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**ICF
Program
Annual
Report**

MASTER

Lawrence Livermore
National Laboratory

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Preface

This report documents the major activities and accomplishments of the Inertial Confinement Fusion (ICF) Program, a part of the Laser Programs at Lawrence Livermore National Laboratory (LLNL), during the 1988–89 calendar years. Prior to 1988, ICF documentation was included in the Laser Programs Annual Report, which was last published in 1987. Beginning in 1991, a new reporting structure was established: the ICF Quarterly Reports and the ICF Annual Report. The ICF Quarterly and Annual Reports began documenting activities by fiscal year rather than calendar year, for ease in tracking progress against project plans and budgets (which are based on the fiscal calendar). Because the fiscal year 1991 Annual Report was a new publication, it covered both calendar year 1990 and fiscal year 1991—the period from January 1, 1990 to September 30, 1991.

This 1988–89 ICF Annual Report is organized into seven sections, beginning with an executive summary overview and followed by six additional sections that describe accomplishments within the major research and development areas of ICF: Target Physics, Nova Experiments, Nova Laser Science and Technology, Target Science and Technology, Advanced Drivers, and ICF Applications.

D. L. Correll

Acknowledgments

As in previous years, scientists, engineers, technicians and other staff members continued to provide incredible effort in support of the LLNL ICF program. Not surprisingly, their aid and cooperation in completing this 1988–89 Annual Report have been invaluable. We are especially grateful to those who wrote and contributed to the articles; their efforts are the backbone of this document.

Members of the Laboratory's Technical Information Department shaped the raw material into a final product. Editors Jason Carpenter, Pete Murphy, Joy Perez and Leslie Wilder, with proofreading assistance from Al Miguel, tackled the task of polishing the articles, fitting them together, and pushing the whole package to completion. For the excellent artwork we thank the production staff: Chuck Chapman, Galen Hazelhofer, C. Sue Lawrence, Kim M. Lanto, Daniel Moore, and Don Ryder. Ron Natali saw the document through the printing process.

Many of ICF's secretarial staff contributed, as well. Those who played a principal role were Gloria Baker, Carol Molkenbuhr, and Nanci Strange. Special thanks go to Pamela D. Dickinson, who not only added to the artwork but also provided the layout and managed the data throughout the project's history.

We sincerely thank all those people above, as well as the many unnamed, for making possible the 1988–89 ICF Program Annual Report.

D. L. Correll

1 Overview

Section editors: E. M. Campbell, D. L. Correll

2 Target Physics

Section editors: J. D. Lindl, M. D. Rosen

3 Nova Experiments

Section editors: R. L. Kaufman, J. D. Kilkenny

4 Nova Laser Science and Technology

Section editors: J. R. Murray, H. T. Powell

5 Target Science and Technology

Section editors: T. P. Bernat, R. C. Cook

6 Advanced Drivers

Section editors: W. H. Lowdermilk, H. T. Powell

7 ICF Applications

Section editors: W. J. Hogan, M. T. Tobin

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Overview

Section editors: E. M. Campbell, D. L. Correll

1. ICF Program Overview

The continuing mission of the Inertial Confinement Fusion (ICF) program is to develop the science and technology base necessary to determine the technical feasibility of achieving significant fusion energy yields in the laboratory, to identify applications utilizing that capability, and to develop the most promising of these applications. This mission has required developments in the fields of high energy density physics and high peak-power/energy laser science and technology, and supporting technologies.

Our near-term goal is to apply ICF technology and facilities to provide a better understanding of nuclear weapons physics issues. A longer-term goal is to explore the feasibility of ICF as a clean and inexhaustible source for commercial electric power production. The potential defense applications of ICF physics and technology include studying basic atomic physics and hydrodynamics, extending capabilities to simulate the effects of nuclear weapons on hardware that must function in a nuclear environment, developing instrumentation and techniques that can be applied to full-scale nuclear tests, and exploring the basic principles of advanced weapons concepts.

The primary scientific requirement for realizing the full defense and energy applications of ICF is the achievement of net fusion gain (fusion energy produced/driver energy required). The ICF Program has made substantial progress in ICF target physics and laser science and technology. In each of these areas, progress required the development of experimental techniques and computational modeling. Our recent experiments and modeling are focused on the physics necessary to validate the basic concept of fusion ignition and gain in the laboratory.

The Nova laser is the latest in a series of ever more powerful laser facilities at LLNL. Nova's principal objective is to demonstrate that laser-driven hohlraums (hollow chambers that trap electromagnetic radiation) meet the conditions of beam/target coupling efficiency, beam irradiation symmetry, beam pulse shaping,

target preheat, and hydrodynamic stability required for ignition and gain. The LLNL Nova laser facility is the primary U.S. facility devoted to the study of the indirect (or hohlraum) drive approach to inertial fusion. In this concept, energy from a laboratory driver is converted by the hohlraum walls to radiation that is used to implode and heat the fusion fuel inside of an inertial fusion capsule that is centered within the hohlraum. We continue to use Nova to study inertial confinement fusion, laser-plasma interaction physics, weapons physics, and laser science.

The major research and development areas of our ICF program fall in the following categories: Target Physics, Nova Experiments, Nova Laser Science and Technology, Target Science and Technology, Advanced Drivers, ICF Applications, and Halite/Centurion.

Our Target Physics advances included significant code development: the new DIMPLE code calculates the evolution of small-amplitude perturbations from spherically symmetric flow; the LASNEX code helped us investigate holographic interferograms; and a new package in LASNEX can calculate electron heat conduction. In atomic physics, we developed a new "atom-in-jellium" condensed-matter model. X-Division's laboratory x-ray laser (LXRL) research included new table-top LXRL design approaches, study of large-laser-pumped LXRL systematics, calculation of LXRL-source coherence for holography, and new escape-probability methods for planar and cylindrical geometries. In addition, we investigated heavy-ion particle accelerators as drivers for inertial fusion power production. We also made advances in the theory and design of targets, in target fabrication and materials science, in laser and optics research, and in diagnostics. With the benefit of experiments at higher drive energies, and using the sophisticated temporal pulse shaping capability of Nova and the advances in diagnostic capability, future work will refine target designs.

Our Nova Experiments have covered much ground. We have learned that a correlation can

exist between second-harmonic emission and density perturbations from an underdense plasma; that plasma densities and temperatures at late time depend on initial target geometry; and that harmonic light images can be significantly refracted. We incorporated SIMPLE, a one-dimensional, three-fluid, local thermodynamic equilibrium, hydrodynamics code for small computers, into our computational studies. Other work in hydrodynamics promises improved computational models. And in 1989, the LLNL laboratory x-ray laser team produced amplified spontaneous emission (i.e. lasing) at 4.48 nm in nickel-like tantalum. This is a significant step toward the achievement of an x-ray laser source for holographic microscopy of living biological cells. In diagnostics, we developed a multichannel Wölter x-ray microscope for time and space resolved high resolution x-ray images of laser fusion experiments; a thick aperture design for neutron penumbral imaging that achieves at least 10- μ m resolution over a 150- μ m field of view without significant first-order distortion; and the MOSFET/NLTL avalanche diode pulser, an alternative picosecond high voltage pulser for new gated x-ray camera structures and ultra-fast transient digitizer topologies.

Supporting the Nova experiments are the Nova Operations and Laser Science groups, whose accomplishments in 1988 and 1989 were many. Installation of platinum-free glass and fused silica lenses brought Nova to its full operational potential. We have evaluated system damage, crystal conversion efficiency, and beam uniformity, making changes to beam apodizers, amplifier configurations, beam fill-factors and pinhole sizes. We began or continued upgrades and maintenance in streak cameras, flashlamps and oscillator capability. Nova continues to operate 12 hours a day, with Monday set aside as a maintenance day. During 1989, we fired 584 high-power shots to provide the drive source for 675 experiments. The energies provided by the laser (at a wavelength of 1.06 μ m) ranged from <1.0 kJ to >120 kJ. Pulse durations varied from 10 ps to several nanoseconds. Temporal pulse shapes included square, Gaussian, and specially shaped pulses designed to improve target performance.

Our Target Science and Technology Program provides targets for fusion experiments, x-ray lasers, and laser matter-plasma interaction studies, as well as experimental activities in

high-density/temperature materials, high-acceleration instabilities, laser driver characterization, and diagnostic development. To achieve this, the program's work includes fabrication and production, development, characterization, and material research.

The Advanced Drivers Program has two objectives: upgrade Nova's performance capabilities, and develop future drivers for the ICF Program. The Precision Nova Project aims to improve target irradiation through better energy and power balance, pointing accuracy, and diagnostics. The beam "smoothing" project looks to develop large aperture phase plates for spatial incoherence and broadband, multifrequency laser oscillators for temporal incoherence. In pursuit of the second objective, we focused on two laser concepts, "Loki" and "Athena." Although we decided to discontinue work on Loki, we nevertheless invented in its support a new laser material, chromium doped lithium calcium aluminum fluoride (LiCaF). Most of our work went toward Athena, with an objective to meet all technical requirements for the Laboratory Microfusion Facility. Our activities, aimed at making Athena less expensive than previous Nd:glass laser systems, included improvements in system efficiency and increases in energy storage density and laser-induced damage thresholds. In addition, the Advanced Drivers Program completed other research in optical materials and laser design and performance.

Investigating ICF applications, we made substantial progress in the development of concepts and technologies for the next higher energy laser facility for ICF. The ICF Applications group revised both the HYLIFE-I and CASCADE power plant designs, including the development of a vacuum system for both concepts.

In a program called Halite/Centurion, a portion of energy from a nuclear device in underground explosions at the Nevada Test Site was used to implode inertial fusion capsules. Halite is the name of the program under the leadership of LLNL; a similar program named Centurion was guided by Los Alamos National Laboratory. Because the ability to fully study and understand the performance of ICF capsules in the laboratory has been limited by the energy and power that can be provided with presently available lasers, the Halite/Centurion program extended the range of inertial fusion research. As

described in a Department of Energy (DOE) Inertial Fusion fact sheet, "the combination of data obtained from laboratory experiments and the Halite/Centurion program has enabled the DOE's Inertial Fusion Program to progress at a faster rate than anticipated. Many issues related to its scientific feasibility have been firmly established. As a result, DOE has begun preliminary planning to determine whether a new facility is feasible and affordable. The new facility under study, called the Laboratory Microfusion Facility (LMF), would be capable of achieving fusion routinely and generating one billion joules (1000 MJ) of energy, the energy equivalent to one quarter ton of TNT. The DOE Inertial Fusion Program Office is currently proceeding with studies aimed at defining the cost basis and performance of the LMF."

While Nova has insufficient energy to achieve ignition and gain, we believe that detailed and quantitative experiments in future Nova hohlraum and capsule physics, coupled with the experimental results achieved with Halite/Centurion, will provide a sufficient data base to confidently predict ignition and gain performance within the laboratory.

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Target Physics

Section editors: J. D. Lindl, M. D. Rosen

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2. Target Physics

2.1. Introduction

The ICF Program's target physics work covers a broad range of topics. Implosion and ignition physics includes Rayleigh-Taylor instability modeling and simulations. Code development work has covered much ground:

- We developed the DIMPLE computer code, which calculates the evolution of small-amplitude perturbations from spherically symmetric flow.
- A new package in the LASNEX code calculates electron-heat conduction using the B.G.K. approximation to the Fokker-Planck equation.
- We derived the escape probability in the Sobolev approximation in a form usable for simulations of radiative transfer in media with cylindrical symmetry, such as x-ray laser experiments.
- LASNEX simulations allowed us to investigate holographic interferograms of laser fusion-type targets.
- We demonstrated that the finite element method is superior to finite-difference based methods for computations using orthogonal meshes.

We completed a variety of projects in atomic physics, such as a new "atom-in-jellium" condensed-matter model and various applications of the Liouville equation. In plasma physics research, we continued to characterize and control the important mechanisms for laser-plasma coupling. X-Division continued important research in laboratory x-ray laser (LXRL) design and analysis. We explored new approaches to table-top LXRL systems, studied the systematics of large-laser-pumped LXRLs, developed approaches to calculating and achieving LXRL-source coherence for holography, and found efficient escape-probability methods for planar and cylindrical geometries.

Heavy-ion fusion (HIF) research progressed significantly through 1989. We have continued to investigate heavy-ion particle accelerators as drivers for inertial fusion power production; this work has been presented to the Department of Energy Fusion Policy Advisory Committee and

the National Academy of Sciences, and their reactions have been encouraging.

2.2. Implosion and Ignition Physics

2.2.1. Introduction

The articles in this section describe some of the work we have been doing on Rayleigh-Taylor instability physics.

2.2.2. LASNEX Simulations of the Classical and Laser-Driven Rayleigh-Taylor Instability

We present the results of two-dimensional LASNEX simulations of the classical and laser-driven Rayleigh-Taylor instability (Fig. 2-1). Our growth rates and eigenmodes for classical two- and three-fluid problems agree closely with the exact analytic expressions. For targets driven by a $1/4\text{ }\mu\text{m}$ laser at $I = 2 \times 10^{14}\text{ W/cm}^2$ our growth rates are 40 to 80% of classical for wavelengths between 5 to $100\text{ }\mu\text{m}$. These are larger than recently reported growth rates.

Author: K. O. Mikaelian¹

2.2.3. Stability and Mix in Spherical Geometry

We consider a spherical system composed of N concentric fluid shells having perturbations of amplitude η_i at interface i , $i = 1, 2, \dots, N-1$. For arbitrary implosion/explosion histories $R_i(t)$, we present the set of $N-1$ second order differential equations describing the time evolution of the η_i which are coupled to the two adjacent $\eta_{i \pm 1}$. We report analytical solutions for the $N = 2$ and $N = 3$ cases. We also present a model to describe the evolution of a turbulent mixing layer in spherical geometry when the interface between two fluids undergoes a constant acceleration or a shock.

Author: K. O. Mikaelian²

2.2.4. Rayleigh-Taylor and Richtmyer-Meshkov Instabilities and Mixing in Stratified Spherical Shells

We study the linear stability of an arbitrary number of spherical concentric shells undergoing

Figure 2-1. Growth rate γ vs. λ . The LASNEX points are indicated by crosses. The continuous curve is the classical growth rate $(gk)^{1/2}$. The dashed curve is an approximate fit $[gk\beta/(k+\beta)]^{1/2} - Vk$ with $\beta^{-1} = 1 \mu\text{m}$ and $V = 2 \mu\text{m/ns}$. These values are close to the density gradient length and ablation velocity respectively. Other fits are also possible.

a radial implosion or explosion. The system consists of N incompressible fluids with small amplitude perturbations at each of the $N-1$ interfaces. We derive the evolution equation for the perturbation η_i at interface i ; it is coupled to the two adjacent interfaces via $\eta_i \pm 1$. We show that the $N-1$ evolution equations are symmetric under $n \leftrightarrow -n-1$, when n is the mode number of the spherical perturbation, provided that the first and last fluids have zero density ($\rho_1 = \rho_N = 0$). In plane geometry this translates to symmetry under $k \leftrightarrow -k$. We obtain several analytic solutions for the $N = 2$ and $N = 3$ cases, which we consider in some detail. As an application we derive the shock timing required to freeze out an amplitude. We also identify "critical modes" which are stable for any implosion/explosion history. Several numerical examples are presented to illustrate perturbation feed-through from one interface to another. Finally, we develop a model for the evolution of turbulent mix in spherical geometry, and introduce a geometrical factor $G.F.$ relating to the mixing width h in spherical and planar geometries via $h_{\text{spherical}} = h_{\text{planar}} \times G.F.$ (Fig. 2-2). We find that $G.F.$ is a decreasing function of $R(t)/R(0)$, where $R(t)$ is the radius of an interface at time t and $R(0)$ is its initial radius. This implies that in our model $h_{\text{spherical}}$ evolves faster (slower) than h_{planar} during an implosion (explosion).

Author: K. O. Mikaelian³

2.2.5. Rayleigh-Taylor and Richtmyer-Meshkov Instabilities in Multilayer Fluids with Surface Tension

Surface tension modifies the evolution of the Rayleigh-Taylor and the Richtmyer-Meshkov instabilities in fluids undergoing a constant acceleration or a shock respectively. We analyzed the general case of N fluids with arbitrary densities and surface tensions, and derived the eigenvalue equation determining the growth rate of the perturbations. For $N = 2$ we recovered the classical case of two semi-infinite fluids and extended it to the case of two finite-thickness fluids between fixed boundaries. The $N = 3$ case is studied in detail; we found universal modes which were independent of the thickness of the intermediate fluid, and we found how surface tension modified Taylor's modes for a single fluid with free boundaries. We also analyzed in detail recent and future two- or three-fluid experiments. Representing a shock as an impulsive

acceleration, we found that post-shock oscillations had frequencies and amplitudes that depended on the wave number k , leading to a nontrivial evolution for the spectrum of perturbations. Finally, we studied turbulence at the interface between two fluids with surface tension and presented specific predictions for the turbulent energy $E_{\text{turbulent}}$ as a function of the surface tension $T(s)$ (see Fig. 2-3). We propose new experiments, physical and/or numerical, to test our predictions.

Author: K. O. Mikaelian⁴

2.3. Code Development and Atomic Physics

2.3.1 A Linear Perturbation Code, DIMPLe

Introduction. We have written a computer code that calculates the evolution of small perturbations from time-dependent, spherically symmetric flow. Arbitrary configurations can be decomposed into sums of terms having the angular structure of spherical harmonics, Y_{lm} . Since the unperturbed state is spherically symmetric, the temporal history of the solution depends only upon l , and is degenerated in m . If the perturbations are small in amplitude, then the evolution of each l -mode is described by partial differential equations in one dimension of space and time, which are linear in the perturbation amplitudes and not coupled to other l -modes. The code DIMPLe calculates the temporal evolution of specified l -modes, given the spatial structure of the initial perturbation.

The model for interface mix developed by Haan⁵ requires a spectrum of modal growth factors. For several years, we have been calculating such growth factors by running 2-D LASNEX calculations, one for each mode. Typically, such calculations need 10 to 20 zones to resolve one-half or one wavelength of the perturbation in the angular direction. If the initial perturbation amplitude chosen is very small, perhaps 10^{-9} cm for a surface finish perturbation, LASNEX is capable of preserving the single-mode structure. Other tests of the numerical method, such as linearity as the initial perturbation amplitude is varied, are also satisfied. However, each such 2-D calculation consumes about 2 to 10 Cray processor hours, and 10 or more such calculations must be run to map out the mode growth

Figure 2-2. Geometrical factors for a shock and for a constant acceleration. These factors connect the mixing widths in plane and spherical geometries: $h_{\text{spherical}} = h_{\text{planar}} \times \text{G.F.}$, where $h_{\text{planar}} = 0.14 A \Delta v t$ and $h_{\text{planar}} = 0.07 A g t^2$ for a shock and a constant acceleration respectively.

Figure 2-3. The turbulent energy $E_{\text{turbulent}}$ in units of Ag/k_0 as a function of k_c/k_0 ($=\lambda_0/\lambda_c$), where λ_c is the cut-off wavelength and λ_0 is the largest wavelength present in the system. The lower curve includes the reduction in the growth rate due to surface tension. The upper curve is the tensionless case.

spectrum. There may be several seed sources for a given capsule, such as surface finish perturbations at different interfaces and shell inhomogeneities. Therefore, it is time consuming to apply the model to a given capsule design, and virtually prohibitive to examine how predictions scale as capsule design parameters are varied.

A linearized code such as DIMPLE needs to perform essentially one 1-D hydro calculation for each mode, and can calculate any number of different modes or initial perturbation structures in parallel. Therefore, such a code can be a much more efficient way to obtain a spectrum of mode growth factors.

Numerical Method. This method for applying linearized hydrodynamics to ICF perturbation growth was developed by McCrory⁶ et al. Morse⁷ wrote a post-processor to LASNEX and RIPPLE that applied this technique. RIPPLE was awkward to use as massive dump files were needed to couple information into RIPPLE from the 1-D LASNEX calculation that described the unperturbed flow. Also, RIPPLE results agreed poorly with those of 2-D LASNEX. We believed that improvements to the RIPPLE implementation of this technique were possible, which would permit application to our mode growth problems.

Numerical hydrodynamics suffers from noise on the scale of the zone size, especially for problems that contain shocks and sharp material boundaries. Such noise can be a source for perturbations. In DIMPLE, we sought to minimize this by differencing the first order hydrodynamic equations identically to the zero order equations. The zero order hydrodynamics is calculated inside DIMPLE at the same time as the first order hydrodynamics, so all zero order flow variables are readily available at every cycle, with the appropriate spatial and temporal centering. However, we have not implemented transport of laser energy or x-rays. Energy deposition from these sources must still be coupled in from another calculation, such as a 1-D LASNEX run. DIMPLE assumes that the energy deposition per unit mass is not perturbed at first order.

We have used the perturbed hydrodynamics equations as written by Munro.⁸ As this reference has not been published, we shall reproduce Munro's equations here. The zero order flow quantities are functions of the spherical radius, r , and time, t . (Planar problems can be calculated by employing a radius which is large compared to any other dimensions in the problem.) Scalar

quantities such as the density are expanded in the form,

$$\rho = \rho_0(r, t) + \rho_1(r, t)Y_{\ell m}(\theta, \psi). \quad (1)$$

Quantities with subscript 0 refer to the unperturbed (spherically symmetric) flow, and those with subscript 1 refer to first order perturbations. The displacement of a fluid element from the unperturbed position is a vector quantity written in the form,

$$r_1 = \xi(r, t)\hat{r}Y(\theta, \psi) + \eta(r, t)W, \quad (2)$$

where

$$W = r \frac{\nabla Y}{\sqrt{\ell(\ell+1)}} \quad (3)$$

We have dropped terms proportional to the third vector spherical harmonic, $r \times \nabla Y$, as such terms, representing toroidal rotation, decouple from the other vector components, and can be explicitly integrated. In this notation, the equation of mass conservation becomes

$$\rho_1 = \rho_0 \left(\Gamma_1 - \frac{1}{r^2} \frac{\partial}{\partial r} r^2 \xi + \frac{\sqrt{\ell(\ell+1)}}{r} \eta \right) \quad (4)$$

Here, the quantity $\Gamma_1(r)$ is constant in time, and permits introducing density perturbations independent of those introduced by displacements from the unperturbed flow. This quantity has the character of a gauge field, and is similar in effect to the freedom of choosing how to zone a 2-D calculation. One could introduce a ripple on the surface of a sphere by displacing the nodes of a uniformly zoned spherically symmetric problem in a Lagrangian fashion, preserving the zonal masses. However, one is more likely to perturb only the radii of the nodes, that introduces variations of the zonal masses. We can see from Eq. (4) that if we initialize a problem with a ξ perturbation, but with no η perturbation, we shall have to specify a Γ_1 to get zero initial density perturbation.

The rY and W components of the first order momentum equation are,

$$\rho_0 \ddot{\xi} + g_0 \rho_1 + \frac{\partial p_1}{\partial r} - \frac{\partial p_0}{\partial r} \frac{\partial \xi}{\partial r} = 0, \quad (5)$$

and,

$$\rho_0 \ddot{\eta} + \frac{\sqrt{\ell(\ell+1)}}{r} \rho_1 - \frac{\partial \rho_0}{\partial r} \left(\frac{\sqrt{\ell(\ell+1)}}{r} \xi - \frac{\eta}{r} \right) = 0, \quad (6)$$

where g_0 is the zero-order radial acceleration, and p_0 and p_1 are the zero- and first-order pressures (include artificial viscosities, applied as pressures). The energy conservation equation, including electron conduction, is

$$\begin{aligned} \rho_0 \dot{\epsilon}_1 = & -\rho_1 \dot{\epsilon}_0 + \frac{1}{r^2} \frac{\partial}{\partial r} \kappa_0 r^2 \frac{\partial T_1}{\partial r} - \frac{\sqrt{\ell(\ell+1)}}{r^2} \kappa_0 T_1 - \frac{2}{r^2} \frac{\partial \xi}{\partial r} \frac{\partial}{\partial r} \kappa_0 r^2 \frac{\partial T_0}{\partial r} \\ & + \frac{\sqrt{\ell(\ell+1)}}{r^2} \kappa_0 \xi \frac{\partial T_0}{\partial r} - \frac{\sqrt{\ell(\ell+1)}}{r^2} \kappa_0 \eta \frac{\partial T_0}{\partial r} + \frac{1}{r^2} \frac{\partial}{\partial r} \kappa_1 r^2 \frac{\partial T_0}{\partial r} - \kappa_0 \frac{\partial T_0}{\partial r} \frac{\partial^2 \xi}{\partial r^2}, \end{aligned} \quad (7)$$

where κ is the thermal conductivity, ϵ is the internal energy per unit mass, and T is the temperature. The specification of the equations of state, $p(\rho, T)$ and $\epsilon(\rho, T)$, and the thermal conductivity, $\kappa(\rho, T)$ complete the definition of the problem. Note that all of the equations, and hence the solution, are independent of m .

DIMPLE uses the modified Thomas-Fermi equation of state, QEOS,⁹ which is smooth, thermodynamically consistent, applicable to cold, dense matter and to hot plasmas. QEOS provides thermodynamic derivatives needed for the first-order dynamical equations. Currently, we use a simple $\kappa \propto T^{2.5}$ conductivity. The zero-order artificial viscosity duplicates the Von Neumann 'Q' in LASNEX, while the first order 'Q' is scaled in proportion to $(\nabla \cdot \mathbf{v})^2$, and is applied as a scalar pressure. Energy deposition from a laser or from radiation is read from an input file as a rate per unit mass for each radial zone number.

DIMPLE is written in the interpreted language, PLAN, which was developed by Munro.¹⁰ The physics equations may be written compactly and executed efficiently owing to the vector syntax of the language. Particularly useful for describing finite difference equations are subscript operators that provide pair-wise differences of vector elements, and that transform between point-centered and zone-centered arrays. PLAN also provides automatic memory allocation, simple commands to generate sequenced output dumps, interactive graphics, and a built-in debugger. If all of these operations had to be coded up in a language such as FORTRAN, it would have taken many times longer to bring DIMPLE into operation.

Test Problems. DIMPLE has been tested on problems with known analytic solutions and

against explicit two-dimensional hydrodynamic calculations of growth of single modes. The classical incompressible Rayleigh-Taylor problem is approximated with an ideal gas with polytropic index γ equal to 10, similar to the test of LASNEX described by Tabak et al.¹¹ DIMPLE reproduces the exponential profile of the displacement versus position and gives exponential growth with a rate within a few percent of the classical value.

A more challenging test is propagation of a planar shock front through a region in which the density is rippled along the direction normal to the shock propagation direction, $\rho_u(x, r) = \rho_{0u} + \rho_{1u} \cos kx$. Here, the subscript u indicates conditions upstream of the shock, and k is the perturbation wave number, regarding the problem as planar. Munro¹² has published the analytic solution to this problem, and has employed it to test LASNEX. We have set up our test calculation almost identically to Munro's.

Figure 2-4 compares results of the DIMPLE test problem to Munro's LASNEX calculation and to the analytic solution. A piston of Mach number 54.8 drives a shock into an ideal gas, from the right as shown in **Figure 2-4**. Zones 200 to 400 are initially unperturbed to avoid wave reflections from the problem boundaries. At zone number 200, the shock enters a region in which the density is periodically perturbed. The figure is a snapshot at a time just before the shock reaches the left boundary of the problem, and shows the ratio of the density perturbation amplitude to the initial amplitude. Both the DIMPLE and LASNEX calculations show lower spatial frequency transients, which are not described by the analytic solution, and ringing at the Nyquist frequency, which we have filtered strongly in the plots. The

Figure 2-4. Density perturbation amplitude divided by initial amplitude, vs. zone number, for rippled shock calculation. Left—DIMPLE calculation and analytic solution (smooth curve), right—LASNEX calculation from Munro.

ringing results from grid-scale noise in the zero-order hydrodynamics that drives oscillations in the first-order solution. Quadratic artificial viscosity damps oscillations in the zero-order hydrodynamics, but is ineffective at damping the first-order noise. Note that the transient oscillations are very similar for both calculations. There are some slight differences in the way the two problems were set up, and the snapshot times may not be quite the same, so the two results should not be expected to be identical. Both calculations show errors in first-order state variables of order 20 to 30%. Munro has shown that the error in the LASNEX calculation scales in a manner commensurate with the effect of finite zone size ($k\Delta r = 0.31$ preshock).

We have also tested DIMPLE against explicit two-dimensional hydrodynamic calculations for direct-drive acceleration of planar foils. One example is the acceleration of a 20- μm -thick CH foil, driven with a 0.26- μm laser at intensity $2 \times 10^{14} \text{ W/cm}^2$. Both calculations employ an energy deposition source based upon a 1-D simulation to represent the laser. The explicit 2-D hydro calculation uses 11 transverse zones to resolve half of one wavelength, and employs an initial perturbation amplitude of $10^{-6} \mu\text{m}$ to assure that the perturbation remains in the linear regime. **Figure 2-5** is a snapshot comparing the mode structure resulting from a 10- μm -wavelength perturbation initially on the surface facing the laser. The ablation front is near zone 58 at the time shown. The oscillations in the explicit hydrodynamic calculation occur far out in the corona, and are probably unimportant. Hydrodynamic calculations tend to become unstable when zones become very elongated along the direction of pressure gradients (i.e., radially for these calculations). DIMPLE suffers from the same problem when its radial zones become large compared to the transverse wavelength being modeled. This instability has been suppressed artificially in DIMPLE by limiting the size of some terms in the equations for large $Q'\Delta r/r$. In other respects, the mode structure in DIMPLE agrees quite well with that of the explicit 2-D hydrodynamic calculation.

Figure 2-6 compares perturbed column density growth factors, as functions of time, for the same two calculations as above. It may be seen that the two calculations agree to within 30% in growth factor, and to much greater accuracy in growth rate, except during a transient oscillation

near zero amplitude at around 0.4 ns. We regard this agreement as satisfactory given that the density gradient near the ablation front is almost never resolved well for simulations of directly driven foils, which makes such calculations difficult.

Conclusions. The DIMPLE code calculates evolution of small-amplitude perturbations away from spherically symmetric time-dependent hydrodynamic flow. This is a more efficient method for calculating the evolution of perturbations whose structure is that of single spherical harmonic modes than performing 2-D hydrodynamic simulations that resolve the modal structure. DIMPLE has been checked with a variety of planar test problems including the classical incompressible Rayleigh-Taylor problem, shock propagation through a region of rippled density, and perturbation growth for directly driven foils.

