-5}$. For Zeus, the $f/\text{No.}$ s of interest range from $f/5$ to $f/40$; this range can be reached at $X = 1 \times 10^{-4}$ with prestress and 1×10^{-5} without prestress.

Figure 6-18 also shows typical operating points for steel and aluminum plates. The pressure differential is assumed to be 10 psi, and the ratio of radius to thickness is 300. Increased pressure moves the operating point to the right on the curves, and increased thickness moves them to the left. For a plate 3 m in diameter, the thickness is 5 mm; this is a thin plate, though not what is normally considered to be a membrane. The use of still thinner plates is possible, provided that a reduced pressure differential is applied. For a foil 50 μm thick, the pressure differential required is only 0.1 psi (7 Torr).

For prestressed plates, the sensitivity of the focal length to pressure fluctuations is linear. To maintain a focal length to an accuracy of 1%, for example, the pressure must be controlled to within 1%. If one side

Fig. 6-19. Relative spot size as a function of $f/\text{No.}$.

of the plate or foil is at ambient pressure, the use of small pressure differentials requires a correspondingly higher control of the absolute pressure. Accurate control of such small differential pressures may be a problem. In the nonlinear regime, the sensitivity is reduced only by a factor of 2.

The ability to reach an appropriate $f/\text{No.}$ is only one criterion for a viable membrane mirror. Sufficiently low aberrations are of equal or greater importance. In the linear regime, the stress distribution is nonuniform, and the shape of the surface of the plate deviates from that of a parabola, as given by Eq. (12). A convenient measure of this aberration is the deviation w of a ray from the edge of the beam as it passes through a plane at the paraxial focus. If the surface were a parabola, this deviation would be zero. The ratio w/a , where a is the size of the beam, is given by the ratio of the first two coefficients in the power-series expansion of the slope of the surface. From Eq. (12), the relative spot size or demagnification of the beam defined in this manner is

$$\frac{w}{a} = \left[\frac{1}{(3-r)/(1-r)} \right] + 256 \frac{\sigma_v l}{(f/\text{No.})^2} \quad (15)$$

where r is Poisson's ratio, again taken to be 0.3.

This spot size is plotted in Fig. 6-19 as a function of $f/\text{No.}$ for the same parameters as in Fig. 6-18; the Zeus operating range is also shown. A 1-cm-diam spot with a 3-m-diam mirror requires a demagnification of 300 ($w/a = 0.03$). A 1-m mirror could produce the same spot with less demagnification or could tolerate more aberrations. In the nonlinear regime, the spot size is large and independent of the $f/\text{No.}$ As prestress is added to the plate, the spot size decreases rapidly with $f/\text{No.}$ However, for the value of prestress chosen ($\sigma/E = 3 \times 10^{-3}$), the spot size is only marginal at high $f/\text{No.}$ s.

In principle, the spherical aberration of a spot size can be reduced to a more acceptable value by application of an arbitrarily large prestress. There is a limit to the stress that can be applied to any material, however, failure usually occurs in tension. For brittle materials such as glass, the failure is usually catastrophic. For ductile materials

References

(70 μm Hg). Neither of these options appears feasible at this time, but further study is required before we reject the use of plastics.

Conclusions

In our study of membrane windows and mirrors to date, we have reached the following conclusions

- Curved surfaces with focal lengths and $f/\text{Nos.}$ in the appropriate range for Zeus can be generated by applying a pressure differential to large circular plates.
- The required plate thickness will be a few millimeters unless very low pressure differentials can be used.
- The deflection of the plates is nonlinear with pressure.
- Reaction stresses in the plate tend to produce a surface with high aberration.

- Prestressing the plate with external forces prior to application of pressure could produce a surface figure that might be marginally acceptable for Zeus.

It is too soon to conclude whether or not membrane mirrors will be useful for fusion laser applications. The surface figure produced is not as suitable for optical use as we would prefer, though there may be solutions to this problem. Subdivision of the mirror into segments may help, and thermal stresses may be useful in modifying the shape. Further, it has been suggested that a radial pressure differential could be established between two thin plates to modify the surface figure. If these or other developments should prove effective, membrane mirrors may lead to lower-cost optical components for Zeus.

Author: L. D. Fleasance

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Section 7

Laser Research and Development

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Laser Research and Development

Introduction

W. H. Lowdermilk

The major objectives of the Laser Research and Development program are two-fold. First, the Nova and Zeus Projects are supported principally by development of improved materials and components for Nd:glass lasers. Second, research on advanced laser systems will have eventual applications to ICF commercial-power generation.

During 1982, we addressed the following broad topics

- Solid-state amplifiers and related technology.
- Optical-materials development, including laser-induced damage studies.
- Nonlinear optical measurements.
- Advanced-lasers research.

We focused on methods of improving the efficiency of Nd:glass amplifiers; we measured the overall radiant efficiency and relative Nd-pumping efficiency of Xe flashlamps of various diameters; we advanced the theory of arc formation and plasma emission; we measured improvements in efficiency obtained by using fluorescent dyes or phosphor materials to convert those portions of the flashlamp spectrum not absorbed by Nd:glass into the absorbed spectrum; and we examined options for reducing the cost of pulsed power to excite the flashlamps, concluding that significant cost reduction can be achieved.

We also made sensitive measurements of aberrations that can be imposed on the laser beam as it passes through the amplifier. Aberrations result from vibrations or temperature variations caused by flashlamp pumping. Understanding and eliminating sources of aberrations is particularly important for multipass amplifier systems.

Our major effort in optical-materials development and damage testing was to extend the application of damage-resistant, graded-index antireflection (AR) surfaces to the fused-silica optics required for laser operation at $0.35\text{ }\mu\text{m}$. Work sponsored by LLNL is proceeding at Westinghouse R&D Center on a graded-index AR coating, which is deposited from solution and is based on sol-gel technology. We developed the method of producing graded-index AR surfaces on fused silica and other oxide glasses by etching nuclear tracks.

Measurements of damage thresholds were made on a wide variety of coatings for operation at both single and multiple wavelengths. We measured damage thresholds with Nd:glass laser pulses at the fundamental wavelength of $1.06\text{ }\mu\text{m}$ (1ω), at the second harmonic of $0.53\text{ }\mu\text{m}$ (2ω), and at the third harmonic of $0.35\text{ }\mu\text{m}$ (3ω), and with a KrF laser at $0.25\text{ }\mu\text{m}$. The principal objectives of our measurements were to select for Nova the most damage-resistant materials and coatings and to provide data on the pulse-width and wavelength dependence of damage thresholds.

We studied damage to KDP and breakdown in the index-matching fluid used in the KDP arrays for Novette and Nova. We used the outcome of the latter study to set the requirements for filtering the index-matching fluid and cleaning the array cells.

As part of the optical-materials research program funded by the Office of Basic Energy Sciences, we studied saturation properties of Nd-doped glass and measured two-photon-solarization losses in borosilicate optical glass. Nonlinear optical measurements were made to define materials requirements and fluence levels at which Novette and Nova can be operated. Our measurements included two-photon absorption, solarization, and the nonlinear refractive index.

Solid-State Amplifier Development

Our continuing research on advanced lasers evaluated the technology and scaling properties of rare-gas halide (RGH) pulse-compressor systems. We completed the pulsed-power system for the Rapier B amplifier, which is a prototype for larger RGH amplifiers that require modularized pulsed-power systems and large gain lengths. We also advanced our understanding of the physics of KrF and Raman amplifiers.

We investigated inner-shell vibronic transitions in NpF_6 as a candidate for a gas-phase storage laser pumped by an efficient long-pulse Nd laser. Our experiments demonstrated that energy-storage densities up to 200 J/l can be maintained for ~ 1 ms in this system. Our energy density is one of the highest ever achieved in a gas medium.

Finally, we participated with M Division of LLNL in preparation for an experiment with a long-wavelength, high-gain, high-extraction-efficiency free-electron laser that will use the ETA accelerator. The experiment is now being assembled and will operate in early 1983.

Solid-State Amplifier Development

Introduction

During 1982, we measured the gain performance of Nova components and began aggressive programs directed toward obtaining a detailed understanding of amplifier pumping. Our ultimate goal is to integrate our flashlamp and amplifier studies to develop a comprehensive amplifier model that will allow us to design the most cost-effective amplifiers for Zeus.

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46-cm Amplifier Tests

We measured gain as a function of input energy of a 46-cm production amplifier with phosphate glass to be used in Novette. Conventional cw gain measurements were made with the 1.05- μm YAG line, which is

extremely close to the operating wavelength of 1 ω . To reject unwanted background fluorescence, we employed spatial filtering (30-cm focal length, 0.2-cm-diam pinhole) of the 1.0-cm-diam probe beam. Gain measurements were made at a position horizontally displaced 2.5 cm from the amplifier centerline and just outside the split region between the disks. (The horizontal displacement was in the direction of the vee formed in the center of the amplifier by the two disk pairs.) The electrical circuit consisted of 16 circuits, each composed of 10 capacitors, giving a total of 145 μF capacitance and 450 μH inductance. Lamp breakdown occurred for voltages ≥ 14 kV without using a reflector bump.

Figure 7-1 shows the measured gain as a function of input bank energy using all 16 circuits and simply varying the charge voltage. The gain value of 1.76 ± 0.02 at 600 kJ is below the design range of from 1.85 to 1.93. We reassessed the modeling previously done with the code GAINPEAK¹ to predict the gain performance. The GAINPEAK predictions were overly optimistic in the choice of current power. In addition, the predictions incorrectly applied geometric-coupling efficiencies and current powers (obtained for amplifiers with 1.5-cm-bore flashlamps) to the present amplifier, which has 2.0-cm-bore flashlamps. These factors explain most of the difference between the expected and measured gains.

We also measured the reduction in gain caused by operating with varying numbers of disconnected lamp circuits. Such operation might be envisioned in the case of a

Fig. 7-1. Small-signal gain vs bank energy for a 46-cm production-type amplifier.

few failed lamps. Furthermore, the variation in gain with number of circuits may provide a clue in explaining the strong gain rollover of our amplifier. Gain results were obtained with 1, 2, 4, and 8 circuits (out of 16) disconnected at 20 kV and with 2, 4, and 8 circuits disconnected at 12 kV.

The curve in Fig. 7-2 shows the measured gain coefficient vs bank energy, independent of whether the number of circuits or the bank voltage was varied. The gain curve appears to be constant for the two different ways of varying input energy. Our result suggests that the blue-shift of the lamps with input is not the major cause of gain rollover. Otherwise, we would expect the gain coefficient at a given voltage to follow a linear curve, intersecting the origin, as the number of circuits is varied. One explanation that is consistent with the observed behavior is that amplified spontaneous emission (ASE) and parasitics are clamping the gain. However, the maximum gain-diameter product for our amplifier is only about 2.6, which is reasonably small. A final possibility is that the rollover results from the decline in cavity Q ; hence the decline in coupling to the Nd:glass caused by either increasing the number of lamps or increasing the lamp loading, both of which increase the absorption of pump light by the lamps.

As a possible method to increase the gain of the 46-cm amplifier, we investigated the effect of reducing flashlamp pulse length, which can increase the gain by reducing the loss of excited Nd ions during pumping. At constant input energy, the benefit is somewhat off-set by the well-known reduction in Nd-pumping efficiency that occurs for increased lamp-current densities. Thus, the net effect is somewhat difficult to predict.

By paralleling two 450- μ H inductors in each of 16 circuits, we easily achieved a circuit inductance of 225 μ H. Given our capacitance of 145 μ F per circuit, the pulse width (3 τ LC) was reduced from 766 to 542 μ s. Unfortunately, at the 20-kV operating point, the change in pulse width increased the gain coefficient only by $2 \pm 1\%$. As discussed in "Flashlamp Pumping of Solid-State Lasers," later in this section, the change is much smaller than that observed for open-lamp pumping of Nd:glass.

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Fig. 7-2. Small-signal gain coefficient vs bank energy for the 46-cm amplifier.

Fig. 7-3. Reflector shapes used in 34-cm amplifier experiments.

Disk-Amplifier Performance as a Function of Reflector Shape

We assembled a simple 34-cm-aperture box amplifier to study the effect of reflector shape on amplifier gain performance. Although many previous tests have been made,² a careful comparison of gain with different reflector shapes for otherwise constant conditions has not been made. We were particularly interested in whether Winston reflectors³ provide an advantage [see Fig. 7-3(c)]. Winston reflectors direct the flashlamp emission so that no directly emitted light is reflected back into the lamp. However, such reflectors have a possible disadvantage in that all incident rays are routed through the lamp and may be absorbed.

We constructed a box amplifier with two 34-cm fluorophosphate disks (3.6-cm-thick LG812 glass, Nd doping = 1.6×10^{20} cm⁻³). Our amplifier was mechanically very similar to the rectangular-disk amplifier (RDA) and used seven Ce-doped quartz-envelope flashlamps on each side, which were excited in pairs by our standard

Fig. 7-4. Gain coefficient vs bank energy for an amplifier with three different reflector shapes.

Fig. 7-5. Transmission spectra of an ED-2 Nd:glass sample and five separate bandpass filters used to generate the curves shown in Fig. 7-6. ▼

electrical circuit (see "Flashlamp Characterization," later in this section). The lamps were spaced 6 cm apart for the three reflector geometries shown in Fig. 7-3. Moreover, the reflectors were designed so that the lamps on a given side did not directly view one another. All reflectors were fabricated from copper and plated with silver of comparable quality. To change reflector shapes, we unbolted one reflector and installed another. Thus, we made accurate comparative measurements over a relatively short time.

Measurements of centerline gain were made at 1.05 μm , with the amplifier terminated at each end by flat, silver-plated reflectors. The measured peak-gain coefficients vs bank energy are plotted in Fig. 7-4. The results are essentially indistin-

guishable for the three reflector shapes. We obtained similar results for the corner reflector painted with white, reflective Kodak 6080 paint. Our results suggest that reflector shape is not a major variable, at least for box amplifiers of this type having large lamp-packing fractions.

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Flashlamp Pumping of Solid-State Lasers

Introduction

Flashlamp plasma is viewed¹ as being in local thermodynamic equilibrium at a temperature near 10 000 K. The temperature increases weakly with input power. Although blackbody emission at this temperature typically peaks in the visible region. The peak occurs because flashlamps are gray bodies; hence, the photon field is not in thermal equilibrium with particle motion.

A fundamental phenomenon of flashlamps is the blue shift in their emission that occurs with increased power loading. As flashlamp loading is increased, emissivity of the plasma increases most dramatically at short wavelengths,⁵ causing the output to blue shift. Because the primary Nd-pump bands are at midvisible and near-infrared wavelengths, the blue shift tends to decrease the pumping efficiency of Nd ions. However, the net decline in pumping is a relatively small effect at reasonable lamp loadings. Our results for open-lamp pumping of Nd:glass cannot explain the large decline in pumping efficiency of amplifiers, which was discussed earlier in this section.

To illustrate the effect of the blue shift on Nd pumping, we used a special set of bandpass filters and a Nd-fluorescence detector. Transmission characteristics of each are shown in Fig. 7-5. The five filters were designed to have a constant transmission over each of the five main absorption bands of Nd:glass. We used the Nd detector to measure the peak Nd fluorescence produced by our standard 1.5-cm-bore

clear-fused-quartz flashlamp (for details, see "Flashlamp Characterization," later in this section).

By comparing the ratio of the peak Nd signals observed with each of the filters in front of the detector to the signal with no filter, we were able to determine the percentage of Nd pumping that occurs via each band, as shown in Fig. 7-6. The percentage

of pumping by the infrared bands generally declines with input to the lamp, while pumping via the 0.53- and 0.58- μm bands generally increases. The shift in pumping to the shorter-wavelength bands allows the overall pumping efficiency to remain relatively constant.

To quantify the spectral output of our lamp, we used a selection of long-pass

Fig. 7-6. Nd-pumping fraction produced by various absorption bands as a function of flashlamp loading.

Fig. 7-7. Transmission of filters used to isolate various components of flashlamp emission.

filters and a spectrally flat pyroelectric detector. By comparing the optical energy transmitted through each of the filters to the energy measured with no filter, we were able to determine the percentage of lamp output that was at a wavelength longer than the cutoff for each filter. Filter transmissions are shown in Fig. 7-7, and the transmitted energy fractions are shown in Fig. 7-8.

The blue shift of the lamp is obvious from Fig. 7-8. At the nominal operating point (80 J/cm), roughly 20% of the lamp output is at wavelengths longer than the Nd-pumping region, and roughly 35 to 40% of the output is at wavelengths shorter than the Nd-pumping region. We hope to use this short-wavelength output in our amplifiers by employing spectral converters, as discussed in "Spectral Converters for Nd Laser Pumping," later in this section.

Author: H. T. Powell

Flashlamp Characterization

Introduction. We are interested in characterizing flashlamps over a broad range of parameters to find the most cost-effective solution for Zeus. Provided their optical performances are equivalent, we favor large (5- to 10-cm bore and 200- to 400-cm arc length) rather than small lamps. Large size is attractive from the standpoint of reducing the number of lamps and connecting wires and because energy can be handled in large units. Additionally, we need quantitative information on the emission and absorption characteristics of flashlamps to implement a predictive amplifier model.

During 1982, we studied the performance of the 1.5- by 112-cm flashlamp used in Nova. We also evaluated the 1.0- by 112-cm and the 2.7- by 112-cm lamps that had been considered for use in Shiva. We developed experimental techniques and implemented computer-interfaced diagnostics for comparing diverse flashlamps. With slight modification, our present methods should suffice for testing the very large lamps being considered for Zeus.

For each flashlamp, we measured the electrical characteristics, overall radiant efficiencies (optical energy out divided by electrical energy in), and relative Nd-pumping efficiencies over a wide range of input energies. We compared two methods for determining the radiant efficiency and obtained reliable values for this parameter. We studied the angular dependence of flashlamp emission and compared two techniques for measuring the angle-averaged emission. For a 1.5-cm-bore lamp, we measured the Nd-pumping efficiency vs pulse duration over the range of 0.2 to 1.2 ms.

Diagnostics. We employed a simple set of diagnostics for recording the time-resolved behavior of flashlamps. Diagnostics include lamp voltage and current monitors, a spectrally flat pyroelectric detector to record the radiant output power, and a Nd-doped fluorescer of ED-2 glass to measure the pumping effectiveness of the lamp. The diagnostic signals were digitized with Tektronix digital-processing oscilloscopes (10-bit resolution, either 6.5- or 1- μ s sampling resolution) and analyzed on-line by an LSI 11/23 computer.

Fig. 7-8. Fractional flashlamp energy at wavelengths longer than the cutoff wavelengths (λ_c) of six filters.

The pyroelectric detector (Molelectron P1-35) had an essentially flat spectral response from 0.2 to 10 μm . The signal from such detectors is caused by the displacement current, which is proportional to the rate of change of the electric polarization, as driven by the temperature change produced by the light intensity. Thus, the low-frequency response of the detector is limited by its thermal equilibration time. Its high-frequency response in our case was limited by the bandwidth of the amplifier built into the detector housing. The electrical bandwidth of our pyroelectric detector was 10 to 400 kHz and its optical sensitivity was $\sim 1 \text{ V/W}$.

The Nd fluorescence detector (historically called the fluorescence-output voltage, or FOV, detector) was previously developed by Holzrichter.⁶ Our detector utilized a 1.27-cm-thick disk of ED-2 glass doped at 3 wt% Nd_2O_3 . The flashlamp light passed through an aperture, which was small compared to the disk diameter, and made one pass through the Nd:glass disk into a black-painted pyrex dump. Nd fluorescence was detected by a silicon photodiode placed behind a 0.01- μm bandwidth filter centered at 1.06 μm . The photodiode was geometrically placed at the edge of the disk, outside the cone of flashlamp illumination, such that it viewed only the Nd fluorescence.

Circuit Parameters. For amplifier pumping, the optical performance of different-sized flashlamps should be compared at energies corresponding to a constant fraction of their single-shot explosion energies. In a given amplifier, the wall area for placing lamps is fixed. Assuming fixed values for the lamp-packing fraction, η , arc length, ℓ , and amplifier width, w , the number of

lamps of diameter d is given by $\eta w/d$. If the input energy to the total lamp array is E , then the input energy per unit arc length is $Ed/(\eta w\ell)$. Thus, the energy per unit length is simply proportional to lamp diameter. Since the single-shot explosion energy, empirically found to be⁷

$$E_{\lambda} = 20\,000 \cdot d\ell \tau^{1/2} \quad (1)$$

is also proportional to d , then, for fixed input, the lamps in the various arrays are loaded at a constant percentage of their explosion energies. The circuit time constant, τ , is defined from the bank capacitance, C , and inductance, L , and is given by $(LC)^{1/2}$. It is also convenient to define the explosion fraction, f_{λ} , the ratio of input energy to a lamp divided by its explosive energy defined by Eq. (1).

A schematic diagram of our flashlamp electrical circuit is shown in Fig. 7-9. For different lamp sizes, we used various numbers of fundamental circuits in parallel to vary the input energy. By varying the number of circuits, the time constant, τ , remains constant because $C = nC_0$ and $L = L_0/n$ for n parallel circuits, where C_0 and L_0 are the capacitance and inductance of the fundamental circuit.

To vary the input energy for a given lamp, we simply varied the charge voltage of the bank. However, because of the nonlinear current-voltage relation for the lamp, the electrical-input pulse shape is a function of charge voltage. The peak Nd fluorescence depends on this pulse shape both by pulse duration and by lamp spectral changes. To counter the problem, we chose our electrical parameters so that the pulse shape was constant for each of the lamps at

Fig. 7-9. Electrical circuit employed to excite flashlamps with 1.0-, 1.5-, and 2.7-cm-bore.

a given explosion fraction. Then, even though the pulse shape varies with input, it should at least vary in the same manner for each lamp, thus allowing a direct comparison.

To obtain a constant pulse shape vs f_e curve, we first assumed the empirically measured I-V relationship for an established arc^{8,9}

$$V = K_0 I^{1/2} = k_0 (\ell/d) I^{1/2}, \quad (2)$$

where $k_0 = 1.3 (P/450)^{0.2}$. The Xe fill pressure, P , is in Torr. Neglecting circuit resistance, Markiewicz and Emmett¹⁰ showed that the electrical pulse shape of a flashlamp driven by an L-C circuit is completely specified by a single parameter

$$\alpha = K_0 (VZ)^{-1/2} = nk_0 (\ell/d) (VZ_0)^{-1/2}, \quad (3)$$

where Z is the source impedance $(L/C)^{1/2}$ and n is the number of parallel circuits in our case. Typically, the design point of amplifiers is near $\alpha = 0.8$, which corresponds to critical damping. Setting the energy stored on the capacitor bank at the voltage V to equal some desired fraction f_e of the explosion energy, we have

$$\frac{1}{2} n C_0 V^2 = f_e (20\,000) G d^2 z^{1/2} \quad (4)$$

We can solve Eqs. (3) and (4) for n and V , assuming the arc length ℓ is specified. Thus, the circuit number and charge voltage scale as

$$n \propto \frac{d^3}{k_0^{4/3} \ell} \quad (5)$$

Fig. 7-10. Pulse width vs explosion fractions for flashlamps with 1.0-, 1.5-, and 2.7-cm bores.

and

$$V \propto \frac{k_{\text{eff}}^{2/3} \rho}{d^{1/3}} \quad (6)$$

Equations (5) and (6) give the scaling of circuit number and charge voltage to maintain a constant α vs f_c curve as the diameter, arc length, or Xe fill pressure are varied.

We employed this scaling in going from our standard 1.5-cm-bore, 300-Torr flashlamps to 1.0-cm, 450-Torr and 2.7-cm, 100-Torr lamps. The standard 1.5-cm-bore case utilized two lamps in series with a total arc length of 224 cm and was excited by single circuit ($C_{\text{eq}} = 87 \mu\text{F}$, $L_{\text{eq}} = 450 \mu\text{H}$) over the charge voltage range of 12 to 28 kV. The corresponding calculated f_c variation was from 0.066 to 0.382, while α varied from 1.08 to 0.70. To investigate higher lamp loadings at a 1.5-cm bore, we used a single lamp [$\ell/2$ in Eqs. (5) and (6)], two circuits, and a voltage range starting at 14 kV. To study the 1.0-cm-bore lamp, we used 0.9 circuits (five capacitors instead of the conventional six in parallel) to excite one lamp. For the 2.7-cm case, we chose three circuits to excite two lamps.

To judge the success of our circuit matching technique, we measured the 50%-50% width and the 10%-10% width of the current pulse and compared their variations with f_c for each diameter. As shown in Fig. 7-10, the match was extremely good when compared to the 0.3-ms storage time of ED-2 glass.

Radiant Efficiency. The overall radiant efficiency of flashlamps is limited by the passband (~ 0.2 to $2.5 \mu\text{m}$) of clear-fused quartz (CFQ) and by the collisional loss of heat to the walls. We believe the latter mechanism is the major loss process for CFQ lamps. For lamps with cerium-doped quartz (CeDQ) envelopes, the short-wavelength cutoff is $\sim 0.37 \mu\text{m}$. Optical absorption then becomes comparable to the collisional wall loss.

To measure the radiant efficiency of flashlamps, we used either a total-absorbing calorimeter, which surrounded the lamp, or we measured the temperature rise of the quartz envelope and thereby inferred the heat deposition in the wall. The total absorbing calorimeter consisted of a 5-cm-diam, thin-wall (0.15-cm) copper tube that was chemically blackened.⁴ The cylinder was terminated by end caps, which made a close-fitting (but not touching) connection around the metal electrode shields. Even though the cylinder was not perfectly black at all wavelengths and incidence angles, the emitted light had little likelihood of either leaking out or being reabsorbed by the lamp.

Since the deposition along the length of the cylinder was not uniform, it was necessary to use several thermocouples to obtain a good thermal average. We used nine thermocouples positioned at equal mass units along the can. We simply calculated the calorimeter response from the known heat capacity of copper and the sensitivity of

Fig. 7-11 Temperature rise of a 1.5-cm-bore CFQ flashlamp envelope vs angular position from a linear trigger wire.

chromel–alumel thermocouples. The decay time of the temperature was quite long (~ 600 s) so that extrapolation to the zero point in time was straightforward.

Measurement of energy deposition in the quartz wall was also complicated by the nonuniformity of the initial temperature rise. Figure 7-11 shows the initial temperature rise, measured by a thermocouple held in contact to the wall by tension, as a function of the azimuthal angle from a longitudinal trigger wire. The temperature rise was significantly larger near the wire, demonstrating clearly that the arc does not uniformly fill the tube bore. Using four thermocouples in series, we obtained an average temperature rise that agreed with the angular average at each charge voltage.

To obtain the overall lamp efficiency from the wall-temperature rise, we assumed that the wall deposition was uniform and confined only to the region between the electrodes. By varying the position of the thermocouple array, we found the uniformity assumption to be good. However, we found the wall was also heated slightly behind the electrodes, presumably because of discharge blooming. In addition, there was undoubtedly some heat deposition in the

electrodes themselves. Thus, the wall-temperature rise provided an upper limit on radiant efficiency, while the calorimeter provided a lower limit on lamp efficiency.

Figure 7-12 displays the radiant efficiencies vs input energy measured for the various lamp diameters. The data from wall temperature and the calorimeter gave very similar results and differed in absolute efficiency values by no more than 10%. The efficiencies generally increased with input starting at very low energies, reached a peak value of $75 \pm 5\%$ for CFQ lamps, and then declined slowly with input. The peak efficiency was in good agreement with the result of 80% found by Holzrichter⁶ for a 1.0-cm bore CFQ lamp. Over the bore range of 1.0, 1.5, and 2.7 cm, there was no obvious dependence of the efficiency on diameter. However, the efficiency of the CeDQ lamp in Fig. 7-12(b) is significantly lower because of optical absorption by the Ce-doped quartz wall. Its efficiency is only about 62% at the typical operating point of amplifiers, $f_s = 0.2$.

Angular Dependence. Spectral measurements of flashlamps are complicated by the fact that the spectrum varies with viewing angle. (The depth of plasma in the line of

Fig. 7-12. Radiant efficiency as a function of explosion fraction for four flashlamps.

Fig. 7-13. Radiant-output efficiency of a 1.5-cm-bore CFQ flashlamp vs input energy and Nd-pumping efficiency vs input energy for various viewing angles.

Fig. 7-14. Radiant-output energy vs viewing angle for two explosion fractions.

light and, hence, the effective emissivity, varies with angle.) To determine the significance of the viewing-angle effect, we mounted a flashlamp on a turntable and measured both the Nd-pumping efficiency (from the peak fluorescence detector) and the radiant output efficiency (from the time-integrated pyroelectric-detector signal) vs angle. The detectors were located behind a slit, the width of which was adjusted with angle so that the detectors always viewed a fixed 18-cm length of lamp. The slit aperture was located about 50 cm from the lamp; the detectors were located 100 cm from the lamp. Thus, the viewing angle was confined to $\pm 5^\circ$.

Figure 7-13(a) shows the radiant output efficiency (in relative units) of a 1.5-cm-bore CFQ lamp as a function of input energy for five viewing angles. An angle of 0° implies that the detectors viewed light emitted normal to the lamp axis. The roll-off with input energy (presumably caused by the blue shift of the emission) is more significant at larger angles, consistent with the longer plasma length that is viewed at larger angles. Figure 7-13(b) shows that the Nd-pumping efficiencies (peak Nd fluorescence divided by the input energy) behave similarly, except that the roll-off is more dramatic. The effect seems reasonable since the shortest Nd-pump bands are at about $0.45 \mu\text{m}$ and, thus, the blue shift of the emission is more significant for Nd pumping.

It is useful to plot the radiant output vs angle at fixed input energy to judge how much the angular pattern deviates from a Lambertian (i.e., black) radiator. Figure 7-14 shows the data for the highest and lowest

input loadings used in Fig. 7-13. Signals were normalized in each case to the value obtained at 0° . The dashed curve in Fig. 7-14 is the expected $\cos \theta$ dependence for a Lambertian radiator. The measured outputs are generally above the $\cos \theta$ curve, when normalized to the 0° value, because the emissivity increases with angle. The plot more closely approaches Lambertian at large inputs because the lamp is closer to being "black."

Angle Averaging. Figures 7-13 and 7-14 clearly demonstrate that, to measure the lamp output, which is representative of amplifier pumping, it is necessary to angle average. In the first of two techniques, we mounted a 1.5- by 112-cm lamp along a

Table 7-1. Data for a 1.5- by 112-cm CFQ flashlamp in an integrating sphere or an open geometry.

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	12- to 26-kV ratios		Wall thermocouple signals (mV)	
	Pyroelectric	Nd fluorescence	12 kV	26 kV
In sphere	3.90 ^a	3.86	0.164	0.675
Sphere 25% open	4.27	4.20	0.155	0.625
Open lamp (scatter plate)	4.50	4.38	0.159	0.600
Through sphere	4.30	4.30	-	-

^aTypical errors = $\pm 2\%$.

central chord of a 1-m-diam integrating sphere. The detectors were outside the sphere and viewed only scattered light rather than light directly emitted by the lamp. In the second technique, the detectors simply viewed light scattered from a small plate located near the center of the lamp. Angle averaging occurs in the second case because rays that strike the plate are emitted from different angles at various positions along the lamp axis. Kodak 6080 paint (BaSO₄ in a polyvinyl alcohol binder) was used as a reflector in both cases because of the whiteness and diffuseness of its reflectivity.

Although EG&G had previously used such an integrating sphere for angle-averaging in similar flashlamp studies,¹¹ we had reservations about using their technique. First, because the reflectivity of the Kodak paint falls off in the UV^{11,12} (different measurements disagree as to the amount), there is little question that multiple reflections in the sphere will modify at least the short-wavelength part of the spectrum. Second, the presence of the lamp in the sphere may represent a significant optical loss, which is then a function of lamp loading. Thus, the input-energy dependence of the cavity *Q* can affect the sampled light level.

To compare the integrating sphere with the scatter plate for angle averaging, we used a Nd-fluorescence detector, a pyroelectric detector, and a thermocouple array. The sphere was hinged in the middle to provide internal access and to allow us to look at the effect of having the sphere cracked open. We obtained data at 12- and 26-kV charge voltages (with a standard 87- μ F, 450- μ H circuit and two lamps in series) corresponding to lamp loadings of 25.7 and 120 J/cm. Table 7-1 shows the 12- to 26-kV signal ratios for a lamp in a closed sphere, for the sphere cracked to give an open area of 25%, and for a single reflection off the scatter plate. The "through-sphere" case corresponds to having the lamp outside the

sphere and letting light enter the sphere through a 10-cm-diam hole.

Table 7-1 clearly shows that the 12- to 26-kV signal ratios are significantly smaller for the lamp in the sphere compared to the scatter-plate data. The ratios with the sphere 25% open are intermediate. Moreover, the thermocouple voltages at 25 kV clearly show additional lamp loading (+10%), presumably caused by optical feedback for the lamp in the sphere. It is reasonable that there was no apparent increase in the thermocouple signal for the lamp in the sphere at 12 kV since the lamp absorption is smaller in that case. The 25%-open case should be approximately midway between open and closed geometries since a 50% crack corresponds roughly to a completely open geometry.

The decline in the signal ratios of the lamp-in-sphere data could be caused either by the blue shift of the lamp and, hence, a reduced reflectivity with multiple reflections off the sphere wall, or by the absorption of light by the lamp itself. Both factors reduce the *Q* of the sphere. However, the through-sphere signal ratios are considerably closer to the scatter-plate data than to the lamp-in-sphere data even though multiple reflections also occurred in that case. This outcome suggests that the primary difficulty of having the lamp in the sphere is the reduction in cavity *Q* caused by lamp absorption.

To determine whether the scatter-plate results do, indeed, properly average the lamp emission, we compared scatter-plate results to results obtained by numerically averaging over the angularly resolved data. Since the absolute signals were not comparable, we simply compared the shapes of the curves with respect to input energy. The pyroelectric and peak-Nd curves agreed closely for the different averaging methods. It should be noted that, since the flashlamps were finite in length, the averaging provided by the scatter plate only extended to about 70°. However, we found the

Fig. 7-15. Nd-pumping efficiency vs explosion fraction for various flashlamps.

Table 7-2. Theoretical dependence of V , L , and C on pulse width T required to hold either L , f_x , or I_p constant at constant α .

Constant	V	L	C	E	f_x	I_p
E	$T^{-1/3}$	$T^{4/3}$	$T^{2/3}$	T^0	$T^{-1/2}$	$T^{-2/3}$
f_x	$T^{-1/6}$	$T^{7/6}$	$T^{5/6}$	$T^{1/2}$	T^0	$T^{-1/3}$
I_p	T^0	T	T	T	$T^{1/2}$	T^0

numerically averaged data were not significantly different when averaged either to 70° or 90°.

Nd-Pumping Efficiency. Figure 7-15(a) shows the Nd-pumping efficiency vs input energy (in units of explosion fraction) for the 1.0-, 1.5-, and 2.7-cm-bore flashlamps. The Nd-pumping efficiency is simply defined as the peak Nd-fluorescence signal divided by the input energy to the lamp. The data were angle-averaged using a scatter plate with 15- by 15-cm dimensions. The plate was located 10 cm from the lamp axis. The Nd-fluorescence detector was positioned about 50 cm from the scatter plate and apertured to observe only a central region of the plate. The optical arrangement was fixed so that, by substituting lamps, we were able to make an absolute comparison of the different flashlamps.

Figure 7-15(a) clearly shows that the flashlamps with different diameters perform similarly in pumping Nd:glass when judged to equal fractions of their explosion energy. Thus, at least for fixed pulse length, explosion fraction may be a good parameter for characterizing the spectral output of flashlamps. The rise in the Nd-pumping efficiency with loading at small inputs follows the rise in the radiant efficiency, while the decline at high inputs again results from the blue shift of the lamp emission.

Figure 7-15(b) compares the Nd-pumping efficiencies of flashlamps with CFQ and CeDQ envelopes, both having a 1.5-cm bore. The Nd-pumping efficiencies of these two lamps are considerably closer than their radiant efficiencies, shown in Figs. 7-12(a) and 7-12(b). The slight advantage of CFQ

lamps in Fig. 7-15(b) may be caused by Ce^{3+} -to- Nd^{3+} energy transfer within the ED-2 glass.¹⁴

A primary conclusion from Figs. 7-15(a) and 7-15(b) is that the roll-off in Nd-pumping efficiency up to the normal operation of amplifiers ($f_x \sim 0.25$ after adding optical feedback to the lamp) is quite small ($\sim 10\%$). Hence, we cannot attribute the observed 50% roll-off in the efficiency of amplifiers to this effect. We would argue that Nd pumping by lamps in a cavity is strongly affected by multiple passes through the disks and multiple bounces in the cavity. Single-pass pumping-efficiency measurements do not directly relate to that case. Viewed in a multipass sense, another way of describing the decrease in amplifier efficiency is to say that the effective cavity Q of the amplifier box declines with input power. The cavity Q declines because lamp plasma absorption (coupled with less than unity reemission) and cavity-wall losses increase with input energy. The situation for amplifier pumping is closely related to the results found with the integrating sphere.

Pump Pulse-Length Variations. The excitation pulse length affects Nd:glass pumping directly because of the 300- μs decay time of Nd^* and indirectly because of change in the radiant efficiency and spectral output of the flashlamp. Using a 1.5- by 112-cm CFQ flashlamp, we measured the

Nd-pumping efficiency as a function of pulse length over the range of from 0.2 to 1.2 ms. We defined the pulse length, T , as the 10%-10% width of the current pulse. We experimentally found this width to be about $3.3 (LC)^{1/2}$ at $\alpha = 0.8$.

The electrical circuit was varied to hold the input energy, the peak current, or the explosion fraction constant as the pulse width was varied. We attempted to hold the pulse shape as constant as possible and succeeded in keeping α within the range of 0.7 to 0.9. As a baseline condition, we adopted the electrical characteristics of two lamps excited by a single circuit charged to 22 kV. This configuration gave $T = 0.71$ ms, $\alpha = 0.78$, $E = 90$ J/cm, $I_p = 6.2$ kA, and $f_x = 0.21$.

To vary the pulse width while holding the other parameters fixed, it was necessary to vary the circuit capacitance, inductance, and charge voltage. The circuit equations that govern the lamp in this case are

$$T = A (LC)^{1/2} \quad (7)$$

$$\alpha = k_0 V^{-1/2} (L/C)^{-1/4} \quad (8)$$

$$E = \frac{1}{2} CV^2 \quad (9)$$

$$f_x = E/[2 \times 10^4 \ell d (LC)^{1/4}] \quad (10)$$

and

$$I_p = B V(L/C)^{-1/2} \quad (11)$$

where A and B are constants for fixed α . To find V , L , and C , three of the five equations must be used. Depending on whether the input energy, the explosion fraction, or the peak current is held fixed, Eqs. (7) and (8) must be solved together with either Eq. (9), Eq. (10), or Eq. (11), respectively.

How V , L , and C vary with T is summarized in Table 7-2. Using the power laws shown, it was easy to scale V , L , and C from our baseline condition (700 ms at 22 kV, 87 μ F, and 450 μ H) to obtain other pulse lengths. In practice, our choice of pulse lengths was largely determined by the available inductors. This set the required capacitance, which was much easier to obtain. We then used the charge voltage for fine-tuning to obtain the precise value of the parameter held constant (E , f_x , or I_p).

The expected dependences of the electrical parameters were closely observed. For example, Fig. 7-16 shows how f_x varied for the cases of constant E , f_x , and I_p . From Table 7-2, we expect f_x to vary as $T^{-1/2}$, T^0 , and $T^{1/2}$ in the three cases, as indicated by the solid lines.

The resulting Nd-pumping efficiencies as a function of pulse length are shown in Fig. 7-17. There is surprisingly little difference in holding E or f_x or I_p constant. We suggest that the primary effect of changing the pulse length is caused by the Nd* decay and not by spectral variations of lamp output. The Nd-pumping efficiency generally peaks at about 400 μ s.

From the thermocouple signals, we attribute the fall-off in Nd pumping for pulse lengths shorter than 400 μ s to a decline in the overall radiant efficiency of the flashlamp. From its baseline value, the radiant efficiency decreased 9% at the shortest pulse width and increased 6% at the longest. We presume the decline in radiant efficiency with decreasing pulse width is

Fig. 7-16. Explosion fraction for a 1.5-cm-bore CFQ flashlamp as a function of pulse width.

caused by the blue shift in the spectra and by increased recombination wall loss because the arc channel spends a larger fraction of the pulse width near the wall where it starts.

The results in Fig. 7-17 are for Nd pumping in an open geometry. In a closed geometry, such as an amplifier, the dependence of Nd-pumping efficiency on pulse width may be much weaker. For instance, at constant input energy, the increased lamp absorption that occurs for shorter pulse widths tends to counter the improvements that might be expected. This agrees with our observation of a relatively small increase in gain in the 46-cm amplifier by shortening the excitation pulse from 800 to 500 μ s. The expected dependence of Nd pumping is, in general, quite complex. The pulse-width dependence should depend on the fill factor of both Nd:glass and flashlamps in a given amplifier head.

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Wall Losses in Flashlamps

We developed a simple model for wall losses in xenon flashlamps and used our model to calculate theoretical radiant efficiency. Our results were then used to predict the scaling of efficiency with diameter.

Other than lamp radiation, the two major loss mechanisms are absorption of radiation by a wall and ion recombination at a wall. Because the two mechanisms have different scalings with diameter and input energy, they were treated separately.

Radiative losses occur because a wall is not perfectly transparent to radiation for the entire flashlamp spectrum. Using the Trenholme-Emmett¹⁵ model for temperature, $T(j)$, and emissivity, $\epsilon(\lambda, j)$, the xenon flashlamp spectra can be calculated as a function of current density and diameter and then integrated vs wall-transmission spectra to determine the loss.

Figure 7-18 shows the calculated overall transmission vs current density for clear-fused quartz and cerium-doped quartz. Both

samples were used in our 112-cm arc-length flashlamps. The transmission of CFQ was high over the range of current densities studied, while that of CeDQ decreased for wavelengths $<0.37 \mu$ m.

For lamps of different diameters, we found that the percentage of loss due to absorption of radiation by a wall of CFQ was a very weak function of diameter at constant current density. The percentage of loss increased with increasing current density due to the blue shift of the flashlamp spectra.

Fig. 7-17. Nd-pumping efficiency of a 1.5-cm-bore CFQ flashlamp as a function of explosion fraction.

Fig. 7-18. Transmission vs current density.

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The second mechanism for energy loss of plasma is energy transfer to the walls due to positive ions recombining at the wall. (The conduction of heat via neutral atoms is a much smaller loss.) For the loss due to recombination, we used the Shcherbakov¹⁶ model and simply calculated the flux of ions to the wall. Due to the larger mass of ions, the wall accumulates a negative charge, which then modifies the flux of particles to the wall until an equilibrium is established so that the flux of ions equals the flux of electrons. Since the wall becomes charged, we defined a sheath potential, V_s , which is established at equilibrium by balancing the ion and electron current to the wall.

Thus, the energy from the ions is

$$E_i + 1/2 kT + eV_s \quad (12)$$

where E_i is the ionization energy the wall gains when the ions recombine, and $1/2 kT$ is the energy from the electrons. The total energy to the wall¹⁶ is

$$E_i = \frac{n_i}{\sqrt{2\pi m_i}} (kT) \left[\frac{1}{2} \ln \left(\frac{m_i}{m_e} \right) + 1 + \frac{I_i}{kT} \right] \quad (13)$$

Fig. 7-19. Collisional wall loss vs current density.

Fig. 7-20. Wall loss vs explosion fraction.

where n_i is the bulk number density and T is the wall temperature. The last and largest term represents the energy gain due to recombination and is equal to the power density to the wall (W/cm²).

Results for our simple model of the collisional wall loss as a function of current density are shown in Fig. 7-19. The loss per unit length due to collisions has a direct, linear dependence on diameter (because of area) as well as an implicit dependence on ion density (a function of temperature, which is, in turn, a function of diameter). Since energy into the lamp scales as current to the power 3/2, and since ion density (and, therefore, loss) does not increase as fast, the percentage of loss by the recombination mechanism decreases with increasing current density.

Comparisons with experimental data require that the loss function be integrated over a current pulse. The current pulse was generated by solving the nonlinear differential equation for the circuit (see Ref. 17 for further information). As discussed in detail by Powell,¹⁵ the correct scaling considers changes in diameter while holding the fraction of explosion energy constant between various lamp sizes.

Figure 7-20 shows the final results for a fixed xenon pressure of 100 Torr. Outcomes are shown for CFQ averaged over the current pulse for the combined losses caused by optical absorption and collisions with the wall. Theory tends to predict a lower overall efficiency than that observed experimentally. However, our results show a peak

in the radiative efficiency near an explosion fraction of 0.25, which is similar to that observed experimentally.

At very low explosion fractions, the theoretical losses increase too quickly because many of the assumptions used to calculate the collisional loss (uniformly filled bore, equilibrium established at the walls, are invalid at low current densities. Our results tend to predict that, at constant pressure, the lamps with smaller diameter are slightly more efficient. The theory also predicts that the collisional loss should increase with xenon fill pressure. Neither effect has yet been observed experimentally.

While our characterization of the absorptive wall loss is accurate, the model used for the collisional loss is inadequate at low current densities. Additional studies of the mechanism within a kinetic framework are in order. Nonetheless, our study has supplied a correct qualitative understanding of the loss mechanism.

Author: A. E. Orell

Flashlamp Arc-Growth Model

Xenon flashlamps have traditionally been analyzed by assuming they are arcs in local thermodynamic equilibrium (LTE). Using the LTE assumption, flashlamps are completely described by specifying the arc temperature, gas density, and spectral emissivity of the plasma. Because this method of analysis has not been completely satisfying, the bulk of our knowledge about xenon flashlamps is still empirical.

Xenon flashlamps are actually dynamic devices and are not totally amenable to steady-state analysis. One example of dynamic behavior is the observation that the rate at which the discharge grows from a small filament to a fully developed arc is a function of gas pressure, tube diameter, and energy loading. The fraction of the bore filled by the arc is also a function of these three values. Such dynamic flashlamp behavior can cause, for example, the explosion limit of the flashlamp to vary in an unexpected manner when the discharge-tube geometry is changed. To investigate the issue of arc expansion and bore filling in

xenon flashlamps, we have developed a detailed arc model.

Consider an expanding arc within a xenon flashlamp. The core of the arc is in LTE. In the core region, the electron, ion, and neutral gas temperatures are approximately equal, and the densities of excited states are given by their Boltzmann values. The transition region, between the arc and the neutral cold gas exterior to the arc, is dominated by complex kinetics and hydrodynamics. The neutral gas density rapidly increases across the transition zone, while the gas temperature and ion density rapidly decrease. The electron temperature, however, changes little from its value in the arc. Within the transition zone, convection is driven by the pressure difference between the arc and the outer cold regions. In the region exterior to the arc, the ion density is small or near zero, the electron temperature is below its value in the arc, the gas temperature is little more than its initial value, and the gas density increases as the bore fills with plasma, a result of convection from the arc. The driving force for the processes is the applied electric field. The electric field is a function not only of the external circuit parameters but also of the arc itself, since the arc is a dynamic, time-varying resistance within the discharge circuit.

The species in the arc model are

- Electron temperature, T_e .
- Ion and gas temperature, T_g .
- Ion and electron density, n .
- Heavy-particle gas density, N .
- Xenon excited-state and molecular ion densities, N_i .
- Heavy-particle convective velocity, u .
- Voltages and currents associated with the external circuit elements.

The differential equations for some of the species are

$$\frac{\partial N}{\partial t} = -\nabla \cdot (\bar{u}N) \quad , \quad (14)$$

$$\begin{aligned} \frac{\partial n}{\partial t} = & -\nabla \cdot \left[(\bar{u} + \mu_e) n + D_n \nabla n \right] \\ & + n \left(\sum_i r_i N_i - \sum_{ij} n''_{ij} r_{ij} N_i \right) \quad , \quad (15) \end{aligned}$$

Fig. 7-21. Gas temperature and density in a fully developed arc are initiated at the wall.

Fig. 7-22. Arc width as a function of time and gas pressure for discharges initiated at the center of the discharge tube.

$$\frac{\partial}{\partial t} \left(\frac{3}{2} nkT_e \right) = - \frac{ne^2 E^2}{m \sum_i v_{mi}}$$

$$- \sum_i \frac{3}{2} nk v_{mi} (T_e - T_R) \frac{2m}{M_i}$$

$$- n \left(\sum_{i,j} \epsilon_{ij} r_{ij} N_i - \sum_{i,j} \epsilon_{ji} r_{ji} N_j \right)$$

$$- \nabla \cdot \lambda_e \nabla T_e \quad , \quad (16)$$

$$\frac{\partial}{\partial t} (\rho c T_R) = \sum_i \frac{3}{2} N_i k \frac{2m}{M_i} v_{mi} (T_e - T_R)$$

$$- \nabla \cdot (\bar{u} \rho c T_R) - \nabla \cdot (\lambda_R \nabla T_R) \quad , \quad (17)$$

and

$$\frac{\partial}{\partial t} (\rho \bar{u}) = - \nabla P - \nabla (\rho \bar{u} \bar{u}) - \nabla \cdot \bar{\tau} \quad , \quad (18)$$

Equation (14) is simply the heavy-particle continuity equation. In Eq. (15), μ_i is the ion mobility, ϵ is the radial ambipolar electric field, D_i is the ion diffusion constant, r_i is the electron impact ionization rate for spe-

cies N_i , and r_{ii} is the ion recombination coefficient for species i by process i . The forces resulting from self-generated magnetic fields in the arc are incorporated into the radial ambipolar electric field.

In Eq. (16), E is the longitudinal electric field, m and M_i are the electron and heavy-particle masses, v_{mi} is the electron momentum-transfer-collision frequency with species i , r_{ij} is the rate of electron impact excitation or relaxation between levels i and j of xenon with energy separation ϵ_{ij} , and λ_e is the electron thermal conductivity. In Eq. (17), ρ is the heavy-particle mass density, c is the heavy-particle heat capacity, and λ_R is the heavy-particle thermal conductivity. In Eq. (18), P is the gas pressure and $\bar{\tau}$ is the viscosity tensor.

As written, Eqs. (14) through (18), coupled to a radiative transport calculation, are more appropriate to nonequilibrium kinetics than an LTE plasma. The complexity of the problem results from the necessity of solving our equations for the transition region between the arc and cold regions where kinetics dominate, while also ensuring LTE conditions within the arc. The compromise between kinetics and LTE was made by assuming that the excited states of xenon and the partitioning between atomic ions and molecular ions are in equilibrium with the species described above. The densities of these states are given by the steady-state solution to the appropriate differential equations. The assumption is quite good, since the time required to reach equilibrium by the excited states is less than a microsecond, while the hydrodynamic time scale is tens of microseconds.

To account for photoionization processes resulting from recombination radiation within the arc, an effective electron-recombination coefficient was used. The coefficient was solved for by comparing the LTE value for ion density with the value

calculated within the arc. As LTE is established in the arc, the recombination coefficient is changed from its kinetic value to its effective LTE value. Our treatment assumes an optically thick plasma within the arc. That is, the absorption length for radiation is less than the characteristic distance over which the plasma properties change. The assumption is not good at the edge of the arc, where the ion density changes rapidly; however, it does enable us to treat the radiation transport as a local property, thereby greatly simplifying the calculation. Despite this imperfection, the results of the model appear to be in qualitative agreement with observed behavior.

The calculation begins by specifying the initial gas pressure, tube diameter, circuit parameters, and the location at which the initial breakdown filament occurs within the discharge tube. The external circuit is triggered, and the circuit equations and partial differential equations are then integrated in time.

Computed gas density and gas pressure in a fully developed arc are plotted in Fig. 7-21 for an initial fill of 300 Torr of xenon in a 1.5-cm-diameter tube (gas density $\approx 9.56 \times 10^{18} \text{ cm}^{-3}$). The breakdown filament was assumed to occur along the wall of the discharge tube. The fully developed arc does not fill the entire bore of the discharge tube. The fraction of the bore that the arc fills is a function of gas pressure, tube diameter, and location of the breakdown filament.

The arc width and filling fraction (defined as an area ratio) for a variety of conditions are plotted in Figs. 7-22 through 7-24. The edge of the arc was assumed to be the point at which the current density is 0.1 of the maximum value at that particular time. Figures 7-22 and 7-23 show that the maximum filling fraction is obtained when the gas pressure is low and that the filling fraction decreases with increasing tube diameter. (The fraction of explosion energy deposited in the discharge is constant for each plot.) Figure 7-24 displays the decrease in filling fraction when the location of the breakdown filament moves away from the center of the discharge tube and towards the wall.

The slowing of the rate of arc expansion and its eventual stalling before filling the bore is, to a first order of approximation, a result of kinetic processes, but is ultimately

Fig. 7-23. Arc width as a function of time and tube diameter for discharges initiated at the center of the discharge tube.

Fig. 7-24. Arc-filling fraction as a function of gas pressure and position of the breakdown filament.

driven by hydrodynamics. As the discharge heats the gas within the arc, an expansion wave pushes gas into the cool region outside the arc, thereby increasing the gas density in that region. The electron temperature, proportional to E/N , decreases in the cool gas region. The ionization rate is proportional to $e^{1/2}$, while the recombination rate is proportional to $1/T_e^2$. Therefore, as the gas density increases in the cool region, the ionization rate decreases and the recombination rate increases.

Eventually, the electron temperature falls below a critical value, at which time the ionization rate is less than the recombination rate. This effect is exacerbated by the fact that the rate at which molecular ions are formed is proportional to the square of the gas density. Therefore, the fraction of molecular ions increases in the transition zone and cool regions. The recombination rate for molecular ions is, however, more than an order of magnitude larger than for the atomic ions. Therefore, the effective recombination rate for ions in the cooler transition zone is even greater than would be predicted based on only the electron temperature. When the breakdown filament

Fig. 7-25. Predicted wall temperature on the inside surface of the discharge tube as a function of time and position.

occurs in the center of the discharge tube, the average gas density outside the arc is smaller than if the filament occurs near the walls. The decrease in gas density allows the arc to fill a larger fraction of the bore.

In addition to the information concerning bore filling, described above, parameters such as flashlamp efficiency, wall temperature, and pressure rise within the tube can be obtained from our model. For example, the inside-surface wall temperature of the discharge tube can be calculated from values obtained from the model for the ion and electron flux to the walls,¹⁰ and the known thermodynamic properties of the quartz wall material. Neglecting radiation, a prediction of the inside-surface wall-temperature distribution for a discharge in 300 Torr of xenon ($D \approx 1.5$ cm), where the breakdown filament was located halfway between the wall and the center of the tube, is shown in Fig. 7-25.

Author: M. J. Kushner

Spectral Converters for Nd Laser Pumping

Xenon flashlamps that pump Nd:glass lasers emit a broad-continuum spectrum.

However, the Nd ions absorb radiation in discrete, albeit broad, bands. The fraction of flashlamp energy absorbed by Nd ions is a function of many parameters, including the optical depth of the glass and the energy loading of the flashlamps. Lightly doped glass is fairly transparent to the pump light, whereas heavily doped glass absorbs about 40% of the flashlamp light into the laser pump bands.

The absorption spectrum for ED 2 glass is plotted in Fig. 7-26(a). The energy fractions in the $0.37\text{-}\mu\text{m} \leq \lambda \leq 0.94\text{-}\mu\text{m}$ band that are not absorbed by ED 2 Nd:glass are plotted in Fig. 7-26(b) as a function of lamp loading and three optical depths.¹¹ In addition, the fractions of flashlamp energy emitted but not absorbed in the two bands $\lambda < 0.37\text{ }\mu\text{m}$ and $\lambda > 0.94\text{ }\mu\text{m}$ are also shown.

Approximately 26% of the flashlamp energy is radiated at $\lambda < 0.37\text{ }\mu\text{m}$ and is not absorbed by Nd ions. Similarly, from 16 to 18% of the flashlamp energy is radiated at $\lambda > 0.94\text{ }\mu\text{m}$ and is not absorbed. The fraction of energy not absorbed between 0.37 and $0.94\text{ }\mu\text{m}$ is greater than 25% for the optical depths of interest.

In principle, it should be possible to capture the unused radiation in the three spectral regions to increase the rate at which Nd ions are pumped. Two types of converters are candidates for the purpose of capturing unused radiation: energy converters and fluorescent converters. Energy converters are probably less generally recognized.

The flashlamp itself is an efficient converter of electrical to optical energy. For our operating conditions, the plasma is in local thermodynamic equilibrium throughout most of the discharge period. Thus, with the exception of the physical location of energy deposition within the volume of plasma, absorbed optical energy will heat the emitting and absorbing xenon plasma, with effects equivalent to electrical-energy deposition. As a result, the flashlamp can be used as an energy converter. The only significant difference between electrical heating and optical heating is the resulting plasma-temperature profile. Thus, we may assume that absorbed optical radiation will be converted to emitted optical radiation with nearly the same overall efficiency with which deposited electrical energy is converted to optical radiation. A detailed analysis of flashlamps as energy converters is

discussed in Ref. 20. The limiting efficiency of pumping Nd ions can be increased by 25% by optimal energy conversion by the flashlamp.

A fluorescent converter is a material external to the Nd:glass medium, doped with an ion or molecule that absorbs in a region not matching Nd-ion absorption bands, and that emits fluorescence that matches Nd-ion bands. The converter medium can be either a solid or liquid.

One problem associated with the use of fluorescent converters to enhance the pumping of a laser medium concerns the placement of the converter in relation to the laser medium to ensure efficient transport of converted radiation to the laser medium. Five possible geometries are illustrated in Fig. 7-27. The geometries with the highest fraction of converted radiation coupled into the Nd:glass place the converter medium in optical contact with the laser medium.

The geometry shown in Fig. 7-27(d) is appropriate for conversion of the short-wavelength radiation that will not pass through the glass. However, the geometry requires pump-cavity reflectors that have excellent reflectivity in the UV portion of the spectrum (not true for Ag reflectors). The geometry shown in Fig. 7-27(e) is appropriate for conversion of the long-wavelength radiation flowing between the Nd pump bands. The geometries shown in Figs. 7-27(a) through 7-27(c) lose a significant fraction of the fluorescence to total internal reflection of the light when passing from the converter medium to air. Some of the loss may be recouped by incorporating

scattering centers into the converter to couple out the otherwise trapped fluorescence.

The obvious requirement for good converter candidates is that their absorption and emission bands be favorably located with respect to those of Nd:glass. In addition, their quantum yields must be large and their radiative lifetimes short, compared to the length of the flashlamp pump. For converter candidates that use the geometry shown in Fig. 7-27(d), the requirement of minimum overlap between the absorption bands of the converter and Nd:glass is especially important. The Nd ions have a presumed quantum yield of unity for the laser

Fig. 7-26. Absorption spectrum for ED-2 glass ($\text{Nd}^{3+} \cdot 10^{21}$ ions/ cm^2) and the fractions of flashlamp energy emitted but not absorbed by ED-2 Nd:glass.