Author: S. V. Weber

2.3.2. A Fokker-Planck Treatment of Suprathermal Electron Heat Conduction

A new package that calculates electron heat conduction via the B.G.K. approximation to the Fokker-Planck equation is installed in LASNEX. The package produces heat fluxes which are more accurate than those produced by local Spitzer-Harm theory in regions of temperature and density where nonlocal electron transport is important. The B.G.K. approximation involves replacing the electron-electron energy exchange terms in the Fokker-Planck equation by a simple relaxation term. However, the electron-ion energy exchange terms are retained except that a large velocity approximation has been applied to them. The resulting equation is linear in the distribution function, but must be iterated to get a self-consistent potential field.

The package has been tested by David Eder. The test case was aluminum foil heated by a 20-ps-laser pulse with a wavelength of 0.53 μm . The conclusion of his study is that for times corresponding to the peak of the laser pulse and greater, the temperature and density profiles from the new package agree very well with calculations using flux limiters. However, during the leading edge of the laser pulse the new package gives quantitatively different results from the flux limited case due to the fact that the new package allows for the possibility of energetic electrons

Figure 2-5. Mode structure for 10- μm -wavelength perturbation on 20- μm -thick CH foil, driven with $2 \times 10^{14} \text{ W/cm}^2$ 0.26- μm -laser. a) DIMPLE, b) explicit 2-D hydro. 2-D hydro calculation shows numerical noise far out in corona.

Figure 2-6. Growth factor vs. time for $\lambda = 10 \mu\text{m}$ on $20 \mu\text{m}$ -thick CH foil, $2 \times 10^{14} \text{ W/cm}^2$ laser intensity, 0.1 ns rise time.

that can reach the center of the target before the heat front.

Author: E. Alley¹³

2.3.3. Cylindrical Escape Probabilities for X-Ray Laser Modeling

The escape probability in the Sobolev limit in 1-D cylindrical symmetry is derived. The result is an integral, a function of the optical depth and another variable determined by the velocity field. Evaluating the integral on the cylindrical axis gives an analytic expression involving Bessel functions, a result which is also valid for self-similarly expanding or contracting media. Off-axis, the integral is approximated by three series expansions. One series is an expansion about the axial result. For velocity fields in which the velocity gradient dv/dR dominates the ratio v/R , the integral reduces to the appropriate formula for 1-D slab symmetry. The second series is an expansion about this slab result. The third expansion is valid for small optical depth τ . As τ increases, the escape probability is asymptotically equal to $1/3 \tau$. The escape probability is used in a simulation of a cylindrical x-ray laser experiment. The simulation is characterized by large velocity gradients. Therefore, it is ideally suited to the Sobolev theory of escape probabilities. The cylindrical escape probability is always greater than the analogous slab result that results in an electron shell population distribution more favorable for lasing. Use of the cylindrical instead of the slab formula results in better values of the laser gain. These higher values are in better agreement with the experiment.

This work is used to simulate radiative transfer in fast-moving media with cylindrical symmetry. Such plasmas occur in x-ray laser experiments. The large velocity gradients imply that photons emitted locally are either locally reabsorbed or escape from the problem altogether since atoms a short distance away have Doppler-shifted and were unable to trap the escaping radiation. The use of escape probabilities greatly simplifies and speeds up the calculation since one then simply multiplies the radiative rate by the escape probability. Before the appearance of the above paper, one knew that escape probabilities were useful in modeling such plasmas but one did not know what expression to use for the probability, much less how to compute it. Some used the expression for a slab multiplied by an

arbitrary factor knowing the slab value to be incorrect. We now have not only the exact expression but also an efficient means of evaluating it. D. Eder has published several papers using our result, to wit:

- D. C. Eder, "Modeling of a Cylindrically Expanding Fluorine X-ray Laser," *Phys. Fluids B*, **1**, 2462 (1989).
- D. C. Eder, "Hydrogen-like Magnesium X-ray Laser Design," Lawrence Livermore National Laboratory, Livermore, CA, UCRL-JC-103375 (1989); *Phys. Fluids B*, **2**(12) 3086 (1990)
- D. C. Eder and H. Scott, "The Calculation of Line Transfer in Expanding Media," Lawrence Livermore National Laboratory, Livermore, CA, UCRL-JC-103769 (1989); *J. Quant. Spectrosc. Radiat. Transf.* **45**(4) 189-204 (1991).

In addition, the cylindrical expression is now used in other codes besides LASNEX and by people at other sites besides LLNL viz: XRASER at LLNL, Rutherford Appleton Laboratory by S. Rose, University of York by G. J. Pert, and Naval Research Laboratory by J. W. Thornhill.

Authors: A. I. Shestakov and D. C. Eder¹⁴

2.3.4. Holographic Interferograms from LASNEX Simulations

Laser fusion-type targets can be probed using a double exposure holographic technique that can provide high-resolution images and plasma electron density distributions at various probe times.¹⁵ Typically, a flat or spherical target is irradiated with a main laser pulse that creates an outwardly expanding plasma from the target surface. In double exposure holography a secondary reference probe beam is passed through a test region before and after the creation of the plasma to interfere and thereby create holographic interferograms. The fringes are contours of the optical path length (OPL) difference between the two exposures of the probe beam and are closely approximated by the isodensity contours in the plasma. The purpose of the imaging optics in such experiments is to bring a selected plane in the plasma to focus on the image. Since the plasma has a locally varying index of refraction, the probe beam will deviate from straight line propagation that can lead to distortions in the interferogram. In particular, the

interferogram should be quite sensitive to the plane of focus.

With a laser fusion code such as LASNEX which has a ray-tracing capability, one can send probe beams through any simulation and construct contours of the OPL difference, i.e., holographic interferograms. In particular we use the ray-tracing package to extract the entry and exit position and direction of a ray, along with the OPL through its curved path. This information is then postprocessed to arrive at an interferogram. First the exiting ray is projected back onto a selected focus plane to arrive at an "apparent" object point Q from which the ray appears to come, insofar as the imaging optics are concerned. Then a parallel straight line ray is sent through Q to simulate the first exposure without the plasma. Assuming the probe beam is collimated, there are no relative changes of OPL beyond a large sphere centered at Q and enclosing the plasma. The OPL difference between the straight and curved rays through Q can then be calculated. This procedure is carried out for a set of probe rays launched from a planar grid and results in a grid of "apparent" object points Q each with an associated OPL difference. Finally, a random mesh contour plotting routine within BASIS is used to plot contours of the OPL difference on the focus plane.¹⁶

We have carried out the above post processing on various 1-D and 2-D LASNEX simulations, some for which we have experimental fringe data. An interesting example is KMS shot 10815: aluminum dot, 270 μm diam and .33 μm thickness, mounted on formvar of thickness .0358 μm . The formvar is mounted on an aluminum washer, 15 μm thick, 500 μm i.d. and 800 μm o.d. The main laser pulse, 53 J in 1 ns and 1500 μm diam-focus, is incident on the formvar. The experimental interferograms are taken with 1/4 μm and 20 ps duration probe at various times from the onset of the main laser pulse. A photographic print taken at 400 ps shows pronounced plasma turbulence on the back side with rather symmetric location of the outer fringes on both sides. The 2-D LASNEX simulation of this shot coupled with the holographic postprocessor agrees very well with the prints on the front side where the laser is incident. The back side, however, shows irregular fringe structure with no apparent symmetry for the outer fringes. This is also clear from the electron isodensity contours of the 2-D simulation which the fringes should approximate. It appears

therefore that LASNEX may not be correctly modeling some essential physics in the energy transport around the target which would give a more symmetric fringe pattern.

Authors: M. K. Prasad, K. G. Estabrook, J. A. Harte, R. A. Bosch, G. E. Busch, and J. S. Kollin

2.3.5. Test Problems for Radiative Transfer Codes

Several test problems are presented for evaluating the radiation diffusion equations. For spatial transport schemes, 1-D problems with known analytic solutions are tested on 2-D domains with nonorthogonal meshes. It is shown that a scheme based on the Finite Element Method is insensitive to grid distortions when the diffusion term is dominant. Other test problems deal with Compton scattering, specifically the 1-D Fokker-Planck equation coupled to an equation describing the change in electron temperature. The test problems model the evolution of a Planckian radiation field as it equilibrates with the electrons. In all cases, the numerical results are compared with the analytic ones.

This paper proves beyond all doubt that if computing diffusion on a nonorthogonal mesh, the finite element method is superior to previously used methods based on finite differences. We may safely claim this on "random" meshes and on grids smoothly distorted away from orthogonality. The claim is also valid for pure and for flux-limited diffusion. The paper was judged sufficiently interesting to be awarded a cash (i.e., real money) prize when it was presented at the International Topical Meeting on Advances in Nuclear Engineering Computation and Radiation Shielding held in Santa Fe, New Mexico, on April 9-13, 1989.

Authors: A. I. Shestakov, D. S. Kershaw, and G. B. Zimmerman¹⁷

2.3.6. Atomic Vibrations in a Self-Consistent Field Atom-in-Jellium Model of Condensed Matter (Fig. 2-7)

The electronic structure of many of the elemental solids is fairly well represented by a single atom embedded in a degenerate electron gas. This is particularly true for bulk properties such as the equation of state of highly compressed matter. Because the atom-in-jellium model is spherically symmetric, it is simpler than band structure models. We study what happens when

Figure 2-7 Transition metal Debye temperatures as calculated by the atom-in-jellium model, o, and derived from data.¹⁸

the nucleus is moved off-center in the atomic cell to form a nonspherical system. This forms the basis of an Einstein model of atomic vibrations. The model is used to calculate Einstein temperatures and Gruneisen constants of simple solids by self-consistent field electronic structure methods.

Authors: D. A. Liberman and B. I. Bennett¹⁸

2.3.7. Modeling Doppler Broadening in Raman Propagation

In December 1986, Karr described an elegant treatment for inhomogeneous broadening of Raman lines, within the context of an integral formulation of the Bloch equation. This formulation relied, for important simplifications, on the assumptions that the molecular population inversion be independent of velocity and that the velocity distribution be Maxwellian. Here we derive an alternative description — a truncated velocity moment series for the differential Liouville equation.

Author: R. A. Sacks¹⁹

2.3.8. Efficient, High-Order Solution Technique for the Liouville Equation

A numerical technique is developed for solving the Liouville equation describing the evolution of Raman-active atomic or molecular states during the passage of a multimode pump beam. Since the approach is a power-series expansion of the exponential time-propagator, it is arbitrarily accurate for arbitrary time step. An iterative technique for generating the individual terms is efficient. Either uncoupled magnetic sublevels or multipole state expansions can be treated.

Authors: R. A. Sacks, B. Shore, M. Hermann and S. Dixit²⁰

2.3.9. Liouville Equation in the Presence of Non-Colinear Light Beams

Explicitly solving the Liouville equation for a Raman active medium on a spatial grid sufficiently fine to represent significant nonparaxial content to the pump beam can be prohibitively expensive. We developed a method for representing the light as small angle cones about a set of discrete propagation angles, and extracting a polarization field consistent with that representation.

Authors: R. A. Sacks, B. Shore, M. Hermann and S. Dixit²¹

2.3.10. Generation of Harmonic Radiation During Electron Scattering from a Piecewise Constant Potential in an Intense Electromagnetic Field

In previous work, we discussed the numerical implementation of an exact solution for the scattering of electrons by a piecewise constant potential in the presence of a strong electromagnetic field. An intimate connection with the physics of multiphoton ionization was demonstrated. Here, that work is extended to include the accurate evaluation of power radiated at harmonics of the applied field frequency. The results do not use the dipole approximation, are nonperturbative and are manifestly gauge invariant. We observe the production of multiple order harmonics whose intensities track well with the intensities of the above threshold ionization peaks. No plateau in radiated intensity as a function of harmonic number is seen. Suppression of even harmonics in a symmetric potential is seen if the incoming state is strictly single parity (even or odd), but is found to be exceedingly sensitive to admixture of a small component of the opposite parity. In scattering from a well with two bound states, we observe a striking minimum in the third harmonic intensity for incident electron energy close to resonance.

Authors: R. A. Sacks and A. Szoke²²

2.3.11. Electron Scattering Assisted by an Intense Electromagnetic Field—Exact Solution of a Simplified Model

Electron scattering by a layered (stepwise constant) potential is calculated in the presence of strong electromagnetic radiation. The numerical solution converges to high accuracy. We study threshold phenomena, resonant emission and absorption of many photons, multiphoton ionization, and the interaction of decaying “bound” states with each other and with the continuum. We demonstrate the close connection between electron scattering and multiphoton ionization. A theoretical treatment in terms of dressed states accounts well for most of our results.

Authors: R. A. Sacks and A. Szoke²³

2.4. Progress in Plasma Physics

2.4.1. Overview of Plasma Physics Research in 1989

Laser plasma coupling is one of the high-leverage issues for laser fusion. Economic constraints on laser energy place a premium on stressing the interaction physics to allow a wide range of target design options. Our strategy is to develop an experimentally-tested understanding of the important coupling processes. This understanding is needed to guide the choice of future drivers and to allow maximum flexibility in the irradiation characteristics. In 1989 substantial progress was made in our ongoing program to characterize and control the important mechanisms for laser plasma coupling. This progress as well as our collaborative plasma research with IGPP and UCLA has been detailed in many articles.²⁴⁻⁵³

A brief overview is given here. Our contributions to heavy ion research are reported in another section of this report.

Collaborations on Laser Plasma Experiments. The close and productive collaboration between X and Y Divisions on the design and interpretation of laser experiments to characterize the major coupling processes continued. There were approximately 12 joint publications³⁵⁻³⁶ with experimentalists in 1989. In experiments with Nova, Stimulated Raman Scattering (SRS) was examined in the largest plasmas to date.⁴³ The targets were exploding foils with a density scale-length of several millimeters. Several of the papers further detailed the suppression³⁸⁻⁴⁰ of SRS by collisional and Landau damping. New results on Stimulated Brillouin Scattering (SBS) were also reported.³⁷⁻⁴⁵ This scattering continues to be rather small for short wavelength laser light. Current data³⁷ emphasize the importance of seeding of SBS by ion waves generated nonlinearly at 1/4 the critical density by the $2\omega_{pe}$ instability. Filamentation continued to be examined in experiments. A correlation between filamentation and second harmonic emission was directly demonstrated for the first time.⁴⁴

Instability Reduction by Temporal and Spatial Incoherence. There has been increasing interest in the use of induced spatial incoherence and related schemes involving combinations of increased laser bandwidth and phase plates or echelons to provide laser beams of greater time-

averaged smoothness. Such schemes are necessary to provide adequate illumination uniformity on directly driven targets, and may have additional benefits from the suppression or reduction of parametric instabilities in the underdense plasma. A comprehensive analysis³² has been made of the most direct effect of increased temporal and spatial bandwidth on plasma instabilities, namely the spoiling of the phase-matching conditions for the resonant instabilities such as SRS and SBS. The effect of spatial incoherence in reducing nonresonant instabilities such as filamentation has also been considered.

Some estimates will illustrate our current understanding. Bandwidth reduces the growth rate provided $\Delta\omega > \gamma$. To illustrate the numbers, consider 0.35 μm light with an intensity of 10^{15} W/cm^2 in a plasma with a temperature of 2.5 keV. For the Raman and $2\omega_{pe}$ instabilities near $n_{cr}/4$, $\gamma/\omega_0 \sim 2 \times 10^{-3}$. For SBS near $0.1 n_{cr}$, $\gamma/\omega_0 \sim 5 \times 10^{-4}$. Hence significant bandwidth ($>> 0.1\% \omega_0$) would be needed to strongly affect the instabilities, unless they were already near threshold. We estimate that a bandwidth of 3 to 4% would raise the collisional threshold intensity for Raman backscatter of 0.35 μm light in a CH plasma to about 10^{15} W/cm^2 . It should be noted that bandwidth is less effective for controlling Raman forward scattering, since the incident and scattered waves then travel together. On the other hand, this weak instability is suppressed by a gentle variation in the plasma density.

Nonresonant instabilities such as ponderomotive filamentation are not directly affected by bandwidth. The growth of this filamentation is reduced by spatial incoherence. Picture the beam as composed of overlapping beamlets. Beamlets whose wave vectors are at too large an angle not only do not contribute to but oppose ponderomotive filamentation of the beam. The angular spread needed for suppression has been calculated to be

$$(\Delta\theta)^2 = \left(\frac{\Delta k}{k_0} \right)^2 \geq \frac{1}{2} \frac{n}{n_{cr}} \frac{I_{16} \lambda_{\mu\text{m}}^2}{\theta_{\text{keV}}}$$

where I_{16} is the intensity in units of 10^{16} W/cm^2 , $\lambda_{\mu\text{m}}$ its wavelength in μm , and θ_{keV} the electron temperature in keV. Physically, this equation corresponds to the condition that the spatial coherence length ($1/k_0 \Delta\theta$) be less than the most unstable wavelength. Bandwidth helps indirectly

by smoothing the interference pattern, i.e., by moving the hot spots around on a time scale $1/\Delta\omega < 1/\gamma$. This reduction of filamentation has been found in simulations at NRL but has not yet been directly observed in experiments.

3-D Calculations of Ponderomotive

Filamentation. The ponderomotive filamentation of a laser beam with finite amplitude intensity modulations has been considered. A 3-D model for light wave propagation including quasi-static plasma response with a saturable nonlinearity is solved numerically. The calculations treat a beam with the imposed intensity modulation used in recent experiments by Young et al.

Figures 2-8 and 2-9 show some illustrative calculations. A strong steepening of the initial sinusoidal intensity modulation is found in both the calculations (Fig. 2-8) and the experiments. Further into the plasma, the intensity modulations collapse into intense narrow filaments. As shown in Fig. 2-9, each filament is only several wavelengths wide. The intensity and width of the filaments are calculated to oscillate in space, with a peak intensity sufficient to significantly deplete the plasma density from the filament. This collapse into intense filaments is a characteristic feature of the 3-D calculations.

These calculations also indicate that small density modulations in the plasma can induce spatial incoherence and suppress ponderomotive filamentation. Figure 2-10 shows the computed spatial growth of the electric field δE associated with the most unstable filament in a plasma with a uniform density of $0.1 n_{cr}$. With no imposed density modulation, δE grows exponentially as predicted. However, when a small density fluctuation with a wavelength of 0.25 of the most unstable wavelength is imposed, the growth is quenched. Physically, the incident wave is broadened in angle by its interaction with the short wavelength density fluctuation before filamentation can develop.

New Targets for Large-Plasma Experiments on Nova. Fairly uniform plasma targets ≈ 0.5 cm in size at 3–5% critical density ($n_e \approx 5 \times 10^{20} \text{ cm}^{-3}$) and $T_e \approx 1$ keV would allow SRS and filamentation experiments under conditions approximating LMF targets. Several concepts have been investigated to create such targets.

One concept uses gas-filled balloons, 0.5 cm in diameter with a 1- μm -CH skin and filled with a large-molecule gas, such as butane at ≈ 0.5 atm. One- and two-dimensional LASNEX simulations have been done with 2 ns laser pulses at 3ω .

Simulations corresponding to the two-beam and 10-beam Nova configurations with up to 4 kJ per beam were done. The laser beams first explode the skin, then burn their way through the gas, ionizing it and bringing it up to an electron temperature of 1 to 2 keV. The initial gas density is chosen such that, after ionization, the electron density has the desired value of $5 \times 10^{20} \text{ cm}^{-3}$. The gas motion during the 2 ns pulse is limited to a compression by a factor ≈ 2 near the interface between gas and exploding skin, and the interior maintains a uniform electron density. Using the two-beam configuration, only a cylindrical core of gas, 0.2 cm in diameter, is ionized and heated. With the 10-beam configuration, the entire balloon (a sphere 0.5 cm in diameter) becomes a plasma at the desired electron density and temperature.

Because of possible fabrication difficulties with 1- μm -skins, additional simulations were done to study balloons with a 2- μm -skin, instead of 1 μm . In this case, the laser energy is not sufficient to explode the entire skin and heat the gas. An alternate lighting scheme was used for this case by irradiating the two opposite 45° polar caps of the balloons with the 10-beam configuration. Only the polar caps are now exploded, requiring much less laser energy. The Nova beams, which are inclined up to 57° from axis, enter through the resulting openings, heat the entire gas to the desired conditions and begin to blow off the interior of the remaining skin. Thus, a plasma sphere 0.5 cm in diameter is still created, but not in a completely open geometry, since a "belt" of cold skin remains intact. A significant fraction of the total solid angle becomes unavailable for diagnostics in this case.

Another concept in which a solid CH disk, 0.2 cm in diameter and 10 μm thick, is exploded with a 200 J prepulse has also been studied. Such a prepulse could come from one of the 10 beams of Nova. The expansion follows approximately the self-similar solution, creating an approximately uniform gas puff, 0.5 cm long and 0.2 cm in diameter. After expansion, this puff would be ionized and heated by a set of 8 beams, and the tenth beam could be used as probe for plasma experiments. This method would create a plasma having a desirable density distribution, but LASNEX simulations show that a delay of 70 ns is needed between the one-beam prepulse and the eight-beam main pulse, a challenging requirement to achieve on Nova.

50-00-0390-1597

Figure 2-8. The intensity modulation at a distance of 1.4 growth lengths into the plasma.

50-00-0390-1596

Figure 2-9. The laser light intensity profile at a distance of 2 growth lengths into the plasma.

Figure 2-10. The calculated growth in space of the amplitude δE of the most unstable filament with and without an imposed density.

Two other concepts, making use of a gas jet and exploding multiple foils, will also be investigated. Potential difficulties due to filamentation of the plasma during heating are also under investigation.

P³—A Plasma Post-Processor Code. A post-processor code is being developed and used to evaluate the role of laser-driven instabilities in target designs. This code uses LASNEX output files to estimate laser light scattering, absorption, and filamentation. Instability thresholds due to collisions, Landau damping, and gradients are included. Above these intensity thresholds, the Raman and Brillouin-induced reflectivities are crudely estimated by fits to short-time particle simulations with an enhanced noise level. This approach is inadequate for quantitative predictions, but is so far found to provide reasonable estimates for the onset of sizable ($\geq 10\%$) reflectivity. The maximum gain due to filamentation is also estimated, and various beneficial effects such as bandwidth and beam divergence are very crudely modeled. This code has compared quite well with measurements of the location and timing of SRS in KMS experiments. Valuable feedback has also been provided to LMF designs.

Collaborative Work with IGPP and UC. In the original work on collisionless, perpendicular shocks, the momentum distribution downstream from the front is isotropic perpendicular to the magnetic field, with no spread parallel to B . We have recently considered the instability of this distribution using linear theory and PIC simulations. For weakly magnetized flows, the Weibel instability dominates, with growth rate $\sim$ the relativistic transverse plasma frequency, and isotropization is rapid. For parameters of interest in our application, the relativistic gyro frequency ranges from $\ll$ the Weibel growth rate, to much larger at low particle energy; this spread alters the instability regime.

We have considered magnetic field reconnection^{50,51} in astrophysical plasmas, in the regime where collisionless rather than resistive processes are dominant. We argue that many astrophysical plasmas are of low enough density and high enough temperature that they are in this collisionless regime. A "Direct Implicit" particle-in-cell plasma simulation code is used to study magnetic reconnection in neutral-sheet geometry and in a sheared magnetic field. In general we observe magnetic reconnection on the Alfvén-wave transit time scale for our system, as seen in

solar flares and in planetary magnetotails. With the ability to use a realistic ratio of the ion to electron mass, we isolated several distinct phases in the reconnection physics: early exponential growth of electron tearing modes in agreement with linear theory, island coalescence, oscillation of island amplitude and trapped flux due to underdamped electron flow, and continued large-amplitude island growth. We found this growth to be linear in time, with slope proportional to the ion acoustic velocity. Based on our computational results and theoretical analysis, we suggest that for forced astrophysical reconnection, rapid collisionless reconnection in sharp spatial gradients may provide important large-amplitude seeds which are already sufficiently nonlinear to initiate rapid macroscopic magnetohydrodynamic reconnection.

A loss-cone plasma inside a theta pinch coil has been suggested by R. A. Post as a possible device for focusing and accelerating ions. Similar in many ways to a Pulsalac, this arrangement is of interest to HIF for ion control in the low energy end of the accelerator. The electrostatic potential function similar to that predicted by Post emerges as the steady-state in Larson's simulations of the time dependence of this device. Larson has added static B fields to FYMNOS for this work and plans to add inductive effects next using the SDF equations with our new DADI solver.⁵³

Finally a collaboration with UCLA has consisted of computer simulations in support of ionospheric modification experiments at the HIPAS Observatory in Alaska. This collaboration has resulted in two publications this year.^{48,49}

In the experiments, a large-scale density modulation was achieved by electromagnetic waves whose frequency matched the plasma frequency at a height where the ionospheric density profile was flat. The density in this region was dramatically clamped during the morning when it normally increases from solar ionization. Electromagnetic wave propagation in the polar magnetic field geometry and strong resonantly enhanced electrostatic fields over a large region of nearly constant density account for the observation.⁴⁸ Zohar simulations were used to illustrate the conversion into electrostatic waves and the subsequent generation of cavitons, which are highly localized regions of electrostatic field.

In other simulations,⁴⁹ recent predictions by Mjølhus (1989) of a large modulation in the coefficient for conversion into the Z mode were confirmed. The absorption varies between zero

and twice the standard mode coupling coefficient due to interference between waves transmitted to the upward (subsequently reflected) and downward propagating Z modes. The simulations also showed an enhanced absorption at the second harmonic frequency, as well as nonlinear features of the heating by the large, resonantly enhanced electrostatic fields.

Author: William Kruer

2.5. X-Ray Laser Design

2.5.1. Introduction

X-Division continues to be the leading center for laboratory x-ray laser (LXRL) design and analysis. In "Plasma Physics Issues in LXRLs" (Sec. 2.5.2), we explored a wide variety of issues not usually dealt with in this field of research. In particular, we discussed plasma instabilities that may lead to inhomogeneities in the XRL-plasma medium that could interfere with XRL propagation. Of course we are always concerned with the accuracy of our atomic physics input into our predictive codes. This is discussed in Sec. 2.5.3 in the context of recombination XRLs.

We continually search for schemes that could lead to cheap table top XRL systems. A new approach is tested, and its promising results presented, in "Studies of High-Z Exploding Foils Irradiated by Combined Long (2 ns) and Short (10 ps) Pulses of 1ω Light" (Sec. 2.5.4). Another approach to tabletop systems is to use femtosecond lasers to create an ultra-short x-ray pulse to photoionize inner shell electrons and thus invert a lasing medium. In "Scaling Laws for Fsec Laser-Plasma Interaction" (Sec. 2.5.5), we explored the expected efficiency of creating such x-ray pulses.

In the meantime, we are using large lasers (such as the Nova two-beam facility at LLNL and the Phoebus Laser at Limeil) as pumps for LXRLs to study the systematics of such systems. For Ne-like collisional excitation XRLs, we have studied the effects of varying the Z of the target (Sec. 2.5.6) and the wavelength of the pump laser (Sec. 2.5.7) on the performance of such systems. The difficulties in using large pump systems, with 100- μ m-high line foci, for driving recombination XRL systems are detailed in Sec. 2.5.8.

Our most likely approach to attain 45 Å lasing in the laboratory is with Ni like Ta. Such a wavelength is optimal for wet biological imaging

via holography. For contact microscopy, a wavelength inside the carbon K edge is desired, such as 43.1 Å which is obtainable from Ni like W. Both lasers are optimized in a design study presented in Sec. 2.5.9.

For holography we require the XRL source to be coherent. An analytic approach to calculating such coherence in LXRLs is described in Sec. 2.5.10 and planned approaches to achieving good coherence are presented in Sec. 2.5.11.

In XRL studies, details count! The ability to accurately account for line transfer is important in the calculation of gain. Efficient escape probability methods for both planar and cylindrical geometries are presented in Sections 2.5.12 and 2.5.13 respectively.

2.5.2. Plasma Physics Issues in Laboratory X-Ray Lasers

Laboratories around the world are producing plasmas with atomic population inversions that are exhibiting gain at wavelengths approaching the carbon K edge near 44 Å, which is optimal for holography of wet biological samples. Often these demonstrations of gain are accompanied by anomalies that may be due to atomic physics effects, but whose possible causes due to plasma physics effects are explored here. The optical pump laser's pulse duration ranges from the relatively long (> 1 ns) to the ultrashort femtosecond regime, and its incident irradiance varies by many orders of magnitude as well. This broad range of results is analyzed here from a unified point of view, via the use of a suite of large simulation codes. The plasma gain media exhibit many laser-plasma instabilities, including stimulated Raman scattering (used to infer the evolution of an exploding foil x-ray laser target's electron temperature and density), and possibly filamentation, radiative instability, and ion acoustic turbulence, all of whose resultant nonuniformities can seriously affect the x-ray laser beam's propagation and achievable gain-length product. Non-Maxwellian distributions of electrons and nontrivial x-ray laser line shapes may also play a role in explaining some of these anomalies.

2.5.3. Sensitivity of Recombination X-Ray Laser Gain Predictions to Atomic Physics

We describe a number of atomic physics issues important for highly ionized ions with moderate atomic number. Examples relevant to

recombination x-ray lasers in sodium, magnesium, and aluminum ions are given. While these ions differ by at most two in atomic number, there are interesting differences in atomic rates for the various ions. We study lasing in hydrogen-like and lithium-like ionization stages following recombination from fully stripped and helium-like ions. We discuss the appropriateness of shell averaging of the various n levels for these ionization stages. In addition, we show the effect of the number of the Rydberg levels on recombination rates and ionization balance.

We have found that $\Delta n = 0$ ion collisions can be important in connecting some sublevels but they do not restore all sublevels to statistical equilibrium. In hydrogen-like magnesium, ion collisions strongly connect the 2s level to the $2p^{1/2}$ level but do not strongly connect the other $n = 2$ levels to each other. The connection of the $2p^{1/2}$ level to the 2s level increases the $2p^{1/2}$ population and greatly reduces the predicted $3d^{3/2} - 2p^{1/2}$ gain coefficient. We discuss how accurate energy levels are critical in determining ion collision rates.

Authors: J. K. Nash and D. C. Eder⁵⁴

2.5.4 Studies of High-Z Exploding Foils Irradiated by Combined Long (2 ns) and Short (10 ps) Pulses of 1ω Light⁵⁵

High-Z exploding foil targets are used in many soft x-ray laser schemes. The foils are typically irradiated with a long (~ 1 ns) optical laser pulse that burns through the foil to produce a large hot plasma with long density scale lengths suitable as an amplifier. While it is expanding, and before the plasma conditions are suitable for stimulated emission, a large fraction of the heating laser energy is lost through radiation and conduction. We report the results of experiments attempting to increase the efficiency of exploding foil amplifiers through the following procedure. A laser beam ($\sim 10^{12}$ W/cm²) is used to heat the foil sufficiently to expand it to approximately 200 μm . A high intensity ($\sim 10^{15}$ W/cm²) short pulse (10-ps) laser beam is then used to raise the plasma to the desired temperature and ion state. Temporally resolved x-ray spectra from Yb foils are presented.