	(a)	(b)	(c)	(d)	(e)
Fluorescent converter					
I_r pump-cavity transfer efficiency					
n = Refractive index of converter	1.5	1.5	1.45	$\approx n$ (Nd:glass)	$\approx n$ (Nd:glass)
Fraction of converted radiation available for absorption in Nd	$0.13 T_r$	$0.18 T_r$	$0.21 T_r$	$0.87 T_r$ (UV) (for $n = 1.5$)	$1.0 T_r$ (for $n = 1.5$)

Fig. 7-27. Five basic optical geometries for the placement of extrinsic fluorescent converters.

pump bands. The net conversion efficiency for the converter (quantum yield $\times$ geometrical coupling efficiency $\times$ absorption probability of Nd^{3+}) is always less than 1. Therefore, a severe penalty is paid if there is substantial overlap between the absorption bands of Nd:glass and those of the converter.

Two classes of candidates for a short-wavelength converter have been identified: divalent rare-earth ions (in glasses or powders),²¹ and organic dyes (in solution or plastics).²² Of the former candidates, divalent europium (Eu^{2+}) is especially attractive; of the latter candidate, several have a high probability of success.

Three classes of phosphors are under consideration: conventional powder phosphors, phosphors incorporated in optically clear glass, and phosphors bound in white polymers. A powder phosphor would be coated on the interior of the amplifier box and would act as a white, diffuse reflector at wavelengths other than its absorption band. Alternatively, the powder might be incorporated into a solid, polymeric sheet material serving as the reflector.

One potential problem in using polymers, either as sheets or as binders for a paint-like coating, is that they tend to degrade under flashlamp illumination. Degradation also applies to organic dyes. Coatings prepared with available inorganic binders do not have the problem of degradation; however, they have relatively poor mechanical properties. Optically transparent phosphors would ideally be located as sheets sandwiched about the amplifier laser glass. Alternative locations are as blast shields or flashlamp envelopes, but the improvement of performance would decrease because of losses associated with total internal reflections, as previously mentioned. Glass is a more attractive host than polymers from the standpoint of lessened optical damage by the laser beam and flashlamp emission.

For phosphors incorporated in glasses, no known ion is ideal for our purpose, but divalent europium is a reasonable candidate. Europium tends to enter alkali silicates in the trivalent state, which has unsuitable spectral characteristics. However, europium is readily incorporated in the divalent state in fused silica or in 96% silica. The generic

term 96% silica is used for a class of glasses manufactured by the Vycor process, which consists of manufacturing the sample, with oversized dimensions, from a certain alkali borosilicate glass. A heat-treatment phase separates the glass to an interconnected two-phase structure. One phase is high in alkali, and the other is approximately 96% SiO_2 , 3% B_2O_3 , and 0.4% Al_2O_3 with traces of alkali and As. The alkali-rich phase is chemically leached out of the glass, leaving a porous body with an interconnected porous structure having 5-nm pores. The porous material is then sintered at moderate temperatures to a solid body, with shrinkage giving the sample its desired size. The overall process is quite economical.

The porous intermediate of the Vycor process is commercially available and can be doped with various ions prior to consolidation. In the case of europium, the porous glass is placed in an aqueous solution of $\text{Eu}(\text{NO}_3)_3$, which is drawn into the pores by capillarity. As the glass is dried and the water evaporates, the salt is deposited in the pore structure. Upon heating, the nitrate decomposes to an oxide; at higher temperatures, the glass sinters and the oxide dissolves in the glass, leaving the Eu ion in the divalent state. Other ions can be added along with Eu^{2+} by simply adding an appropriate salt to the solution. The salt permits some adjustment of the Eu^{2+} spectrum.

We have prepared samples of 96% silica doped with Eu^{2+} in this way. Our samples fluoresce blue under UV illumination, with the emission band centered at $0.44\ \mu\text{m}$. The maximum absorption band is at $0.37\ \mu\text{m}$. The emission band is somewhat too short a wavelength for ideal Nd pumping. We will attempt to shift the emission spectrum to somewhat longer wavelengths by codoping with the other ions and by varying the heat-treatment conditions to maximize overlap with the Nd pump bands.

We have developed a model to aid in evaluating converter candidates. Our model calculates the change in Nd pumping rate as a function of converter type and geometry, glass type, and doping level using the xenon flashlamp spectrum calculated from the Trenholme-Emmett model.²³ For evaluating candidates, our results from the

model are discussed in terms of ED-2 glass and the double-sided surface-converter geometry shown in Fig. 7-27(d).

An important parameter in evaluating converter candidates is the optical depth (ions/cm²) of Nd:glass. Results for the change in Nd pumping rate as a function of optical depth are plotted in Fig. 7-28. The results in Fig. 7-28(a) are for organic-dye candidates,⁹ and those in Fig. 7-28(b) are for Eu²⁺ converters in different hosts. The coupling efficiency for surface converters was calculated to be 0.85 because of the benefit of total internal reflection. Increases in the Nd³⁺ pumping rate in excess of 20% are possible for our best converter candidates. The amount of increase, however, is a function of the optical depth of the glass, as discussed below.

For one class of converters shown in Fig. 7-28 the change in pumping rate decreases as optical depth of the Nd:glass increases. The decrease results from the fact that, as the optical depth of Nd:glass increases and its absorption bands broaden, the overlap between the absorption bands of the converter and Nd:glass increases. As discussed above, this is not a favorable situation for efficient conversion. For the second class of converters shown in Fig. 7-28, the change in pumping rate increases as optical depth of the Nd:glass increases. The increase results from the broadening absorption bands of

Nd:glass, increasing their overlap with the emission bands for the converter.

Our laboratory evaluation of converter candidates has begun. We have given special attention to organic-dye phosphors doped with Eu²⁺ for short-wavelength converters. The converter candidate is placed in contact with a disk of Nd:glass to simulate the geometry shown in Fig. 7-27(d), and the assembly is exposed to a pulsed flashlamp. A photodiode equipped with a narrow bandpass filter monitors fluorescence from the upper laser level of Nd. Comparing the Nd fluorescence for a converter material with and without the active ingredient (undoped vs doped Vycor) yields the change in Nd pumping resulting from the conversion process.

Typical experimental results for the change in Nd pumping with the dyes coumarin 522 and 9,10-bisphenylethynyl anthracene are plotted in Fig. 7-29. Experimental results are compared to our model calculations. Increases in Nd pumping (>10%) have been demonstrated when using single sided conversion and pumping. The dyes were in a liquid solvent and contained in a reservoir that was in contact with the Nd:glass. The coupling efficiency assumed in our calculations was 0.45 because total internal reflection in the outer window is not useful. In contrast, a coupling efficiency of 0.85 would be obtained

Fig. 7-28. Calculated change in Nd pumping with organic dyes and with Eu²⁺ in different hosts as converters.

Laser-Induced Damage

Fig. 7-29. Experimental results for the change in Nd pumping using two organic dyes as converters.

Fig. 7-30. Absorption and emission spectra of Eu^{2+} in Vycor manufactured by two different processes.

with the dye dissolved in a low-index plastic and optically coupled to the glass. (The calculated curves in Fig. 7-28 are for a coupling efficiency of 0.85.) Practical implementation of dyes as converters for disk glass lasers requires the development of suitable plastic solvents to obtain the required optical-damage limit and high coupling efficiency.

Samples of Vycor glass doped with Eu^{2+} , using the process described above, were tested as converter candidates. The absorption and emission spectra of a typical sample are plotted in Fig. 7-30. The spectra are shifted to the blue from the ideal; nevertheless, the sample yielded an increase in Nd pumping in excess of 8%. Also shown is the spectra of Eu^{2+} in a Vycor host processed in a somewhat different manner,²⁵ illustrating the discrete tunability of the method. This result is especially encouraging, since glass converters would be more easily implemented than organic-dye-doped plastics.

Increases in Nd pumping of more than 10% have been measured for extrinsic fluorescent converters. Our modeling results indicate the increases would be in excess of 20% for conditions with a larger coupling efficiency.

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Laser-Induced Damage

Introduction

The Nova laser requires coatings that are either antireflective (AR) or highly reflective (HR) at some combination of the 1ω , 2ω , and 3ω wavelengths. Reflective-focusing optics that are highly reflective for wavelengths between 1.06 and 0.25 μm may be needed for future lasers. Having studied laser damage at several frequencies during 1982, we first review our experiments at 1ω , 2ω , and 3ω that were performed to support the design of Nova. In the final section, we review our experiments at 0.25 μm.

Damage at 1ω , 2ω , and 3ω for AR Surfaces

Neutral-solution processing²⁶ (NSP) produces damage-resistant AR surfaces on components made of borosilicate glasses that can be used to transmit high-intensity pulses at either 1ω or 2ω . A physical plant for NSP treatment of surfaces on lenses up to 1 m in diameter was installed at LLNL.²⁷

The damage thresholds measured with 1-ns, 1ω pulses on our NSP surfaces ranged from 12 to 14 J/cm². Our range is comparable to thresholds of the NSP surfaces initially made at Schott Optical Glass, Inc. (Duryea, Pa.)²⁶

Recent experiments with NSP-treated surfaces indicated that the threshold for damage induced by 0.7-ns, 2ω pulses was the same as that for damage induced by the parent 1-ns, 1ω pulses.²⁷ The damage thresholds were the same for entrance and exit surfaces of windows with NSP surfaces,²⁸ and the threshold of NSP surfaces scaled as the square root of laser-pulse duration (see Fig. 7-31 and Ref. 28).

We also determined that the damage threshold of an NSP-treated surface could be increased by repetitively irradiating the surface at fluences below the damage threshold (*N-on-1* irradiation). In a typical test, during which each test site is irradiated only once (*1-on-1* irradiation), laser-induced emission of light occurred on test sites irradiated at 8 to 10 J/cm² in 1 ns. Damage visible by Nomarski microscopy was produced on sites irradiated at fluences exceeding 12 J/cm². If, however, a given test site was repetitively irradiated at fluences of from 8 to 10 J/cm², laser-induced emission of light ceased after a few shots. The fluence could then be increased to a level exceeding the typical *1-on-1* threshold before either light emission or damage occurred. The median *N-on-1* threshold measured with 1-ns pulses was 17 J/cm²; the maximum was 25 J/cm². Increase of thresholds by *N-on-1* irradiation was observed at pulse durations from 1 to 20 ns (see Fig. 7-31), although fractional increases in rear-surface thresholds were less for the longer pulses.

The mechanism for the threshold increase is unknown, but we assume that the surface is cleaned by low-fluence irradiations. We were unable to measure a change in surface reflectivity on sites subjected to *N-on-1* irradiations. Evidently, the porous surface is not collapsed by the irradiations. Some photographs taken with a scanning electron microscope showed that areas subjected to *N-on-1* irradiation were free of small particles that existed on adjacent areas of the surface.

The NSP process does not produce AR surfaces on fused silica, which is the only material presently capable of transmitting

Fig. 7-31. Damage threshold 1ω vs laser-pulse duration for NSP gradient-index AR surfaces of BK-7 glass.

high-intensity 3ω laser pulses. Therefore, we continued our evaluation of several techniques for producing AR coatings on fused silica for use at 3ω . Figure 7-32 shows thresholds measured with 0.6-ns, 3ω pulses on 16 AR films made by electron-beam evaporation at Optical Coating Laboratories, Inc. (OCLI, Santa Rosa, Calif.); thresholds ranged from 2 to 3.5 J/cm². The films with the greatest resistance to damage were made of scandia/magnesium fluoride, a combination of materials that was tested at 3ω because of its resistance to damage by 0.25- μ m KrF laser pulses.²⁹

We investigated several techniques for producing graded-index AR surfaces on silica. Scientists at Westinghouse Research and Development Center, Pittsburgh, Pa., demonstrated the production of a porous silica layer by spreading an organometallic solution on a silica substrate and baking it to drive off the volatile compounds.³⁰ The coated substrate was then held at a temperature of from 400 to 700°C for several hours to partially sinter the coatings and reduce the median size of the coating pores to ~ 40 Å. Etching was done in weak hydrofluoric acid. The resultant porous layer had tapered pores, a graded refractive index, and behaved as a single-layer AR coating. Transmission $>99\%$ at 3ω , and some reduction of surface reflectivity at all wavelengths from 3ω to 1ω , was obtained for silica substrates coated on both surfaces. The damage thresholds measured with 0.6-ns, 3ω pulses on 20 of the porous silica surfaces ranged from 0.6 to 4.1 J/cm²; the median threshold was 2.2 J/cm².

We also evaluated porous coatings made at Owens-Illinois, Inc. (Toledo, Ohio) and at

Fig. 7-32. Damage thresholds with 0.6-ns, 3ω pulses on four-layer AR coatings made with four combinations of materials.

Fig. 7-33. Cones formed by ENT damage in silica.

Battelle Research Laboratories (Columbus, Ohio). Their coatings were made by sol-gel technologies, similar to that described above, in that a homogeneous surface layer of borosilicate glass was deposited on a polished silica surface. A porous layer, composed primarily of silica, was then produced by leaching the glass in an acid that preferentially removed constituents other than silica. The thresholds of two films from each vendor, measured with 0.6-ns, 3 ω pulses, ranged from 1 to 2 J/cm². The thresholds of 15 sol-gel films made by Owens-Illinois, Inc., measured with 1-ns, 1 ω pulses, ranged from 7 to 15 J/cm².

We believe that damage thresholds in coatings made by sol-gel techniques are

limited by inclusion of microscopic contaminants in the solutions. However, clean sol-gel films have potential as damage resistant, broadband AR coatings. Therefore, we have begun a major effort to produce porous coatings on large components under clean-room conditions.

As a backup to both electron-beam evaporation and sol-gel techniques, we studied production of graded-index surfaces by etching-nuclear-track (INT) damage in silica. The substrates were covered with a foil of enriched uranium and bombarded with thermal neutrons. Fission fragments entering the silica left trails of damaged material, which etched at a rate greater than the etch rate of the undamaged silica. Etching a single track produced a conical pit. The cone angle of about 25° was estimated by using a scanning electron microscope to photograph replicas of cones that were etched to a depth of several micrometers. An example of such a photograph is shown in Fig. 7-33. A surface that bears a high spatial density of conical pits has a graded refractive index. However, to serve as an AR film at a wavelength of $1\text{ }\mu\text{m}$, the graded layer must have a minimum thickness (cone depth) of $\sim 0.3\text{ }\mu\text{m}$. The problem is that cones with the necessary depth have diameters so large that they are efficient scatterers of light. Our experiments, aimed at producing a larger differential in etch rates and smaller cone angles, are continuing.

Dielectric Reflectors

In studies of HR coatings, our first priority was the testing at 1ω , 2ω , and 3ω of coatings made of common oxide materials that will probably be used to fabricate reflectors for Nova. We tested 24 quarterwave-stack 3ω HR coatings made by electron-beam evaporation at OCLI. The coating designs were (HL)³HL, where H and L are the quarterwave thicknesses of high-index and low index materials, respectively. The low-index material for all reflectors was silica; the high index materials were zirconia, tantala, and hafnia.

Thresholds for the coatings shown in Fig. 7-34 were measured with 0.6-ns, 3ω laser pulses. Thresholds of the zirconia/silica and tantala-silica coatings ranged from 1.7 to 3.3 J/cm^2 . Thresholds of the hafnia/silica coatings were lower ($\sim 1 \text{ J/cm}^2$), but the outcome is probably a characteristic of the particular coatings tested and not a general indication that hafnia is inferior to either zirconia or tantala.

An unusual result was that thresholds of overcoated reflectors were less than or equal to thresholds of nonovercoated reflectors. The outcome is contrary to results obtained in previous experiments at both $0.25 \text{ }\mu\text{m}$ (Ref. 31) and 1ω (Ref. 32). Addi-

tional tests will be necessary to determine whether the result is generally true. We are presently testing 3ω HR coatings made with scandia/silica and scandia/magnesium fluoride, combinations that were previously found to be resistant to damage at $0.25 \text{ }\mu\text{m}$. Preliminary results indicate that thresholds for these films will routinely exceed 3 J/cm^2 .

We tested 24 trichroic mirrors, highly reflective at 1ω , 2ω , and 3ω , made by depositing one multilayer reflector over another. Some of the individual reflectors were quarterwave stacks, which are reflective at the primary design wavelength, λ , and at $\lambda/3$. Using this design alone, a trichroic reflector can be made by first depositing a primary 1ω mirror (also reflective at 3ω) and then depositing over that mirror a primary 2ω reflector. The trichroic reflector is designated as $Q(2\omega)/Q(1\omega-3\omega)$.

The individual reflectors can also be made using a nonquarterwave (NQ) design in which the optical thickness of the low-index layers is twice that of the high-index layers. Nonquarterwave 2:1 coatings are reflective at a primary design wavelength, λ , and at $\lambda/2$. Some trichroic reflectors were made by placing a conventional quarterwave-stack 3ω reflector over a 2:1 coating that was reflective at 1ω and 2ω . Such a trichroic reflector is designated as $Q(3\omega)/NQ(1\omega-2\omega)$.

Fig. 7-34. Damage thresholds with 0.6-ns, 3ω pulses on HR coatings made with three combinations of materials.

Fig. 7-35. Damage thresholds at 1ω vs pulse duration for HR coatings.

Trichroic reflectors fabricated by both OCLI and Spectra Physics, Inc. (SP, Mountain View, Calif.) were tested with 1-ns, 1ω pulses; 0.7-ns, 2ω pulses; and 0.6-ns, 3ω pulses. None of the 24 trichroic reflectors had good resistance to damage at all three wavelengths.

The most successful coating supplied by OCLI was a $Q(3\omega)/NQ(1\omega-2\omega)$ design fabricated of zirconia/silica. The average thresholds at 1ω , 2ω , and 3ω were 6, 4, and 2 J/cm^2 , respectively. The best trichroic reflector supplied by SP was a $Q(2\omega)/Q(1\omega-3\omega)$ design fabricated of hafnia/silica. The average thresholds at 1ω , 2ω , and 3ω were 9, 10, and 0.3 J/cm^2 , respectively. The thresholds observed for the outer reflectors in each design were 2 J/cm^2 for the $Q(3\omega)$ reflector and 10 J/cm^2 for the $Q(2\omega)$ coating.

The results for outer reflectors are comparable to, or better than, thresholds observed in corresponding monochromatic coatings. With the exception of the 9 J/cm^2 threshold measured at 1ω on the SP reflector, thresholds of the buried reflectors in each design were lower than would be expected had the buried reflector been tested individually. The low thresholds for buried reflectors may arise from a systematic difficulty that limits our ability to make damage-resistant trichroic reflectors.

We investigated one new technique for depositing thin-film coatings. Ion-beam sputtering was developed at the Honeywell Corporate Technology Center (Bloomington, Minn.) as a technique for producing low-scatter mirrors. Fourteen 3ω tantalum/silica HR films supplied by Honeywell were tested with 0.6-ns, 3ω pulses. The damage thresholds ranged from 0.5 to 2.6 J/cm^2 ; the median threshold was 1.4 J/cm^2 . The coatings had a spatially uniform damage morphology that suggested that the coating absorption was large. Honeywell is investigating production of coatings with lower absorption.

Pulses with durations of 1, 3, 6, 9, and 20 ns at 1ω were used to measure damage thresholds of a titania/silica quarter-wave-stack reflector designed for use at 1ω and a zirconia/silica trichroic reflector of the design $Q(2\omega)/Q(1\omega-3\omega)$. Thresholds for the 1ω reflector were measured by 1-on-1 tests; those for the trichroic were measured by N-on-1 tests.

The measured thresholds shown in Fig. 7-35 scaled with pulse duration, τ , as $\tau^{0.4}$. The trichroic reflector was tested N-on-1 because it exhibited a behavior that is highly unusual for multilayer dielectric films. Its 1-on-1 threshold was low (4 J/cm^2 at 1 ns), but the damage only consisted of slight enlargement of a few existing pinholes. When test sites were repeatedly irradiated (N-on-1), fluences larger than the 1-on-1 threshold were required to produce either additional damage or laser-induced emission of light. For both reflectors, thresholds measured with 20 ns pulses fell within the distribution of thresholds measured with 23 ns, $0.69\text{-}\mu\text{m}$ pulses.³³

Bare Polished-Silica Substrates

Since fused silica is the only material known to be free of two-photon absorption and color-center formation when irradiated by intense 3ω pulses, we continued to study damage on bare polished-silica substrates. Three materials were tested with 0.6-ns, 3ω pulses: Suprasil II and Optica I (Amersil, Inc.) and a recently improved code-7940 silica (Corning Glass Works, Inc.). Test substrates were polished by Zygo Corporation. Thresholds for the three types of silica ranged from 8 to 11 J/cm^2 .

We used 1-ns, 1ω pulses to test nine silica substrates that were polished at the Honeywell Corporate Technology Center on dry powders: either colloidal silica or strontium carbonate, or a mixture of nickel and manganese oxides. Material removal in this polishing process occurs primarily by chemical attack, rather than by abrasive fracture, so that little subsurface fracture is induced. Damage thresholds for the nine silica substrates ranged from 5 to 17 J/cm^2 ; the median threshold was 12 J/cm^2 . The thresholds were somewhat lower than the median threshold of 16 J/cm^2 for substrates polished in conventional abrasive ceria slurries.

We tested two glasses, NG-1 and NG-4 (Schott Optical Glass, Inc.), that will be used in absorbing-glass calorimeters. NG-1 thresholds measured with 0.7-ns, 2ω pulses and with 0.6-ns, 3ω pulses were 9 and 4 J/cm^2 , respectively. The corresponding thresholds for NG-4 were 13 and 6 J/cm^2 .

Plastics

We measured bulk-damage thresholds of several plastics with single 1-ns, 1ω pulses. The plastics were PMMA, CR-39, and polycarbonate. PMMA is polymethyl methacrylate (Lucite, Perspec, and Plexiglas) with outstanding transparency. Two PMMA samples were of general-purpose grade; one was of clinical grade (used to make lenses for human eye implants). CR-39 has the common chemical name allyl diglycol carbonate and is used in spectacles because of its exceptional abrasion resistance. Two of the CR-39 samples were of a general grade; two were fabricated (PPG Industries) under documented conditions from monomer filtered to $1.2 \mu\text{m}$. Polycarbonate (Lexan, Tuffak) is the polymer with the repeating unit

and is known for its impact resistance. While none of the polymers had special additives (UV absorbers or flame retardants) or surface treatments, the clinical-grade PMMA might plausibly contain a UV absorber. The measured damage thresholds are listed in Table 7-3.

Material	Number of samples	Damage threshold (J/cm^2)
CR-39, filtered ^a	2	$2.8 \pm 0.6, 4.1 \pm 0.4$
CR-39, general ^b	2	$2.2 \pm 0.2, 2.6 \pm 0.3$
PMMA, clinical ^c	1	3.4 ± 1.0
PMMA, general ^d	2	$4.2 \pm 0.4, 4.6 \pm 0.5$
Polycarbonate ^e	1	<1.4

^aSupplied by S. D. Jacobs (Univ. of Rochester) and fabricated by Richard Schwarz (PPG Industries).

^bSupplied by Signet Optical Corp.

^cPerspec CQ (Imperial Chemical Industries).

^dPlexiglas V920 (Rohm and Haas).

^eTuffak A (Rohm and Haas).

Table 7-3. Damage thresholds in bulk-plastic materials tested with 1-ns, 1ω pulses.

Table 7-4. Damage thresholds for thin sheets of several plastics.

Material	Test wavelength	Pulse duration (ns)	Damage threshold (J/cm^2)
Sun Gain	1ω	6	9.2 ± 0.9
	1ω	1	6.0 ± 0.6
	2ω	0.7	<0.7
	3ω	0.6	<0.55
Teflon FEP			
4 mil ^a	1ω	6	12.0 ± 1.3
5 mil ^a	1ω	1	6.7 ± 0.8
4 mil ^a	2ω	0.7	7.4 ± 1.1
4 mil ^a	3ω	0.6	3.2 ± 0.5
1 mil ^a	1ω	1	3.2 ± 0.3
1 mil ^b	1ω	1	3.0 ± 0.3
Mylar (1 mil, extra clean)	1ω	1	0.9 ± 0.1
Polyvinyl alcohol (unstretched)	1ω	1	1.5 ± 0.5

^aNo surface treatment.

^bTreatment on one side.

In all samples, damage consisted of isolated internal microfractures that are presumed to result from heating of included absorbers. While absorbers might be excluded by careful manufacturing procedures, it may not be possible to produce plastics with damage thresholds equal to that of glass. Failure depends on the material strength and the temperatures of the inclusions, so that, for comparable levels of cleanliness, glasses should be more resistant to failure. Another potential damage mechanism in polymers is thermal degradation around absorbing inclusions. The result would be an accelerating degradation with successive shots due to local carbonization.

We tested thin sheets of four polymeric materials. Sun Gain is a UV-stabilized polyester (similar to Mylar) made by 3M, available in sheets bearing broadband AR coatings, and reportedly³⁴ consists of a layer of acicular $\text{Al}(\text{OH})_3$ crystals. Teflon FEP is a

copolymer of tetrafluoroethylene and hexafluoropropylene and is made by Dupont. Extra-clean Mylar was made by Dupont. Unstretched polyvinyl alcohol was provided by Polaroid Corp. (this material is normally stretched to make it birefringent and is sandwiched between plastic sheets for use as quarterwave plates). Teflon FEP is available either with a surface treatment (necessary if the material is to be cemented to another substance) or without the treatment. Teflon FEP has a single-surface specular reflectance of $\sim 2\%$ and a rather nonuniform birefringence.

The results of our damage testing are given in Table 7-4. Teflon FEP might be usable at all three Nova wavelengths as a debris shield. At values equal to and just above threshold, the damage observed in Teflon FEP is minor. Sun Gain is potentially interesting at 1ω .

KDP Harmonic-Generation Cells with HC-56 Index-Matching Liquid

In the harmonic-generation cells for Nova and Novette, the internal surfaces of KDP crystals and windows are covered by thin layers of a halocarbon index-matching liquid, HC-56. In January 1982, a general prototype third-harmonic cell was assembled and installed on the Nova laser. The cell consisted of a borosilicate input window, the outer surface of which was AR

coated for incoming 1ω light; an 18-mm-thick Type II KDP doubling crystal; a 9-mm-thick Type II KDP tripling crystal; and a bare fused-silica output window. All internal surfaces were covered by very thin layers of HC-56. The cell produced a spatially uniform 3ω beam and survived for several hundred shots when used with 0.2-ns, 1ω input pulses with fluences of from 0.4 to 0.6 J/cm². However, small bubbles were created in the HC-56 layers when the cell was exposed to 1.0-ns input pulses at 1 to 2 J/cm².

We then tested 1-cm thick cuvettes filled with several liquids, including HC-56, water, and other index matching liquids. We observed that 1ω irradiation of all the liquids produced bubbles both at the cell windows and at randomly distributed sites in the bulk of the liquid. At fluences slightly above that necessary to produce bubbles, carbon was produced by decomposition of halocarbon liquids. However, carefully cleaned cuvettes filled with HC-56 filtered through 500- or 1000-Å filters were usually undamaged when irradiated by 1-ns, 1ω pulses at fluences up to 10 J/cm². The experiments suggested that the liquid boiled and decomposed at local sites when light was absorbed by contaminants on the windows and in the liquids.

A test program requiring more than 100 damage tests was begun to determine acceptable techniques for cleaning the cell bodies, windows, and crystals; for filtering

Fig. 7-36. Damage thresholds with 20-ns, 0.25- μ m pulses on four-layer AR coatings.

the liquid; and for maintaining the required cleanliness during cell assembly. The Novette second-harmonic cells and two associated small certification cells were assembled by procedures developed during our tests. The certification cells were damage tested with 1-ns, 1-w pulses in a collimated beam 1 cm in diameter and with the cells aligned to optimize conversion to the second harmonic. The threshold for randomly selected sites in either cell was 5 J/cm², a level that should ensure successful operation of the Novette cells.

Damage Studies at 0.25 μ m

Our principal objectives in damage studies at 0.25 μ m were to identify or develop AR and HR coatings that could withstand 5 J/cm² and to measure thresholds in UV-window materials. All planned studies of damage induced by 20-ns, 0.25- μ m KrF laser pulses have been completed. The thin-film coatings discussed in this section were fabricated at OCLI by electron-beam evaporation.

Initial experiments at 0.25 μ m indicated that damage-resistant AR coatings could be fabricated by using the high-index/low-index combinations scandia/magnesium fluoride or scandia/silica or by using coating designs that required only low-index materials. Both approaches were evaluated.

Figure 7-36 shows the thresholds for 24 four-layer scandia/magnesium fluoride and

scandia/silica AR coatings on fused silica.²⁰ A halfwave-thick undercoat layer of either magnesium fluoride or silica was placed between some AR coatings and the substrate. Comparably large thresholds ($\sim$ 6 J/cm²) were observed in coatings of either design with silica undercoats.

We then tested AR coatings of three designs requiring only low-index materials: lanthanum fluoride/silica (four layers), thorium fluoride/silica (six layers), and silica/magnesium fluoride (three layers). We also tested additional coatings made with scandia. Damage thresholds measured in this experiment are shown in Fig. 7-37. The three-layer silica/magnesium fluoride coatings and the four-layer scandia coatings had the largest threshold, $\sim$ 6 J/cm². Thresholds for films of both designs using scandia were the same as those observed in the first experiments (see Fig. 7-36). Therefore, our testing program identified three coating designs that provide AR films with thresholds exceeding our goal of 5 J/cm².

We conducted two major experiments with HR coatings. The first was a study of the relationships between damage thresholds and the deposition and design parameters of scandia/magnesium fluoride HR coatings.²⁹ We tested 12 variations of HR coatings made by depositing coatings of three designs (19 layers without overcoat, 19 layers plus a halfwave magnesium fluoride overcoat, and 19 layers plus a halfwave silica overcoat) on substrates made of two materials (BK-7 and fused

Fig. 7-37. Damage thresholds with 20-ns, 0.25- μ m pulses on four-layer AR coatings.

Fig. 7-38. Damage threshold with 20-ns, 0.25- μ m pulses on HR coatings.

Fig. 7-39. Damage thresholds with 20-ns, 0.25- μ m pulses on halfwave-thick films of 15 coating materials.

silica) at two substrate temperatures (150 and 250°C). A few 31-layer coatings were also tested. Thresholds measured on the subset of films deposited on BK-7 substrates are shown in Fig. 7-38. The presence or absence of an overcoat layer was the parameter that most strongly affected thresholds. Overcoated HR coatings had thresholds of from 5 to 7 J/cm². Comparable results were obtained in the films deposited on silica substrates.

To identify other materials for use in 0.25- μ m HR coatings, we tested halfwave-thick films of 15 materials, and HR coatings made from 13 combinations of these materials.¹⁵ Thresholds for the single-layer films are shown in Fig. 7-39. With a few exceptions, thresholds were greatest for films of low-index fluoride materials, which have UV band edges well below 0.25 μ m (low absorption). Thresholds were lowest for films of high-index oxide materials, which tend to be absorptive because their UV band edges are only slightly below 0.25 μ m. Thresholds of the reflectors in Fig. 7-40 were generally greater than those of high-index, single-layer films and less than those of low-index films. Further, the threshold rankings of the various HR coatings could not be predicted from the threshold rankings of constituent materials. Therefore, neither magnitudes nor rankings of thresholds of single-layer films were particularly instructive in selection of materials for HR

coatings. Our study of reflectors did, however, identify the combination of magnesium oxide/lithium fluoride as a candidate for production of damage-resistant 0.25- μm HR films. We believe, therefore, that HR coatings made from several combinations of materials will have thresholds exceeding the program goal.

Finally, we tested bare, polished samples of several materials that are potentially useful at 0.25 μm (Ref. 36). In press-forged lithium fluoride samples (Honeywell Corporate Technology Center) and in single crystals of lithium fluoride (Harshaw Chemical Corp.), the median thresholds for rear-surface damage and for bulk-inclusion damage were comparable ($\sim 16 \text{ J/cm}^2$). All 19 samples withstood at least 10 J/cm^2 .

We tested four fused-silica windows. The safe operating fluence in fused silica was set by the threshold for rear-surface damage, which ranged from 5 to 9 J/cm^2 . Bulk damage was not a limiting problem in fused silica.

Four potassium dideuterium phosphate (KD^*P) crystals and one crystal each of potassium dihydrogen phosphate (KDP) and ammonium dideuterium phosphate (ADP) were tested. Safe operating fluence in these

crystals was set by the threshold for entrance-surface damage, which ranged from 3 to 6 J/cm^2 . Nonlinear absorption in the bulk of the crystals prevented damage to the exit surfaces. Therefore, for substrates made of lithium fluoride or fused silica, damage thresholds are greater than, or equal to, thresholds of the AR and HR coatings with the greatest resistance to damage, whereas thresholds for isomorphs of KDP are generally below coating thresholds.

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Fig. 7-40. Damage thresholds with 20-ns, 0.25- μm pulses on HR coatings made with 13 combinations of materials.

Nonlinear Optical Measurements

Introduction

In 1982, we conducted two major measurements of nonlinear optical properties in support of the ICFO movement towards short wavelengths. First, we expanded the study of nonlinear absorption and solarization in potential UV-transmitting glasses. Our work led to a numerical model for nonlinear loss in window materials that incorporated four newly measured parameters.

Second, we used two methods to measure the nonlinear refractive index at 3ω in fused silica. Because time-resolved interferometry proved difficult and uncertain at 3ω , we developed the technique of direct measurement of small-scale self-focusing. We found that, in fused silica, the nonlinear refractive index, n_2 , is the same at both 3ω and 1ω .

Author: W. L. Smith

Fig. 7-41. Two-photon-induced solarization.

Two-Photon Absorption and Solarization in UV Glasses

During 1981, we showed³⁷ that fused silica would remain loss-free under high-intensity illumination at 3ω . However, fused silica is not amenable to the phase-separated, gradient-index treatment that provides Λ -surfaces with high damage resistance. In addition, fused silica is several times more expensive than borosilicate glasses. Hence, our goal in 1982 was to investigate alternatives to fused silica for optics that must transmit high-intensity 3ω laser pulses.

The processes involved in two-photon-induced solarization are shown in Fig. 7-41. Two-photon-absorption (2PA) transitions deposit energy in the solid, thereby populating excited electronic states in the glass. During the lifetime of the 2PA-induced state, it may absorb an additional photon by excited-state absorption (ESA). Upon de-excitation, a fraction of the excited population relaxes into midgap energy states rather than into the lower-energy initial state. The midgap states in glasses are long-lived ($\geq$ months) at room temperature and give rise to absorption bands in the one-photon spectrum, extending well into the visible. The induced absorption resulting from prior exposure to light is solarization.

Equations (19) through (23) express the attenuation of a laser wave of intensity $I(z)$ by the above processes. In addition to the usual linear (one-photon) attenuation term, we identify four other terms, each described by a coupling coefficient— σ_1 , β , σ_2 , or σ_{21} —so that

$$\frac{dI(z)}{dz} = -\alpha I(z) \quad (19)$$

for one-photon absorption,

$$-\sigma_1 \frac{\alpha}{h\nu} \int_0^z I(t') dt' dz \quad (20)$$

for one-photon solarization,

$$-\beta I(z) dz \quad (21)$$

for two-photon absorption,

$$-\sigma_2 \frac{\beta}{2h\omega} \int_{-\infty}^{\infty} I^2(t') dt' dz \quad (22)$$

for ESA following 2PA, and

$$-\sigma_2 \frac{\beta}{2h\omega} \int_{-\infty}^{\infty} I^2(t') dt' dz \quad (23)$$

for two-photon solarization.

Our apparatus for measuring the five coefficients is shown in Fig. 7-42. The 1ω pulses from the Cyclops laser were frequency tripled to provide up to 3 J of energy at 3ω . The beam quality at 3ω , shown in Fig. 7-43, was excellent. The beam was reduced (~ 10 mm ϕ diam) and directed along the beamline defined by crosswires X_2 and X_3 . Beam splitters S_1 , S_2 , and S_3 reflected energy to the incident- and transmitted-beam diagnostics (IBD and

TBD). The energy detectors (absorbing-glass-aperture calorimeter), spatial-shape detectors (vidicons and multiple-image cameras), and temporal-waveform detectors (apertured biplanar photodiodes) were located both before and after the sample. The IBD were located at a position of propagation equal to that for the input surface of the sample. The TBD, except for photodiodes, were located with lenses L_1 and L_2 in image planes of the output surface of the sample. Silicon diodes and the associated radiometer were used with the 3ω line of an argon laser to measure the passive transmission of the sample. The cw beam was also aligned to crosswires X_2 and X_3 . Apertures of 3-mm diameter were centered on the beamline in front of the calorimeter and all four diodes, so that only the central 3 mm of the pulses and cw argon beams

Fig. 7-42. Apparatus for measuring two-photon absorption and solarization at 3ω .

Fig. 7-43. Photographs of the Cyclops-beam spatial profile at 3ω .

Fig. 7-44. Nonlinear absorption for a 4.95-cm sample of BK-10 glass.

were used for the measurements. Over the central 3 mm, the pulsed beam was spatially flat to an excellent approximation.

A typical run began with an initial measurement of passive transmission. Next, a pulsed shot was made for which the incident intensity and instantaneous transmission were recorded. Finally, the passive transmission was remeasured. The sequence was repeated as the laser intensity was varied over a desired range or as a number of shots were accumulated at a fixed intensity.

In 1982, we recorded solarization data of two types. We first made time-integrated measurements of the energy transmission of 1.1-ns pulses through the samples, as in our earlier work with subnanosecond pulses. Typical results from our 3ω measurements for BK-10 glass are shown in Fig. 7-44. Figure 7-44(a) shows the intensity-dependent absorption of the pulses by the glass. Figure 7-44(b) shows a 20% reduction in passive transmission (i.e., solarization) caused by a series of pulses ranging from 0.7 to 2.3 GW/cm². As should be the case, solarization varies quadratically with laser intensity, as shown in Fig. 7-44(c). The major source of scatter in the data shown in Fig. 7-44(c) arises from the fact that the amount of solarization at a given pulse intensity depends on the number and intensity of pulses previously transmitted through the sample. That is, the solarization saturates.

From time-integrated data, we extracted 2PA coefficients, which were considerably larger than those we measured earlier with shorter pulses. The variation confirmed that ESA by the two-photon-created states is significant in UV glasses. We then modified our apparatus to record with 0.4-ns resolution the time-dependent transmission of 7-ns pulses.

Fig. 7-45. Procedure for extracting the 2PA coefficient, β , and the excited-state cross section, σ_2 , from time-resolved transmission data.

Sample	BK-10	Reflectivity	0.0422
Sample size (cm)	5.0	Alpha (1/cm)	0.0067
Sample steps	10	β (cm/TW)	5.5 (final)
Input fluence (J/cm ²)	4.1	σ_2 (cm ²)	3.2×10^{-18} (final)

Fig. 7-46. Incident, transmitted, and calculated waveforms from which β and σ_2 are extracted for BK-10 glass.

Nonlinear Optical Measurements

Our procedure, shown in Fig. 7-45, began with an oscillogram bearing the incident waveform (of approximately trapezoidal shape) and the transmitted waveform. The waveforms were digitized, entered into a

Table 7-5. Measured parameters of nonlinear absorption at three pulse lengths.

Glass	β (cm/TW)			$\sigma_2(10^{-18} \text{ cm}^2)$	
	0.7 ns ^a	1.4 ns ^a	7 ns ^a	7 ns ^b	
BK-10	4.5 ± 1.5	17	5.5 ± 1.5	3.2 ± 0.8	
BK-7	6.0 ± 3.0	40	30 ± 10	10.3 ± 2.5	
BK-3	—	—	1.3 ± 5	30 ± 7	
BSC3A	—	—	0.5 ± 0.2	212 ± 70	

^aPrevious measurements.

^b1982 measurements.

Fig. 7-47. Variation of effective ESA cross section with the 3ω laser pulse for four glasses.

Fig. 7-48. Procedure for extracting the two-photon-induced soliarization coefficient, σ_{21} .

computer, and overlapped on the time axis, as indicated in Fig. 7-45(b). We then used the input waveform and the known linear transmission of the sample as input for a computer program based on Eqs. (19) through (23). We searched for the β - σ_2 pair, which gave a computed-output waveform that agreed with the measured transmitted waveform. Figures 7-45(c) and 7-45(d) illustrate the fitting procedure. In Fig. 7-45(c), the value of σ_2 is too large by a factor of 2. In Fig. 7-45(d), there is poor agreement on the leading section of the pulse because β is too large by a factor of 2. However, the fit over the trailing section of the pulse is good because the product $\beta\sigma_2$ is well chosen. In Fig. 7-46, both β and σ_2 are properly chosen, and the overall fit is good.

Values for β and σ_2 are collected in Table 7-5 for BK-10, BK-7, BK-3, and BSC3A glasses. For BK-10 glass, we found that a single value of σ_2 could be used to fit shots of low and high fluence. For the other three glasses, a variation in σ_2 was required for a fixed β . Figure 7-47 shows the fluence dependence of the effective cross section, σ_2 .

The final parameter measured in our work was two-photon-induced soliarization, σ_{21} . We determined σ_{21} from plots, such as that shown in Fig. 7-48, of the observed change in linear attenuation coefficient ($\Delta\alpha$) vs the fluence-related integral (Ψ) from Eqs. (19) through (23).

The measured parameters³⁸ are collected in Table 7-6 along with the coefficient for one-photon-induced soliarization, σ_1 . Our materials were borosilicate glasses BK-10, BK-7, BSC3A (Hoya), and fused silica (Corning 7940). We conclude that none of the borosilicate glasses is an acceptable alternative to fused silica for high-power transmitting windows or focusing lenses for Nova. However, such glasses may still be used in components in beam diagnostics, which experience lower intensity. In addition to the materials already discussed, soliarization data for a number of other materials are listed in Table 7-7.

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Fig. 7-49. Experimental apparatus for time-resolved interferometric measurement of n_2 at 3ω .

Material	One-photon solarization, σ_1 (10^{-24} cm 2)	Two-photon absorption, β (cm/TW)	Two-photon-induced ESA, ^a σ_2 (10^{-18} cm 2)	Two-photon-induced solarization, σ_{21} (10^{-18} cm 2)
BK 10	1.8 ± 0.3	5.5 ± 1.5	3.2 ± 0.8	0.82 ± 0.20
BK 7	3.4 ± 0.5	30 ± 10	10.3 ± 2.5	1.2 ± 0.25
BK 3	—	15 ± 5	30 ± 7	0.74 ± 0.20
BSC3A	—	0.5 ± 0.2	212 ± 70	2.4 ± 0.6
Silica ^b	0	<0.5	0	0

^aAt 1 J/cm^2 fluence.

^bComing 7940.

Table 7-6. Summary of measured nonlinear absorption parameters at 3ω .

Nonlinear Refractive Index at 3ω in Silica

Time-Resolved Interferometric Measurement of n_2 . During 1982, we obtained the first measurements in the UV region³⁹ of the nonlinear refractive index, n_2 , in fused silica and BK-10 glass. Figure 7-49 shows our experimental layout, which was similar to the configuration of previous n_2 measurements. Table 7-8 shows the measured values of n_2 for three samples using time-resolved interferometry at 3ω with pulses obtained by third-harmonic generation in KDP crystals.

Glass³⁰ predicted that the value of n_2 at 3ω would be 1.5 times that of n_2 at 1ω in fused silica. However, based on 28 measured data points, we found that n_2 was no greater at 3ω than at 1ω . Unfortunately,

Borosilicates	Silica
Hoya ARG2	Hereaus Ultrasil
Schott BK-10	Hereaus Homosil
antisolarant	Hereaus Suprasil II
Fluorophosphates	Hereaus Optosil
FK51	Hereaus Optosil MS
FK54	Filter glasses
FCD10	UG-11
Schott FK54/2	TFF-1
Schott FK5	TFF-2
Schott FK5/11	Other materials
Schott K-11A/1	KDP crystal
Schott K-11A/2	Halocarbon 56 fluid
Schott FK5/8	NaCl crystal
	KI crystal
	H $_2$ O
	CR-39 polycarbonate plastic

spatial irregularities associated with the harmonic-conversion process limited beam quality and contributed a large uncertainty

Table 7-7. Additional materials tested for solarization.

to our results. Table 7-9 compares our results at 3ω to those from previous measurements at 1ω .

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Direct Measurement of Small-Scale Self-Focusing. To overcome the difficulties of the interferometric technique in the UV, we began a second approach for measuring n_2 . In 1974, Bliss et al.⁴³ demonstrated that spatial modulation deliberately imposed on a

Table 7-8. Measured values of n_2 at 3ω .

Sample	2-pass length (cm)	Number of measurements	n_2 (10^{-13} esu)
Dynasil 4000	30.8	9	2.2 ± 0.5
Suprasil II	20.0	—	—
Suprasil II	6.0	19	1.0 ± 0.85
BK-10 glass	20.0	9	0.6 ± 0.6

Table 7-9. Comparison of n_2 at 3ω and 1ω .

Material	n_2 at 3ω (10^{-13} esu)	n_2 at 1ω (10^{-13} esu)
Fused silica	$0.9 \pm 85\%$	0.95^a
BK-10 glass	$0.6 \pm 100\%$	0.49^b

^aValues from Ref. 41.

^bValues determined with the empirical method of Ref. 42.

laser pulse grew in a predictable fashion as the pulse traveled through a glass sample (see Fig. 7-50). Bliss et al. used a shearing-plate interferometer to generate modulation of a single spatial frequency. Although only a limited numerical analysis was available at the time to compute the expected nonlinear growth, the observed growth was consistent with their computation using the known value of n_2 at 1ω .

For applying the spatial modulation to the UV pulse, we investigated four different apodizers

- A filament (soft apodizer).
- A wire.
- A razor blade.
- Two knife edges crossed at 90° .

The two-dimensional diffraction pattern from the knife edges was the most successful in terms of enabling the imposed modulation to grow faster than the ambient noise modulation.

Our experimental layout for studying the modulation growth is shown in Fig. 7-49. The crossed-knife-edge apodizer was located 40 cm before the first surface of beam splitter S_1 . Multiple-image cameras (MIC) with 1Z film were used to record the spatial profiles for computer analysis.

Fig. 7-50. Nonlinear growth of spatial modulation on a laser pulse due to propagation through a glass sample.⁴³

Fig. 7-51. Vidicon spatial profile of the apodized beam.

Fig. 7-52. Fluence vs position without and with a silica sample.

Figure 7-51 shows the appearance of the apodized laser pulse both without a silica sample and after traveling through the 31-cm sample of Corning 7940 fused silica. Fast diodes recorded the 1.1-ns (FWHM) temporal waveform on each shot. Figure 7-51(a) shows the checkerboard pattern of peaks and valleys on the beam without silica formed by diffraction from the crossed-knife-edge apodizer. With silica, the peaks are enhanced by nonlinear growth, as shown in Fig. 7-51(b).

Several experimental checks were necessary to rule out beam-shape effects from sources other than n_2 in the sample. With the sample and apodizer removed, we ran shots to check that identical beam profiles were recorded at the two film locations. With the sample removed and the apodizer in place, we checked that plasma formation at the knife edges did not distort the beam.

Fig. 7-53. Exponential growth of a peak induced by fused silica.

With the sample in place and the apodizer removed, we documented the amount of modulation growth that occurred on the unapodized laser pulse. The latter was negligible compared to the growth that occurred on the apodized pulse.

Information regarding film calibration was recorded on sheets of film taken from the same box as the film bearing the n_2

data. The film was developed in the same batch of processing chemicals. The patterns of optical density were digitized with 57- by 57- μm resolution on a Photometrics Data Systems 1010A microdensitometer.

The conversion of the digitized film to peak fluence without vs with a sample followed normal procedures. The digitized film was carefully calibrated to provide relative fluence values for both without-sample and with-sample data. Using the energy into the calorimeter and its location with respect to the without-sample beam, absolute fluence values were determined, and total beam energies were fixed. The latter allowed conversion of the relative-fluence data to absolute fluence. In Fig. 7-52, we show cuts across the fluence profile of one pulse both without and with the sample. Data for the growth of one of the peaks are plotted in Fig. 7-53. These data exhibit the anticipated exponential growth induced by the nonlinearity of the sample.⁴⁴

Our computational model consisted of two parts. In the first part, MALAPROP^{45,46} generated a data base consisting of the results of laser profiles with a range of intensities propagating through the experimental setup. In the second part, a MATHSY⁴⁷ program integrated MALAPROP data for comparison with experimental data.

The MALAPROP model began with a spatially shaped intensity profile striking the crossed razor blades. For the path without the sample, the beam propagated to the film plane. Due to the relatively long propagation path, high frequencies were removed before propagation. The resulting profile showed the same spacing between peaks as the experimental data. For the path with the sample, the beam propagated to the entrance to the sample. Alternate steps of nonlinear phase adjustments and free-space propagation modeled propagation through the glass. As MALAPROP models single time slices, a series of results can be integrated to represent fluence measurements. However, for the path without the sample, the relative intensity at the film plane is independent of the initial intensity, and a single result is sufficient.

Figure 7-54 shows the four primary peaks induced by the razor blades for the without-sample path. Since the peaks

Fig. 7-54. MALAPROP model without a sample.

Fig. 7-55. MALAPROP model with self-focusing.

narrow with increasing distance from the blade, nonlinear effects are different for different locations. In this way, the experiment sweeps a range of source terms. Figure 7-55 shows the extreme self-focusing that is possible on the with-sample path. In the presence of nonlinear self-focusing, the innocuous peaks of Fig. 7-51 become sources for damage. (Our experiment did, in fact, damage the fused silica.) The peaks low on the exponential-growth curve are just visible.

The second part of our computational model consisted of converting instantaneous results into fluence values representing a particular value of n_2 . The experimental-beam pulse shape was almost indistinguishable from a 1.1-ns Gaussian pulse, and this shape was used for integration. Since MALAPROP results were based on an n_2 equal to the $1\text{-}\mu\text{m}$ value, the results were mapped to equivalent results for other values of n_2 by multiplying the intensity by the ratio of the $1\text{-}\mu\text{m}$ n_2 to the new value. Thus, the MATHSY portion of the analysis provided results for a range of intensities and gamma factors. Figure 7-56 shows the exponential growth in the MALAPROP model.

The analysis of our data, particularly of the error limits, remains unfinished. However, preliminary results indicate that experimental data can be successfully matched by calculations. Figure 7-57 shows the agreement for one peak. (The final hook on the curve results from the resolution limits of the model.) More importantly, for the prospects of ICF research at UV wavelengths, the agreement between experimental data and calculations does not require an increase in the value of n_2 for wavelengths down to 3ω .

A constant n_2 is currently being used in MALAPROP for simulations of 1ω , 2ω , and 3ω laser beams. Our final results await completion of refinements to the data reduction of the film and analysis of the sensitivity of the model to experimental errors.

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Fig. 7-56. Exponential growth of a peak in a MALAPROP computer model.

Fig. 7-57. Comparison of experimental results and model growth.

Optical Materials Research (OBES Programs)

Spectral and Polarization Hole-Burning in Nd:Glass Laser Amplifiers

Dopant ions in laser glass reside in a variety of environments and experience different perturbing local fields. In principle, we may use fluorescence-line-narrowing (FLN)

to measure energy levels and transition probabilities over the ensemble of local environments in glass, thereby providing the information necessary to predict the gain saturation for amplifier glass. In practice, for the technologically important ${}^4F_{3/2} \rightarrow {}^4I_{11/2}$ transition of Nd^{3+} , precise determination of the gain saturation is made difficult by site-to-site variations in homogeneous line widths, cross-section anisotropies, and nonradiative transition

rates, as well as the possibility of accidental coincidences of transitions between pairs of Stark states. Nevertheless, simple models that allow for the physical mechanisms of inhomogeneity provide the potential for optimization of glass composition with respect to energy extraction, without the need for construction and testing of actual amplifiers.

During 1982, we developed such a model to predict the saturation behavior of Nd:glass short-pulse amplifiers. Spectral inhomogeneity is modeled by assuming that the broadened fluorescence line of the laser transition is the result of a superposition of Lorentzian lines of equal strength and homogeneous line width ($\Delta\nu_h$), having center frequencies distributed about the observed fluorescence peak with a Gaussian probability of width $\Delta\nu_{ih}$. The polarization inhomogeneity treats the ions as an isotropically oriented collection of anisotropic oscillators, with cross sections σ_s and σ_p for light polarized perpendicular and parallel, respectively, to an ion site's principal axis.

The effect of the calculated spectral inhomogeneity on energy extraction is shown in Fig. 7-58 for several values of $\Delta\nu_{ih}/\Delta\nu_h$. The extraction efficiency, η , is defined as the ratio of energy gain of the amplified beam to the extractable energy initially stored in the amplifier. Fluence (energy/area) is normalized by the saturation parameter, $h\nu/(\sigma K)$, where h is Planck's constant and K is the effective degeneracy factor. Note that all the curves approach an extraction efficiency of unity with increasing input fluence; this indicates that the inhomogeneity does not limit the extractable energy. To get the same fraction of stored energy out of an amplifier, however, we must operate at higher fluence levels—an important consideration when component damage is the limiting design factor.

The effect of cross-section anisotropy on extraction efficiency for a plane-polarized input beam is shown in Fig. 7-59. The behavior is similar to that of the spectral model, except that, in the polarization case, there is a limit to the decrease of the efficiency curve with increasing inhomogeneity. Whereas the width of a Gaussian distribution of Lorentzian-line center frequencies can conceptually be increased

Fig. 7-58. Extraction efficiency vs normalized input fluence for a homogeneous and a spectrally inhomogeneous amplifying medium having different homogeneous-to-inhomogeneous ratios, $\Delta\nu_h/\Delta\nu_{ih}$.

Fig. 7-59. Extraction efficiency vs normalized input fluence for amplifying media having isotropic (homogeneous) and anisotropic (σ_s/σ_p) stimulated-emission cross sections.

Fig. 7-60. Comparison of measured and calculated normalized output fluence vs input fluence for a model ED-2 laser-glass amplifier and a homogeneous (Frantz-Nodvik) amplifier.

indefinitely, a system with $\sigma_{\infty} = 0$ (a system of pure dipoles) is the most anisotropic case possible.

We have combined our models of spectral and polarization inhomogeneity to compare the computational results with experimental amplifier data. Scaling parameters required to predict a specific glass performance include the effective degeneracy factor, K , the average cross section, σ , the ratio of polarized cross sections, σ_p/σ_{np} , and the ratio of homogeneous line widths, $\Delta\nu_h/\Delta\nu_{inh}$.

- The effective degeneracy ratio is related to the average rate of relaxation of the lower laser level and can, in principle, have the range $1 < K \leq 2$ for Nd^{3+} in glass; experimental results⁴⁸ indicate $K \approx 1$ for extraction pulses of duration >1 ns in a wide variety of glass compositions.
- Judd-Ofelt analyses of absorption and emission spectra have provided values of σ for hundreds of glasses; variations in σ for more than a factor of 8 have been observed in various oxide, halide, and sulfide glasses.
- By the end of 1982, average values of σ_p/σ_{np} had been determined by a polarized resonant FLN experiment⁴⁹ for two silicate glasses (ED-2 and LG-650); the ratio was 0.42 ± 0.05 for both glasses, within experimental error.
- The homogeneous line width, $\Delta\nu_h$, has been measured⁵⁰ for a single host glass, ED-2; variations of other, more easily measured, transitions have shown that $\Delta\nu_h$ can vary with the host composition by a factor of 4 (Ref. 51).

Because of the multistate character of the relevant Nd^{3+} transition, a precise value for $\Delta\nu_h$ is difficult to glean directly from FLN experimental results. FLN results do indicate that the inhomogeneous broadening is approximately equal to the average Stark splitting of states of the lower laser level. We have used this observation to obtain a value for $\Delta\nu_h$.

Figure 7-60 compares measured and calculated output vs input fluence behavior, using parameters appropriate for ED-2. The experimental data points⁵⁰ are scaled by $h\nu/(\sigma K) = 6.9 \text{ J/cm}^2$. The excellent experimental agreement with our model contrasts with that of the homogeneous Frantz-

Nodvik treatment, which overestimates the output fluence by more than 12% for inputs greater than $0.5 h\nu/\sigma K$.

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Stimulated-Emission Cross Sections of Rare Earths in Glass

One of the most important parameters for laser action is the stimulated-emission cross section, σ , of dopant ions in laser glass. The value of σ for paramagnetic-ion lasers depends upon the chemical composition of the host glass. Systematic studies of rare-earth ions in glasses conducted during the past decade have demonstrated that large variations in σ can be achieved by varying the glass-forming anions and cations. This knowledge has provided the laser designer with a much wider range of parameters and makes possible the tailoring of σ for specific applications.⁵²

Figure 7-61 summarizes the general types of known inorganic glass-formers. These are grouped by the glass-former anion into oxides, halides, and chalcogenides. Simple glass-former cations are denoted by a circle; conditional glass-formers—those that require the presence of one or more compounds to form a glass—are denoted by a half-circle. The color of the circle indicates the possible glass-former anions.

The compositional space of inorganic glass-forming systems is obviously vast and has by no means been explored exhaustively. We have, however, investigated a sufficiently large number of selected glasses that we can broadly define the relative ranges of stimulated-emission cross sections that are possible for two trivalent rare-earth ions: neodymium and ytterbium.

Neodymium. The stimulated-emission cross section of a transition between states J and J' is given by⁵³

$$\sigma = \left[\frac{8\pi^3 c^2}{3hc(2J+1)} \right] \left[\frac{(n^2 + 2)^2}{n} \right] \times \left(\frac{\lambda_p}{\Delta\lambda_{eff}} \right) S(J, J') \quad (24)$$

where n is the refractive index, λ_p is the wavelength of the emission peak, S is the line strength, and $\Delta\lambda_{\text{eff}}$ is the effective line width of the emission. Since the emission bands of rare earths in glass are usually asymmetric and consist of inhomogeneously broadened transitions between several Stark levels, we obtained the effective

line width in Eq. (24) by integrating over the band and dividing by the peak intensity.

Table 7-10 summarizes the ranges of effective stimulated-emission cross sections for the $^4F_{3/2} \rightarrow ^4I_{11/2}$ transition of Nd^{3+} observed for different glass types; also given are the associated ranges for n , $\Delta\lambda_{\text{eff}}$, λ_p , and

Table 7-10. Range of spectroscopic properties for the $^4I_{11/2}$ transition of Nd^{3+} in different glasses at 295 K.

Glass	Refractive index	Cross section, σ (pm^2)	Wavelength, λ_p (nm)	Line width, $\Delta\lambda_{\text{eff}}$ (nm)	Radiative lifetime, T_R (μs)	Ref.
Oxides						
Silicate	1.46-1.75	0.9-3.6	1057-1088	34-55	170-1090	54,55
Germanate	1.61-1.71	1.7-2.5	1060-1063	36-43	300-460	54
Tellurite	2.0-2.1	3.0-5.1	1056-1063	26-31	140-240	54,56
Phosphate	1.49-1.63	2.0-4.8	1052-1057	22-35	280-530	54,57
Borate	1.51-1.69	2.1-3.2	1054-1062	34-38	270-450	54
Halides						
Fluoroberyllate	1.28-1.38	1.6-4.0	1046-1050	19-29	460-1030	54,58
Fluorozirconate	1.52-1.56	2.9-3.0	1049	26-27	430-450	59
Fluorohafnate	1.51	2.6	1048	26	520	54
Fluoroaluminate	1.41-1.48	2.2-2.9	1049-1051	30-33	420-570	54
Chloride	1.67-2.06	6.0-6.3	1062-1064	19-20	180-220	52
Oxyhalides						
Fluorophosphate	1.41-1.56	2.2-4.3	1049-1056	27-34	310-570	54
Chlorophosphate	1.51-1.55	5.2-5.4	1055	22-23	290-300	60
Chalcogenides						
Sulfide	2.1-2.5	6.9-8.2	1075-1077	21-22	64-100	61,62
Oxysulfide	2.4	4.2	1075	27	92	62

Fig. 7-61. Inorganic glass-forming systems grouped by anion type.

IA	IIA	IIIB	IVB	VB	VIB	VIIIB	VII	IB	IIB	IIIA	IVA	VA	VIA	VIIA	O		
1 H															2 He		
3 Li	4 Be																
11 Na	12 Mg																
19 K	20 Ca	21 Sc	22 Ti	23 V	24 Cr	25 Mn	26 Fe	27 Co	28 Ni	29 Cu	30 Zn	31 Ga	32 Ge	33 As	34 Se	35 Br	36 Kr
37 Rb	38 Sr	39 Y	40 Zr	41 Nb	42 Mo	43 Tc	44 Ru	45 Rh	46 Pd	47 Ag	48 Cd	49 In	50 Sn	51 Sb	52 Te	53 I	54 Xe
55 Cs	56 Ba	57 La	72 Hf	73 Ta	74 W	75 Re	76 Os	77 Ir	78 Pt	79 Au	80 Hg	81 Tl	82 Pb	83 Bi	84 Po	85 At	86 Rn
87 Fr	88 Ra	89 Ac	104	105													

Oxide glasses

Halide glasses

Chalcogenide glasses

Glass-former

Conditional glass-former

58 Ce	59 Pr	60 Nd	61 Pm	62 Sm	63 Eu	64 Gd	65 Tb	66 Dy	67 Ho	68 Er	69 Tm	70 Yb	71 Lu
90 Th	91 Pa	92 U	93 Np	94 Pu	95 Am	96 Cm	97 Bk	98 Cf	99 Es	100 Fm	101 Md	102 No	103 Lr

the calculated radiative lifetime, τ_R . In the cases of silicate, phosphate, and fluoroberyllate glasses, we explored a large compositional space; for other glasses, we examined only a few compositions. Therefore, the values in Table 7-10 are representative, but may not be the most extreme values possible.

The ranges of cross sections are shown graphically in Fig. 7-62. For silicate glasses, σ varies by a factor of 4 with a change of host composition from an alkali silicate to a bismuth silicate having a high refractive index. The peak cross sections obtainable for phosphates are larger than for silicates. This is not because the Judd-Ofelt intensity parameters (Ω_i) are significantly larger, but because the line widths are narrower. Some tellurite glasses have even larger cross sections than phosphates, not because S or $\Delta\lambda_{eff}$ are more favorable, but because of their large refractive indexes.

As expected, the halide glasses have narrower line widths than the narrowest oxide glass (phosphates) in Table 7-10. For the fluoroberyllate glasses, this does not result in larger cross sections than for phosphate glasses because the Ω_i parameters are smaller. Chloride glasses combine a small $\Delta\lambda_{eff}$ with a larger Ω_i and a high refractive index to yield values for σ larger than those for any oxide glass. The high-index, narrow-line-width, covalently bonded chalcogenide glasses result in the largest Nd^{3+} cross sections observed thus far.

The above results are for the principal lasing transition, ${}^2F_{3/2} \rightarrow {}^1I_{1/2}$, but a similar trend with composition has been observed for the ${}^4F_{3/2} \rightarrow {}^4I_{13/2}$ transition for a few glasses.^{26,61,63} Note that in all cases, the above results are the stimulated-emission cross sections, not the net cross sections that include contributions due to losses such as excited-state absorption (net cross sections also vary with glass composition).

Ytterbium. We investigated the variation of the effective stimulated-emission cross section of Yb^{3+} with glass composition, using absorption and emission measurements taken for 41 different oxide, fluoride, and oxyfluoride host glasses. Whereas silicate, phosphate, and fluoroberyllate glasses have been investigated extensively as hosts for Nd^{3+} , borate glasses received less attention because of the low quantum efficiency and

rapid decay of the Nd^{3+} fluorescence in these glasses—features that are deleterious to the efficiency of flashlamp-pumped lasers. In the case of Yb^{3+} , nonradiative decay from the ${}^2F_{5/2}$ state by multiphonon processes is negligible even in borate glasses. Therefore, we have emphasized borate glasses in these studies.

We also made measurements of selected silicate and phosphate glasses to compare the behavior of the spectroscopic properties of Yb^{3+} and Nd^{3+} in these hosts. Since large amounts of fluorine can be added to several glasses, we included some mixed-anion fluoroborate and fluorophosphate glasses to complete the survey. Details of the glass compositions and preparation procedures (glasses were melted at the National Bureau of Standards) are given in Ref. 64.

Due to the low symmetry at rare-earth sites in glass, the degeneracy of the ${}^2F_{5/2}$ and ${}^2F_{7/2}$ levels is completely removed, except for Kramer's degeneracy; thus, the optical spectra consist of transitions between the four states of ${}^2F_{7/2}$ and the three states of ${}^2F_{5/2}$. As evident from the fluorescence spectra in Fig. 7-63, only the transition between the lowest Stark levels of the ${}^2F_{7/2}$ and ${}^2F_{5/2}$ manifolds at ~ 974 to 978 nm (the 1-1 transition) is clearly resolved in all glasses. Inhomogeneous broadening and vibronic transitions combine to reduce the resolvability of the other transitions.

Fig. 7-62. Ranges of peak stimulated-emission cross sections reported for the ${}^4F_{3/2} \rightarrow {}^4I_{13/2}$ transition of Nd^{3+} in glasses at 295 K.

We calculated σ from the corresponding absorption spectrum via the relation

$$\sigma_p = \frac{4\Sigma}{3\Delta\lambda_{\text{eff}}} \quad (25)$$

where σ_p is the peak cross section for the 1-1 line; Σ is the integrated absorption cross section, $\int \sigma_{3\lambda} d\lambda$, of the ${}^2F_{7/2} \rightarrow {}^2F_{5/2}$ transition; and $\Delta\lambda_{\text{eff}}$ is the effective emission line width obtained by integrating over the ${}^2F_{5/2} \rightarrow {}^2F_{7/2}$ fluorescence band and normalizing the peak intensity to unity.

Low-threshold, efficient lasing of Yb^{3+} occurs for transitions terminating on higher-

lying Stark levels of ${}^2F_{7/2}$ or on vibronic states and not for the intense zero-phonon 1-1 transition, which forms a three-level lasing scheme. The cross section, σ_s , of the broad secondary peak in the fluorescence spectrum (which occurs in the region of 0.99 to 1.02 μm for all the glasses we examined) was determined from the fluorescence intensity, I , and the relation⁶⁵

$$\frac{\sigma_s}{\sigma_p} = \left(\frac{n_p}{n_s} \right)^2 \left(\frac{\lambda_s}{\lambda_p} \right)^5 \left[\frac{|dl/d\lambda_s|}{|dl/d\lambda_p|} \right] \quad (26)$$

where λ_p and n_p are the wavelengths and refractive indexes at the primary and secondary peaks in the fluorescence spectra. The radiative lifetime, τ_R , of the ${}^2F_{5/2}$ level was calculated from⁵⁵

$$\tau_R = \frac{3\bar{\lambda}_d}{32\pi c n^2 \Sigma} \quad (27)$$

Although there is small difference in the mean wavelength for absorption and emission at finite temperatures, we simply used the wavelength of the ${}^2F_{7/2}(1) \rightarrow {}^2F_{5/2}(1)$ transition for λ throughout.

The variation of σ with glass composition is illustrated in Table 7-11, which summarizes the effects on Yb-doped borate glass of changing the alkali species M in a series of alkali diborates ($\text{M}_2\text{O} \cdot 2\text{B}_2\text{O}_3$).

A summary of the cross sections determined thus far for different glass types is presented in Fig. 7-64; glass compositions and cross-section values are tabulated in Ref. 64. The effective peak cross sections for transitions to Stark levels above the ground state range from 0.24 to 0.84 pm^2 . The largest values occur in borate and phosphate glasses; the smallest values occur in silicate and low-refractive-index fluoride glasses. Radiative lifetimes calculated from integrated absorption spectra range from 0.6 to 2.7 ms.

As discussed in Ref. 64, systematic variations in spectroscopic properties with changes in modifier ions can be used to tailor stimulated-emission cross sections and fluorescence lifetimes in Yb^{3+} . As noted above, when the category of inorganic glasses is expanded to include the heavier-anion halide glasses and chalcogenide glasses, we found even larger values of

Fig. 7-63. Comparison of the ${}^2F_{5/2} \rightarrow {}^2F_{7/2}$ fluorescence spectra of Yb^{3+} in Ba-La borate, Na-Ba-Zn silicate, and Li-Ca-Zn phosphate glasses at 295 K.

Fig. 7-64. Ranges of stimulated-emission cross sections for the secondary peak of the ${}^2F_{5/2} \rightarrow {}^2F_{7/2}$ transition of Yb^{3+} in glasses at 295 K. ►

Table 7-11. Variation of physical and spectroscopic properties of Yb^{3+} -doped borate glass at 295 K with changes in alkali modifier M. Composition (mol%): $66\text{B}_2\text{O}_3 \cdot 33\text{M}_2\text{O} \cdot 1\text{Yb}_2\text{O}_3$. ▼

Alkali ion M	Lithium	Sodium	Potassium	Rubidium
n_d	1.567	1.522	1.504	1.508
ν	66	66	65	67
n_2 ($10^{-20} \text{ m}^2/\text{W}$)	3.4	3.1	3.0	2.9
d (g/cm^3)	2.41	2.47	2.38	2.91
${}^2F_{5/2} \rightarrow {}^2F_{7/2}$ fluorescence				
σ (pm^2)	0.70	0.52	0.37	0.33
λ_s (nm)	1010	1018	1010	1013
$\Delta\lambda_{\text{eff}}$ (nm)	49	46	45	50
τ_R (ms)	0.82	1.28	1.72	1.87

Nd^{3+} cross sections. Therefore the stimulated-emission cross sections reported here are not necessarily the extreme values possible for Yb^{3+} .

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Molecular-Dynamics Simulations of Glass Structure

Molecular-dynamics simulations have proved to be a powerful technique for investigating local structure and interactions at dopant-ion sites in laser glass.⁶⁶⁻⁷² In this article, we present calculations of vibrational spectra and glass-relaxation phenomena, after comparing simulated glass structures with various experimental measurements of the structure of rare-earth-doped glasses.

Comparison of Simulated Glass Structure with Experiment. To compare our simulated structures of BeF_2 glass with average ion separations and coordination numbers obtained using x-ray^{73,74} and neutron diffraction,⁷⁵ we very slowly relaxed a system of 130 BeF_2 . The repulsive parameters used in the ion-ion potential were $A(\text{Be-F}) = 25.5$, $A(\text{F-F}) = 31.2$, and $A(\text{Be-Be}) = 15.6$, with a constant softness factor $\sigma = 3.45$. A continuous random network of BeF_4 tetrahedra joined at corners was formed. The Be coordination number is 4.0 at 2.0 Å, but increases to 4.04 at 2.4 Å, indicating the presence of a few "defects" in the structure.⁶⁹

Table 7-12 compares the peaks of the calculated pair-correlation functions and average coordination numbers (averaged over five configurations involving a total of 650 Be ions) with x-ray diffraction data. The overall agreement is good, demonstrating that simulations of simple glasses using systems of a few hundred particles and slow relaxation are capable of yielding physically meaningful structures.

Although we have tested our simulated site to site structural variations using optical and FLN spectra of Eu^{3+} in BeF_2 glass, it is difficult to relate optical spectra and local glass structure directly because of the approximate treatment of the local field. We

have therefore supplemented the FLN studies with tests based on measurements of

- The extended x-ray absorption fine structure (EXAFS) spectrum of Nd^{3+} .
- Neutron-diffraction studies of glasses containing mixtures of Dy isotopes selected to yield information about the rare-earth environment.⁷⁷

Both methods yield information only about distributions of ion-pair distances and average coordination numbers.

The results of the neutron-diffraction studies, using a glass of (mol%) 60 BeF_2 -29 NaF -11 DyF_3 , are summarized in Table 7-13 and compared with predictions for a simulated glass consisting of five configurations of a 380-particle system of the same composition. In this table, the numbers in parentheses after the simulated coordination numbers are the radii of the coordination spheres, in Å. The simulation coordination number of Dy by F, for $r = 3.0$ Å, is slightly higher than that determined experimentally. The EXAFS tests, using a BeF_2 glass containing 4% NdF_3 , yielded an average Nd-F separation of 2.43 ± 0.09 Å and a coordination number of 7.1; these values are in good agreement with the simulated values of 2.35 Å and 7.2, respectively.

Beryllium Sites in Complex Glasses.

For simple BeF_2 glass, simulations predict

i	j	Average separation (Å)		Coordination number	
		Measured	Simulation	Measured	Simulation
Be	F	1.554 ± 0.004	1.55	3.8 ± 0.3	4.0
F	F	2.537 ± 0.004	2.47	5.7 ± 0.3	6.6
Be	Be	3.037 ± 0.005	3.05	3.8 ± 0.3	4.1

Table 7-12. Comparison ▲ of simulated and measured (x-ray diffraction⁷⁴) average separations and coordination numbers for BeF_2 glass.

Table 7-13. Comparison of simulated and measured (neutron diffraction⁷⁷) average separations and coordination numbers for 80 BeF_2 -40 NaF -15 DyF_3 glass. ▼

i	j	Average separation (Å)		Coordination number	
		Measured	Simulation	Measured	Simulation ^a
Be	F(I)	2.31	2.32	7.7	8.5 (3.0)
Dy	F(II)	4.65	4.6	—	32 (3.0-5.4)
Dy	Be	3.75	3.7	—	7.7 (4.5)
Dy	Na	—	3.8	—	5.6 (5.3)
Dy	Dy	—	4.2	—	2.7 (5.9)

^aNumbers in parentheses are radii of the coordination spheres in Å.

fourfold coordination of Be by F within 2.0 Å. For more complex fluoroberyllate glass compositions, simulations predict that the coordination number at 2 Å is significantly greater than 4. The simulations are illustrated in Fig. 7-65, which compares the Be-F radial distribution function, $g(r)$, and average coordination number from five simulated configurations each of 80BeF₂-40NaF-15DyF₃ and 130BeF₂. The peak of $g(r)$ for both glasses is at ~ 1.55 Å, but the distribution of Be-F distances is broader and more asymmetric for the latter and approaches zero only at $r \approx 2.4$ Å. The associated average coordination number, $n(r)$, becomes greater than 4 at $r > 1.8$ Å.

The definition of a coordination shell for Be in multicomponent glasses is not as straightforward as for simple BeF₂. There are a number of BeF₃ complexes with four short and one long Be-F distance; these complexes might be considered perturbed tetrahedra. There are also a significant number of complexes for which all five Be-F distances are less than 2 Å, and no such simple physical interpretation is possible. The number of nontetrahedrally coordinated Be complexes in well-relaxed molecular-dynamics computer glasses depends on the presence of modifiers and increases with the concentration and valence of the modifier.

Whether such sites really exist in laboratory glasses or are a consequence of the rapid quenching and other approximations

involved in making computer glasses is subject to experimental verification. Neutron-diffraction studies of BeF₂, BeF₂-NaF, and BeF₂-NaF-DyF₃ glasses show an increase in the average Be coordination number (by ~ 0.2 F ions).²⁷ The predicted Be coordination number for the simulated BeF₂-NaF-DyF₃ glass is 4.3 at $r = 2.0$ Å. This value is slightly higher than the neutron diffraction result, but the latter was obtained by fitting the correlation function peak with a Gaussian distribution. The calculated pair correlation function is skewed, as shown in Fig. 7-65; therefore, an asymmetric distribution should be used to avoid truncating the contributions at large r , where the average coordination number becomes greater than 4.

Another potential test of the presence of nontetrahedrally coordinated Be sites is the observation of a change in the quadrupole broadening of the ^{9}Be ($I = 3/2$) NMR spectrum. If Be ions are perfectly tetrahedrally coordinated by nearest-neighbor F, then by symmetry there will be no electric field gradient (efg) at the Be site and no quadrupole splitting of the nuclear resonance due to these ions. To search for this effect, Bray and Dell²⁸ have compared the NMR spectra of a BeF₂ glass and a 65BeF₂-35NaF glass. We have simulated the ^{9}Be -modified glass using a system of 95BeF₂-2NaF; we find that about 25% of the Be sites have odd coordination. We also find that there is a distribution of F-Be-F angles for both simple and modified BeF₂ glasses, indicating deviations from perfect fourfold tetrahedral coordination. Because of the rapid quenching and lack of angle-dependent bonding in the potential used for the simulations, however, the predicted geometric arrangements of ions about the different Be ions are not sufficiently meaningful to calculate distributions of dipolar and quadrupolar splittings for comparison with the observed NMR spectra.

Vibrational Spectra of Simulated Glasses.

During 1982, we employed our glass structure simulation to perform ab initio calculations of the vibrational density of states and of the spectra of infrared absorption, Raman scattering, and inelastic neutron scattering. The force constants required for these calculations are obtained directly from the interaction potential used in the

Fig. 7-65. Comparison of the radial distribution function and coordination number for two computer-simulated Be glasses.

structure simulations. The spectra are obtained with no free parameters, apart from an overall multiplicative constant.

Vibrational spectra can be calculated in two ways. The first method is direct diagonalization of the dynamical matrix for the glass and using the resulting eigenfunctions and frequencies to evaluate the formulas for the scattering intensities. The second method is to evaluate appropriate auto-correlation functions computed directly from a low-temperature molecular-dynamics simulation where the potential is harmonic. We used both methods and found they gave the same result.

The results of the calculations are summarized and compared with experimental results⁸⁰ in Fig. 7-66 which also includes the calculated density of states (DOS). The agreement with experiment is generally fair. The shapes of the calculated spectra are similar to those of the experimental spectra, but the calculated band maxima are from 100 to 150 cm^{-1} too small. The reasons for this discrepancy are not known, but it is not due to the small size of the sample.

Thus far, we have treated only the case of simple BeF_2 glass, for which experimental data are available. The details of the lattice dynamics of harmonic systems and calculations of vibrational spectra have been described⁸⁰⁻⁸²; the explicit expressions used here are to be published in Ref. 83.

Phenomenology of Glass Relaxation.

During the course of our simulations of rare-earth-doped glasses, it became apparent that the simulated glasses were more disordered than laboratory-produced glasses. The discrepancy is due to the high quenching rate of computer-simulated glasses, necessitated by the limitations on computer time and the speed of programs. While we could reproduce qualitative pre-

dictions and trends in the optical properties of rare-earth ions in glass, quantitative agreement with experiment was elusive.

It became clear that extrapolation of the results from computer-simulated glasses to real glasses requires knowledge of glass relaxation. With the support of DOE's Office of Basic Energy Sciences (OBES), we have compiled extensive information on glass relaxation,⁸³ developing the following main points

- The phenomenology of the glass transition is consistent with the idea that glass is a frozen liquid.
- The physical properties of any particular sample of glass depend on its thermal and mechanical history and require more than a one-parameter description.
- The relaxation process is governed by a broad distribution of relaxation times.
- Immediately following a temperature jump, the apparent activation energy for the relaxation time is about half that for the equilibrium fluid (the relaxation time increases with time during the relaxation).
- Relaxation at temperatures below the glass transition temperature, T_g , can be explained by extrapolation of relaxation observed above T_g and there are anomalous excitations at temperatures much below T_g giving rise to unusual thermal, acoustic, dielectric, and optical properties.

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Thermal-Lensing Studies: Measurement of Two-Photon-Induced Color Centers in Borosilicate Glass

During 1982, we used thermal lensing, a sensitive calorimetric technique for

Fig. 7-66. Comparison of simulated and measured (a) reduced inelastic neutron-scattering cross section, (b) infrared absorption, and (c) reduced Raman-scattering intensity.

measuring small optical absorbances, to measure nonlinear absorption in a borosilicate glass (BK-7) at 2ω ($0.53 \mu\text{m}$). BK-7 glass is generally considered a suitable material for transmitting high-intensity 2ω laser beams. Experiments at 2ω , however, show that borosilicate glasses such as BK-7, BK-10, and BK-3 are unsuited for high-intensity applications due to sizable two-photon absorption and UV-induced linear absorption (solarization).⁸⁴ We have now observed solarization in BK-7 at 2ω , for one sample irradiance of several hundred to one thousand laser pulses at intensities of $>1.5 \text{ GW/cm}^2$.

Figure 7-67 shows the experimental configuration and apparatus used for our thermal-lensing measurements. The laboratory itself is pictured in Fig. 7-68 and described in the 1981 Annual Report.⁸⁵

One of several data sets is illustrated in Fig. 7-69, where we plot the photo-induced linear absorption coefficient at four pump-fluence levels vs the number of laser shots into the sample. The computed peak intensity for these runs ranged from 1.5 to 3.8 GW/cm^2 . At the 10-Hz laser rate, each data run of 1.8×10^4 laser shots lasted 30 min. There is a slight offset in the computed value of the absorption coefficient at

zero laser shots because this coefficient actually corresponds to an "average" linear absorption induced by approximately the first 100 laser shots of the sampling period of each data run. At sufficiently high intensity ($>1.5 \text{ GW/cm}^2$), we observed that even the first laser shot into a "fresh" region of the sample showed a finite lensing signal; this signal may arise from residual linear absorption, two-photon absorption, or solarization induced by the first pump pulse.

From a log-log plot of the data, we found that the induced absorption depended quadratically on the laser intensity. This dependence is indicative of two-photon-induced color-center formation, which is energetically allowed. [As a test, we irradiated a sample of fused silica of equal thickness (1 cm) at an intensity of $\sim 4.0 \text{ GW/cm}^2$. As shown in Fig. 7-69, we observed no absorption for SiO_2 ; the outcome is expected, because the fundamental band gap is $>4\omega$.] Note that, at the higher intensity runs in Fig. 7-69, the absorption appears to saturate. The color-center band was stable at ambient temperature, but could be bleached at elevated temperatures.

The two-photon-induced solarization per shot at 2ω is very small after transmission of a few laser shots, and no discoloration is

Fig. 7-67. Apparatus for transient thermal-lensing measurements of nonlinear absorption and photo-induced color centers in materials.

evident. But after an accumulation of $\sim 10^4$ shots at intensities $> 3 \text{ GW/cm}^2$, a brown discoloration track was clearly visible against a white-light background. The spectrum of the photo-induced absorption is

similar to that observed in glasses due to UV and ionizing radiation.

The probability per shot for three-photon-induced color-center formation at intensities near the laser damage limit of

Fig. 7-68. Laboratory and apparatus used for thermal-lensing studies.

Fig. 7-69. Photo-induced color-center absorption coefficient at 2ω for BK-7 glass, measured by transient thermal lensing.

optical crystals and glasses is estimated to be lower than that for two-photon processes by $\sim 10^{-4}$. For laser applications requiring high repetition rates or long lifetimes, however, these higher-order processes may contribute significant losses unless corrected by materials development.

Transient thermal lensing, in addition to its usefulness for detecting multiphoton-induced color-center formation, can also be used to measure transient color centers, small absorption due to dilute impurities or excited-state absorption, fluorescence quantum efficiency, rates and yields of chemical reactions, and thermal properties of materials. In many of these applications, measurements of materials of known absorbance, introduced either in series or as substitutes for the sample, can be used to calibrate the absorption strength.

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Fig. 7-70. Experimental setup for measuring excited-state absorption in laser glass.

Fig. 7-71. Change in absorption induced by dye-laser excitation.

Excited-State Absorption in Neodymium Glass

At high excitation densities, absorption of flashlamp photons by the Nd^{3+} ions in the upper laser level could seriously impair amplifier efficiency. During 1982, we performed a direct measurement of excited-state-absorption cross sections in the spectral region of interest.

The experimental setup is shown schematically in Fig. 7-70. Typically, we fired three shots with both the dye laser and flashlamp, three with the flashlamp only, and two with the dye laser only, to measure the background at each wavelength. The region from 350 to 800 nm was covered in steps of 10 nm. Energy output of the dye laser was monitored on each shot and kept constant within 10% during a run. The electrical signals from the photomultiplier (PM) tubes were fed into a pair of gated integrators, and their outputs were plotted on a chart recorder. The change of transmission due to the excitation was measured and averaged for each wavelength.

Figure 7-71 plots the change in the exponential absorption coefficient thus obtained for a 1-wt% Nd-doped silicate glass, ED-2. Qualitatively similar results were obtained for a phosphate glass, LGH-8. The large positive signals occur at the wavelengths corresponding to strong absorption from the $^4I_{9/2}$ ground state. This "gain" is due to bleaching of the absorber and can be used to calculate the excited-state density. Approximately 20% of the Nd ions were excited to the upper level. The negative signals in the spectrum can be identified as absorption from the $^4F_{3/2}$ state of Nd^{3+} ions (the upper laser state).

Figure 7-72 shows the same data normalized to the known ground-state absorption cross section of the peak at ~ 750 nm. Note the good agreement with the ground-state absorption at 800 and 530 nm and the strong excited-state absorption at ~ 600 nm. The cross sections for excited-state absorption are clearly small enough here not to cause trouble in the existing Nd-glass lasers and excited-state populations.

Table 7-14 compares the experimentally observed absorption with line strengths calculated using the Judd-Ofelt model.⁸⁶ The data again were normalized to the ground-state-absorption line strength at 750 nm,

Fig. 7-72. Cross sections for the ground-state and the excited-state absorption.

Table 7-14. Comparison of predicted and observed excited-state absorption and line strength.

Initial state	Final state	Wavelength (nm)	Line strength (Judd-Ofelt)	Experiment
Ground	$^4F_{7/2}$	753	2.40	3.62
Ground	$^4S_{3/2}$	748	1.22	
Ground	$^4G_{5/2}$	590	5.06	
Ground	$^4G_{7/2}$	585	1.25	6.31
Ground	$^2K_{13/2}$	531	0.185	
Ground	$^4G_{7/2}$	530	0.500	
Ground	$^4G_{9/2}$	518	0.500	1.92
Ground	$^4D_{3/2}$	360	1.005	
Ground	$^4D_{5/2}$	357	0.412	
Ground	$^2I_{11/2}$	354	0.102	2.88
Ground	$^4D_{1/2}$	352	1.195	
Ground	$^2I_{13/2}$	346	0.165	
$^4F_{3/2}$	$^2F_{5/2}$	361	0.029	0.65 ± 0.25
$^4F_{3/2}$	$^2D_{5/2}$	437	0.0784	0.135
$^4F_{3/2}$	$^2H_{11/2}$	452	0.0311	
$^4F_{3/2}$	$^2D_{3/2}$	459	0.0253	
$^4F_{3/2}$	$^2L_{11/2}$	524		0.4 ± 0.25
$^4F_{3/2}$	$^4D_{7/2}$	533	0.4528	
$^4F_{3/2}$	$^2I_{13/2}$	543	0.0096	
$^4F_{3/2}$	$^2L_{15/2}$	567		1.4 ± 0.25
$^4F_{3/2}$	$^4D_{5/2}$	598	1.400	
$^4F_{3/2}$	$^4D_{3/2}$	604	0.481	

assuming for simplicity a constant line width for all transitions. Reasonable qualitative agreement can be seen between the model and ground-state absorption; excited-state absorption in the visible region is also in reasonable agreement. The discrepancy at ~ 360 nm, with the observed excited-state absorption an order of magnitude larger than expected, is explainable. First, this absorption band is in the vicinity of the absorption edge and overlaps possible emission from Ce^{3+} , which is present as an antisolarization agent. The absorption band at 360 nm also overlaps a region of strong ground-state absorption.

In conclusion, we have qualitatively measured the excited-state absorption in Nd glasses; its magnitude is not large enough to cause problems for existing laser systems.

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Laser System and Component Evaluation

Optimization of Multipass Systems

Introduction. The next generation of fusion lasers will operate at the multimegajoule level with pulse durations in the 10-ns range. High energy-storage and extraction efficiencies are necessary to reduce cost. To evaluate and optimize various architectures, a methodology to trade off different design features and to obtain a relative performance scaling is essential. Actual performance and costs will be obtained from more detailed beam-propagation analysis and costing of all components.

Performance-to-cost ratios of various architectures and options are typically well within a factor of 2. Thus, it is important to maintain high accuracy of system simulations while reducing the complexity of the various designs for better side-by-side comparisons. Difficulties arise in comparing systems that differ in dynamic range and isolation characteristics. Since we have not yet found a single figure of merit that will completely determine the quality of a

system, cost effectiveness, Q ($\text{J}/\$$), at a particular pulse duration remains our best measure. To avoid biased conclusions, it is important to compare only equivalent characteristics and side effects (e.g., risk and operating cost) while selecting for systematic evaluations an average laser-pulse duration and pulse shape that covers most experimental requirements. An optimum system is then usually simultaneously limited at the selected pulse duration by nonlinear effects, damage, and energy extraction.

Simulation of complete laser systems is complex and too slow for optimization procedures. However, it is possible to subdivide the system into components with weak interactions. Thus, we can optimize the amplifier on a per-disk basis without concern about the impact of other components. The only system impact on the amplifier can then be reduced to a specific criterion for either high power, high energy, or high overall efficiency. Amplifier design differs significantly for these various tasks. The high-power amplifiers for Shiva had thin, highly pumped disks. While a high-efficiency amplifier for a fusion drive requires thick disks pumped at low flashlamp energies. For megajoule lasers, high-energy-storage amplifiers are of primary interest.

Laser System Design. Our system optimization procedure is simplified by separating the cost of the dominating output stage, C_o , from that of the driver stages C_D , in the form

$$Q = \frac{E_o}{\sum C_i} = \frac{E_o}{C_o + (C_D/G_o)} \\ = \frac{E_o}{C_o} \left(1 - \frac{k}{G_o} \right) \quad (28)$$

where E_o is output energy, $k = Q/Q_D > 1$ is the ratio of cost effectiveness of the complete system Q relative to the driver $Q_D = E_D/C_D$, and G_o is saturated gain of the output stage.

Our approach eliminates the need for detailed design and simulation of the drive stages. Thus, we reduce both complexity and potential distortions for relative comparisons. The driver can be scaled in relation to the output stage because any improvements made on the output stage

should also apply for the drive stages. Hence, a simple k -weighted scaling with the required drive energy, E_D , is sufficient.

Amplifier-stage cost is best described by two components: amplifier cost, C_A ; and fixed cost, ΔC , associated with isolation, spatial filtering or beam expansion, and alignment. The amplifier cost is

$$C_A = \frac{I_s}{Q_s} + \frac{\Delta C}{\eta Q_s} \quad (29)$$

where I_s is stored energy, Q_s is the cost effectiveness of storing energy, ΔC is extracted energy, η is the fill factor, and η is extraction efficiency.

The cost effectiveness of a high-energy laser system can then be described by

$$Q = \frac{I_s}{C_{A0} + \Delta C_{A0}} \left(1 - \frac{k}{G_{A0}} \right) = \eta Q_{A0} \frac{1 - \Delta k / (G_{A0} - 1)}{1 + \Delta} \quad (30)$$

where $\Delta k = k - 1$ and $\Delta = \Delta C_{A0} / C_{A0} < 1$. Both correction terms vanish for isofluent, isoquality stages ($Q_{A0} \approx \text{constant}$) and low fixed costs. The laser system is then dominated by the cost effectiveness of extracting energy $Q_t = \eta Q_s$ which allows the separation of amplifier optimization, Q_s , from system optimization, Q .

By separately optimizing the amplifiers and the chain assembly, we reduce the number of independent variables from about 20 to 10 each and, hence, make the process more tractable. The most important variables for disk amplifiers operating at long pulse durations are

- Clear aperture
- Disk thickness.
- Nd doping.
- Stimulated-emission cross section.
- Internal ASE control.
- Fluorescence lifetime.
- Pump-pulse duration.
- Number of lamps.
- Bank energy.
- Specific costs for ME, EE, and optics.

The above variables lead to specific values for the cost effectiveness of storing energy, Q_s . A detailed analysis shows that the cost of storing energy depends on the

clear aperture, D , of the amplifier and the stimulated-emission cross section, σ , of the laser material. For Q_s -optimized disk amplifiers, the relative change in cost effectiveness can be approximated by

$$Q_{s1} = \frac{Q_{s0}^*}{\sqrt{\sigma}} \frac{1 + b}{1 + b(D_0/D_1)^2} \quad (31)$$

where $Q_{s0}^* \approx 55 \text{ Jpm/k\$}$ for disk amplifiers. Size scaling can be derived from a model containing a cost component that is independent of D . For laser glasses with $\sigma \approx (2 \pm 1) \text{ pm}^2$, a scale factor of $b \approx 0.025$ and maximum segment size $D_0 \approx 0.5 \text{ m}$ provides a reasonable match to the amplifier simulations shown in Fig. 7-73. Glasses with higher stimulated-emission cross sections are limited to smaller segment sizes to avoid internal ASE losses.

The laser-chain design can be optimized by applying these generalized performances and cost scalings for various amplifier sizes and evaluating the remaining variables for the chain assembly. For comparison of different system architectures, the focusing and target requirements are usually quite similar and, hence, do not affect such a side-by-side comparison. It is necessary to add other correction terms to Eq. (30) only if there are distinctly different requirements with respect to the target chamber or facilities. The comparison of different laser architectures can then be accomplished by applying the same basic amplifier characteristics for the evaluation of the most costly output stages, while the drive stages

Fig. 7-73. Normalized cost effectiveness of storing energy, Q_s , for various amplifiers and laser materials.

and other specific requirements are added in the form of appropriate correction terms.

Cost Scaling of Driver Stages. For linear-chain designs with low stage gains of ~ 2 to 3, a relative drive-cost factor of $k \approx 1.2$ was obtained from systematic chain optimizations and applied for most previous evaluations. However, for multipass systems with high-gain stages, the relative drive-cost factor, k , must be appropriately scaled with the gain of the output stage. Such a scaling can be obtained from the relative cost of various-sized amplifier stages by evaluating

$$k = \frac{Q}{Q_D} = \frac{E_o C_D}{E_D(C_o + C_D)} = \frac{G_{so}}{1 + C_o/C_D} \quad (32)$$

The stage costs include fixed costs for spatial filters, isolation, and alignment. The costs are, to the first order, proportional to D . We approximated the fixed cost by $\Delta C \approx \xi D$, where $\xi \approx 2.3$ k\$/cm for spatial filters.

The relative drive costs are determined by

$$\frac{C_D}{C_o} = \frac{\sum_{i=1}^I (C_{Ai} + \Delta C)}{C_{Ao} + \Delta C_o} = \frac{\sum_{i=1}^I \frac{Q_{so} C_i^2 \ln G_i}{Q_{so} D_o^2 \ln G_o} + \Delta \sum_{i=1}^I \frac{D_i}{D_o}}{1 + \Delta} \quad (33)$$

where G is the small-signal gain of the amplifiers and i is the stage number, counting from the output stage with $i = 0$.

For evaluation of the drive-cost factor, k , three additional considerations were included in our analysis. First, for the assembly of an amplifier chain, it is desirable to reduce the output fluence of the driver stages, F_D , in order to reduce the risk of damaging lower-cost components. Fluence reduction was scaled with the saturated gain in the form

$$\frac{F_i}{F_o} = \frac{1}{G_{oi}^\delta} = \left(\frac{E_i}{E_o} \right)^\delta \quad (34)$$

with $\delta \approx 0.1$; G_{oi} is the saturated net gain, counting from the output stage to the i th driver stage. Size scaling of the driver stages is then

$$(D_i/D_o)^2 = G_{oi}^{\delta-1} \quad (35)$$

Second, for complete system optimization, Q must be evaluated for various gains of the individual stages. Complexity was reduced by scaling the saturated amplifier gains by

$$G_{oi} = G_{oi}^m \quad (36)$$

and optimizing Q for various values of G_{oi} and $m \geq 1$, where $m = 1$ corresponds to isogain stages. Saturated net gain is then determined by

$$G_{oi} = T^i G_{oi}^m \quad (37)$$

where T is transmission of the interstage components and

$$m^* = \sum_{i=1}^I m^{i-1} = \frac{m^I - 1}{m - 1} = I \quad (38)$$

for $m = 1$.

Third, to incorporate the cost reductions from building fewer stages with more identical components, a learning factor of ~ 0.9 was applied. We scaled the total cost with $N_s^{-0.1}$, where N_s is the number of stages required to provide a total-chain gain of $G_{tot} > 60$ dB. An arbitrary limit on the number of stages was necessary because the sum in Eq. (33) does not converge for a fixed-cost component per stage.

The results of our systematic chain optimization are shown in Fig. 7-74 for various saturated gains of the output amplifier. The number of stages, N_s , the gain scaling factor, m , and the fluence ratio

$$\sigma = \frac{F_o}{n F_s} = \frac{F_o \lambda \sigma}{n h \nu} \sim \sigma \quad (39)$$

were optimized for fixed values of the output fluence $F_o = 20$ J/cm², refractive index $n = 1.53$, effective degeneracy ratio $\kappa \approx 1.3$, and photon energy $h\nu \approx 1.87 \times 10^{-19}$ J. The fluence ratio is proportional to the stimulated-emission cross section, σ , of the glass or inversely proportional to the saturation fluence, F_s .

It is clear from Fig. 7-74 that k is primarily determined by G_{s0} and b , but is relatively insensitive to all other variables. The approximation

$$\Delta k \approx \frac{3.1b^{0.7} \frac{G_{s0} - 1}{G_{s0}^{0.3}} + \frac{1 - T}{1 + 150b^2}}{T} \approx 0.25 \frac{G_{s0} - 1}{G_{s0}^{0.3}} \quad (40)$$

matches the results quite well for typical cases and reduces the comparison of different laser architectures to the evaluation of

$$Q \approx \frac{\eta Q_{s0} (1 - 0.25/G_{s0}^{0.3})}{1 + \Delta} \quad (41)$$

Generalized Results. For linear chains, the highest performance-to-cost ratios are again observed for saturated-stage gains of about 2 to 5, depending primarily on the fixed cost ΔC , as shown in Fig. 7-75. Lower-gain stages are more expensive, while higher gains lead to a loss in extraction efficiency. For multipass systems, the optimum stage gain would be shifted toward higher gains because high extraction efficiencies can be maintained at high gains.

Evaluating peak values of Q leads to the approximation

$$Q \approx \frac{\eta Q_{s0} T^{1/3} \sqrt{1 - \delta}}{(1 + c_0 b) / (1 + c_1 a^{0.7} T^{1/3})} \quad (42)$$

with coefficients

$$c_0 = 10T^2 (1 - \delta)^{2/3},$$

$$c_1 = 1/(1.1/T + 15b),$$

and

$$a = \frac{\Delta \ln G_{s0}}{\sqrt{\alpha}} = \frac{4\xi Q_{s0}^*}{\pi \sqrt{\kappa nh\nu} D_{s0} \sqrt{F_0}} \approx 0.23 \frac{1}{\text{k\$}} \xi / \sqrt{F_0} \quad (43)$$

The relatively small impact of both the fluence-reduction factor, δ , and the variations in coefficients c_0 and c_1 leads to a rougher but more practical scaling of

$$Q \approx \frac{f Q_{s0} T}{(1 + 6b)/(1 + 0.7 a^{0.8})} \quad (44)$$

The cost effectiveness of extracting energy Q_{t0} can be approximated by

$$Q_{t0} = \eta Q_{s0} \approx \frac{1 - e^{-\beta \bar{v}}}{\kappa \sqrt{\alpha}} Q_{s0}^* \quad (45)$$

where $\bar{v} = \bar{F}_0 \kappa \sigma / nh\nu$ is the average fluence ratio in the amplifier, and

$$\beta = \frac{\gamma}{\alpha} \approx \frac{-\ln T}{\ln G} \quad (46)$$

is the loss-to-gain ratio, with γ the loss coefficient and α the gain coefficient. The maximum Q_{t0} for a single-pass system is obtained at

$$\bar{v}_0 \approx \frac{1.26}{1 + 2.5 \beta} \quad (47)$$

which leads to

$$Q_{t0} = \frac{Q_{s0}^* \bar{F}_0}{1 + 1.9 \beta} \quad (48)$$

where

$$Q_{s0}^* = \frac{Q_{s0}^* (1 - e^{-1.26})}{\sqrt{1.26 \kappa nh\nu}} \approx 5.8 \sqrt{\text{J cm/k\$}} \quad (49)$$

Therefore, the peak cost effectiveness of a linear chain is primarily determined by

$$\dot{Q} \propto \frac{Q_{t0}^* F_0 f T}{(1 + 1.9 \beta)(1 + 6b)(1 + 0.7 a^{0.8})} \quad (50)$$

where Q_{t0}^* reflects both design and cost of the disk amplifiers and largely depends on the specific costs of ME, EE, and optics. Improved amplifier designs have a direct impact on $\dot{Q}$ via Q_{t0}^* . The fill factor, f , and interstage transmission, T , are essentially proportional to $\dot{Q}$, while the loss-to-gain ratio, β , drive-scaling factor, b , and relative fixed costs, a , are somewhat less sensitive. The final leverage item is the output fluence, which will significantly reduce the cost per joule of Zeus in comparison with Nova and Shiva.

We may now compare multipass systems and other architectures with our generalized

analysis. The advantage of multipass systems is in the reduction of the output-fluence ratio, ϕ_N , since a reduction of this ratio increases the value of Q_{E0} by $\sqrt{\phi/\phi_N} \approx N_p^{0.4}$. In addition, the extraction efficiency of multipass systems remains high even for large saturated gains, which dilute the cost and transmission losses of interstage components, even if the multipass systems are relatively more costly and complex. The weaknesses of multipass systems are their increased effective loss-to-gain ratios, the potential reduction in the fill factor, and the higher fixed-cost components that reduce the $N_p^{0.4}$ improvement to somewhere between 1.5 and 2. Our general analysis does not yet include the miscellaneous costs associated with buildings and maintenance that can be added quite easily but are difficult to quantify.

Our method of scaling the performance-to-cost ratio of various laser architectures provides some indication of the relative sensitivity of individual variables and can

Fig. 7-74. Relative drive-cost factor, k , for various amplifier-scaling factors, b .

Fig. 7-75. Normalized cost effectiveness for linear-chain designs.

help determine where the largest potential improvements exist in order to maximize R&D efforts.

Selection of Laser Glasses. The most cost-effective, fluence-limited system design does not operate at maximum extraction efficiency. Rather, it operates at the maximum cost effectiveness of extracting energy [Eq. (44)], which includes the improved energy-storage capability of low-cross-section glasses. In Fig. 7-76, single and multipass systems are evaluated for various saturation fluences of laser glass. With amplifier gain and transmission losses held constant, the illustration shows that, for a given fluence limitation, the highest cost effectiveness is at about 83% of peak extraction with a higher-storage material (dashed curves in Fig. 7-76).

A laser designed for maximum extraction efficiency could not take advantage of future improvements in damage limitations, whereas a cost-optimal laser could operate at higher energy. It is also better to select a material with a higher saturation fluence than the optimum saturation fluence in order to reduce pulse distortions, ASE on target, and isolation requirements. Figure 7-76 also shows that a significant improvement in cost effectiveness derives from multipass systems and that the optimum saturation fluence for multipass systems is higher than for single-pass systems, assuming similar output-fluence limitations. It is also clear

that the most cost-effective multipass systems do not operate at maximum extraction efficiency. From a more accurate and detailed analysis, a rough approximation of the optimal saturation fluence, F_s , can be obtained from

$$\alpha_s = \frac{F_s}{nI_s} \approx \frac{1.6}{N_p^{0.8}} \quad (51)$$

or

$$I_s \approx 0.4 I_p N_p^{0.8} \quad (52)$$

For example, a fluence limitation of $F_s \approx 25 \text{ J/cm}^2$ would lead to optimal saturation fluences of 10 and 40 J/cm^2 for a single- and six pass system, respectively. Hence, our glass-development program should be directed toward high saturation-fluence glasses with stimulated-emission cross-sections in the range of $1 \pm 0.5 \text{ pm}^2$.

Reduction of Beam Modulations in Multipass Systems. Multipass systems have a very flat output response for large variations in input fluence, as shown in Fig. 7-77. It should be possible to maintain relatively smooth beam profiles with multipass systems.

An inflection point in the output fluence exists near maximum extraction efficiency and, hence, does not require any compromise with respect to energy extraction. However, this condition with the first and second derivative being zero is only satisfied for particular combinations of gain and loss in various multipass cavities.

The conditions of an inflection point for different numbers of passes, N_p , can be approximated by

$$G_o^* = \exp \left[\frac{12.3}{(N_p - 0.4)(1 + 2.9\sqrt{\beta})} \right] \quad (53)$$

$$i^* = \frac{0.52}{(N_p + 3\beta - 0.3)(1 + 0.24/\beta)} \quad (54)$$

and

$$o^* = \frac{10}{(N_p + \beta - 0.4)(1 + 13\beta)} \quad (55)$$

where i^* and o^* are the input and output fluence ratios at maximum energy extrac-

tion. Our approximations are within a few percent of the actual values of an inflection point for the range of $0.05 < \beta < 0.2$ and $2 < N_p < 6$.

The required amplifier gains for an inflection point differ from those selected for maximum cost effectiveness. However, the peak-to-average intensity ratio for our earlier analyses was assumed to be constant, and variations due to saturations have been ignored. On the other hand, the inflection-point conditions do not take into account differences in cost effectiveness. Combining

$$Q = \frac{IQ_s^*(o - ki)}{\lambda \sqrt{\sigma}(1 + \Delta) \ln G_o} \quad (56)$$

with the stimulated-emission cross section corresponding to a damage-limited peak output fluence, F_s ,

Fig. 7-76. Highest cost effectiveness is obtained at ~80% of peak extraction for single-pass systems.

$$\sigma = \frac{nh\nu}{\kappa \hat{F}_o} \frac{\delta}{\delta} \quad (57)$$

leads to

$$Q = \frac{fQ_{so}^* (\nu - ki)}{\kappa nh\nu (1 + \Delta) \ln G_o} \frac{\sqrt{\hat{F}_o}}{\delta} \quad (58)$$

which takes into account the peak-to-average ratio of the output fluence δ/ν for variations of the input fluence.

The optimization of this quantity for various multipass systems results in an output fluence ratio $\delta/\nu \approx 1.25 \pm 0.05$ for a 2-to-1 modulation of the input beam. In other words, the most cost-effective multipass system does not require an inflection point in the output fluence. Instead, a reduction of input modulations to about a 20 to 30%

output modulation is sufficient to provide high cost effectiveness.

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Evaluation of Cassegrainian Geometries

Cassegrainian geometries are a relatively low-risk option that is intermediate between single-pass and multipass systems. Relative to single pass systems, the improvement in performance-to-cost ratio is from 30 to 40%.

We have addressed the optimal amplifier gains, fluences, transit number, cavity length, and laser materials for Cassegrainian geometries. We distinguish the number of transits from the number of passes that are, in general, considered as transits with the

Fig. 7-77. Normalized output fluences and extraction efficiencies for various multipass systems.

same laser volume. In our analysis, the performance-to-cost ratio, Q , from Eq. (7-41) was evaluated for nominal 0.4-m-aperture amplifiers.

The evaluation of different laser architectures requires a reasonable assessment of fixed costs. Our analysis was based on the geometries shown in Fig. 7-78, where the single-pass stage is represented by a beam expander and an amplifier that is equivalent in its beam transport characteristics to the Cassegrainian geometry.

In our simplified approach, costs of the beam expander and Cassegrainian mirrors were assumed to be identical at 0.1 M\$/m² and correspond to Nova turning-mirror costs, including mounts and controls. The higher material costs for transmitting optics are assumed to be compensated by the higher fabrication and coating costs of the reflective optics for equivalent beam quality and pointing accuracy.

The effective cost for housing was assessed at 2 k\$/m², which corresponds to cost of the Nova laser bay and spaceframe. The length of the Cassegrainian cavity is

$$L = f_t \left(1 + \frac{1}{\sqrt{M}} \right) \approx \frac{\left(1 + \frac{1}{\sqrt{M}} \right) \left(\frac{1 + \sqrt{4f_t^2 - 1}}{3} \right) \ell}{2(1 - t_d)} \quad (59)$$

where f_t is the focal length of the large Cassegrainian mirror, M is the area magnification, ℓ is the length of the amplifier, and t_d is the effective beam-overlap fraction defined in Eq. (69), below.

The length of the disk amplifier was approximated by

$$\ell \approx n D N_D \left(1 + \frac{2t}{D \sqrt{1 + n^2}} \right) \approx 1.8 D N_D \quad (60)$$

where N_D is the number of disks and t is disk thickness.

The length of the single-pass system was based on an $f/10$ beam expander with

$$L_1 = \ell + f_{\infty} D \left(1 + \frac{1}{\sqrt{M}} \right) \quad (61)$$

The area-normalized fixed costs are then determined by

$$c_1 \approx 0.1 \left(1 + \frac{1}{\sqrt{M}} \right) + 0.002 \frac{L}{D} \frac{\text{M\$}}{\text{m}^2} \quad (62)$$

which, to a first order of approximation, are independent of the aperture size, since L/D depends primarily on N_D , M , and f_{∞} . Our simplifications make it possible to optimize the different architectures independently of aperture size. The maximum dimensions of the disk are then primarily determined by considerations of the suppression of internal amplified spontaneous emission (ASE).

The area-normalized cost of the amplifier can be approximated by

$$c_A^* = \frac{h\nu}{\sigma} n g_0 \frac{N_D}{Q_A} \quad (63)$$

where g_0 is the logarithmic gain of the disk. From the evaluation of different amplifiers and laser materials, the general approximation

$$g_0 = g_0^* \alpha \approx 0.15 \text{ pm}^{-1} \sqrt{\alpha} \quad (64)$$

was obtained, which leads to

$$c_A^* = n h \nu g_0^* \frac{N_D}{Q_A} \approx 0.78 \text{ M\$/m}^2 N_D \quad (65)$$

Fig. 7-78. Comparison of a five-transit Cassegrainian geometry and a single-pass geometry.

(a) Five-transit Cassegrainian geometry

(b) Single-pass geometry

Laser System and Component Evaluation

Thus, the area-normalized cost per disk is essentially constant for all Q_1 -optimized laser glasses, so long as material costs are of the same magnitude.

Quantity discounts were based on a learning factor of 0.9 and a normalized disk count (10 per size and chain), so that

$$c_A \approx c_A^* \left(\frac{10}{N_D} \right)^{0.1} \approx 1 \text{ M\$/m}^2 N_D^{0.9} \quad (66)$$

The fill factors of various geometries differ due to beam obstructions and the need for an "uncertainty" rim, ΔR , to avoid clipping of the beams. The fill factor for the Cassegrainian geometry is

$$f = \frac{(R_0 - \Delta R)^2 - (R_1 + \Delta R)^2}{R_0^2} \\ = 1 - \frac{1}{M} - 2 \frac{\Delta R}{R_0} \left(1 + \frac{1}{M} \right) \quad (67)$$

and for the single-pass stage is

$$f_1 = \left(1 - \frac{\Delta R}{R_0} \right)^2 \quad (68)$$

where $\Delta R/R_0$ was arbitrarily selected at 2.5% for a 1-cm-wide rim and a 40-cm clear aperture. The fill factor does not include the reduction in effective beam area due to apodization, which is assumed to be equal in all cases.

The performance analysis for the diverging beams of the Cassegrainian geometry requires a spatially resolved determination of energy extraction. The determination is obviously very slow and impractical for an

optimization procedure, where many geometries with various gains, magnifications, lengths, and fluences must be evaluated. To simplify the multidimensional problem, an effective beam-overlap fraction, τ , was defined as

$$\tau = \frac{A_{\text{eff}}}{A_{\text{out}}} \approx \int_0^1 \left(1 - \frac{z}{f_1} \right)^2 \frac{dz}{f_1} \\ = 1 - \frac{f_1}{3} + \frac{1}{3} \left(\frac{f_1}{f_1} \right) \quad (69)$$

with the assumption that the extraction of the diverging beam can be approximated by a parallel beam with an equivalent beam area that is reduced by τ , compared to the output-beam area.

The saturated-energy gain is calculated only for the double-pass region because it exhibits the lowest gain and determines the size of the driver stage. The saturated gain of outside rays is significantly higher because those rays experience single-pass gains. To provide a uniform output beam, the input beam to a Cassegrainian geometry must be apodized by reducing the intensity at the edge of the beam. Because apodization depends on the output energy, the complexity of operation is increased.

With our assumptions, the performance-to-cost ratio, Q_1 , can be optimized for particular values of the external transmission $T_o \approx 0.98$ and loss-to-gain ratio $\beta \approx 0.1$ by varying the amplifier gains, fluence, and cavity length for various numbers of transits. Our results for iso-fluent designs are summarized in Table 7-15. The optimal small-signal amplifier gains are in the range

Table 7-15. Optimized Cassegrainian geometries.^a

Number of transits, N_T	Amplifier gain, G_o	Saturated gain, G_s	Magnification, M	Fill factor, f	Saturation fluence, F_s (J/cm ²)	Nonlinear effects, ΔB_1	Amplifier cost, c_A (M\\$/m ²)	Fixed cost, c_D (M\\$/m ²)	Drive cost, c_D (M\\$/m ²)	Cost effectiveness, Q_1 (J/kS)	Q_1/Q_1
1	4.4	2.2	2.2	0.95	9.4	0.71	6.5	0.18	8.8	12.3	1.00
2	12.3	18.5	18.5	0.88	9.5	1.6	10.5	0.40	2.0	13.7	1.12
3	7.4	31.8	9.4	0.83	11.3	2.9	9.5	0.35	1.3	15.2	1.24
4	6.1	58.2	7.6	0.80	12.3	3.6	8.8	0.34	0.9	15.9	1.30
5	5.1	100	6.0	0.76	13.1	5.1	8.2	0.33	0.7	16.6	1.35
6	4.5	150	5.3	0.74	13.9	5.9	7.9	0.33	0.5	16.9	1.38
7	4.2	280	4.8	0.72	14.2	7.9	7.6	0.32	0.4	17.2	1.41

^aConditions for iso-fluent design: average output fluence $F_o = 20 \text{ J/cm}^2$, loss-to-gain ratio $\beta = 0.1$, external transmission $T_o = 0.98$, rim $\Delta R/R_o = 2.5\%$, mirror costs = $0.1 \text{ M\$/m}^2$, and space costs = $2 \text{ k\$/m}^2$.

of 2 ± 0.5 Np. Lower-gain amplifiers lead to costly drivers and reduced energy extraction; high-gain amplifiers are too expensive, as shown in Fig. 7-79, and the cost cannot be recovered by lower driver costs. The range of near-optimal gains decreases with increasing number of transits. Therefore, it is more important to select the correct amplifier gain for a five-transit system than for a single-pass system.

Our results depend importantly on the relative drive-cost factor, k . However, even a doubling of Δk would increase the small-signal gains for Cassegrainian geometries by only ~ 0.4 Np, while the optimal single-pass gain increases by ~ 1 Np.

The primary advantages of Cassegrainian geometries are the higher saturated gains and the ability to use higher-storage materials, thereby reducing the cost of drive energy and energy storage. In these respects, Cassegrainian geometries provide some of the advantages of multipass systems, even though their effective pass number is less than 2.

Clearly, the cost of drive energy dominates the single-pass cost, while it is less than 20% of the total cost for a two-transit

Cassegrainian system. However, the higher amplifier cost reduces the improvement to about a 12% net advantage of the two-transit Cassegrainian over a single-pass system. More transits certainly help to improve the performance-to-cost ratio, even at a reduction in the fill factor. However, the nonlinear effects of many transits became appreciable even at 10 ns for typical silicate glasses, as shown in Table 7-15. For laser fusion pulses, the saturated B value at the end of the pulse was approximated by

$$\Delta B_1 = \frac{0.7 - t}{1 - t} \quad (70)$$

where

$$B_0 = \frac{2\pi n_2^2 F_0}{\lambda n^2 \tau \alpha_0} \approx 0.6 \quad (71)$$

for nominal pulses of $\tau = 10$ ns, an output fluence of $F_0 \approx 20 \text{ J/cm}^2$, a nonlinear refractive index of $n_2 \approx 1.4 \times 10^{-13} \text{ esu}$, and a gain coefficient of $\alpha_0 \approx 5 \text{ m}^{-1}$.

The nonlinear effects of high-gain multitransit systems are of a serious magnitude for isofluent designs. The high ΔB

Fig. 7-79. Highest cost effectiveness for Cassegrainian geometries is obtained for amplifier gains in the range of 5 to 12.

values can be reduced by lowering the fluences in the driver transits and compromising Q values.

The optimal length of all geometries is about six times greater than the amplifier, which corresponds to an effective beam-overlap fraction of $f_0 \approx 90\%$. Longer geometries become too costly; shorter geometries have poorer overlap fractions and, hence, lower saturated gains. The total length-to-diameter ratio of all Cassegrainian geometries is in the range of 120 ± 20 , which corresponds to 50 m for a 0.4-m-aperture system.

The optimal extraction efficiency of all the geometries is in the range of 47%, which corresponds to 75% of maximum extraction efficiency. The compromise in extraction efficiency is exceeded by savings in the cost of storing energy.

Our general analysis allows us to determine the sensitivity of various system constraints. For example, the fixed cost for mirrors and space is typically less than 4% of the amplifier cost. With no significant impact on the performance-to-cost ratio, we are led to relatively long Cassegrainian geometries. However, more important variables are the external transmission and the loss-to-gain ratio, β , which leads to a scaling of the performance-to-cost ratio by

$$Q_B = fF_0 \frac{1 - k/G_{\text{sat}}}{c_A + c_F} \approx \frac{Q_0}{1 + 2\beta} \quad (72)$$

which is essentially the same as Eq. (48). A reduction of β from 1/10 to 1/30 leads to about a 13% improvement in Q , which is primarily due to better energy extraction. Our comparison of single-pass and multitransit systems is insensitive to changes in the absolute costs for ME, EE, and optics and to the absolute value in output fluence and loss-to-gain ratio.