Authors: L. B. Da Silva, B. J. MacGowan, D. L. Matthews, M. D. Rosen; H. A. Baldis, G. D. Enright, (LLNL); B. LaFontaine, and D. M. Villeneuve, (NRC of Canada)

2.5.5. Scaling Laws for Femtosecond Laser-Plasma Interactions

We present scaling laws for the interaction of femtosecond high power lasers with solid targets. We solve for scale lengths and temperatures of the hot dense plasma in both the classical and flux limited electron thermal conduction regimes, and apply those results to the prediction of radiation output from these targets. The radiative yield and duration are given as functions of both laser and target conditions. Our results compare quite favorably with sophisticated numerical simulations and data.

Authors: M. D. Rosen⁵⁶

2.5.6. Study of Soft X-Ray Amplification in a Laser-Produced Strontium Plasma⁵⁷

Amplification of soft x-rays in Sr ($Z = 38$) plasmas produced by second harmonic laser irradiation of exploding foil targets is studied. Gains of 4.4 cm⁻¹ and 4.0 cm⁻¹, respectively, have been measured for the 164.1 and 166.5 Å $J = 2-1$ transitions in Ne-like Sr at laser intensities of 1.3×10^{14} W/cm². Numerous x-ray and soft x-ray transitions were identified in order to demonstrate typical spectra from x-ray laser plasmas and infer information regarding the plasma ionization balance. The $J = 0-1$ transition at 159.8 Å analogous to the well known Se $J = 0-1$ line at 182.4 Å was not definitively observed due to a wavelength overlap with a Na-like Sr line. Two other $J = 0-1$ transitions of interest at 84.9 and 133.0 Å were not observed. Reduced gain on the $J = 2-1$ lasing lines was seen for several shots taken at a lower pump laser intensity of 7.0×10^{13} W/cm². Possible reasons for this are considered. The results of this work imply the feasibility of a high output power cavity geometry Ne-like Sr x-ray laser operating at 164.1 and 166.5 Å.

Authors: C. J. Keane, D. L. Matthews, M. D. Rosen, T. W. Phillips, B. J. MacGowan, B. L. Whitten; M. Louis-Jacquet, J. L. Bourgade, A. DeCoster, S. Jacquemot, D. Naccache and G. Thiell, (Centre d'Etudes de Limeil-Valenton, St. Georges, France)

2.5.7. Collisional Excitation X-Ray Laser Experiments in Plasmas Produced by 0.53-mm and 0.35-mm Laser Light⁵⁸

Recent Ne- and Ni-like x-ray laser experiments carried out at the Centre d'Etudes de

Limeil-Valenton (CEL-V) are reviewed. A variety of experiments in Ne-like x-ray lasers were performed; here we discuss measurements of soft x-ray amplification in Ge ($Z = 32$) and Sr ($Z = 38$) plasmas. In Ge plasmas produced by 0.53- μm laser light at an irradiance of $6.0 \times 10^{13} \text{ W/cm}^2$, gains between 2.2–2.5 cm^{-1} on the 232.2 and 236.2 $\text{\AA}$ $J = 2-1$ lines and a gain of 1.0 cm^{-1} on the 196.1 $\text{\AA}$ $J = 0-1$ line were measured. In addition, gains of 4.4 cm^{-1} and 4.0 cm^{-1} have been demonstrated on the $J = 2-1$ transitions at 164.1 and 166.5 $\text{\AA}$ in Ne-like Sr at laser intensities of $1.3 \times 10^{14} \text{ W/cm}^2$. The effects of pumping the Ne-like Se x-ray laser with 0.35- μm laser light have also been investigated; the Se lasing spectra is similar to that obtained with 0.53- μm light. Experiments have also been carried out to optimize the gain of the 50.3 $\text{\AA}$ Ni-like Yb ($Z = 70$) $J = 0-1$ line. For Yb, no significant increase in gain over that previously reported was seen, but the time history of the Ni-like Yb x-ray laser was measured for the first time. Finally, attempts to extrapolate the Ni-like results to shorter wavelength were made using Ta ($Z = 73$), W ($Z = 74$), and Re ($Z = 75$). No definitive observation of the Ni-like $J = 0-1$ lasing lines was made in these experiments.

Authors: C. J. Keane, B. J. MacGowan, D. L. Matthews, R. A. London, S. Maxon, M. D. Rosen; J. L. Bourgade, A. DeCoster, S. Jacquemot, M. Louis-Jacquet, and D. Naccache (Centre d'Etudes de Limeil-Valenton, St. Georges, France)

2.5.8. Hydrogen-like Magnesium X-Ray Laser Design

We present design calculations for a H-like Mg recombination x-ray laser driven by a 20-ps optical laser. Results for narrow free-standing strip targets and the more realizable situation of a strip of lasant material supported on a thin plastic backing are compared. General principles of target design are discussed, as well as specific targets suitable for the Nova two-beam laser. The modeling of composite targets is improved by using distinct grids in different regions and allowing these regions to slide past each other. We show that the major effect of the plastic backing is to partly contain the outer portions of the expanding strip. During the expansion of a free-standing strip, the outer portions have the largest fractional expansion and therefore are the source of the majority of the expansion cooling. Since the plasma is nearly isothermal due to very fast electron conduction, this cooling is

transferred throughout the plasma. For the interior region, where the gain is the largest, the result is a decrease in the cooling rate for the strips backed with plastic with only a small effect on the electron density. The plastic backing has a significant effect on the calculated gain for strips that are both thin and narrow. These effects are shown in Fig. 2-11 for a 25- μm -high strip of magnesium fluoride with a thickness of 2000 $\text{\AA}$ on a 500 $\text{\AA}$ plastic backing. For this comparison, the plastic behind the strip is not included, giving horizontal and vertical symmetry so that only one quadrant of the cross section needs to be modeled. The initial mesh is shown in frame (a) and the mesh 100 ps after the peak of the 20 ps driving pulse is shown in frame (b) where the tamping by the plastic is clearly seen. The mesh for a free-standing 25- μm , 2000- $\text{\AA}$ strip is shown in frame (c) also 100 ps after the peak of driving pulse. The large expansion of the initially small outer zones is clearly seen in this case.

The calculated total energy requirements are 160 to 320 J per centimeter of lasant material depending on the thickness of the strip and using a 100- μm -line focus. If a 25- μm -line focus is used, the energy requirements would drop by a factor of four. Our calculations predict that the gain for pure magnesium and magnesium fluoride strips should be comparable assuming that our model treats photon trapping appropriately. If trapping is not important, the gain for pure magnesium should be significantly higher. An increase in predicted gain of comparable value for pure magnesium and magnesium fluoride is found if the cooling rate is artificially increased. There is not a widely accepted reason for why experimentally measured gains in recombination systems tend to be larger than predicted by modeling. Experiments using both magnesium and magnesium fluoride strips might aid in understanding this difference.

Author: D. C. Eder⁵⁹

2.5.9. New Calculations for Ni-Like Soft X-Ray Lasers: Optimization for W (43.1 $\text{\AA}$)⁶⁰

In order to test our 2-D hydrodynamic calculations, holography experiments were carried out at KMS in which a thin foil target containing 90 $\mu\text{g/cm}^2$ Re was irradiated from one side by a 0.53- μm , 700-ps-FWHM laser pulse with spot size 150- μm -diam. The total intensity was $1 \times 10^{15} \text{ W/cm}^2$. Interferograms were taken at four probe times. Each interferogram is Abel

Figure 2-11 (a) The initial mesh for one quadrant of a 25- μm -high and 2000 Å thick magnesium fluoride strip with 500 Å of plastic above and below the strip. The last column of zones on the far right have very low density and do not affect the expansion. We give positions in an xy coordinated system with origin being at the center of the strip. **(b)** The same mesh as (a) but 100 ps after the peak of the 20 ps driving pulse. The vertical scale is the same as (a) but the horizontal scale is increased by a factor of 10,000 and matches the vertical scale. **(c)** The mesh for a free-standing strip at same time as in (b).

inverted to produce an electron density profile. The predicted electron density profile taken 100 ps before the peak of the laser pulse matches the measured profile as shown in Fig. 2-12. The Z axis is along the direction of the incoming laser.

We approach the problem of designing a Ni-like W experiment by first carrying out an optimization procedure using a series of 1-D problem. One dimensional LASNEX, XRASER, and SPECTRE calculations were done for a 90- $\mu\text{g}/\text{cm}^2$ W target and 2ω irradiation of several intensities. The space and time averaged gain is plotted versus I in Fig. 2-13. The dominant effect on the rising part of the curve is ionization to the Ni-like state while the lower gain at higher intensities results from the rapid decrease in density due to expansion. The intensity scale in Fig. 2-13 should be multiplied by a factor of 1.5 due to 2-D expansion so that the peak in α for 2-D calculations occurs at $I \sim 2 \times 10^{14} \text{ W}/\text{cm}^2$. Previous experiments were done at intensities of $1\text{--}1.4 \times 10^{14}$ and gave $\alpha \sim 1 \text{ cm}^{-1}$.

Based on the optimization results, design calculations showed that a 90 $\mu\text{g}/\text{cm}^2$ W target irradiated in double-sided geometry with 2ω , 550 ps FWHM trapezoidal pulse with a total intensity $3.0 \times 10^{14} \text{ W}/\text{cm}^2$ should give a space and time averaged gain $\sim 5 \text{ cm}^{-1}$. Peak gain is generated for $n_e \sim 2.5 \times 10^{21} \text{ cm}^{-3}$ at an electron temperature of 860 eV.

Recent experiments were carried out by B. MacGowan at the Nova 2-beam facility using intensities of $2.5\text{--}3.0 \times 10^{14} \text{ W}/\text{cm}^2$, yielding $\alpha = 2\text{--}2.5 \text{ cm}^{-1}$. Eight gain lengths were obtained for Ta. Similar results have been obtained for W.

Author: Stephen Maxon

2.5.10. Modal Analysis of X-Ray Laser Coherence⁶¹

The property of transverse coherence is important for eventual applications of x-ray lasers (XRLs), such as holography. Past work on x-ray laser propagation theory does not address the question of coherence. In this paper we have addressed the coherence question, presenting a steady state, wave optics model for mirrorless lasers, such as XRLs. The model is based on a modal decomposition of the paraxial wave equation. It describes the development of the radiation from spontaneous emission into all order transverse modes (rather than just the lowest mode as considered by others previously) for arbitrary transverse profiles of gain and

electron density. Using semi-analytical calculations and derived scaling laws, we have identified for the first time the basic mechanisms which influence the degree of coherence, namely the effects of gain guiding and refractive defocusing. Because the transverse paraxial equation is non-Hermitian, the modes have "excess spontaneous emission," particularly in the higher order modes. We have identified the importance of controlling this excess spontaneous emission in order to achieve a high degree of coherence. We have derived scaling laws for the number of guided modes and their gains for several plasma profiles relevant to x-ray lasers, in terms of three parameters: the gain-dependent Fresnel number, the strength of refraction relative to gain, and the amplification length. We have identified the regions of parameter space where the excess spontaneous emission can compromise coherence, and those in which it can be controlled to achieve high coherence. We have presented detailed results for the output intensity and the coherence function for several examples of plasma profiles (see Fig. 2-14). Finally, we have suggested designs for highly coherent x-ray lasers and an experimental effort to achieve this goal. There is a large experimental effort at many laboratories including Livermore which will be influenced by this work.

Authors: R. A. London, M. Strauss (NRC-N, Israel), and M. D. Rosen

2.5.11. The Development of Coherent X-Ray Lasers for X-Ray Holography⁶²

We describe several of our research paths to develop an x-ray laser capable of producing the flux required to obtain x-ray holograms of biological structures in live cells. Current predictions are that we will need at least 20 μJ of spatially-coherent radiation at a wavelength longer than 43.7 Å in a pulse length of 50 ps and focused to a beam of 10 μm diam at the sample. We provide a detailed description of how we hope to meet these requirements. We describe proposals to optimize our current Ni-like Yb laser that operates at 5.02 nm and to create a new Ni-like Ta laser that operates at 4.48 nm, near the optimum wavelength for holographic imaging of biological objects. Several schemes to achieve spatial coherence are described. In addition, some comments are given on future trends in x-ray

Figure 2-12. Density profile of an exploding Ni-like Re foil. Solid line is LASNEX simulation, squares are data points.

Figure 2-13. Space and time averaged gain vs. incident irradiance.

Figure 2-14. Intensity (dashed lines) and coherence function (solid lines) versus transverse position at the exit of x-ray lasers with square gain and refraction profiles. The gain-length product is 20. The effective Fresnel number is $F_e \equiv kga^2$ where k is the XRL wave number, g is the gain and a is the width of the profile. The refraction strength parameter is $\eta \equiv \omega_p^2/(kc^2g)$, where ω_p is the plasma frequency.

laser development, namely towards more efficient "table top" systems.

Authors: D. L. Matthews, J. Trebes, R. London and M. Rosen

2.5.12. Application of Escape Probability to Line Transfer in Laser-Produced Plasmas⁶³

The escape probability method has been widely applied in the calculation of optically thick line transfer because of its simplicity. This method is accurate when line reabsorption is relatively local in space. In this paper we demonstrate the usefulness and accuracy of the escape probability method in modeling laser-produced plasmas in plane-parallel geometry. We have investigated the effect of trapping on charge-state distribution, the ion level populations, and on the gain for x-ray lasers. Because of the large velocity gradients typical of laser produced plasmas, the Doppler shifting of the line profile is an important process in line transfer. In certain cases it represents the dominant effect in line transfer, and because it is a local effect, the escape probability method is expected to be quite accurate.

We have developed a simple and accurate algorithm for calculating the escape probability for plasmas with monotonic velocity profiles, that spans the limits from low to high velocities. We have incorporated this algorithm into DCA, LASNEX and XRASEF. We have illustrated the utility of the escape probability method by considering two examples: a selenium exploding foil x-ray laser and an ablating aluminum slab, proposed as a pump for a photo-resonant x-ray laser scheme. In general, the results of the escape probability calculations agree well with results using the detailed multifrequency, multiangle numerical line transfer package in XRASER (see Fig. 2-15). Only in the high (super-critical) density region of the ablating slab is there a significant difference, due to the low velocities and associated nonlocal radiation transfer at that location. Our calculations demonstrate that the escape probability method is accurate for a wide range of conditions in laser plasmas. Owing to its simplicity and calculational speed, it is an extremely valuable tool.

Authors: Y. T. Lee, R. A. London, G. B. Zimmerman (LLNL); P. L. Hagelstein (MIT)

2.5.13. The Calculation of Line Transfer in Expanding Media⁶⁴

An efficient method for calculating 1-D line transfer that allows for bulk Doppler motion is described. The line transfer is done using a ray-based algorithm derived from the formal integral solution to the transfer equation. Consistency between atomic populations and line intensities is achieved with a complete linearization procedure, that can handle an arbitrary number of overlapping and interacting lines. The main advantage of the ray-based integral formalism is its generality. Geometry (and even dimensionality) enters only implicitly in the choice of rays. The main disadvantage is the large number of rays which can be required to achieve good angular coverage in nonplanar geometries. The computational time required by the line transfer calculation scales almost directly as the number of rays, so it is much more expensive to do calculations in cylindrical geometry than in spherical geometry, which is much more expensive than in planar geometry.

We compared an alternative method of calculating photon escape for cylindrical geometry in situations where there are steep velocity gradients in the medium. This method of using local escape probabilities can give very accurate results at a fraction of the cost of doing nonlocal line transfer calculations. A complication that can affect the accuracy of the escape probability treatment is the presence of neighboring lines that can be coupled via bulk Doppler motions. We give an example of using these methods in the modeling of an x-ray laser in cylindrical geometry. We calculate the gain for the $3d_{5/2} - 2p_{3/2}$ and $3d_{3/2} - 2p_{1/2}$ transitions in H-like fluorine. The lower levels of the lasing transitions are depopulated by rapid resonance transitions to the ground state. The $2p_{3/2} - 1s$ line is 0.3 eV higher in energy and has twice the strength of the $2p_{1/2} - 1s$ line if the $n = 2$ levels are populated statistically. We use both line transfer and escape probabilities to determine if photons emitted from these transitions escape from the plasma. We first show the results of the comparison for the $3d_{5/2} - 2p_{3/2}$ line in Fig. 2-16, where the gain is plotted as a function of radius at the time of maximum gain. The peak gain for the escape probability calculation is 8% higher than the value obtained from doing full line transfer. The peak gain at this time if one assumes no trapping is 4 per cm or 150% higher.

Figure 2-15. The accuracy of the escape probability method is illustrated in this plot of the calculated gain of the 206 Å line in Ne-like Se. The calculation is for a standard selenium exploding foil laser and the gains are shown for the central zone of the plasma.

Figure 2-16. The calculated gain coefficient for the $3d_{5/2} - 2p_{3/2}$ line in H-like fluorine using escape probabilities and line transfer to determine photon trapping. The gain is plotted as a function of radius at the time of maximum gain.

The predicted maximum gain for the $3d_{3/2} - 2p_{1/2}$ line using escape probabilities is 1.5 per cm while the line transfer calculation gives a value of 0.6 per cm. The reason for the large difference in this case is due to the bulk motion of the plasma during which photons emitted from one line can be shifted in energy and can be absorbed by a neighboring line at a different portion of the plasma. Since the velocity of the plasma is monotonically increasing as a function of radius, a frame co-moving with the plasma will see all other portions of the plasma moving away. This means that photons emitted by the stronger $2p_{3/2} - 1s$ line will be red-shifted into the $2p_{1/2} - 1s$ line and affect the $3d_{3/2} - 2p_{1/2}$ gain but the converse does not happen. The conclusion is that neighboring lines are important only if they have higher energy and the energy difference corresponds to the red shift due to the bulk Doppler motions.

Authors: D. C. Eder and H. A. Scott

2.6. Heavy Ion Fusion Research

2.6.1. Introduction

The articles in this section describe research on heavy-ion particle accelerators as drivers for inertial fusion power production. This approach has received increasing attention over the past year, both nationally and at LLNL.^{65,66} The LLNL HIF work has been presented before two major review panels: a committee of the National Academy of Sciences and the Department of Energy Fusion Policy Advisory Committee. At this writing interim reports are available, and we are encouraged by the committees' findings.

Heavy ion fusion research at LLNL has recently enjoyed active participation by personnel of the Laboratory's beam research effort, now a part of the Laser Program. The collaboration is funded in part by a Director's Initiative grant.

The articles following this introduction describe our work in a number of areas:

- We continued our drift-compression and final focus studies. Recent developments include techniques for focusing a variable-current beam and for correcting third-order geometric aberrations.
- Our work on low-density chamber propagation using a new fully kinetic model led to more optimistic estimates of the fraction of the beam that can be expected to hit the target.

- Our source simulations, using both discrete-particle and new smooth-flow methods, led to an understanding of the plasma-grid interactions.

- We completed a rather exhaustive study of possible beam-plasma interactions in the corona surrounding the HIF target, including use of realistic mass ratios and a realistic density gradient. We have satisfied ourselves that no deleterious effects will be present.

- A new multidimensional discrete-particle simulation code WARP, capable of both 3-D and axisymmetric simulations, has been designed and constructed for HIF accelerator studies.

- We have devoted considerable effort to understanding the properties of recirculating induction accelerators, which may represent a path to a more economical driver.

Besides the work described in the articles, other areas have seen significant progress:

- We have further optimized indirectly driven targets using two-sided illumination. New target gain curves have been derived; these have a higher associated confidence level but are essentially as favorable as the previous curves.

- We are developing the ion beam physics and target simulation capabilities of LASNEX,⁶⁷ and of a variety of smaller codes for heavy ion target studies. We have completely rewritten the ion-beam ray trace package in LASNEX; the new package is faster by a factor of 2.5, and should be a real help on 2-D hydrodynamics-3-D ion beam tracing problems.

- We have developed a new way⁶⁸ of combining three (or any odd number) of linacs into a recirculating accelerator. The new arrangement combines these sections in such a way that the beam goes through each straight linac in alternate directions. This cuts the frequency for inductive core operation in half because at the end of an accelerating pulse, the core is already "reset" for the next pulse in the opposite direction. This may be useful for HIF accelerators if the reduction of inductive core efficiency with frequency becomes an issue.

- We have continued our switch development efforts and recently took delivery of a solid-state (FET) switch. We hope that such elements will prove to be useful for repetitive pulsers and induction modules with favorable impedance properties. The latter may be important to the longitudinal stability of the beam.

Three U.C. Davis Department of Applied Sciences Ph.D. students have continued their work with the HIF group. Debra Callahan has built an r,z physics subset into the WARP code; it shares the particle loading algorithm and many diagnostics with the 3-D branch, but is faster and smaller when 2-D is applicable. It will be applied to problems of longitudinal stability and thermal equilibration. David Grote has been involved in the development of WARP from the beginning. He has built particle loading schemes, diagnostics, and field solution methods for more realistic geometries. David Larson has been providing computational support to Richard Post on the Mirrortron project.⁶⁹ This is a high-current heavy ion accelerator concept employing a quasi-neutral plasma for acceleration, focusing, and space charge neutralization. He has also been working to develop improved numerical solution techniques for strongly coupled 2-D partial differential equations.⁷⁰

Author: A. Friedman

2.6.2. Drift Compression and Final Focus Studies

During the past year we have made progress on five areas in the physics of heavy ion fusion accelerators.

1. We completed research on and writing of the report on longitudinal beam compression.^{71,72} This paper has been accepted for publication in Particle Accelerators.

2. The first-order design of focusing systems for magnets with linear fields is a classical problem. However, for high-current beams, even some basic problems of first-order design have not been explored. We have discovered that, for a parallel-to-parallel (beam envelope starts with zero slope and then being focused to a spot with zero slope) focusing system, the initial waist radius of the beam has a profound effect on the final spot size.⁷³ For a large initial waist radius of 3.4 cm, we are able to focus more than 95% of the beam on a 3 mm focal spot. However, if the waist is reduced to about 1.3 mm (with the emittance held constant), only about 80% of the beam will hit the spot. In both cases, the initial profiles are flat and the distribution is semi-Gaussian.

3. We identified three necessary conditions for the focusing of a variable-current beam using constant-current quadrupoles:⁷⁵

- a. The upstream beam manipulation must be slow enough that every slice of the beam is matched.

- b. The beam must be space charge-dominated, at least over most of its extent, so that the density profile is uniform when the beam exits from the FODO (focusing-drift-defocusing-drift) transport system.

- c. The velocity tilt must be completely removed before the variable-current beam enters the final focusing system.

We showed that these conditions can be achieved in particle accelerators. However, the conditions are necessary but not sufficient. We have identified a fourth condition: the pressure force from the emittance at the spot should be roughly equal to the space charge force (also at the spot). Taken together, these four conditions are necessary and sufficient for effectively focusing variable-current beams using constant-current quadrupoles.

4. Just before leaving the FODO system, the beam is highly compressed longitudinally. It is therefore extremely tune-depressed, with $\sigma/\sigma_0 < 0.1$ (σ/σ_0 is the single particle phase advance per period with/without space charge). One might therefore expect that instability (of the type observed in a periodic lattice) would occur and lead to emittance growth. Such instability does not, in fact, occur. This is best understood by noting that the motion is time-reversible; in the time-reversed sense, the rarefaction waves will reach the beam center very rapidly and cause the line-charge density to fall. Thus, the high line-charge density regime is not sustained long enough for instability to grow.⁷⁴

5. We have derived a systematic method for determining the location and strength of octupoles needed to remove the third order geometric aberrations that can arise in the final focus of a space charge-dominated beam.⁷⁴ By including the octupole terms in the envelope code TRACE,⁷⁵ we demonstrated that the strengths and locations of the octupoles can be determined within a few percent by solving a three-by-three matrix equation.⁷⁶ The exact strengths and locations are then obtained via a small number of iterations. These results are now being checked using the particle-in-cell code SHIFT-XY.⁷⁶ We expect this work to be completed and submitted for publication within a few months.

We find the third-order geometric aberrations to be very important for the final focusing systems of both scaled experiments and full-sized drivers. For a scaled experiment such as might be done using the SBTE (single-beam transport

experiment) apparatus, we find that these aberrations can lead to as much as 70% of the beam falling outside the desired spot. For a driver lens focusing 3 kA of beam current, third order aberration causes about 20% beam loss.

Authors: D. D.-M. Ho, I. Haber (NRL), K. R. Crandall (AccSys Tech., Inc.), and S. T. Brandon

2.6.3. Chamber Propagation

Development continued on the PIC code for chamber propagation.^{77,78} A kinetic (particle-in-cell) model for photo- and co-injected electrons supplements the earlier neutralization model. The plane perpendicular to propagation now has a two-dimensional description, instead of being restricted to azimuthal symmetry. This enables more realistic simulation of drive by a number of beamlets. Other improvements include adaptive variable-step time integration to handle the strongly-varying PIC electron density.

Earlier results (with an axisymmetric version of PIC and pessimistic electron modeling) indicated possible loss of half the beam energy.⁷⁹ With the current code and with less ambitious parameters (seven beamlets, single ionization, 8 m standoff distance, and particle-in-cell electron modeling) less than 10% of the beam energy is lost. This result shows that beam photoionization need not be serious even for propagation in vacuum, though it is a factor to retain in modeling.

Author: A. B. Langdon

2.6.4. HIFI Ion Source Simulations

In the past year, we have completed a study guided by simulation of the plasma grid interaction of the LBL HIF Arc Source.⁷⁹⁻⁸⁶ The principal result has been to characterize the plasma parameters that the arcs must produce and predict, for diagnostic purposes, the source behavior when these parameters deviate from the desired values.

HIF ion sources may require ion arcs confined behind wire meshes in order to provide sufficient ion optical stability with acceptable ion thermal spread. The physics of the plasma-grid interaction and implications for its use as a source for HIF accelerators has been investigated numerically. We have studied these interactions with two models: a time-evolutionary electrostatic PIC code, GYMNOS, with an internal boundary configuration representing the grid wire(s); and a

new method for finding time-independent configurations with flow.

PIC code development has led to provisions for field emission and particle injection. The PIC algorithm easily and correctly describes the evolutionary behavior of any reasonable sort of initial plasma configuration subject to the applied voltages on the grid and the external extraction surface on the accelerator side. Coupled with the appropriate plasma evolution properties on the other side, the system is allowed to evolve to a self-consistent time-asymptotic configuration.

The second approach, computing a time-independent state with flow was motivated by the difficulty in finding proper wire mesh size, spacing, and voltages that provide maximum brightness of the emerging ions. The most challenging problem is to find those configurations for which the ions do not develop large transverse velocities as they "wrap around" the wire mesh. The second approach begins by postulating a desired ion flow without "wrap around" such that no ions pass through a line of symmetry passing through each wire in the confining mesh. The problem then becomes one of finding the other parameters needed to provide this flow. The further assumptions that the ion current has no shear, that the ions have relatively small thermal velocity from the arc, and that the electrons have no flow and constant temperature lead to a representation for both electron and ion densities in terms of the potential. Substitution into Poisson's equation leads to a nonlinear equation for the electrostatic potential that can be solved to provide the necessary parameters needed for consistency with the desired ion flow.

Examples have been calculated using both techniques. The picture that has emerged is a "road map" of ion source performance as a function of plasma arc parameters. We now understand what bounds the regions where acceptable source behavior is found and what to expect from the source if these bounds are exceeded.

Authors: D. W. Hewett and H. L. Rutkowski⁸¹⁻⁸⁶

2.6.5. Ion Beam /Target Corona Interaction

We have completed our survey of the possible mechanisms that might prevent heavy ion beams from fully depositing their energy into ICF targets. We have now submitted a revised manuscript for publication^{79, 80} that gives evidence for our belief that the ion beam will not

be diverted by any plasma instabilities. In the past year we worked to characterize the plasma gradient length more carefully. LASNEX calculations have yielded estimates that made our earlier concern with the longer gradient lengths unnecessary. We were ultimately able to show that, for all reasonable parameters, it is unlikely that filamentation modes will destabilize the ion-electron streaming modes. Furthermore, if they were to become unstable, density gradient stabilization will suppress growth for the expected warm limit of this mode. Code limitations were finally overcome by a large-scale 1-D simulation in GYMNOS, a new RZ PIC code currently being developed for HIF ion source work.

Authors: D. W. Hewett, W. L. Kruer, and R. O. Bangerter

2.6.6. 3-D Code Development

Over the past year, the development of our numerical simulation capabilities has been guided by a need for flexible tools which are applicable to both the mainline driver configuration and to more aggressive attempts at cost reduction based upon advanced concepts. We have developed our discrete-particle simulation code, WARP,⁸¹ tailored to the needs of the HIF program. This code embodies a family of descriptions of a high-current ion beam, ranging from a simple "envelope" model to a fully 3-D particle-in-cell representation. Construction of the code to the point where it is more usable than other 3-D codes was relatively straightforward. Learning to use the code effectively is a larger task, and so our applications have been biased toward those that increase our understanding of the relations between physics and numerics in 3-D simulations.

Our initial applications of WARP have included studies of beams in 3-D that focus on the exchange of thermal energy between transverse and longitudinal motions; the examination of the behavior of finite-length beams undergoing drift-compression; and the search for "quiescent" beam configurations which evolve minimally with time. In contrast with our earlier observations using ARGUS,⁸² no transverse "jets" are observed in our WARP simulations of beams with tapered ends; this is a result of the tapered-end loading and improved numerics, which include a "momentum-conserving" particle advance and adequate axial resolution.

We plan to apply WARP to study the effects of three-dimensionality on the final focus process (where the beam is not truly "long and thin"); to examine limits on the rapidity with which drift compression can be carried out in the presence of strong focusing field errors; and to study space charge dominated beams in bent systems, particularly chromatic effects, images, and the transition from straights to bends. To this end we plan to install a newly developed bent-beam algorithm⁸³ in the code.

The r,z physics subset of WARP will be available for a variety of applications. These could include longitudinal stability studies in 2-D (at first with an axially averaged gap impedance, then with discrete gaps), and/or compression studies to extend our earlier CONDOR results.⁸⁴ The r,z package will be a convenient testbed for extensions to the electrostatic field solver, including the capability of modeling nonsmooth pipes (for example, acceleration gaps). A 1-D axial-only model (assuming incompressibility) is also under development. Aside from their direct utility on problems with simplified geometry, these lower-dimensional models will serve as "spotting scopes" for the 3-D simulation effort, which must choose its runs carefully.