Both Cassegrainian and reflex-ring geometries have unexpected energy-extraction features due to beam obstruction. For low-gain amplifiers and few transits, the hole in the beam becomes quite large, so that a significant portion of the amplifier is not used. For example, a three-transit system with an area magnification of $M = 2$ has an unused portion of the amplifier proportional to $1/M^2 = 25\%$. In other words, the maxi-

mum extraction efficiency is reduced by 25%. For Cassegrainian geometries, a further reduction of the effective extraction efficiency exists at locations where the laser beams do not overlap and the laser medium experiences only single-pass extraction. Losses in effective aperture are of the same magnitude as the vignette losses for off-axis multipass systems.

The maximum improvements in cost effectiveness for Cassegrainian geometries are only obtainable at increased ΔB values. A fixed ΔB limit of 3 would limit the potential improvements of Cassegrainian geometries to $\sim 30\%$, and there would be little incentive to increase the number of transits beyond five. Thus, Cassegrainian geometries are only applicable for long pulse durations. The additional cost for correction of astigmatic aberrations and the difficulties in beam relaying make Cassegrainian geometries less attractive than harmonically switched multipass systems.

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Nonlinear Characteristics of Saturated Amplifiers

For laser systems operating at high energy extraction, the nonlinear effects and system design largely depend on the laser pulse shape. Fusion targets require maximum power at the end of the pulse, when the gain of the amplifiers is depleted. Because of pulse distortions due to saturation, the highest beam intensities and nonlinear effects are experienced at that time. The system must be designed to provide sufficient intensity gain at the end of the pulse to avoid damage to optical components. Isointensity staging differs significantly from isofluence staging for fusion pulse shapes. Isointensity chains are much larger and more costly than isofluence chains.

The nonlinearity for a saturated disk amplifier element can be defined by

$$Y_c = \frac{dB_c}{dl} = \frac{2\pi n_2}{\lambda n^2 \alpha_c} = \frac{2\pi n_2}{\lambda n^2 \alpha_0} e^R = Y_0 e^R$$

where $dl = l\alpha_s d$, is the incremental gain in beam intensity, α_s is the saturated-gain coefficient, and R is the ratio of the beam fluence, F , relative to the effective saturation fluence, nF_s , defined by

$$R = \frac{\bar{I}}{nF_s} = \frac{F_A \sigma}{nh\nu} = F^* \sigma, \quad (74)$$

where F^* represents an effective photon fluence.

For short-pulse lasers, saturation is insignificant, and the nonlinear effects are proportional to $dB_s/dg_1 = Y_0 dl/n$. For long-pulse lasers and large amplifiers with $\alpha_0 \approx \text{constant}/D$, the nonlinearity of laser glasses is proportional to $n_2 e^{I^* n}/n^2$ and increases exponentially with the stimulated-emission cross section that dominates over n_2 considerations for $F^* da > dn_2/n_2$.

For saturated amplifiers, the nonlinear contribution dB_s per incremental peak-intensity gain $dg_1 = \alpha_s dz$ is a more practical measure of nonlinear effects

$$\frac{dB_s}{dg_1} = Y_0 \bar{I} = Y_0 \bar{I} e^{I^* n} \propto \frac{1}{\sqrt{\tau}} e^{\sqrt{\tau}/\tau_0}, \quad (75)$$

where $\tau = F/\bar{I}$ is an effective pulse duration. The evaluation of dB_s/dg_1 for the amplifiers of Table 7-16 and a fluence scaling proportional to $\sqrt{\tau}$ is shown in Fig. 7-80(a). It is interesting to note that dB_s/dg_1 has a minimum at

$$\tau_0 = \left(\frac{nF_s}{\bar{I}_1} \right)^2 \approx 3 \text{ ns} \quad (76)$$

for ED-2, and a minimum value of

$$\begin{aligned} \frac{dB_s}{dg_1} &= \frac{Y_0 \bar{I}_1}{\sqrt{\tau_0}} e^1 = \frac{Y_0 \bar{I}_1^2 e}{nF_s} \\ &= \frac{2\pi \epsilon_0 n_2}{\lambda h \nu \alpha_0} \bar{I}_1^2 \approx 2.6, \quad (77) \end{aligned}$$

where $\bar{I}_1 = 5 \text{ J/cm}^2/\text{ns}$ is the average fluence limitation at 1 ns. In other words, the nonlinear contributions for the same saturated intensity gain have a minimum at τ_0 and increase for longer pulse durations despite the reduction in beam intensity with $\sqrt{\tau}$. The minimum value depends on the laser glass and is proportional to n_2/n^3 .

Parameter	ED-2	LHG-8	LG-650
Stimulated-emission cross section (μm^2)	2.7	4.0	1.05
Index of refraction	1.556	1.523	1.512
Nonlinear index of refraction (10^{-13} esu)	1.41	1.05	1.3
Clear aperture (mm)	315	317	335
Disk thickness (mm)	48	54	32
Volume of laser glass per disk (L)	9.2	10.4	6.8
Nd-doping (wt%)	2.1	1.7	4.7
Absorption efficiency	0.41	0.37	0.42
Effective fluorescence lifetime (ms)	0.29	0.32	0.53
Pump-pulse duration $\sqrt{\text{LC}}$ (ms)	0.19	0.19	0.28
Pump energy (kJ)	240	210	270
Energy-storage efficiency (%)	2.35	2.23	3.52
Energy storage per disk (kJ)	2.8	2.3	4.8
Energy storage density (J/cc)	0.30	0.22	0.70
Gain coefficient (1/m)	4.6	5.0	4.1
Small signal gain	1.68	1.91	1.37
Saturated gain at 30 J/cm^2	1.16	1.13	1.19
Y_0	0.314	0.224	0.344
dB/dI , short pulse (relative units)	1.00	0.72	1.10
dB/dg_1 , long pulse (relative units)	1.00	1.08	0.44
ΔB for 30 J/cm^2 , 10 ns, $G_s = 2.5$	1.48	1.49	0.96
Cost of disks (k\$)	68	74	54
EE costs (k\$)	49	46	55
ME costs (k\$)	30	30	30
Total costs (k\$)	147	150	139
Cost of stored energy (\$/J)	26	32	14

Table 7-16. Rectangular-disk amplifier specifications^a for minimum cost of stored energy.

^aAll 14 standard lamps with 4 in. arc length, 1.5-cm bore. Disk material cost = $\$3/\text{cm}^3$, disk polishing = $\$3/\text{cm}^2$, fixed cost = $\$300$. Electrical storage = $\$0.15/\text{J}$ and 2 k\$/circuit.

Thus, our figure of merit for laser glasses at τ_0 is

$$M_s = \frac{n^3}{n_2 \sigma} \quad (78)$$

for saturated, parasitic-limited disk amplifiers. For pulse durations longer than τ_0 , the figure of merit becomes inversely proportional to the exponential of the stimulated-emission cross section σ ; for unsaturated amplifiers, it is independent of σ . The outcome is in contrast to small unsaturated and pump-limited amplifiers, where the figure of merit is proportional to $M_s = \sigma n^2 / n_2$.

Of 240 laser glasses evaluated, the highest figure of merit ($M_s = 1.31$) was obtained for fluoroberyllate B816; the best silicate glass ($M_s = 0.63$) was 3669A. Samples of LG-650, ED-2, and LHG-8 yielded values of 0.55, 0.21, and 0.20, respectively. Data are insufficient for other high- M_s glasses.

While our analysis applies strictly for amplifier elements, it also provides a good approximation for amplifier stages, so long as the average fluence, $\bar{F}$, is properly scaled to the output fluence, F_o . The ΔB value for

a disk-amplifier stage providing a saturated gain, G_s , is

$$\Delta B = Y_o I_o \frac{e - i}{1 - e^{-e}} = \frac{2\pi n_2 k F_o^2}{\lambda n^2 e \tau} \frac{1 - 1/G_s}{1 - e^{-e}} = \frac{\Delta B^*}{1 - e^{-e}} \quad (79)$$

where $I_o = F_o/\tau$ is the peak output-beam intensity at the end of the pulse, $e = \alpha_o h\nu/\sigma$ is the stored energy density.

$$e = \frac{F_o}{n l_s} \quad \text{and} \quad i = \frac{F_i}{n l_s} = \frac{e}{G_s} \quad (80)$$

For fusion-laser pulses, the effective intensity gain at the end of the pulse is

$$G_I = e^e = G_o e^{e^o} = \frac{1 - e^{-e}}{1 + e^{-e}} \quad (81)$$

while the small-signal gain, G_o , is

$$G_o = \frac{e^o - 1}{e^i - 1} \quad (82)$$

Fig. 7-80. Saturated nonlinearity of (a) amplifier element and (b) amplifier stage, where the fluence is scaled with $\sqrt{\tau}$.

The nonlinear contributions of ΔB per logarithmic peak-intensity gain, g_p , for amplifier stages providing a saturated gain of 2.5 are shown in Fig. 7-80(b) for a scaling of

$$F_o = F_i \sqrt{\tau} = 7 \frac{J}{\text{cm}^2 \cdot \text{ns}} \sqrt{\tau} \approx \bar{F} G_s^{0.37} \quad (83)$$

The differences in comparison with Fig. 7-80(a) could be further reduced with a more elaborate gain scaling.

It is important to distinguish between the saturated-fluence gain, G_s , and peak-intensity gain, G_i . The amplifier stages must provide not only sufficient saturated gain for energy but also enough intensity gain to amplify the peak power sufficiently to overcome the beam-area expansion ratio between stages. For fusion-laser pulses, the following general condition exists

$$G_s > G_i \geq \left(\frac{D_{i+1}}{D_i} \right)^2 \quad (84)$$

An isofluence design is not necessarily an isointensity design, and both requirements must be fulfilled. In other words, we must consider both a fluence limit and an intensity limit.

For nominal Gaussian pulses, the fluence and intensity remain essentially proportional to each other so that the two limits

are equivalent. In Fig. 7-81(a), a 3-ns Gaussian input pulse is amplified in five isofluence LHG-8 stages, each providing a saturated gain of $G_s = 2.5$ at an output fluence of 15 J/cm^2 . The peak intensity does not change dramatically despite the large temporal advancement.

However, as shown in Fig. 7-81(b), the situation changes rapidly for fusion-laser pulses. An output pulse shape was specified with a 5-ns foot at 5% of the peak intensity, followed by a 3-ns exponential rise to the peak intensity, which was then maintained for 2 ns. The required input pulse shape was then calculated for the same five isofluence stages, providing a total fluence gain of ~ 100 . The peak intensity at the end of the pulse for the driver stages would have to exceed the peak intensity of the output stage by over an order of magnitude. This is obviously not possible, and the output of Nova for a fusion-laser pulse will be much lower than for a nominal Gaussian pulse.

As illustrated in Fig. 7-81(c), the situation is less severe for low-cross-section materials with the same isofluence stages but different laser materials. Our examples show only a general trend and do not include two-ion saturation characteristics, which have not yet been measured for LG-650. It is, however, quite obvious that pulse distortions for

Fig. 7-81. Pulse distortions for (a) Gaussian input pulse, (b) fusion input pulse, and (c) various laser glasses.

Nova and Zeus will be of considerable concern, requiring pulse-shapers with a large dynamic range. To obtain the high contrast ratios of $\sim 10^6$, two or three Pockels cells in series would be required with extinction ratios of from 10^3 to 10^5 . For small Pockels cells, such high extinction ratios should be feasible by reducing the active volume of the crystal and, hence, the residual stress birefringence and optical scattering.

The examples shown in Fig. 7-81(c) are only for five stages providing a fluence gain of about 100, while a total saturated gain of $\sim 10^6$ is required for a system. Since we cannot afford to saturate all the driver stages and preamplifiers without severe pulse distortions, larger driver stages and preamplifiers are required than for an iso-fluence design. In this respect, multipass systems are much more favorable because they essentially provide large-aperture driver stages at no extra cost. The pulse distortions in multipass systems should be much lower than for linear chains.

It is clear from our comparisons that pulse shape has a large impact on the B of laser systems operating in saturation. The range in $Y_s = dB_s/dI$ extends from $Y_s = Y_{s,0}$ for pulses with the peak intensity at the leading edge to $Y_s = Y_{s,0}^R$ for pulses with peak intensity at the end of the pulse. An average value of $Y_s = Y_{s,0} (1 + e^{-R})$ was used for most evaluations of Shiva and Nova.

The selection of laser glasses with low cross section with respect to ASE, isolation, and pulse sharpening was recognized earlier in general terms. However, the large impact of pulse shape on the nonlinear beam distortions was not fully appreciated. In other words, selection of low-cross-section materials for Zeus is essential. For fusion-laser pulses, it is not sensible to operate any material much above its saturation fluence without a significant increase in nonlinear distortions. This will limit a high-cross-section material to relatively low output fluences and will require much larger output apertures or more chains than for low cross-section materials.

Our examples provide only an indication of general trends, rather than results for a fully optimized system. A critical assessment of pulse shape with respect to laser and target performance is important for lasers operating deep in saturation. In addition, the effect of pulse shape on damage limitations requires further examination. Since the high-intensity ripples on the beam due to nonlinear effects exist only for a short portion of the laser pulse, it may be possible to transmit higher fluences than those measured for a uniform beam.

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Two-Photon Color-Center Formation for Apodizers

Apodizers are required for high-power lasers to avoid diffraction patterns on hard apertures that can cause damage-limiting intensity modulations. Apodization of a beam without adding components could be accomplished by selectively solarizing a window or other optical component. The introduction of color centers by two-photon absorption (2PA) has interesting features

Fig. 7-82. Axial and radial transmission for a uniform $f/1$ beam at five different fluences.

that provide a variety of options. Color centers could be produced inside glass without causing significant discoloration at the surface by focusing a laser beam or other source inside the medium. Surface-damage limitations of bead-blasted apodizers could then be reduced.

As shown in Fig. 7-82, 2PA can produce a localized region of color centers. Since the radial and axial dimensions are about twice the beam waist for an $f/1$ beam, we could fabricate any desired three-dimensional apodization function on a $1\text{-}\mu\text{m}^3$ scale. For an information storage medium, such high storage density may be of interest.

The density of color centers, N , is assumed to be proportional to the energy absorbed by two-photon processes, such that

$$N = k\alpha_2 I^2 \quad (85)$$

where α_2 is the 2PA coefficient, and I is the beam intensity. Figure 7-82 shows the calculated transmission functions for a uniform $f/1$ beam of various fluences, F_0 , but equal intensities, I_0 , at the beam waist with diameter d_0 . The radial transmission, T_r , is defined as the transmission at different radii of the beam integrated along the beam direction. The axial transmission, T_z , is the transmission transverse to the beam direction for different locations along the beam axis, z . The absorption at the excitation wavelength and the dimension of the absorbing region increases with increasing fluence. The peak absorption also moves from the focal point toward the focusing lens due to self-absorption, which limits the formation of color centers beyond the focal point.

Self-termination and limited absorption can be overcome by moving the focal point during the process from inside the sample toward its surface. In this manner, opaque columns or lines can be obtained, as shown in Fig. 7-83(a). Enhancement may be necessary to compensate for the lower absorption of the color centers at $1\text{ }\mu\text{m}$ in comparison to the excitation wavelength, λ .

It is also more energy-efficient to scan the focal spot compared to a stationary case where self-absorption consumes a large fraction of the excitation energy. Figure 7-83(b) shows the axial transmission for beams with various $f/\text{Nos.}$ but identical

focal-spot diameters, d_0 . The radial transmission is identical in this case, where the intensity and fluence at the beam waist were scaled inversely proportional to $\sqrt{f/\text{No.}}$.

A detailed analysis indicates that a series of short, high-power pulses with good optical quality is most effective for color-center formation via 2PA. Of course, it is also possible to generate color centers by linear processes with a continuous UV source. The major difference is that the formation of color centers by UV cannot be concentrated as much as with 2PA. In addition, it is not possible to obtain as high a density of color centers due to increased self-absorption. The reduction in the localization of color centers is not important for a typical apodization width of a few millimeters, but

Fig. 7-83. Radial transmission for 10 axial locations at five different fluences, and axial transmission at constant beam energy.

Fig. 7-84. Off-axis layout for measurements of flashlamp-induced aberrations.

the obtainable absorption at $1\ \mu\text{m}$ is questionable for linear processes.

Since it is possible to control the absorption within the medium, it may be desirable to absorb the laser beam uniformly within the apodizer to distribute the thermal loading. Such an absorption profile can be obtained by scanning the focal spot toward the lens at an increasing rate.

Issues concerning the stability and absorption of color centers at $1\ \mu\text{m}$ and the damage limitations require further investigation. In general, it appears possible to produce a damage-resistant apodizer by introducing color centers in a window or another optical component. Such an apodizer could be combined with a bead-blasted apodizer if the absorption of the color centers alone were insufficient for apodization, but sufficient to avoid plasma formation at the bead-blasted surface.

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Flashlamp-Induced Aberrations in Disk Amplifiers

An accumulation of aberrations with pass number limits the useful number of passes for multipass-system architectures. Previous measurements with a multipass rod amplifier⁸⁷ showed severe aberration in the rod induced by flashlamp pumping. To aid simulations of multipass architectures with disk amplifiers, we measured the induced aberrations of a Shiva β amplifier. This amplifier consists of six ED-2 disks, each with a clear aperture of 9.2 cm, and is pumped by 16 flashlamps made of 1.5-cm-bore Ce-

doped quartz tubing with an arc length of 112 cm and a xenon fill pressure of 300 Torr. A continuous, cylindrical fused-quartz shield separated the disks from the flashlamps. All tests were done at an input energy of 174 kJ in a 600- μs pulse (3, LC), which produced a nominal small-signal gain of 4 at $1.06\ \mu\text{m}$.

We sampled amplifier aberrations with a collimated He-Ne probe beam. An off-axis multipass layout allowed the 8-cm probe beam as many as eight passes through the amplifier without an optical switch. Figure 7-84 shows the layout for two passes through the amplifier. Multipass mirrors R₁ and R₂ were maximum reflectors at $0.63\ \mu\text{m}$ but reflected less than 20% at $1.06\ \mu\text{m}$ to avoid self-lasing. Attenuation in ED-2 at $0.63\ \mu\text{m}$ made a photomultiplier necessary for adequate signals with six or more passes. A polarizer and bandpass filter helped discriminate against flashlamp light. Beam tubes throughout the layout were essential to stabilize the beam against turbulence in the air.

Our primary means of detecting aberrations induced by flashlamps was to monitor the transmission of the He-Ne probe beam through a spatial filter when the flashlamps were fired. We used pinholes and slits of various sizes, as well as a knife edge in the focal plane of the spatial filter, to characterize the induced aberration. We chose spatial-filter transmission rather than interferometric techniques because it was easier to obtain time-resolved data and to interpret the results with respect to energy on target.

Initial measurements showed both an aberration component that oscillated in time and the anticipated component, which increased with the time integral of flashlamp output. Figure 7-85 shows oscillations starting

Fig. 7-85. Flashlamp-induced aberrations include both oscillations and a thermal-optic component that follows the time integral of flashlamp output.

less than $200\ \mu\text{s}$ after the lamps and continuing for at least 10 ms, long after the lamp output stopped. The lower two traces in Fig. 7-85 show $1.06\text{-}\mu\text{m}$ fluorescence from the amplifier and transmission through a pinhole with a $50\text{-}\mu\text{rad}$ acceptance angle, both at $0.2\ \text{ms/div.}$ The upper traces show the same pinhole-transmission signal at one-fifth the sweep speed ($1.0\ \text{ms/div.}$). Since the absorption of pump light causes a temperature rise in the disks, we anticipated that nonuniformities in pumping would cause aberrations through linear expansion, the stress-optic effect, and direct refractive index variations with temperature. These thermal-optic aberrations would follow the temperature rise in the disks and, therefore, the integral of flashlamp output.

The oscillating aberrations in Fig. 7-85 cannot be caused by temperature changes in the disk because the diffusivity of ED-2 glass ($0.006\ \text{cm}^2/\text{s}$) is far too small to allow changes on the observed time scale. Therefore, oscillation must result from acoustics or mechanical motion of the disks. The signal amplitude from the oscillating aberration was comparable to or larger than that from the thermal-optic aberrations and often completely masked the latter. Thus, we first identified and suppressed the major sources of oscillation before attempting to characterize the more typical thermal-optic aberration.

Oscillating Aberrations

Several independent contributions caused the complicated temporal structure of the oscillating signal shown in Fig. 7-85. The major components were an acoustic disturbance in the atmosphere surrounding the disks and two types of disk vibration.

We isolated the influence of flashlamp-induced atmospheric disturbances by testing an amplifier without disks. The continuous, cylindrical blast shield of the Shiva β ampli-

fier completely isolated the disk cavity from the flashlamps. Installing windows against the end flanges allowed evacuation of the disk cavity to as low as 10 Torr. By varying gases and pressure we found that

- Aberration amplitudes qualitatively followed the refractive index of the gas. Pure N_2 and air produced comparable effects. SF_6 produced a signal several times larger and He produced a signal several times smaller than either N_2 or air.
- The dominant oscillation frequencies for N_2 (3.3 kHz) and SF_6 (1.5 kHz) corresponded to an acoustic-path length comparable to the amplifier aperture, suggesting a cylindrical pressure wave starting at the edge, converging through the center, and reflecting back from the far edge.
- N_2 and air gave the same amplitude within our measurement accuracy ($\pm 15\%$). We saw no indication of the O_2 absorption problem reported earlier,⁸⁸ probably because our Ce-doped flashlamp envelopes blocked the UV radiation absorbed by oxygen.
- He at atmospheric pressure or a reduced pressure of N_2 or air (≤ 50 Torr) dropped the aberration amplitude by a factor of 4.
- The aberration amplitude with N_2 or air at atmospheric pressure corresponded to either a full-angle steering of $6 \pm 2\ \mu\text{rad}$ per pass or an increase in beam divergence of $\sim 10\ \mu\text{rad}$ per pass.

The amplitude of the atmospheric effect seems relatively large, but it was measured without disks in the amplifier. With disks, the atmospheric contribution was apparently smaller, since the total oscillating aberration was reduced only slightly by evacuation. We reduced atmospheric aberrations to insignificant levels for all subsequent measurements by evacuating to ~ 20 Torr.

Disk vibrations resulted from excitation of natural vibrational modes of the disk and from driven flexure of the disk by the

$$\delta = d\alpha - d\beta = \frac{n^2 - 1}{n} \frac{t}{R}$$

n = index of refraction

t = disk thickness

R = radius of curvature

Fig. 7-86. Bent disks cause steering in the plane of incidence because the direction of the surface normal changes at the exit surface.

disk-support structure. In both cases, the disk bent into nonplanar shapes with parallel surfaces. Bending caused the aberration because the disks were not mounted perpendicularly to the transmitted beam, as shown in Fig. 7-86. Bending to a constant radius of curvature over the disk caused steering of the entire beam; a variable curvature caused higher-order aberrations. Experiments clearly showed the aberration to be predominantly in the plane of incidence of the beam and disks, as required by our explanation.

Spectrum analysis of transmission signals such as those in Fig. 7-85, showed a strong, well-defined component at 4.2 kHz and weaker components at 5.9 and 9.8 kHz. Mechanically striking either the disks or their support structure also produced the same components. Eigenmode calculations performed by R. Murray and D. Field (Nuclear Test Engineering Division) showed the components to be the first three vibrational modes of a β -size ED-2 disk. The calculated shape for the 4.2-kHz mode showed mostly bending to a relatively constant radius of curvature about the minor axis, which caused steering of the entire beam, as discussed above. Interpretation of our transmission data in terms of steering indicated an oscillating steering of $\pm 2 \mu\text{rad}$ for a single disk in the amplifier. This steering magnitude corresponds to a bending radius of 11 km and a sag at disk center of only $0.45 \mu\text{m}$. A fully loaded six-disk amplifier produced a steering of only $\pm 3 \mu\text{rad}$, indi-

cating considerable cancellation of the contributions from individual disks.

We also observed oscillating aberrations at frequencies from 200 to 400 Hz, well below the lowest natural frequency of the disks. Mechanical vibration of the disk-support structure, which, in turn, forced the disks to bend at the lower frequencies, produced these aberrations. Figure 7-87(a) shows oscillating aberrations from a single disk with the 4.2-kHz natural frequency riding on a 200-Hz frequency driven by the support structure. In Fig. 7-87(b), the disk mounting was modified to reduce excitation of the disk's natural frequencies, leaving only the driven frequency. We attribute the phase and frequency change between Figs. 7-87(a) and 7-87(b) to a change in location of the single disk in the six-disk holder.

To determine whether disk oscillations were initiated mechanically or via absorbed flashlamp light, we tested single ED-2 and BK-7 disks, either mounted or freestanding, in the amplifier. In the freestanding case, the disk stood on its edge on the blast shield without touching the metal disk holders or support structure. We saw no oscillations with either disk when freestanding and no significant difference between their oscillations when they were mounted. We concluded that thermal shock from absorption of flashlamp light in the disk does not initiate the oscillations and that they must be initiated by a mechanical shock from the support structure.

Accelerometer measurements showed that the mechanical shock to the support structure was not caused by magnetic fields associated with current flow through the lamps. For our test, the flashlamps were replaced by wires and external loads to eliminate light output but retain the current flow. An accelerometer attached to the end of the amplifier showed more than 10 times the axial acceleration with lamps than with wires. Since the current pulses for the two cases were nearly identical, the associated magnetic fields apparently did not cause the mechanical shock.

The shock responsible for the vibrations was most likely caused by a rapid heating of the disk-support structure by the flashlamps. Accelerometer measurements on the two end flanges of the amplifier

showed an axial acceleration of 13 g away from the amplifier midplane. Since the measured steering of $\pm 2 \mu\text{rad}$ from a single disk could be produced by an acceleration as small as 4 g, the source of the end-flange acceleration most likely caused the vibration as well. Calculations of flashlamp light absorption by the stainless-steel support structure indicated a sufficient average temperature rise and thermal expansion during the 600- μs flashlamp pulse to account for the measured acceleration. The rapid heating and high temperature reached by the surfaces of the stainless steel also provide a potential source for the atmospheric aberrations.

We suppressed the oscillating aberrations by isolating the disks from their holders and by minor modification to the disk-support structure. We used 0.1- by 0.5- by 0.05-in. silicon sponge-rubber foam pads (Kirkhill Rubber Co., Brea, Calif.) to prevent direct contact between the disk and its stainless-steel holder. An aluminum foil wrapping protected the pads from indirect flashlamp light. Figures 7-87(a) and 7-87(b) show pinhole transmission (50- μrad acceptance angle after four passes through a single disk) without and with foam-pad isolation, respectively. The reduction in 4.2-kHz aberration with isolation is clearly shown.

We reduced forced bending of the disks by cutting two of the four weldments between disk holders and their support structure. Figure 7-87(c) shows the total reduction in oscillating aberrations accomplished with both modifications. Quantitatively, the modifications reduced oscillating aberrations by about a factor of 3, enough that they no longer dominated the thermal-optic aberrations.

Thermal-Optic Aberrations

We used several measurements and techniques to resolve different components in the thermal-optic aberration. We relied most heavily on measurements of spatial-filter transmission using a knife edge or slit in place of the usual pinhole to allow direct determinations of the magnitude and direction of induced steering and induced

(a) Single disk, normal mounting.

(b) Single disk with foam isolation.

(c) Two disks with foam isolation and reduced connection between disk holder and support structure.

lensing. Pinhole measurements proved to be ambiguous because they could not distinguish steering direction or astigmatic lensing. Two additional techniques, a Smartt interferometer and monitoring of the probe-beam pattern through a slit, verified the knife-edge and slit-transmission results.

Figure 7-88 illustrates the use of a slit to determine induced lensing. The curve shows the best fit to slit-transmission data at different axial distances from the lens, measured before firing the amplifier. At each data point, the transverse slit position was optimized for maximum transmission. The two oscilloscope-trace inserts show the effect of firing the amplifier at slit locations ahead and beyond best focus. For the case

Fig. 7-87. Oscillating aberrations are reduced by foam isolation and by eliminating two of four weldments between disk holders and support structure.

Fig. 7-88. Transmission changes through a slit show a lensing component in the induced aberration.

shown, the initial slit position beyond best focus gave an increase in transmission on firing the amplifier, whereas transmission decreased for a slit positioned ahead of best focus. These changes show an axial shift in best focus away from the lens, indicating a negative cylindrical-lensing component in the total induced aberration. Rotating the slit by 90° shows the orthogonal cylindrical component.

Fitting a new transmission curve through several "after" data points, such as the two in Fig. 7-88, typically showed an overall decrease in transmission, as well as the axial shift, indicating induced steering or higher-order aberrations, as well as lensing. We characterized all additional aberrations with an increased beam divergence, $\Delta\theta$, defined by

$$\Delta\theta = \sqrt{(\theta_a^2 - \theta_b^2)/n} \quad (86)$$

where θ_b and θ_a are probe-beam divergences before and after firing the amplifier for n passes through the amplifier. We determined θ_b and the focused beam shape from measured slit transmission as a function of transverse position without firing the amplifier. We inferred θ_a from the decrease in transmission at best focus, assuming a constant beam shape.

Figure 7-89 shows two examples of the use of a knife edge to determine beam steering. The knife edge was initially positioned at best focus, but at various transverse positions relative to the beam axis. The solid curves indicate transmission before firing the amplifier, and the dashed

curves show a best fit to the "after" transmission values. Pure beam steering translates the after-transmission curve without changing its shape, as shown in Fig. 7-89(a). A slope change together with translation, as shown in Fig. 7-89(b), indicates an increase in beam size as well, caused by induced lensing or higher-order aberrations.

The two values of residual beam divergence obtained from our two different measurements contain different information. The value from knife-edge measurements includes lensing, while the slit value includes steering. We combined the results by calculating the influence of steering (determined from knife-edge measurements) on divergence (determined from slit measurements) and vice versa. The two corrected $\Delta\theta$ values, excluding both lensing and steering, should have been equal, since both characterized the same residual aberration. However, the corrected beam divergence measured with the knife edge was typically 50% larger than the slit data. Such variations are caused by the large uncertainties of our measurements.

We independently confirmed steering and lensing results with two additional techniques. We found that the spherical aberration added to the probe beam by the multipass optics created a distinctive beam pattern when spatially filtered by a slit just ahead of best focus. The pattern changed with axial and transverse position of the slit, providing a means to measure induced steering and lensing. In addition, a Smartt, or point-diffraction,⁸⁹ interferometer, which obtains its reference locally from the

focused probe beam, gave an interference pattern sensitive enough to resolve our magnitude of lensing and steering. Both techniques confirmed the lensing and steering results of the knife-edge and slit measurements, but neither technique was sensitive enough to show higher-order aberrations.

Table 7-17 shows a summary of the measured thermal-optic aberrations per pass in the amplifier's plane of incidence. The first row shows results for the full complement of six disks in the amplifier. Our measurements for two, four, and eight passes demonstrated that, within our measurement accuracy, both steering and lensing accumulate linearly with pass number. The $\Delta\theta$ value in Table 7-17 represents the average residual beam divergence of the knife-edge and slit results.

To determine the scaling of induced aberrations with the number of disks per pass, we also measured aberrations in the same amplifier with only two disks mounted in the central disk holders. We found significantly more steering in the opposite direction than for the fully-loaded amplifier and much less lensing (i.e., a longer induced focal length). With two disks, the two corrected $\Delta\theta$ values from slit and knife-edge measurements did not indicate any significant beam divergence within our accuracy of measurement. Apparently, pumping differences at the different disk locations and cancellation of individual disk contributions prevent a simple scaling relationship between induced aberration and number of disks per amplifier.

Measurements perpendicular to the plane of incidence showed all aberration components to be several times smaller. Magnitudes were not well-defined because of the smaller signals. Lensing for the two-disk case was negative, the same as in the plane of incidence; lensing for the six-disk case changed to positive.

The most consistent explanation for our results is that pumping differences for the end disks in the amplifier caused most of the aberration. For the two-disk case, the centrally located disks see a relatively uniform pumping field, except for pump light reflected from each other. The enhanced pumping of the interior surfaces due to reflections would heat those surfaces more than the exterior surfaces, bend the disks slightly, and steer the beam in the mea-

Number of disks	Steering (μrad)	Cylindrical lensing (kni)	$\Delta\theta$ (μrad)
6	2 ± 1	-4 ± 1.5	20 ± 10
2	-3 ± 1	-12 ± 12	0 ± 5

Table 7-17. Thermal-optic aberrations per pass in the plane of incidence of a β amplifier after a 174-kJ flashlamp pulse.

Fig. 7-89. Transmission changes past a knife edge show beam steering and beam spreading.

sured direction. A pumping asymmetry of $\sim 10\%$ would be sufficient to explain the observed steering. For the six-disk case, the central disks would see more uniform illumination and produce much smaller aberrations. The end disks, however, would see a large transverse gradient from reduced pumping caused by the ends of the flashlamps. The resulting thermal-optic wedge would steer the beam in the direction opposite to that of the two-disk case, as measured. A temperature gradient between 1 and 2°C across the disk would explain the observed steering magnitude and direction.

In summary, oscillating and thermal-optic aberrations of comparable magnitude can be produced by flashlamp pumping of a disk amplifier. Mechanical precautions can reduce oscillating aberrations for future amplifier designs, and our measurements suggest the order of magnitude to be expected of thermal-optic aberrations. Without efforts to reduce them, flashlamp-induced aberrations would limit multipass architectures to

relatively few passes in order to maintain adequate beam quality for harmonic generation.⁹⁰

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Advanced Lasers

Vanadium-Doped Solid-State Laser

We have previously determined that most of the spectroscopic and physical properties of vanadium-doped magnesium fluoride (V:MgF₂) are suitable for use in a fusion driver.⁹¹⁻⁹⁴ We have discovered, however, that an unfavorably strong excited-state-absorption (ESA) transition spectrally overlaps the ⁴T₂ — ⁴A₂ laser transition peaked at 1152 nm. This ESA process significantly lowers the gain coefficient of the laser medium, increases the thermal heat load in the material, and degrades the laser efficiency.

To give us some insight into the origin of the ESA transition, and to determine possible countermeasures, we examined the 1954 Tanabe-Sugano energy-level/crystal-field diagram... for a 3d³ electronic configuration.⁹⁵

We tentatively assigned the ESA transition to the high-energy wing of the ⁴T₂ — ⁴T₁ transition. The Tanabe-Sugano diagram was calculated for an assumed C/B parameter ratio most appropriate to the Cr³⁺ ion (and 3d³), but it is indicative of energy-level ordering in V²⁺. This diagram also indicates that the laser transition energy is proportional to the octahedral crystal-field strength, Dq, whereas the ESA transition energy is essentially independent of Dq. Thus, placement of the V²⁺ ion in an octahedral site with a larger Dq than exists in MgF₂ should shift the laser transition spectrally away from the ESA transition. (The validity of this simple argument is impaired as a result of lattice-relaxation phenomena associated with a triply degenerate ⁴T level in an octahedral field; however, there is at present no quantitative theory available to guide our thinking beyond the simple considerations cited.)

The octahedral, divalent lattice site for V²⁺ ions provided by the fluoride-perovskite crystal KMgF₃ has a larger Dq than that provided by MgF₂ (Ref. 95). We therefore speculated that V:KMgF₃ would exhibit a reduced amount of ESA and correspondingly improved characteristics for an efficient fusion laser. Accordingly, we have prepared and begun characterization of V:KMgF₃ single crystals. We have also implemented codes capable of calculating energy-level/crystal-field diagrams of specific ion/host systems, free of the restrictions on atomic parameters adopted in producing the prototypical Tanabe-Sugano diagrams (see the following article).

Crystal Growth. Using the crystal-growth techniques described in the 1980 *Laser Program Annual Report*, we have grown two V:KMgF₃ single crystals, each measuring 1.5 in. in diameter by 1.5 in. long and doped with 1 wt% V₂O₅. Figure 7-90 is a photograph of one of these crystals, showing aggregation of melt impurities at the end of the boule. The upper portion of the boule exhibits crack-free crystalline material with a relatively uniform distribution of V²⁺ along the growth axis. The crystalline material also exhibits a nonuniform "cloudy" appearance in the bulk. The top of the second crystal exhibited white crystals, which we found to be KHF₂. We assumed that the KF starting material, which is hygroscopic in nature, contained enough H₂O

Fig. 7-90. Boule of KMgF₃ crystal doped with 1 wt% V₂O₅, grown using the gradient-freeze technique.

to cause the cloudiness and allow formation of the KHF_2 crystals found in one of the growth runs.

We pulverized most of the material from the first two runs and used it to charge a third growth run in the hope that some zone refining from previous runs would have reduced the H_2O contamination in the charge. The crystal resulting from the third run was colorless (no discernible vanadium content) and was at least as cloudy as the two earlier crystals. In addition, the inside of the growth furnace was covered with a yellowish white deposit.

At this point we assembled an apparatus to remove H_2O from the commercially obtained KF starting material and performed several experiments to determine the proper operating procedure for purification. The final procedure consisted of putting the KF powder in a vitreous carbon boat, placing the boat inside a quartz tube, and placing the tube in a clamshell resistive furnace through which dry HF gas was passed. When HF gas was injected over the liquid KF a considerable amount of water appeared in the exit line of the system, giving us reasonable assurance that we now have

a reliable method for removing H_2O from KF.

Spectroscopic Properties. Absorption, emission, and lifetime measurements have been carried out using available V:KMgF₃ crystals. Figure 7-91 shows the spectral-absorption cross sections of V:KMgF₃ and V:MgF₂, compared with the relative spectral emission intensity of a xenon flashlamp. The cross-section values for V:KMgF₃ are only approximate at present, since the absolute V-ion concentration in these crystals has not yet been established and because scattered light at shorter wavelengths obscures precise definition of the baseline. Nevertheless, these preliminary data show that V:KMgF₃ exhibits the expected blue shift in the 4T_2 upper laser level and that V:KMgF₃ is well-coupled spectrally to xenon-flashlamp emission.

Figure 7-92 shows the fluorescence lifetime of the $^4T_2 \rightarrow ^4A_2$ emission of V^{2+} in KMgF₃, MgF₂, and MnF₂, as functions of temperature. The V:KMgF₃ medium exhibits a low-temperature fluorescence lifetime of 1.3 ms, compared to lifetimes of 2.48 and 3 ms for V:MgF₂ and MnF₂, respectively. Note also the strikingly different

Fig. 7-91. Spectral-absorption cross sections of V:KMgF₃ and V:MgF₂, and spectral emission of a xenon flashlamp.

temperature-dependence of the fluorescence lifetime of V:KMgF₃ compared with the other crystals; the V:KMgF₃ fluorescence lifetime remains essentially constant for temperatures up to 350 K, whereas the lifetimes of MgF₂ and MnF₂ fall rapidly above temperatures of 200 and 150 K, respectively.

An earlier attempt to theoretically predict the temperature-dependence of the V:MgF₂ fluorescence lifetime, using the single-configuration-coordinate model of Struck and Fonger, was unsuccessful.⁹² While there is no quantitative basis for anticipating the temperature-dependence of the fluorescence lifetime of V:KMgF₃, we cannot rationalize the observed dependence in terms of the

Fig. 7-92. Fluorescence lifetimes of the ${}^4T_2 \rightarrow {}^4A_2$ level of divalent vanadium ions in KMgF₃, MgF₂, and MnF₂ crystal vs temperature.

Fig. 7-93. Fluorescence spectrum of the ${}^4T_2 \rightarrow {}^4A_2$ transition of V:KMgF₃ at room temperature.

relatively small change in crystal-field strength (see the following article). The observed difference may be dominated by the structural differences between host crystals; KMgF₃ is a cubic perovskite with perfect octahedral symmetry at the Mg²⁺ site, whereas both MgF₂ and MnF₂ have the rutile tetragonal structure with D_{2h} symmetry at the Mg²⁺ site (see Ref. 96).

Figure 7-93 shows the room-temperature fluorescence spectrum of V:KMgF₃, corrected for the spectral response of the detection system (but not for possible reabsorption of fluorescence emission at the shorter wavelength). As expected, the peak fluorescence wavelength (1048 nm) of V:KMgF₃ is blue-shifted from the 1152-nm fluorescence peak of V:MgF₂. Given that the fluorescence line widths for V:MgF₂ and V:KMgF₃ are virtually the same, and that the fluorescence (radiative) lifetime of V:KMgF₃ is a factor of 2 smaller than that of V:MgF₂, we anticipate that the stimulated-emission cross section of V:KMgF₃ will be about twice as large as that of V:MgF₂ (or in the range of 8×10^{-21} cm²; see Ref. 93). This larger cross section will provide a more comfortable margin between the saturation fluence of the gain medium (estimated at 24 J/cm²) and the damage fluence. The preservation of a millisecond-long fluorescence lifetime at temperatures at least 50 K above ambient is also a major improvement in properties of V:KMgF₃ compared with V:MgF₂.

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Energy Levels of 3d^N Configurations in Crystals

To interpret spectroscopic and laser-kinetics data, and to help guide the selection of useful laser ions and crystal hosts for laser-fusion drivers, it is necessary to be able to describe the nd^N energy levels of the ions in various high-symmetry (octahedral and tetrahedral) crystal-field sites. During 1982, we adapted and developed a series of codes that allow us to

- Calculate, in the weak-field basis, matrix elements of F^2 , F^4 , ζ , α , and any crystal-field parameters, B_q^k (Ref. 97).

- Diagonalize the matrices for any desired values of the parameters.
- Perform least-squares fits to experimental energy levels.
- Generate Tanabe-Sugano diagrams for any set of parameter values.

We have applied our codes initially to the analysis of V^{2+} energy levels in $KMgF_3$, MgF_2 , and MgO .

There are two problems associated with previous analyses of spectra by Sturge et al.¹⁶ First, the calculated energy levels were compared to the centers of wide vibronic bands; the theory calculates the zero phonon transition that is somewhere near the onset of these bands. Second, F^2 and F^4 were allowed to vary independently, even though the data are insufficient to determine both of them; the result was ratios, F^4/F^2 (or C/B), that are less than the free-ion values. In the lanthanides and actinides, F^4/F^2 is always greater in crystals than in the free-ion case.

By choosing the onset rather than the peak of the absorption bands, and by fixing F^4/F^2 in the expected range, we have obtained a number of least-squares fits to the data; the best results are summarized in Table 7-18. While our results are preliminary, it is clear that it is possible to obtain reasonable agreement between theory and experiment with parameter values that behave as expected with increasing ligand polarizability, i.e.

- F^2 decreases.
- F^4/F^2 increases.
- $f_2 = F^2/F^2(\text{free ion}) > f_4 = F^4/F^4(\text{free ion})$.
- $f_2, f_4(\text{fluoride}) > f_2, f_4(\text{oxide})$.

It is not yet clear whether the poorer fit for $V^{2+}:\text{MgO}$ than for V^{2+} in fluorides is an indication that the model is less appropriate for oxides.

A Tanabe-Sugano diagram (including spin-orbit coupling) with the free-ion parameters of $V^{2+}:\text{KMgF}_3$ is shown in Fig. 7-94; the positions of the crystal fields of $KMgF_3$ and MgF_2 are indicated by vertical lines. In both cases, the crystal-field strength is such that we are close to the avoided crossing of the 2E and 4T_2 Γ_8 levels ($B_0 \sim 2.9 \times 10^4 \text{ cm}^{-1}$). If the crystal field of $V^{2+}:\text{MgO}$ were plotted on this diagram, it would be just to the right of the avoided crossing, but the decrease in F^2 in going from fluoride to oxide ligands moves the

crossing down to $B_0 \sim 2.4 \times 10^4 \text{ cm}^{-1}$ so that the crystal field is actually well to the right, and the order of the quartet and doublet levels is reversed from the fluoride case.

Our preliminary results may improve with feedback from analyses of other $3d^N$ systems. It appears possible, however, to calculate the V^{2+} levels with parameters consistent with the intuition built up from study of the lanthanides. In the next year, we will continue this analysis with emphasis on nd^2 and nd^3 ($n = 3$ to 5) systems, which have the greatest potential for laser applications.

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Advanced Storage Laser: NpF_6

Energy-storage laser media are attractive for use in fusion lasers because they allow for slow pumping and short-pulse extraction. Beitz et al.⁹⁹ recently observed that NpF_6 vapor absorbs 1.06- μm radiation and fluoresces at 1.35 μm with a quantum efficiency near unity (see Fig. 7-95); the measured fluorescence lifetime was 4 ms, and there was no measurable collisional deactivation by unexcited NpF_6 molecules. Stimulated by these observations, we have analyzed NpF_6 gas as an energy-storage laser medium.¹⁰⁰

The vapor pressure of NpF_6 at room temperature is 100 Torr, a value ideally matched to its estimated stimulated-emission cross section of 10^{-20} cm^2 (Ref. 100). The undesirable properties of NpF_6 , such as alpha-decay radioactivity (with a half-life of 2.2 million years), and dissociation by visible light ($\lambda < 600 \text{ nm}$), can be handled with proper system design.

We set up an experiment to measure the lifetime of NpF_6 under conditions of high excited-state density, as would obtain in laser-fusion applications; Fig. 7-96 shows the experimental layout. A $\text{Nd}^{3+}:\text{YAG}$ laser was operated in a long-pulse ($\sim 300 \mu\text{s}$) regime, with typical energy output of $\sim 1 \text{ J}$. The beam was reduced to a spot size of $\sim 1.5 \text{ mm}$ diam to give fluences of up to 40 J/cm^2 in the excitation region. Calibrated attenuators were used to vary the fluence during the experiment. Fluorescence from the excited NpF_6 was collected with a lens having a 10-cm focal length and imaged

	Free-ion		KMgF ₃		MgF ₂		MgO	
	LLNL	Sturge et al.	LLNL	Sturge et al.	LLNL	Sturge et al.	LLNL	Sturge et al.
F^2	58 718 ± 302	57 397	51 828 ± 660	53 362	51 775 ± 622	53 362	41 233 ± 1371	47 625
F^4	35 174 ± 424	32 634	33 170	28 065	33 136	28 065	29 687	25 455
F^4/F^2	0.60	0.57	[0.64] ^a	0.53	[0.64]	0.53	[0.72]	0.53
f_2	—	—	0.88	0.93	0.88	0.93	0.70	0.83
f_4	—	—	0.94	0.86	0.94	0.86	0.84	0.78
ζ	[170]	—	[150]	0	[150]	0	[135]	0
α	51.5 ± 5	79	[45]	[79]	[45]	[79]	[45]	[79]
β	-266 ± 73	0	[-200]	0	[-200]	0	[-200]	0
B_0^4	—	—	23 605	24 150	21 942	25 515	26 165	29 400
σ	173 ^b	—	184	—	167	—	428	—
B	800	801	682	771	681	771	521	683
C	2792	2590	2632	2227	2630	2227	2323	2020
C/B	3.49	3.23	3.86	2.89	3.86	2.89	4.45	2.96
Dq	—	—	1124	1215	1045	1150	1238	1400

^aValues in square brackets were held fixed.^bValues for σ = rms deviation.Table 7-18. V^{2+} parameters. Fig. 7-94. Tanabe-Sugano diagram with the free-ion parameters of V^{2+} :KMgF₃.

Fig. 7-95. Energy levels of NpF_6 . Fig. 7-96. Experimental apparatus for measuring the fluorescence lifetime of NpF_6 under high excited-state density.

through a monochromator; we used a germanium photodiode as the detector.

NpF_6 was produced by the reaction of molecular F_2 with solid neptunium metal at a temperature of $\sim 923 \text{ K}$; the product was then distilled, kept in a trap cooled to 77 K , and protected from room light. For the optical experiment, we filled small quartz cells (having a volume of $\sim 1 \text{ cm}^3$) with a sufficient amount of NpF_6 to leave an observable speckle of orange solid in the cell; this assured us of known vapor pressure. Fluorine and SiF_4 gas were added to some of the cells to study the effect of buffers. The quartz cells were hermetically sealed and placed inside a sealed steel container. All the windows in the steel container were covered with Corning 2-58 filters, which transmit IR and cut out visible radiation. The quantity of NpF_6 handled at any given time was sufficiently small that no observable ionizing radiation was detectable outside the steel container.

Figure 7-97 shows typical experimental data. At low incident intensity, we observed long decay times, as expected. Raising the input fluence, however, increased the decay rate. Though we obtained the best results with 1 atm of SiF_4 buffer gas, there was little improvement in increasing the buffer

pressure from 0.5 to 1 atm; thus, going to even higher pressures is not likely to provide a solution.

To estimate the fraction of molecules excited, we plot in Fig. 7-98 the peak fluorescence vs the input energy measured in saturation fluences ($E_{\text{sat}} = h\nu/\sigma \sim 10 \text{ J/cm}^2$). By fitting the low-energy part of the curve to a straight line ($N^*/N_0 = E/E_{\text{sat}}$), we obtain the absolute calibration for the vertical scale. As can be seen from Fig. 7-98, approximately half of the molecules are in the excited state. This fraction corresponds to a stored energy density of $\sim 300 \text{ J/l}$.

RAPIER-B Amplifier

which is a record for high-density energy storage in a g is medium.

The decay rate of NpF_6 is a function of the degree of excitation. Plotting the decay rate, $-N^*/N^*$ vs N^* , we see in Fig. 7-99 that, after the laser pulse is over, the decays

tend to fall on a straight line. This indicates that the quenching rate is proportional to the excited-state density and that $(\text{NpF}_6)^*$ is most likely its own quencher. Thus, if a storage time of 1 ms is required, the energy density must be reduced to 200 J/l.

We observe a somewhat higher decay rate immediately after the termination of the laser pulse, indicating the possibility of direct dissociation of NpF_6 , followed by quenching by the fragment of dissociation. More work is required, however, to fully understand the kinetics of such high excitation densities, since very little is known about the dissociation products.

In summary, we have shown that NpF_6 molecules can store energy at a very high density of excited states, some undesirable quenching processes can be explained, however, but we can fully evaluate the potential of NpF_6 as an energy storage medium.

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RAPIER-B Amplifier

During 1982, we completed the assembly, installation, and characterization of the pulsed power system for the electron beam pumped RAPIER-B amplifier. The system consists of two groups of three modules arranged on opposite sides of a single laser cell, together with several supporting subsystems. Each module consists of a Marx generator, a pulse forming line, an output switch, and an electron-beam diode. The supporting systems include output switch triggers and diagnostics packages for each module; a Marx trigger generator, charging system, and diode vacuum system for each group of three modules; and a water system and insulating-oil system for the complete pulsed-power system. This setup is shown schematically in Fig. 7-100.

Supporting Systems

The RAPIER-B amplifier's charging system, water system, and diagnostic system are the most notable of the supporting systems. The charging system consists of a 100-kV

Fig. 7-97. The 1.35- μm fluorescence intensity of NpF_6 gas containing 1 atm of SF_6 as a buffer.

Fig. 7-98. Peak fluorescence power of NpF_6 vs 1.06- μm input energy, measured in saturation fluences.

Fig. 7-99. Fluorescence decay rate of NpF_6 vs excited-state density.

Fig. 7-100. Block diagram of one of six modules of the RAPIER-B amplifier.

60 MA power supply with a controller-sequencer that charges the Marx generators at constant current (in ~ 20 s), stops charging at a preset voltage, and disconnects the power supply from the system before giving permission to fire. The water system provides continuously circulating water in separate circuits for each of the pulse-forming lines and prevents the lines from draining if a power failure occurs. A separate flow circuit continuously circulates water through a vacuum chamber, to deaerate the water, and then, through ion-exchange resin columns, to maintain the water resistivity (>10 M Ω -cm).

Three diagnostic signals are provided for each module, and we also monitor the output-switch trigger generator. The 18 diagnostics signals are multiplexed onto three dual-beam oscilloscopes, using analogue delay lines for the signals and digital circuitry that retriggers the oscilloscopes and provides a "staircase" signal for separating the traces. The system sacrifices some signal fidelity in exchange for compact displays.

Single-Module Tests

We assembled each module and installed it in a test stand, separate from the amplifier facility, to test the generation of single elec-

tron beams. We tested each Marx generator by connecting it to a CuSO_4 resistive load and firing a sequence of pulses at voltages up to 4/3 of the designed operating point. We then connected the Marx generator to its pulse-forming line and tested the complete module, including the formation of the electron beam at its design size. We then fired a minimum of 25 shots at 450 kV (the design output voltage), a minimum of 50 shots at 15% above 450 kV, and finally, 10 more shots at 450 kV. We considered operation of a module to be acceptable only when we could complete this sequence without any component failures.

During the shot tests, we monitored the electron-beam energy using a calorimeter located at the anode plane. We also monitored the output voltage of the Marx generator; the voltage pulse, which triggered the output switch; and the output voltage and current of the pulse-forming line. After testing the first three modules, we installed them in the B amplifier and tested them, as a group before completing the assembly of the remaining three modules.

Three-Module Tests

We tested each of the three-module groups separately to simplify the checkout procedure.

Fig. 7-101. Output-switch trigger system.

Each set of three output switches is triggered from a signal generated by a single spark gap located at the rear of each set; we allowed this gap to self-break in the three-module tests. The trigger inputs for the three Marx generators are electrically coupled so that they can be triggered from a common source; if one prefires, it will trigger the other two in the group and thus maintain the integrity of the amplifier. The output-switch trigger gap has a midplane electrode driven by the RC circuit shown in Fig. 7-101; with this design, the output-switch trigger will break down and trigger the output switches even if it does not receive a command from an external source.

As soon as three modules were operable, we addressed the issue of interactions between parallel modules. The current delivered by three modules (270 kA) is twice the critical current for beam self-pinching in this geometry. To prevent the spatial collapse of the beam, we provided each module with a separate cathode and used 3-mm-thick stainless steel plates to connect the anode to the pulse-forming lines between the separate cathodes. These current paths effectively isolated each module from magnetic fields produced in adjacent modules.

To check their parallel operation, we exposed a piece of lithographic paper to the entire electron beam produced by three modules. The resulting image showed no signs of beam collapse or, indeed, any interaction between adjacent beams. The lithographic paper records only the time-integrated beam profile; to determine if the beams were collapsing individually in time, we sampled the local current density with resistive collectors at several locations, including points near the beam edge. The measurements showed little change in

beam size during the 150-ns output pulse, though there was some indication of beam expansion early in the pulse near the edge of the beam.

During the checkout of three-module groups, we measured the efficiency of the conversion from electrical energy stored in the Marx generator to electron-beam kinetic energy by comparing the energy delivered to a full-aperture carbon calorimeter (located at the diode anode plane) with the stored energy computed from the Marx charge voltage and measured capacitance. Using this method, we measured 19.1 ± 0.3 kJ of beam energy at an efficiency of $70 \pm 2\%$. This efficiency is fairly good; the overall "wall-plug" efficiency is partly due to the pumps and heaters for the vacuum and water systems, which dominate power consumption for the complete power system.

The charging system operates at constant current rather than constant voltage, so losses in the Marx charging circuit are small ($<2\%$); we made no attempt to measure or optimize the efficiency of the charging power supply. A large part of the lost 30% of the Marx energy is residual charge in the Marx generators due to a capacitance mismatch between the Marx and water lines. The losses in voltage profiles are also unnecessarily large. The conversion efficiency of Marx to electron-beam energy could therefore be increased to 80 to 85% at little risk, if that were required.

Full-System Tests

We had originally intended to synchronize the two groups of three modules by laser-triggering the two spark gaps that produce the output-switch trigger pulses. This

technique is complicated and unnecessary for pulsed-power system tests that do not require synchronization to a laser system. Instead, we connected the two midplane electrodes provided for fault protection to a common overvolted spark gap, with the cables replacing the capacitors of the RC circuit shown in Fig. 7-101. This approach not only permitted pulsed-power tests, but worked so well that we decided to use it in the planned laser experiments.

We measured the timing of all six modules with respect to the common trigger gap. The rms jitter of the poorest module was 3.9 ns; the remaining five modules had jitters of less than 2.5 ns. Previous measurements of the rms jitter of a single output switch yielded 1.9 ns under these conditions. The average time difference between the earliest module output and the latest output was 11 ns, which is considerably greater than the 5 ns delay introduced by the distribution of output-switch trigger pulses. We do not know if this additional delay is the result of physical differences in the switches or of electrical differences in the output-switch trigger circuits.

Gas Experiments

In anticipation of laser experiments, we performed several measurements involving the interaction of the electron beams of the full system with gas in the laser cell. Our principal intent was to verify that sufficient plasma return current would be induced in the gas ionized by the beam to cancel the beam current, thus reducing the self-generated magnetic field and preventing beam self-pinching within the laser cell. We filled the laser cell with mixtures of argon and oxygen and photographed the fluorescence from the auroral transition in oxygen. The photographs showed no evidence of serious pinching, though the spatial distribution of fluorescence agreed only qualitatively with the results of calculations of energy-deposition profiles. The outcome is shown graphically in Fig. 7-102; the values plotted have been normalized to the average value over the entire volume. The total energy deposited in the gas, as inferred from the gas pressure rise, Δp , measured with an electronic transducer, agreed with calculations for pressures of 1.0 and 1.5 atm, but we had difficulty interpreting the results

Fig. 7-102. Energy-deposition profiles taken inside a gas-filled laser cell.

at 2 atm, which require further study. At 1.5 atm, the measured Δp yields 11.5 ± 1 kJ of energy deposited in the gas. Doubling the partial pressure of oxygen from 10 to 20 Torr, to eliminate the possibility of burnout, did not change the results.

Summary

We have completed the construction and characterization of the pulsed-power system for the RAPIER-B amplifier. We have verified the generation of >36 kJ of electron-beam kinetic energy at an efficiency of 70%. System timing is more than adequate for laser operation, and the energy deposited in the gas-filled laser cell appears to be sufficiently uniform.

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Rare-Gas-Halide Kinetics

Kr₂F⁺ Trimer Kinetics in Rare-Gas-Buffered Kr-F₂ Laser Mixtures

Last year, we reported studies of the formation and quenching kinetics of the Kr₂F⁺ and Ar₂F⁺ rare-gas-halide (RGH) trimers. During 1982, we extended those studies with an examination of Kr-F₂ gas mixtures buffered with lower-Z rare gases. Using the

Fig. 7-103. Peak KrF^* and Kr_2F^* fluorescence signals vs calculated energy deposition for Kr-F_2 , He-Kr-F_2 , Ne-Kr-F_2 and Ar-Kr-F_2 gas mixtures.