We have developed a "damped" particle mover to filter out noise associated with discrete particle effects, and we plan to test it in WARP. This explicit scheme is a variation of an implicit mover⁸⁵ developed for multiscale simulations.⁸⁶ We have also studied fundamental phase space properties of particle advancement schemes,^{87, 88} and use the results of this work as guidance in the development of algorithms for practical codes such as WARP.

Authors: A. Friedman, D. P. Grote, D. A. Callahan, A. B. Langdon, I. Haber (NRL)

2.6.7. Preliminary Recirculating Linac Studies

Heavy ion induction accelerators hold the promise for meeting all of the driver requirements for an inertial confinement fusion reactor for the production of electrical power. These requirements include efficient beam-target coupling, high repetition rate (~1-10 Hz), high reliability, and long stand-off focusing. Present cost projections for the linear induction accelerator place the cost for such a machine comparable to other drivers, all of which lie in the billion dollar range. A significant reduction in driver cost (by a factor of two or more) would

make the case for heavy ion fusion very compelling. A substantial cost savings over a linear machine may be achieved in an induction accelerator in which a heavy ion beam makes many (≤ 100) passes through one or more racetrack shaped accelerators. We are examining a point design for such an accelerator, consisting of three racetracks, each of which increments the energy by a factor of ten. Since the induction cores used to accelerate the ions are used up to a hundred times, the size of the accelerator and the number of cores used can be much smaller, resulting in the possible large cost reduction. However, the introduction of bends, varying magnetic fields, varying voltage pulse formats, and longer ion residence times in the accelerator leads to new complexities that must be studied before the feasibility of the recirculator can be evaluated.

The size of the recirculator is roughly determined by the bending radius R_b of the ions in the ~ 1 Tesla magnetic fields. The circumference of the machine can be shown to be $\sim 2\pi R_b \sim 2 \text{ km}/q$, for a 10 GeV heavy ion of charge state q . This is in contrast to the linear machine, the length of which is determined by the maximum accelerating gradient achievable without breakdown, leading to a length of $\sim 12 \text{ km}/q$. Further reductions in cost occur since the total induction-core cross section is inversely proportional to the number of times the cores are reused on each ion bunch. Since the costs of the HIFSA design⁸⁹ of the linear machine are predominantly in cores, quadrupoles, structure and pulsers, it is evident that the cost of the recirculator has the potential for being significantly less than for a linear machine.

A relatively conservative straw man design is now being constructed for a recirculating induction linac. In its current manifestation, an ion bunch is accelerated through three racetrack-shaped accelerators, each multiplying the energy of the ion bunch by a factor of ten. By limiting each racetrack or ring to one decade in energy, modest dynamic ranges are required in bending and focusing magnet strengths, and voltage pulse formats.

The first racetrack is about 200 m in circumference. Each successive racetrack is roughly a factor of three larger than the previous, reflecting the factor of three increase in velocity and momentum per racetrack. The ion pulse is injected into the small ring at an energy of about 10 MeV, with a total current of about 10 A, distributed over 4 beams, and pulse duration of

about 40 μs . One hundred circuits and 4 ms later, the ion bunch enters the second ring. Upon injection into the second ring, the total current has been increased to 100 A, pulse duration decreased to 4 μs , and the energy has been raised to 100 MeV. About 50 circuits and 2 ms later, the 1 GeV, 1 kA, 400 ns pulse leaves the medium-sized racetrack and is injected into the largest ring which has a 2 km circumference. About 30 circuits and 1 ms later, the 10 GeV, 4 kA, 100 ns pulse leaves the last ring and enters a drift compression region. Here, as in the linear machine, a head-to-tail velocity tilt on the pulse allows the beam to shorten to a 40 kA (total), 10 ns heavy-ion pulse which then bombards a fusion target.

Each racetrack is divided into three types of region: the bends, the transition regions, and the linac portions. The bend regions (about 60% of the circumference) consist of bending dipole magnets and focusing quadrupole magnets. The bending magnets must (at least in traditional designs) ramp up in field strength as the ions ramp up in momentum. In order for the bends to be achromatic, which is desired because of the imposed head-to-tail energy variation, the focusing magnets must also ramp in time maintaining a constant tune in the bend. (The tune is proportional to the number of radial oscillations an ion undergoes as it travels through the fields of the focusing magnets.) The linac portion (about 25% of the circumference), consists mainly of induction modules and focusing magnets. The focusing magnets in the linac section need not be ramped since the requirement of constant tune does not exist in this section. The transition region (about 15% of the circumference) allows the magnetic field to adiabatically taper from the linac to the bend region minimizing radial oscillations of the beam.

The choice of ion charge state greatly influences the configuration of the machine. In the linear machine of Ref. 71, a charge state of +3 was chosen to reduce the overall length of the machine. The increased space-charge forces however required the final current to be carried by 16 beams rather than 4. In the recirculator, the charge state of +1 results in a machine still significantly smaller than the linear version with charge state +3, but still large enough so that the ramping times for the magnets are not too short. Further, the lower charge state allows easier final focus onto the target. The smaller number of beams does not require the beams to be combined with the associated possibility of increased

emittance growth. Furthermore, the longitudinal resistive instability, which would seriously threaten the HIFSA design appears to be relatively benign in this recirculator design, owing to the smaller number of beams and the lower average impedance. (Since there are no modules in the bends the impedance there is negligible relative to the linac section, yielding a much lower average impedance.)

Other recirculator issues currently under investigation include the emittance growth from the repetitive transport of a space-charge dominated beam around a bend. Centripetal forces lead to nonlinear forces in the equation of motion, which in principle can lead to growth of the emittance. Preliminary estimates indicate this should not be a serious problem. Another area of concern is the vacuum requirement needed to avoid beam loss. The larger residence time in the accelerator and the possibility of beam induced interactions with the walls place severe requirements on the vacuum pumping,

particularly on the smallest ring where the interaction cross sections are largest. Finally, the switching requirements for the voltage pulses applied to the induction cores are quite severe. Since the voltage waveform changes in amplitude, duration, and spacing as the ion pulse is accelerated, programmable switches with rapid rise, fall and recovery times are required. The requirements are most stringent in the largest ring where they may force technological development of the required switches.

We are in the process of investigating these, and identifying and evaluating other critical recirculator issues. In the process, we will iterate on our straw man design and perhaps demonstrate the feasibility of a cheaper way to inertial confinement fusion.

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3

Nova Experiments

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3. Nova Experiments

3.1. Introduction

Nova experiments for 1988 and 1989 were numerous. The four main areas investigated were laser-plasma interactions, hydro-dynamic theory and experiments, laboratory x-ray lasers (LXRLs), and diagnostic developments.

Our laser-plasma interactions research has been aimed at understanding the plasma evolution of laser-heated matter, characterizing the plasma, and investigating laser-plasma instabilities. We have found that a correlation can exist between second-harmonic emission and density perturbations produced by a laser beam interacting with an underdense plasma. We have learned from low-density, long-scalelength plasma experiments that plasma densities and temperatures at late time depend on the initial target geometry, and may be caused by lateral heat transfer effects. Experiments done on the Chroma laser at KMS Fusion, Inc. have shown that harmonic light images can be significantly refracted, especially in the near region, where a steep density gradient exists. LASNEX analysis of underdense plasma evolution from gold disks shows that the plasma can be separated into two regions, the transition and the corona. Using exploding-foil targets on Nova, we have demonstrated that production of both blue- and red-shifted components is associated with the existence of a quarter critical plasma density, and that densities higher than $n_c/4$ are responsible for producing backscattered blue photons. We have also presented evidence for seeding of stimulated Brillouin scattering by the two-plasmon decay, which may help explain differences between theory and experiments in this field.

In our hydrodynamics work, we have progressed in both theory and experiments. We developed, installed, and tested SIMPLE, a new one-dimensional, three-fluid, local thermodynamic equilibrium, hydrodynamics code for small computers. It not only helps us design and understand simple ICF experiments, but also helps ICF researchers understand how physical processes interact. In addition, a series of planar foil mix experiments is under way that,

when completed, will shed more light on plasma/hydrodynamic instabilities in a radiation-driven and preheated environment; our results should lead to improved computational models.

1989 was an outstanding year for x-ray laser (XRL) research at LLNL. We developed XRLs that operate at 44.8 Å, the optimum wavelength for x-ray holographic imaging of living cellular materials. We also achieved amplification inside the "water window," at wavelengths of 43.2 and 35.6 Å. In addition, we have identified optimum samples to study and have calculated the XRL coherent energy needed for an x-ray hologram. This work promises to provide a currently missing tool for macromolecule-resolution analysis.

We have begun experiments on the Nova laser to better understand the physical mechanisms working in a XUV amplifier setup with neon-like ions. We made a silver amplifier, the highest-Z neon-like amplifier to date. We also determined that short-pulse laser-pumped recombination x-ray lasers appear feasible, and that they would allow very efficient scaling to short wavelength and high gain in the quasi-steady-state regime. With our development of a facility that can irradiate solids with x-rays, we aim to facilitate the study of higher-density regions of x-ray emission from conventional laser-matter interaction experiments.

Much of our progress has been in the area of diagnostics. We developed instruments that utilize electronic gating of microchannel plates, which take gated images with very short aperture times (<100 ps). Anticipating that future ICF implosion experiments will have much higher fuel and pusher areal densities, we have begun looking for second generation line-broadening species, namely neon-like xenon. We also determined that diamond PCDs should make effective x-ray detectors because of their speed, simplicity, and insensitivity to visible radiation.

By modifying an existing Wölter x-ray microscope, we developed a multichannel x-ray microscope for time- and space-resolved high-

resolution x-ray images of laser fusion experiments. The new diagnostic can be applied to studies of implosion symmetry, hydrodynamic instability, and plasma physics. A new thick-aperture design for neutron penumbral imaging achieves at least 10- μm resolution over a 150- μm field of view with no significant first-order distortion. We also developed a general expression for the signal-to-noise ratio in a reconstructed penumbral-aperture image; this expression allows a fast prediction of an imaging system's performance with a given target.

We improved our ability to reliably measure and analyze, with Nova's current mode neutron detectors, ion temperature for DT-filled capsules. Another neutron based technique is activating of trace material in the pusher. Previous techniques rely on collecting aluminum from laser-fusion target with an estimated collection fraction; our new technique addresses this weakness by previously irradiating the target with thermal neutrons to produce a rubidium tracer. Another advance in diagnostics is in the generation of ultrafast, high-voltage pulses. Vacuum triode drivers, bulky and expensive, are currently used. We have developed the MOSFET/non linear transmission line avalanche diode pulser, an alternative for use with new gated x-ray camera structures.

3.2. Laser Plasma Interactions

3.2.1. Introduction

Activities in laser-plasma interactions during this last year included efforts in understanding the plasma evolution of laser-heated matter and methods for characterizing the plasma as well as investigating laser-plasma instabilities. Much of the work in this section has been done using other lasers since much of Nova's activities during the past year have emphasized hydrodynamic and implosion target experiments, and characterizing and improving the laser.

In the article entitled "Filamentation and Second Harmonic Emission in Laser-Plasma Interactions," the origin of second harmonic emission ($2\omega_0$) from underdense plasmas is studied. In experiments done on the Chroma laser at KMS Fusion Inc., one beam of light is used to explode a foil to produce an underdense plasma. A second beam is focused to a narrow line in the underdense plasma. Density depressions due to

filamentation of the high intensity line focused beam are observed in the plasma using a holographic interferometer. The $2\omega_0$ emission is correlated with the position of the density depressions.

Laser-plasma interactions have been studied using exploding foil targets on Nova. In the article entitled "Three-halves Harmonic Emission from Two Color Experiments," the emission spectrum of the three-halves harmonic ($3\omega_0/2$) is correlated with the plasma density using an exploding foil target. The $3\omega_0/2$ spectrum contains both a blue and red shifted component. This experiment demonstrated that production of both components is associated with the existence of a quarter critical plasma density. The scattered light spectrum near the laser wavelength due to stimulated Brillouin scattering (SBS) was also measured in these experiments. In the article "Scattered Light Near the Laser Wavelength from Nova Two-Color Experiments," evidence for seeding of SBS by the two-plasmon decay instability is presented.

In the article entitled "Initial Results from Low-Density, Long-Scalelength Plasma Experiments," initial results are presented of the characterization of exploding foil plasmas with densities of $n_e/n_c < 0.1$. After these initial characterization experiments, the intensity scaling of stimulated Raman forward scattering (SRFS) will be measured in these very low density plasmas. SRFS has already been observed in one shot of these experiments. These initial experiments show that the plasma densities and temperatures at late time depend on the initial target geometry and are probably due to lateral heat transfer effects. These effects have been confirmed using LASNEX modeling.

The use of harmonic light imaging for characterizing laser-produced plasmas is examined in the article entitled "An Experimental Investigation of Refraction-Induced Distortions in Harmonic Light Images of Laser-Produced Plasmas." The second harmonic ($2\omega_0$) and three-halves harmonic ($3\omega_0/2$) light are produced in the plasma near the critical density (n_c) and quarter-critical density ($n_c/4$) surfaces, respectively. Imaging these frequencies of light could provide information about the location of these plasma densities, but the images could be distorted by the surrounding plasma. This experiment examines the level of distortion by imaging a known optical pattern

through underdense plasma. The experiments performed on the Chroma laser at KMS Fusion Inc. showed that the images can be significantly refracted, especially in the near region with a steep density gradient.

In "Plasma Evolution from Laser-Driven Gold Disks," analysis of experiments of the underdense plasma evolution from thick Au disks is presented. The experiments, done in 1988 using the Chroma laser at KMS Fusion Inc., measure the underdense plasma evolution from Au disks irradiated with $0.53\text{ }\mu\text{m}$ light using a variety of pulse shapes. The plasma evolution was studied using a four-frame holographic interferometer for $n_e/n_c < 0.25$ and Au M-band emission for $0.25 < n_e/n_c < 0.5$. Analysis using LASNEX show that the plasma can be separated into two regions, the transition and the corona. The transition region where $n_e/n_c > 0.25$ is close to steady state and expands away from the disk relatively slowly. The corona region where $n_e/n_c < 0.25$ initially expands in nearly a self-similar expansion. After about 1 ns the corona reaches a steady-state regime although longer pulse experiments are needed to confirm the steady-state behavior.

Author: Robert Kauffman

3.2.2. Filamentation and Second-Harmonic Emission in Laser-Plasma Interactions

The filamentation instability¹ in laser-produced plasmas is of great interest because of its possible role in degrading the coupling of lasers to inertial-confinement-fusion targets. Filamentation has recently been studied extensively in numerical simulations²⁻⁸ and experiments.^{9, 10} Filaments have been invoked to explain numerous experimental phenomena, one of which is the observation^{11,14} of emission at the second harmonic of the incident laser frequency, ω_0 , from underdense plasmas at angles perpendicular to the incident laser beam. No direct link between filaments and underdense $2\omega_0$ emission has been demonstrated to date, in part because there are numerous theoretical explanations for harmonic emission,^{14,15} and in part because the density perturbations produced by filamentation have only recently been clearly identified.⁹

In this experiment, a short-pulse, high-intensity laser beam interacts with an underdense plasma formed by another laser pulse, and produces a density perturbation which

can be detected interferometrically. A fast-framing camera simultaneously records the image of the $2\omega_0$ emission. This allows a simple comparison between the spatial positions of the density perturbation and the second-harmonic emission. The main result is that the $2\omega_0$ emission is spatially correlated with the density perturbations, but is not observed for the entire length of the filament.

The experiment was conducted using two beams of the Chroma laser at KMS Fusion. One beam irradiated a thin foil target ($0.35\text{-}\mu\text{m}$ -thick CH) at an angle 50° from the target normal to form a plasma. This laser pulse was nominally square with a 500-ps duration, and had an energy of 65 J and a wavelength of $1.06\text{ }\mu\text{m}$. The laser spot on target was circular with a diameter of $450\text{ }\mu\text{m}$. The foil burned through approximately 300 ps into the laser pulse and eventually formed a plasma that was nearly symmetric about both the plane of the target and the laser axis.

A second, short-pulse beam called the interaction beam, was delayed relative to the falling edge of the first pulse by 300 ps. The interaction-beam wavelength was also $1.06\text{ }\mu\text{m}$; the pulse was Gaussian with a 100-ps FWHM, and had a laser energy of 35 J. The interaction beam was focused to a line by a positive spherical ($f/6$) lens with a negative cylindrical lens placed just before it. At best focus, the beam had dimensions of $50 \times 50\text{ }\mu\text{m}$, which was verified by imaging the transmitted beam at the target plane. The position of best focus was varied between target irradiations by moving the spherical focusing lens. The polarization of this beam was linear and tilted with respect to the horizontal plane by 45° .

A single-frame holographic imaging system¹⁶ recorded the line integral of the phase shift of a probe beam which passed through the plasma, tangential to the target plane. The probe beam had a wavelength of $0.25\text{ }\mu\text{m}$, was a Gaussian pulse of 20-ps duration, and was aligned parallel to the line focus on the interaction beam. The probe-beam pulse arrived at the target at the same time as the interaction beam. Because the interaction beam is focused to a line, rather than a circular spot, it is possible to see a large phase shift [see Fig. 3-1(a)] even if the density perturbation produced by the interaction beam is small.^{9,10}

The other important diagnostic in this experiment is a gated optical imager (GOI)¹⁷ which is a single-frame framing camera with

Fig. 3-1. Comparison of (a) an interferogram and (b) a 120-ps gated image of a plasma with the best focus of a line focus located at the target plane. The pictures are recorded at the same time and only the side of the plasma facing the laser (incident from the right) is shown. The white lines in (a) show the estimated path of the $1/e$ level of the intensity profile of the interaction beam without the plasma.

120-ps gate width. The target was imaged onto the camera by a single-element lens ($f/6$) which viewed the plasma from the side opposite to the holographic imager, and at an angle of 5° in front of the target relative to the probe beam. The light from the plasma was spectrally filtered to provide $\sim 10^{-16}$ attenuation of emission outside a wavelength window of $\lambda = 5320 \pm 100 \text{ \AA}$. The GOI was gated at the same time that the interaction beam reached the target. Because of this short gate width, all of the $2\omega_0$ emission is due to the interaction beam, not the plasma-forming beam. The time resolution of the GOI simplifies the interpretation of the images compared to time-integrated measurements¹¹ and minimizes blurring due to plasma motion.

In Figs. 3-1 and 3-2, we compared the simultaneous interferograms and the GOI pictures from two target irradiations. The object plane for both diagnostics was located midway into the plasma. By imaging a 400- μm -diam steel ball onto both diagnostics we can reconstruct images with the same magnification. The three small spots seen in the lower half of the GOI pictures are defects in the microchannel plate which reference to the position of the second-harmonic emission on any shot relative to the target plane. The target plane in the interferograms is well defined by the shadow of the washer that supports the CH foil [left side of Figs. 3-1(a) and 3-2(a)].

For the target irradiation of Fig. 3-1, the best focus of the interaction beam was positioned at the target plane. The interferogram shows a density perturbation which is centered vertically and extends between 150 and 400 μm in front of the target. Subsequently, the spherical focusing lens was moved 400 μm away from the target to produce the data shown in Fig. 3-2. Once again there is a density perturbation centered vertically but with a smaller horizontal extent. These two figures show that the peak of the $2\omega_0$ emission moves as well.

From the interferogram, we can deconvolve the background density with the assumption of axial symmetry, by using an Abel-inversion routine while neglecting the phase shift due to the line-focus density perturbation. This gives the density profile in the central plane of the plasma which passes through the laser axis. The profile is then corrected for the density perturbation produced by the line-focused interaction beam by noting that the phase shift,

$\Delta\phi$, of the probe beam through the line focus is related to the density perturbation δn by

$$\Delta\phi \approx 9 \times 10^{-22} \pi L_{p,\mu} \lambda_{p,\mu} \delta n, \quad (1)$$

where $\lambda_{p,\mu}$ is the wavelength of the probe beam in μm , $L_{p,\mu}$ is the path length through the plasma in μm , and δn has units of cm^{-3} . We will take $L_{p,\mu}$ to equal the length of the line focus ($= 250 \mu\text{m}$) to get $\delta n = 6 \times 10^{18} \Delta\phi$, which we subtract from the unperturbed density profiles (dotted lines) to get the density profiles (dashed lines) shown in Fig. 3-3.

The density perturbation given by simple pressure balance between the laser and the plasma can be calculated from Eq. (2) of Ref. 9; for plasma parameters of $T_e \sim 500 \text{ eV}$ (from simulations), $n/n_c = 0.1$, and $I_L = 1.3 \times 10^{15} \text{ W/cm}^2$, we calculate $\delta n/n \sim 30\%$. The observation in this experiment that the perturbations have a finite axial extent suggests, however, that the laser light is self-focusing (and, subsequently, defocusing). The size of the initial density perturbation is large enough that predicted filamentation focusing lengths based on small initial perturbations^{10,18} are not applicable. One can integrate the equation of refraction of a light ray¹⁹ and define the self-focusing length l_{SF} to be the distance required to deflect a ray at the radius of the hot spot, $r = r_0$, to $r = 0$. Using the equation for ponderomotive pressure balance we find

$$l_{\text{SF}} = \alpha \frac{\omega_0}{\omega_{pe}} \frac{v_e}{v_{os}} r_0 \quad (2)$$

where ω_{pe} is the electron plasma frequency, v_e is the electron thermal velocity, and v_{os} is the electron oscillatory velocity in the light wave. The dimensionless parameter α is selected, based on numerical simulations of the nonlinear Schrödinger equation,²⁰ to be $\alpha = 1.7$. For the plasma parameters listed above, with an initial beam radius of 50 μm , we find $l_{\text{SF}} \approx 250 \mu\text{m}$, which is in good agreement to the observed focusing distance (from the edge of the plasma to the maximum density perturbation). The inverse bremsstrahlung absorption length is calculated to be (for $Z = 3.5$) $> 4 \text{ mm}$, which is too large to explain the finite axial extent of the density perturbation.

Our calculations with a saturable (exponential) non-linearity suggest that as the

Fig. 3-2. Comparison of (a) an interferogram and (b) a 120-ps gated image of a plasma with the best focus of a line focus located $400\mu\text{m}$ in front of the target plane.

Fig. 3-3. Comparison of density profiles (obtained by unfolding the interferograms) with relative intensity contours of $2\omega_0$ light (solid lines). Shown are density profiles along the laser axis with (---) and without (....) the perturbation produced by the interaction beam. The profiles are obtained from the data for the cases shown (a) in Fig. 3-1, and (b) in Fig. 3-2.

laser beam self-focuses, it develops substructure and collapses into a number of filaments. In 3-D calculations, these filaments become very intense and small (a radius of only several wavelengths). The resulting strong gradients in intensity and density enhance the $2\omega_0$ emission. However, ion waves are expected to be generated in the final, nonadiabatic stage of the formation of such intense, narrow filaments. These ion waves can provide density modulations which act to delocalize or spray, the filaments.^{3,5}

We converted the recorded density of the $2\omega_0$ emission to relative intensity and compared these profiles along the laser axis to the density profiles for the two cases discussed. Figure 3-3(a) is for the interaction-beam focus at the target plane, as seen in Fig. 3-1, and Fig. 3-3(b) is for the interaction-beam focus 400 μm in front of the target plane, as seen in Fig. 3-2. The peak of the $2\omega_0$ emission occurs in or near the density perturbation in both cases, probably at the position of the peak incident intensity. The difference in position of the $2\omega_0$ emission relative to the filament is probably due to the dynamics of the self-focusing process; in Fig. 3-1 the beam is converging in the plasma while in Fig. 3-2 the beam is diverging. Refraction of $2\omega_0$ rays propagating through the filament, combined with the large extent of the radiating object compared to the depth of focus of the imaging lens, is probably responsible for the vertical and horizontal extent of the $2\omega_0$ emission seen in the images [particularly Fig. 3-1(a)].

While our results are generally consistent with previous explanations^{11,13} of $2\omega_0$ emission from underdense plasmas and with the belief that this happens in filaments, these time-resolved results show the importance of the filamentary focus as the chief source of emission, rather than emission along the length of the filament. The source term is proportional to $E_0(E_S \cdot \nabla)n + E_S(E_0 \cdot \nabla)n$, where E_0, E_S are the incident and scattered laser lights, respectively. Hence, in a uniform filament one expects the $2\omega_0$ source to be localized at the filament edge along the length of the filament. Experimentally one sees a laterally diffused $2\omega_0$ source concentrated near the end of the measured filament, presumably the focus. There one expects ∇n no longer to have strict transverse symmetry, and to be quite large, with E_S likely not in a backscatter collimated direction. All of this should allow a considerable relaxation of constraints on $2\omega_0$

emission direction and polarization.

Investigation of these and other spectral features is left to future work. It is likely that the z-extension observed in time-integrated images¹¹ (for long laser pulses) is largely due to focus motion.

In conclusion, we found that there can indeed be a correlation between second-harmonic emission and density perturbations produced by a laser beam interacting with an underdense plasma. Although we have demonstrated experimental conditions in which $2\omega_0$ is associated with filaments, the use of $2\omega_0$ as a diagnostic for filamentation in underdense plasmas is problematic, because (1), the $2\omega_0$ emission does not show the complete extent of the filaments (see Figs. 3-1–3-3), (2), significant $2\omega_0$ emission is observed only at high intensities (we do not observe significant $2\omega_0$ from filaments for the conditions of Ref. 9), and (3), the $2\omega_0$ emission does not give a quantitative measurement of $\delta n/n$. The unambiguous way to study filamentation is with direct measurement of the density perturbation, i.e., using interferometry of Schlieren.

Authors: P. E. Young, H. A. Baldis (National Research Council of Canada), T. W. Johnston (Institut National de la Recherche Scientifique-Energie), W. L. Kruer, and K. G. Estabrook

3.2.3. Initial results from Low-Density, Long-Scalelength Plasma Experiments

I. Introduction. In 1989 we began a series of experiments to study the scaling of low-gain instabilities in laser-produced plasmas. Specifically, these experiments will address the scaling of stimulated Raman forward, Brillouin, and Compton scattering in long-scalelength plasmas where the characteristic electron density is less than $0.1 n_c$, and the electron temperature is about 1 keV. These mechanisms are of concern to laser fusion because they can lead to decreased laser absorption, and can generate hot-electrons that degrade the gain of ICF targets. Quantitative understanding of the scaling for these processes will allow us to calculate their levels in a high-gain target design.

The objective for the first phase of experiments is to produce a well-diagnosed, reproducible plasma environment with an electron temperature of about 1 keV, an electron density $n_e/n_c < 0.05$ (n_c for 527-nm light), and a

density scalelength >2 mm. Upon completion of this phase of experiments, our experiments will focus on the intensity scaling of stimulated Raman forward scattering in such a plasma. To date we have made significant progress toward our first objective, and have gained insight into the experimental parameters that affect creating a reproducible plasma. Additionally, we have performed one experiment to demonstrate our ability to observe SRFS from plasma densities as low as $n_e/n_c = 0.02$. In this article, we will present our strategy for performing laser-interaction experiments relevant to high-gain ICF targets, will describe our experiments to produce and characterize plasmas with the desired conditions, and will report preliminary results. The evolution of the peak-electron density from these experiments will be discussed, as well as the effects of radial heating and expansion on the reproducibility of the density evolution. These observations are compared to calculations using 2-D LASNEX.

II. Experimental Strategy. For most high-gain ICF target designs, there will be very large regions (a few millimeters or more) of very underdense plasma ($n_e/n_c < 0.1$), with electron temperatures on the order of 1 keV or more. Under these conditions, SRS backscatter will be strongly Landau damped ($k_{epw}\lambda_{De} > 0.3$ at these densities and temperatures) so that its role in producing hot-electrons is significantly reduced. However, given sufficient density scalelengths, low-gain instabilities such as SRFS can grow to significant levels to be a source of hot-electrons. It is widely held that Raman forward scattering is a convective instability, with a gain that increases exponentially with scalelength and intensity. The SRFS scaling experiments will allow us to make quantitative comparisons with this model. Several papers have been published that describe the model for SRFS in detail,^{21,23} so only a brief summary will be given here.

Raman scattering is the decay of a laser light wave into a scattered light wave and an electron plasma wave, and obeys the frequency and wave number matching conditions

$$\omega_0 = \omega_s + \omega_{epw} \quad (1)$$

$$k_0 = k_s + k_{epw} \quad (2)$$

where the subscripts (0, s, epw) refer to the laser, scattered, and electron plasma waves,

respectively. For SRFS, the wave numbers k_0 , and k_s are both positive and roughly the same magnitude, thus making k_{epw} small. For the plasma conditions stated above, $k_{epw}\lambda_{De} \ll 0.3$ so that SRFS is not limited by Landau damping, but rather by plasma density inhomogeneity. For a parabolic density profile, $n_e/n_c = n_p/n_c (1 - x^2/L^2)$ with peak density of n_p and scalelength L , the intensity of the SRFS light would be

$$I_{SRFS} = (I_{noise}) \exp(G_{SRFS}) \quad (3)$$

where the gain coefficient G_{SRFS} is given by²²

$$G_{SRFS} = 7 \frac{\gamma_0^{3/2}}{(v_s v_{epw})^{3/4} (k'')^{1/2}} \quad (4)$$

The term γ_0 is the homogeneous growth rate for SRFS, v_s and v_{epw} are the group velocities for the scattered and electron plasma waves, and k'' is the wave number mismatch factor due to the density profile. Expressed in terms of laser intensity and wavelength, plasma density, temperature and scalelength, the gain coefficient G_{SRFS} scales as

$$G_{SRFS} \sim (I\lambda^2)^{3/4} (T_e)^{-1/4} (n_e/n_c)^{7/8} L/\lambda, \quad (5)$$

or

$$G_{SRFS} \sim I^{3/4} \lambda^{1/2} L. \quad (5b)$$

To perform scaling experiments for SRFS in the exact plasma conditions expected in high-gain ICF targets would require a few megajoule laser. With energy limitations on present laser facilities, one can achieve density scalelengths, L , about a factor of 2–5 lower than those expected for high-gain ICF targets, but similar or greater gain in SRFS levels can be realized by performing experiments at higher intensities and longer wavelengths. For our scaling experiments we will achieve SRFS gains greater than those expected in high-gain ICF targets based on the model above. Quantitative measurements of the SRFS levels will allow us to place an upper-bound on the level expected in high-gain ICF targets. Our experiments will use a 527 nm, high-intensity (up to 4×10^{15} W/cm²) interaction beam into a plasma preformed by several 351 nm beams. By using a 527 nm beam instead of 351 nm

to study SRFS, the gain coefficient G_{SRFS} (Eq. 5b) will be increased by about 25% for the same size plasma, which corresponds to over 3 orders of magnitude increase in SRFS signal at the highest intensities. Thus, choosing to use a 527-nm beam for the scaling experiments will increase our confidence in the absolute signal level measured and will increase the likelihood of observing adequate signal at the lower laser intensities.