Fig. 7-104. Kr_2F^* decay rate, τ^{-1} , vs pressure for a 11 600:100:5 He:Kr:F_2 mixture at low current density ($\sim 33 \text{ A/cm}^2$).

experimental technique described last year,¹⁰¹ we measured the KrF^* and Kr_2F^* fluorescence signals from 20:1 Kr:F_2 mixtures buffered with helium, neon, or argon and excited at low current density ($\sim 33 \text{ A/cm}^2$). To simplify the comparison of the performance using different buffers, we chose the partial-pressure ratios to yield the same energy deposition per unit volume from the 600-keV electron-beam pulse as a "standard" 150:1 Kr:F_2 mixture would yield at identical F_2 partial pressures, as calculated using the stopping powers of Berger and Selzer.¹⁰² Our procedure resulted in partial-pressure ratios of $\text{He:Kr:F}_2 = 1897:20:1$, $\text{Ne:Kr:F}_2 = 422:20:1$, and $\text{Ar:Kr:F}_2 = 246:20:1$.

Both the KrF^* and Kr_2F^* fluorescences from all of the buffered mixtures were qualitatively similar to those from the Kr-F_2 binary mixtures described last year.¹⁰¹ Figure 7-103 shows the peak dimer and trimer signals as a function of the average energy-deposition rate for each three-component gas mix, along with the signals from the standard Kr-F_2 mixture. We

assumed the energy deposition to be twice that indicated by the stopping powers of Berger and Selzer,¹⁰² for a single pass, based on detailed numerical calculations made using SANDY1, a three-dimensional deposition code.¹⁰³

The low-current-density analysis for binary gas mixtures that we reported last year¹⁰¹ was applied this year to the fluorescence signals from the buffered mixtures to determine the changes in the KrF^* - Kr_2F^* kinetics with the various buffer gases. We determined best-fit values for the trimer decay rate, τ^{-1} , for each of the gas mixtures at a variety of pressures. The results for the helium-buffered case are plotted in Fig. 7-104; the solid line is the result of a least-squares fit of the decay rate to a form linear in the gas pressure. In the helium-buffered case and all other cases, the intercepts were consistent with a 220-ns radiative lifetime for Kr_2F^* .

For our quenching studies, we assumed that the pressure-dependence of τ^{-1} was given by

$$\tau^{-1} = \tau_{\text{rad}}^{-1} + k_{\text{Kr}} n_{\text{Kr}} + k_{\text{buf}} n_{\text{buf}} + k_{\text{F}_2} n_{\text{F}_2}, \quad (87)$$

where n represents the concentration of the various gas species. Using the values of k_{Kr} and k_{F_2} given last year,¹⁰¹ linear fits to the values of the decay rates yielded values for the rate constants for the quenching of Kr_2F^* by the buffer gases. In all cases, the pressure-dependent quenching was dominated by its fluorine contribution ($> 80\%$), allowing us to obtain only rough estimates for the values of k_{buf} . For helium, we

calculated $k_{\text{He}} \sim (2 \pm 1) \times 10^{-14} \text{ cm}^3 \text{ s}^{-1}$; for neon and argon, we could only place an upper bound of $10^{-13} \text{ cm}^3 \text{ s}^{-1}$ on k_{Ne} and k_{Ar} .

The values of the trimer-formation rate, α , determined in the analysis of the fluorescence data from all of the buffered-gas mixtures, as well as that from the standard Kr-F₂ mix, are plotted vs the total gas pressure in Fig. 7-105. These data were used to calculate estimates of the relative values of the three-body trimer-formation rate constants in the various buffer gases. For helium and neon, we determined that $k_{\text{He}} \sim (0.25-0.6)k_{\text{Kr}}$ and $k_{\text{Ne}} \sim (0.5-0.7)k_{\text{Kr}}$, respectively. The argon data yielded no such simple result due to the increased complexity of the dimer-trimer gas kinetics.

KrF Laser Performance in Buffered-Gas Mixtures

The performance of a KrF laser is, in general, a function of the electron-beam current density, the gas pressure and mixture, and the cavity output coupling. As a partial solution to determining the maximum possible system efficiency, we chose to examine the variation of the KrF laser output with the Kr and F₂ partial pressures for gas mixtures buffered either with 1500 Torr of argon or with 11 600 Torr of helium at two current densities (160 and 440 A/cm²). With argon as the buffer, the laser cavity consisted of a maximum-reflectivity mirror and

a 50%-reflectivity mirror. When the buffer gas was helium, however, a 70%-reflectivity output mirror was required to produce reasonable gain and thus ensure a sufficiently rapid intensity buildup in the cavity.

Figure 7-106 presents the variation of laser output with krypton and fluorine pressure in a helium-buffered mixture pumped at 160 A/cm². Also shown are the maximum

Fig. 7-105. Kr_2F^* formation rate, α , vs gas pressure at low current density ($\sim 33 \text{ A/cm}^2$).

Fig. 7-106. KrF* laser output and peak KrF^* and Kr_2F^* fluorescence signals vs (a) F_2 pressure and (b) Kr pressure. ∇

Buffer gas (Torr)	Current density (A/cm ²)	Kr pressure (Torr)	F ₂ pressure (Torr)	Maximum efficiency (%)
Ar (1500)	160	100	2.5	0.9
Ar (1500)	440	300	6	4.5
He (11 600)	160	50	4	0.4
He (11 600)	440	200	7	1.3

Table 7-19. Maximum calculated efficiencies of buffered KrF.

KrF* and Kr₂F* sidelight-fluorescence signals measured under identical conditions, but with the mirrors removed from the gas cell. In Fig. 7-106(a), we see that the laser output and KrF* fluorescence reach a peak at approximately the same F₂ pressure. The lower value of the fluorine pressure at the peak Kr₂F* fluorescence is presumably due to the greater significance of fluorine quenching for the RGH trimer than for the dimer due to the longer radiative lifetime of the former. In Fig. 7-106(b), the peak of the laser output occurs at a higher value of the krypton pressure than of the KrF* fluorescence signal. Such an effect is probably the result of the presence in the medium of optical absorbers (perhaps the helium 2³S metastable state), whose density is reduced at the larger krypton pressures. The Kr₂F* fluorescence rises linearly with krypton pressure; this is consistent with a three-body mechanism for the formation of the RGH trimer from the dimer, as discussed above.

The variation of the laser output with the gas pressures for the helium-buffered mixtures at 440 A/cm² and the argon-buffered mixtures at both current densities were qualitatively the same as that shown in Fig. 7-106. These results are summarized in Table 7-19, which lists the maximum calculated KrF laser efficiencies (based on our estimate that the energy deposition of the electron-beam equals twice the single-pulse value), along with the krypton and fluorine pressures at which these values were attained. The efficiencies using helium as the buffer gas were lower than those achieved using argon, as the result of two causes: the reduced efficiency in transferring the energy from the excited helium atoms to the KrF* dimer (due to the fact that the excitation and ionization energies for helium are greater than those for argon); and the lower output coupling used for the helium measurements.

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Free-Electron Laser

The free-electron laser (FEL) directly converts the kinetic energy of a relativistic electron beam to laser radiation and, therefore, has the potential for becoming a very efficient, minimum-maintenance ICF driver. System designs for a fusion-class FEL were given in the 1979 and 1980 Annual Reports.^{104,105} During 1982, in cooperation with LLNL's M Division (Magnetic Fusion) and the DOE Office of Basic Energy Sciences, we initiated an experimental program to gather preliminary data on FEL performance. This program includes construction of an electron laser facility that will be connected to and will use the electron beam produced by the experimental test accelerator (ETA).¹⁰⁶

An ICF driver must operate at high pulse energies. In view of this requirement, we have restricted our studies to master-oscillator, power-amplifier (MOPA) laser architectures. To achieve high extraction efficiency, we have adopted the tapered-wiggler concept suggested by Kroll, Morton, and Rosenbluth.¹⁰⁷

The basic scaling relations for a high-gain tapered-wiggler FEL are given by

$$\lambda_L \propto \frac{\lambda_w}{\gamma^2} \left[1 + \frac{1}{2} \left(\frac{e B_0 \lambda_w}{2 \pi m c} \right)^2 \right] \quad (88)$$

$$G \propto (B_0 \lambda_w \lambda_L)^2 \quad (89)$$

$$\eta \propto \frac{I^{1/2} \gamma}{\lambda_w B^{1/2}} \quad (90)$$

$$r \propto \lambda_w^2 \left(\frac{I}{B} \right)^{1/4} \quad (91)$$

and

$$P \propto \frac{I^{5/2}}{\lambda_w^3 \lambda_L^2 B^{5/2}} \quad (92)$$

where λ_L is the laser wavelength; λ_w is the wiggler wavelength; γ is the electron energy in units of the electron rest mass; η is the deceleration, $(\gamma_{\text{initial}} - \gamma_{\text{final}})/\gamma_{\text{initial}}$; I is the electron-beam current; B is the

electron-beam brightness $= I/(\gamma\epsilon)^2$; B_w is the peak wiggler magnetic field; ϵ is the electron-beam emittance (no π); G is the gain; r is the electron-beam radius; and P is the required laser input power.

If we desire laser operation at short wavelengths ($\sim 0.25 \mu\text{m}$), Eq. (88) restricts us to high electron-beam energies for reasonable wiggler periods (10 cm) and magnetic fields (1 to 5 kG). Equation (89) indicates that a high-gain, short-wavelength FEL requires a very bright (high current, low emittance) electron beam. Equation (90) states that large extraction efficiencies also require high currents, although Eq. (92) implies that, at high currents, the required input power may become unreasonable. Equation (91) relates the required electron-beam radius to the wiggler and electron-beam parameters. The FEL operates best at high intensities and small beam radii.

As the above equations show, lower-energy electrons are required for longer-

wavelength FELs. We also expect higher gain at longer wavelengths. These considerations, plus the accessibility of the ETA, led us to design our initial experiments as a millimeter-wavelength FEL operating at 35 to 140 GHz. The long wavelength necessitates the use of a waveguide to guide the microwave radiation; otherwise, the system should be an excellent test bed for FEL physics.

The electron-laser facility (ELF) is attached to the end of the ETA beam-propagation region and consists of three sections: a beam-conditioning region, an interaction region, and a microwave-diagnostic section (Fig. 7-107). The beam-conditioning section reduces the current in the 4.5-MeV ETA beam from 8 to 1 kA and reduces the emittance from roughly 100 to 30 mrad-cm. The beam-conditioning section also has an energy filter that allows an energy band (4.5 to 4.2 MeV) to pass into the interaction region. In the interaction region,

Fig. 7-107. ELF beamline.

a periodic wiggler field transverse to the beam direction is established with electromagnets; the electron beam propagates in a rectangular waveguide. The beam and radiation execute a single pass through the interaction region, and the electron beam is then swept out of the rectangular

waveguide by a transverse magnetic field. The microwave radiation propagates down the waveguide into the microwave-diagnostics region, where the pulse shape, total energy, and frequency spectrum are measured. Each of these ELF regions are discussed in detail below.

Beam Conditioning

To reduce the ETA current and emittance to a level acceptable for FEL operation, the beam is passed through a 2-m-long tube inside a solenoidal magnetic field. The tube has an inner diameter of 2.5 cm and continuously removes particles having an unacceptably large transverse velocity. At the end of the 2-m tube, the beam current is reduced to 1 kA with an emittance of 30 mrad-cm. The beam then passes through two quadrupole doublets separated by 2 m of vacuum; a slit aperture placed midway between the two doublets passes only electrons with energy between 4.5 and 4.2 MeV. This filtering is based on the energy-dependent focal properties of the quadrupole fields.

The behavior of the beam during conditioning is monitored using five beam bugs¹⁰⁸ to measure current and beam position in the beam tube. The beam profile is monitored with a TV monitor and an x-ray wire probe. Two vacuum-pumping stations maintain the tube pressure at better than 5×10^{-9} Torr. The total length of the beam-conditioning section is approximately 9 m.

Before it can enter the interaction region, the beam tube must be made rectangular to mate with the FEL waveguide. In the circular-rectangular transition piece at the end of the beam-conditioning region, a 4- μm -thick aluminum foil is placed at 45° to the beam axis. This foil is transparent to the electron beam, but serves as a reflector to the microwave input signal injected from below (Fig. 7-108). From this foil on, the microwave signal and electron beam propagate together.

Interaction Region

The interaction region is composed of a rectangular waveguide immersed in a transverse, spatially periodic magnetic field

Fig. 7-108. Transition piece between the ELF beam-conditioning and interaction regions.

Fig. 7-109. Schematic of an individual power supply, and a typical current pulse, for the ELF wiggler magnetic field.

called a wiggler. The wiggler, which provides the pump wave for the FEL, has a period of 9.8 cm and is 3 m long (with 30 full periods). Fringe fields of rectangular solenoids placed in a cusp configuration provide the transverse magnetic field in the interaction region. An input signal of the desired laser output frequency is injected at the beam entrance to the interaction region and copropagates with the beam through the interaction region; this input signal is the perturbation that is amplified in the laser cavity.

With the use of separate power supplies, the wiggler field strength can be varied along the interaction region; in particular, it can be tapered to maximize the output energy of the FEL (Ref. 109). Figure 7-109 shows one of the 15 pulsed-power supplies that drive the 30 wiggler periods (2 periods per power supply). The current pulse through the wiggler coil, which is essentially DC to the 40-ns electron-beam pulse, lasts for 1 ms; like the electron beam, the pulse can be run at a 1-Hz repetition rate (a typical current pulse is shown in Fig. 7-109). The shot-to-shot variation in pulse amplitudes is less than 1% to ensure reproducibility of the experiment.

The amplitude of the magnetic field on the input and output periods of the wiggler is tailored to match the beam into and out of the interaction region. This matching forces the electron beam to oscillate about an axis that coincides with the center of the interaction waveguide. The field from the wiggler magnets focuses the electron beam in the vertical dimension parallel to the wiggler field; this focusing is due to the z -component of the magnetic field from the wiggler coils as the beam moves off axis. Quadrupole magnets surround the wiggler coils to provide focusing in the plane perpendicular to the wiggler field.

The interaction waveguide immersed in the wiggler field is rectangular and oversized. The walls of the waveguide cavity are made of thin stainless steel to ensure good penetration of the magnetic field. We will use two sizes of interaction waveguide in the ELF: a standard WR-229 (S-band) waveguide; and a nonstandard rectangular guide measuring 2.91 by 9.83 cm. A schematic of the magnet of the wiggler coil structure is shown in Fig. 7-110. As can be seen from the orientation of the magnetic

field, the electrons oscillate parallel to the wide dimension of the waveguide. Hence, the electric field of the radiation is parallel to the waveguide width, and the transverse electric mode, TE_{10} , is excited by the microwave pulse injected at the beginning of the interaction region.

Fig. 7-110. Cross section of the interaction region showing orientation of the wiggler field and electron oscillation.

Diagnostic Section

At the exit of the interaction region, the electron beam passes through a rectangular current monitor and is swept out of the waveguide by a transverse magnetic field. For initial experiments, the electron beam will be dumped on the waveguide wall. Future experiments call for an electron spectrometer to determine the energy lost by the electron beam to the radiation field.

Beyond the current monitor at the wiggler exit, the waveguide is a standard WR-229 copper guide; a laser calorimeter located at the end of this 6-m waveguide section measures the total microwave energy. A sidewall coupler permits a fraction of the radiation to pass to a crystal detector and microwave spectrometer. The crystal detector monitors the temporal behavior of the microwave pulse and, coupled with the information from the calorimeter, gives the instantaneous radiated power. The microwave spectrometer gives the frequency

bandwidth of the radiation and, with interchangeable gratings, the harmonic context. Simulations predict that the ELF amplifier, in an optimized configuration, will have a gain of roughly 10^4 and a millimeter-wave output power of up to 500 MW.

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Atomic and Molecular Theory

Coherence Effects in Charge-Transfer Collisions

There has been much interest recently^{110,111} in collisional processes involving atoms excited by a laser that is tuned to a resonance transition of the target atom. If the atomic transition is $s \rightarrow p$, for example, and the incident ion is in an S state, the products of the collision will depend on the polarization of the laser light relative to the collision axis due to different initial fluxes in the excited Σ and Π electronic states. This behavior has been observed by Schmidt et al.¹¹⁰ in $\text{Na}^+ + \text{Na}^* \rightarrow \text{Na}^* + \text{N}_2$.

In these experiments, it is assumed that an incoherent superposition of substates with a known¹¹⁰⁻¹¹² ratio of populations is generated by pumping the target atoms for times that are very long compared to the spontaneous lifetimes of the excited states. In this case, effects related to Rabi oscillations in the state populations are expected to be damped out. If shorter pumping times could be utilized, however, it would be possible to create a coherent superposition of the initial sublevel populations. We would then be able to observe additional changes in the various final product states of the scattering intensities; these additional variations will be due to the coherent properties of the excited radiation.

MacGillivray and Standage¹¹³ have observed directly the oscillations corresponding to Rabi oscillation in the resonance fluorescence of Na atoms excited by an intense single-mode laser. We believe it

should be possible also to observe oscillations in charge-exchange cross sections if the excited atoms collide with incident ions before they can decay radiatively and if the effects are large enough. Since the duration of a collision is very short compared to the period of the Rabi oscillations for moderate laser fields (10^{-13} vs 10^{-8} s), in our scattering calculations we ignored changes in the excited-state populations due to the laser.

We studied the $(\text{NaLi})^+$ collisional system, with either the Li or the Na atom as the target atom. The Na target atom is more attractive experimentally due to the availability of a laser at its resonance frequency; it is somewhat less attractive theoretically, however, since a higher degree of electronic excitation is involved and more molecular states are required to accurately describe the transitions involved. The initial state of the $(\text{NaLi})^+$ system consists of a Na ion incident on a Li atom that is being resonantly pumped by a laser to its 2p state. In our calculations we ignored spin-orbit couplings and the spin of the nucleus, so that the excited p state correlates to a single Σ molecular state and a single Π molecular state. The generalizations required to include these effects are straightforward^{111,112} and not crucial to the present discussion.

We define the initial preparation of the initial target state in terms of two angles, θ_R and θ_p . The first of these, θ_R , is defined such that $\sin^2 \theta_R$ is the fraction of target atoms in the excited state and $\cos^2 \theta_R$ is the fraction of target atoms in the ground state. This is the phase of the Rabi oscillation between the ground and excited atomic states in the case of adiabatic, coherent laser pumping of the resonance transition and in the absence of damping processes such as collisions or spontaneous emission by the upper state. For an incoherent pumping process, θ_R cannot exceed $\pi/4$ —that is, no more than half the target atoms can be maintained in the upper states.

The second angle, θ_p , defines the relative populations of p states with projections parallel to or perpendicular to the space-fixed z axis. Thus, we define the fraction of excited atoms that have an electron in the p_z orbital to be $\cos^2 \theta_p$; those having an electron in the p_x orbital are defined by $\sin^2 \theta_p$. We assume that the orbital perpendicular to the

scattering plane, p_z , is not occupied in the excitation process; since this orbital is antisymmetric with respect to reflections in the scattering plane, it plays no part in the scattering process either.

Finally, we must allow for the fact that the coupling matrix elements have been determined in terms of the body-fixed (rotating) coordinate system. In the asymptotic region, the excited state will propagate freely (in the space-fixed frame) up to some radius at which the angular coupling term between the p_z and p_x states in the rotating frame becomes smaller than the difference between the excited Σ and Π adiabatic potential curves. For internuclear distances smaller than this radius, the electronic angular momentum will be coupled to the molecular axis. Grosser¹¹⁴ has recently provided an analysis of this phenomenon, establishing that the integration begins outside the molecular region; consequently, we must include an impact-parameter-dependent factor in defining our initial state in terms of the molecular wave functions.

Figure 7-111 shows the coordinate system for our scattering calculations using the straight-line impact-parameter method. The arrow through the target atom (designated P) indicates the orientation of the laser-excited initial state; the z axis is the space-fixed direction of the incident ion, and b is the impact parameter. The figure shows that, for a finite impact parameter and a finite initial value of z at which we begin the integration of the coupled equations, the polarization of the excited state created in the lab frame must be rotated through the angle α , defined by $\tan \alpha = b/z_0$, to define the polarization of the excited state in the molecular frame.

Including all the factors described above, and defining the four molecular states by their asymptotic limits,

$$|1\rangle = \text{Li } (2s) + \text{Na}^+, \quad (93)$$

$$|2\rangle = \text{Li}^+ + \text{Na } (3s), \quad (94)$$

$$|3\rangle = \text{Li } (2p_y) + \text{Na}^+, \quad (95)$$

and

$$|4\rangle = \text{Li } (2p_x) + \text{Na}^+, \quad (96)$$

Fig. 7-111. Coordinate system for the scattering calculations with the straight-line impact parameter method.

we can express the initial state of the laser-excited target as

$$\begin{aligned} |i\rangle = & \cos \theta_R |1\rangle |n\rangle \\ & + \sin \theta_R [\cos (\theta_P + \alpha) |3\rangle \\ & + \sin (\theta_P + \alpha) |4\rangle] |n-1\rangle, \quad (97) \end{aligned}$$

where $|n\rangle$ is the number density of photons in the laser field. Of course, since α depends on the impact parameter, there is an implicit dependence of this initial state on b .

Since the coupled equations to be solved are linear, we can solve them for four linearly independent initial conditions and then find the cross sections for arbitrary combinations of θ_R and θ_P by constructing the proper linear combinations of the separate amplitudes. Under conditions such that the initial state includes a coherent superposition of substates, the appropriate interference terms must be included in evaluating the cross sections. We have defined a_{ij} to be the j th solution of the coupled equations [Eq. (98)] with initial condition $a_{ij}(-\infty) = \delta_{ij}$. Then the total cross section from a coherently excited initial state is

$$\begin{aligned} \sigma_{12}^L(E) = & 2\pi \int_0^{\infty} b db \left\{ \cos^2 \theta_R |a_{12}|^2 \right. \\ & + \sin^2 \theta_R [\cos^2 (\theta_P + \alpha) |a_{32}|^2 \\ & + \sin^2 (\theta_P + \alpha) |a_{42}|^2 + \cos (\theta_P + \alpha) \\ & \times \sin (\theta_P + \alpha) (a_{32}^* a_{42} + a_{42}^* a_{32})] \left. \right\}. \quad (98) \end{aligned}$$

In the case that the excitation is not coherent, the last cross term in Eq. (98) will not appear. We obtained accurate potential-energy curves and nonadiabatic-coupling matrix elements from ab initio all-electron configuration-interaction calculations.

Atomic and Molecular Theory

Due to the long range and strength of the coupling between the lowest two states, the charge-exchange cross section from the ground state is very large over the energy range of interest (300 to 3000 eV). The coupling between the charge-exchange state, $|2\rangle$, and the two states that correlate to the excited target, $|3\rangle$ and $|4\rangle$, is smaller, and we find that the cross sections from these states are approximately an

Fig. 7-112. Charge transfer cross sections from the ground and excited states of Li to Na^+ as functions of the collision energy.

Fig. 7-113. Energy dependence of the cross sections from the (laser) excited p-states of Li to Na^+ for both coherent and incoherent initial states.

Fig. 7-114. Energy dependence of the superelastic cross sections for $\text{Li}^* + \text{Na}^+ \rightarrow \text{Li} + \text{Na}^+$.

order of magnitude smaller than those from the ground state. Figure 7-112 shows the charge-transfer cross sections from these three states. Our results agree reasonably well with the previous calculations of Melius and Goddard¹¹⁵ and with the experiments of Daley and Perel.¹¹⁶

It is obvious from these results and from Eq. (98) that, if the laser pumps a substantial fraction of the target atoms into the excited state, the observed total cross section will decrease dramatically. When the ground-state and excited-state populations are approximately equal ($\theta_R \approx \pi/4$), the total charge-exchange cross section is smaller almost by a factor of 2; for a complete inversion, the cross section will be smaller by an order of magnitude. The cross section for the excited state, which can be measured by a time-of-flight analysis of the scattered particles, is obtained from Eq. (98) by setting $\theta_R = \pi/2$, such that

$$\sigma_{12}^*(E) = 2\pi \int_0^\infty b db \left[\cos^2(\theta_p + \alpha) a_{32}^2 + \sin^2(\theta_p + \alpha) |a_{42}|^2 + \cos(\theta_p + \alpha) \times \sin(\theta_p + \alpha) (a_{32}^* a_{42} + a_{42}^* a_{32}) \right] \quad (99)$$

where again the last term only appears when the initial state is a coherent superposition. Figure 7-113 shows the excited-state cross sections, σ^* , both including and excluding the interference term, as functions of collision energy for a particular value of θ_p . The laser light is assumed to be linearly polarized, with the polarization vector at an angle, θ_p , relative to the direction of the incident ion beam.

We also calculated the superelastic scattering cross section for the process $|\text{i}^*\rangle \rightarrow |\text{i}\rangle$, where $|\text{i}^*\rangle$ indicates the initial state defined in Eq. (97) by $\theta_R = \pi/2$. To measure this cross section, we performed a time-of-flight analysis of the scattered ions to select those whose velocities increase due to the quenching of the excited target. The superelastic cross section is given by

$$\sigma_{11}^*(E) = 2\pi \int_0^\infty b db \left[\cos^2(\theta_R + \alpha) a_{31}^2 + \sin^2(\theta_R + \alpha) |a_{41}|^2 + \cos(\theta_R + \alpha) \times \sin(\theta_R + \alpha) (a_{31}^* a_{41} + a_{41}^* a_{31}) \right] \quad (100)$$

where the last term is appropriate only for coherently prepared excited states. Figure 7-114 shows the superelastic cross sections for both coherently and incoherently excited initial states as functions of the collision energy.

In conclusion, we have verified the occurrence of large, observable differences in total and partial cross sections for coherent and incoherent initial states. Although we have not discussed the effects here, we further expect that coherent preparation of the initial state would lead to pronounced differences in differential cross sections and in the angular dependence and polarization of light emitted from excited product species.

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Electronic Structure of Metal Hexafluorides

In support of research on advanced solid-state lasers, we are investigating the electronic structure of metal and actinide hexafluorides. We have so far carried out ab initio calculations for the clusters VF_6 , UF_6 , UF_6^+ , PuF_6 , PuF_6^+ , and PuF_6^{2+} . Our goal is to determine the excitation energies and transition moments for these clusters, as well as the ground- and excited-state potential surfaces. Our results will allow us to assign the optical spectra, predict band intensities, and provide theoretical data for the improved parameterization of our crystal-field-fitting programs.

The calculations for VF_6 are designed to model the optical properties of the vanadium ion in a MgF_2 crystal host. For VF_6 , we have used a large basis set and have considered each electron explicitly. The results to date show the importance of including the effects of the lattice in stabilizing the VF_6 complex.

For UF_6 and PuF_6 , the high nuclear charge required that the core electrons be replaced with an effective potential that included the relativistic effects.¹¹⁷ For UF_6 , we have used the uranium core potential and basis set given by Hay et al.¹¹⁸ The occupied one-electron orbitals were determined for the LS coupled 1A_1 ground-state wave function from Hartree-Fock calculations on the valence electrons; the lowest unoccupied orbitals were determined using an

improved virtual orbital method.¹¹⁹ The molecular orbitals were used to calculate electric dipole transition probabilities between the one-electron levels.

Of course, a proper description of the electronic states requires that we describe the significant electrostatic and spin-orbit effects among the f-electrons. Both of these interactions produce a strong mixing of the $1a_{2u}$ and $2t_{2u}$ f-like orbitals. At this time, we have included the spin-orbit coupling using a one-electron model.¹²⁰ This approach yields a $4\gamma_{7u}$ relativistic orbital as a linear

Transition	f
$9\gamma_{8g} \rightarrow 12\gamma_{8u}$	—
$11\gamma_{8g} \rightarrow 4\gamma_{7u}$	3.8×10^{-3}
$(10\gamma_{8g}, 11\gamma_{8g}) \rightarrow 13\gamma_{8u}$	2.8×10^{-2}
$(10\gamma_{8g}, 11\gamma_{8g}) \rightarrow 5\gamma_{7u}$	7.1×10^{-2}

Table 7-20. Oscillator strengths for UF_6 .

Fig. 7-115. One-electron model of the UF_6 ligand-to-5f transitions. ▼

combination of the a_{2u} and $2t_{2u}$ orbitals and, similarly, the $13\gamma_{8u}$ one-electron wave function as a mixture of the $2t_{2u}$ and $5t_{1u}$ orbitals. The ligand-metal charge-transfer transitions of interest originate from the t_{1g} ($10\gamma_{6g}$, $11\gamma_{8g}$) ligand orbital (the very small spin-orbit splitting of this orbital was ignored). The results for the one-electron model are shown in Fig. 7-115. The oscillator strengths for UF_6 are summarized in Table 7-20, which roughly reproduces the experimental ratio of the transitions into the $4\gamma_{7u}$ and $13\gamma_{8u}$ orbitals.

Many areas require further investigation before we can have confidence in our metal hexafluoride calculations. Among these are the following issues

- Proper inclusion of spin-orbit effects for d^n and f^n configurations when $n > 1$.
- Accurate calculation of the relative energies of the various electronic states; the calculation requires the inclusion of electron correlation by configuration interaction.
- Extension of the basis sets to include functions appropriate for negative ions, and the inclusion of lattice effects for our investigation of the VF_6^{-1} complex.

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Reduction of Degenerate Two-Level Excitation to Independent Two-State Systems

The nondegenerate two-level atom (or, more precisely, two-state atom) provides an important model in treating time-dependent excitation¹²¹ and serves as the basic starting point for more elaborate studies of atomic and molecular excitation by lasers. The presence of magnetic sublevels can complicate the appearance of the dynamical equations,¹²²⁻¹²⁶ but, during 1982, we demonstrated that, even with arbitrary, elliptically polarized light and with degenerate sublevels, the coherently excited near-resonant decay-free two-level atom can be treated as a set of independent two-state systems.

Figure 7-116 illustrates the reduction of interest. Figure 7-116(a) shows the dipole-transition linkages present with elliptically polarized excitation (a combination of ΔM

$= +1$ and $\Delta M = -1$, where M is the magnetic quantum number) for transitions between levels having angular momenta $J = 2$ and $J = 1$. In the absence of radiative decay, the dynamics comprise two independent systems: a three-state Lambda system (dashed lines) and a five-state M system (solid lines). Figure 7-116(b) is a system dynamically equivalent to Fig. 7-116(a), obtained by introducing new degenerate basis states for the five-state linkage; Fig. 7-116(c) is also dynamically equivalent to that of Fig. 7-116(a), but it is obtained by introducing new basis states for both the Lambda and M systems. We can always introduce a new basis of upper-level and lower-level states (replacing the basis of magnetic sublevels), within which the dynamics follow a pattern appropriate to the one- and two-state linkages of Fig. 7-116(c).

The linkage patterns of Figs. 7-116(a) and 7-116(c) are familiar from expressions of linear polarization in terms of a quantization axis along the propagation direction [Fig. 7-116(a)] or along the electric vector [Fig. 7-116(c)]. The present, more general, transformation applies to excitation with arbitrary multipolarity and polarization as well as to arbitrary degeneracy. Although we term the system of interest a degenerate two-level atom, the transformation applies to a more general class of multilevel systems.

The problem of interest requires solution of a time-dependent Schrödinger equation,

$$i \frac{d}{dt} c_i(t) = \sum_j H_{ij} c_j(t) \quad (101)$$

for the probability amplitudes, $c_i(t)$. The Hamiltonian matrix, H , has the structure

$$|H| = \begin{bmatrix} h_1 & V \\ V^\dagger & h_2 \end{bmatrix} \quad (102)$$

where h_1 and h_2 are diagonal submatrices, and V denotes the linkages between sublevels of level 1 (having angular momentum J_1) and sublevels of level 2 (with angular momentum J_2) originating in the laser-driven excitation.

In the rotating-wave approximation (RWA), the elements h_i represent detunings, and V is a matrix of (slowly varying) Rabi

Fig. 7-116. Treatment of a coherently excited near-resonant decay-free two-level atom as a set of independent two-state systems.

frequencies, proportional to the electric field strength and to appropriate Clebsch-Gordan coefficients. For our method, we need only assume that h_1 and h_2 are (possibly time-dependent) scalar multiples of unit matrices, and that the ratio of any two elements of V is time-dependent.

We can always introduce a unitary transformation on the basis states,

$$c' = U c \quad \text{with } U^\dagger U = 1 = U U^\dagger \quad (103)$$

such that, in the new basis, the amplitudes satisfy the matrix equation,

$$i \frac{d}{dt} c' = H' c' \quad \text{with } H' = U H U^\dagger \quad (104)$$

This is equivalent to introducing a new set of basis states.

We now describe our system for reducing the degenerate two-level system to independent two-state systems. First we obtain unitary matrices, A and B , that diagonalize $V V^\dagger$ and $V^\dagger V$

$$A V V^\dagger A^\dagger = a^2 \quad \text{and} \quad B V^\dagger V B^\dagger = b^2 \quad (105)$$

We then remove the null subspaces to obtain $\bar{A}$ and $\bar{B}$. Finally we form the matrix $\bar{B}^\dagger$, such that

$$\bar{B} = \bar{B}^\dagger \bar{B} V^\dagger (\bar{a})^{-1} \quad (106)$$

Now the new basis states are

$$c' = \bar{A} c \quad \text{for level 1} \quad (107)$$

and

$$c' = \bar{B} c \quad \text{for level 2} \quad (108)$$

The nonzero linkages (i.e., Rabi frequencies) between these new basis states are the elements, $\bar{a}$. The transformation of basis states does not alter the resonant nature of the ex-

citation; the matrices h_i remain null matrices after the transformation.

The resulting dynamical simplification can be considerable. Instead of dealing with an $N \times N$ Hamiltonian matrix, where $N = (2I_1 + 1) \times (2I_2 + 1)$, we need only deal with n independent two-state systems, where n is at most the lesser of $2I_1 + 1$ and $2I_2 + 1$. In general, each two-state system has a different Rabi flopping frequency, so that the overall population dynamics do not consist of a simple sinusoidal Rabi flopping.

The two-state reduction discussed here differs from the reduction of a sequentially linked multilevel system to an equivalent two-level system; that reduction occurs when intermediate levels are not resonant or when intermediate Rabi frequencies are very large.¹²⁷

Although the foregoing discussion deals with linkage patterns arising from excitation of degenerate magnetic sublevels by polarized light, the proof of reduction of a two-state system makes no use of the rotational-symmetry properties of the interaction, V . In particular, it does not require that the elements of V be related in any way to angular momentum operators. Furthermore, h_i need not be real (and hence Hermitian), but may include any imaginary term to express probability loss due to ionization or dissociation. Thus, our analysis applies to the well-known¹²⁸⁻¹²⁹ Lambda system of Fig. 7-117(a), as well as to the other multilevel linkages shown in Fig. 7-117.

In the illustrated systems, a set of n distinct lasers excite a set of n lower levels into a common upper level. Each laser is resonantly tuned to the appropriate transition frequency, so that, in the RWA (Ref. 130), the elements of h_i are degenerate. Such systems are dynamically equivalent to the system of Fig. 7-117(c), in which only a single lower sublevel connects with the excited state. Thus, if the ground sublevels initially have equal population, then only $1/n$ of the ground level population can ever reach the

Fig. 7-117. (a) Three-state and (b and c) other multilegged linkage patterns for resonantly tuned lasers.

excited state and subsequently ionize. The remainder of the population remains trapped in a coherent superposition of ground levels.

If the lasers are detuned from resonance, the elements of h , no longer are equal (to zero), and the lack of degeneracy prevents transformation to a two-state system—the ionized fraction can exceed $1/n$. Alternatively, if different lasers have different temporal behavior, then we no longer have the conditions necessary for application of our transformation (i.e., time-independent ratios of V -matrix elements); again, the ionized fraction can exceed $1/n$ (sequential access can yield complete ionization.) For further details of this analysis, see Ref. 131.

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Optical Gating in Resonant Multiphoton Excitation

The behavior of N -state systems excited coherently by near-monochromatic radiation has been studied for several years¹³²; the relevant dynamics obtain from solutions of the time-dependent Schrödinger equation in the rotating wave approximation (RWA). It is now generally recognized that the coherent flow of probability along a chain of N coupled, near-resonant transitions proceeds most efficiently when the successive Rabi frequencies appropriate to the transitions are roughly equal. Einwohner, Wong, and Garrison have pointed out¹³³ that the regular flow of probability along a chain can be severely curtailed by the presence of a strong resonant branching transition, and they presented analytical expressions for the special case, $N = 2$, to support this observation. During 1982, we considered this branched system from the viewpoint of dressed states. We observed that the hin-

dered flow of probability is a consequence of the familiar AC-Stark effect. Further, we saw consequences not noted in the analysis of Einwohner et al.

Consider the chain of excitation transitions shown in Fig. 7-118: a sequence of states (nondegenerate energy levels) linked by near-resonant excitation. Population flows along this chain, beginning at time $t = 0$, from an initial concentration in state 1. We can imagine monitoring an ionization or fluorescence signal to register time-dependent arrival of population into the terminus of the chain, state N . We assume near-resonant, relaxation-free excitation; at each step the cumulative detuning is much less than the Rabi frequency for the step. At some point in the chain (state X), a competing transition or chain of transitions provides an alternative excitation route, to state Y . If all the main-sequence transitions (denoted by solid lines in Fig. 7-118) have comparable Rabi frequencies (i.e., coupling strengths), and if the branched transition (shown in dashed lines) has a much smaller Rabi frequency, then the excitation dynamics of the N -state chain are essentially unaffected by the branch. Population flows out of (and subsequently returns into) the initially populated level, state 1, arriving at (and departing from) the terminal state, N , in a manner appropriate to N strongly coupled (or dressed) states.

When the Rabi frequency, b , of the branch to state Y becomes comparable to the Rabi frequency, c , of the principal chain [Fig. 7-118(a)], then we can no longer regard the branch linkage as a negligible perturbation; we must then explicitly diagonalize the Hamiltonian of the $(N + 1)$ -state system. In so treating the $(N + 1)$ -state system, we recognize that linkages c and b form a generalized Lambda system¹³⁴⁻¹³⁶ in which a portion of population will become locked into a coherent state, excluding excitation into states $X + 1, X + 2, \dots, N$. If $c = b$,

Fig. 7-118. Linkage patterns for branched multistep excitation chains; full lines form the principal chains; dashed lines show side branches.

then, as in the Lambda system, only half the population can ever reach state N .

A particularly interesting situation occurs when the branching Rabi frequency far exceeds the adjacent chain Rabi frequency ($b \gg c$). Under this condition, almost no population at all flows into states X , $X+1$, $\dots$, N ; the strong transition, b , acts as a switch to turn off population flow into states above the linkage.

The nature of this switching can be understood easily if we regard states X and Y , the two strongly coupled states, as forming a dressed two-state system upon which the remaining chain acts as a perturbation. In this view, the (relatively) weak linkage, c , serves as a probe of the dressed X - Y system. The absorption rate as a function of probe frequency shows the well-known double-peaked Autler-Townes or AC-Stark

doublet structure of two peaks separated by the Rabi frequency, b . Thus, if the c transition remains tuned to the resonant Bohr frequency while the b transition is turned on, we find that the tuning falls midway between the Autler-Townes doublet, and, hence, the effective absorption rate is quite small. Because the absorption-profile wings fall exponentially with detuning, b need only exceed c by a factor of 2 or so to diminish the chain excitation rate by an order of magnitude.

The branch transition X - Y thus acts as a gate to population flow along the chain. There need never be any actual population in state X ; the dressed X - Y states occur even in the absence of population. The blockage occurs when we impose a branch Rabi frequency that appreciably exceeds the adjacent chain Rabi frequency. The gating

action does not occur instantaneously, of course; roughly one complete X - Y Rabi cycle must be completed to define the Autler-Townes doublet. Furthermore, the population blockage is only temporary, as Einwöhrner et al. point out; the imposed detuning of state X does not prevent resonant two-photon transitions between states $X - 1$ and $X + 1$, thereby bypassing the gate. The two-photon Rabi frequency decreases in direct proportion to the imposed detuning, so that the gating action can readily be made to hold for many Rabi periods of the X - Y transition.

Obviously the gating action of a single branch extends to multiple branches, as is shown in Fig. 7-118(b). Here, any one of the branches, if strongly driven, can exclude population from states further along the chain. Note, however, that the presence of a second strong, resonantly tuned transition in the side chain [Fig. 7-118(c)] negates the gating action. In other words, gating occurs for an odd number of steps in the side chain, whereas an even number of steps induces no strong gating effect. The explanation of this curious odd-even effect, noted earlier in a slightly different context,¹³⁷ appears quite natural if we regard the side chain as a strongly coupled M -state system. When we sweep the frequency of the major-chain transition, c , now acting as a weak probe, we observe an M -peaked spectrum. When M is an odd integer, one of these peaks occurs at zero detuning, and, hence, the major chain remains on reso-

nance" at the unperturbed Bohr frequency, even in the presence of the side chain. Of course with an odd-state chain, the population flow does not simply move without impediment along the original chain—we have a modified Lambda system.

As the dressed-atom picture makes clear, the gating action of a side chain only holds if the main-chain excitation frequencies remain tuned to the Bohr frequencies, as is appropriate to excitation in the absence of a side chain. If we instead retuned the excitation frequency to the relevant Autler-Townes multiplet, then population can flow upward (as well as into the side chain).

The Autler-Townes splitting of the branch transition provides a limiting bound on the intensity of amplified stimulated emission (ASE) into any branch of the excitation chain; if the branch Rabi frequency exceeds the chain Rabi frequency, then the branch transition no longer satisfies the resonant condition necessary for continued stimulated emission of population out of the chain sequence.

Author: B. W. Shore

Coherence in the Quasi-Continuum Model

In treating excitation of molecules by a laser, we often encounter a situation in which the laser excites one particular mode that, in turn, couples via nonradiative interactions to numerous other modes lying close in energy to the excited mode. Figure 7-119(a) shows the interaction linkage pattern for an excited state linked by the laser transition to a bound state (the solid line) and by internal couplings to numerous excited states (the dashed lines).

When we diagonalize the nonradiative Hamiltonian, all these excited states form a quasi-continuum¹³⁸ (QC) of densely spaced states, all radiatively linked to the ground state. Figure 7-119(b) shows the linkage pattern after diagonalization of internal couplings. Note that the ground state now has laser interactions with each level of the quasi-continuum. Since the dipole transition moments linking the QC states with the

Fig. 7-119. Linkage patterns (a) before, and (b) after, diagonalization of internal couplings.

ground state are obtained by diagonalization, the phases of adjacent states may vary quite irregularly, even though the transition moments may differ very little in magnitude.

It has sometimes been assumed that the quasi-randomness of matrix elements would act to destroy coherence. We have now shown that the assumption is not correct. The dynamics of excitation are completely unaffected by these phases and, in fact, we can always assume that all transition dipole moments are real and positive.¹³⁸ Thus, diagonalization of internal molecular couplings to produce a photoabsorption quasi-continuum does not produce an incoherent absorption coefficient, but maintains coherence in dipole moments.

Author: B. W. Shore

Low-Cost Pulsed Power

Future laser systems will require significant reductions in the cost of providing millisecond pulses of very high power to laser flashlamps. To achieve the low-cost goal, we have explored two energy-storage methods in detail

- Conventional capacitive energy storage, such as we have used for Novette.
- Rotational (inertial) energy storage with high-power mechanical-to-electrical conversion. Rotational stores include the active rotary flux compressor¹³⁹⁻¹⁴² (ARFC)

invented at the University of Texas (UT) in Austin, and homopolar generators.¹⁴³

We have also briefly studied a third option: chemical energy storage with explosive magnetohydrodynamic energy conversion. Table 7-21 summarizes the results of our studies by showing estimated hardware costs for these competitive systems.

Capacitive storage technology is the most highly developed system and enjoys the lowest risk. Furthermore, we believe that a considerable cost saving can be realized for capacitive storage through the use of large system elements that reduce the parts count. Capacitors are also competitive with the best projected costs of rotational stores; thus, further development of large ARFCs or homopolar generators is not advantageous for lasers that require only occasional discharges.

Nevertheless, we believe that the ARFC will be the most cost-effective future solution for powering a large repetitively pulsed laser; for example, ARFCs are projected to have higher mechanical-to-electrical power-conversion capabilities (up to 600 kW/kg) than any other known device. During 1982, we began computer modeling of large-scale ARFCs, and our work has yielded two optimum designs, which we discuss later in this article.

As part of our analysis of rotational stores, we briefly studied a circuit utilizing the homopolar generator (HPG). Its discharge time (~ 0.1 s) is not fast enough, however, to directly power flashlamps; the energy would have to be transferred to a large inductor before it could be delivered

	Capacitive store		Rotational store		Explosive store, XMHD
	Modular system (Nova)	Monolithic system (Zeus)	ARFC	HPG	
Energy storage	4.3	2.5	3.5	2.0	0.5
Peripheral hardware	6.4 ^a	3.7 ^a	2.5 ^b	3.5 ^c	1.5 ^d
Flashlamps	2.7	1.0	1.0	1.0	1.0
Installation	0.9	0.8	0.5	0.5	0.5
	14.3	8.0	7.5	7.0	3.5
Development required	None	Low	Moderate	High	High

^aIncludes bank circuitry, switch system, power supply, pulse-forming inductor, and energy transmission.

^bIncludes prime mover, start-up bank and associated hardware, and energy transmission.

^cIncludes prime mover, inductive energy store, commutated switch with pulse-forming bank, and energy transmission.

^dIncludes explosive isolation, energy-storage and transmission cable, and pulse-forming and start-up bank.

Table 7-21. Specific cost estimates for installed hardware to power very large Nd:glass amplifiers (c/j - 1982 dollars).

Low-Cost Pulsed Power

to the flashlamps with a commutated fast-recovery switch. Thus, the energy transfer and switching technique must be developed before such a system could become operational.

The third area we examined, explosive magnetohydrodynamics (XMHD), is potentially the lowest-cost method yet proposed for producing a single millisecond pulse of very high power. A study by ARTEC Associates, Inc.,¹⁴⁴ suggests that the initial capital hardware cost for an XMHD system may be less than one-third that for any competitive power-producing method. Immature development, and therefore high risk, is its chief drawback; the XMHD technique will not likely be considered unless the laser is to be built at a remote site.

Low-Cost Capacitive Energy Storage

The monolithic capacitive energy-storage method, utilizing large blocks of energy and large system components, has been adjudged to be the optimum method to power very large Nd:glass lasers. We are now designing large monolithic systems that will succeed the modular capacitor bank array presently designed for Nova.

Discussions with vendors revealed that the lowest-cost capacitors will probably store 30 to 50 kJ in each package and will sell for less than \$0.03/J. These capacitors will provide 0.5-ms, nonreversing discharges with a life expectancy of $\sim 10^4$

shots. Capacitors carrying up to 250 kJ can be bussed together and connected to a main bus with a single fuse (see Fig. 7-120). The energy attached to the main bus would depend upon the switch. For example, the present size-D ignitrons can switch at least 1 MJ. Size-E ignitrons are as yet untested but they should readily handle 2 to 4 MJ. Solid-dielectric switches, which operate by multichannel breakdown through a replaceable dielectric element, can be made in virtually unlimited size.

We are seeking cost reductions in conventional capacitive energy-storage by assembling the elements of the system into larger packages. The modular approach, while convenient, will be replaced with megajoule-size capacitor banks, large switches, power supplies with a rating of several MVA, and large interconnecting cables or buswork. Work is also under way to increase the size of each flashlamp load, because fewer connections and fewer parts will mean a lower cost.

Rotational Energy Storage

During 1982, we made a preliminary investigation of gigajoule-size ARFCs because the cost for large electrical equipment tends to scale as the 0.8 power of machine mass. In other words, the specific cost of one large machine would be roughly 60% of the cost of 10 smaller ones having the same total mass. The problem with drum-type ARFCs, however, is that the mass scales nearly as

Fig. 7-120. Design for a monolithic capacitor bank.

d^3 (where d is the rotor diameter), but the output power scales as only d^2 , since the power is taken out as shear on the rotor surface. The rotor length is also fixed because the rotor becomes vibrationally unstable if it is longer than two or three times its diameter. Because of this limitation, we decided that gigajoule-size drum-type ARFCs are not feasible.

We also performed cursory calculations for disk-type ARFCs because, with such machines, both the mass and the power output scale as d^2 . Therefore, the power-to-mass ratio would stay approximately constant with size, and we might take advantage of the lower specific cost for larger mass scaling. However, immature development has diverted us from this path; no disk machine has ever been designed, much less modeled. Furthermore, no computer model exists, and a model will be difficult to develop because of asymmetry in the windings.

As a result of these preliminary studies, we decided to look at large drum-type ARFCs in the energy range of 1 to 20 MJ. We discovered that, within this range, larger is probably better, provided that the rotor laminations could each be made in one piece and not segmented (segmented rotor laminations are not as strong as those stamped from a single sheet). A thorough canvas of the suppliers of sheet silicon steel of the grade we would require revealed that this material can be obtained in widths up to 42 in.; thus, 42 in. became the maximum dimension in our calculations.

Our computer modeling of rotary flux compressors culminated in optimized designs for two machines. The first would weigh ~ 13.3 tonnes and would provide enough energy to power the flashlamps in nine of Nova's 46-cm amplifiers with a 600- μ s pulse. For comparison, each of those nine amplifiers presently requires a 600-kJ capacitor bank. The other machine design would yield the largest ARFC that can be made without segmenting the rotor; it would weigh ~ 30 tonnes and drive 24 of the large Nova amplifiers with a 900- μ s pulse.

The larger machine would feature a 42-in.-diam rotor and a larger, segmented stator. The smaller machine has a stator

with a maximum diameter of 42 in., so that both rotor and stator could be stamped from one piece. Each of these two machines was optimized to pump the most possible energy into Nd:glass, via flashlamps, in a single pulse. A decay constant of 0.33 ms was assumed for the glass, so that the resulting pulse widths were less than 1 ms. The features of the two optimized designs are summarized in Table 7-22; the power output pulses and current pulses are given in Figs. 7-121 and 7-122, respectively.

Fig. 7-121. Calculated power output pulses for two drum-type ARFC designs.

Fig. 7-122. Calculated current pulses for two drum-type ARFC designs.

Optical Switch

	Large machine (42-in. rotor)	Small machine (42-in. stator)
Equivalent capacitor energy	14.44 MJ	5.40 MJ
Start-up capacitor energy	2.13 MJ	0.80 MJ
Rotor mass	15.23 tonnes	6.84 tonnes
Machine mass	29.54 tonnes	13.27 tonnes
Specific energy	488.8 J/kg	406.3 J/kg
Specific power	539.0 kW/kg	600 kW/kg
FWHM current pulse width	0.897 ms	0.589 ms
Peak flashlamp power	15.92 GW	7.96 GW
Average shear stress	26.71 MPa (3874 psi)	27.04 MPa (3922 psi)
rpm	2706	3600
Tip speed	151.2 m/s	150.8 m/s
Poles	8	8
Conductors per pole	2	2

Table 7-22. Parameters of optimum ARFCs.

Note that the equivalent capacitor energy is 14.44 MJ for the larger machine and 5.4 MJ for the smaller machine. These are the sizes of the capacitor banks that would be required to produce the same pump level in Nd:glass. Note also that a capacitor bank of almost 15% of this equivalent output energy is required to produce start-up current in the machines. Thus, our ARFC designs are "capacitive energy amplifiers" and will not entirely replace capacitors. The same is true for any efficient alternative scheme to power flashlamps; all must use a capacitor bank of from 5 to 25% of the size needed to completely replace the conventional capacitor energy-storage method scheme (this includes both the HPG and XMHD methods).

The smaller of the two machines is our choice for the optimum device because of its compact size, transportable mass, and simplicity of construction (stator segmentation is not required). We are confident that this machine can be configured to produce a continuous repetition rate when the need arises for this application.

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Optical Switch

Introduction

Transient, volumetric thermal gratings are formed by the interference of two coherent beams in an absorbing medium.¹⁴⁵⁻¹⁴⁹ Ab-

sorption of a modest amount of energy from the intersecting beams can drastically change the optical properties of the medium. A probe beam of a different wavelength, for which the medium is transparent, is strongly scattered by volumetric phase grating. Some authors refer to the effect as forced scattering.^{148,149} The probe beam can be aligned to the Bragg angle of incidence, with respect to the volumetric grating, so that a large fraction of probe power is scattered into the first order of the grating. For the proper choice of grating parameters, the fraction approaches unity.

The effects of forced scattering by the laser-generated gratings can be used to construct a large-aperture switch with low insertion loss and high switching efficiency. Such a device could find uses in a large high-power laser system for prepulse suppression, target retropulse isolation, parasitics control, and, possibly, output switching from regenerative amplifiers.

Figure 7-123 illustrates a possible layout of the switch. In our example, we discuss a case in which UV beams are used to switch a visible or near-infrared beam, since there are many fluid media that have absorption in the near-UV and are transparent at longer wavelengths. A cell filled with an appropriate fluid is placed in the path of the visible beam. With AR coatings on the cell windows, the insertion loss can be minimized. When the cell is illuminated by crossing UV beams, a volumetric grating is formed, and the visible beam is efficiently deflected.

Hargrove¹⁴⁷ proposed the use of a reflecting volumetric thermal grating in a liquid as an optical switch. However, our deflecting switch has several important practical advantages, including a lower switching laser-energy requirement and the possibility of using gaseous media.

Our experimental demonstration of an optical switch used a KrF laser to generate thermal gratings in liquid chloroform and in a gas mixture of xenon and CF₃I. The media were chosen because they have appropriate absorption depths at the KrF laser wavelength, transparency in the visible region, and relatively large ratios of index of refraction to heat capacity. We used continuous He-Ne and pulsed dye lasers to monitor the transient phase gratings.

Previous research on thermal gratings has concentrated on liquid media because the physics are reasonably well understood, and no problems arise in demonstrating the switching effect in chloroform. However, most liquids require significantly more UV energy input than gases, and nonlinear effects limit the peak intensity that they can transmit. A switch with a gaseous medium would be more useful, but gases have undesirable effects as well. Nonlinear acoustics and finite thermalization length severely limit the useful operating regime. Nonetheless, with a proper choice of working parameters, high-quality switching was observed in gases.

Optical Properties of Volumetric Phase Gratings

Light scattering by volumetric gratings is a classical problem that has been solved with different approximations by investigators using ultrasound,¹⁵⁰ thick-emulsion holograms,¹⁵¹ phase conjugation by degenerate four-wave mixing,^{152,153} and stimulated Raman scattering (SRS) with multiple pump beams.¹⁵⁴ Our general approach uses coupled-wave theory.^{150,151} Thus, we can calculate the parameters for which the thermal grating acts as an efficient switch.¹⁵⁵

Figure 7-124 shows the allowed regions in the parameter space for a 2-cm-thick cell. Optical properties of the switch are due to practical alignment tolerance and switching into higher grating orders.

Fluid Dynamics

The phase grating that is used as a switch is formed in the medium following the deposition of energy by the heating pulse. We calculate the energy required to create a grating with $\nu = \pi/2$.

The equations of motion for a fluid with small perturbations (neglecting heat conduction) can be combined into a linear differential equation for local density¹⁵⁶

$$\frac{\partial^2 \rho}{\partial t^2} - v_s^2 \nabla^2 \rho = - \frac{v_s^2}{\rho_0 c_p} \frac{\partial \rho}{\partial T} \int_p^{\lambda^2} dt Q(r,t), \quad (109)$$

Fig. 7-123. Beam-deflecting switch based on a volumetric thermal grating.

Fig. 7-124. Allowed and forbidden regions in the parameter space for transient volumetric thermal gratings.

where v_s is the speed of sound and $Q(r,t)$ is the local power deposition. In this case, we can write

$$Q(r,t) = q(y,t) (1 - \cos Kx), \quad (110)$$

and by defining

$$\rho(r,t) = \rho_0 + R(y,t) \cos Kx, \quad (111)$$

Optical Switch

we obtain a simple equation for a driven harmonic oscillator

$$\frac{\partial^2 R}{\partial t^2} + \omega^2 R = A \int_0^t q|t'|dt' \quad (112)$$

with

$$\omega = Kv_\lambda \quad (113)$$

and

$$A = \frac{\omega^2}{\rho_0 c_p} \left(\frac{\partial \rho}{\partial T} \right)_p \quad (114)$$

Equation (112) can be readily solved for two extreme cases: short- and long-pulse excitation. For a long heating pulse ($\omega\tau_p \gg 1$), the first term on the left can be neglected, and

$$R(t) = \frac{\int_0^t q|t'|dt'}{\rho_0 c_p} \left(\frac{\partial \rho}{\partial T} \right)_p \quad (115)$$

which is the expected answer for isobaric heating.

After a short-pulse excitation, the solution is

$$R(t) = \frac{\int_0^t q|t'|dt'}{\rho_0 c_p} \left(\frac{\partial \rho}{\partial T} \right)_p (1 - \cos \omega t) \quad (116)$$

The peak magnitude of R will be reached at time $t = \pi/\omega$, and, for the same input energy, R will be twice as large as that of a long pulse.

Now, the density change must be related to the change in the index of refraction. In general,

$$\Delta n = \left(\frac{\partial n}{\partial T} \right)_p \Delta T + \left(\frac{\partial n}{\partial \rho} \right)_T \Delta \rho \quad (117)$$

However, for the media we considered, only the second term is important. The coefficient can be obtained by differentiating the Lorentz-Lorenz formula

$$\left(\frac{\partial n}{\partial \rho} \right)_T = \frac{(n-1)(n+1)(n^2+2)}{\rho_0 6n} \quad (118)$$

For gases ($n \sim 1$), only the first factor needs to be considered.

For an efficient switch,

$$\left| \int \Delta n(y) dy \right| = \frac{\lambda}{2} \quad (119)$$

Therefore, the energy per unit area required in short-pulse excitation is

$$E = \frac{\lambda}{4} \frac{-\rho_0 c_p}{\left(\frac{\partial \rho}{\partial T} \right)_p \left(\frac{\partial n}{\partial \rho} \right)_1} \quad (120)$$

For xenon at atmospheric pressure,

$$\rho_0 c_p = 0.8 \text{ J/cm}^3 \text{ K} \quad (121)$$

$$\left(\frac{\partial \rho}{\partial T} \right)_p = \frac{-\rho_p}{T} \quad (122)$$

and

$$\left(\frac{\partial n}{\partial \rho} \right)_1 = \frac{7 \times 10^{-4}}{\rho_0} \quad (123)$$

Thus, for $\lambda = 0.6 \mu\text{m}$, the required energy is $E = 5.5 \text{ mJ/cm}^2$.

For liquid chloroform,

$$\rho_0 c_p = 1.54 \text{ mJ/cm}^3 \text{ K} \quad (124)$$

$$\left(\frac{\partial \rho}{\partial T} \right)_p = -1.3 \times 10^{-3} \rho_0 / \text{K} \quad (125)$$

$$\left(\frac{\partial n}{\partial \rho} \right)_1 = \frac{0.51}{\rho} \quad (126)$$

and $E = 37 \text{ mJ/cm}^2$.

Not all of the absorbed laser energy is converted into heat. For xenon gas with the CF_3I absorber, the following reaction occurs

Approximately half the energy is liberated in the form of kinetic energy of the fragments. Some additional heat is liberated at later times as fragments recombine.¹⁵⁷ The hot fragments ($\sim 1 \text{ eV}$) undergo approximately 10 collisions with xenon atoms before the energy is thermalized. Assuming a random walk, the range for each fragment

is $\sim 0.3 \mu\text{m}$. Thus, formation of gratings with $\Lambda < 1 \mu\text{m}$ will be very inefficient in gas at STP. Although a higher gas density can be used in the cell, such use is not practical for large-area optical switches. In liquids, the conversion of absorbed laser energy into heat has comparable inefficiency; however, the thermalization range is negligible due to the higher density.

It is possible, especially for gaseous media, to drive the system into a nonlinear regime, where Eq. (109) is no longer valid. Nonlinear optical effects (such as scattering of the UV beams by the thermal grating) and nonlinear acoustics cause nonsinusoidal deposition of energy. Such deposition, in turn, further drives the system into the nonlinear regime. Although rigorous treatment of the problem is beyond the scope of our work, we can make a qualitative estimate of boundaries of the undesirable part of the parameter space. For a relatively large grating period ($\Lambda \gg 2\pi v_s \tau_p = 22 \mu\text{m}$ for xenon), nonlinear optics manifest themselves through generation of higher-order UV beams. Numerical solutions of nonlinearized fluid-dynamic equations¹⁵⁸ with slightly nonsinusoidal initial-temperature profiles start showing significant deviation from linear solutions approximately when

$$\frac{\Delta \rho}{\rho} > 5\% \quad (128)$$

The above criterion limits the maximum allowed absorption coefficient, as shown in Fig. 7-124.

For a small grating period, the fraction of energy scattered into higher orders drops off sharply with Λ . In addition, because of the finite pulse length, the excitation efficiency of acoustical waves drops off. Thus, larger perturbations of density are tolerable in this regime. The price is that more energy is required, by up to a factor of 2, as can be seen from Fig. 7-125.

Experiments

The KrF laser that was used to generate thermal gratings consisted of an oscillator and a regenerative amplifier¹⁵⁹ with typical output energy of 0.5 J in $\sim 20\text{-ns}$ pulses and with a frequency bandwidth of $\sim 0.1 \text{ cm}^{-1}$. The laser output was geometrically split

into two beams, which were recombined on the absorption cell, as shown in Fig. 7-123. The input energy was varied with calibrated attenuators and measured with a Scientech calorimeter.

A He-Ne laser aligned to the Bragg angle was used to observe the time-dependent behavior of the thermal grating formed. A small fraction of KrF laser output was split off from the main beam and delayed by $\sim 70 \text{ ns}$ to generate a short-pulse probe beam. The delayed beam was used to pump a simple dye laser, operated free-running with Rh6G at $\sim 0.58 \mu\text{m}$.

In our first experiment, a 1-cm-wide absorption cell was filled with liquid chloroform and probed by the He-Ne laser. The UV beams were crossed at $\sim 4 \text{ mrad}$. A pulsed red beam was readily observed exiting the cell in the expected direction, corresponding to the first order of the thermal grating. A photomultiplier was used to detect the switched beam. Figure 7-126 shows typical oscilloscope traces for different UV input-energy densities.

Fig. 7-125. Relative grating-formation efficiency for laser pulses of finite duration.

Fig. 7-126. Deflection of a $0.63\text{-}\mu\text{m}$ beam by a thermal grating formed by a KrF laser in chloroform.

Optical Switch

Approximately 50% of the input energy was absorbed by the chloroform in the cell.

The switched beam was strongly modulated in time. Modulation was due to excitation of a standing acoustical wave by the short heating pulse. The period of modulation was equal to the grating spatial period divided by the speed of sound. At the UV incident fluence ($\sim 150 \text{ mJ/cm}^2$), the grating switching efficiency reached maximum and locally approached 100%. The overall switching efficiency is determined by the uniformity of UV illumination.

The behavior of the switch in chloroform was reproducible for tens of shots, after which transmission of the visible beam started to degrade due to buildup of some unknown absorbing and scattering suspension in the liquid. It is likely that the chloroform was being decomposed by the UV radiation.

To demonstrate the switching effect in a gas, a 2-cm-wide absorption cell was filled with a mixture of 1-atm xenon and 50-Torr CF_3I gas; the latter is a thoroughly studied

absorber of UV light because of its application in photolytic iodine lasers. The 50 Torr of CF_3I in our experiment gives an absorption of e^{-1} in the cell. The UV beams were crossed at 8.4 mrad. High-quality switching was again observed. Significantly less UV energy was required than in the previous case, as shown in Fig. 7-127. The peak switching efficiency approached unity, and the signals were reproducible.

A dye-laser beam, timed to the peak of switching efficiency, was also deflected by the switch. To test the angular sensitivity of the switch, the dye-laser beam was set to expand, covering a total angle of ~ 8 mrad. Figure 7-128 shows the resulting output-beam profiles. Consistent with theoretical predictions, part of the expanding beam (~ 1 -mrad wide) that was close to the Bragg angle was switched out, while the rest of the beam was transmitted. The dark patch in the transmitted beam confirms high switching efficiency ($>90\%$).

The energy required to switch the visible beams was $\sim 15 \text{ mJ/cm}^2$ of UV absorbed in the cell. Considering the inefficiency of conversion of absorbed photons to thermal energy ($\sim 50\%$), $\sim 8 \text{ mJ/cm}^2$ of heat was delivered by the pulse, a value close to the theoretical prediction of 6 mJ/cm^2 .

Analysis of the hydrodynamics associated with formation of a thermal grating (see "Fluid Dynamics," earlier in this article) shows that, for a certain choice of parameters, nonlinear behavior is expected. To verify our prediction, the experimental parameters were changed from point A to point B on Fig. 7-126. In practice, we increased CF_3I pressure to 100 Torr and decreased the crossing angle to 4 mrad.

As before, a switched beam exited the absorption cell after each KrF pulse. However, the switching efficiency was poor, and the time dependence, shown in Fig. 7-129, was irregular and nonreproducible. Attenuation of UV beams eventually led to more regular modulations, which occurred for very low switching efficiencies. The nonlinear acoustics render operation of the switch in this regime impossible.

Discussion

We have demonstrated sufficient understanding of transient thermal gratings to

Fig. 7-127. Intensity of a $0.63\text{-}\mu\text{m}$ beam deflected by a thermal grating.

Fig. 7-128. Propagation of an expanding 0.58-μm beam through a volumetric phase grating.

design an efficient optical switch. A proper choice of operating parameters allows the use of a gaseous medium that can transmit high peak powers at visible and near-infrared wavelengths. Assuming that the high-power laser will operate at fluences of $\sim 10 \text{ J/cm}^2$, the total UV laser energy required for the switch is 3 orders of magnitude lower than the main laser-output energy.

We can further reduce that fraction by providing another source of energy for the switch (e.g., electrical or chemical) and by using the UV laser to initiate striated, spatial energy deposition. More research is required before such techniques can be perfected.

The UV laser requires a fairly high degree of spatial and temporal coherence for the geometry used in our experimental demonstration. High coherence increases the complexity and cost of the switch; however, different geometries can reduce the coherence requirement. For example, the necessity for crossing two beams to generate an interference pattern can be eliminated by using a flat-transmission grating, which can be illuminated by a single beam and the image of the grating then projected into an absorption cell, as illustrated in

Fig. 7-129. Intensity of a 0.63-μm beam deflected by a density grating in a gas.

Fig. 7-130. Thus, the laser is only required to provide uniform average illumination of the grating, and the divergence is $\sim 1 \text{ mrad}$. This can be accomplished with a simple multimode oscillator and a divergence several hundred times greater than the laser used for our experiments. The bandwidth requirement is determined only by the achromatism of the imaging optics. It may

Fig. 7-130. Alternative geometry for generation of thermal gratings in a less coherent UV laser.

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Fig. 7-131. Volumetric thermal-grating switch used for target isolation.

Fig. 7-132. Deflecting switch in a regenerative amplifier.

be possible to design a system that uses a bright flashlamp as the source of UV light in such a scheme.

An illustration of how the switch could be used for target isolation is shown in Fig. 7-131. The UV absorbing cell, which acts as a deflector, is placed just before the final focusing lens. Without UV illumination, the laser output (ASE) is focused away from the target. Even for a small switching angle of 10 mrad with the focal length from 10 to 20 m, the displacement of the focal spot is from 10 to 20 cm. The UV illumination is turned on at the appropriate time, causing peak grating strength at the time of the main-laser output. The beam is deflected onto a target. Since the grating strength oscillates in time with a period of ~50 ns, it is easy to arrange the flight time so that the pulse reflected from the target is not deflected, but, rather, is scraped off by the laser system's spatial filters.

Another application is illustrated in Fig. 7-132. The volumetric grating can be used as a switch out of (or into) a regenerative amplifier. It is reasonable to extrapo-

late switching angles up to 200 mrad. With a 2-m laser aperture, 10 m of propagation are required to clear the optics. The distance is small compared to the length of the entire laser system.

In our experiment, the beams were switched only by small angles (10 to 20 mrad). Larger switching angles may be needed for practical applications. The regime of small Λ is very close to the nonlinear regime. Further work is needed before we can predict how much of the parameter space is usable.

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Temperature and Stress in a Pumped Glass Slab

Temperature rises and stresses are produced when Nd-doped glass slabs are used as the gain medium in a repeatedly pulsed laser amplifier. Because the absorption of light and deposition of heat are largest near the faces and smallest in the center of each slab, the heat transport is more efficient than in a uniformly heated slab, and the stress is lower. The stress per unit of deposited energy thus depends on the heat-deposition profile, which, in turn, depends on the doping. We have therefore calculated the temperature-induced stress as a function of doping for doped-glass slabs.

We approximate the glass as an infinite plane-parallel slab, with the heat deposition depending only on the dimension through the slab. We assume the pump pulse repetition rate is rapid compared to the thermal relaxation time, and we calculate steady-state results. Pumping is assumed to be equal on both sides, so that the deposition is symmetric about the slab center. If the elastic and thermal constants of the glass are independent of temperature, the standard rules for stress¹⁶⁰ give

$$\sigma_{xx} = \sigma_{yy} = \frac{\alpha E}{1 - \nu} \left(-T + \frac{1}{2h} \int_h^h T dx \right. \\ \left. + \frac{3x}{2h^3} \int_h^h T x dx \right), \quad (129)$$

where α is the coefficient of linear thermal expansion, E is Young's modulus, ν is Poisson's ratio, T is the temperature, $2h$ is the slab thickness, and x is the distance from the center of the slab. By the assumption of symmetry, the third term inside the parentheses of Eq. (129) is zero. The first and second terms inside the parentheses combine to give the difference between the temperature at a single point and the average temperature of the slab as a whole. We know the slab is warmer in the center and, thus, is in compression in the center and in tension near the faces. Since glass is strong in compression but weak in tension, we concentrated on the area of largest tension, at the surface of the slab.

To find the surface tension, we need to find the temperature distribution, $T(x)$. Since the problem is linear, we superimpose the effect of many thin slices of heat deposition. Heat deposited in a thin slice at position x^1 produces (in steady state) a uniform thermal gradient to both faces, and the resulting temperature distribution is

$$dT(x) = \begin{cases} \frac{(h + x^1)(h - x^1)}{2kh} dQ, & x \geq x^1 \\ \frac{(h + x^1)h - x^1}{2kh} dQ, & x \leq x^1 \end{cases}, \quad (130)$$

where k is the thermal conductivity and Q is the heat-deposition rate.

To find $T(x)$, we integrate the heat-deposition profile, $Q(x^1)$, multiplied by Green's function above, to find

$$T(x) = \int_h^{-h} \frac{(h + x^1)(h - x^1)}{2kh} Q(x^1) dx^1 + \int_h^0 \frac{(h + x^1)h - x^1}{2kh} Q(x^1) dx^1, \quad (131)$$

assuming $T(\pm h) = 0$.

The one-sided deposition in glass to an optical depth, u , is conveniently described by the standard double-exponential expression,

$$d(u) = A[B e^{-Cu} + (1+B) e^{-Du}], \quad (132)$$

where $u = \rho x$ is the number of ions/cm² in the optical path (ρ is the ionic density), and A , B , C , and D are constants that depend on the specific glass type. Differentiating this

expression, and adding a symmetric term from each side of the slab, we have

$$Q(x) = \rho A \left\{ B C e^{-Cu} + e^{-Du} + (1+B) D e^{-Cu} + e^{-Du} \right\}. \quad (133)$$

We then normalize the results by the total deposited heat,

$$H = 2A \left\{ B(1 - e^{-2\rho h}) + (1+B)(1 - e^{-2\rho h}) \right\}. \quad (134)$$

Carrying out the integration, we find that the temperature is

$$T(x) = \frac{A}{K\rho C} \left[1 - e^{-Cu} + e^{-2\rho h} - e^{-C(2h-u)} \right] + \frac{A}{K\rho} \left[\frac{(1-B)}{D} (1 - e^{-Du}) + e^{-2\rho h} - e^{-D(2h-u)} \right], \quad (135)$$

so that the stress is

$$\sigma = \frac{\alpha E}{(1-\nu)K} \frac{A}{\rho} \left\{ \frac{B}{C} \left[(1 + e^{-2\rho h}) - \frac{(1 - e^{-2\rho h})}{\rho Ch} \right] + \frac{(1-B)}{D} \left[(1 + e^{-2\rho h}) - \frac{(1 - e^{-2\rho h})}{\rho Dh} \right] \right\}. \quad (136)$$

The temperature rise to the center of the slab is

$$T(0) = \frac{A}{K\rho} \left\{ \left[\frac{B}{C} (1 - e^{-\rho h})^2 \right] + \left[\frac{(1-B)}{D} (1 - e^{-\rho h})^2 \right] \right\}. \quad (137)$$

Figure 7-133 shows the deposition profiles and temperature profiles for the case of ED-2 glass, for which $B = 0.7127$, $C = 0.1343 \times 10^{-20}$ cm², and $D = 1.1167 \times 10^{-20}$ cm². Figure 7-134 shows the

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Fig. 7-133. Deposition and temperature profiles for LD-2 glass.

Fig. 7-134. Relative surface stress per unit of deposited heat as a function of total ions/cm².

relative surface stress per unit of deposited heat as a function of the total ions/cm². Also shown are numeric results found by summing Green's functions instead of doing the integrals. Note that the surface stress drops fairly rapidly at first, but more slowly later; it takes large optical depths before the stress has dropped to 70% of the uniform-slab result. The problem, as we see from Fig. 7-133, is that there is appreciable deposition in the slab center even at high doping. Because diffusion from the slab's center is slow, the temperature remains high and, thus, so does the surface stress.

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Energy Applications

Introduction

M. J. Monsler

In searching for suitable solutions to future world energy problems, we anticipate that fusion will play a major role in providing a stable source of sustainable energy. Indeed, we can foresee the development of fusion power plants that operate without the air pollution and radioactive waste problems found in current coal-fired plants and fission reactors. Fusion power plants should also combine freedom from escalating fuel costs with low-risk utility investment and public safety. In view of these considerations, the potential future contribution of fusion energy to world economic and political security cannot be overemphasized. However, significant technological challenges confront us; accordingly, we must first demonstrate that fusion is both scientifically and economically feasible.

The scientific principles of inertially confined fusion have long been demonstrated, albeit in impractically large sizes, in the technically sophisticated and mature field of nuclear explosives. Still required is proof of the scientific feasibility of inertial fusion achieved with the use of lasers on a practical scale under laboratory conditions. Accordingly, the bulk of the Laser Fusion Program at LLNL is engaged in experimental and theoretical studies of fusion physics and laser driver development. To assess the eventual economic feasibility for energy production, we are conducting conceptual design studies of inertial fusion power plants for the production of electricity, fissile fuel, and synthetic chemical fuels. These studies have provided valuable guidance for planning the engineering development phase of fusion energy.

Previous studies have led to the definition of a very attractive reactor system known as HYLIFE (high-yield lithium-injected fusion energy). The characteristics of this design, one of a class of reactors referred to as liquid-metal-wall (LMW) reactors, have been reported in detail in previous annual reports.¹⁻⁴

The HYLIFE study demonstrated that a fusion power plant could be built with today's materials and reactor technology, provided that adequate targets and lasers had also been developed. Even so, two design improvements would still be needed: an increase in the allowed repetition rate (to ≥ 5 Hz) and a reduction of the estimated capital cost (to about that of a light-water reactor).

This year we studied new reactor concepts that address the shortcomings of HYLIFE and that also exploit some fresh ideas. In each case, our goal was to reduce the future cost of fusion electricity in one or both of two ways: (1) by reducing the capital costs of generating heat; or (2) by increasing the efficiency of conversion to electricity, either through direct conversion or by operating at a higher temperature. We developed three concepts, all more adventurous than HYLIFE in that they attempt to test and define the asymptotic potential of fusion power plants, rather than using currently available technology.

The first concept called Pulse*Star, is an LMW reactor that uses a lead-lithium eutectic as the tritium-breeding and heat-transfer medium. The reactor cavity is formed from a porous-walled spherical structure immersed in a pool of liquid metal. A double-walled heat exchanger between the lead-lithium and steam cycles eliminates the complexity and cost of an intermediate loop. Because we do not have the stored energy of a fission core, and because lead-lithium has sufficiently low flammability with air, water, and concrete, we can explore the possibility of drastically reducing or eliminating the conventional nuclear

containment building. Instead of a 2-m-thick concrete containment structure, we suggest a far less expensive strategy of multiple tritium-confinement barriers, along with tritium inventory minimization and segregation. Also, to reduce cost, we envision that a power plant would consist of up to four reactors driven by one time-shared laser and serviced by a common reactor services building. The sum of these innovations points toward the first fusion reactor that may be comparable in cost to a light-water reactor: an exciting prospect for fusion!

Sunburst, the second concept, differs from our other approaches in that we seek to maximize the kinetic energy delivered to a slug of liquid metal. Since the liquid metal is an electrical conductor, we can anticipate the efficient conversion of kinetic energy directly to electricity through magnetic induction, through magnetohydrodynamic generators, or both. A bottoming steam cycle can further increase efficiency. In exploring the wide range of novel physics and engineering constraints in this study, we hope to provide the physics basis for an entirely new kind of power plant that would operate at 60% efficiency.

The third concept, Cascade, was the result of our search for better ways to use a solid tritium-breeder material, while still preserving the benefits of the successful LMW designs. The initial concept appears exceptionally promising; all our goals for repetition rate, fire safety, wall protection, and efficiency seem attainable. The Cascade reactor concept features an approximately elliptical chamber rotating slowly on its horizontal long axis in a way that a 1-m bed of lithium oxide particles is held centrifugally against the wall. While this bed of loose particles cascades slowly toward its outlet at the largest radius, it simultaneously breeds tritium, absorbs the fusion energy, and protects the wall from both radiation and shock damage. In one option, a combined gas turbine and steam cycle would provide 50%-efficient electricity generation. Because of the superior fire safety and low tritium inventory of this design, it also may be operable without a conventional containment building, thus reducing capital cost.

In designing fusion reactors, many subsidiary studies and computational tools must be developed to support conceptual design studies. We also report on our efforts to develop better analytical and computational methods to calculate the induced radioactivity of complex fusion reactor designs constructed of diverse materials. This is an area in which we have had to lead the fusion community in order to quantify the unusually low activation of several of our fusion reactor designs.

Similarly, we have both developed our own expertise and funded outside contracts for behavior and compatibility studies of lithium and lithium-bearing compounds and alloys when used with steel structures. We therefore report on small experiments undertaken to understand the corrosion and embrittlement of liquid metals and their fluid-mechanical interactions with structures. We also present an assessment of electromagnetic pumps for pumping liquid lithium in fusion reactors.

Finally, returning to the bigger picture, we present the results of two studies performed to better understand the road to commercialization. Because we are currently uncertain as to whether fusion targets will be direct-driven or radiation-driven, we conducted a trade-off study of the reactor implications of the direct-illumination geometry. The most significant impact seems to be in the severe restrictions on cavity and heat-transfer system design that are due to the wide angular distribution of required beam directions. High $f/\text{No.}$ illumination from significant distances (~ 25 m) is required for all illumination schemes in a commercial reactor.

In a commercialization study, we contrast some options for engineering development of inertial fusion reactors and show a success-oriented plan for the development of fusion power. We show that a working demonstration reactor (300 MW_e) is possible by the year 2015 if we can begin the development of average-power drivers and reactor systems in the near future.

It is clear that, if we are ever to see the commercialization of fusion power, we must develop a very strong program to reduce the cost and complexity of eventual reactors. If not, fusion may suffer the same fate as solar-electric or fission breeders: demonstrably feasible, but rejected in the marketplace. We must direct our research toward innovative solutions for producing simple and cost-effective fusion power plants. If we can show with our reactor concepts that the physics and engineering constraints do allow an economically and environmentally competitive fusion plant, then public policy will demand that the technology be accelerated to hasten the era of commercial fusion energy.

New Inertial-Fusion Reactor Concepts

Pulse*Star

Pulse*Star: Ideas for a Commercial Inertial-Fusion Power Plant. If we were to suppose that inertial fusion eventually approaches its highest feasible performance potential, could a power plant then be designed that the utilities would buy and that people would accept? We believe that, for fusion to be accepted, technical feasibility alone is insufficient; a fusion power plant must also be safer and cheaper than any other alternative. With these ideals as guides, we set the following goals for our pure fusion power plant, Pulse*Star. This power plant must (1) have the same electricity cost as an advanced fission reactor (probably fueled by a fusion-fission hybrid), (2) be inherently safe for public health, and (3) be a sound utility investment.

Fulfilling these goals required a new approach. First, we began with a top down definition of goals for required cost, safety, and performance. Second, we searched for a reactor concept that was extremely simple in form and function, since complexity invariably leads to cost. Finally, we did not accept any features of fission reactor design or regulation without question or re-examination. The Pulse*Star concept described here is our first expression of these goals and approaches.

We first consider the safety aspect of the Pulse*Star reactor. We define a reactor as inherently safe when it is impossible for a catastrophic internal accident to occur that could either release significant quantities of harmful nuclides to the environment or destroy the utility's investment. To satisfy this goal, we cannot have any significant internal energy sources other than the normal fusion energy release, the normal thermal content of the reactor vessel, and the reactor-vessel coolants. We eliminate liquid lithium because, in the event of an accident, it can release large amounts of energy if it combines chemically with air, water, or concrete. Two attractive alternatives to liquid lithium can breed adequate tritium: a solid, such as lithium oxide; and the liquid lead-lithium eutectic Pb₈₃Li₁₇. These compounds also have the low vapor pressure required for beam propagation of either lasers or heavy ions.

For Pulse*Star, we chose to use lead-lithium in a liquid-metal-wall reactor design.³ Lead-lithium is very effective at protecting the steel structure from radiation damage; it has an exceptionally low tritium inventory, and it also has an exceptionally high energy multiplication. Because lead-lithium is simultaneously the energy-absorption medium, the tritium-breeding medium, and the heat-transfer medium, we hope to achieve the simplicity that will lead to low cost. Finally, by using a thick layer of lead-lithium to protect the steel vacuum vessel, we can expect that the major structural components of the reactor will not have to be replaced during the lifetime of the power plant.

New Inertial-Fusion Reactor Concepts

On the negative side, lead-lithium is 20 times heavier than lithium, thus greatly aggravating plumbing and piping design. In addition, since tritium has such a low solubility in lead-lithium, we also will have to take special measures to control tritium permeation through piping and component walls.

The Pulse*Star reactor pool is shown in Fig. 8-1. The lead-lithium pool, at 350°C, contains the reactor cavity, the liquid-metal pumps, and the steam generators. The cavity is formed by a hemispherical 4-m-radius first wall, open at the bottom. This first wall is simply a dense screen, shaped to form a canopy, made of a refractory metal such as a vanadium alloy. The screen canopy is surrounded by a 5.5-m-radius perforated-steel reactor vessel. In the 1.5-m-thick blanket zone between the first-wall screen and the perforated outer vessel, a major fraction of the neutron energy is absorbed in an extremely dense spray of lead-lithium. The outer perforated vessel acts as the spray manifold and is fed by inflow from the pool. The two-phase fluid in the

blanket zone both provides effective shock isolation from the very rapid neutron energy deposition and rewets and cools the inner screen canopy. The screen canopy is supported inside the reactor vessel by six vertical ribs (not shown). This entire inner structure will be replaced every five years when the reactor is down for normal turbine maintenance.

On each shot, the x rays and debris from the target are absorbed in the first few micrometers of lead-lithium film on the inside of the screen canopy. The extremely hot vapor ablates, radiates much energy back to the wetted screen, and flows to the lower blanket region, where it is very rapidly cooled and condensed by the spray. We have determined condensation rates by scaling detailed computer calculations for the University of Wisconsin HIBALL⁶ reactor. These calculations show that the lead-lithium plasma should condense to 10^{-2} Torr in less than the 200 ms required by the 5-Hz repetition rate. This rapid condensation occurs even if the gas can only cool and condense on the wetted surface area of

Fig. 8-1. Elevation drawing of the Pulse*Star pool-type reactor. The reactor vessel, four pumps, and eight steam generators are all located in the 20-m-diam pool of lead-lithium eutectic.

the screen; the lower spray zone should ensure an even faster cooling and condensation time.

Heat transfer to the screen is not a serious concern. Even if the entire 200 MJ of x-ray and debris energy was transferred from the hot plasma to the vanadium-alloy wire mesh by radiation and convection, the increase in temperature would be less than 600°C, comfortably less than the melting temperature of 1900°C.

The fast ablation resulting from x-ray and debris energy deposition causes a reaction force on the fluid filling the screen and, thus, also on the metal mesh itself. For a 600-MJ yield at 4-m radius, the momentum per unit area is quite moderate—equivalent to about 10 cm of water dropping a distance of 1 m. We believe that the first-wall screen can support an impulse of this magnitude.

The blanket spray zone is filled with lead-lithium droplets at approximately one-half the density of the liquid. Every 0.2 s, newly injected streams shatter into more droplets to relieve the pressure created by the neutron energy addition, and an equal volume of droplets falls to the hot pool between shots. Thus, the zone is kept filled with a dense distribution of spray. The momentum created by neutron energy addition is delivered to the wall over very long times (about 300 ms) because the interaction time is on the order of the layer thickness (1 m) divided by the induced fluid velocity (~ 3 m/s). Because this time is over 10 times greater than the wall response time, the structure is well protected against transient impact stresses from the liquid set into motion by the neutron energy addition. Furthermore, the induced droplet velocity is sufficiently low that droplets will not be propelled upward into the beam path.

The inner screened canopy is essentially unprotected from radiation damage and will have the shortest projected life of the cavity components. Fortunately, it does not have to meet the standards of a pressure vessel code, and there is no swelling constraint. We are now trying to calculate the lifetime of this first wall in a consistent manner.

Pulse*Star has some features that are similar to two other ICF reactor concepts. In the original "wetted wall" concept of Los

Alamos,⁷ the first wall and blanket functions were accomplished by a dense, porous wall backed by a plenum of lithium at full density. A thin film of lithium was reestablished by liquid oozing from the plenum through the porous wall. Technical concerns in the wetted wall design included the condensation rate to the minimal wetted surface area, the uncertainty of keeping the porous wall wetted, and the shock magnitude resulting from isochoric heating of the lithium in the plenum. In Pulse*Star, neither cooling and condensation of the hot vapor nor the neutron-energy-induced shock is a problem because of the effectiveness of the extensive spray zone. However, assurance of effective rewetting remains an issue in any porous-wall reactor.

In the Eagle concept of Bechtel,⁸ a thin region of lithium spray cools the cavity gas used for ion-beam propagation and x-ray absorption. In Pulse*Star, there is no need to cool large volumes of noncondensable gas. The Eagle concept neither employs a screen to prevent droplets from reaching the center of the chamber nor uses a spray zone sufficiently thick to affect neutrons. Unlike Eagle, the spray region in Pulse*Star is large enough and dense enough to absorb the neutron energy and protect the pressure vessel.

In evaluating costs for Pulse*Star, our goal is to design a fusion reactor where costs (exclusive of driver) are comparable to the capital cost of a pressurized water reactor (PWR), even though our design has more in common with a pool-type liquid-metal fast-breeder reactor (LMFBR). The portion of the electricity cost attributable to the fission fuel cycle would, in inertial fusion, be approximately balanced by the laser system.

In all of these studies, we assume that a fully established fusion industry exists and that licensing procedures are favorable for the innovations that we suggest here. We also borrow ideas for relevant cost-cutting techniques from two recent major efforts to reduce LMFBR costs: the Westinghouse/Burns & Roe effort,⁹ and the French "Superphenix II" studies.¹⁰

In comparing Pulse*Star to an LMFBR, we can first eliminate all the fission-related components and systems—the core, the

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rotating-plug refueling machine, redundant control-rod drive assemblies, sodium-cooled spent-fuel assemblies, emergency core-cooling systems, and decay heat-removal systems. In addition, we can eliminate the intermediate loops, along with the substantial containment volume required to house them, and substitute duplex-tube steam generators. Piping complexity is minimized by the use of pipe bellows and siphon breakers and by the pool design.

The need for expensive steel-lined concrete compartments with inert gas atmospheres is essentially eliminated by the use of lead-lithium in a pool geometry; there are no liquid metals outside the pool. Recovery of tritium permeating through piping structures is accomplished by having the entire pool surrounded by an outer liner, with a tritium-recovery gas flow between the pool structure and liner. Because tritium confinement is accomplished at the pool, and because there is no possibility of a core melt or other large-scale energy release, the entire building is technically not a containment building. It is strong enough for seismic and tornado forces and tight enough for confining most of the tritium in the event of an internal release during maintenance.

The confinement building volume and cost can also be dramatically reduced by

eliminating the internal crane loft; we are investigating rooftop access for repair or replacement of all major components. Square buildings are used to reduce cost. The sum of these features leads to a reactor building that is $34 \times 48 \times 30$ m, substantially smaller than either of the low cost LMFBRs being designed (SuperPhenix II is $72 \times 31 \times 64$ m, not counting separate steam generator buildings, and the Westinghouse/Burns & Roe LMFBR is $51 \times 51 \times 75$ m, including the steam generators).

The layout of the Pulse*Star confinement building is shown in Fig. 8-2. The reactor cavity, the four liquid-metal pumps, and the eight duplex-tube steam generators are all in the central pool. Vacuum pumps flank the laser beam ports. The vacuum system consists of Roots blower-type pumps and a cryogenic collection system for the exhaust gases, primarily helium, deuterium, and tritium. Tritium is recovered both from these exhaust gases and from the purge-gas flows (He plus a trace of O_2), which prevent tritium permeation to the outside environment through the steam generators and the pool surface area.

The final optics for directing the laser beam to the target are located 30 m from the center of the fusion chamber. A low-pressure gas flow is directed from the mirrors down the beam corridors to the

Fig. 8-2. Plan drawing of the Pulse*Star reactor pool and confinement building. The reactor first wall, pumps, steam generators, and final optics can all be replaced by access through the roof.

vacuum pumps to protect the mirrors from the damaging x rays and debris from the microexplosion. Beam tubes penetrate the pool and the screen first-wall structure.

The power plant consists of two or more 1290 MW_e reactors, serviced by a common laser driver located in a separate building of ordinary industrial construction. Underground concrete tunnels carry the laser beams to the reactors. The laser driver, reactor services building, and pellet factory are shared among reactors; costs for these facilities are shared as well. The power plant layout is illustrated schematically in Fig. 8-3.

When the porous wall of the reactor chamber must be replaced (every five years), an outdoor crane transfers components between the reactor and the reactor services building. The reactor vacuum vessel, the lead-lithium pumps, and the steam generators are designed to last the lifetime of the plant; should they need to be replaced, they too can be removed through the roof of the reactor building with minimum downtime. The turbine generators and their buildings are of standard design and construction; we expect that 1500 MW_e turbines will be available after about the year 2020.¹⁰

The power plant performance is summarized in Table 8-1. Essentially, we have a twin 1300-MW_e unit plant driven by a single shared short-wavelength laser system. Although driving up to four reactors with one laser would be more cost effective, there are diminishing returns after two units. Our cost-control strategy is evident: to operate at high gain, at a large unit size and at a high operating temperature and to operate with minimum containment structure and with multiple reactor units driven by a single laser. The absolute costs have not yet been determined because sufficiently detailed engineering designs do not yet exist.

Substantial technical issues arise with the Pulse*Star conceptual reactor. The most severe issues are associated with the use of lead-lithium as the coolant and with the control of tritium. Lead-lithium is far more corrosive than either lithium or sodium, and it will become more radioactive than lithium, thus (in conjunction with dissolved corrosion products) becoming a source of radioactivity and afterheat. The extremely low tritium solubility in lead-lithium will drive tritium through pipe and component walls, increasing tritium confinement and recovery costs when compared with lithium.

Fig. 8-3. Layout of the Pulse*Star power plant, showing two 1300-MW_e reactors driven by a common laser.

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Table 8-1. Performance parameters for the Pulse*Star twin-unit inertial-fusion power plant.

<u>Each reactor of a two-reactor system</u>	
Target gain	= 530
Yield	= 600 MJ
Pulse repetition frequency	= 5 Hz
Fusion power	= 3000 MW _{th}
Blanket energy multiplication	= 1.26
Thermal power	= 3800 MW _{th}
Thermal to electric efficiency	= 39.5%
Gross electric power	= 1500 MW _e
Required auxiliary power	= 67 MW _e
Net electric power	= 1290 MW _e
<u>Short-wavelength laser driver</u>	
Pulse energy	= 2 MJ
Pulse repetition frequency	= 10 Hz
Driver efficiency	= 7%
Required electrical power	= 286 MW _e
<u>Power plant (one laser plus two reactors)</u>	
Net electric output	= 2580 MW _e
Fusion energy gain	= 21
Recirculated power fraction	= 14%
Net electrical efficiency	= 34%
Tritium breeding ratio	= 1.4

The improvements in fire safety and the gain in energy multiplication may not be adequate benefits; we may need to return to an intermediate loop (either water or sodium) to ensure safety. These safety and cost evaluations will be part of our ongoing studies. In the coming year, we will design a lithium-cooled version of Pulse*Star, perform some accident-risk analyses, and estimate capital costs. In doing so, we can assess more accurately whether our ambitious cost and safety goals for Pulse*Star are attainable.

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Tritium: Flow and Heat-Transfer Considerations for Pulse*Star. The Pulse*Star conceptual reactor will reduce plant cost by transferring heat from the working fluid to a steam loop within a pool surrounding the reaction chamber. Safety considerations encourage the use of nonreactive Li₁₇Pb₈₃, rather than pure lithium, as the working fluid. This article examines the pumping-power, heat-transfer, and tritium-recovery aspects of the Pulse*Star concept.^{11,12}

To determine the minimum number of liquid-metal pumps required, we first calculate the minimum coolant flow rate, which is set by the temperature rise, the thermal power, and the specific heat. The maximum temperature of 500°C is set by the steel strength, and an absolute lower bound is the coolant melting point of 235°C.

Radiation-induced embrittlement of the structural steel sets a more restrictive lower temperature limit of ~400°C; therefore, the system ΔT is ~100°C. For 1000 MW_e-net, about 3300 MW_{th} must be added to the Li₁₇Pb₈₃, and the specific heat is 0.16 kJ/kg·K. Thus, the flow rate is ~2.2 m³/s. Since the chamber diameter is ~11 m, a gravity head of 16 m is probably a minimum, and the pumping power is ~33 MW(mech).

The hydraulic design of a liquid-metal pump is influenced by the volumetric flow rate and the head; it is independent of the fluid density. If four pumps are used in Pulse*Star, the flow rate per pump of 5.5 m³/s is 70% of the HYLIFE pump flow rate and is ~3 times the Clinch River breeder reactor pump flow. A mixed flow (between axial and radial type) impeller pump with two impeller stages of 1.2-m diam can produce the desired head and flow at a suction head (NPSH) of 3.4 m.

The mechanical design of liquid-metal pumps depends on fluid density, and Pulse*Star pump components must be considerably stronger (and more massive) than their counterparts in Li or Na pumps of the same design. Also, the drive motor for the Li₁₇Pb₈₃ pump must be 19 times more powerful than the motor for a Li pump of the same design. Motors of this size are commercially available.

Two geometries have been considered for steam generators within the liquid-metal pool: horizontal cross-flow tubes arranged above the reaction vessel in four steam generators, and vertical counter-current tubes in eight steam generators between the pumps.

In the horizontal cross-flow geometry, the pumps will deposit the liquid metal on a free surface above a wide array of heat-exchanger pipes covering a significant portion of the pool area. The pipes will be spaced far enough apart to reduce the pressure drop while maintaining a reasonable heat-transfer coefficient. The optimum design point has 18 000 m² of total exchanger area arrayed in a 4.9-m-deep layer with a pitch-to-diameter ratio of 2.0. The pool radius is 10 m.

In the vertical counter-current steam generator option, the pumps inject the hot liquid metal into the top of the steam generators. The down-flowing liquid metal exits the steam generator into the pool bottom. The pump head required is the sum of the

pool depth and the friction loss in the steam generator tubes. The total heat exchanger area is 8500 m², the steam generator diameter is 1.7 m, and the pool radius is 9 m.

Heat loss from the hot-leg piping to the cold pool is effectively eliminated by insulating the piping runs with 1 cm of calcium silicate insulation sandwiched between two concentric pipe walls. The counter-current design was selected over the cross-flow design because it reduced the steam generator size, the pool radius, the liquid metal volume, and the pumping power.

We wished to determine the best time to recover bred and unburned tritium: from the liquid Li₁₇Pb₈₃ from the chamber vapor, or through the heat exchanger tubes. Before considering these options, we first made some calculations of breeding rates and required recovery rates.

The recovery scheme will be economically successful if >95% of the tritium is recovered by the desired system. However, the tritium fraction that is not recovered must be prevented from reaching the general environment because the acceptable release rate of ~3 Ci/day represents only 3.5 ng/s, or 0.00002% of the production rate.

With 30% burnup and a 1.1 tritium-breeding ratio, a 1-GW_e Pulse*Star produces ~16.4 mg/s of bred and unburned tritium. At a chamber pulse-rate of 5 Hz, 3.3 mg of tritium, 2.6 mg of helium, and 1.5 mg of deuterium are produced per pulse. In the 520 m³ Pulse*Star sphere, these incremental masses represent partial-pressure increments per pulse of 51 μTorr tritium, 60 μTorr helium, and 35 μTorr deuterium, assuming that all of the vapor is added to the 500°C chamber. On a steady-state basis, this is reasonable, since the tritium and helium bred in the liquid metal will have an opportunity to join the chamber vapor when the liquid metal enters the chamber void.

In steady state, the removal rates of tritium, helium, and deuterium must be exactly the same as their production rates. However, if their concentrations have been allowed to build up during a transient period, the steady-state removal fraction of the total inventories can be very small. The upper limit on the inventory is set by either safety or beam propagation constraints.

The concentration of tritium in the 500°C Li₁₇Pb₈₃ liquid metal determines the tritium

partial pressure in the vapor, $P(\text{Torr}) \sim 2.2 (\text{wppm})^2$; and the concentration is found from the solubility ($C = \sqrt{P/K_s}$). We need a solubility [$K_s = 1800 \text{ atm}^{1/2}$ (mole tritium/mole alloy)] that is very low.

The tritium added to the liquid metal during its 33-s exposure to the neutron flux is

$$\frac{(16.4 \text{ mg tritium/s})}{(10.93 \text{ m}^3 \text{ liq/s})(9.35 \text{ Mg/m}^3 \text{ liq})} = 0.0002 \text{ wppm} \quad (1)$$

Using the solubility data, this corresponds to only 6×10^{-8} Torr tritium-pressure rise. The concentration rise is further reduced by dilution with the remainder of the pool. Thus, many passages are required before the steady-state inventory of tritium in the liquid is reached and before tritium recovery by any method is practical.

Three tritium recovery options¹¹ have been considered: from the liquid, from the vacuum system, and through the heat exchanger.

First, recovery from the liquid using a molten salt is impractical because leakage through the vacuum system requires processing the liquid metal to an uneconomical 0.02 wppm. Alternatively, if 1 wppm processing is assumed, we must be able to propagate the laser through 260 Torr helium plus 2 Torr tritium.

Second, recovery through the vacuum system requires low leakage through the heat exchanger, and this leakage is a function of the concentration of tritium in the liquid. If the tritium concentration in the liquid is kept low by heroic vacuum pumping down to twice the liquid-metal vapor pressure, about 3% of the tritium leaks through the heat exchanger. This leakage must be reduced by a factor of 10^5 by using double-wall tubes or diffusion barriers.

Third, recovery through the heat exchanger requires low leakage through the vacuum system and the pool wall. The leakage through a 10-cm-thick steel pool wall is ~0.5% of the diffusion into a 25 000 m² heat exchanger. If ~90% of the tritium is forced to diffuse through the heat exchanger, the required tritium pressure is 0.05 Torr (0.17 wppm). For a chamber pressure that is the maximum possible for laser propagation (say, ~1 Torr, most of which is helium), we would pump 6×10^{-5} of the

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vapor ($60 \mu\text{Torr}/1 \text{ Torr}$) per pump cycle, and 6% of the tritium would escape through the vacuum system.

The tritium can then be recovered from either the vacuum system or the heat exchanger by setting the background chamber pressure; i.e., by design of the vacuum system. If the heat-exchanger recovery option is chosen, diffusion barriers (10^5 reduction of leak rate) on the pool wall will be required, as well as some recovery from the vacuum system; however, these requirements can be met. In addition, the vacuum-system design is very simple and economical. If vacuum-system recovery is chosen, we would use diffusion barriers on the pool wall and heat exchangers (10^5 and 10^7 reductions, respectively) and would operate the vacuum system at $\sim 10^{-2}$ Torr.

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Sunburst—A Direct-Energy-Conversion ICF Reactor

Introduction. In inertial-confinement fusion (ICF), thermonuclear burn occurs in very short pulses and in a very small volume. We have designed a new reactor concept that exploits these two characteristics of ICF. This conceptual reactor, Sunburst, features

- High net electric efficiency—electricity generated both by direct conversion of kinetic energy and by thermal conversion of heat energy, yielding a net plant efficiency of $\sim 50\%$.
- High power density—high yields contained in chambers less than 5 m in radius.
- Low neutron-induced radioactivity—the blanket geometry makes possible a neutron induced radioactivity several orders of magnitude less than other DT fusion reactor concepts.

Most other fusion reactor concepts convert fusion energy to electricity primarily by a thermal cycle, have relatively large blankets, and, except for liquid-metal-wall ICF

reactors (such as HYLIFE¹³), have highly radioactive first walls that must be replaced every few years because of neutron-induced damage. In addition, those designs usually require neutron multipliers (such as lead or beryllium) or isotopic enrichment of ⁶Li to achieve acceptable tritium-breeding ratios.

The Sunburst concept, first described in Ref. 14, most closely resembles the proposal of Burke and Cutting,¹⁵ who also proposed to absorb most of the fusion neutron energy in a lithium sphere contiguously surrounding the ICF plasma and to convert most of the energy to electricity by work done on magnetic fields. In their scheme, however, the lithium is initially solid, is entirely *brought to the plasma state with an extremely large fusion yield*, and is contained by a correspondingly large containment vessel.

Reactor Operation. In the Sunburst concept, we maximize the shock-induced kinetic energy in a liquid-lithium blanket. The kinetic energy does PdV work against an applied magnetic field to directly generate electric energy with high efficiency by induction. The Sunburst concept invites a wide range of imaginative configurations. Although many engineering problems are associated with each approach, we will examine one of them—the rotating conducting chamber—to see how a reactor might actually be constructed.

A liquid-lithium sphere, which we will call the “compact blanket,” is dropped into the reaction chamber while carrying a solid lithium tube that contains a fusion target. This is illustrated in Fig. 8-4(a). The tube provides access to the fusion target from two sides for laser-beam energy. (A removable bullet-shaped cap prevents liquid lithium from entering the tube as it is inserted into the liquid lithium.)

The compact blanket is formed by a mechanical “ice cream scoop” apparatus. With the apparatus itself in free fall for the short time required to form the sphere, deformities can be avoided. Similarly, because the compact blanket is in free fall in vacuum, no forces deform its spherical shape as it falls.

The spherical chamber, Fig. 8-4(b), rotates at constant angular velocity about the axis of the two cylindrical armatures to capture a new compact blanket every half rotation.

The chamber acts as two half-turn electro-magnets each separated by a continuous insulation break and driven by current supplied to the armatures. Helium gas is circulated inside the chamber wall for cooling.

When the compact blanket reaches the center of the chamber, the magnetic field is pulsed on in a time scale for which the sphere is effectively motionless. Lasers then irradiate the target and the fusion neutron yield is absorbed in the lithium in a time that is short compared to a hydrodynamic response time, raising central pressures to over 1×10^6 atm and launching a strong shock wave.

The deposited neutron energy brings the lithium that is within a few centimeters of the target to extra cohesive internal energies—it is a gas at liquid density. This gas fills the access tube so quickly that energy loss occurring from blackbody radiation and free-streaming ions through the access tube is negligible.

The energy behind the shock wave is in the form of both kinetic energy with outward-directed velocity and internal en-

ergy or pressure. This energy reaches the surface in less than 1 ms, and a pressure relief wave moves back to the center, nearly doubling the outward-directed fluid velocity to more than 160 m/s.

The kinetic energy in the liquid lithium originates from high internal pressures caused by the deposition of neutron and x-ray energy in the lithium, as well as from the work done by the expansion of both the postfusion plasma and the central portion of lithium, which has been heated to the gaseous state. Therefore, the neutron and x-ray energy, in addition to the ion energy, take part in the direct generation of electricity. This is in contrast to electrostatic direct-conversion schemes in which only the charged-particle energy may be directly converted to electricity.

After the liquid lithium is brought to rest by the magnetic field pressure, it drains to the chamber floor, Fig. 8-4(c), and is pumped through heat exchangers, where residual heat is converted to electricity in an ordinary thermal cycle. A fraction of the lithium is then diverted for impurity removal, tritium

Fig. 8-4. Conceptual drawing of the Sunburst reactor.

Not shown:

- Inert-gas cooling veins in armature and magnet
- Glass fiber and epoxy or SiC-fiber binding to contain magnet expansive forces

harvesting, and casting for new laser-beam access tubes; the remainder is recycled for subsequent compact blankets.

Liquid lithium—rather than solid lithium—must be used because efficient thermal-to-electric conversion requires heat-exchanger outlet temperatures to be well above the melting point of lithium. In addition, liquid lithium, a very efficient heat-transfer agent, is its own coolant and can be pumped electromagnetically without wear or corrosion to seals or impellers. Moreover, it need not be pumped across strong magnetic fields, as is necessary in magnetic confinement schemes.

Magnet Configuration. The magnets may be configured in several ways. Even though a stationary reactor chamber with fixed holes at top and bottom and with exterior magnetic-field-shaping coils would be simpler, the rotating conducting chamber has the advantage of reducing the total magnetic field energy required to restrain the liquid lithium by confining the magnetic flux lines within the chamber. Alternatively, moving flux shields in electrical contact with a fixed chamber could be used to cover entrance and exit ports in a variety of configurations.

With proper chamber shape design, a homogeneous magnetic field strength between the conducting chamber and the lithium sphere can be approximated, an important point for minimizing fluid instability growth times. On the other hand, it may be desirable to shape the magnetic field so that the field strength is either greater or weaker at the throats of the chamber—like a magnetic mirror—to minimize or maximize the portion of energy converted by MHD generation.

Because the pulsed-field time is short compared with the time between pulses, the averaged ohmic power losses are a small fraction of the total plant electric power. For this reason, the magnets need not be superconducting, greatly simplifying the design and reducing the cost.

The magnetic field serves the additional purpose of distributing the total impulse imparted to the reaction chamber vessel over time, acting as a cushion. If the impulse time is sufficiently long, the chamber wall will not be stressed beyond acceptable fatigue life limits.

In addition, during the deceleration of the lithium, the applied magnetic field can be maintained at a constant level by decreasing the electric potential of the external energy-storing load, so that the lithium suffers a constant deceleration, comes smoothly to rest, and the Lorentz stress on the chamber walls is at a constant, acceptable value.

Operating Design Space. The Sunburst reactor design is limited by the following general constraints

- A minimum lithium sphere radius is required to achieve a tritium-breeding ratio greater than unity.
- A maximum lithium sphere radius is limited by the pumping power required to recirculate the liquid lithium.
- A minimum fusion yield is required to displace the magnetic field faster than the field can diffuse into the lithium.
- A maximum fusion yield is limited by the strength of materials.

We will now discuss some of the major considerations for each of these parameters. Concurrently, we will use computer calculations from the LASNEX, ORLIB, TART, and ACTL codes to see how one point in this design space—3-m radius, 1200-MJ yield—determines other design parameters. Finally, we will apply these parameters to estimate a net electric efficiency for that operating point in this design space.

Tritium Breeding and Energy Multiplication. With natural lithium, the tritium-breeding ratio, in compact blankets of radii greater than 1 m, is well above unity, as shown in Fig. 8-5. No neutron multiplier or lithium isotopic tailoring is needed, and no external tritium breeding blanket is required. With the tritium content of the liquid lithium kept to 1 wppm by molten-salt extraction techniques, the reactor tritium

Fig. 8-5. Tritium-breeding ratios and blanket gain to total fusion energy using natural lithium.

inventory and the tritium doubling time will also be low, an advantage over solid-breeder reactors for both environmental safety issues and start-up inventory expense. The gain of fusion energy from nuclear reactions in the blanket is also shown in Fig. 8-5. Using spherical reactor models with radii of from 1 to 3 m in the TART neutron transport code, reaction-probability-rate constants were obtained with ORLIB, which averages the ACTL library cross sections with the flux provided by TART, to give tritium breeding ratios. The gains, taken directly from TART outputs, are the ratio of net energy per fusion event to 17.6 MeV.

The reactors were modeled with a typical ICF compressed-fuel target at the center of liquid-lithium spheres of varying radii, surrounded by a vacuum of one-half the lithium radius, followed by a structural shell composed of 8 cm of steel. A typical MFE fusion reactor was also modeled, having a homogeneously distributed 14.08-MeV neutron source to a radius of 2 m in a DT plasma at 0.01 mbar, followed by vacuum to 3 m, a 2-cm first wall, a 70-cm lithium blanket, and a 2-cm outer structure. Natural lithium and a low-ferritic steel (2-1/4 Cr-1 Mo) were used in all cases.

Neutronics. With lithium sphere radii on the order of 3 m, the fusion neutrons are required to traverse many more mean free paths of lithium than they do in any other fusion reactor designs. As a result, a Sunburst reactor can have many more exponential attenuations in both population and energy content of the escaping neutron flux. For this reason, the neutron-induced

radioactivity in the reaction chamber walls will be several orders of magnitude less than in typical fusion reactors.

As with other liquid-metal-wall reactor designs, there is no structural material interposed between the fusion neutron source and the lithium absorbing medium. In targets designed to minimize induced radioactivity, and with ongoing impurity removal from the circulating lithium by eutectic condensation, the induced radioactivity of extracted target debris can be undetectable after a few weeks' holding time.¹⁶

The normalized neutron flux spectra in the reactor chamber wall for the reactors modeled with TART are shown in Fig. 8-6(a). The neutron-path-length distribution shown there may be converted to a flux distribution by multiplying by the number of source neutrons per second and dividing by the wall area.

Figure 8-6(b) shows the neutron energy deposition profile calculated with TART using a 3-m sphere. A two-dimensional model was used to account for collisionally enhanced neutron leakage through a two-beam laser access tube, each side of which described a cone of 0.0153 sr. (This is the same as currently assumed for the HYLIFE reactor design.) The data points are fit closely by $\exp(-\alpha r^2)/r^2$, where $\alpha = 1.2/\text{m}^2$, and this curve is plotted for fusion yields of 300, 600, 1200, and 1900 MJ. The cohesive energy (above which liquid-density lithium behaves as a gas) is shown as the horizontal line at $\approx 24 \text{ MJ/kg}$.

Hydrodynamics. After the neutron energy is deposited in the lithium, a strong shock wave propagates outward to relieve

Fig. 8-6. (a) Neutron flux spectrum in the reactor chamber wall. (b) Neutron energy deposition profile.

the pressure. The shock response of a 3-m-radius lithium sphere to a 1200-MJ fusion yield was modeled with LASNEX and is shown in Fig. 8-7. Because of the low temperatures, a two-phase equation of state was used. Energy loss from escaping ions and gases through the access tubes was modeled with a homogeneous energy loss rate in the extracohesive region equal to the product of the local thermal velocity, the density, the specific energy, and the area of the beam cones at the extracohesive radius.

Despite the velocity difference during the shock propagation, the outer surface of lithium moves less than a centimeter by the time that shock-wave relief takes place. Since the velocity of the outer surface after that time is essentially the same as that of the lithium behind it, the lithium sphere does not spall.

Fluid Instability. The shock wave transports energy outward into the lithium before significant fluid motion of the lithium can occur. Thus, Rayleigh-Taylor instability on the surface between the inner gas and the liquid lithium does not inhibit the isotropic coupling of most of the central internal fusion energy to the rest of the lithium. Furthermore, the instability growth rate initially is zero, since the gas and the liquid have equal densities, and the pressure gradient provides stability until the time that the inward-moving relief wave arrives at the gas-liquid surface. After that time, the acceleration of the fluid is in a direction that will stabilize this surface.

The outer surface of the liquid metal will, however, develop flute instabilities. This surface is decelerated by interaction with the magnetic field, which can be thought of as the light fluid in a Rayleigh-Taylor unstable interface. For small amplitudes, the

modes grow exponentially in time with the square root of the product of the wave number k and the acceleration a ; they saturate, growing quadratically in time, after they have grown so large that they decouple hydrodynamically from their surface of origin, approximating free fall. Saturation begins when the instability departs from the sinusoidal—when the amplitude exceeds the inverse of the wave number.

Although such instabilities would be catastrophic in a steady-state system, in a pulsed system they can be dealt with by choosing a design such that the critical instability growth times are long compared to the time required to stop the motion of the liquid lithium. The wave number of the critical instability mode is about equal to the inverse of the final average lithium shell thickness. For a compact blanket with an initial radius of 3 m, the critical instability has a wavelength of 4.9 m. For a constant magnetic field energy equal to the initial kinetic energy density, instability "punch-through" is avoided if initial perturbations are less than 19 cm in amplitude. For higher magnetic fields, greater perturbations can be tolerated.

Energetics, Kinematics, and Continuity Considerations. If a magnetic field of energy density equal to the initial kinetic energy density of the lithium is chosen, and if the field is constant and homogeneous, a volume of magnetic field equal to that of the lithium must be displaced to bring the lithium to rest. By the time the lithium has displaced a volume equal to the original volume of the sphere, it will have formed a shell whose average thickness is $(2^{1/3} - 1)$, or about one-fourth of the initial radius, and whose average radius is about 1.25 times that of the initial radius.

Fig. 8-7. Time-resolved velocity, pressure, and density distribution in a 3-m-radius lithium sphere.

With a constant applied magnetic field, the deceleration is approximately constant. The time required to stop the lithium motion is, therefore, one-fourth the initial radius divided by one-half the initial velocity.

Because the lithium is approximately an incompressible fluid, the volume of undisplaced magnetic field contained between instability peaks after the lithium is arrested is equal to the volume displaced. If the critical instability grows to an amplitude of $1/k$, the reaction chamber wall must be located at a radius of about 1.5 times the initial lithium sphere radius to avoid contact with critical instability peaks.

With a homogeneous and constant applied magnetic field of 3.3 T, a lithium surface, originally at 3 m and with an initial velocity of 135 m/s, moves to an average radius of 3.8 m, and the stopping time is 12 ms. The chamber wall radius necessary to avoid contact with instability peaks originating from 19-cm perturbations is 4.6 m.

Electrical Coupling. With a conductivity of 2.5 Mmho/m (MS/m) for liquid lithium at 500°C, the applied magnetic field has a classical diffusion depth over the stopping time of 6.1 cm, less than 8% of the lithium thickness at the time it is arrested. (The square of the diffusion depth very nearly equals the diffusion time divided by π .)

Using a solenoidal magnetic field model (appropriate since it is also homogeneous), 3.3 T corresponds to a current of 21 MA (with the solenoid length chosen so that its surface equals that of the lithium sphere). By Ampere's law, a current of the same magnitude is induced in the lithium, which can be thought of as the secondary of a transformer. Using one-half the magnetic and diffusion depth as an average current conductance depth, the coupled joule heating energy loss in the liquid lithium over the expected lithium stopping time is 93 MJ. The joule heat loss in the magnet conductors is roughly this number reduced by the ratio of the conductivity of copper to that of lithium, or about 3 MJ.

The diffusion depth, the coupled energy loss, and (as we will see in a subsequent paragraph), the stress response of the chamber vessel all implicitly assume that the magnetic-field rise time is much shorter than the lithium stopping time. Using our solenoidal model from above, the induc-

tance between the magnet conductors and the lithium is $5\mu_0\pi R/8$, or, for a 3-m radius, 7.4 μH . The rise time is roughly given by the product of the current and the inductance, divided by the applied voltage. With an applied voltage of 500 kV, the rise time is 0.31 ms, which is negligible compared to the 12 ms stopping time.

Materials Considerations. Protected by 1 to 3 m of lithium, the reactor chamber wall experiences negligible damage from neutrons or x rays; in addition, the vacuum quality that is typical of magnetic fusion containment vessels is not required for laser transport. Therefore, we have more freedom in the choice of structural materials for our reactor. However, because of power consumption and Lorentz forces due to induced eddy currents that accompany rapidly changing magnetic fields, it will not be desirable to use electrically conducting materials except for the conducting elements of the magnet coil and the flux restrainers.

With these considerations in mind, the reactor chamber wall might be constructed of a layer of yttria (an electrically nonconducting, high-temperature material that is compatible with hot liquid lithium and is a good thermal insulator), followed by aluminum (for current conduction) followed by layers of silicon-carbide-fiber epoxy or glass-fiber epoxy (for tensile strength against expansive forces). Two other ceramic materials besides yttria that are nonconducting and compatible with hot liquid lithium are zirconia and thoria.¹⁷

Glass-fiber epoxies, with enhanced toughness and crack-growth inhibition, were used with the 2XII-B baseball magnet with tensile strengths of 46 700 psi, or 332 MPa, and a Young's modulus of 2.7×10^6 (Ref. 18). Using these values with the momentum and stopping time from the previous section, we have solved the equation of motion for spherical shells of 5-, 8-, and 11-cm thickness to determine the time-dependent maximum stress in the wall (see Fig. 8-8). We conclude that it is possible to contain a 1200-MJ yield with a fiberglass chamber about 8 cm thick. (Because epoxies rapidly lose strength when subjected to neutron flux, they are generally not considered in other reactor concepts.)

Thermal stress arises from differential expansion due to a thermal gradient. Since

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the inside temperatures are larger than those toward the outside of the shell, the yttria would be in compression and the fiberglass in tension—the strongest modes for the two respective materials.