III. Experimental Design and Plasma

Diagnosis. We now describe our experiments to produce and characterize a low-density, long-scalelength plasma that will be used in the scaling experiments. Using the 1-D London-Rosen model,²⁴ a CH exploding foil with a 2 μm thickness was selected in order to give the maximum scalelength at the desired density by the end of a 2 ns heating pulse. A large laser spot is necessary to minimize 2-D expansion effects for this pulse length, so several beams are needed to preform the plasma. The details for the specific experiments performed to date are as follows:

We use 7 beams of Nova (351 nm, $E_{\text{total}} = 15$ kJ, $\tau = 2$ ns square) to preform a plasma by irradiating a 2.0 ± 0.1 μm -thick-CH foil. The foil is 1.5–2.0 mm wide, 4–5 mm in length, and is supported at the long ends by a wire frame. The beams illuminate both sides of the foil in an overlapping spot, 950 μm in diameter, with a total average intensity of 1×10^{15} W/cm², and an intensity of about 1.4×10^{14} W/cm² for each beam. After the foil has become transparent to the preform beams and expanded to the desired conditions, the SRFS scaling experiments will use an additional 527 nm, high-intensity beam to drive the instability. For the present experiments, this beam is used as a low-intensity probe for diagnosing the plasma conditions. A schematic for these experiments is shown in Fig. 3-4 illustrating the beam and target geometry relative to the principal diagnostics.

We use a 527 nm, low-intensity beam (which we call the probe beam) to drive SRS in this plasma. The probe beam was focused to a 400- μm -diam spot located in the central region of the 950- μm -diam spot formed by the overlapping 351 nm plasma-forming beams. The average intensity of the probe beam was less than 2×10^{14} W/cm² and was delayed by 0.25 ns so that there would be little effect on the hydrodynamics of the exploding foil (this assumption was confirmed using the 1-D hydro-code SIMPLE.²⁵ The evolution of the peak density and

temperature of the foil is obtained by examining the backward-directed SRS spectrum from either the plasma-forming beams or from the probe beam. The data reduction methods for these type of measurements are described in other publications.^{26, 27} Time-resolved SRS spectra are observed at an angle of 27° from the probe beam and 39° from the nearest plasma-forming beam (as shown in Fig. 3-4). A discrete-channel spectrometer observed time-resolved emissions at 3/2 and 1/2-harmonics from the 351 nm plasma forming beams. A measure of the disappearance of these signals yields the time when the plasma drops below $0.25 n_c$ for the 351 nm beams (that correspond to $0.56 n_c$ for 527-nm light). An example of time-resolved SRS spectra from these experiments is given in Fig. 3-5.

Future diagnostics are planned to obtain 2-D time-resolved images of the UV-bremsstrahlung emission looking edge-on to the plasma in order to measure the plasma density profile. In past experiments using CH exploding foils,²⁶ the radius of the foil was much larger than the laser-spot, leaving regions of cold, dense plasma outside of the laser spot. This cold, dense plasma occludes imaging the UV-emission from a large region of the plasma, so our targets were designed to be only slightly larger than the laser-spot. Initial difficulties in target fabrication prevented us from building a small circular foil, so a rectangular foil was built with the width slightly larger than the laser spot as an approximation to a circular target. Later experiments will use a circular target currently being developed. In the next section, we present results from the rectangular foils using two different target widths.

IV. Experimental and Simulation Results.

Two target widths, 1.5 and 2.0 mm nominally, were used in our characterization experiments. Time-resolved SRS spectra and half-harmonic spectra were obtained and used to infer the peak-density and temperature evolution for each experiment. After 1 ns into the plasma-forming pulses, the electron temperature inferred from the Landau damping cutoff of the spectra is about 1.3 keV (assuming $k_{epw}\lambda_{De} = 0.3$) and is consistent for all experiments. The density measurements reported here use this electron temperature in their calculations.

Measurements of the temporal evolution of the peak-electron density from a 2-mm wide target is shown in Fig. 3-6a. Collected in this figure are measurements from the 351-nm SRS

Fig. 3-4. The target irradiation geometry and the primary diagnostics are shown in this schematic diagram.

Fig. 3-5. Time-resolved spectra of backscattered light showing SRS from the preform beams and from a 527 nm probe beam that was delayed 0.25 ns after the start of the 351 nm beams. A wavelength and timing fiducial is shown at 527 nm. Spectra such as this are used to infer the peak-plasma density and temperature in the interaction region.

Fig. 3-6a. Peak-density evolution of a 2.0-mm wide target. The density information is derived from various measurements: diamonds for 351 nm SRS and $w_0/2$ emission, and triangles from 527 nm SRS. The solid line shows the peak density calculated by 2-D LASNEX for an infinite-diameter target and the dashed line shows the peak density calculated for a 2.0-mm diam target.

spectrum, the 527-nm SRS spectrum, and from emission originating at the $n_c/4$ density surface from 351-nm. The data points shown are when the emission corresponding to the indicated density fell to 20% of its maximum value. The solid and dashed lines are the results from 2-D LASNEX simulations and are discussed later in this section. Initially, the density fell rapidly as the foil burned-through after about 0.4 ns. After 1 ns, the peak density decreased less rapidly. The lowest density measured was about $0.05 n_c$, which occurred at 2.2 ns. The probe beam turned off after 2.2 ns, so it may be possible in the future to extend this measurement for this target.

The targets that were only 1.5-mm wide gave somewhat different results, as shown in Fig. 3-6b. The initial, rapid decrease in peak density was about the same as the 2.0-mm wide target. After this, the density decay was more rapid than the 2.0-mm target. The lowest measured density was also about $0.05 n_c$, but occurred at 1.2 ns. The solid line indicates a 2-D LASNEX calculation to model the 1.5 mm wide target.

The interaction of the plasma-forming beams and the plasma was modeled using 2-D LASNEX.²⁸ Because the actual focal-plane distribution of the overlapping preform beams was not measured, the distribution was approximated by assuming that the far-field intensity patterns of all the beams were identical, and then a measured pattern was used to compute an azimuthally-averaged distribution in the overlap region. Because the overlapping focal spot is generated by several different $f/4$ beams, the interaction was modeled as a cone of $f/4$ rays at 50° to the target normal. All calculations were done using a heat transport flux-limiter of 0.03.

Three different cases were considered. In the 1.5-mm case, the initial foil diameter was 1.5 mm, while the 2.0 mm case used a 2.0-mm diam foil. The "infinite" case used a 5.0-mm diam foil. The laser energy, 15 kJ, and nominal intensity, 1×10^{15} W/cm², were the same for all three cases. While the targets used in the experiments are not radially symmetric as assumed in the simulations, the modeled plasma evolution compares well with the rectangular targets. A comparison of both the infinite and 2.0 mm case is shown in Fig. 3-6a. No arbitrary timing offsets have been added to improve the fit between experiment and simulation; the absolute timing of the data is accurate to ± 0.075 ns and no

changes have been made to the simulation. The agreement is very good throughout the density range measured. The infinite case appears to be in better agreement with the experimental data except at about 2.0 ns, where the error bars may allow agreement with either simulation.

Good agreement between the density evolution of the 1.5 mm target and the 1.5 mm foil simulation was also found, as seen in Fig. 3-6b. The simulation accurately reproduced the data over the range of densities. No significant discrepancy between the measured and calculated burn-through times was seen in any of the cases, in contrast to earlier, single-beam experiments where a substantial discrepancy was observed.^{6, 29}

The difference in the foil widths is most evident in the decay to lower density. This difference arises because in the 1.5 mm case, the target was mass-limited, i.e., the laser heated and ablated all of the target material in the radial direction before the plasma-forming pulse ended. This was not true in either the infinite case or the 2.0 mm case where the target material was not heated all the way to the edge of the target. A mass of cold, dense material remained at the edge of the plasma and "tamped" or inhibited the radial expansion of the plasma. As an example of the presence of the higher-density material, at 3.5-ns into both the 2.0 mm and infinite simulations, the $n_c/4$ surface was at a radius of 1 mm from the target axis. This surface had completely disappeared after 0.7 ns in the 1.5 mm foil simulation. It appears that a finite extent of the target in only one direction is sufficient to induce this behavior since both targets used in the experiments could be considered to be very large in the other direction.

The electron temperature calculated by LASNEX for all three cases, while consistent between simulations, is not in agreement with the experiments. The electron temperature given by LASNEX was a factor of about 1.6 higher than measured by the Landau-damping cutoff for all experiments. This result is in agreement with previous single-beam experiments,²⁹ where a factor of about 1.5 was found between the temperatures calculated in LASNEX and those measured in the experiment.

A single experiment was also performed in which a 527 nm, high-intensity beam ($\sim 4 \times 10^{15}$ W/cm²) irradiated the preformed plasma with a spot diam of about 200 μ m. This beam was

Fig. 3-6b. Peak-density evolution of a 1.5 mm wide target. The LASNEX calculations are for a 1.5-mm-diam target. A dramatic difference in the density evolution compared to the 2.0-mm-wide target can be seen.

delayed 1.1 ns after the beginning of the plasma-forming beams, and had a 2.0 ns square pulse shape. Streaked spectra from SRFS (shown in Fig. 3-7) were observed at an angle of 40° from the forward-direction which corresponded to densities of about $0.02 n_c$ at a time of 3.0 ns after the plasma-forming beams began. The spectra early in time are backscattered SRS from the 351-nm-plasma-forming beams. The detector sensitivity was increased for this shot to insure that the SRFS signal could be observed, so the detector is saturated in this photo. This experiment demonstrated that the dynamic range of our instruments should be sufficient to observe SRFS in the scaling experiments.

V. Conclusions. A method for producing low-density, long-scalelength plasmas has been demonstrated. The temporal evolution of the peak-plasma density was measured and is in good agreement with 2-D LASNEX simulations. Both experiment and simulation show that the late-time plasma-density evolution is sensitive to the initial target size, and the simulations show this to be an effect of radial heating and expansion. Additionally, an experiment was performed where SRFS from densities as low as $0.02 n_c$ was observed in this plasma.

Authors: D. S. Montgomery, S. H. Batha, R. P. Drake, K. G. Estabrook

Major contributors: D. W. Phillion and B. A. Remington

3.2.4. An Experimental Investigation of Refraction-Induced Distortions of Harmonic Light Images of Laser-Produced Plasmas

In this paper the results of a simple experiment designed to study refraction-induced distortion of images of narrow-band-light emitting regions in a laser produced plasma are described. Source regions were formed by relaying the image of a backlit pinhole array into the plasma (25- μm -diam pinholes; 30-ps, 0.25- μm backlighter beam). Images of these source regions were recorded on film and compared, with and without plasma present. Refraction effects, evidenced by translation of the apparent pinhole locations and blurring of the pinhole spots, are compared for different plasma sizes and for different delays of the backlighter pulse with respect to the plasma-production laser pulse. These results were then compared with expectations based on numerical

raytraces through a plasma density profile derived from a hydrodynamic simulation of our exploding-foil plasma. The significance of these results in the imaging of harmonic and half-harmonic light as a plasma density diagnostic is discussed.

Introduction. The characterization of the electron density profile produced when an intense laser beam impinges on the surface of a solid material is of fundamental interest in such fields as inertial confinement fusion,³⁰ x-ray lasers,³¹ and high-intensity laser-matter interactions.³² In hot, dense plasmas non-perturbative, in situ measurements of such plasma parameters as electron density, temperature, and expansion velocity are impractical. Instead experimenters have relied heavily on external probing (e.g., holographic interferometry,³³ Thomson scattering^{34, 35} or plasma emission spectroscopy³⁶ for estimates of these quantities.

Many interesting laser-plasma interaction phenomena occur at plasma densities equal to or greater than the quarter-critical density for the interaction laser³⁷ ($n_{\text{critical}} = 1.1 \times 10^{21} \lambda_{\text{laser}}^{-2} \text{ cm}^{-3}$ where λ_{laser} is the laser wavelength in microns). To effectively probe this high-density region of the profile, the probe wavelength must be sufficiently short so as to minimize refractive effects. Although in principle one can choose to work with arbitrarily short wavelengths for plasma diagnosis, it is most often the case that the wavelengths of practical significance are sufficiently long that the ray trajectories of the collected light undergo significant refractive bending as they emerge from (in the case of plasma emissions) or pass through (as in the case of external probing) the sample plasma. This fact has several consequences. In experiments where the light is simply collected by a spatially-integrating detector, some of the emitted light may never reach the detector. In experiments in which it is the goal to form an image of the light-emitting region in the plasma, refraction can distort the image to the point where little or no useful information can be inferred. In external probe experiments refractive bending of the light rays can lead to anomalous spatially resolved data.

The primary motivation for this work was the fact that in past laser-plasma interaction experiments much has been learned about the critical and quarter-critical density surfaces through the study of the second (2ω) from the

Fig. 3-7. Forward SRS spectra of the 527 nm probe beam showing scattering at low densities. Backscattered SRS from the 351-nm beams begins shortly after the beams turn on. The probe beam has been delayed 1.1 ns after the preform beams begin. The instrument is saturated on this particular experiment during both the 351-nm SRS emission and the 527- nm SRS emission.

critical-density surface) and three-halves ($3/2 \omega$ from the quarter-critical density surface) laser-harmonic emitting regions imbedded in the plasma density profile. In particular, by imaging the sources of these emissions, it is possible to spatially and temporally resolve the evolution of the plasma density profile.³⁸ The success of this technique has been due largely to the fact that small, relatively low power lasers, in order to reach the intensities relevant to the excitation of stimulated processes in plasmas, were focused to small spots at the target. Because refractive bending of light rays is a path-integrated effect, the shorter the path and the lower the density, the smaller the effect. For this reason refraction has typically presented few difficulties in probing plasma densities below quarter-critical. However, given the capabilities of present-day lasers and the relevance of large irradiation spots to ICF, the dimensions of the larger plasmas now attainable can exceed several thousand laser wavelengths.³⁹ Under such conditions it is necessary to reevaluate the impact of refraction when making optical probe measurements likely to be adversely affected by ray bending.

Description of Experiment. The experiments were performed at the Chroma Laser Facility at KMS Fusion. A 150-J, 1.5-ns, 0.53- μm laser beam, normally incident on a 600- μm -diam, 22- μm -thick gold disk target in a 200–400- μm focal spot, was used to produce a test plasma. Our goal was to simulate the conditions under which one is interested in obtaining images of a harmonic-light source region embedded in the plasma. To accomplish this the image of a pinhole array, positioned outside the vacuum vessel and backlit by a low-intensity ($I_{\text{probe}} < 10^6 \text{ W/cm}^2$), 0.25- μm probe beam, was relayed (1:1 magnification) perpendicular to the plasma (symmetry) axis (parallel to the disk face) to a plane lying on the plasma axis as shown in Fig. 3-8. The pinhole array image formed in the plasma volume was then relayed by an $f/2$ catadioptric reflector system³³ to a second image plane outside the vacuum vessel where the image, properly filtered for 0.25- μm light and satisfactory signal to noise ratio, was recorded on negative film. By varying the time of arrival of the short (30 ps) 0.25- μm backlighter probe with respect to the 0.53- μm plasma production beam (delay variable from 0–1200 ps), refraction-induced image degradation could be studied at different times during the evolution of the plasma profile.

For reference, prior to each shot, an image of the target was superimposed on the target-plane image of the pinhole array at the film plane. This “preshot” image served both as a final check on the positioning of the target with respect to the array and as a reference with which to compare when quantifying distortions observed in the images recorded when the plasma was present. After the preshot image was recorded, the film was translated several centimeters to the side and the “shot” image was recorded (plasma-forming beam fired).

For the four shots taken, we investigated both low- and high-refraction cases by varying the 0.53- μm laser focal spot size and the time at which the target-plane image of the pinhole array was recorded (as described above). Small spot sizes and early probe pulses were expected to induce little refractive image distortion because a small focal spot minimizes the path-integrated refractive bending and early probe times allow insufficient time for appreciable plasma density buildup in the vicinity of the pinholes. For large spot sizes and late probe pulse times, the probe witnesses a more highly developed plasma and was expected to undergo significant refraction. The spot sizes and probe delays used in the shots taken are summarized in Table 3-1.

Results. In discussing the results, we will use the pinhole labeling shown in Fig. 3-8. Figure 3-9 shows “shot” and “preshot” images from two experiments in which weakly refracting and highly refracting conditions were investigated. The different intensities of the various pinholes are due to imperfect centering of the backlighting beam on the back side of the array and is a normalizable effect. Shown particularly well in Fig. 3-9a is the positioning of the target with respect to the array where P1 and P2 lie about 10–30 μm behind the target, P3 and P4 are centered at about 30 and 50 μm from the front side and P5 – P10 are situated at different points in the blowoff region out to about 300 μm (the apparent thickness of the target, about 50 μm , is larger than the actual thickness, 22 μm , because of imperfect target rotation, target curvature, and diffraction.) The exact details of the relative positioning of the target with respect to the pinhole array vary from shot to shot as seen in Fig. 3-9.

As expected, for small irradiation spots and early probe pulse times (as in Shot 8919, 200 μm , 200 ps, Fig. 3-9b) distortion of the resulting

Fig. 3-8. Schematic of experiment showing the gold disk target, normally incident $0.53\text{-}\mu\text{m}$ plasma production beam, transversely incident $0.25\text{-}\mu\text{m}$ backlighter probe beam, and image relay optics. The inset depicts the positioning of the pinhole array image with respect to the planar target and the labeling of the pinholes.

Table 3-1. Summary of the spot sizes and probe delays used in the shots taken.

Shot No.	Spot Size	Probe Delay
8918	200 μm	600 ps
8919	200 μm	200 ps
8920	200 μm	600 ps
8921	400 μm	1200 ps

Fig. 3-9. Data from two shots. The "preshot" images were taken without firing the 0.53- μm plasma production beam for the purpose of establishing the relative positioning of the target and pinholes. The "shot" images were taken when the 0.53- μm beam was fired and the plasma was present. Weak refraction case ($D_{\text{spot}} \sim 200 \mu\text{m}$ and $\tau_{\text{probe}} \sim 200\text{ps}$) (a) preshot and (b) shot. Strong refraction case ($D_{\text{spot}} \sim 400 \mu\text{m}$ and $\tau_{\text{probe}} \sim 1200\text{ps}$) (c) preshot and (d) shot.

image of the pinhole array is undetectable. The most noticeable difference between the shot and preshot images is the apparent reduction in intensity and size of the point (pinhole) sources. This effect is greater for the axial pinhole row (the lower of the two rows in Fig. 3-9c and 3-9d). However, the fact that this attenuation is seen both in front and in back of the target, and does not significantly change for the pinholes farthest from the target surface leads us to believe that this is neither a refractive nor absorptive effect. Similar results were seen for Shot 8920 (no shot data obtained for Shot 8918).

For a large spot size and late probe pulse delay (Shot 8921, 400 μm , 1200 ps, Fig. 3-9d), several qualitative observations can be made. One sees that P4 is absent, while both P3 and P6 have apparent positions roughly 50 μm down the density gradient relative to their preshot positions. Also, P7 and P8 show significant distortions which are not present in their respective preshot images. This suggests that, as expected, refraction-induced image distortion is greater along the plasma-formation laser focal axis (even p's where the density is higher and the ray optical path is longer) than on the wings of the focal spot (odd p's). In particular, for perfect alignment of the array (even p's lying exactly on axis) one expects the apparent displacements on the pinhole images to be along $\hat{z}$ for axial positions and along both $\hat{r}$ and $\hat{z}$ for wing positions. This signature is indeed seen in Shot 8921 for pinholes P3 and P6.

These results were compared to expectations based on modeling of probe ray propagation through the laser-produced plasma. A 2-D LASNEX⁴⁰ simulation was run for the high refraction case (as in Shot 8921, Fig. 3-9d). The laser pulse was modeled as a 0.53- μm , 1.5-ns Gaussian pulse, super-Gaussian in space with a 400- μm spot diam, $E_{\text{laser}} = 150\text{J}$, and $I_{\text{peak}} = 10^{14}\text{W}/\text{cm}^2$. The simulated gold-disk plasma density profile was recorded at a simulation time corresponding to the experimental time of arrival of the 0.25- μm probe pulse ($t_{\text{delay}} \sim 1200\text{ps}$) and is shown in Fig. 3-10a. For simplicity in calculating probe-ray trajectories, an approximate analytic representation of the simulated profile was generated. The analytic profile, shown in Fig. 3-10b, matches the simulated profile quite closely for densities

where refraction effects are significant

$\frac{n}{n_{\text{crit}}} \geq 0.01$, where n_{crit} is the critical density for the probe light.

To model the probe light imaging system, we assumed point-pinholes and traced the marginal rays of the converging $f/10$ focal cone as they propagated into and through the cylindrically symmetric test plasma (we considered only rays which began and ended in a plane containing the symmetry axis of the plasma). Two examples are shown in Fig. 3-11 for the nominal $z = 50$ and $z = 150\text{-}\mu\text{m}$ axial pinholes (P4 and P6). Upon exiting the plasma boundary (chosen such that the plasma density was many orders of magnitude below the critical density for the probe light), the rays' final positions and angles were used to calculate the apparent axial source points for each ray. The separation of these apparent source points was taken to be the point spread (i.e., blurring) error of the imaging system (optics plus plasma relay medium) at that axial location. Roughly speaking, blurring of the pinhole image at recording-film plane becomes significant when the point spread exceeds the actual pinhole diam (25 μm). In Fig. 3-12 we plot (for the assumed profile) the calculated on-axis point spreads as a function of the axial position of the pinhole. The point spread is seen to decrease gradually with z for $z > 50$ but increases abruptly in the vicinity of the steepened region of the density profile, $z = 45\text{ }\mu\text{m}$. For $z < 40$, the refractive distortion is so severe it makes little sense to speak of a focus at all. This is not terribly surprising as the $-\hat{z}$ -propagating marginal ray witnesses a much more non-uniform plasma and is correspondingly more severely refracted than its $+\hat{z}$ -propagating complement. In the case on pinhole P4 ($z \sim 30\text{--}40$), the image spot diameter calculated to blur to at least 60 μm . Thus the fact that P4 is missing in the shot image of Fig. 3-9d might be explained by the factor of 1/3 decrease in exposure (factor of three increase in the exposure area; little blurring occurs transverse to $\hat{z}$ -axis) on the film. Less clear is the apparent reduction in the apparent spot size of P6 (Fig. 3-9d), as Fig. 3-12 predicts less than 10 μm of blurring at the axial position of this pinhole ($z = 100\text{--}150\text{ }\mu\text{m}$). Image degradation is also expected when rays to be collected by the imaging optic are

Fig. 3-10. (a) LASNEX-simulated plasma density profile for the high-refraction case of Shot 8921 (Fig. 3-9d). (b) Analytic representation of LASNEX-like profile used in ray trajectory computations. Note that profile steepening begins at about $z \sim 50 \mu\text{m}$.

Fig. 3-11. Ray trajectories of the central and two marginal rays of the $f/10$ relay optic cone for the nominal $z \sim 50 \mu\text{m}$ and $z \sim 150 \mu\text{m}$ pinholes (P4 and P6 in Fig. 3-8b) superimposed on the model density profile contours.

Fig. 3-12. Calculated point spreads along the plasma symmetry axis of the relayed pinhole images. Image blurring can be expected when the point spread errors approach the pinhole diameter ($25\ \mu\text{m}$).

refracted (i.e., bent) to the extent that they miss that optic altogether. In Fig. 3-13 we have plotted the calculated z -dependent refraction angles for the central and two marginal rays of the incident $f/10$ focal cone for a pinhole located at $z = 40 \mu\text{m}$. Roughly speaking, about two thirds of the image light for the P4 pinhole ($z \sim 30\text{--}40 \mu\text{m}$) is calculated to be refracted out of the $f/2$ collection cone of the imaging optic. Note however that this effect does not explain the poor quality of the P6 pinhole image (cf. Fig. 3-9d and Fig. 3-12 for $z \sim 100\text{--}150 \mu\text{m}$). We also note that the apparent displacement of pinhole P6 (about $50 \mu\text{m}$ along $+\hat{z}$ in Fig. 3-9d) is substantially more than the few microns predicted by our analysis. In general the experimentally observed distortion, blurring, and displacements appear greater than predicted for the assumed analytic profile.

We qualitatively note that a third possible image-distortion mechanism is scattering of light off of laser driven ion acoustic fluctuations or turbulence.⁴¹ These disturbances can be driven when hot spots in the high intensity beam (the $0.53\text{-}\mu\text{m}$ -beam in our case) filament and locally drive such processes as streaming instabilities or stimulated Brillouin scattering. Such filaments were observed to give rise to large refractive effects for the plasma-production beam due to its long copropagation distance with the filaments and their attendant ion acoustic perturbations. However in the case of low intensity probe light incident transverse to such filaments, we believe that, in a path-integrated sense, the distortions due to refraction in the background profile dominate those due to scattering from ion fluctuations, particularly in the vicinity of the steepened region. Quantitatively, the relative roles of this and the above described image distortion mechanisms are as yet unknown.

Of more specific interest in laser-plasma interaction experiments is how refraction will affect 2ω and $3/2 \omega$ emissions images. The densities at which these emissions are born, expressed as fractions of their respective critical densities, are quarter-critical for 2ω and ninth-critical for $3/2 \omega$. It is tempting to infer the extent of image degradation for $3/2 \omega$ and 2ω light from the present data by looking at the image quality in the data for axial positions with calculated densities corresponding to these fractions of the critical-density for the backlighter probe. This, however, neglects the

importance of density profile steepening⁴² in the vicinity of the quarter-critical layer of the plasma-producing light ($0.53\text{-}\mu\text{m}$ in this case). In the experiment the ninth- and quarter-critical surfaces for the probe ($n_e \sim 2 \times 10^{21}$ and $n_e \sim 4 \times 10^{21} \text{ cm}^{-3}$ respectively) lie well within the steepened region of the profile (cf. Fig. 3-10). However, although the source region for 2ω emissions ($n_e \sim 4 \times 10^{21} \text{ cm}^{-3}$) also lies within the steepened region, the bulk of the $3/2 \omega$ emission source ($n_e \sim 10^{21} \text{ cm}^{-3}$) lies just outside the steepened region in a relatively flat portion of the profile ($z \sim 50 \mu\text{m}$ for the assumed profile of Fig. 3-10). Raytraces for this case show moderate blurring (point spread $\sim 12 \mu\text{m}$) for the $3/2 \omega$ source point and significant for the 2ω source point (point spread $\approx 60 \mu\text{m}$). It should be noted that blurring of $3/2 \omega$ images can be a large effect in the limit of long transverse plasma dimensions ($L_{\text{trans}} \sim \text{few} \times 10^3$). However, for the plasma profile assumed here, comparison of Figs. 3-12 and 3-13 suggests that, even for the relatively fast $f/2$ imaging system used in these experiments, image degradation due to light loss (when rays are refracted out of the collection optic acceptance cone) dominates that due to focal blurring.

Summary. A simple experiment was performed in an attempt to assess the extent to which refraction-induced image distortion can be predicted and understood in situations where harmonic emissions can provide information useful in the characterization of laser-produced plasma density profiles. The data, in the form of $0.25 \mu\text{m}$ light images of a well defined spatial structure (pinhole array), exhibits the expected qualitative features in the limits of large and small plasma dimensions. More detailed analysis of ray trajectories in LASNEX-simulated density profiles indicates that source-region image quality is sensitive to the details of the plasma profile when the source is located near the steepened region of the profile (typically near the quarter-critical density layer). While this in principle presents no problems for probing-diagnostic measurements, it can be a serious limitation when imaging $3/2 \omega$ emissions in the increasingly common case where the plasma dimensions exceed several thousand laser wavelengths. Future work will aim to observe these effects with $0.35 \mu\text{m}$ irradiation and to investigate the possibility of making use of refractive distortions as a density profile diagnostic.

Fig. 3-13. Refraction angles of the central and two marginal rays of the $f/10$ relay optic cone plotted against the axial location of the pinhole defining that cone. Areas which lie between the marginal ray curves (dot and dash-dot) and above the $f/2$ half-angle line represent image light which is lost due to refraction out of the acceptance cone of the $f/2$ imaging optic.

Authors: C. Darrow, D. S. Montgomery, K. Estabrook, and R. P. Drake; G. E. Busch and E. F. Gabl, (KMS Fusion Inc.).

3.2.5. Plasma Evolution from Laser-Driven Gold Disks

I. Introduction. Analytic models give rough approximations of the plasma evolution from laser-driven disk targets.⁴³ More accurate modeling is provided by computer codes; here we use the LASNEX code and model the plasma blowoff from laser-driven gold disks in a series of simulations. The modeling of the plasma shows the presence of two regions, the *transition* and the *corona*. The transition region is characterized by electron densities $n_e > n_c/4$, where the so-called critical density $n_c \equiv m\omega_L^2/4\pi e^2$, with ω_L the laser frequency; $n_c \approx 4 \times 10^{21} \text{ cm}^{-3}$ for 0.53- μm laser illumination. Computational results show that the transition region is close to steady-state and is driven away from the disk relatively slowly. It is the source of (and can be experimentally identified by) the strongest M-band x-ray emission (2–3.5 keV). The corona region is the underdense plasma where $n_e < n_c/4$, and is predicted to consist of both steady-state and self-similar segments. Corona temperature can be inferred from the high-energy x-ray continuum. (We assume throughout this article that the electron velocity distribution is Maxwellian.) The preliminary modeling suggested diagnostic techniques [time- and space-resolved imaging of M-band x-ray emission and time- and space-averaged measurements of high-energy (3.5–20 keV) x-ray spectra] that were later used (along with holographic interferometry) in experiments at KMS Fusion.

In the experiments, gold disks were irradiated with 0.53- μm laser light. Several laser temporal profiles were used, including 0.85-ns square pulses, 2-ns square pulses, picket pulse trains, and pulses with graduated leading edges; we refer to these pulse shapes as *square*, *2-ns square*, *pickets+square*, and *t-squared*, respectively. The experiments produced time-resolved measurements of the underdense ($<10^{21} \text{ cm}^{-3}$) corona plasma electron-density profile with four-frame holography, time- and space-resolved images of the M-band x-ray emission, and time- and space-averaged measurements of the high-energy (3.5–20 keV) x-ray spectra. We compared the LASNEX results with the experimental data.

For the electron-density profiles, agreement between simulation and experiment was generally good during the early portion of all pulse shapes studied, and the electron-density profiles contain both steady-state and self-similar segments as expected. During long-pulse illumination, however, the simulation increasingly overestimates the electron density, because LASNEX appears to obtain an inaccurate solution for the steady-state electron-density profile in the corona. The trajectory of the peak M-band x-ray emission is used to infer the motion of the transition region (in the electron-density range $n_c/2 > n_e > n_c/4$); simulation and experiment show excellent ($<10 \mu\text{m}$) agreement for this motion for picket pulse shapes, for which the comparison is straightforward. The corona electron temperature is obtained in a time- and space-averaged sense by examining high-energy x-ray spectra. Simulations that do not limit the electron heat conduction (below classical values) are in good agreement with the experiment.

II. Computational Model. The modeling uses a one-dimensional zoning scheme used extensively in previous work.^{44–48} Figure 3-14 shows the disk that is modeled, a spherical cap of gold with a radius of curvature twice the laser spot $1/e$ diam. The fixed curvature of the model approximates the three-dimensional effects of plasma evolution from the center of the actual planar disk target. The cap is 30- μm -thick and is finely resolved with many zones in the radial direction. Particularly fine resolution (80 zones) is used in the outermost 0.5- μm -thick region, which is also modeled with non-local thermodynamic equilibrium (non-LTE) atomic physics.

X radiation from the disk is calculated with a modified hydrogenic average-atom model (XSN), which is similar to that used in previous work.^{46–49} The radiation frequency structure is defined with 120 photon bins nonuniformly spaced from 0 to 20 keV. Radiation is transported spatially by summing over a set of spherical-harmonic basis functions (P_n transport). The bound-bound and bound-free radiative rates in the usual XSN model were reduced by factors of 2 and 4, respectively, to give a better match to experimental measurements of gold M-band conversion efficiencies.

The laser illumination was approximated in one dimension with a single incident ray. For the gold-disk blowoff experiments, the spatial

Fig. 3-14. Zoning and geometry used in one-dimensional model of a laser-illuminated disk.

beam profile of the KMS Chroma laser was approximately Gaussian and had a FWHM diameter of 200 mm. Large-amplitude, fine-scale variations in the profile make it difficult to assign a nominal intensity to the irradiation. To bracket possible experimental uncertainties in the illumination, simulations were carried out at three intensities: nominal (3×10^{14} W/cm²), one-third nominal ("low"), and three times nominal ("high").

The flux limiter f affects the balance between electron thermal energy in the corona plasma and x-ray generation in the high-density gold. Simulations were therefore carried out with two flux limiters to determine the effects of this parameter on the computational results: at nominal intensity, a flux limiter $f = 3.02$ corresponds to strongly flux-limited diffusion; $f = 0.1$ corresponds to little or no reduction in the thermal flux. A flux limiter of 0.1 is believed to approximate the heat flow actually produced by Coulomb collisions within the plasma.⁵⁰ A recent study⁵¹ of microwave-heated plasma obtains the best match between LASNEX and experimental temperature results with $f = 0.15$. However, previous computational work uses smaller values ($f = 0.01$ – 0.05) to obtain a better match to total x-ray conversion efficiency and laser absorption at high intensities ($>10^{15}$ W/cm²) and/or shorter pulse lengths (~ 100 ps).^{44, 48, 49}

Typical examples of four laser pulse shapes were modeled. Figure 3-15 shows the square, pickets+square, and t -squared pulse shapes; each has a peak intensity $I_0 \approx 3 \times 10^{14}$ W/cm². For each pulse shape, five code runs were carried out: two at nominal intensity, with $f = 0.02$ and 0.1 ; two at high intensity, with the same flux limiters; and one at low intensity, with $f = 0.1$ only. (At low intensity, flux inhibition is inconsequential at both $f = 0.1$ and $f = 0.02$.) To determine some of the steady-state features of the computational model, an experimental 2-ns square pulse and generic (flat-topped) 5-ns pulse were modeled at intensities of 1.5 and 3×10^{14} W/cm², respectively, both with $f = 0.1$.

III. Electron-Density Profiles. We now examine the electron-density profiles produced by laser irradiation of gold disks. The general features of the computational electron-density spatial profile, and how it evolves in time, are discussed first. Next, computational results are compared with experimental measurements for

three pulse shapes. Finally, the trajectories of $n_c/10$ and $n_c/40$ surfaces are examined to reveal the transition from self-similar to steady-state behavior; LASNEX results are again compared with experimental data. LASNEX obtains a good estimate of the early-time self-similar behavior, but it tends to overestimate the electron density in the later, steady-state segment of the profile.

General Features. Computational results show that during an idealized flat-topped laser drive, the plasma blowoff beyond the original disk surface can be divided into the two regions shown in Fig. 3-16. The transition (or steepened) region occurs near the critical density. It is characterized by rapidly varying electron density and temperature, and the entire region moves away from the initial surface of the disk relatively slowly. Other work has used steady-state analytic solutions of Euler's equations to approximate the fluid behavior in the subsonic portion of the transition region, in which $n_e > n_c$.⁴³ Realistically, the transition region must be considered to extend to lower densities than n_c , where inverse-bremsstrahlung absorption plays a dominant role; simulation results indicate an approximate outer boundary of the transition region to be near $n_c/4$.

The corona region is the hot, isothermal, lower-density plasma beyond the transition region. During steady laser illumination, the corona electron-density profile is initially well described by a single self-similar solution, as expected in a planar isothermal expansion. Later, as the plasma evolution continues, the three-dimensional nature of the expansion becomes important. A steady-state segment of the corona develops, extending outward from the transition region and connecting to a self-similar segment. The spatial extent of the steady-state segment increases with time.

Qualitatively similar behavior has been described theoretically for one-dimensional spherical isothermal rarefactions,⁵² but these results do not give an accurate match to the computational electron-density and fluid-velocity spatial profiles. The self-similar parts of the computational solution are approximately described by

$$\log \frac{n_e}{n_0} \approx -\frac{\sqrt{2} (r - r_0)}{C_s(t - t_0)}, \quad (1)$$

Intensity (10^{14} W/cm²)

Fig. 3-15. Laser pulse shapes used in disk simulations. Arrows indicate holographic frame times.

Fig. 3-16. During steady laser illumination the blowoff plasma consists of transition and corona regions. An overlay of electron-density profiles from various times shows that the corona contains both steady-state and self-similar segments.

where n_0 and r_0 are electron density and radius of the break between steady-state and self-similar solutions, and t_0 corresponds to the start of the laser-driven expansion (chosen here to be the time at which the laser drive intensity reaches half of its maximum value). The sound

speed $C_s = \sqrt{\bar{Z} T_e / M_i}$ where $\bar{Z}$ is the average charge state, T_e the corona temperature, and M_i the ion mass.

Equation (1) deviates from the self-similar solutions of a planar isothermal rarefaction by the factor of $\sqrt{2}$ in the logarithmic slope. Theoretical calculations⁵² indicate that the logarithmic slope for a one-dimensional spherical isothermal rarefaction is increased by a factor of $\sqrt{3}$ over the planar case, thus suggesting that the curvature of our quasi-spherical geometry causes the observed deviation. Because of this deviation and the presence of both steady-state and self-similar segments in the corona, it is difficult to determine the sound speed, and thereby the corona temperature, from the logarithmic slope of the electron-density profiles.

Spatial Profiles. Turning now to the experimental results shown in Fig. 3-17, we first examine electron-density profiles produced by the square laser pulse shape of Fig. 3-15(a). Each experimental profile corresponds to a holographic snapshot ("frame") taken at a particular time; Fig. 3-15(a) shows these times relative to the laser pulse. In addition, the location of peak M-band emission is plotted at an electron density $n_e/n_c = 0.4 \pm 0.2$ (see Sec. IV). In general, the experimental data are best fitted by the nominal or 1/3-nominal intensity simulations. In the computational results, thermal flux limiting has the effect of increasing the steepening of the transition-region profile: In the corona, flux limiting causes only a slight reduction in the slope of the steady-state profile and does not affect the self-similar segment. As a result, the experimental electron-density data do not conclusively distinguish between the flux-limited and non-flux-limited models. The corona temperature data discussed in Sec. IV indicate that the non-flux-limited model is the better match to the experiment, however.

Figure 3-17(a) shows all five simulation results; in the other comparisons in Fig. 3-17 (and in the comparisons of Figs. 3-18 and 3-19) we plot only the nominal and 1/3-nominal intensity

modeling, both without artificial flux limiting ($f = 0.1$). Consider now a frame-by-frame comparison of the LASNEX calculations with the data. Just after the leading edge of the drive pulse, the first electron-density profile shows a good match between experiment and the nominal modeling. However the next two frames show an increasing electron-density offset. At the trailing edge of the laser pulse [Fig. 3-17(c)], the nominal simulation is too high by a factor of 5. Despite the offset, the slopes of the experimental and nominal-simulation profiles agree well in all three frames. The comparison suggests that LASNEX is obtaining an inaccurate result for the steady-state corona electron-density profile, but is managing the self-similar isothermal blowoff correctly.

Figure 3-18 shows the experimental and computational electron-density profiles produced by the pickets+square laser drive of Fig. 3-18(b). In Fig. 3-18(a), immediately after the first picket, the LASNEX result is too high by a factor of 2 to 4. Then, in Fig. 3-18(b), after a period of undriven coasting, the code results come into agreement with the experiment. The transition region is further from the disk surface than with the square laser drive. The code continues to track the experiment in Fig. 3-18(c). The transition region has moved closer to the disk surface between the second and third frames, as the electron-density profile once again steepens because of the renewed laser drive. This retrograde motion is discussed further in Sec. IV. In Fig. 3-18(d) it is again apparent that the LASNEX model tends to exaggerate the rate of plasma blowoff during steady laser illumination but in this case the 1/3-intensity model agrees well with the experiment. The experimental data confirm the presence of a kink in the calculated electron-density profile; these kinks form as a result of the discontinuous pickets in the drive.

Figure 3-19 shows that the agreement between experiment and theory for the t -squared data [Fig. 3-15(c)] is better. In Figs. 3-19(a-c), experiment and simulation agree to within the experimental uncertainty. The more continuous t -squared laser drive produces electron-density profiles that are comparatively smooth, without the kinks seen in the data of Fig. 3-18. Figure 3-19(d) shows the data from a lower intensity shot. Shot #8845 had 1/3 the nominal intensity and should therefore be compared with the 1/3-nominal intensity code result (dotted

Fig. 3-17. Experimental and computational electron-density profiles at the three times shown in Fig. 3-15(a) during the square laser pulse.

Fig. 3-18. Measured and calculated electron-density profiles at the four times shown in Fig. 3-15(b) during the pickets+square laser pulse. Nominal-model data are shown with the heavy, solid line (—); low-intensity-model data are shown with the light, dotted line (-----). Experimental data from two shots are indicated with markers: shot 8840 — $\square$; shot 8856 — $\blacklozenge$.

Fig. 3-19. Experimental and computational electron-density profiles at the four times shown in Fig. 3-15(c) during the t -squared laser pulse shape: (a) 520 ps, (b) 920 ps, (c) 1320 ps, (d) 1570 ps (1/3-intensity experimental data). Nominal-model data are shown with the heavy, solid line (—); low-intensity-model data are shown with the light, dotted line (-----). Experimental data from two shots are indicated with markers: shot 8850 — $\square$; shot 8845 — $\blacklozenge$.

line) in the figure. The tendency for LASNEX to increasingly overestimate the blowoff electron density during steady laser illumination is again apparent.

Electron-Density-Profile Trajectories. Figure 3-20 shows the trajectories of the $n_c/10$ and $n_c/40$ surfaces. Experimental measurements for square and 2-ns square pulses were obtained from holographic data compiled over several similar shots with different probe times; LASNEX data are shown for comparison. In this data, the transition from self similar to steady state is manifest as a saturation in the velocity of a particular electron-density-surface trajectory, and the experiments qualitatively verify the expected transition from self-similar to steady-state behavior. The computational results agree with the experiment only during the initial self-similar period. Later, theory and experiment diverge as LASNEX overestimates the time it takes for particular electron-density surface to reach a steady state. This observation is consistent with the discrepancies in the spatial profiles: LASNEX overestimates the electron density in the steady-state segment of the corona.

Points in the steady-state segment of the profile are not entirely stationary, but move outward slowly. The modeling predicts the velocity of the steady-state segment to be about 3×10^6 cm/s. The experimental velocity of $3\text{--}4 \times 10^6$ cm/s, obtained from the peak M-band trajectories, is in good agreement with this value.

IV. High-Energy X-ray Emission. The laser-illuminated gold produces an easily measurable x-ray flux at energies above 2 keV. Such radiation is useful as a diagnostic in two ways. We will first discuss the two methods and then give examples of the data.

The first diagnostic technique uses the location of peak M-band (2–3.5 keV) emission to determine the trajectory of the transition region. The spatial correlation between peak emission and the transition region can be discerned immediately from the simulations, but it is useful to examine the M-band emission in a semi-analytic fashion as well. Simulations show that the M-band emission opacity E_k scales as

$$E_k \propto \frac{n_e^{3/2}}{T_e^2 (h\nu)^{5/2}}, \quad (2)$$

where $h\nu$ is the photon energy. The differential photon emission rate is related to the emission opacity by the expression⁵³

$$d\epsilon \propto n_e E_k (h\nu)^2 \exp\left(-\frac{h\nu}{T_e}\right) d(h\nu). \quad (3)$$

By combining Eqs. (2) and (3) and integrating over the M-band, we obtain a scaling for the M-band photon-emission rate:

$$\dot{\epsilon} \propto \frac{n_e^{5/2}}{T_e^{3/2}} \left[\operatorname{erf}\left(\sqrt{\frac{3.5 \text{ keV}}{T_e}}\right) - \operatorname{erf}\left(\sqrt{\frac{2.0 \text{ keV}}{T_e}}\right) \right] \quad (4)$$

Equation (4) and the more detailed LASNEX calculations indicate that peak emission occurs at the leading edge of the transition region and at an electron density somewhat greater than critical (see Fig. 3-20). As one approaches the disk surface through the isothermal corona, the emission rises because of the increasing electron density. Then, close to the disk surface the temperature drops sharply so that $T_e \ll 2$ keV, and we then have the approximation

$$\dot{\epsilon} \approx \frac{n_e^{5/2}}{T_e^{3/2}} \exp\left(-\frac{3.5 \text{ keV}}{T_e}\right)^2. \quad (5)$$

Thus, it can be seen that the temperature drop dominates the electron-density rise, and the M-band emission rate falls off near the disk surface; the peak emission rate occurs in the transition region between the disk surface and the corona, where $n_e \approx 2n_c$. The plasma is optically thick in the M-band at this electron density, however, so the net emission parallel to the disk surface peaks at a lower electron density, typically between $n_c/4$ and $n_c/2$ depending on the laser intensity.

In the second diagnostic technique, the corona temperature can be obtained from the slope of the high-energy x-ray spectrum. At energies beyond the M-band ($h\nu > 3.5$ keV), calculations show that the spectrum produced by the corona region is dominated by the bound-free continuum (see Fig. 3-21). Care must be taken to shield the measurement to obtain the spectrum only from the corona plasma and not from the much brighter transition region, in which the spectrum is distorted by broadened M-band emission; it is also necessary to ignore the region of spectrum dominated by the L-band emissions (9–11 keV). In our simulations, most of the emission is produced during the period of peak drive intensity, so there is little difference

Fig. 3-20. Experimental and computational trajectories for $n_c/10$ and $n_c/40$ surfaces.

Fig. 3-21. The M-band x-ray emission rate peaks in the transition region.

between spectra integrated over a 200-ps interval at peak drive and one that is fully time-integrated. The temperature obtained from the shielded spectrum agrees well with the peak (equilibrium) corona temperature. The calculated corona temperature in Fig. 3-22 is 2.3 keV; the corresponding temperature inferred from the spectrum is 2.2 keV.

M-band X-ray Emission. It is interesting to compare the trajectories of observed and predicted peak M-band x-ray emission. Figure 3-23(a) shows the comparison for the pickets+square laser drive; the agreement is excellent for the nominal LASNEX model. Note the distinct retrograde motion between the onset of the second picket and main pulse, as observed in Sec. III. In the pickets+square waveform, the first picket produces a localized burst of x-ray emission that indicates the time-space origin on the streak-camera output; it is more difficult to determine the origin for the other laser waveforms. We can match the peak emission to the time of peak laser illumination, but no spatial fiducial is available, so an arbitrary spatial offset must be allowed. Figure 3-23(b) shows that even with arbitrary offset for the t -squared laser drive, the agreement between the experimental and nominal-model trajectories is poor. The retrograde motion predicted by the code is not observed. The 1/3-intensity model is in good agreement with the experiment, however; maximum deviations are about ± 10 mm.

Corona Temperature. Table 3-2 shows the simulated (at nominal intensity) and experimental corona temperatures for the various laser drive types. The code values are obtained from a simple exponential fit to a simulated time- and space-integrated corona-region x-ray spectrum in the ranges 5–9 and 12–20 keV (see Fig. 3-22). Experimental values were obtained similarly with filtered x-ray diodes; the estimated uncertainty is ± 0.5 keV. The LASNEX model that does not strongly limit the electron heat conduction (i.e., $f = 0.1$) provides the best match to the experiment in all cases. Because the measurements are time-integrated, there is little difference between the temperature estimates for the various laser pulse shapes; the high-energy spectrum is dominated by x-ray emission during the period of highest-intensity illumination in each case.

V. Summary and Conclusions. Computational modeling of the plasma blowoff from a gold disk

by laser illumination shows the presence of two regions: the transition and the corona. The transition region, approximately in steady state, is driven away from the disk relatively slowly. It contains, and can be experimentally identified by, the strongest M-band x-ray emission. The corona region is predicted to consist of both steady-state and self-similar segments. The equilibrium corona temperature T_e can be determined directly from the time-integrated high-energy x-ray continuum.

The computational predictions are in fairly good agreement with the experimental data. Information regarding both the spatial profile of the electron density and the equilibrium corona temperature has been obtained.

The trajectory of the peak M-band x-ray emission was used to infer the movement of the steepened electron-density profile. Simulation and experiment agree to within 10 mm for the pickets+square laser drive, and the measurements confirm details such as retrograde motion of the transition region late in the drive. Data from all the waveforms confirm that the transition region moves outward quite slowly, at about 3×10^6 cm/s. Computed electron-density profiles agree well with the experiment during the early, low-intensity portion of the laser drive, but LASNEX overestimates the plasma electron density during steady drive. Although experiments confirm the presence of both steady-state and self-similar segments in the corona electron-density profile, LASNEX obtains an inaccurate result for the steady-state segment. It would be useful to conduct additional long-pulse experiments to characterize the steady-state corona in more detail. Drake et al.⁵⁴ have observed that LASNEX simulations overestimate the rate of plasma blowoff in experiments with CH exploding-foil targets. It is likely that the code overestimates the mass-ablation rate, resulting in an erroneous boundary condition at the edge of the corona region. Further code calculations will be carried out to elucidate the source of the discrepancy.

For the corona electron temperature, LASNEX simulations that do not strongly limit the electron heat flux (i.e., calculations with $f = 0.1$) agree with the experiment to well within the $\pm 20\%$ accuracy of the experimental observation. The measurements correspond to the corona temperature during the period of peak laser illumination and therefore do not vary much for the three laser drive pulse shapes.

Fig. 3-22. A shielded measurement of the high-energy x-ray spectrum yields the corona temperature. For the square laser pulse shape shown here, similar spectra are obtained for time-resolved and time-integrated measurements.

Fig. 3-23. Predicted and observed peak M-band x-ray emission trajectories for (a) pickets+main laser drive (also shown for comparison is the 1/2-critical surface trajectory); (b) *t*-squared laser drive, including both nominal and low-intensity models.

Table 3–2. Plasma production beam spot diameters and probe delays for the experiments discussed in the text.

Pulse shape	LASNEX temperature (keV)		Experimental temperature (keV)
	f = 0.02	f = 0.10	
Square	3.3	2.3	2.4
Pickets + square	3.3	2.1	2.4
T-square	3.3	2.1	2.3

Future experiments will improve on these results with more accurate, time-resolved measurements of the high-energy x-ray spectrum.

Authors: D. Ress, L. J. Suter, E. F. Gabl, and B. H. Failor

3.2.6. Three-Halves Harmonic Emission From Two-Color Experiments

Introduction. Emission from laser-produced plasmas at half-harmonics of the incident laser frequency, ω_0 , is usually attributed to the two-plasmon decay (TPD) instability^{55–61} at one-fourth the critical density, where the critical density, n_c , is the electron plasma density at which the plasma frequency equals ω_0 . The emission at $3\omega_0/2$, in particular has been investigated for a number of years in both glass^{62–68} and CO₂^{69,70} laser-produced plasmas. The spectral feature at $3\omega_0/2$ typically appears as a pair of peaks, one blue-shifted and the other red-shifted with respect to $3\omega_0/2$. Although the production of $3\omega_0/2$ emission is still a subject of investigation, it is generally agreed that the blue-shifted peak is associated with TPD plasmons which propagate up the density gradient. Experiments generally see, however, a blue-shifted peak even in the direct backscatter direction; theories have been developed to explain this in terms of either plasmon or photon reflection at their critical surface.

In this article, we present evidence of backscattered blue-shifted photons in the absence of a critical surface. This is achieved by using a 0.53- μm wavelength laser to irradiate a plasma pre-formed by a 0.35- μm wavelength laser; the experiment has been designed such that the longer wavelength interaction beam sees a peak plasma density less than n_c , but greater than $n_c/4$ for 400–500 ps. We refer to these experiments as being “two color” because two different wavelength lasers were used. The main result of this experiment is that we observe both peaks from the underdense plasma, and they disappear at the same time, presumably when $n_c/4$ disappears. This is evidence that densities much higher than $n_c/4$ are responsible for the production of backscattered blue photons. This result is important because it constrains the theories used to predict the angular dependence of the $3\omega_0/2$ intensity and wavelength shift.

This problem was addressed earlier in the work of Aboites et al.⁶⁶ They monitored forward

and backscattered $3\omega_0/2$ light from thin burn-through foils; backscattered blue-shifted photons continued to be seen after the foil burned through and this result was interpreted as evidence that the $3\omega_0/2$ were not reflecting from higher densities. One problem with that experiment is that it is difficult to eliminate contributions from the edge of the laser spot where the target has not burned through. In fact, $3\omega_0/2$ emission was observed long after laser light was transmitted through the foil; the emission terminated only at the end of the laser pulse. It is also difficult to interpret the evolution of densities above $n_c/4$ in single-beam irradiations of a burn-through foil because the peak density (seen by the laser) drops, after the target burns through, from n_c to approximately $n_c/4$ on time scales short (~ 100 ps) compared to available diagnostics. These difficulties are avoided in the present experiment.

Additional evidence for a source of blue-shifted photons near $n_c/4$ was evidenced in recent experiments⁶⁸ with a short pulse, 1.06- μm laser-produced plasma. Three-halves emission was simultaneously spectrally and spatially resolved at 90° to the incident laser. Low frequency modulation in the peaks suggested that ion acoustic turbulence could provide a mechanism that could produce the broad angular distribution observed in other experiments. In addition, the blue peak emission appeared to be coming from well in front of the target, that provides additional evidence that reflection of the blue photon at high densities is not an important mechanism for backscattering blue photons. The present experiment shows more direct evidence for the location of the scattering location of the backscattered blue photons.

Description of the Experiment. The experiment used two beams of the Nova laser system.⁷¹ The first beam has a wavelength of 0.35 μm and forms a plasma by irradiating a thin CH foil target. The pulse has a 100-ps rise and fall time, a flat top temporally (within 20%), and a full width at half maximum of 2 ns. The angle of incidence is 50° from the target normal. The spot area was kept at $6.7 \times 10^{-3} \text{ cm}^2$ throughout the experiment. The energy of the laser beam was chosen (based on earlier experiments) to be 700–800 J; this gave an intensity of $\sim 1 \times 10^{14} \text{ W/cm}^2$ which was sufficient to burn through a 3- μm -thick CH foil 100 ps before the end of the laser pulse.

The plasma-forming pulse is followed by a second pulse of 1 ns (square) duration and a wavelength of 0.53 μm . This beam was incident normal to the target. The spot diameter was kept at 350 μm throughout the experiment; the laser spot from the second beam fit within that of the first beam. The energy was 1800 J for a laser intensity of $2 \times 10^{15} \text{ W/cm}^2$. The density (relative to its critical density) seen by the second beam is higher than that of the first pulse by the square of the ratio of the laser wavelength ($= 9/4$). The start of the second pulse is delayed 100 ps relative to the end of the first pulse. As we will see later, the peak plasma density will take 400–500 ps to decay to quarter-critical (for the second beam) after the foil burns through. Thus there is a substantial amount of time (relative to the temporal resolution of the diagnostic) in which there is a quarter-critical surface but no critical surface.

The backscattered $3\omega_0/2$ light from the green beam was collected by a lens located 153° from the beam and 19° from the plane of polarization. Scattered Raman light from the green beam was scattered at 153° from the laser beam and 52° from the plane of polarization. Each signal was passed through 1/4-meter Czerny-Turner spectrometer to a streak camera located at the exit plane. The temporal and spectral resolution for both diagnostics are 50 ps and 10 $\text{\AA}$ respectively. Fiducials for both wavelength and time were provided for both diagnostics; calibrated streak camera sweep speeds then allow us to compare the time behavior of the different spectra.

After the foil burns through during the blue laser pulse, the peak density begins to decay with time. The time dependence of the peak plasma density seen by the green laser is inferred from the Raman spectrum. The Raman light is produced by incident laser light scattering from electron plasma waves with frequencies given by

$$\omega_e^2 = \omega_p^2 + 3k_e^2 v_e^2 \quad (1)$$

where ω_p is the electron plasma frequency, k_e is the electron plasma wave wave number, and v_e is the electron thermal velocity. Knowing the frequency of the scattered Raman photon, ω_R , the matching conditions $\omega_0 = \omega_e + \omega_R$, and $k_0 = k_e + k_R$ for direct backscatter⁷² we can infer the electron density at which the Raman photon is produced.

In Fig. 3-24, we compare Raman spectra from a 3- μm -thick CH foil and a 7- μm -thick CH foil. The two spectra are similar with the exception of the long wavelength cutoff of the Raman band. For the thicker foil the Raman emission nearly extends to $\omega_0/2$ indicating the presence of densities up to $n_c/4$ (for the green laser) throughout the pulse. For the thinner foil, the long-wavelength cutoff of the Raman moves towards shorter wavelengths which we interpret as the decrease of the peak plasma density below $n_c/4$ at late times.

We interpret the lowest contour in the upper right-hand corner of Fig. 3-24a as the peak plasma density reached by the laser at a given time. This gives a density versus time shown in Fig. 3-25 which compares favorably with the density evolution predicted by a 2-D LASNEX⁷³ simulation. It also correlates well with the abrupt drop in the $3\omega_0/2$ emission (at $t = 550$ ps in Fig. 3-24c).

There continues to be a small $3\omega_0/2$ signal at later times (see Fig. 3-24c) which must originate from low intensity regions near the edge of the laser spot, although the signal is essentially zero by 800 ps into the laser pulse. Our conclusions from Fig. 3-25 are, however, that the bulk of the plasma, and therefore the source for most of the Raman and $3\omega_0/2$, has a density below n_c for almost 500 ps before $n_c/4$ disappears.

If a substantial amount of the backscattered blue-shifted photons required a reflection from densities $>n_c/4$, then, based on our previous interpretation of the density evolution we would expect the level of the blue-shifted peak (relative to the red-shifted peak) to be substantially different from the 3- μm -thick foil (which is underdense) than from the 7- μm -thick foil (see Fig. 3-26b), the relative amplitudes of the two peaks are essentially the same. In Fig. 3-26, we have chosen to show the spectra at the time of peak emission, but we will see momentarily that the ratio of intensities is essentially constant for all times.

We next investigated the relative intensities of the two peaks as a function of time. In Figs. 3-27a and 3-27b, we present the data in the form of contour plots. Because of the large range in the relative intensities, the contours represent constant levels of film density. Spectrally integrating over the two peaks separately, and multiplying the relative intensity of the blue peak by a constant, gives Figs. 3-27c and 3-27d which show that the ratio

Fig. 3-24. Comparison of Raman spectra from a) a 3- μm -thick CH foil and b) a 7- μm -thick CH foil. The contours are levels of constant relative intensity and the increment between contours is constant and the same for both figures. Also shown are the wavelength-integrated $3\omega_0/2$ signals from the green laser for c) a 3- μm -thick CH foil and d) a 7- μm -thick CH foil. The time scale is chosen to start at an arbitrary interval before the laser pulse begins.

Fig. 3-25. Comparison between predicted (solid line) and experimental (data points) density decay after a 3- μm thick CH foil target has burned through (at time τ_{bt}). The target is irradiated with a $3 \times 10^{13} \text{ W/cm}^2$ and $\lambda_0 = 0.35\text{-}\mu\text{m}$ laser pulse. The experimental density decay ($\diamond$) is derived from the long wavelength cutoff of the Raman band and the datapoint ($\square$) comes from the $3\omega_0/2$ signal. The right-hand vertical scale shows the relative density seen by the 0.53- μm beam.

Fig. 3-26. Comparison of spectra from a) a 3- μm thick CH foil and b) a 7- μm -thick CH foil at the same times ($t = 550$ ps) during the laser pulse.

Fig. 3-27. Comparison of time-resolved three-halves harmonic emission from the 0.53- μm wavelength laser. We compare spectra versus time for a) a 3- μm -thick CH foil, and b) a 7- μm -thick CH foil. Each contour is a constant film density. We also compare the spectrally-integrated relative intensity of the red and blue-shifted components for the same foil thicknesses in c) and d). In both cases, the relative intensity of the blue-shifted component tracks the red-shifted intensity for the entire pulse duration.

of the relative intensities of the two peaks is similar for all times and for both foil thicknesses. This is consistent with a single instability (two-plasmon decay) as the source of both spectral features.

The important point is shown in Fig. 3-27c. For the 3- μ m-thick foil, the blue peak drops when the red peak drops. This is inconsistent with a production mechanism for the backscattered blue-shifted photons which requires reflection of the photons from a density near or above n_c .