Because of the low neutron flux in the reaction chamber wall, replacement of the wall due to helium embrittlement or lattice displacement damage is not necessary. The intervening lithium eliminates the problem of x-ray-induced wall erosion. Sunburst is probably the only conceptual reactor in which robotic maintenance might not be required—a possible great savings in operating and capital costs.

Energy Balance and Plant Efficiency.

The energy flow diagram of Fig. 8-9 is unusual because the blanket internal energy is divided into kinetic and thermal energy, each with different conversion efficiencies. Also, there are three cases in which the energy losses add heat to the liquid lithium, and that energy is partially recovered in the thermal cycle.

About 41.6% of the yield is converted to kinetic energy. This is the value inferred from Fig. 8-7 after shock and relief. In practice, a small amount (up to 21 MJ) of the yield is converted to kinetic energy by free

expansion of the internal gas through the access tubes and by PdV work done by the gas on the inner liquid surface. Inclusion of these effects should only increase the results for net electric efficiency, since MHD generators are typically more efficient than thermal conversion cycles.

In calculating the data shown in Fig. 8-9, we use the same laser beam energy, fusion yield, heat exchanger outlet temperature for liquid lithium, electromagnetic pump efficiency for liquid lithium, and thermal-to-electric conversion efficiency as those currently proposed for HYLIFE. A greater pump lift height (20 m) than that for HYLIFE was used, and a 10% loss due to randomized kinetic energy and an 8% loss due to pulsed-power conditioning were assumed. (At 1 Hz, and with lithium spheres with 3-m radii, the liquid-lithium pump rate is about the same as that of HYLIFE.) The overall plant efficiency is 49.9%. This is roughly 50% more than typical fission power plants.

Considerations Remaining. Liquid lithium raises serious fire safety concerns. However, if the levels of radioactivity pose a negligible environmental threat, then the fire hazard of this reactor concept requires

Fig. 8-8. Sunburst chambers of 8-cm-thick fiberglass can contain 1200-MJ yields.

only those precautions afforded other large fire hazards, such as municipal natural gas storage. If that were the case, then it may be possible to eliminate the very expensive containment structures typical of fission reactors, resulting in great capital savings.

Assuming that the compact blankets are dropped from rest into the reaction chamber, the repetition rates in each reaction chamber will be limited to less than about 1 Hz. Generally, low repetition rate is viewed as a disadvantage from the point of return on the capital investment in the laser. If this is an issue of economic importance, it could be overcome with the use of multiple reaction chambers served by the same laser. On the other hand, high repetition rates at the same plant may be unac-

ceptable if target gains are not as large as currently hoped for: to achieve economic gain-efficiency products, the yields may have to be so large that a plant of typical electric power cannot accommodate a high repetition rate.

Some other areas requiring investigation are electrical load matching and power conditioning, magnetic field coil design, thermal and fatigue wall stress analysis, the behavior of the outer portion of liquid lithium in which the magnetic field reaches large values, and a full two-dimensional MHD model of the compact blanket in response to the fusion yield.

Summary. Our study to date has revealed no major obstacle to an ICF reactor concept that maximizes the kinetic energy

Fig. 8-9 Sunburst reactors have very high net electric efficiency.

Gross electrical efficiency = 61.1%

Net electrical efficiency = 49.9%

Recirculating power fraction = 18.6%

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in a mass of liquid lithium with an ICF yield at its center and that converts that kinetic energy directly to electricity by induction or by MHD means. The neutron-induced radioactivities in such a concept can be exceptionally low, and the net electric efficiency can be unusually high. It is possible that these two major advantages will lead to an economic competitiveness available only to inertial-confinement fusion.

Author: G. P. Lasche

Cascade

Cascade: A Centrifugal-Action, Solid-Breeder Reaction Chamber. The Cascade concept¹⁹ represents a different approach for developing an ultrasafe, easily built reaction chamber in which the energy released by inertial-confinement fusion can be captured and used for producing commercial power. In the Cascade concept (see Fig. 8-10), a football-shaped reaction chamber rotates along its horizontal axis. Solid Li_2O pebbles are injected tangentially through each end of the chamber at a pe-

ripheral speed approximately equal to that of the adjacent pebbles inside the chamber. The pebbles are held against the wall of the chamber by centrifugal action and cascade axially from the smaller radii at the ends to the larger radius at the center, where the pebbles are then ejected from the chamber into a stationary catcher. Heat and tritium are extracted, and the pebbles are re-injected into the chamber.

As the pebbles cascade axially toward the center, they follow a helical pattern and experience an increasing centrifugal force that is proportional to the chamber radius. Because the chamber wall slope is greatest at each end, the velocity of pebbles along the chamber wall in the axial direction is greatest at each end, and a constant pebble thickness adjacent to the wall exists to satisfy continuity of pebble flow. This produces a constant-thickness blanket (~ 1 m) that offers nearly uniform protection of the wall from the high-energy neutrons generated in the fusion process.

A vacuum is maintained inside the chamber, and fusion-fuel pellets—each with a nominal yield of 600 MJ—are injected horizontally from one end at a rate of five pellets

Fig. 8-10. Cascade: a centrifugal-action, solid-breeder reaction chamber.

per second. Laser or ion beams, focused from each end on the center of the chamber and pulsed as each pellet reaches the focal point, are used to ignite the fuel pellets.

Other reactor concepts have used centrifugal blankets. Fraas²⁰ used swirling liquid lithium in the BLASCON design, whose central cavity was vortex-shaped. He proposed injecting inert gas into the lithium to reduce stresses in the chamber structural wall resulting from pressure-wave propagation. Baird and Anderson²¹ suggested a modification wherein both the liquid lithium and the chamber rotate together, forming a paraboloid-shaped cavity. Burke²² also suggested using centrifugal force to hold the blanket against the wall of the chamber. Concurrently with our study, Madaram et al.²³ proposed a horizontal rotating chamber employing liquid lithium and using wetted thermal shields for end closures. We wished to retain the concept of holding the blanket against the wall with centrifugal action, but looked for a means to make a more uniform blanket without top- or end-closure problems.

We also wished to eliminate the fire hazard of liquid lithium, so we incorporated the moving-solid-blanket concept documented previously by Meier and Maniscalco.²⁴ They suggested using falling Li_2O spheres, but their flow rate was large and, again, there were top-closure problems. Combining centrifugal action with a solid Li_2O pebble blanket capitalized on the best features of all the previously proposed chambers.

A key advantage of the Cascade concept is its inherent safety: it does not use liquid lithium. Fusion reactors using liquid lithium must be designed to mitigate potential fire hazards. The Cascade concept also offers the potential for improved power-plant efficiency because the high maximum Li_2O pebble-blanket temperature (1200 K) permits the use of a helium gas-turbine cycle. In addition, the inherent safety of Cascade may allow it to be located close to population centers, which would make more feasible the use of waste heat for low-grade industrial processes and space heating in a cogeneration mode.

The Cascade chamber may allow for some cost reductions in the balance-of-plant. First, we are investigating ways of

eliminating the intermediate heat-transfer loop required for tritium isolation in other fusion power plants: removing most of the unburned and bred tritium by the vacuum system, then removing the remaining tritium in fluidized-bed heat-exchangers using helium gas, which fluidizes the Li_2O pebbles. Second, we may be able to reduce the cost of the reactor containment building because lithium oxide does not burn when exposed to air, water, or concrete, and, therefore, cannot be a mechanism for releasing either tritium or radioisotopes in the activated structure. Hence, we may be able to construct the reaction chamber containment building with standard industrial requirements, rather than with the massive concrete structures commonly envisioned with nuclear power plants.

Shock propagation to the chamber structural wall, resulting from the deposition of x-ray, debris, and neutron energy in the solid Li_2O pebbles, produces only moderate stresses (~ 10 MPa). Stresses produced due to the weight of breeding material and chamber rotation (~ 30 rpm) are only 8 MPa. These low stresses permit flexibility in the design of the chamber wall. Presently available steel may be used with nominal external cooling, or more exotic materials may be used without external cooling.

Some breakup of the pebbles due to deposition of x-ray, debris, and neutron energy is expected, but this will produce only small effects on the operation of the chamber and will actually ease the tritium removal. Although Li_2O was selected because of its high tritium breeding ratio (TBR) of 1.35, other solid breeders (e.g., LiAlO_2 with a TBR of ~ 0.9) are feasible if a neutron multiplier is incorporated.

A number of issues still need to be resolved

- We need to establish the actual shape of the chamber wall.
- We may have to conduct experiments to quantitatively determine erosion of the chamber wall by the solid Li_2O pebbles; however, we are encouraged by the fact that large gravel-crushing and ore-grinding chambers have reasonable lifetimes.
- We need to determine limits to the chamber repetition rate.

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- We need an economical tritium recovery method to minimize the tritium inventory in the solid Li_2O blanket.
- We are examining potential options for power conversion equipment to take advantage of the high-temperature potential of the chamber.

In summary, the Cascade concept offers a new option for developing safe, efficient, and flexible inertial-fusion reaction chambers for future commercial-power generation. In the following six articles, we describe some initial calculations associated with the Cascade chamber and its balance-of-plant. Four studies address the viability of the Cascade chamber. First, we investigated whether a TBR greater than unity could be obtained. We selected Li_2O as the blanket material because this study showed it gave a TBR of 1.35. Second, we evaluated the buildup of carbon-14 in the Li_2O blanket to assure ourselves that activation from this source would be low. Third, we conducted a stress analysis of the Cascade chamber wall and found a consistent design could be made using materials that are available today. Fourth, we designed a vacuum system that could be used to satisfactorily pump the chamber in the 0.2-s time span available between consecutive pulses.

In the last two articles, we consider the conversion efficiency of thermal energy to electricity. Cascade's high-temperature chamber gives us the potential for thermal efficiencies of ~60%. The trade-off occurs

in higher equipment costs to reach the higher values of efficiency. We studied three 1000-MW_e alternatives for the balance-of-plant that resulted in efficiencies ranging from ~30 to 55%. We also documented a parametric analysis on the highest efficiency cycle, including a high-temperature helium turbine and a steam-bottoming cycle.

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Neutronics Studies. Several solid breeding materials were initially considered for use in the Cascade chamber. Neutronics calculations²⁵ were run to compare the nuclear characteristics of five different materials: Li_2O , Li_2CO_3 , LiAlO_2 , LiF , and Li_8ZrO_6 . Table 8-2 gives the densities and atom fractions for each compound. Natural lithium (7.42 at.% ^6Li , 92.58 at.% ^7Li) is used throughout.

The neutronics model used for these scoping studies is shown in Fig. 8-11. The source consists of 14.1-MeV neutrons born in a compressed DT region with a ρr of 3.0 g/cm². A spherical shell of breeding material extending from 4.0 to 5.0 m, surrounds the fuel and is subdivided into five zones as shown. The 1-cm-thick innermost zone is included to obtain the peak neutron heating rate. The entire blanket is at one-half normal density to represent the packing of particles in the Cascade chamber. The effective blanket thickness is thus 0.5 m. The steel chamber wall and beam penetrations are not included in these calculations.

The results for the five different breeding materials are summarized in Table 8-3; Li_2O has the highest tritium breeding ratio (TBR) at ~1.36, and LiAlO_2 has the lowest at 0.87. If beam ports were included, the breeding performance of Li_2CO_3 might be marginal. The neutron leakage, N_L , gives an indication of the adequacy of a 0.5-m-thick blanket. A thicker blanket would retain some of the leakage neutrons and increase the TBR. The number of neutrons that reenter the central vacuum region per DT neutron, N_R , gives an indication of the seriousness of neutron leakage through chamber ports.²⁶

Table 8-2. Breeding material characteristics.

Material	Density (g/cm ³)	Element	Atom fraction (%)
Li_2O	2.01	^6Li	4.95
		^7Li	61.72
		O	33.33
Li_2CO_3	2.11	^6Li	2.47
		^7Li	30.86
		C	16.67
		O	50.00
LiAlO_2	2.55	^6Li	1.86
		^7Li	23.15
		Al	25.00
		O	50.00
LiF	2.63	^6Li	3.71
		^7Li	46.29
		F	50.00
Li_8ZrO_6	2.99	^6Li	3.96
		^7Li	49.38
		Zr	6.67
		O	40.00

Fig. 8-11. In the one-dimensional neutronics model for Cascade, the breeding blanket is subdivided into five zones.

Table 8-3. Summary of neutronic results.

Material	T_6	T_7	T_T	N_L	N_R	E_n	M_s
Li_2O	0.950	0.407	1.357	0.066	1.364	16.09	1.113
Li_2CO_3	0.870	0.202	1.072	0.103	1.545	15.65	1.088
LiAlO_2	0.742	0.123	0.865	0.157	1.751	15.61	1.086
LiF	0.987	0.266	1.253	0.040	1.517	17.28	1.181
Li_8ZrO_6	0.963	0.293	1.256	0.077	1.593	16.72	1.149

The following per DT neutron:

$T_6 = {}^6\text{Li}(n, T)\alpha$ reactions

$T_7 = {}^7\text{Li}(n, n'T)\alpha$ reactions

$T_T = T_6 + T_7$

N_L = neutron leakage from blanket

N_R = neutrons reentering vacuum region

E_n = total neutron energy deposited plus energy leakage

M_s = system energy multiplication = $(E_n + 3.5)/17.6$

Zone	Volume ^a (m ³)	Energy deposition (MeV)				
		Li_2O	Li_2CO_3	LiAlO_2	LiF	Li_8ZrO_6
1	2.016	0.453	0.390	0.414	0.562	0.450
2	51.46	8.170	7.272	7.441	9.454	8.271
3	60.15	3.779	3.778	3.656	3.741	4.020
4	67.22	1.366	1.628	1.521	1.235	1.523
5	74.68	0.414	0.560	0.542	0.367	0.487

^aThis volume contains the breeding material at one-half normal density.

All these values of N_R are comparable and will result in port leakage on the order of 2 to 3 times the port solid-angle fraction. The value of E_n includes the neutron energy deposited in both target and blanket plus the neutron energy leakage. The resulting system energy multiplication factors range from ~ 1.09 to 1.18, with LiF giving the highest value.

The energy deposition by zone is given in Table 8-4 for the different breeding materials. (The neutron energy deposited in the target zone is approximately the same in each case and is equal to ~ 1.86 MeV.) These numbers can be used to calculate the specific energy deposition per pulse as follows

$$e = \frac{E_D}{17.6} \times \frac{2}{\rho V} \times Y \times 10^6 \frac{\text{J}}{\text{g}}, \quad (2)$$

where

E_D = neutron energy deposited in the zone MeV/DT reaction

ρ = material density (g/cm³)

V = volume of the zone (cm³)

Y = fusion energy yield per pulse (MJ).

The first term ($E_D/17.6$) is the fraction of the fusion energy deposited in the zone; the second term ($2/\rho V$) is the mass of material in the zone, and the factor 10^6 converts megajoules to joules. Li_2O has the highest specific energy deposition and Li_8ZrO_6 has the lowest.

After selecting Li_2O as the breeding material, we carried out a two-dimensional

neutronics calculation for the Cascade chamber. The model is shown in Fig. 8-12. The Li_2O blanket and steel vessel wall are concentric, ellipsoidal shells. The Li_2O region is at one-half normal density and is 1.0 m thick (radially) at $Z = 0$. The Fe vessel wall is 5.0 cm thick. The elliptical surfaces intersect the axis of rotation at 7.00, 8.00, and 8.05 m. Two conical penetrations along the axis of rotation represent the beam ports. The enclosed half-angle is 8.11° , which gives a total solid angle fraction of 1% for the two beams. A ρr -3 target is located at the origin. Material characteristics are the same as listed in Table 8-2.

Table 8-4. Distribution of energy deposition.

Fig. 8-12. The two-dimensional neutronics model for Cascade represents the chamber geometry and includes beam ports.

Tritium breeding	
T_6	0.947
T_7	0.402
Total	1.349
Energy deposited (MeV)	
Target	1.86
Li ₂ O blanket	14.06
Fe wall	0.04
Total	15.96
Energy leakage (MeV)	
Beam ports	0.15
Vessel wall	0.04
Total	0.19
Neutron leakage	
Beam ports	0.027
Vessel wall	0.039
Total	0.066

Table 8-5. Neutronic results for two-dimensional calculation.

The results are given in Table 8-5. The chamber has a tritium breeding ratio of ~ 1.35 , with ~ 0.95 from ${}^6\text{Li}$. The total energy deposition is 15.96 MeV per DT reaction, which gives a system energy multiplication of 1.11. The energy lost through the beam ports is dominated by the direct loss of 14.1-MeV neutrons. The number of neutrons leaking through the ports is 2.7 times the solid angle fraction, as expected.

Author: W. R. Meier

Carbon-14 Generation in the Li₂O Blanket of Cascade. Lithium oxide (Li₂O) as a fusion reactor blanket material will generate radioactive carbon-14 (${}^{14}\text{C}$) through neutron reactions with oxygen. If the quantities of ${}^{14}\text{C}$ thus produced are large, there obviously would be serious safety implications for Cascade and other fusion reactor concepts in which Li₂O is contemplated as a safe blanket material.

Using the method and symbology of Ref. 27 to estimate the time-dependent ${}^{14}\text{C}$ radioactivity from neutron reactions with the Li₂O blanket of a Cascade reactor on a per-GW_e basis,²⁷ we found that ${}^{14}\text{C}$ generation in Li₂O is possible by one first-ordered sequence (sequence 1) and four second-ordered sequences (sequences 2 through 5)

1. ${}^{17}\text{O}(n,\alpha){}^{14}\text{C}$
2. ${}^{16}\text{O}(n,\gamma){}^{17}\text{O}(n,\alpha){}^{14}\text{C}$
3. ${}^{16}\text{O}(n,\alpha){}^{13}\text{C}(n,\gamma){}^{14}\text{C}$
4. ${}^{18}\text{O}(n,2n){}^{17}\text{O}(n,\alpha){}^{14}\text{C}$
5. ${}^{18}\text{O}(n,\alpha){}^{15}\text{C}(n,2n){}^{14}\text{C}$.

Of these reactions, only the ${}^{17}\text{O}(n,\alpha)$, the ${}^{16}\text{O}(n,\alpha)$, and the ${}^{13}\text{C}(n,\gamma)$ reactions have been computed and entered into ACTL, the evaluated neutron-activation cross section library.²⁸

To obtain the reaction-rate constants λ (the reaction probabilities per nucleus per unit time), we used a spherical neutronics model of Cascade with a homogeneous 14-MeV neutron source in a compressed target at the center of the chamber; a spherical shell of 50% solid-density Li₂O at a radius of 4 to 5 m; and a 5-cm steel shell beginning at a radius of 5 m. The steel shell is necessary to account for reactions from reflected neutrons; if it is not included, the number of reactions is about 2% lower. Because Cascade is oblong in shape, the actual ${}^{14}\text{C}$ activities will be on the order of 20% greater than calculated.²⁹ A source of 9.59×10^{20} (fusion) neutrons per second was assumed, which corresponds to a fusion power of 2700 MW_f, or a net electric power of ~ 1 GW_e. The reaction rates scale directly with fusion power because the sequences of second order in λ have little significance.

For a reactor operation time of 30 years, we may ignore burnup of all nuclides because, of all carbon and all oxygen isotopes, no total reaction rate exceeds $1.71 \times 10^{12}/\text{s}$, which corresponds to a $1/e$ burnup time of 18 844 years. In addition, since the half-life of ${}^{15}\text{C}$ (2.45 s) is much less than the smallest burnup time, we may equate the ${}^{15}\text{C}$ loss-rate constant, Λ , of ${}^{15}\text{C}$ with its radioactive-decay constant, $\lambda^{15}\text{C}_{\text{eff}}$. With a half-life of 5730 years, and a 30-year reactor operation time, the loss by decay of ${}^{14}\text{C}$ can also be neglected to an accuracy of less than 30/5730 (less than 1%).

Reaction rates for which the cross sections do not exist in the ACTL library are contained in Table 8-6. We have also listed there the substitutions that we used. It will be seen, after computation, that an error here of 2 or 3 orders of magnitude for the first three listed reactions will make little difference; an error many orders of magnitude will also make little difference for the fourth reaction. Examination of the λ values for all nuclides shows such large variances to be unlikely.

Using the method described in Ref. 27, and with these assumed λ values, the contributions to the population of ^{14}C from each sequence are:

1. $N(t=0)_{^{17}\text{O}} \lambda_{^{17}\text{O}(n,\alpha)} T = 2.01 \times 10^{24}$ nuclei,
2. $N(t=0)_{^{16}\text{O}} \lambda_{^{16}\text{O}(n,\gamma)} \lambda_{^{17}\text{O}(n,\alpha)} T^2/2$
 $= 9.69 \times 10^{21}$ nuclei,
3. $N(t=0)_{^{16}\text{O}} \lambda_{^{16}\text{O}(n,\alpha)} \lambda_{^{17}\text{C}(n,\alpha)} T^2/2$
 $= 2.11 \times 10^{20}$ nuclei,
4. $N(t=0)_{^{18}\text{O}} \lambda_{^{18}\text{O}(n,2n)} \lambda_{^{17}\text{O}(n,\alpha)} T^2/2$
 $= 5.18 \times 10^{20}$ nuclei,
5. $N(t=0)_{^{18}\text{O}} \lambda_{^{18}\text{O}(n,\alpha)} \lambda_{^{15}\text{C}(n,2n)} T/\Delta^{15}\text{C}$
 $= 1.62 \times 10^{12}$ nuclei.

This gives a total of 2.02×10^{24} nuclei of ^{14}C , which corresponds to a ^{14}C radioactivity of 210 Ci/GW_e after 30 years of reactor operation.

We can expect that sequences of each higher order will, in sum, make negligible contributions relative to those of the next lower order (of the order of the ratio of the sum of the contributions of sequences 2 through 5 to the contribution of sequence 1, or 0.0052). Neglect of all other sequences therefore causes an error on the order of $0.0052 + 0.0052^2 + 0.0052^3 + \dots = 0.0052/(1 - 0.0052)$, or approximately 0.5%.

We can therefore conclude that, even on time scales of 30 years, ^{14}C generation from a pure Li_2O blanket in a spherical Cascade reactor is on the order of 210 Ci. Accounting for approximately a 20% increase because of Cascade's oblong shape, we may expect 250 Ci. Compared to the probable hazards from activated impurities in Li_2O , from tritium, and from activation of surroundings (steel, etc.), the hazard from ^{14}C radioactivity in the Cascade reactor would

Unlisted reaction	Substitute reaction	Notes
$^{16}\text{O}(n,\gamma)$	$^{12}\text{C}(n,\gamma)$	Both even-even. Both tight-shell closures (integer number of α particles in nucleus). No isotones of ^{16}O listed in library.
$^{18}\text{O}(n,\alpha)$	$^{16}\text{O}(n,\alpha)$	Isotopes, both even-even. (The Coulomb barrier, which is primarily a function of the number of protons, is important here, as well as the pairing potentials.)
$^{18}\text{O}(n,2n)$	$^{19}\text{F}(n,2n)$	Isotones. (Same neutron energy levels and shell closures.)
$^{15}\text{C}(n,2n)$	$^{13}\text{C}(n,2n)$	Nearest available similar reaction.

be negligible. (The latter two sources are on the order of 10 and 100 MCi, respectively, in the HYLIFE reactor at 30-year shutdown, which, in turn, will have more than an order of magnitude less radioactivity than any magnetic fusion reactor of the same power output and having a steel first wall.³⁰)

At long decay times (over 100 years), ^{14}C may become a significant relative contributor of total radioactivity, but it emits no gamma rays (and is thus easily shielded), and it does not concentrate in any particular organ or gland of the human body. Also, it is easily formed into stable solid compounds, and its energy of decay is relatively low. (Accounting for neutrino energy loss, the average β^- energy is less than 0.1 MeV.) Nevertheless, since all waste must meet a shallow burial criterion based on current congressional and federal regulations after some specified period of decay, ^{14}C will probably remain a concern.

Because the major contribution of ^{14}C is from sequence 1 and varies linearly with time, its contribution relative to those that vary quadratically in time (sequences 2, 3, and 4) is even greater at reactor operation times of less than 30 years. Therefore, our assumptions of neglecting oxygen burnup, ^{14}C decays, and transmutation sequences of order >2 are only strengthened. If fresh Li_2O is cycled into the Cascade reactor, the contributions of the quadratic sequences are again diminished. We can therefore expect the total ^{14}C radioactivity generated by the Cascade reactor in its Li_2O blanket to be approximated by $250(T/30)(P/2.7)$, or $3.17P$ Ci, in which T is the total time of full-power operation in years, and P is the fusion power in GW_f.

Author: G. P. Lasche

Table 8-6. Reaction rates and substitutions for cross sections not included in ACTL.

Stress Analysis of the Cascade Chamber Wall.

We calculated the stress³¹ in a 25-mm-thick 2-1/4 Cr-1 Mo steel wall with a 5-m maximum radius perpendicular to the axis of rotation, and we found the magnitude to be satisfactory. The design can withstand the environment of the 600-MJ, 5-Hz Cascade chamber, including exposure to a 1200 K Li₂O pebble blanket, for the entire 30-year lifetime and can be built with techniques available today.

We made our calculations based on the methodology for a consistent wall design developed for the HYLIFE chamber.³² We combined the steady or quasi-steady wall stresses—caused by centrifugal action and temperature differences—with the transient stresses caused by x-ray, debris, and neutron energy deposition. All stresses are shown to comply with the intent and methodology of the American Society of Mechanical Engineers (ASME) code.³³

To evaluate centrifugal stress, we note that the rotational speed of the chamber is set so that Li₂O pebbles at the minimum radius perpendicular to the axis of rotation will be held against the wall. The centrifugal acceleration, a , at the minimum radius must equal that of gravity. Our minimum radius, $R_{\min}$, is 1 m, so that the rotational speed, ω , becomes

$$\omega = \sqrt{a/R_{\min}} = \sqrt{\frac{9.8 \text{ m/s}^2}{1 \text{ m}}} = 3.13 \text{ rad/s (30 rpm)} \quad (3)$$

Maximum stress due to centrifugal action occurs at the maximum radius, R_a , perpendicular to the axis of rotation. To calculate this stress, we first find the wall pressure, P , caused by centrifugal action, and next the stress, σ , in the prolate spheroid-shaped chamber caused by the pressure³⁴

$$p = \frac{F}{A} = \frac{ma}{A} \quad (4)$$

where

F = force due to centrifugal action (N)
 A = area (m²)
 m = mass (kg)
 a = acceleration (m/s²)

Consider a strip at this maximum radius of unit length in the direction of the axis of rotation, then

$$P = \frac{\omega^2 \left[\rho_{\text{Li}_2\text{O}} (R_a^3 - R_i^3)/3 + \rho_{\text{steel}} R_a^2 t \right]}{R_a} \quad (5)$$

$$\begin{aligned} \sigma &= \frac{P R_a (2R_b^2 - R_a^2)}{2t R_b^2} \\ &= \frac{1}{2t R_b^2} \left\{ \omega^2 \left[\rho_{\text{Li}_2\text{O}} (R_a^3 - R_i^3)/3 \right. \right. \\ &\quad \left. \left. + \rho_{\text{steel}} R_a^2 t \right] (2R_b^2 - R_a^2) \right\} \\ &= 8.0 \text{ MPa} \quad (6) \end{aligned}$$

where

$\rho_{\text{Li}_2\text{O}}$ = density of Li₂O pebble blanket (1000 kg/m³)
 ρ_{steel} = density of steel (7800 kg/m³)
 R_a = maximum radius perpendicular to the axis of rotation (5 m)
 R_b = maximum radius parallel to the axis of rotation (8 m)
 R_i = maximum inner radius of Li₂O pebble blanket perpendicular to the axis of rotation (4 m)
 t = wall thickness (0.025 m).

Calculation of thermal stress requires a heat-transfer analysis to determine the temperature difference across the wall. We selected a Li₂O pebble-blanket temperature of 1200 K, and an inside-wall surface temperature of 800 K, and we then used Wunschmann and Schlünder's correlation³⁵⁻³⁷ to determine heat transfer. We found that the heat-transfer coefficient between the Li₂O pebble blanket and the wall is dominated by thermal radiation. Heat flux into the wall is substantial—equal to 63 kW/m². The total heat-transfer rate through the Cascade wall is 28 MW.

The temperature difference, ΔT_w , across the wall and the thermal stress were found to be 55 K and 132 MPa, respectively. Of the 60 K temperature difference, 55 K is caused by heat transfer from the Li₂O pebble blanket to the wall, and 5 K is due to internal heat generation. The outside surface of the wall needs to be maintained at

740 K by convective heat transfer. This can be accomplished with the use of liquid metals, water, or perhaps even helium gas.

In summarizing, the hoop stresses in the wall may be divided into transient and steady (or slowly varying) portions. Transient stresses occur following each fusion reaction as a result of energy deposition from x rays, debris, and neutrons. These transient stresses³² are shown in Fig. 8-13. At the design point of 600 MJ and 5 Hz, the total is between 5.4 and 12.7 MPa, depending on the assumption of the type of geometry. In addition, we showed³² that, if we make the most conservative assumption that there is no heat transfer between vaporized materials and the Li₂O pebble blanket, the vaporized material would exert a pressure on the wall resulting in a stress of 14.4 MPa for a prolate spheroid with major and minor semiaxes of 8 m and 5 m, respectively. However, by the time the pressure builds up, the stresses shown in Fig. 8-13 would have decayed. Hence, they are not additive, and do not occur simultaneously.

Coupling all of the stresses together results in the values shown in Table 8-7. In using these stresses to determine allowable limits, we need to separate primary stresses caused by a force or pressure from secondary thermal stresses. Secondary stresses relieve themselves with plastic deformation, whereas primary stresses do not.

Allowed stresses are based on the ASME pressure-vessel code,^{32,33} which cites four criteria.

The first criterion involves creep-fatigue evaluation. Determining the number of cycles necessary to cause a failure under the fluctuating stress conditions present in the Cascade reaction chamber. If the time when a steady creep load exists is negligible (as in our case), it is required that

$$\sum_{i=1}^n \left(\frac{n}{N_{D_i}} \right) \leq 1 \quad (7)$$

where n is the number of applied cycles of loading condition i , and N_{D_i} is the number of design allowable cycles of loading condition i permitted for an equivalent strain range, $\Delta\epsilon_{\text{equiv}}$. Following the methods used

Fig. 8-13. Cascade chamber wall hoop stress due to x-ray, debris, and neutron energy deposition increases with pellet yield.

Table 8-7. Summary of the inside-surface wall hoop stress.

Source of stress	Level of stress (MPa)
Centrifugal action ^a	8.0
Thermal stress ^b	132.0
Pressure from vaporized material ^{c,d}	14.4
Shock caused by x rays, debris, and neutrons (600-MJ, 5-Hz chamber) ^e	
Cylindrical geometry limit	12.7
Spherical geometry limit	5.4
Total stress	
Cylindrical geometry	153 to 155
Spherical geometry	145 to 15 ⁺

^aPrimary stresses.

^bSecondary stress (relieved with plastic deformation).

^cStress caused by pressure from ablated material assumes no heat transfer. The value included here is twice the static value, which accounts for dynamic loading.

^dPeak stress caused by pressure from ablated material occurs at a later time 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 two limits—cylindrical and spherical. Our value for Cascade is between these two limits.

in Ref. 32 with Booker's fatigue data,³⁸ and accounting for a nonzero average stress, we show that the allowable number of cycles exceeds 10^{12} , using worst-case conditions shown in Table 8-7. The Cascade reaction chamber has a lifetime of 5×10^9 cycles, which is orders of magnitude below the allowed value. Hence this criterion offers no limitation. However, because of the logarithmic nature of the data, a small increase in magnitude of the fluctuating loads

substantially decreases the allowed number of cycles. If the stress due to pressure is increased from 14.4 to 46 MPa, the actual lifetime would equal the allowed lifetime.

The second criterion is associated with maximum strain limits. If an elastic analysis is used, the sum of the primary and secondary stresses must be less than the yield strength. The yield strength³³ at the inside wall temperature of 800 K is 166 MPa, which exceeds the worst-case conditions shown in Table 8-7 by only 7%. It is this criterion that limits the design of the Cascade chamber.

The third criterion applies to the allowable time a primary stress can be present under operating conditions. Again following the methods of Ref. 32, we find that actual times the primary stress levels exist are below the allowable times, so that this criterion offers no restriction.

A fourth criterion applies to steady primary loading conditions. With the exception of the stress resulting from centrifugal action, all primary loads on the Cascade chamber are fluctuating. The steady 8.0-MPa stress due to centrifugal action is well below the allowed 63 MPa at our inside wall temperature of 800 K and offers no restriction.

In conclusion, we find that a consistent design for the Cascade chamber that conforms to the intent and methodology of the ASME code results in the following limits:

- Radius at midsection—5 m.
- Wall thickness—25 mm.
- Li₂O pebble-blanket temperature—1200 K.
- Steel used—2-1/4 Cr-1 Mo.

A number of variations are possible to reduce the size and cost of the chamber. Decreasing the wall thickness would reduce thermal stresses and permit a smaller-radius chamber. Lower blanket temperature would also reduce thermal stresses below the calculated value of 132 MPa and permit a smaller-radius chamber. Different materials, such as silicon carbide, could be used to eliminate cooling the chamber wall. We are examining these and other options. However, the present design—while not necessarily optimum—is a workable option that is conservative and could be built today with compliance to the ASME code.

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Vacuum Pumping of the Cascade Reaction Chamber. There are two convenient locations for vacuum pumping the Cascade chamber—the aperture at each end of the horizontal axis. Calculations³⁹ show that these apertures are adequate for pumping a 3000 MW_t chamber to a pressure of at least 10^{-2} Torr, while pressures of 3×10^{-3} Torr can be reached, but with some difficulty. This pressure range is more than adequate for laser or light-ion-beam drivers.⁴⁰

On a steady basis, noncondensable gas that must be pumped from the Cascade chamber comes from unburned DT gas in the pellets, from helium generated in the fusion reactions, and from helium and tritium produced in the blanket by neutron capture with lithium. It can be argued that the noncondensable gas produced in the blanket is trapped in the solid Li₂O pebbles. However, because the temperature of these pebbles is high, diffusion of both helium and tritium out of the pebbles is expected to be rapid. At any rate, inclusion of the noncondensable gas produced in the blanket offers an upper limit for vacuum pumping requirements.

To calculate the noncondensable gas throughput, we must first know (1) the burn fraction for DT gas injected into the chamber, (2) the chamber thermal power, and (3) the tritium breeding ratio. The burn fraction, f_b , can be calculated⁴¹ by

$$f_b = \frac{\rho r}{6.3 + \rho r} \quad (8)$$

where, for reactor pellets, the expected product of density and radius (ρr) is about 3 g/cm². Hence, a typical burn fraction is 0.32.

From the DT reaction

we see that for each D + T burned, a total of 17.6 MeV is released and one He atom (from the α particle) is produced. The total DT burned for a 3000-MW_t chamber is 8.9×10^{-6} kg/s. The injected and unburned DT gas and the He gas produced in the DT reaction are then

$$\begin{aligned} \text{DT gas injected} &= \frac{\text{DT gas burned}}{f_b} \\ &= 2.8 \times 10^{-5} \text{ kg/s} \end{aligned}$$

DT gas unburned = 1.9×10^{-5} kg/s

He gas produced in the DT reaction
= 7.1×10^{-6} kg/s.

Noncondensable gases generated in the Li_2O blanket can be calculated easily if the TBR is known. Meier calculated a TBR for Li_2O in the Cascade chamber as ~ 1.4 . (See the article, "Neutronics Studies," earlier in this section.) There are, then, ~ 1.4 atoms each of tritium and helium generated in the blanket for each tritium atom burned, or

Tritium generated
in the Li_2O blanket = 7.5×10^{-6} kg/s
Helium generated
in the Li_2O blanket = 9.9×10^{-6} kg/s.

Using the perfect gas law, we calculate that the throughput of DT and helium gas is 0.39 and 0.26 Torr-l/s K. If we assume the worst case, where the gas temperature is equal to the Cascade pebble outlet temperature of 1200 K, and where the typical chamber pressure for laser drivers is 0.1 Torr, an effective pumping speed of 4700 l/s for DT and 3100 l/s for helium is required.

The effective pumping speed is decreased by the conductance of the duct connecting the vacuum pump to the chamber by the relationship

$$\frac{1}{S_e} = \frac{1}{S_p} + \frac{1}{C} \quad (10)$$

where

S_e = effective pumping speed, l/s

S_p = vacuum pump speed, l/s

and the duct conductance, C , in l/s is found (Ref. 42) with

$$C = \frac{3.64 A (T/M)^{1/2}}{1 + (3/16)(HL/A)} = \frac{\pi a^4 P_{\text{ave}}}{8\eta l} \quad (11)$$

molecular flow viscous flow

Here, A is the area of the duct in cm^2 , H is the perimeter in cm, L is the length in cm, M is the molecular weight of the gas, T is the absolute temperature in K, a is the duct radius in cm, R is the viscosity in poise, and P_{ave} is the average pressure in Torr.

The relationship between effective speed, pump speed, and duct conductance is plotted in Fig. 8-14. Note that, if the duct con-

Fig. 8-14. Relationship between effective speed, pump speed, and duct conductance.

ductance is greater than 10 times the required effective pumping speed, then the pump speed needed is essentially equal to the effective speed. As the duct conductance is reduced, more and more pumping speed is needed to obtain the required effective pumping speed. When the duct conductance is twice the effective speed, the pump speed needed is twice the effective speed.

Typical pumping systems used for removal of hydrogen isotopes and helium from conceptual fusion reactor designs contain both Roots blowers and cryogenic surfaces operating at a liquid helium temperature of 4.2 K. The University of Wisconsin⁴³ suggests that cryopumps have speeds of $5 \text{ l/s} \cdot \text{cm}^2$ of cryosurface for DT and $2 \text{ l/s} \cdot \text{cm}^2$ for helium.

For a duct temperature of 1200 K, we need a pumping system capable of continuously removing 310 Torr l/s of helium and 470 Torr l/s of DT. If the chamber pressure is 10^{-2} Torr (a value well below that required for a laser driver), the effective pumping speeds needed are 31 000 l/s for helium and 47 000 l/s for DT. Consider two 1-m-radius, 2.5-m-long ducts emanating from the Cascade chamber. For our conditions, the ducts operate in the viscous flow regime, and, for an average duct pressure of 10^{-2} Torr, the conductance per duct, based on calculations using Eq. (11), is 4×10^6 l/s for DT and 7.5×10^6 l/s for helium. The ratio of duct conductance to effective speed is large, and the ratio of pumping speed to effective speed, $S_p/S_e \approx 1.0$. Hence, the necessary pump speeds are about 31 000 l/s for helium and 47 000 l/s for DT.

If we conservatively assume that no Roots blowers are used, pumps with a total

cryosurface area of $31\,000/2 + 47\,000/5 = 25\,000\text{ cm}^2$ would then be needed. Even with this conservative assumption, such an area is feasible from an engineering standpoint.

If the chamber pressure is reduced to 10^{-3} Torr, molecular flow dominates, and the effective pumping speed required exceeds the total conductance of the two ducts discussed above; the necessary pumping speed then becomes infinite. If these two ducts only are used, a practical lower limit to chamber pressure would be 3×10^{-3} Torr. Reduction of chamber pressure below this number would require a larger total duct area.

Two conclusions emerge from this study. First, a vacuum system capable of reaching chamber pressures of at least 10^{-2} Torr can be constructed without difficulty using the two apertures in the ends of the 3000-MW_e Cascade chamber. At chamber pressures less than 3×10^{-3} Torr, additional duct area to the vacuum system pumps will be required. Laser drivers require about 10^{-1} Torr chamber pressures. Second, throughput for a 3000-MW_e Cascade chamber is 310 Torr l/s of helium and 470 Torr l/s of DT if the gas temperature is 1200 K. Additional transient pumping to accommodate outgassing during startup may be necessary; this can be easily incorporated at chamber pressures above 10^{-2} Torr.

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Three Alternatives for the Cascade Balance-of-Plant. To evaluate the net energy production of the Cascade reactor chamber, we have considered several balance-of-plant alternatives. We report here on three 1000-MW_e alternatives: (1) a steam cycle, (2) steam with a helium-topping cycle, and (3) helium with a steam-bottoming cycle (Figs. 8-15 through 8-17, respectively). The steam cycle alternative adopts the technology developed for the Solase conceptual laser fusion reactor design⁴⁴ and combines it with the Cascade chamber. Net plant efficiency (electrical power to the buss bar

divided by the reaction chamber thermal power) is 30%. No attempt is made to capitalize on the high-temperature potential of the Li₂O. The other two alternatives add a helium cycle and use the high-temperature potential of the Li₂O. In each case, we have modified Li₂O temperature in the chamber to accommodate balance-of-plant requirements.

In the steam cycle alternative (Fig. 8-15), Li₂O exits the Cascade chamber at 870 K, thereby easing cooling problems that could be encountered with the chamber wall. After leaving the chamber, the Li₂O flow is split and falls through both steam generators and reheaters. The designs use a standard counterflow shell-and-tube with an overall heat-transfer coefficient⁴⁴ of 770 W/m²K and a log mean temperature difference⁴⁵ of 100 K. The controlling heat-transfer conductance⁴⁴ occurs between the Li₂O particles and the outside tube wall; it equals 850 W/m²K. Total heat transfer areas are 34 000 m² and 13 000 m² for the steam generators and reheaters, respectively. Energy is extracted from the steam in turbine-generators that produce 1300 MW_e gross and 1000 MW_e net, yielding a net plant efficiency of 30%.

Tritium outgasses from the Li₂O and is removed by the vacuum system. Some small quantities of tritium will diffuse through the steam generator and reheater tubes to form tritiated water in the secondary loop. Should this be a problem, double-walled tubes or alternative ways of reducing tritium permeation into the secondary loop would be required.

The second alternative, steam with a helium-topping cycle (Fig. 8-16), has two distinct advantages over the straight steam cycle. First, it utilizes the high-temperature potential of the Li₂O particles, and, second, it includes fluidized-bed heat-exchangers that increase the overall heat-transfer coefficient. One fourth of the energy produced in the Cascade chamber is removed in the helium-topping cycle, and the remainder is removed in the steam cycle. The Li₂O particles exit from the chamber at 1200 K and then pass through a batch valve separating the chamber, where the pressure is 10^{-2} Torr, from a holding tank, where the pressure is near atmospheric.

Fig. 8-15. Cascade chamber — steam cycle.

A separate flow of helium gas is used to fluidize the Li₂O particles in the heat exchanger. The velocity of the helium must be enough to fluidize the bed, but below that which would result in carryover of particles from the bed. These two limits are calculated⁴⁶ as:

$$u_{ml} = \frac{d_p^2 (\rho_s - \rho_g) g}{1650 \mu} = 0.26 \text{ cm/s} \quad (12)$$

and

$$u_t = \left[\frac{4 (\rho_s - \rho_g) g^2}{225 \rho_g \mu} \right]^{1/3} = 37 \text{ cm/s} \quad (13)$$

where

- u_{ml} = minimum fluidizing velocity, cm/s
- d_p = particle diameter, taken as 0.01 cm
- ρ_s = particle density, taken as 2 g/cm³
- ρ_g = gas density, equal to 0.0029 g/cm³ at 6.1 MPa and 1000 K (Ref. 47)
- g = gravitational constant equal, to 980 cm/s²
- μ = gas viscosity, equal to 4.6×10^{-4} g/cm-s at 6.1 MPa and 1000 K (Ref. 47).

The corresponding flow-rate range is from 2×10^{-4} to 2×10^{-2} kg/s.

Heat transfer for the minimum fluidizing velocity⁴⁶ was calculated to be 1800 W/m²K.

Fig. 8-16. Cascade chamber—steam with a helium topping cycle.

Conductance across the tube wall was retained as established in the Solase study⁴⁴ at 10 000 W/m²K. The inside heat-transfer coefficient for helium⁴⁵ and condensing steam⁴ were found to be 5000 and 59 000 W/m²K, respectively. This results in overall heat-transfer coefficients of 1200 W/m²K in the helium section and 1500 W/m²K in the steam section of the heat exchanger. Heat-transfer areas are 2300 m² for the helium-topping cycle and 4900 m² for the steam cycle. In addition, the steam-cycle heat exchanger has an area of 2700 m², which makes a grand total of 9900 m². Net efficiency for the helium-topping cycle⁴⁸ is 44%, yielding, when combined with the steam cycle, a net plant efficiency of 36%.

An additional advantage to the fluidized-bed approach is that this small flow rate

will contain essentially all of the tritium that diffuses out of the Li₂O as it passes through the heat exchanger. Tritium can be separated from the fluidizing helium, thereby offering additional protection to the secondary loop fluids.

The third alternative proposed by Meier⁴⁹ uses helium with a steam-bottoming cycle (Fig. 8-17). This cycle, with a net plant efficiency of 55%, obtains 42% from the helium cycle and 13% from the steam-bottoming cycle. Total heat-transfer area is 103 000 m², including the recuperator. If the steam bottoming is eliminated to save capital expense, the heat-transfer area required drops to 61 000 m².

A comparison of the three alternatives is shown in Table 8-8. The heat-transfer area of the steam cycle (alternative 1) could be

Fig. 8-17. Cascade chamber—helium with steam-bottoming cycle.

Table 8-8. Comparison of three alternatives for a 1000-MW_e Cascade chamber balance-of-plant.

Alternative	Net plant efficiency (%)	Total ^a heat exchanger areas (m ²)	Chamber thermal power (MW _t)
(1) Steam cycle	30	47 000	3340
(2) Steam with He-topping cycle	36	9 900	2800
(3) He with steam-bottoming cycle (Meier cycle)	55	103 000	1820

^aNot including condenser to remove waste heat.

decreased if a fluidized bed is used; however, it is limited in the maximum temperature permitted by the steam cycle. Adding a helium-topping cycle (alternative 2) reduces the total area of the heat exchangers and increases the efficiency with added benefits of lower chamber power, lower pellet yield, and lower chamber-wall stresses. Maximum efficiency is obtained with alternative 3, but, with larger heat exchangers and more equipment, cost is increased.

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Parametric Analysis of a Power-Conversion System for Cascade. To evaluate efficiency for the high-temperature

power conversion of the Cascade reactor, we conducted parametric analyses of a closed-cycle helium turbine system. This system contains a regenerator and is water cooled; if the system proves cost effective, a steam-bottoming cycle can be added. Our parametric analyses of the gas cycle

Fig. 8-18. A steam-bottoming cycle can be coupled to a high-temperature gas cycle.

Fig. 8-19. Gas-cycle efficiency increases with the cycle peak temperature.

examine the effects of peak temperature, regenerator effectiveness, and system pressure loss on the cycle efficiency.

The power-conversion cycle is illustrated schematically in Fig. 8-18. Low-pressure, low-temperature He is first compressed, then preheated in the regenerator before entering the $\text{Li}_2\text{O}/\text{He}$ heat exchanger. High-temperature (T_2), high-pressure He exits the heat exchanger and expands through the gas turbine, providing the shaft work. After leaving the turbine, the He is initially cooled in the regenerator and then passes through a $\text{He}/\text{H}_2\text{O}$ heat-exchanger that reduces the He to the lowest cycle tempera-

ture (T_1). The $\text{He}/\text{H}_2\text{O}$ heat exchanger could be a steam generator, and the steam could be used as a bottoming cycle, as illustrated.

This cycle configuration is relatively standard^{50, 52}; other components can also be added to increase the efficiency of the cycle. Two ways of increasing efficiency would be to add intercooling in the compression stage or to reheat between the two expansion stages in the turbine. Since the intercooler is a low-temperature component, it is more likely to be cost effective than a reheater. Neither are examined here.

The gas-cycle thermal efficiency (without the bottoming cycle) is given by⁵³

$$\eta_1 = \frac{\eta_c T_2 [1 - |R_p(1 - \delta)|^\alpha] - T_1 (R_p^\alpha - 1)/\eta_r}{\Delta T} \quad (14)$$

where

$$\Delta T = T_2 - T_1 (1 - \eta_r) [1 + (R_p^\alpha - 1)/\eta_r]$$

$$- T_2 \eta_r (1 - \eta_r) [1 - |R_p(1 - \delta)|^\alpha]$$

η_c = turbine efficiency

η_r = compressor efficiency

η_r = regenerator effectiveness

R_p = compressor pressure ratio

δ = fractional system pressure loss

$\alpha = (\gamma - 1)/\gamma$

$\gamma = 5/3$ = ratio of specific heats

T_1 = cycle minimum temperature, K

T_2 = cycle peak temperature, K

In the parametric studies, the pressure ratio was used as the independent variable. In all cases,^{51, 52} the compressor and turbine efficiencies were held constant at 90%, which is typical. The minimum cycle temperature, which is set by the available heat sink (in this case, water), was held at 330 K. Increasing the minimum cycle temperature above 330 K would decrease the cycle efficiency, but also would reduce the required heat-transfer area for the $\text{He}/\text{H}_2\text{O}$ heat exchanger.

The gas-cycle thermal efficiency as a function of the pressure ratio for three different peak temperatures ($T_2 = 1100, 1200$, and 1300 K) is shown in Fig. 8-19. The fractional pressure loss is taken to be 7.5%, which is typical of values cited in the

literature.⁵⁰⁻⁵² The peak thermal efficiency increases from ~38% at 1100 K to ~42% at 1200 K and to ~45% at 1300 K. The pressure ratio at which the peak efficiency occurs increases slightly with increasing peak temperature; it goes from 2.2 at 1100 K to 2.5 at 1300 K. In going to higher temperatures, the trade-off for greater efficiency appears in higher costs and added material problems.

The thermal efficiency curves for different values of regenerator effectiveness are shown in Fig. 8-20. Here, the peak temperature is held at 1200 K, and the system pressure loss is fixed at 7.5%. As seen, the peak thermal efficiency is quite sensitive to the regenerator effectiveness. Here the trade-off is the additional heat-transfer area required to achieve high values of η_r for increased thermal efficiency. The UWMAK-III study⁵² showed that, for its particular design, increasing the regenerator effectiveness from 85% to 90% required nearly 50% more heat-transfer area. Moreover, increasing η_r from 90% to 95% more than doubled the required heat-transfer area. Considering this increase in requirements, the UWMAK-III study selected a value of 88%. Other reports use 90% as typical.^{50,51}

Figure 8-21 shows the thermal efficiency curves for various values of the fractional system pressure loss, where δ = pressure loss in the circuit divided by the compressor exit pressure. Clearly, δ should be minimized. Careful design may be required to minimize pressure loss and operation at a high system pressure; both of these conditions are likely to increase system cost.

For the specific representative parameters given in Table 8-9, the gas-cycle thermal efficiency is 42%. We find that ~58% of the heat is rejected in the He/H₂O heat exchanger, where the temperature falls from 532 K to 330 K. Adding a steam-bottoming cycle to utilize this heat can increase the overall system efficiency. However, the corresponding increase in costs would have to be evaluated.

If we assume that the steam cycle operates between 500 K and 310 K, and if we use a typical estimate for the steam cycle efficiency (60% of the Carnot cycle), then its efficiency will be ~23%. The overall system efficiency is then 55% [i.e., $\eta = 0.23(0.58)$

Fig. 8-20. Gas-cycle efficiency increases with increasing regenerator effectiveness.

Fig. 8-21. Gas-cycle efficiency decreases with increasing system pressure loss.

Compressor efficiency	$\eta_c = 90\%$
Turbine efficiency	$\eta_e = 90\%$
Regenerator effectiveness	$\eta_r = 90\%$
Pressure ratio	$R_p = 2.5$
Fractional pressure loss	$\delta = 7.5\%$
Thermal cycle efficiency	$\eta_t = 42\%$
Temperatures	$T_1 = 330\text{ K}$ $T_2 = 1200\text{ K}$

Table 8-9. Gas-cycle parameters for a representative example.

+ 0.42]. The power plant cost could, therefore, be increased by 31% for the same cost of electricity based on gross electric output. When laser recirculating power requirements are included, the allowable cost difference is greater.

In summary, we can state that the closed-cycle high-temperature helium-power conversion system that we have examined could be coupled to the Cascade chamber. We carried out parametric analyses to evaluate the trade-offs involved in the

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configuration and design of such a system. We found that the He cycle alone can provide a gross thermal efficiency of 42% at a peak-cycle temperature of 1200 K. If a steam-bottoming cycle is added, the overall conversion efficiency then rises to 55%. Cost information on various components and configurations would allow us to assess more completely the economic implications of the parametric variations examined here.

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Neutron-Activation Studies

Influence of Steel Type on Fusion Wall Activation

Introduction. Conceptual designs for fusion reactors need to address the possibility of radioactive buildup in the reactor vessel caused by neutron bombardment. Therefore, we developed a set of computer

codes combined with hand calculations⁵⁴ to determine the potential radioactivity of materials in our conceptual fusion reactors. These codes, illustrated in Fig. 8-22, are also being used by the magnetic fusion and fission waste communities.

This article summarizes a larger study⁵⁵ of the time-dependent radioactivity of several steels that have been proposed for use in fusion reactors. The results are reported in terms of activity and afterheat, of biological hazard potential for inhalation and ingestion, and of feasibility for shallow-burial disposal. Further considerations on the influence of reactor geometry on radioactivity are described in the next article, "Influence of Fusion Reactor Configuration on Neutron-Induced Radioactivity."

Steel Composition. Several steels have been considered for fusion reactor designs. In Table 8-10, we show the composition of the five steels we studied

- **Case N**—The primary candidate alloy (PCA) is an advanced titanium-modified type-316 stainless steel under consideration for first walls in magnetic fusion reactors. It was the material selected for the STARFIRE commercial tokamak design study.⁵⁶

Fig. 8-22. Flow chart for neutron activation and decay calculations.

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Table 8-10. Materials investigated.^a

Element	Case				
	N PCA	O HT-9	P Nb- stabilized 2-1/4 Cr	Q Thermally stabilized 2-1/4 Cr	R Case Q with V replacing Mo
B	0.005	—	—	—	—
C	0.05	0.22	0.09	0.11	0.11
N	0.01	—	—	—	—
Al	0.03	—	—	—	—
Si	0.5	0.38	0.32	0.25	0.25
P	0.01	0.019	0.012	0.005	0.005
S	0.005	0.006	0.007	0.011	0.011
Ti	0.3	—	—	—	—
V	0.1	0.27	—	—	0.94
Cr	14	11.3	2.19	2.22	2.22
Mn	2	0.52	0.46	0.41	0.41
Fe	64.88	85.435	94.261	96.054	96.054
Co	0.03	—	0.0011	0.002	0.002
Ni	16	0.5	0.5489	0.0998	0.0998
Cu	0.02	—	—	—	—
As	0.02	—	—	—	—
Nb	0.03	—	1.13	—	—
Mo	2	0.85	0.98	0.94	—
Ta	0.01	—	—	—	—
W	—	0.5	—	—	—

^aAll compositions are in wt%.

- Case O—A ferritic steel, HT-9 (Ref. 57) contains 12% Cr; it has been used in the HIBALL heavy-ion ICF design,⁵⁸ the EAGLE light-ion ICF design,⁵⁹ and the WITAMIR tandem-mirror design.⁶⁰
- Case P—Low-chromium ferritic steel, 2-1/4 Cr-1 Mo, is extensively used in steam generators, and two variations have been investigated for use with liquid lithium.⁶¹ The steel used in Case P has a small amount of Nb to stabilize the carbides against lithium corrosion.
- Case Q—The other version of 2-1/4 Cr-1 Mo steel, which stabilizes the carbides with a heat treatment, is used in the HYLIFE laser or heavy-ion ICF design.⁶²
- Case R—A 2-1/4 Cr-1 V steel is a modification to thermally stabilized 2-1/4 Cr-1 Mo in which the Mo is replaced by V to reduce the long-term activity. Although this steel has not yet been evaluated from other viewpoints (e.g., weldability or corrosion resistance), we include it to quantify its benefits from the viewpoint of activation.
- Case S—The low-activating alloy, 2-1/4 Cr-1 V, was also used with an ICF source and a 75-cm liquid-lithium wall protection.

Reactor Geometry. The neutronics calculations required to determine the neutron

flux were based on a one-dimensional reactor geometry, as illustrated in Fig. 8-23. The 14.1-MeV point source was enclosed by the 2-cm-thick spherical wall at 5-m radius. The neutron blanket/reflector consisted of 78 cm of lithium, 18 cm of graphite, 2 cm of boron carbide, and a 10-cm-thick outer wall. In each case, the neutron source was on for 30 years at a flux corresponding to 1 GW_e chamber power and 70% capacity factor, and then the wall material was allowed to decay.

Activity. The activity is shown as a function of time in Fig. 8-24(a), using a linear format. The linear-linear form is used to present an undistorted visual comparison because the apparent relative magnitudes in logarithmic form can be made large or

Fig. 8-23. Geometry used for comparison of materials.

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small by the choice of the lower cutoff decade. However, the linear scale has the disadvantage of not showing relative magnitudes at long decay times; therefore, we have expanded that portion of the linear scale in the insets in Fig. 8-24. A semilog plot was also generated (but is not shown) to allow easy extraction of numerical values

Table 8-11. Activity 50 years after shutdown.^a

Case	Description	Activity 50 y after shutdown (kCi)	Reduction since shutdown	Reduction from PCA
N	PCA	650	2 400	1
O	HT-9	30	52 000	22
P	Nb-stab. 2-1/4 Cr	44	36 000	15
Q	Thermally stab. 2-1/4 Cr	16	100 000	41
R	2-1/4 Cr-1 V	5.5	293 000	119
S	LMW protected 2-1/4 Cr-1 V	1.6	600 000	412

^aThe final two columns show the reduction of activity compared with shutdown and compared with PCA (50 y after shutdown).

at all times of interest and to allow estimation of the effective decay half-life as a function of time from the slope of the decay curve. The gap in the curves just after shutdown corresponds to the first year of decay.

At shutdown, all five unprotected walls have roughly equal activities, ~ 1.6 GCi, and the liquid-lithium-protected wall has a factor-of-16 less activity. Ten years after shutdown, the activity of each wall has dropped by a factor of ~ 30 , and the five unprotected walls still have roughly equal activities.

Fifty years after shutdown (80 years after start-up), a considerable spread has developed between the six cases, as shown in Table 8-11. The PCA (case N) is by far the worst with an activity of 650 kCi, a factor-of-2400 reduction from shutdown. The two versions of 2 1/4 Cr-1 Mo and the HT-9

Fig. 8-24. (a) Activity; (b) afterheat, as a function of time; (c) inhalation bhp, as a function of time; (d) ingestion bhp.