Brief consideration shows that the blue-shifted TPD plasmons cannot propagate far from $n_c/4$ so we expect both peaks to disappear at about the same time. The plasmon propagating up the density gradient (and which produces the blue-shifted photon) has a reflection point near $n_c/4$. To calculate the exact density we need to estimate the wave number of this plasmon. From Ref. 60, the most unstable plasmon has a wave number component k_y perpendicular to the density gradient given by

$$k_y = k_0 \left(\frac{1}{4\alpha\beta} \right)^{1/3} \quad (2)$$

with

$$\alpha = 0.08 L_\mu I_{14}^{1/2}, \text{ and } \beta = 1.41 T_{keV}^2 / I_{14} \lambda_\mu^2,$$

where L_μ is the density gradient scale length in microns, I_{14} is the laser intensity in 1×10^{14} W/cm², T_{keV} is the electron temperature in keV, and λ_μ is the laser wavelength in microns. From the LASNEX simulation we estimate $L_\mu \sim 400$, and $T_{keV} \sim 1$, and the experiment gives $I_{14} \sim 10$ and $\lambda_\mu = 0.53$; these numbers give $k_y \sim 0.2k_0$. The component of the wave number parallel to the density gradient is given by:⁵⁸

$$k_x = \left(\frac{k_0}{2} \right) \left\{ \left[1 + 4 \left(\frac{k_y}{k_0} \right)^{1/2} \right] \pm 1 \right\}. \quad (3)$$

Referring to Eq. 1, we note that the plasmon begins at $n_c/4$ with $\omega_e = .52 \omega_0$. The usual assumption is that, as the plasmon propagates up the density gradient, the frequency and k_y are conserved; reflection of the plasmon occurs when $k_x = 0$. From the electron plasma wave dispersion relation we then find that the reflection point is

near $0.27 n_c$ which will disappear within 100 ps (from Fig. 3-25) before $n_c/4$ (and the red-shifted peak) disappears.

Wave number- and frequency-matching conditions require the TPD plasmons to propagate some distance before the scattering of the incident light photons from the plasma waves become efficient enough to produce observable amounts of $3\omega_0/2$ photons. Reflection of photons or plasmons can increase the range of angles over which the $3\omega_0/2$ light is scattered. Our results make it unlikely that the backscattered blue photons are produced by reflection of the photons near $9n_c/4$ or by interaction between the plasmons and laser light reflected from densities above $n_c/4$. This leaves either reflection of the blue plasmons near $n_c/4$ or scattering of blue plasmons or photons by ion turbulence.

In earlier work,⁶⁸ we showed evidence, in high-intensity irradiations of plasmas, of ion turbulence scattering either TPD plasmons or $3\omega_0/2$ photons. This picture of the connection between ion turbulence and the backscattered blue peak is consistent with the lower ratio of the blue-to-red intensities observed at shorter wavelengths (this article) and lower laser intensities.^{74,75} One would expect that if simple reflection of the blue photons were involved that the ratio of the intensities of the blue and red peaks would remain constant regardless of the laser intensity. Ion turbulence, on the other hand, will have a threshold laser intensity below which the density fluctuation level is small. We can estimate density fluctuation required for a 20% reflectivity (the ratio of the blue to red peak intensities in Fig. 3-26) from:⁷⁶

$$r = \tan^{-2} \left[\frac{\pi}{2} \frac{L_c}{\lambda_0} \frac{\delta n_e}{n_c} \left(1 - \frac{n_e}{n_c} \right)^{1/2} \right].$$

For fluctuations with a coherence length of $L_c = 400 \mu\text{m}$ and $n_c = 0.25n_e$, we can solve this equation of $\delta n_e/n_e$ to find $\delta n_e/n_e \sim 10\%$, which is a modest fluctuation level. If the fluctuations are considered to be incoherent, then the fluctuation level needs to be higher to achieve the same r , but then the extent of the fluctuations will be larger than L_c .

In summary, we have presented data from an experiment in which we compare $3\omega_0/2$ emission

from an overdense and an underdense plasma. Comparison between the Raman spectra and LASNEX simulations shows that for 3- μm -thick CH foils, the peak density should be below n_c throughout the second laser pulse. We find that for all cases and for all times, both the blue-shifted and the red-shifted features are proportional in intensity, and for the underdense target, they drop in intensity at the same time. This observation makes it likely that the backscattered blue-shifted photons are produced near $n_c/4$ and we have suggested earlier⁶⁶ that ion acoustic turbulence is a possible mechanism to explain this observation as well as other observations of the angular dependence of the $3\omega_0/2$ emission.

Authors: P. E. Young, H. A. Baldis, B. A. Remington, K. G. Estabrook

3.2.7. Scattered Light near the Laser Wavelength from Nova Two-Color Experiments

The stimulated Brillouin scattering instability in a plasma is the decay of an electromagnetic wave into an ion acoustic wave and a light wave.^{77,78} This instability is of concern to inertial confinement fusion targets because it can potentially scatter a substantial fraction of the incident light energy. Although observed scattered light levels in experiments with 0.35- μm -wavelength lasers have been small (<3%),^{79,81} prediction of the scattered light levels due to SBS in reactor targets is difficult because there are significant discrepancies between the observed scattered light levels and spectra, and those calculated assuming convective growth from thermal fluctuations. For example, experiments generally observed blue-shifted SBS light from laser-irradiated low Z targets which theory predicts should give red-shifted light.⁸⁰⁻⁸² The difference between theory and experiment could be resolved if an enhanced noise source was seeding SBS at densities different from the location where theory predicts the maximum convective gain will occur. Enhancement of ion fluctuations by light reflected from the critical surface, n_c , (the surface at which the laser frequency equals the electron plasma frequency) has been suggested theoretically^{83,84} and enhanced SBS with a seed beam has been observed in CO₂ laser-plasma interactions.⁸⁵ Another possibility is the two-plasmon decay instability^{86,87} at $n_c/4$, which can produce ion

waves of the proper frequency and wave number to enhance backscattered fundamental light.⁸⁸

We report evidence of seeding of SBS by TPD. We have studied scattered light produced by the interaction of a laser beam with a plasma which contains a quarter-critical density for ~ 400 ps, but no critical density. This condition is obtained by pre-forming a plasma with a beam of shorter wavelength than the interaction beam; we refer to these experiments as "two color" because two different wavelengths of laser light are used. Comparison of the time history of the spectrally-resolved backscattered light near λ_0 with the evolution of the peak plasma density inferred from the Raman spectrum and $3\omega_0/2$ emission, leads us to conclude that the ion wave fluctuation levels are highest in the density range $0.1n_c$ to $0.25n_c$. For target experiments where n_c is present, we see evidence of seeding by reflection of laser light from the critical density surface as well.

The present experimental technique allows us to correlate backscatter signal levels with the presence or absence of a particular plasma density; this is difficult with single laser beams incident on thin foil targets⁸¹ because, when the foil burns through, the peak density reached by the laser rapidly drops (within 100 ps) from n_c to $\sim n_c/4$. Since the wavelength shift, $\delta\lambda$, is related to the sound speed, c_s , and the expansion velocity, u , by^{79,89}

$$\frac{\delta\lambda}{\lambda_0} = 2 \left(\frac{c_s}{c} \right) \left(1 - \frac{n_c}{n} \right)^{1/2} (1 - M) \quad (1)$$

where $M = u/c_s$ is the Mach number. It has been difficult in the past to experimentally determine the location of the source of SBS in the density profile from scattered light spectra due to the lack of measurements of the coronal electron temperature and the expansion velocity profile. The location of the SBS source (for lasers with $\lambda_0 \leq 1 \mu\text{m}$) has previously been estimated by comparing the observed SBS wavelength shift with Mach number and density profiles predicted by computer simulations.⁸⁹ In another approach,⁹⁰ a solid target was tilted by an angle θ which prevented the incident laser light from penetrating beyond $n = n_c \cos^2 \theta$; however, that experiment, which found SBS originating at low densities, did not analyze the SBS spectrum.

The experiment that we performed used two beams of the Nova laser.⁹¹ One of the beams

forms a plasma by irradiating a thin CH foil target. The laser pulse is 2 ns (square) in duration, has a wavelength of $0.35\ \mu\text{m}$, and an energy of 700–800 J. The angle of incidence is 50° from the target normal. The average intensity is $1 \times 10^{14}\ \text{W}/\text{cm}^2$ and is chosen (based on earlier Nova experiments⁸¹) so that the foil burns through 100 ps before the end of the pulse.

The plasma-forming pulse is followed by a second, interaction pulse of 1 ns (square) duration, a wavelength of $0.53\ \mu\text{m}$ and an energy of 1800 J. The laser spot for the interaction beam fit within that of the plasma-forming beam had an intensity of $2 \times 10^{15}\ \text{W}/\text{cm}^2$, and was incident normal to the target. The density (relative to its critical density) seen by the interaction beam is higher than that of the first pulse by the square of the ratio of the laser wavelengths ($= 9/4$). The start of the second pulse is delayed 100 ps relative to the end of the first pulse. After the foil burns through, the peak plasma density will take 300–400 ps to decay to quarter-critical for the second beam (see Fig. 3-28). Thus there is a substantial amount of time (relative to the temporal resolution of the diagnostic) in which there is a quarter-critical surface but no critical surface.

Temporally and spectrally resolved data were obtained for several wavelength bands of light scattered by the second (green) pulse. A portion of the plasma light backscattered into the focusing lens was collected by a beamsplitter and analyzed with a spectrometer in a spectral region near the wavelength of the incident laser ($0.53\ \mu\text{m}$). The backscattered $3\omega_0/2$ light from the green beam was collected by a lens located 153° from the beam and 19° from the plane of polarization. Scattered Raman light from the green beam was collected at 153° from the laser beam and 52° from the plane of polarization. These latter two signals were each optically relayed through a 1/4-meter Czerny-Turner spectrometer to a streak camera. Fiducials for both wavelength and time were provided for all diagnostics; calibrated streak camera sweep speeds then allow us to compare the time behavior of the different spectra.

We will first discuss data from 3- μm -thick CH foils. These targets burned through near the end of the first laser pulse, so that the second pulse saw a plasma with a peak density less than n_c . The density time-dependence of the 3- μm foils was confirmed from the Raman spectrum; each wavelength of the spectral band

of Raman emission can be related to a unique density.⁹² If we assume that the long wavelength cutoff at late times (when the peak density is less than $n_c/4$) is indicative of the peak density seen by the laser,⁹³ then we infer the density time dependence plotted ($\bullet$) in Fig. 3-28. The density decay is consistent with the disappearance of $n_c/4$ (inferred from the $3\omega_0/2$ signal) which is also shown ($\square$) in Fig. 3-28; the experimental results are in agreement with the predictions of the 2-D LASNEX⁹⁴ code (solid line in Fig. 3-28) and confirms our belief that the second pulse does not see critical density.

An intensity contour plot of time and wavelength resolved SBS backscattered light from a 3- μm -thick foil, corresponding to the experimental data of Fig. 3-28, is shown in Fig. 3-29a. The backscattered light that is observed at the start of the laser pulse is correlated with the $3\omega_0/2$ emission. In Fig. 3-29b we compare a trace of the backscatter relative intensity at a fixed wavelength ($\delta\lambda = +3\ \text{\AA} \pm 1\ \text{\AA}$) with the wavelength-integrated time history of the $3\omega_0/2$ emission; both signals drop abruptly at the same time ($\tau_1/4$). This time correlation shows that the early time backscatter signal is seeded by TPD at $n_c/4$. Using $c_s = 3.3 \times 10^7\ \text{cm}/\text{sec}$ (from LASNEX predictions), $n_c/n_c = 0.4$, and Eq. 1, we find that the observed wavelength shift implies $0.4 < M < 1$ which brackets the LASNEX prediction of $M \approx 0.5$ at $n_c/4$, 200 ps into the laser pulse. Simulations have suggested previously⁸⁷ that plasmons produced by TPD (which also produce the $3\omega_0/2$ photons⁹⁵) at $n_c/4$ can beat together to produce ion fluctuations which can seed SBS. The ion waves produced at $n_c/4$ in CO₂-laser-plasma interactions were observed directly in an experiment using a Thomson scattering diagnostic.⁹⁶ This mechanism was invoked to explain a similarity in the dependence on laser intensity of time-integrated SBS and $3\omega_0/2$ emission as a function of laser intensity from CH foam targets.⁹⁷ This is also consistent with our earlier conclusions⁸¹ of the importance of seeding of SBS by TPD when $\lambda_0 = 0.35\ \mu\text{m}$.

The signal at late times can be explained by SBS growing at the peak of the density profile where the convective gain is maximized. The calculated wavelength shift for SBS at the peak of the profile is $\delta\lambda \approx 9\ \text{\AA}$, where we have used $M \approx 0$ (from LASNEX) and Eq. 1; this agrees with the wavelength shift of the signal observed at

Fig. 3-28. Comparison between predicted (solid line) and experimental (data points) density decay after a 3- μm -thick CH foil has burned through at time τ_{bt} . The experimental density decay ($\bullet$) is derived from the long-wavelength cutoff of the Raman band. The data point ($\square$) comes from the $3\omega_0/2$ signal and marks the time ($\tau_{1/4}$) at which the signal finally drops. The right-hand vertical scale shows the relative density seen by the 0.53- μm beam.

late times. We also note from Fig. 3-29a that the signal ends approximately 800 ps into the laser pulse; from Fig. 3-28 we see that this corresponds to the time when the peak plasma density drops below $0.1 n_c$.

Convective growth from thermal ion fluctuations could be responsible for the scattered light that originates at the peak of the density profile. The backscattered flux per unit solid angle per unit wavelength of light scattered from thermal ion fluctuations is given by Eq. 25 of Ref. 98, which we simplify to:

$$F_L = \frac{ck_B T_L}{\lambda_s^4} E'' \frac{\omega_s}{2\omega_{ia}} \frac{1}{(1 - \omega_p^2/\omega_s^2)^{1/2}} [e^G - 1], \quad (2)$$

where $E'' \approx 1$ is the transmission coefficient for light propagation from the interaction region to the vacuum; $T_L \approx T_e$, k_B is the Boltzmann constant; ω_s (λ_s) is the scattered light frequency (wavelength); ω_{ia} is the ion acoustic frequency; G is the convective gain; and we have chosen the case of direct backscatter. In our experiment, $\lambda_s \approx \lambda_0 = 0.53 \times 10^{-4}$ cm, $\omega_s/\omega_{ia} \sim 750$, $n_e/n_c \sim 0.15$, and $T_e \sim 2$ keV (from LASNEX). At the peak of the density profile, the gain is limited by the velocity gradient scale length, L_u . Midway into the green pulse, the velocity gradient scale length is large enough to produce substantial convective gain. LASNEX predicts $L_u \sim 300 \mu\text{m}$ in the latter half of the green pulse, for which we find a convective gain^{78, 81} of $G \approx 16$. This is in contrast to the case⁸¹ of $\lambda_0 = 0.35 \mu\text{m}$ in which the maximum gain was $G = 2\pi$. Using the numbers we have just listed, the flux through a solid angle of 0.1 steradian (solid angle of the focusing lens) and 10 Å bandwidth is approximately 1×10^{14} W/cm², which agrees well with the estimated experimental backscatter level of 5% ($\sim 1 \times 10^{14}$ W/cm²). We conclude that, at least at later times during the pulse, there is sufficient gain for SBS to grow from thermal noise at the peak of the density profile.

We also compared our results from 3- μm -thick foils with identical experiments except for the target thickness which was increased to 7 μm . In this latter case, a critical surface exists at the start of the green pulse. A contour plot of the backscattered SBS light is shown in Fig. 3-30a where we see that the signal now lasts for the entire laser pulse. Figure 3-30b shows that, at the same fixed wavelength as in Fig. 3-29, there

is a strong SBS signal as well as a $3\omega_0/2$ signal, for the entire duration of the laser pulse.

If light reflected from the critical surface were playing a dominant role in seeding SBS, then there should be a substantial difference in both the backscattered signal levels and wavelength shifts between the case with (7- μm -thick foil) and without (3- μm -thick foil) a critical surface exists at the start of the green pulse. The matching condition for this seeding mechanism requires that the difference between the incident laser and the reflected light frequencies equal the ion acoustic wave frequency. The reflected wave is frequency shifted from ω_0 by both the motion of the critical surface and the plasma flow between the critical surface and the matching region.^{81,99} LASNEX simulations indicate that the critical surface is moving away from the laser with $M_c \approx 0.3$; the red shift ($\delta\lambda \approx 4$ Å) produced by this motion dominates the blue shift produced by plasma flow. As shown in Fig. 3-31, which compares SBS light from both targets which is red-shifted by 7 Å, the thick target signal is much stronger and turns on much earlier than for the thin target, even though the gradient scale lengths in the first case should be much smaller.

In summary, our results demonstrate that noise levels have an important influence on SBS spectra and scattered light levels. For a plasma with a peak density less than n_c , the backscattered light near the incident laser wavelength originates in a density range $0.1 n_c < n < 0.25 n_c$. A portion of that signal is the result of seeding by the two-plasmon decay instability; the enhanced noise level allows significant back reflection even when the convective gain is low. The remainder of the signal can be attributed to convective growth of SBS from thermal noise at the peak of the density profile as the convective gain increases late in the pulse. For a plasma with a critical surface, we also see evidence that light reflected from critical density seeds SBS.

Authors: P. E. Young, H. A. Baldis, K. G. Estabrook

3.3. Hydro Theory

3.3.1. SIMPLE—A Radiation-Hydrodynamics Code for Small Computers

With the ever increasing power and availability of small computer systems, work

Fig. 3-29. a) Contours of relative intensity of time and wavelength resolved backscattered light near the laser frequency from a 3- μm -thick CH foil. b) Relative intensity of backscatter (—) versus time at a wavelength shift $\delta\lambda = +3 \text{ \AA} \pm 1 \text{ \AA}$, and three-halves harmonic emission (—).

Fig. 3-30. a) Contours of relative intensity of time and wavelength resolved backscattered light near the laser frequency from a 7- μm -thick CH foil. b) Relative intensity of backscatter (—) versus time at a wavelength shift $\delta\lambda = 3 \text{ Å} \pm 1 \text{ Å}$, and three-halves harmonic emission (—).

Fig. 3-31. Relative intensity of backscatter versus time at a wavelength shift of $\delta\lambda = 7 \text{ \AA}$ from a 3- μm -thick (—) and a 7- μm -thick (- - -) target.

stations, and personal computers, the possibility of doing simple yet instructive calculations at one's "desktop" has emerged. It was with this realization in mind that a code relevant to laser fusion was developed specifically for small computer systems. The SIMPLE hydrodynamics code¹⁰⁰ was constructed with two purposes in mind: (1) as an educational tool for researchers in inertial confinement fusion to learn how the many and complicated physical processes interact, and (2) as a tool for designing and understanding simple ICF experiments. It is not the intent of this simulation code to be used for detailed target design since many of the important physical processes are treated in a "simple" fashion. Indeed, it is due to this simplicity that the code is ideally suited for a desk-top system such as a Macintosh II. SIMPLE was designed from the "ground up" with the user in mind; hence, the input files are easy to create and the output files are readily interpreted. A casual user can begin to do meaningful calculations almost immediately. For the more serious user, there is available a variety of options to explore the details of the physics models and the effect of their change on target behavior.

The SIMPLE code is a one-dimensional, three-temperature, Lagrangian hydrodynamics and energy transport code. The electron, ion, and radiation components are treated separately in a fluid approximation, and are loosely coupled to each other. Each is in local thermodynamic equilibrium, namely, the electrons and the ions are described by Maxwell-Boltzmann statistics, and the radiation field is approximately Planckian. For the electron and ion fluids, heat conduction is handled through (single-group) diffusion. Radiation flux is transported by multigroup diffusion, which allows for slight departures from a pure Planckian. The present version of the code does not include realistic optical coefficients. The degree of ionization is described by a Saha model, while the thermodynamic quantities are derived either from an ideal gas model or from realistic equation-of-state tables. The hydrodynamic motion is treated in a Lagrangian formalism, namely, the mass in each zone is conserved.

The code has the ability to restart from a saved restart file. Thus, a problem may be run in successive segments, or terminated and then restarted after adjustment of selected physics parameters. The user must specify the time

dependent nature of the energy source(s) interacting with the target. The major external source of energy is one or more lasers. The user may also specify temperature, pressure, or velocity sources separately or in combination. To illustrate possible uses of the code, we next turn to two calculations relevant to typical experiments done at the Nova laser-fusion facility: (1), an exploding CH foil, and (2) a direct drive implosion.

The first calculation, the results of which are presented in Fig. 3-32, is for an exploding CH foil of the type used in recent low density, large scale length plasma instability experiments. In the experiments and the calculations, three 2 ns square, blue (pump) beams of combined intensity of 6×10^{14} W/cm² irradiate a CH foil on each side at an angle of 50° to the normal, exploding the foil and preforming a large scale length, underdense plasma. A 2 ns square green beam of intensity 2×10^{14} W/cm² is then introduced normally to the target after a 250 ps delay to probe the peak density decay through SRS backscattering. The question addressed with the calculations was whether the delayed probe beam was significantly perturbing the preformed plasma, as opposed to simply probing. In Fig. 3-32 the calculated electron density profiles, $n_e(z,t)$, are given in 250 ps intervals. (Note: the profiles are artificially offset one from another by factors of 10 for clarity of display.) The peak n_e drops below critical density after 250 ps, indicating that the foil has burned through. The figure shows both the results from the "2-color" calculation overlaid with the results of a calculation with the probe beam omitted. The conclusion reached was that the green beam was not significantly perturbing the density profiles, and indeed was acting only as a probe. This calculation took 8-1/2 min of CPU on the LEAF VAX 8650, and about a factor of 5 longer on a Mac II.

The second calculation, the results of which are shown in Fig. 3-33, represents a high yield, direct drive implosion of the type recently used in neutron diagnostics development shots. The experiments and calculation used 30 kJ of blue laser light in ten 1 ns square pulses incident on a 1-mm-diam glass capsule with a 2.5-μm wall thickness and filled with 25 atm of DT. In the figure, the calculated pressure profiles, $P(r,t)$, are shown in 50-ps intervals. The solid curves represent the fuel, and the dotted curves the pusher. The time, average fuel ion temperature,

Fig. 3-32. Electron density profiles, $n_e(z,t)$, in units of 10^{21} cm^{-3} at intervals of 250 ps for an exploding CH foil calculation. The initial foil thickness was $2 \mu\text{m}$. For clarity, the profiles are artificially offset one from another by factors of 10. The dotted curves correspond to a calculation with six blue pump beams, three incident on each side, which quickly burn through the foil and preform the plasma. The solid curves represent a similar calculation, only with an additional green probe beam introduced after a 250 ps delay. The peak density, $n_e^{\text{peak}} (\text{cm}^{-3})$, and time, $t (\text{ns})$, of each profile are noted on the figure.

Fig. 3-33. Pressure profiles, $P(r,t)$, in intervals of 50 ps for a direct drive implosion of a DT-filled glass capsule. The capsule had an initial diameter of 1 mm, wall thickness of 2.5 μm , and 25 atm fill pressure. For clarity, the profiles are artificially offset one from another by factors of 100. The dotted curves represent the pusher and the solid curves correspond to the fuel. The time, average fuel temperature, convergence, and neutron yield of each profile are noted on the figure.

convergence, and neutron yield are noted on each profile. The yield was calculated by incorporating a simple $\langle \sigma v \rangle_{DT} = f(n_D, n_T, T_i)$ prescription¹⁰¹ into SIMPLE. The burn was turned off in the inner 50% of the fuel to eliminate the large neutron emission due to the high temperature spike at the core due to the perfect spherical symmetry of a 1-D calculation. The strong shock front generated by the intense laser has just reached the pusher-fuel interface at about 0.55–0.60 ns, the origin at 0.75–0.80 ns, and has reached the pusher on the rebound at about 0.85 ns. If the calculation is truncated when the calculated neutron yield matches the observed yield of $Y_n = 2 \times 10^{13}$, then the predicted bang time is about 850 ps, which is not too far off from the observed bang time of 750 ps. Perhaps not coincidentally, at 850 ps in the calculation is when the shock on its rebound has reached, and presumably destroyed the pusher. This calculation took 22 min of CPU on the VAX 8650 (and a factor of 5 longer on a Mac II).

In conclusion, we have developed, installed, and tested a 1-dimensional, 3-fluid, LTE, hydrodynamics code for small computers. The code, SIMPLE, was written primarily as an educational tool, and as such, was designed specifically with the user in mind. Input files are easy to set up and calculated results are easy to interpret. Enough generality has been built into the code, though, that more in-depth studies of specific problems are possible. Due to the modular structure of SIMPLE, modifications and enhancements to the code are generally straightforward. Due to its ease of use, transparent physics, and availability on small computer systems, the SIMPLE code should be a useful tool for experimentalists and theorists alike.

Author: Bruce A. Remington

3.4. Hydro Experiments

3.4.1. Planar Foil Mix Experiments in Preheated Laser Accelerated Packages

When an interface between materials of different density is accelerated, small perturbations can grow into an intermixing of major proportions. For continuous acceleration this mixing is due to Rayleigh-Taylor (RT) instabilities,¹⁰² while in the case of shocked acceleration it is due to Richtmyer-Meshkov

(RM) instabilities.¹⁰³ We have undertaken and nearly completed a series of experiments investigating shock induced planar foil mix at the Nova laser where two distinct layers of material were driven together and the effects of instabilities induced at the preheated interface were measured. In these experiments we measure the penetration of each material into the other after a smooth planar interface is accelerated by the passage of a shock. This interpenetration is referred to as the mix width. The purpose of these experiments is to increase our understanding of plasma/hydrodynamic instabilities in a radiation driven and preheated environment. Under these conditions it is usually quite difficult to obtain quality data with sufficient statistics, but these measurements strive to do just that. Such data can be effectively used to improve the present computational models to the point where a high level of predictive capability and reliability is achieved.

Usually the density difference between the two materials is quoted in the context of the Atwood number α , where $\alpha = (\rho_H - \rho_L)/(\rho_H + \rho_L)$. Here ρ_H and ρ_L refer to the materials at the interface with the acceleration being in the direction from material L to material H. A large α implies a high degree of instability, whereas a small value of α implies negligible instability. Negative or reverse α situations should result in no instabilities at all for RT, but yield significant instabilities for RM cases!

Method. The planar foil package for a high α experiment must be carefully constructed of a lower density material pushing a higher density one. **Figure 3-34** shows the cross section of a typical disk package used in these experiments. The crucial interface is between the thin Mo layer and the thicker PPS (polyphenylene sulfide) layer. This interface is kept as smooth and distinct as possible for these measurements. The layered PPS/Mo section is fitted into a gold guard ring which is covered by a thin layer of paralyene-N. This plastic acts as a tamper that helps contain the gold. In some cases the tamper was also extended over the Mo layer.

The intense radiation drive that accelerates the mix package normal to its surface is produced by using 6 to 10 kJ of 0.35- μ m light in a 1.0 ns pulse width from eight of the ten beams of the Nova laser. As the PPS ablates, the ablation front is preceded by a shock front of about 60 Mbar in magnitude and 50 ps in rise time. This

Fig. 3-34. Cross section of the cylindrically symmetric final mix package design for the PFM1 series. This design exhibits no obvious dishing effects.

planar shock propagates through the interface driving the PPS ablator into the Mo payload. With the passage of the shock wave the package is accelerated to about half its final velocity and instabilities are induced at the PPS ablator/Mo payload interface. These instabilities result in perturbation growth at the interface which allow the materials to penetrate each other. This interpenetration results in a region of mixed materials. The time at which the mix width is measured is usually in reference to the point at which the drive begins unless otherwise stated. The other relevant time is when the shock reaches the material interface and drives the instability. In this experiment shock breakout usually occurs about 1.2 ns after the drive begins.

After the drive is initiated the initial decrease in density is due to preheat and expansion of the package. The passage of the shock then compresses the Mo density to its maximum value with any later recompression being due to reflected compression waves. After shock breakout at the rear of the foil, the package rapidly decompresses. At this point the package expands as it moves away from the drive source. Due to this decompression, the density rapidly decreases. After about 4 ns the density ($<0.2 \text{ g/cm}^3$) and opacity are such that the package can be probed with a 2 to 3 keV x-ray backlighter. Timing constraints on the Nova backlighter beams limit the time window in which we can probe the mix to about 9.0 ns. Figure 3-35 shows a calculation¹⁰⁴ of the mean density of the Mo payload as a function of time that reflects these features.

From the above it is clear that the mix width evolves in time in a fairly complicated fashion. One-dimensional hydrodynamic calculations shed some light on this evolution. Figure 3-36 is an example of a LASNEX calculation¹⁰⁵ of the mix development from the time the drive is turned on to 10 ns later for a 58.0- μm PPS / 1.8 μm Mo package. The interpenetration due to RM instabilities occurs at about 1.2 ns. The mix width is frozen in after shock breakout and subsequently expands with the package. This interpenetration is shown in the figure and can be compared with our measurements.

Experimental Technique. To measure the mix width we use the point projection spectroscopy.¹⁰⁶ After a suitable delay from the drive pulse a Nova beam with a pulse length of

about 150 ps produces a point-like x-ray source (backlighter) which images the mix package onto a Bragg crystal spectrometer. The crystal acts as a mirror and combined with the point-like source projects an image of the mix package onto the focal plane as shown in Fig. 3-37. There is energy dispersion along one axis of the focal plane while the other axis is purely spatial. The two-dimensional image allows us to spatially analyze the mix, while the absorption features of the ablator and payload quantify the material interpenetration.

The absorption features used in the experiment are shown schematically in Fig. 3-38. From these features we determine penetration values from the detectable limits of the S K-edge into the Mo and of the Mo L absorption lines into the PPS. The extent of this interpenetration defines the total mix width. On each shot there are two data channels like that shown in Fig. 3-37. With this double data channel/double backlighter capability we can obtain two views of the mix package at the same time, or by staggering the backlighter delay, two views at different times can be obtained. The option of modifying one of the spectrometers to get a different perspective of the mix also exists (as is explained below). With double backlighting, these experiments are one of the most complicated performed at Nova.

This series of planar foil mix experiments (PFM1/LoRo) took many Nova shots to develop into a useful experiment. The early developmental shots dealt with such issues as tilting and dishing of the mix package, effects which tend to mimic real instability induced the mix. However, the double backlighting technique presently used and the fact that point projection spectroscopy yields a two-dimensional image of the package, made these effects obvious in one or both data channels if they existed. Refinements of the mix package design (to that shown in Fig. 3-34) minimized the above effects and they are no longer considered a problem.

Since that point we have successfully investigated high mix packages, low mix packages, and lateral expansion (edge) effects. It became necessary to reduce the drive energy to obviate possible effects of recompression of the mix region due to reflected shocks in the package, so there are data at higher and lower mean drive temperatures. The PFM1/LoRo experiment is nearly completed. We performed 18 shots this year alone, nearly completing the

Fig. 3-35. Calculation¹⁰⁴ of the mean density of the Mo payload for a PyN tamped target as a function of time.

Fig. 3-36. Dynamic simulation of the experiment using the LASNEX computer code for an untamped mix package.¹⁰⁵ Note the interpenetration limits of the PPS and Mo.