N = PCA
O = HT-9
P = Nb-stab. 2-1/4 Cr

Q = Th-stab. 2-1/4 Cr
R = 2-1/4 Cr-1 V
S = LMW-protected 2-1/4 Cr

(cases O, P, and Q) have activities of a few tens of kCi. The two 2-1/4 Cr-1 V walls (cases R and S) have activities of only a few kCi. Also, the liquid-lithium-protected 2-1/4 Cr-1 V wall (case S) activity is now over 400 times less than the PCA activity; it was only 16 times less at shutdown.

In all six cases, tritium has been eliminated from consideration because the tritium that diffuses into the steel during operation is much more significant than the tritium produced in the steel by neutron reactions.⁶³ Inspection of the FORIG output revealed the importance of accounting for the activation of daughters and for changes in parent nuclide populations. In one example, the ⁵⁵Mn inventory tripled during 30 years of activation, and its activation daughters would be underestimated if the initial population was used in a hand calculation.

Afterheat. The afterheat in the six walls is shown as a function of time in Fig. 8-24(b). If the temperature limit is 1300 K, the spherical chamber can be buried in earth when its afterheat is ~14 kW if it is first crushed, or ~61 kW if it is buried as a hollow sphere.⁶⁴ Assuming hollow sphere burial, the PCA wall cannot be buried until 8 years after shutdown, the other four unprotected walls until 4 years after shutdown, and the liquid-lithium-protected wall until ~1 year after shutdown. For crushed sphere burial, the PCA wall must be held for 17 years after shutdown, the other four protected walls for ~7 years, and the liquid-lithium-protected wall for 2 years.

Inhalation Hazard. The inhalation biological hazard potential (bhp) is shown for the six walls in Fig. 8-24(c). This index represents the amount of air required to dilute a fully vaporized wall such that the dose received by continuously breathing the mixture will not exceed 0.5 rem per year. The inhalation bhp is calculated for each nuclide by dividing the activity by the radioactivity concentration guide (RCG) taken from column 1 of Table II of 10CFR20 (Part 20, Title 10, Code of Federal Regulations). Since structure vaporization is only possible by applying large amounts of energy, inhalation bhp is only appropriate at shutdown or shortly thereafter. Three of the steels (HT-9, thermally stabilized 2-1/4 Cr, and 2-1/4 Cr-1 V) have roughly equal shutdown inhalation bhp of $\sim 3 \times 10^8 \text{ km}^3$. The PCA and Nb-stabilized 2-1/4 Cr walls are about 25%

and 50% higher, respectively, while the liquid-lithium-protected 2-1/4 Cr-1 V is a factor of 15 lower.

Ingestion Hazard. The ingestion bhp is the amount of water required to dilute a fully dissolved wall to the level that the dose received by using only this water will not exceed 0.5 rem per year. The ingestion bhp is calculated by using the RCG values taken from Column 2, Table II, 10CFR20. Dissolution of the wall into water is most likely after the material is buried (as waste). The ingestion bhp is shown for the six walls in Fig. 8-24(d). The data are plotted as a ratio of the radioactive ingestion bhp of the activated wall to the amount of water required to dilute an unactivated wall, from a chemical hazard viewpoint. Thus, when the ratio falls to 1.0, dissolution of the activated wall is similar to dissolving clean, unirradiated steel. After only 15 and 28 years, the liquid-lithium-protected and unprotected 2-1/4 Cr-1 V walls, respectively, reach the level where the radioactive hazard of shallow land burial is similar to the routinely accepted chemical hazard of burying an unirradiated wall. The other four walls reach roughly steady bhp values at ~35 years after shutdown. The HT-9 and thermally stabilized 2-1/4 Cr level is ~4 years, the Nb stabilized 2-1/4 Cr steady level is ~9 years, and the PCA wall reaches a level of ~20 years.

Disposal. The ratios of radioactive to chemical ingestion bhp at ~35 years after shutdown indicate that shallow burial might be feasible at that time. However, the Nuclear Regulatory Commission (NRC) has recently addressed the shallow burial issue in a new regulation, 10CFR61. This regulation lists the maximum allowable concentration of each nuclide in the solid waste (rather than in water or air). The actual concentrations are divided by the allowable concentrations, and the ratios are summed over all nuclides. The resulting sum (called the shallow burial index, or SBI) must be less than 1.0 before burial is allowed. This technique has been applied to our six steels at a time of 50 years after shutdown. The SBI values of PCA, HT-9, and both thermally and niobium-stabilized 2-1/4 Cr steels are above 100 times the proposed limits (due to the low allowable ⁹⁴Nb limit). The unprotected and lithium-protected 2-1/4 Cr-1 V walls have SBI levels of 0.14 and

0.06, respectively. Thus, only the 2-1/4 Cr-1 V walls could be buried 50 years after shutdown under the proposed rules. Since liquid-lithium protection would also reduce ^{94}Nb production in the other steels due to the softer and less intense neutron flux, liquid-metal-wall ICF reactors will have a significant advantage over magnetic fusion reactors from the viewpoint of structural waste disposal.

Conclusion. The PCA steel appears to be a poor choice from both afterheat and disposal viewpoints. The HT-9 and both thermally and Nb-stabilized 2-1/4 Cr steels are better from an afterheat viewpoint, but still unacceptable with respect to waste disposal. The unprotected 2-1/4 Cr-1 V steel has about the same time required as the latter two steels to reduce the afterheat to an acceptable level, and it is acceptable for shallow burial. The liquid-lithium-protected 2-1/4 Cr-1 V steel wall is by far the best of the six cases, with a short cooling time and legal disposal by shallow burial.

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Influence of Fusion Reactor Configuration on Neutron-Induced Radioactivity

Introduction. Although the neutron-induced radioactivity in a fusion reactor is a nonlinear problem requiring, for its solution, the use of neutron transport codes and neutron activation and decay codes, simple arguments can be made that give useful scaling laws for the total radioactivity in a fusion reactor.⁶⁵ Because these scaling laws rely heavily on assumptions of linearity and the smallness of second-order effects, we have compared them to the results of computer experiments designed to investigate their validity over the range of operating parameters typical of fusion reactors.

We varied the following factors to compare activation and decay

- Reactor power level.
- Neutron source distribution (point/homogeneous).
- First-wall radius.
- First-wall thickness.
- Reactor configuration (spherical/cylindrical).

We then investigated the

- Total radioactivity of the first wall compared with a structural wall located behind a lithium blanket.
- Effect of protecting the first wall with a thick liquid-lithium wall (as in the case of some ICF reactors).
- Effect on total radioactivity of neutron collisions in a typically dense ICF fusion plasma.
- Effect on total radioactivity of the neutron flux reflected from the blanket.

We address here only the totals for radioactivity, hazard potential, and thermal decay power. (Specific activity can be made arbitrarily small by increasing the total mass of waste and is therefore not a good figure of merit.)

Effect of First-Wall Thickness and Reactor Power. There are only two fundamental parameters with which activation in a typical fusion reactor varies linearly (if all else remains unchanged): total reactor thermal power, and the thickness of the first wall. Total radioactivity in a fusion reactor is not a direct function of first-wall radius, nor is it one of first-wall configuration. The derivation for this linear scaling law is described elsewhere.^{66,67}

Essentially, this law is an observation that, for the same fusion power, the number of neutrons impinging on the first wall is constant. Thus, the larger the area of the first wall, the lower the incident neutron flux, and total radioactivity is, to first order, independent of the area of the first wall. If the probability for any given nucleus in the first wall undergoing a neutron reaction over the lifetime of the irradiation is small, then secondary activations can be neglected, and the isotopic composition is then approximately unchanged. In this case, the total radioactivity in the first wall after some period of irradiation is directly proportional to the activation rate. If we may also assume that average neutron path lengths in typical fusion reactor first walls are small compared to a neutron mean free path, we may conclude that the total radioactivity in any fusion reactor is directly proportional to

- The number of fusion neutrons per second, which is in turn proportional to the reactor thermal power.

- The average neutron path length in the first wall, which is in turn proportional only to the first-wall thickness.

The degree to which these scaling laws hold valid depends on the extent to which secondary activations, parent nuclide depletion, and neutron absorption in the first wall can be neglected. We will investigate these scaling laws by comparing them with computational results for ranges of first-wall thicknesses and first-wall radii typical of fusion reactors.

The Effect on Total Radioactivity of First-Wall Configuration and Plasma Distribution. The argument for the linear scaling of total radioactivity with first-wall thickness and reactor power holds for reactors with distributed neutron sources, as well as for point sources; for cylindrical or toroidal wall shapes, as well as for spherical wall shapes; and for first walls protected by a thick layer of liquid lithium, as well as unprotected first walls.

However, because the characteristic flux spectrum in the first wall is different among these configurations, they will exhibit different total radioactivities. The direction of change in total activity due to first-wall configuration and plasma distribution is predictable in the following cases

- If the neutron source is distributed throughout the volume enclosed by the first wall, the averaged neutron path lengths in the wall will be longer because some of the neutrons will enter the wall obliquely. This leads to greater radioactivity than that expected from a point source of neutrons at the center of an enclosed volume.
- Greater induced radioactivity can be expected as the reaction chamber shape departs from that of a sphere toward that of a cylinder or a torus. This is again because neutrons will have greater averaged path lengths in the wall due to a greater averaged neutron angle of incidence upon the wall.
- For inertial-confinement fusion, the fuel must be compressed to very high densities. Since the collision probability increases with density, a significant portion of the neutrons will have lost sufficient energy through nuclear collisions in the fusion plasma to be below the thresh-

old for some neutron reactions in the first wall. The result is less induced radioactivity.

- Finally, for some inertial-confinement fusion reactor designs, a thick liquid-lithium wall is placed between the neutron source and the fusion chamber walls. In this case, the neutrons available to activate the first wall will be less in number and lower in average energy. The result again is less induced radioactivity.

Computational Experiments to Test the Scaling Laws. To investigate the validity of the scaling laws, we performed a number of computational experiments that calculated the magnitude of each first- and second-order effect listed above. Computational evidence is necessary for assurance that the nonlinear nature of neutron-reaction cross sections, combined with the after-scattering changes in population, direction, and energy of the neutrons, does not produce significant nonlinear effects.

In each experiment, one parameter of the base case was changed while all others were held equal. A neutron-transport code was then run to determine the neutron flux spectrum, from which flux-averaged cross sections for neutron activation were calculated. Next, a neutron activation and decay code was run to determine the activity during 30 years of reactor operation and for the first 100 years following decommissioning. (The computational procedure is described in detail elsewhere.⁶⁸) The parameters chosen for each experiment are shown in Table 8-12.

Table 8-12. Parameters for the computational experiments.

Case	Description
A	The base case. Point source of neutrons corresponding to 1 GW _e at the center of a sphere with 5-m radius and 2-cm first-wall thickness, followed by a blanket of 80-cm Li, 18-cm graphite, 2-cm ³ He, and 10-cm steel outer structure.
B	One-half fusion power (0.5 GW _e).
C	Double fusion power (2 GW _e).
D	Thinner wall (1 cm).
E	Thicker wall (3 cm).
F	First wall at 3 m radius.
G	First wall at 7 m radius.
H	Cylinder with line source and reflecting ends.
I	Homogeneously distributed neutron source (0 to 5 m).
J	Neutron-collisional ($\rho r = 3 \text{ g/cm}^2$) ICF source.
K	First wall protected by a lithium wall at 50% liquid density from 0.5 to 2.0 m.
L	Neutron reflections from blanket neglected. (Blanket replaced by a perfect absorber of neutrons.)
M	Activity in the outer steel structure of the base case.

The effects of neutrons scattered by the blanket were included by placing a typical blanket behind the first wall in each case. The neutron flux corresponded to 2700 MW fusion power at a 70% capacity factor. Each first wall was irradiated for 30 years to maximize the nonlinear effects that could invalidate the scaling laws.

In each case, the first-wall material (2-1/4 Cr-1 Mo ferritic steel) remained unchanged. We also compared and analyzed other steels that will yield different characteristic curve shapes.⁶⁸ (See also the previous article.) However, the neutron flux spectrum for all steels will be approximately the same, since it is primarily determined by collisions with iron, the major component of steel. Therefore, although the magnitudes and compositions of the radioactivity in other steels will be quite different, they will nevertheless obey the same scaling laws to

roughly the same degree that they are obeyed by 2-1/4 Cr-1 Mo steel.

The results of the computer experiments conform closely with the scaling laws based on first-order theory and are presented in Fig. 8-25 in linear-linear form. (The linear-linear form is necessary in order to present a good intuitive visual comparison of relative magnitudes. In log-log or semilog form, relative magnitudes can appear to be large or small depending on the choice of the cutoff decades.) To show radioactivities at long decay times, we have expanded that portion of the linear scale in the inset of Fig. 8-25. For comparison, the magnitudes of each curve of Fig. 8-25 relative to the magnitude of the base case are shown in Fig. 8-26 at four sample times.

Departures from the scaling laws become more pronounced at long decay times. This is because most of the original activity,

Fig. 8-25. Activity curves for reactors with parameters described in Table 8-12.

which was produced by direct activation of stable nuclides, has decayed away at long decay times, leaving secondary activations in greater relative importance. Differences in curve shapes are most pronounced for cases K and M because moderation of the neutron flux spectrum by an interceding thick layer of lithium produced a different set of averaged cross sections.

Conclusion. We have found that previously described scaling laws for total neutron-induced radioactivity in typical fusion reactors are in good agreement with computed results for a range of typical reactor power levels, operation times, and wall thicknesses. Except at long decay times, the total radioactivity in the first wall

- Varies in almost direct proportion (to within about 5%) to the first-wall thickness, and the rate of increase decreases slightly as thickness increases.
- Varies in almost direct proportion (to within about 5%) to the reactor thermal power.

- Is almost independent (to within about 5%) of the first-wall radius.
- Increases (on the order of 35%) if the first-wall geometry is changed from a sphere to a long cylinder.
- Increases (on the order of 35%) for a distributed neutron source instead of a point source.
- Decreases (on the order of 10%) if the effect of the neutrons reflected from collisions in the blanket are neglected.
- Decreases (on the order of 30%) with very high-density fusion plasmas typical of inertial-confinement fusion.
- Is greater (by an order of magnitude) than the radioactivity in typical structural material behind the blanket. (This justifies approximating the total induced radioactivity with that of the first wall.)
- Is greater (by an order of magnitude) than the radioactivity in a similar first wall protected by an effective thickness of 75 cm of liquid lithium (as in the case of liquid-metal-wall ICF reactors).

Fig. 8-26. Total radioactivity of each case in Table 8-12 compared with that of the base case.

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In summary, all fusion reactors, regardless of geometrical shape or first-wall radius, having unprotected first walls of the same thickness and made of similar material, and operated for the same amount of time at the same power levels, will have roughly equal total radioactivities (but distinctly less in the case of ICF due to chamber sphericity, neutron collisions in the dense plasma, point-source effects, and, in some cases, liquid-metal-wall protection). Their total radioactivity scales only with the product of their thermal power and their first-wall thickness.

Because other quantities of interest—biological hazard potential, thermal decay power, and gamma activity—are the scalar product of a set of constants and the corresponding radionuclide activities, it follows that the totals of these other quantities of interest also scale in direct proportion only to thermal power and to first-wall thickness to the extent to which secondary activations, parent nuclide depletion, and absorption of neutrons by the first wall can be neglected.

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Liquid-Metal Systems

Temperature and Velocity Limits for Liquid-Lithium Solutions

Fusion power plant designs are very sensitive to limits imposed on liquid-metal temperature and velocity. Temperature limits affect plant efficiency; velocity limits determine pipe diameters and, consequently, influence plant costs. System designers, therefore, frequently request quantitative corrosion rates—data as a function of temperature and velocity—for structural alloys immersed in either lithium or lead-lithium liquids. Such data are available for flowing sodium systems, but the liquid-metal community is reluctant to give even qualitative data on lithium and lead-lithium. In this article, we discuss the reasons for this situation, and we present some very qualitative data based on studies done for LLNL at the Energy Technology Engineering Center (ETEC).

Temperature Limits. Screening tests for liquid-metal corrosion usually start with isothermal capsule or crucible tests using static liquid metal surrounding specimen tabs. Initial weight loss is large as highly soluble elements leave the structural metal, diffuse through a static subboundary layer, and begin to mix into the boundary layer and bulk liquid. Carbides or impurity precipitates in the grain boundary may be preferentially attacked. After hundreds of hours, the structural alloy surface reaches a relatively steady-state composition, and the rate of dissolution in the specimen is determined by the activity gradient through the subboundary layer of the liquid in contact with the specimen. As the elements dissolved from the specimen build up in concentration within the bulk liquid metal, the activity gradient across the subboundary layer becomes shallower, and the corrosion rate slows, reaching zero as the bulk liquid saturates. A curve of weight loss vs time for static tests resembles the curve shown in Fig. 8-27.

Much of the static-test liquid-metal corrosion data are "slope A" results. ("Slope A" lines are drawn between the origin and points on the curve; "slope B" lines are drawn between two points on the curve, eliminating the early behavior before the liquid becomes saturated in the dissolved elements of the steel.) The weight loss between time zero and some arbitrary number of thousands of hours is measured. The weight loss is converted to surface regression and this value—measured in typical units of corrosion rate: mil/y—is divided by the exposure time. Corrosion rates obtained from slope A data are quite dependent on the test time chosen, the volume of liquid metal, and the liquid-solid surface area. Perhaps the most useful data that can be obtained from static capsule tests are slope B measurements. The difference in weight loss between 1000 and 2000 h of exposure is measured and expressed in terms of mil/y. In general, a slope B value of less than 0.2 mil/y at a particular temperature is seen as an encouraging indication that the structural alloy may be used at that temperature. Values above 0.2 mil/y are a warning sign that the alloy may not perform adequately at the test temperature. The most reliable static test results are those

that follow the behavior in Fig. 8-28, where the slope B value plotted as a function of temperature is a low value over a wide temperature range.

Thermal convection loops and pumped loops give data that are far more useful than static capsule tests, but complications preclude the assurance that quantitative corrosion data are valid for a particular design. Deposition of dissolved solute in regions of supersaturation and thin boundary layers is the dominant characteristic observable in the flowing loops. Flow blockage by deposited crystals becomes much more important than the excessive removal of a structural-wall material.

Saturation effects that cause static-capsule corrosion rates to approach zero at long times also contribute to a downstream effect in loop tests. If the hottest part of a loop is a long isothermal section, the bulk liquid moves toward saturation, and the activity across the subboundary layer becomes a function of distance from the beginning of the section. Additionally, the thickness of the subboundary layer is a function of liquid velocity and flow patterns, and the activity gradient is the ratio of the activity across the layer to the layer thickness. Liquid flow characteristics influence corrosion rate even in a completely isothermal loop. When ΔT conditions are imposed on the loop, temperature gradients along the flow interact with the other test parameters to complicate the situation. In general, deposition of dissolved solute is favored in regions that are cold, supersaturated, and have high velocity. As the deposited solute grows, the subboundary layer thins at the tip of the growth, and the deposition rate increases until plugging occurs.

The behavior of each loop is different, and even the same loop can change in crystal deposition behavior as certain elements are leached away. Deposited crystals begin to change in stoichiometry as manganese, for example, is no longer available from hot surfaces. Partially closing a valve can create a high-velocity flow and cause crystallization from the liquid to occur locally at the valve seat. Introduction of an impurity into the liquid metal can cause major changes in behavior; e.g., nitrogen at a level of hundreds of ppm causes lithium to preferentially attack grain boundaries of austenitic

stainless steel. Thus, corrosion-rate data obtained in a given loop using lithium with 50 ppm by weight of nitrogen will not predict the corrosion rate data in the same loop with lithium that contains 500 ppm nitrogen.

All these phenomena are viewed by the metallurgist as valid reasons for not furnishing the system designer with quantitative corrosion rate data as a function of temperature for structural alloys in lithium or lead-lithium. For lithium, we can say qualitatively that properly heat-treated ferritic steels and austenitic stainless steels should not give problems at 400°C. However, at

Fig. 8-27. Steel weight loss vs time for a static test.

Fig. 8-28. Slope B corrosion rates depend on the temperature.

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500°C, ferritic steels are near the upper boundary of the temperature regime where these alloys may be used; at 550°C, ferritic steels are well into the temperature regime where these alloys may cause excessive mass transfer. By comparison, the austenitic stainless steels are well into the temperature regime of excessive mass transfer at 500°C.

We are now investigating the temperature limits for lead-lithium liquids. In addition to corrosion data, lead embrittlement is a serious concern that will affect heat-treatment practice. This embrittlement has been noticed at 1 at.% Pb in Li. We do not know as yet whether the effect extends down into the ppm range or up to the high-lead eutectic ($\text{Pb}_{43}\text{Li}_{17}$). The embrittlement phenomena may create a low-temperature limit, creating a temperature window in which alloys may be safely used. System design considerations require a temperature window of at least 50°C.

Velocity Limits. Velocity effects fall into two categories: polishing and etching (Fig. 8-29). Systems operating in the etching regime have rough surfaces where grains have etched out, while systems operating in the polishing regime usually have smooth surfaces. Liquid-metal systems should be operated in the polishing regime so that grains of structural alloy cannot etch out and fall into the moving liquid metal.

If the steel has not been properly heat-treated, the lithium can attack the carbides in the grain boundaries, leading to an etched appearance even though the velocity

is low. Thus, an etched appearance could indicate either a poor heat treatment, or corrosion in the etching regime, or both.

The critical subboundary layer thickness at which etching stops and polishing begins is a function of structural alloy, liquid-metal, temperature, and, particularly, liquid-metal solute concentrations. The critical subboundary layer thickness is usually expressed in terms of the critical liquid-metal velocity through a pipe. In LMFBF sodium systems operating at 500°C with an oxygen content below 2 wppm, the critical velocity is just under 7 m/s for austenitic stainless steels. Pipe diameters are specified with this limitation in mind. However, for lithium systems, no experiments have successfully established the critical velocity, and there is no theoretical reason to expect the lithium critical velocity to approximate that of sodium. Preliminary experiments have indicated that both the nitrogen content of the lithium and the stress state in the structural alloy will affect the critical velocity for lithium flowing through piping. Wilkinson⁶⁹ rotated a 3-cm-diam disk of 2-1/4 Cr-1 Mo steel at speeds up to 270 rpm in 500°C lead-lithium (1 at.% Pb-99 at.% Li) liquid for 400 h and found that the maximum velocity of ~0.42 m/s did not exceed the critical velocity. System designers are presently using the 7 m/s sodium limit on lithium designs, even though the actual limit may be more than an order of magnitude lower or several orders of magnitude higher than this limit.

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Fig. 8-29. Corrosion rate depends on liquid-metal velocity.

Liquid-Jet Impact Experiments

The HYLIFE reactor chamber uses a cylindrical steel shell as the first structural wall to absorb the radial momentum imparted to the liquid-metal jet array. The liquid must impact the wall of HYLIFE in a two-phase flow to keep transient stresses within acceptable limits. Monsler proposed an alternative first-structural-wall (FSW) concept consisting of multiple layers of steel bars wound in a helical configuration. The steel

bars in each layer are separated by a gap that allows a fraction of the impinging liquid to flow past that layer. The liquid momentum is absorbed over a longer time period than is possible with the shell configuration.⁷⁰ In this way, the peak stresses in the bars can be kept small, even if the liquid impacts at full density. The concept is illustrated in Fig. 8-30.

The University of California at Davis has performed a study to evaluate the transient and steady-state drag of a liquid on a single bar—and some selected arrays of bars—as

well as to determine the momentum removed from impacting slugs of liquids.^{1,72} Data from relevant published papers were cast in terms of dimensionless time, pressure, and impulse to obtain dimensionless pressure and impulse time-histories. The scatter band in the dimensionless time-histories was sufficiently small (considering the wide variety of data) to suggest near universal dimensionless curves for pressure and impulse.

Concurrently, experiments were performed at Davis to study the impact of

Fig. 8-30. Tapered HY11FE reaction chamber with alternate-bar FSW.

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liquid jets on single bars and several arrays of bars. The drag coefficient of a jet impacting on a single bar was found to be 0.54 for a Reynolds number of 10^4 . This is about half the classical drag coefficient of an infinite fluid over a cylindrical bar. The reduction in the drag coefficient is due to the three-dimensional effects and the splitting of the round jet once it strikes the bar, which causes a region of near-ambient pressure on the downstream side of the bar.

The steady-state drag of a jet on both a coplanar array and a staggered array of bars was measured. The drag forces measured are about an order of magnitude less than those predicted for infinite flows (of the same velocity as the jets) over the geometrically identical arrays. We used the single-bar drag coefficient for jet impact to develop a method of estimating the drag force of jets impacting on an array of bars.

This method can be coupled with the near-universal dimensionless pressure and impulse time-histories to estimate the momentum removal of a jet impacting on an array of bars. As an example, a jet slug with a thickness-to-diameter ratio of 1:5 impacting on a staggered array of four rows of bars with bar diameters 8.3% of the jet diameter will transfer about 13% of the momentum to the array.

Summary. The problem of momentum removal from slugs of liquids using arrays of cylindrical bars has been investigated, and small-scale scoping experiments have been performed. These experimental results permit a first estimate to be made of the magnitude of the momentum removal possible with typical staggered arrays of bars four rows deep or with in-line arrays three rows deep. These estimates, which must be considered tentative until more extensive experiments are performed, are much better than estimates based on the pressure drop of infinite flows over tube banks. Based on these tentative results, we can state that a first structural wall consisting of an array of steel bars will absorb the liquid momentum over a longer period of time than will a cylindrical shell.

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Passage of Lithium through Stainless Steel Mesh

Some inertial fusion power plant designs require that lithium or lead-lithium

Fig. 8-31. BLiP loop flow paths.

solutions flow through a porous material such as steel mesh, woven cloth, or a porous-metal wall. These designs often rely on reestablishing 100% liquid coverage of the porous material surface almost immediately after a fusion energy pulse has evaporated liquid from the surface. In lithium-flow studies done for LLNL, ETEC developed a bimetallic lithium pumped loop (BLiP loop) that can operate isothermally at 500°C. This loop was designed so that flow could be directed axially through a cylindrical vessel packed with a spirally wound screen. The rolled screen was wound around a vertical axis so that the main flow was between tightly packed sheets of screen and the flow perpendicular to the general flow direction (i.e., radial flow) was through the fairly wide mesh of the screen.

The screen was so tightly wound that radial flow through the mesh was easier than axial or circumferential flow between the sheets of screen. The top of the screen was covered by a plate with many holes which evenly distributed the liquid from a pool that formed on the plate. An electromagnetic (EM) pump supplied liquid to the top of the pool, and argon pressurization was used to force the liquid down into the mesh.

The BLiP loop at ETEC is shown in simplified form in Fig. 8-31. Lithium flowed for ~4000 h from mid-September 1981 to mid-April 1982. We found that the flow through the screen could not be initiated without argon pressurization. Argon pressure in the 10-psi range pushed the lithium through the screen and the flowmeter showed several gpm. The start-up argon overpressure was then reduced to the nominal system pressure, and the flow continued without an increase in the height of the lithium pool above the screen. This indicated that, once the flow was started, the gravity head was sufficient to maintain flow.

At the end of the experiment, the test vessel was cut open, and the screen was examined to see if the lithium had flowed uniformly in the screen section and whether the entire screen surface was covered by lithium. Since the total length of wire in the screen was on the order of 610 km, the huge surface area of lithium was of concern as a fire hazard; hence, the examination was performed in an inert gas, rather than in air.

The examination showed that the vast majority of screen wire was completely bare

Fig. 8-32. Cutaway view showing lithium stalactites and the lithium piston formed at bottom of rolled screen.

Fig. 8-33. Temperature of the test tank at midnight during the draining process.

of lithium. Preferred lithium paths had developed, and, like water tracing down a windshield, the entire lithium flow was down the wetted paths. The result was five lithium stalactites with circular cross sections coming down the rolled screen and merging into a 3-cm-thick "piston" of lithium at the bottom of the screen pack. The screen away from the stalactites and the piston were never wet by the lithium that flowed for thousands of hours at ~500°C. The bottom of the lithium piston was very smooth, with about 1 mm of coated screen wires sticking through it. Evidently, lithium droplets had rained down from the bottom surface of the suspended lithium piston into the lower pool. The stalactites and piston are shown schematically in Fig. 8-32.

To make sure that the piston was not an artifact of sudden cooling of lithium during the final drain process, we monitored the temperature of the test vessel at screen midheight. Draining took 0.5 h, while the vessel did not cool to the lithium freezing point (~180°C) until 15 h after draining (Fig. 8-33). Since the liquid-lithium piston at the screen bottom was in good thermal

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contact with the vessel wall, the existence of the piston was not an artifact caused by solidification during draining.

The lack of complete wetting of lithium on stainless steel wire mesh during 4000 h of time at 500°C may not be as serious a problem to fusion reactor designers as it would appear at first thought. The lithium did wet the screen within the stalactites and in the piston. Unwet areas initially rejected lithium and then did not have their surfaces exposed to the liquid during the test. Some prewetting techniques may be required in these types of applications.

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Electromagnetic Pumping of Liquid Lithium

Free liquid metal is proposed for use within the fusion reaction chamber of several ICF conceptual reactors, including HYLIFE, JADE, Pulse*Star, EAGLE, and the modified wetted wall. In each case, the reaction chamber is maintained at a vacuum to allow beam transport to the target. Because of the vacuum within the liquid-metal pumping circuit, and the cyclic pressure pulses added to the liquid by neutron deposition, ICF reactor pumps are prone to cavitation, which could damage the driving surfaces in mechanical pumps. Electromagnetic (EM) pumps, however, use the Lorentz Force ($\mathbf{F} = \mathbf{I} \ell \times \mathbf{B}$) and require no moving surfaces within the fluid. Consequently, EM pumps can more easily cope with a low suction head and cyclic pressure pulses.

In Ref. 73, we describe 10 distinct EM pump configurations and the design procedure for one promising geometry, the helical-rotor EM pump (HREMP). Our design procedure was shown to be valid when applied to an experimental sodium HREMP. The design procedure was then used to specify a design point for a lithium flow rate of 8.08 m³/s at a head of 82.5 kPa (the requirements for HYLIFE). This design point includes a complete set of 70 variables for a HREMP with either a normally conducting or a superconducting winding.

To design a HREMP, three dimensionless variables were varied parametrically until optimum performance was achieved. The

first variable was the ratio of pole pitch to the magnetic gap (see Fig. 8-34); the second dimensionless variable was the ratio of rotor length to pole pitch; and the third variable was the ratio of the field velocity to the circumferential component of the fluid velocity.

While a more complete development of our design equations is contained elsewhere,⁷³ the two most important equations are those for the required magnetic-flux density and for the pump efficiency. The peak magnetic-flux density may be expressed as

$$B = \left[\frac{PQ/N}{v_f(v_b - v_f) L^2 R} \right]^{1/2}, \quad (15)$$

where

$Z = (R^2 + X^2)^{1/2}$ = the impedance of the liquid-metal eddy-current path

R = resistance

X = inductance

PQ/N = the pressure times flow rate per pole

v_f = the fluid circumferential velocity

v_b = the field velocity

l = the rotor length.

The pump efficiency is

$$\eta = \frac{PQ}{1.1} + \frac{PQ(v_b - v_f)}{v_f} + 2N \left(\frac{V_w}{Z_w} \right)^2 R_w + 0.168 M_{\text{trp}} + (N I_w)^2 R_w, \quad (16)$$

where

PQ = the gross hydraulic power added to the liquid metal (10% friction loss in the pump)

$PQ(v_b - v_f)/v_f$ = the liquid-metal slip loss due to the difference between the circumferential fluid and field velocities

$2N(V_w/Z_w)^2 R_w$ = the total $I^2 R$ loss in the duct walls

$0.168 M_{\text{trp}}$ = the combined eddy-current and hysteresis loss in the flux-return path

$(N I_w)^2 R_w$ = the $I^2 R$ loss in the rotor winding (the factor of $N = 2$ accounts for the two poles on the rotor).

Fig. 8-34. (a) Longitudinal cross section of a HREMP. (b) Horizontal cross section of a HREMP.

The slip loss mentioned above is also the liquid-metal I^2R loss, numerically equal to $N(V/Z)^2$. (Calculation of this loss from the two independent formulae was used as a check on the design procedure.) The total I^2R loss in the duct walls is also a slip loss, but with $v_t = 0$ and $Q = 0$, the slip formula would have to be calculated in the limit, rather than directly.

For a superconducting winding, the load on the cryogenic cooling system must be kept low because the cooling efficiencies are in the 0.1 to 0.2% range. For example, consider a 60%-efficient pump operating at 82 kPa and 8.08 m³/s. From Fig. 8-35, the heat load on the cryogenic system should be limited to $\sim 0.01\%$ of the pump losses, i.e., to ~ 40 W. This goal can be met by using insulation consisting of a 6.35-mm Fiberfrax insulation layer, a 12.7-mm gap filled with circulating argon, a 6.35-mm Fiberfrax insulation layer on a thin Dewar outer wall, and five polished stainless-steel radiation shields inside the Dewar; the resulting heat load is then ~ 21 W.

Figures 8-36(a) and 8-36(b) show the pressure and efficiency of the normal and superconducting designs as a function of flow rate. For the HREMP with a superconducting winding, the zero-flow pressure is higher than the HYLIFE gravity head, and start-up will not require additional winding current or reduced rotor velocity. In contrast, for the HREMP with a normally conducting winding, the zero-flow pressure is

Fig. 8-35. The effect of cryogenic refrigeration efficiency on overall HREMP efficiency.

insufficient to overcome the HYLIFE gravity head (which is 91% of the head at the design flow rate). Thus, for start-up, the pump with a normal winding must be temporarily operated at a higher winding current or lower rotor velocity to develop the pressure required for static lift at zero flow rate.

If the magnetic flux density is increased to 143% of the design value to allow start-up of the normal-winding pump, the flux density increases to 3 T in the rotor core, well above the 2-T saturation limit. Alternatively, if the flux density and rotor dimensions are left at their design-point values, the zero-flow pressure can be increased by decreasing the rotor speed and, hence, the field velocity. At a field velocity of 43% of the design value, the zero-flow pressure equals the HYLIFE gravity head. Start-up can be accomplished by using a three-speed

Fig. 8-36. Pressure as a function of flow rate for a HREMP with (a) a normally conducting winding, and (b) a superconducting winding.

Fig. 8-37. A HREMP with a normally conducting winding can be started by varying the rotor speed.

gear box between the drive motor and the rotor, as shown in Fig. 8-37.

The HREMP designs for both a normally conducting and a superconducting winding were selected from a large number of potential designs in which the three dimensionless variables were varied parametrically. The information contained in Ref. 73 shows the sensitivity of the design figure of merit as the three dimensionless variables are changed. The ratio of efficiency to mass-diameter product was selected as a figure of merit because our overall objectives are to maximize the pump efficiency and to reduce the power-plant capital cost. Minimizing the pump mass reduces pump

cost, and minimizing the pump diameter reduces building cost. Because both mass and diameter appear in the denominator of the figure of merit, the optimal values for the three dimensionless variables may be derived from the points where the figure of merit is maximized. In terms of cost, the pump height is relatively unimportant because it is less than the height of both the reactor vessel and the steam generators. In fact, there is sufficient height in the building to allow for mounting the rotor drive motor either above or below the pump.

Our design values for the three dimensionless variables (and for the efficiency, mass, and diameter) are shown in Table 8-13 for each pump type. At first glance, the pump with a superconducting winding seems far superior because its figure of merit is 11.5 times larger than for the pump with a normal winding. However, the superconducting pump has additional costs because of the cryogenic cooling system and the decreased plant availability resulting from reduced reliability; hence, the pump with a normal winding may prove to be the best choice. For example, if efficiency alone were the figure of merit, the maximum pump efficiency would increase from 50 to 54% for the normal winding, and from 60 to 69% for the superconducting winding. Increasing the HYLIFE pump efficiency from 50 to 69% increases the plant output by 0.36 MW (per pump). At 5¢/kW-h and a 70% capacity factor, this energy has an annual value of \$110 000. If

reliability is unchanged, and a 10% interest rate is used, we could invest up to \$1.1 million in each wider or heavier pump (and larger building) before the efficiency savings are lost. However, if the efficiency is increased by using a more complex pump (such as one with cryogenic windings), the pump is likely to be less reliable, and the plant will be operational less of the time. Decreased plant availability of only 0.3% would cancel the benefit of increasing the efficiency of 11 HYLIFE pumps from 50 to 69% if that efficiency improvement could be obtained with no capital cost.

For a given winding type, trade-off of efficiency and size can be made when the pump cost per kilogram of pump mass and the building cost per meter of pump diameter are known. The figure of merit to be optimized will be a function of these cost factors, in addition to the interest rate and electric energy value. When the normal and superconducting windings are compared, additional consideration must be given to the effect of each pump type on the plant reliability.

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Embrittlement of 2-1/4 Cr-1 Mo Steel by a 1 at.% Pb-99 at.% Li Liquid

Introduction. The steel alloy 2-1/4 Cr-1 Mo is a major potential candidate for the first structural wall in a laser fusion reactor because of its low strategic alloy content, low swelling, and resistance to liquid metal corrosion. Data generated by the Colorado School of Mines (CSM) in previous LLNL-funded programs illustrated that 2-1/4 Cr-1 Mo steel could provide good resistance to lithium corrosion, provided the carbides were stabilized either by the addition of niobium or by careful heat treatment.^{74,75}

In inertial fusion reactors such as HYLIFE, a heavy metal from the fuel capsule, such as lead, will be mixed with the lithium. Data generated by CSM, also under the auspices of LLNL funding, indicated that the addition of lead to the liquid lithium

Variable	Normal winding	Superconducting winding
Pole pitch/magnetic gap	6.5	3.0
Rotor length/pole pitch	1.77	0.55
Field/fluid velocity	1.13	1.40
Efficiency (%)	50.2	59.6
Mass ^a (Mg)	105	12.5
Diameter (m)	2.54	2.21
Figure of merit	0.188	2.19

^aNeither design includes the mass of the drive motor or the heat-removal systems.

did not alter in a major way the lithium corrosion of 2-1/4 Cr-1 Mo steel at temperatures from 400 to 500°C.^{76,77}

A major concern under study at CSM, again under LLNL auspices, is the potential embrittlement of 2 1/4 Cr-1 Mo steel by a lead-lithium liquid. Several earlier studies have indicated that lithium can be an embrittling agent for metals under load. The tensile properties (especially elongation) of Armco iron have been shown⁷⁸ to be adversely affected by exposure to lithium, particularly when the iron has been slowly cooled, then strained at a slow strain rate in lithium at 200 to 600°C. Other studies⁷⁹⁻⁸² have shown that lithium embrittles not only iron, but silver, copper, nickel, palladium, stainless steels, and high-strength steels.

Previous research by Rostaker et al.⁸² on the lithium embrittlement of AISI 4130 steel found that the steel was catastrophically embrittled by lithium only when the steel was heat treated to a tensile strength of 150 ksi (1 GPa) or above. The ductility of the steel, measured after tensile testing in molten lithium at 205°C, was less than expected even when the steel had been heat treated to low strength levels. These workers further found that the 210°C fracture strength of AISI 4130 steel wetted by lithium increased inversely with the square root of the mean free path between carbides. Grain size has also been found to be an important microstructural parameter in lithium embrittlement. Rostaker et al.⁸² showed that the fracture strength of both copper and mild steel wetted by lithium at 205°C increased inversely with the square root of grain diameter.

Lead also has been identified as an element that embrittles other metals, including zinc, germanium, copper, nickel, iron, and low-alloy steels.^{81,83} This phenomenon has been found to be sensitive to second-element content. For example, 0.1% zinc

Table 9-13. Design values for helical-rotor EM pumps.

Fig. 8-33. Ultimate tensile strength of 2-1/4 Cr-1 Mo steel, determined in both argon and lead-lithium liquid.

was found to extend the temperature of embrittlement of AISI 4140 steel by lead to above 425°C (Ref. 83).

Results and Discussion. The only known study of the potential embrittlement of 2-1/4 Cr-1 Mo steel by lithium, lead, or by a lead-lithium liquid is the LLNL-funded preliminary study in progress at CSM. In this program, two microstructural conditions of 2-1/4 Cr-1 Mo steel are being strained in tension at fixed deflection rate (2 $\mu\text{m/s}$) while submerged in either pure lithium, pure lead, or the 1 at.% Pb-99 at.% Li (17.6 wt% Pb, 82.4 wt% Li) liquid at eight different temperatures. The two microstructures being studied are obtained by the following heat treatments: (1) austenitized at 1300°C, oil-quenched; and (2) austenitized at 1300°C, oil-quenched, tempered at 750°C for 10 h. These microstructural conditions represent the extremes in potential embrittlement susceptibility. The

as-quenched condition is the highest strength and, therefore, most susceptible to liquid-metal embrittlement. The high austenitizing temperature produces a microstructure simulating a weldment that has not been postweld heat treated. In contrast, the quenched and tempered microstructure, representative of an extensively tempered weldment, is sufficiently ductile to likely avoid liquid metal embrittlement, although such a heat treatment significantly lowers the allowable stress to which a structure could be designed.

The results of testing these microstructures at fixed deflection rate while submerged in a 1 at.% Pb-99 at.% Li liquid are shown in Figs. 8-38 and 8-39. Figure 8-38, which shows the ultimate tensile strength vs temperature for both microstructures, tested in both argon atmosphere and the lead-lithium liquid, vividly illustrates the significance of the embrittlement problem. While the heavily tempered specimens display little or no difference in fracture strength whether tested in argon or in the lead-lithium liquid, the untempered specimens display significantly reduced fracture strengths over the entire range of temperatures when tested in the lead-lithium liquid. The tensile strength data also suggest another difficulty with utilizing an extensive temper to solve a lead-lithium embrittlement problem. The heavily tempered specimens are significantly softened by the heat treatment; base plate given a normal (900 to 950°C) austenitizing treatment would conceivably be excessively weakened by such an extensive temper. Measurements of ductility, expressed as percent elongation in a nominal 1-in. gage length, are shown in Fig. 8-39 for the same specimens described in Fig. 8-38. Note that the as-quenched steel exhibited limited ductility even when tested in argon. In the argon tests, the percent elongation of the as-quenched steel was low, unless the test temperature was greater than 400°C, although the reduction in area was found to be 10 to 30%, depending on test temperature. The as-quenched steel exhibited nil ductility at all temperatures when tested in the 1 at.% Pb-99 at.% Li liquid.

Ductility measurements of the tempered steel tested in argon were significantly greater than comparable values for the untempered steel, as expected. (The apparent

minimum in percent reduction in area at about 400°C is not yet understood; further study of this phenomenon is in progress.) Ductility measurements of the tempered steel tested in the 1 at.% Pb-99 at.% Li liquid were also significantly greater than comparable values for the untempered steel, again attesting to the effectiveness of a tempering heat treatment to reduce liquid metal embrittlement susceptibility. Ductility measurements at 250°C provide a cautionary note; an apparent loss of ductility at this temperature, even after an extensive temper, implies that tempering may not completely alleviate embrittlement by a lead-lithium liquid. These results are consistent with results by Kustaker et al.⁸² showing ductility losses in heavily tempered 4130 steel when tested in lithium at 205°C. Further testing in this temperature regime is in progress to clarify the extent of embrittlement, if any, in heavily tempered 2-1/4 Cr-1 Mo.

Scanning electron fractographs of both as-quenched and quenched and tempered 2-1/4 Cr-1 Mo steel samples were made after testing in both argon and the 1 at.% Pb-99 at.% Li liquid at 348°C. In the argon tests, the fracture surfaces for both the as-quenched and quenched and tempered specimens were characterized by fine microvoid coalescence, indicating that the steel failed by ductile rupture at this temperature. By contrast, the as-quenched 2-1/4 Cr-1 Mo steel, when tested in a 1 at.% Pb-99 at.% Li liquid at 348°C, failed by intergranular fracture—a phenomenon commonly observed in liquid-metal embrittlement. Tempering the steel for 10 h at 750°C restored the ductile-rupture mode of failure to specimens tested in the lead-lithium liquid.

Conclusions. Specimens of 2-1/4 Cr-1 Mo steel, heat-treated to simulate the near heat affected zone of a weldment, were severely embrittled by a 1 at.% Pb-99 at.% Li liquid at temperatures from 250 to 525°C. Embrittled specimens failed by intergranular fracture. An extensive temper (10 h at 750°C) significantly reduced the embrittlement susceptibility of the steel. Tempered specimens generally exhibited good ductility and failed by ductile rupture. However, tempering significantly reduced the strength of the steel. If applied to base plate metal, given a conventional austenitizing heat treatment, such tempering could reduce the

strength to an undesirable level. Also, measurable loss of ductility was observed in tempered specimens tested in the lead-lithium liquid at 250°C.

Fig. 8-39. The ductility of 2-1/4 Cr-1 Mo steel, determined in both argon and lead-lithium.

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Power-Plant Considerations

Evaluation of Direct-Illumination-Driven ICF Reactors

Inertial-fusion reactors for power generation require high-energy drivers to compress and ignite the fusion fuel pellet. There are two techniques for irradiating the fuel pellet: radiation-driven and direct-illumination-driven. The radiation-driven targets have a high-Z radiation case that converts the driver beam energy to x rays. This radiation

case provides a hohlraum-like region that gives relatively uniform irradiation of the pellet with as few as two driver beams. On the other hand, the direct-illumination-driven pellets absorb the driver energy in the expanding plasma from the fuel capsule. To absorb the driver energy uniformly, a direct-drive pellet requires a large number of driver beams evenly distributed over its spherical shell. The final optics that turn the driver beam into the reaction chamber can be either "defocused" high $f/\text{No.}$ or low $f/\text{No.}$ optics. In either case, the velocity asymmetries must be kept to less than about 4%.

Recent conceptual designs for inertial-fusion power reactors are suitable for radiation-driven targets. These reactor designs usually feature high $f/\text{No.}$ driver beams entering the reactor cavity from essentially two sides. In a recent study,⁸⁴ we examined the possibilities of using directly-driven targets, even though the energy gains are expected to be less than in radiation-driven targets. For the purpose of this study, we took them to be equal for the same driver wavelength and energy; more accurate data from ongoing physics experiments will be available later. In this examination of direct-illumination effects, we studied four parameters

- Mirror damage constraints.
- Reactor neutronics performance.
- System energetics.
- Cost.

Mirror Damage Constraints. In analyzing laser-light damage to the final-focusing mirror, we find that the maximum $f/\text{No.}$ is limited by the laser damage threshold of the optics. This damage can be minimized by 10 Torr-mm xenon gas, which will protect the mirror surface from both the pellet debris and the x rays generated by the reactor

target. Analysis also shows us that neutron damage is greater on the low $f/\text{No.}$ optics required for direct-illumination targets than on the high $f/\text{No.}$ optics associated with radiation-driven targets. Because of this higher damage, we see an increase in optics costs and in power plant downtime, both factors reducing net revenue from a power-producing plant.

Reactor Neutronics Performance. In a fusion power plant, a liquid-metal blanket in the reactor vessel absorbs neutrons and breeds tritium. For this purpose, we analyzed two blanket materials: lithium and $\text{Li}_{17}\text{Pb}_{83}$. We found that the maximum solid angle of beam penetrations through the blanket will be limited by the tritium breeding ratio. The fusion energy lost through large-solid-angle beam ports represents a substantial loss of revenue from the power plant. When we reviewed reactor design concepts, we also found that the complex irradiation geometry for a directly-driven plant eliminates several of our cavity design options (e.g., HYLIFE) for protecting the first wall from the fusion reactions. Further studies are required to design a cost-effective direct-illumination-driven reactor.

System Energetics. In Table 8-14, we compare the performance of two power plants: one a direct-illumination-driven power plant with 24 beams and $f/3$ final optics, the other a radiation-driven power plant with 2 beams and $f/50$ final optics. Each of them would use either a lithium or a $\text{Li}_{17}\text{Pb}_{83}$ blanket. The data show that the $\text{Li}_{17}\text{Pb}_{83}$ blanket has an unacceptable tritium breeding ratio (<1) in the direct-illumination system. For a fixed fusion energy gain (ηQ) of 10 (where η is the laser system efficiency and Q is the pellet gain), the ratio of the net power to the gross power produced is about 10% less for this direct-illumination system than for the radiation-driven system.

Cost. In assessing costs, we see that the target factory of a direct-illumination plant will be slightly less expensive because the targets are less complex. However, this power plant will also require a more complex and, therefore, more costly reactor and containment structure to accommodate the globally distributed laser beams and related optics. If we consider the total system efficiency discussed in the previous paragraph, for the same net cost of electricity produced,

Table 8-14. Comparison of direct-drive reactor system with radiation-driven reactor system.

	Direct drive	Radiation drive
Optics	$f/3$	$f/50$
Number of beams	24	2
Tritium breeding ratio		
Lithium	1.0	1.3
$\text{Li}_{17}\text{Pb}_{83}$	0.6	1.1
Blanket multiplication		
Lithium	0.90	1.20
$\text{Li}_{17}\text{Pb}_{83}$	0.87	1.30
Net/gross power ratio ($\eta Q = 10$)		
Lithium	0.59	0.66
$\text{Li}_{17}\text{Pb}_{83}$	0.58	0.67

the allowable capital cost of this direct-illumination-driven power plant must be 10% less than that of a radiation-driven power plant. For a \$1.5 billion plant, this represents \$150 million.

In conclusion, we note that the capital costs for low $f/\text{No.}$, direct-illumination-driven inertial-fusion power plants are required to be significantly less than those for the radiation-driven power plants in order to produce electricity at the same cost. High $f/\text{No.}$ direct illumination solves the neutron damage problems for the final optics and will reduce the tritium production and blanket performance problem of low $f/\text{No.}$ direct illumination. The globally distributed optics, however, will result in extra costs, as well as the elimination of several feasible reactor concepts relative to the radiation-driven power plants. Detailed self-consistent reactor design and costing studies would be required to quantify these differences. However, it is clear that high $f/\text{No.}$ illumination of radiation-driven targets, from the fewest possible directions, will maximize reactor performance and minimize plant cost.

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Commercialization of Inertial Fusion for Electric Power Production

Introduction. We have devised a plan to develop inertial fusion for power production.⁸⁵ First, we determined what large facilities we need for ICF development; then, we estimated the time schedule required for construction of these large facilities. Our initial plan included goals and a development schedule for each of the three main development facilities. We have recently proposed⁸⁶ an operating point (target gain, driver energy, pulse rate, etc.) for each facility as a function of the commercial-reactor laser-driver type and pulse rate. We also considered the effect of developing fusion as a source of fissile fuel rather than as thermal power. In setting a schedule and in planning operating points for the major engineering facilities, we also generated milestones for development of the driver, the target factory, and many of the other components of an ICF power plant. This article summarizes our development plan.

Facility Goals. We believe that three integrated facilities will be needed to commercialize inertial fusion

- An engineering test facility (ETF) that tests several potential reactor concepts at the lowest practical fusion power for only a few pulses at a time. The goal of ETF is to verify modeling and simulation results for several reactor first-wall concepts (e.g., dry wall, wetted wall, liquid-metal wall) and then to use the experimental data to choose a first-wall concept for further development. The ETF is assumed to be an add-on to the target physics demonstration facility (TPDF), which will demonstrate the reactor physics required for commercialization. The goal of TPDF is to demonstrate target yields of a few hundred megajoule or more for both military and energy applications.
- An experimental power reactor (EPR) that integrates the selected driver and reactor concept at a low fusion power. The EPR will demonstrate integrated operation of most power plant systems.
- A demonstration power reactor (DEMO) that operates reliably for extended periods. The DEMO fusion power level may be lower than required for economic power production, but it must be large enough, and provide a reliable enough cost data base, for utilities to risk construction of full-size commercial plants.

The timely operation of each facility will require extensive development programs for supporting systems (such as target-fabrication factories, target-injection hardware, tritium-processing equipment, etc.).

The thermal power of each facility can be quite low and still provide energy densities comparable to a commercial reactor because, in ICF, we can move the energy-absorption materials closer to the point source. (Energy density is roughly proportional to the target yield divided by the square of the wall radius.) This design freedom allows the reactor portion of ICF development facilities to be much smaller than magnetic fusion facilities with similar goals. However, the ICF facilities still include large, expensive drivers, and the cost of ICF and MFE development facilities will be of the same magnitude. The high cost of the driver could be offset by the separability of the driver and reactor (allowing more reactor design freedom) and by the

Power-Plant Considerations

possibility of reusing the driver in a subsequent facility.

We have not proposed a materials test facility (MTF) in the time line because we believe that ICF reactor materials selection will be possible using analytical studies, small-scale experiments, and results from other energy programs (LMFBR, coal gasification, and MFE). If we are successful, DEMO will validate our choice of materials; if not, DEMO becomes the MTF, and the schedule must be extended.

The three facility descriptions (ETF, EPR, and DEMO) are based on the viewpoint of reactor development, and the schedule is used to provide milestones for the driver development program. A scenario based on driver development may not distinguish the facilities in the same way. For example, the EPR and DEMO both require a repetition-rated, multimegajoule driver, and it is possible that the same driver will be used in both facilities; thus, DEMO could be considered an upgrade of EPR. Similarly, if a heavy-ion-beam or repetition-rated laser driver is chosen for the TPDF/ETF, that driver could possibly be reused in the EPR and DEMO.

Schedule. The selection of three facilities and the lead-time constraints for large facilities allow the formulation of a schedule. The nominal time line shown in Fig. 8-40

begins in 1985 after some results from Nova and PBFA are available. A preliminary Schedule 44 is not submitted to DOE for any facility (TPDF/ETF, EPR, or DEMO) until the preceding facility has operated for two years (one year start-up, one year of experimentation). As Fig. 8-40 shows, the DEMO would operate in 2016, using these assumptions.

The preliminary driver decision for TPDF/ETF will be made in 1987 under the nominal schedule, and the final decision will be made in 1989, when Title I is complete. Thus, several laser and heavy-ion-beam options will be available for consideration. The preliminary driver decision for EPR will be made in 1998 under the nominal schedule. Therefore, if a single-pulse driver is used for TPDF/ETF, 1998 is the key date for development programs for the driver, target factory, and other high-repetition-rate systems.

This schedule could be compressed in the unlikely prospect that a Manhattan Project-style development of ICF is implemented. For example, if the preliminary design Schedule 44 is submitted for construction, rather than for design funding, the conceptual design and Title I steps are eliminated. Double-shift construction could cut the 4- to 6-year construction time to 3 to 4 years. Results from Nova/PBFA could be made

Fig. 8-40. Engineering development schedule.

available earlier. The accelerated time line has DEMO operating around 2006, ten years earlier than the nominal schedule. Since the EPR driver decision point is moved up four years to 1994, development programs for repetition-rated systems require earlier funding to meet the schedule. This rather unrealistic accelerated scenario is of interest primarily because it defines the earliest possible date for the DEMO.

If the ICF development program continues at the present pace through 1990 and then becomes an accelerated program, the nominal schedule could be compressed with DEMO operating in 2009. In this case, development programs for repetition-rate systems must succeed by 1997.

The schedules described are success-oriented; i.e., they assume availability of funds and timely success in each technology development program and in each of the three integrated facilities. Clearly, unanticipated delays in the schedules are always possible. For example, because of funding limits, only a few candidate structural materials will be developed for use in the EPR and DEMO facilities. If the synergistic effects of fatigue, pulsed radiation, and liquid-metal corrosion drastically shorten the life of the structural components in EPR or DEMO, a considerable delay could result while alternative materials are developed for the next facility.

Characteristics of the Development Facilities. To specify the characteristics of each development facility, we must first consider the characteristics of potential commercial plants as a function of chamber pulse-rate and driver efficiency. Then, the size of each development facility can be chosen to satisfy its design goals while generating a minimum of fusion power. The accomplishments for this procedure have been described in Ref. 86 for five potential chamber-pulse-rate/driver-efficiency combinations. The TPDF/ETF has the same characteristics in each case, but the commercial chamber and the driver types must be selected prior to EPR construction.

The target-gain level demonstrated in TPDF will allow elimination of those driver/reactor combinations with a recirculating power fraction that is too high (product of driver-efficiency and target-gain that is too low). The ETF experiments, coupled with economic and safety analyses, will be

used to select the driver/reactor combination leading to a commercial reactor low in cost and high in safety. Then the EPR and DEMO can be sized as intermediate points leading to the commercial plant.

Fissile Fuel Breeder Development. The high efficiency of using 14-MeV neutrons to convert fertile material (^{232}Th or ^{238}U) to fissile fuel (^{233}U or ^{239}Pu) allows one fusion breeder reactor to support as many as 20 light-water reactors of the same thermal power as the breeder. The high support ratio allows the fusion-breeder cost to be shared among many light-water reactors, as part of the fuel cycle. Hence, the breeder can cost considerably more than a fission reactor and have little impact on the cost of electricity. In addition, the product of target gain and driver efficiency (also called the fusion gain) of a fusion breeder can be much lower (1 to 5) than for a pure fusion reactor (>15).

While economic considerations may permit an earlier application for fusion breeders than for pure fusion, the relaxed fusion gain requirement will probably not allow the development schedule to be accelerated, since the time lines shown in Fig. 8-40 require success in driver and target designs for either the pure fusion or the fusion breeder application. The target gain for a pure fusion commercial plant is scheduled to be demonstrated in TPDF. If TPDF only demonstrates target gains that are suitable for fusion breeder reactors, the pure fusion schedule will be delayed or terminated, and the breeder schedule will proceed. If TPDF demonstrates target gains for both breeders and pure fusion reactors, then both programs can proceed on the same schedule. Therefore, relaxation of target gain will probably not accelerate the schedule for fusion breeders beyond the schedule described in Fig. 8-40. However, if funding is very limited, relaxed target gain could mean a smaller driver that costs less, thus enhancing the probability of obtaining funding for breeder development facilities.

If a heavy-ion-beam or high-efficiency laser is developed for TPDF or EPR, relaxation of the driver efficiency requirement for breeders will have no impact on the development schedule. If a low-efficiency laser is chosen, a fusion breeder program would require less driver development than for a pure fusion driver, but, again, this

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difference cannot accelerate the breeder schedule. Rather, problems with the driver development program could delay the pure fusion time line.

The ETF will probably be able to serve both fusion-breeder and pure fusion development because of its inherent design flexibility. The EPR and DEMO facilities might be able to test both technologies simultaneously by allocating a portion of the reactor chamber to each application, or, more likely, by sharing the driver between two separate chambers.

Although a fusion-breeder schedule might proceed while a pure fusion schedule might be delayed because of target gain or driver efficiency problems, each fusion-breeder development facility will include a

more complex blanket and balance-of-plant than would its pure fusion counterpart. The development of the fuel processing equipment, the afterheat removal system, and other safety systems will make it more difficult for fusion breeder development to stay on schedule, as compared to pure fusion.

The fusion breeder represents an attractive application for fusion. In addition to its favorable economics, the fusion breeder can have a rapid impact on a large fission reactor energy grid because each new breeder will supply makeup fuel for as many as 20 fission reactors. But we do not believe that the development time can be faster than the plan proposed for pure fusion.

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