Point projection spectroscopy (PPS) is the technique used to backlight and probe the unstable ablator foil interface region in the NOVA foil mix experiment

Fig. 3-37. Schematic of Nova planar foil mix experiment. Point projection spectroscopy is the technique used here to backlight and probe the unstable ablator/foil interface region in the shock accelerated mix package. The actual experiment actually uses two of these diagnostic arms during each shot.

Normal

Fig. 3-38. Schematic of the point projection spectroscopy image for the normal spectrometer. Here the spectral features are used to quantify the mix region which is experimentally defined as the detectable limit of penetration of the Mo to that of the S.

PFM1 series of experiments. What remains to be done are several definitive null shots at the lower drive temperature and undispersed imaging of the mix package. The issues discussed below are preliminary results.

High Mix Measurements. The quantity α is large for the PPS/Mo package, and these targets are the basis for our high mix experiments. The package was accelerated by radiation drive in a high and low range corresponding to total mean laser energies of 5.3 and 6.9 kJ respectively. We have high mix data in the 5.0 to 8.5 ns regime for the lower drive and from 6.0 to 8.9 ns for the higher drive. This data consistently shows over 100 μm of mix as defined by the material interpenetration discussed above. Of this we can attribute about 30 to 40 μm to experimental effects (motion blurring, finite backlighter source size, digital smoothing, etc.), which leaves about 60 μm of real measurable mix. To first order, the calculated mix width^{104, 105} agrees with our measurements, but to constrain various models it is important to systematically determine the mix width within the time window where we can take data. This is also necessary to understand the physics and make a comparison with the calculations.

Our present mix data clearly shows the accelerated package and the strong absorption features of the Mo and S. The data even show the energy shift of the S line as it crosses the hot ablation front that changes the opacity of the PPS. **Figures 3-39a and 3-39b** are reduced image data from two shots at early (6.0 ns) and late (8.5 ns) times. The bent object in the image is a fiducial used to check the space scale. These images have been enhanced with false color set so that dark blue represents a transmission of one. **Figures 3-39a and 3-39b** are the corresponding areal density distributions of the ablator and payload plotted as a function of the distance from the original position of the package (displacement). The conversion from transmission to areal density is made assuming cold opacities. The limits of detection for the Mo and PPS are indicated by the circular and diamond data points respectively. The low density tail of PPS extending into the Mo seen for both shots is probably due to edge effects¹⁰⁷ and not shock induced mix. The measured mix widths are indicated in the figure.

Control Experiments. Low Mix Experiments. In order to determine the experimental resolution we also performed low mix experiments using low

α packages where the mix due to instabilities is calculated to be less than 20 μm . In these targets the Mo layer shown in **Fig. 3-34** is replaced with a thicker (15 μm) chlorinated plastic payload (polychlorostyrene) of density only 6% from that of the PPS ablator. As the low mix experiments should show a fairly clean interface (negligible mix overlap), they are also useful in confirming the value of the radiation drive. As the interface is smeared into a mix region, a cleaner interface provides a better defined reference on the upper limit of the PPS displacement. This can be used in a calculation to determine the drive it would take to accelerate the material to that location in time. As this can only be done to the accuracy of the location of the interface, the low mix experiments are important.

The measured mix width in some of the low mix shots is less than 40 μm , which is close to the system resolution. However, most of the low mix data is presently being analyzed and we still require data at the lower drive to make a definitive statement. If all the mix measured in these shots can be attributed to system resolution, then this control experiment suggests that the high mix data is not an artifact of edge or other spurious effects. However, due to drastic differences in the hydrodynamic behavior of the low and high mix package payloads, a quantitative study of edge effects must be done in order to exclude edge effects. Such an investigation was pursued and is described below.

Lateral Expansion Measurements. There are three major confounding effects in these planar foil mix investigations: dishing, payload trailing and ablator blow-by (see **Fig. 3-40**). Dishing was discussed above, but trailing and blow-by can be very serious effects if they are of sufficient density and they are not as easily identified as dishing. These can compromise the experiment and require detailed measurements of the lateral expansion of the mix package to ascertain their effects. During the past year we have performed such measurements by modifying one of the point projection spectrometers to operate when rotated by 90° while using a normally oriented spectrometer in the other data channel. By this method we could analyze the lateral expansion and measure the mix in the same shot.

For the normal and rotated data channels the two-dimensional image is essentially the same, but the spectral absorption features run

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Fig. 3-39. Typical high mix point projection spectrometry image data at the lower drive for backlighter delay times of (a) 6.0 and (b) 8.5 ns. (a') and (b') show the density distribution and material interpenetration derived from a and b respectively. The measured mix widths are indicated for both shots.

Fig. 3-40. Package effects which could mimic mix.

parallel to the interface in the rotated channel instead of perpendicular to it. **Figure 3-41** in comparison to **Fig. 3-38** schematically illustrates this. As the strength of the absorption feature depends on the density (and opacity, in addition to the chord length through the mix package) of the material, ablator material blowing by the payload would show up as a strong line extending well beyond the limits of the package as would payload material trailing the package.

Figure 3-42a shows the result of a typical high mix shot measuring the lateral expansion (the other channel measured the image of **Fig. 3-39a**). In this shot the S absorption edge is set to be approximately in the center of the mix region and the lateral expansion of the S shows no gross blow-by of ablator material. **Figure 3-42b** shows the PPS areal density distribution in addition to a transmission lineout on the S absorption line. The symmetry of both the transmission and density distribution affirm the cylindrical symmetry of the driven package in addition to the lack of any significant blow-by. Other shots in this series analyzed the PPS and Mo lateral expansion at different positions. In all cases there is no indication of significant PPS blowing by the Mo payload or Mo trailing about the PPS.

Several shots were performed taking data on the lateral expansion of the PPS ablator and/or Mo payload at different places in the package, but none of these data show any significant blow-by or trailing to confound the measurement and all of the lateral expansion data is consistent. In fact, the only evidence of any blow-by at all is the low density tail of the PPS distribution into the Mo shown in **Fig. 3-39**. These tails show up on some of the high mix data and one explanation for them is a low density sheath of PPS blowing by the Mo payload. Two-dimensional modeling¹⁰⁷ of the mix experiment shows such a sheath of PPS that causes a similar low density tail in the resultant distribution, but the density is low enough that it has negligible influence on the mix measurement.

Discussion of the Data. As there is some shot to shot changes in the total laser energy, the data should either be normalized to the mean energy or separate calculations should be done for each shot. To constrain the calculation it is better to normalize the data as then it effectively represents the same drive.

Figure 3-43 is a very preliminary version of the data which attempts to map out the penetration

and hydrodynamics. Only a subset of the data is plotted as not all of it has been analyzed at this time. Furthermore, these data have not been normalized to the mean laser drive energies. However, the figure does show that the measured interpenetration follows a pattern similar to that calculated with some exceptions. A more quantitative statement will have to wait for the final data set where the appropriate corrections and normalizations have been applied.

The solid circles and diamonds are the lower laser energy data while the open circles and diamonds are the higher drive data. Here the errors in time are taken to be the width of the backlighter pulse and the position errors due to system resolution. The higher laser energy data includes the addition of a 3.3 μm Paralyene-N (PyN) tamper over the Mo foil. By comparing this final data set to the untamped low drive data we hope to be able to make a statement on the effect of a tamper and changes in drive temperature on mix. At present it is not clear whether the higher drive data have serious recompression problems as is indicated by some of the modeling.¹⁰⁴

X-ray Preheat. In laser driven plasma/hydrodynamic experiments such as these, one can vary the degree of preheat suffered by the interface to ascertain how it affects the instability. The present experiment has a reasonable degree of preheat which causes expansion at the foil interface prior to the passage of the shock and results in a smoother density gradient from one material to the next as shown in **Fig. 3-44a**.

Figures 3-44a and b are calculations¹⁰⁴ of the hydrodynamics out to 1.0 ns and 2.0 ns respectively of the high drive data with the inclusion of a tamper. According to this calculation, at about 1.2 ns the shock breaks out and the Mo foil has expanded by over a factor of two. By decreasing the density of the Mo and changing the density gradient from an abrupt step to a smoother transition, the mix caused by the shock passage is greatly reduced.¹⁰⁸ The amount of preheat also changes the hydrodynamics and causes the mix package to decompress more with increasing preheat. However, the effect of preheat on the experiment is a rather open question as other calculations¹⁰⁵ do not show such a drastic change in the measurable mix width as a function of preheat. In any event, in addition to possibly

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Fig. 3-41. Schematic of the point projection spectroscopy image for the modified rotated spectrometer. While the normal spectrometer uses the spectral lines to quantify the mix region (as in Fig. 3-38), the rotated spectrometer uses these lines to quantify the lateral expansion.

Fig. 3-42. (a) Typical point projection spectrometry image for a rotated spectrometer data channel. Here the S absorption line measures PPS lateral expansion in the center of the mix region. As the Mo lines are well into the PPS ablator region, they are not seen. (b) Transmission on the S line and derived PPS density distribution.

MM-025-U-9345-3

Fig. 3-43. Preliminary penetration/hydrodynamic map for all presently analyzed data. Here penetration limits are plotted for comparison to calculated values. Data have not been normalized to a particular drive. Indicated errors are due to backlighter duration (200 ps) and system resolution (40 μm). The dashed line is a least squares fit to the Mo data while the solid line is a fit to the PPS data.

MB-25-U-9345-06

Fig. 3-44. Hydrodynamic calculation¹⁰⁸ of the radiatively driven mix package. The calculation to 1.0 ns (a) show the effect of the preheat on the Mo and PPS/Mo interface prior to the shock break-out. At shock break-out (b) the RM instabilities induce interpenetration and a mix region develops. The mix region and package then decompress. The calculation does not account for any mixing of the PyN tamper with the rest of the package.

affecting the amount of mix, the preheating in this experiment clearly complicates any study involving mix as a function of interface surface texture. A lower preheat variation on the present experiment is planned which should settle the modeling discrepancy mentioned above in addition to allowing study of the effect of surface finish, seeded perturbations, and modulated perturbations as these would not be washed out.

Conclusions. The present series of planar foil mix experiments has measured about 100 μm of material interpenetration in high Atwood number targets. As 30–40 μm of this can be attributed to system resolution, there appears to be ≥ 60 μm of real mix resulting from RM instabilities. Analysis is still in the preliminary stages, but the high mix data seem consistent and a map of the hydrodynamics and interpenetration from 5–9 ns after the drive is initiated can be determined from the data. Low Atwood number control experiments have not yet yielded conclusive results and are being pursued. Calculations by LLNL and Atomic Weapons Establishment (AWE) have shown agreement with the high mix experimental results and are in progress. We have investigated a host of effects which could mimic mix in the experiment and have determined them to be negligible. In particular, the lateral expansion of the mix package was investigated at great length and no significant blow-by of material was measured. Based on these studies, the measured mix is due to accelerated material intermixing and not edge effects. These experiments have been performed with a significant amount of x-ray preheating. In order to assess the effect of preheat the next step is to perform a series of experiments where it is lowered so that the interface is colder and more representative of a true density jump.

Beyond this we are starting to look at the design of a next generation of plasma/hydrodynamic instability experiments which would radically change the experiment. These experiments would study the mix in compressed matter. The crucial element is of course the target and mix package, but a higher energy backlighter of sufficient intensity would improve the experiment and help in the target design. Such a backlighter and perhaps different diagnostics would have to be developed. LLNL (A-Division) and AWE are already starting to look into new target/mix package designs. LLNL (L-Division) is also working on developing a

backlighter continuous over a 1 keV range in the x-ray energy region from 8.5–10 keV. If this is successful it could have a major impact on the next generation experiment.

Author: John D. Molitoris

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3.5. Laboratory X-ray Lasers

3.5.1. Overview of X-ray Laser Research

FY89 was a banner year for x-ray laser research at LLNL. The highlight of our research has been the development of x-ray lasers which operate at the optimum wavelength (44.8 Å) for x-ray holographic imaging of living cellular materials. Progress did not stop there, however, since we also demonstrated amplification at even shorter wavelengths (43.2 and 35.6 Å), well inside the so-called "water window" often emphasized in x-ray microscopy. A detailed review of this progress is given in the following section on Ni-like lasers.

We also have spent considerable effort on better understanding the kinetics of x-ray laser formation. In section 3.5.3 we discuss our recent progress, particularly on Ne-like schemes, in understanding the population mechanisms for Ne-like Ge amplifiers. We also were able to make silver—the highest Z, Ne-like amplifier to date.

New schemes for producing smaller, more efficient x-ray lasers are also becoming an important focus of our research. In section 3.5.3 we discuss our plans for studying multiphoton-ionized H-like recombination lasers. In addition to identifying potential Balmer alpha laser schemes that will lase at modest wavelengths, we are also studying the necessary conditions to produce Lyman alpha lasers that can achieve very short wavelengths.

Efforts to develop applications of the x-ray laser continue to make considerable progress. In Sec. 3.5.4 we will discuss x-ray holography. This year we have identified optimum samples to study and we have calculated the x-ray laser coherent energy necessary to make a usable

hologram. In Sec. 3.5.6 we discuss another novel application of the x-ray laser, that being to focus it to produce plasmas, probe plasmas or to harmonically generate even shorter or perhaps tunable x-ray laser wavelengths.

Author: D. L. Matthews

3.5.2. Nickel-Like X-ray Laser Source Development

Since 1984 when XUV amplification was demonstrated near 200 Å,^{109,110} there has been much effort to develop an amplifier to produce significant power at wavelengths below the carbon K absorption edge at 43.76 Å.¹¹¹ One motivation was to produce a coherent, high brightness source suitable for holography of biological specimens within the "water window" between the K edges of carbon and oxygen.¹¹² More recently, a detailed study of the power requirements specific to x-ray holographic imaging by London et al.,¹¹³ has shown that while for x-ray microscopy it can be advantageous to be inside the water window, for x-ray holography¹¹²⁻¹¹⁵ the dosage received by a cell is minimized for a given image resolution if the illumination wavelength is slightly to the long wavelength side of the carbon K edge.

Nickel-like x-ray lasers¹¹⁶⁻¹²⁰ were first demonstrated in a laser produced plasma of Eu in 1987.¹¹⁷ Four gain lengths of amplification were observed at 71.0 Å in Ni-like Eu. Subsequently the scheme was isoelectronically extrapolated to 50.26 Å in Ni-like Yb.¹¹⁹ The result reported here extrapolates the system further to the edge of the water window, at 44.83 Å, and provides the basis for the design of an x-ray laser source that should fulfill some of the requirements for holography of living cells.^{114, 115} Also reported here is the further extrapolation of the scheme to 43.18 Å, within the water window, a wavelength more suited perhaps, to short pulse x-ray microscopy.

Collisionally pumped Ni-like x-ray lasers are 4d-4p transitions in high-Z Ni-like ions. Electrons are elevated to the 4d levels through a combination of direct collisional excitation from the ground state and cascading from upper levels. The population inversion is maintained between the 4d and 4p levels due to fast radiative decay from the 4p levels to the 3d¹¹⁹ Ni-like ground state while the 4d levels are metastable to radiative decay to the ground state. The largest gain is expected on the J = 0-1 transition

$3d3/24d3/20-3d3/24p1/21$, at 44.83 Å in Ta⁴⁵⁺. The upper state of this transition is populated mainly by collisional excitation from the 3d¹¹⁹ ground state. The same upper state has another J = 0-1 lasing transition, to the $3d5/24p3/21$ level, (at 50.97 Å in Ta⁴⁵⁺). This second, longer wavelength transition has more gain in lower-Z ions¹¹⁷ but in higher-Z ions its stimulated emission cross section is smaller than that of its shorter wavelength partner.¹²¹ For Ta the stimulated emission cross sections of the two transitions are in the ratio 1.3:1.

The creation of the conditions for gain on these transitions requires irradiating a thin foil of Ta with a high intensity optical laser, thus causing the heated foil to become a plasma and explode.¹¹⁰ The plasma is heated by laser absorption, to an electron temperature of about 1 keV, such that the Ta is ionized to Ni-like as the foil expands to a density of order 10²¹ electrons per cm³. A complication in the design of the plasma amplifier is that transverse electron density gradients will cause the x-rays to refract out of the plasma before traveling its length, unless transverse scale lengths are sufficiently long.¹¹⁰ At a density of 10²¹, transverse scale lengths larger than 50 μm are needed for a 44.83 Å x-ray to travel the length of a 3-cm amplifier. A further constraint is set by trapping, i.e., self-absorption of the 4p-3d dump lines. The typical line center optical depth of these lines is ~3-6 with the $3d5/24p3/21-3d^{118}$ having almost twice the opacity of the other dump line and hence being more sensitive to trapping. Bu¹k Doppler shifts due to velocity gradients in the explosion of the foil lead to a reduction of the opacity of both lines in the direction perpendicular to the original foil surface, hence radiation should escape in this direction. Calculations utilizing the LASNEX, XRASER and SPECTRE, hydrodynamics, time dependent kinetics and line transfer models¹²⁰ predict that the two J = 0-1 lines should have similar gains. (See Table 3-3.)

The target design used in the Ta experiment was a foil of 127-μg cm⁻² Ta on a 24-μg cm⁻² Lexan substrate irradiated simultaneously with two of Nova's beams, superposed with a total irradiance of 4.8×10^{14} W cm⁻² in a 120-μm wide line focus up to 2 cm in length. The 2ω pulse was a 500 ps (full width at half power) Gaussian.

The foil was set up in a target chamber vacuum vessel into which two of Nova's ten beams were directed. Each beam was focused

Table 3-3. Wavelengths (in angstroms) and gains of the Ni-like Ta and W lines. The estimated uncertainties are given in parentheses. Wavelengths are calibrated against known wavelengths for the Cu-like Ta (Ref. 14) and W (Ref. 20) resonance lines.

	Transition	Wavelength (Å)		Gain (cm ⁻¹)			
		Theory ¹⁵	Observed	SFFS (peak)	MCPIGS	SFFS (Int.)	XRASER
Ta ⁴⁵⁺	$3d_{3/2}4d_{3/2}0-3d_{3/2}4p_{1/2}1$	44.76	44.83(0.02)	2.3(0.2)	1.4(0.3)	1.6(0.2)	2.6
	$3d_{3/2}4d_{3/2}0-3d_{5/2}4p_{3/2}1$	50.84	50.97(0.02)	0.5(0.4)	-1(1)	-0.9(1)	2.0
	$3d_{5/2}4d_{5/2}2-3d_{5/2}4p_{3/2}1$	74.31	74.42(0.02)		-1(1)		3.9
	$3d_{5/2}4d_{5/2}1-3d_{5/2}4p_{3/2}1$	77.40	77.47(0.02)		1.3 (0.4)		1.8
W ⁴⁶⁺	$3d_{3/2}4d_{3/2}0-3d_{3/2}4p_{1/2}1$	43.08	43.185(0.01)		2.6(0.2)		5.5
	$3d_{3/2}4d_{3/2}0-3d_{5/2}4p_{3/2}1$	49.24	not observed				4.2
	$3d_{5/2}4d_{5/2}2-3d_{5/2}4p_{3/2}1$	72.25	72.40(0.015)		-0.6(1)		4.9
	$3d_{5/2}4d_{5/2}1-3d_{5/2}4p_{3/2}1$	75.23	75.35(0.015)		0.8(0.3)		2.2

with a combination of two cylindrical negative lenses and a spherical focusing lens to produce a line focus of variable length. The target chamber has a large number of XUV and x-ray spectrometers and imaging diagnostics.¹²² In particular the foil is viewed from one end of its axis (the preferred axis for stimulated emission) by a gated grazing incidence XUV spectrometer with a micro-channel plate detector (the MCPIGS) and from the other end by an XUV spectrometer coupled to an x-ray streak camera (SFFS-Streaked Flat Field Spectrometer). Another MCPIGS spectrometer viewed the plasma from an off-axis direction while various time resolved and time integrated x-ray spectrometers recorded the x-ray transitions near 2 keV as a signature of the charge state of the plasma.

Time integrated spectra from the off-axis MCPIGS spectrometer showed optically thick Cu-like 4 – 4 resonance lines that have been previously observed and their wavelengths well documented.^{123,124} The on-axis spectra were dominated by strong line emission at 44.83 Å; weaker Ni-like emission was visible at 50.97, 74.42, and 77.47 Å. These last two lines are J = 2–1 and J = 1–1 transitions which are expected to have gain and share the same lower state as the 50.97-Å J = 0–1 line. The complete list of wavelengths both calculated and measured (using the Cu-like wavelengths as fiducials from Ref. 128), is given in Table 3-3. The calculations are Grant's Multi-Configurational-Dirac-Fock code with an optimized level scheme.¹²⁵

Figure 3-45 shows spectra from the other on-axis diagnostic, the streaked flat field spectrometer which is a Harada grating¹²⁶ spectrometer coupled to an x-ray streak camera with 70 ps continuous time resolution. Spectra at one instant in time are shown from two different length foils. Figure 3-45 illustrates the non-linear growth of the line intensity with foil length. The spectrometer viewed 2nd, 3rd and 4th order to obtain good dispersion from the 1200-mm⁻¹ grating. The aluminized-Mylar filter bandpass restricted its spectral range to the region between the carbon K edge and slightly to the long wavelength side of the 50.97-Å J = 0–1 line.

The line intensities from the SFFS are plotted in Fig. 3-46(a) as a function of foil length. They are fit to the formula for the axial emission from a distributed source of amplified spontaneous emission from Linford et al.¹²⁷ The

data shown is for the time of peak gain. The time resolved values of the gains on the two Ni-like J = 0–1 lines are shown in Fig. 3-46(b) illustrating the 250 ps duration of the 44.83-Å gain. The time integrated MCPIGS line intensities were fit in the same way; the results are summarized in Table 3-3. The SFFS data for the 44.83-Å line, peaks at a gain of $2.3 \pm 0.2 \text{ cm}^{-1}$ while the MCPIGS gain is $1.4 \pm 0.3 \text{ cm}^{-1}$. The value of the gain obtained by integrating the SFFS data in time is $1.6 \pm 0.2 \text{ cm}^{-1}$ which is consistent with the MCPIGS value and, as expected, shows evidence that our time integrated gain measurements may underestimate the peak gain.

As an attempt to extrapolate this scheme to shorter wavelengths and try and demonstrate a higher gain-length product, tungsten foils of thickness 89 $\mu\text{g cm}^{-2}$ on 20 $\mu\text{g cm}^{-2}$ Lexan, up to 3 cm in length were irradiated with a total of $3.2 \times 10^{14} \text{ W cm}^{-2}$ of 2 ω light. The Ta experiment had shown significant amounts of (Co-like) Ta⁴⁶⁺ 4 – 3 and 5 – 3 x-ray emission indicating that we could reduce the irradiance and still ionize to Ni-like. The foils were thinner as our experience with optimizing the gain in Yb and earlier work with Eu⁸ showed that thinner foils produced higher gain provided that they could be made hot enough. This may be due to less trapping or reduced refraction in thinner foils that burn through faster. Figure 3-47(a) shows on-axis MCPIGS spectra from 2.5 and 1.7 cm long W foils with the J = 0–1 laser line at 43.185 Å in 1st and 2nd order and weaker J = 2–1 and J = 1–1 Ni-like lines at 72.40 and 75.35 Å respectively. In Fig. 3-47(b) the intensity of the 43.185-Å line is plotted as a function of foil length from the MCPIGS spectrograph. The time integrated gain from the fit to the same formula as before is 2.6 cm^{-1} with a maximum gain length of 7 recorded. Time resolved XUV spectra were not obtained from the W shots but the peak gain may be higher than the time averaged MCPIGS value. The other J = 0–1 line (predicted to be at 49.24 Å) was not observed. The gains on the J = 1–1 and 2–1 lines were qualitatively the same as for Ta and are summarized in Table 3-3.

Table 3-3 also shows calculated (time integrated) gains from simulations similar to those of Ref. 120. Note that there is qualitative agreement with the theoretical gain on the 44.83 and 43.185-Å lines and also their respective J = 1–1 lines (at 77.47 and 75.35 Å). The calculated gains on the other lines terminating on the

Fig. 3-45. On axis spectra, from the streaked Harada grating spectrometer, (SFFS) from (a) 1.7 and (b) 0.6 cm long Ta foils. Data obtained at the peak of the 44.83-Å gain.

Fig. 3-46. (a) Line intensity as a function of foil length from SFFS at the peak of the gain. The fit is to the equation $I = \epsilon/\alpha (\exp^{\alpha L} - 1)^{1.5} / (\alpha L \exp^{\alpha L})^{1/2}$ for a distributed source with emissivity ϵ , small signal gain α and length L ¹⁸. (b) The time history of the gain deduced from the SFFS data.

Fig. 3-47. (a) On axis spectra (MCPIGS) from 2.5 and 1.7 cm long W foils. The 2.5 cm foil had a gain-length product of 7 on the 43.18-Å $J = 0-1$ laser line. (b) Intensity of the 43.18-Å $J = 0-1$ transition of Ni-like W as a function of foil length, from the MCPIGS spectrometer.

$3d5/24p3/2)_1$ level are inconsistent. In particular, the model predicts more gain on the $J = 2-1$ line at 74.42 Å (72.40) than the $J = 1-1$ line at 77.47 Å (75.35) while the experiment shows no gain on the $J = 2-1$. (In Ref.117 a low gain was seen on the 2-1 and 1-1 lines of Eu^{35+} , in disagreement with calculation.) The ratio of the gains on the two $J = 0-1$ lines is predicted to be the same as the ratio of their stimulated emission cross sections while the Ta $J = 0-1$ at 50.97 Å has little or no gain and the analogous line in W was not even observed. These observations may be explained by errors in the population mechanisms for the $3d5/24p3/2)_1$ level, either through an underestimate of the significance of trapping on the 4p-3d dump lines, or through an error in the coupling of the 4p levels to other levels such as the 4s states. Postulating a reduced electron temperature in the plasma (which might explain the lower than expected short wavelength $J = 0-1$ gains) does not account for the ratio of the $J = 0-1$ gains. This inconsistency of the gains of two $J = 0-1$ lines that share the same upper state, points to a shortcoming in our modeling of Ni-like lasers.

We estimate an energy output from the 44.83-Å Ta amplifier of 1 μJ for 4 gain lengths of amplification. The 7 gain lengths of the brightest W laser gave an estimated energy output of order 20 μJ. The gain duration of order 250 ps is long enough to allow the use of normal incidence multi-layer x-ray mirrors to double pass the gain medium.¹²⁸ Using a 5% reflectivity mirror we should be able to produce a net amplification equivalent to 11 gain lengths by double passing a 3-cm foil. This amplifier would have an energy output of order 1 mJ. Optimization of the gain in the target through varying its thickness could increase this output further. The saturation intensity of the XRL is estimated to be of order $10^{11} \text{ W cm}^{-2}$ which would occur at 12-6 gain lengths and an energy of 4 mJ depending on the laser divergence. If the divergence is restricted due to the geometry of the target or through use of apertures in combination with the mirror, the amplifier will saturate at a higher gain-length product. Estimates of the x-ray laser intensity needed for holography vary. For single frame holography of a cell with resolution of 300 Å an energy of 300 μJ has been estimated in Ref. 114, while more sophisticated experiments have been proposed using as little as 10 μJ.¹²⁸ This energy should be in one transverse mode of the laser. The number

of modes emitted by the 44.83-Å x-ray laser will be studied experimentally now that we have passed the initial hurdle of generating gain and significant energy. A worst case estimate for a 3-cm Ta amplifier is of order 10^4 modes, but this may be overly pessimistic as it ignores any coherence in the beam beyond that expected from its geometric divergence and source size. Now that we have produced gains of order 2.5 cm^{-1} at 43.18 and 44.83-Å, the optimization of that gain may allow the saturation of a Ta laser (with $\alpha_1 > 12$) using two beams of Nova, and provide enough energy per transverse mode to perform holography experiments.

In more recent experiments up to 8 gain lengths of amplification have been demonstrated with Ta targets shot at a similar irradiance to that used with the W foils. The gain coefficients were up to 3 cm^{-1} at 44.83-Å with a gain duration of order 250 ps. In order to develop a holography x-ray source, our program will now concentrate on fine tuning the gain optimization through varying the target design and also performing experiments with multi-layer mirrors. Experiments with 1ω light and also solid targets will also be attempted in order to try to enhance the efficiency of the Ni-like Ta laser to the point where a smaller facility than that of Nova could be used as a pump source.

Experiments with gold foil x-ray lasers have succeeded in showing amplification at 35.60-Å in Ni-like Au. The foils, $140 \mu\text{g cm}^{-2}$ thick and up to 1.5 cm in length were irradiated with a total of $6 \times 10^{14} \text{ W cm}^{-2}$. A gain coefficient of order 2 cm^{-1} was estimated based on comparing the output of 1 and 1.5-cm-long foils. This gain is consistent with Yb, Ta, and W results and is an illustration of how the Ni-like system can be extrapolated to even shorter wavelengths. Future experiments are planned with Bi foils that should extend the Ni-like laser line spectroscopy to 30.6 Å.

Another interesting result from the source development experiments is the observation of cobalt-like analogs to the Ni-like lasers. These analogs have been searched for in Ne-like x-ray laser schemes. In 1984, XUV amplification of $2p53p-2p^53s$ transitions was demonstrated in a laser produced plasma of neon-like selenium.^{109,110} Gains of order 4 cm^{-1} were observed on several Ne-like Se lines near 200-Å. At the conclusion of these experiments two major inconsistencies with theory remained. One was the low gain on the $J = 0-1$ collisionally pumped 182-Å line. The other anomaly was the non-

appearance of fluorine-like analogs to the Ne-like lasers.¹¹⁰

Ne-like ions have 10 bound electrons, with a $2p^6$ ground state. Electrons in the $n = 3$ shell will usually decay rapidly back to the ground state by one of the fast $3s-2p$ or $3d-2p$ dipole transitions. However $3p$ electrons can be metastable to radiative decay and hence $3p-3s$ population inversions can occur. Current understanding of Ne-like x-ray lasers is that population of the $3p$ levels is occurring through a combination of collisional excitation from the ground state and cascading from higher n states and dielectronic recombination from F-like. The line fed most efficiently by collisional excitation from the ground state is the $2p_{1/2}3p_{1/2}0 - 2p_{1/2}3s_1$ $J = 0-1$ (at 182.4-Å in Se) although other lines are also predicted to have gain. Simulations of the Ne-like Se laser agree quite well with the measured gains on the brighter lines although the $J = 0-1$ gain is a factor of two less than calculated. These calculations generally predict an F-like population that is ~40% at the time of peak Ne-like gain.¹¹⁰ Fluorine-like ions have a ground state with either a $2p_{3/2}$ or a $2p_{1/2}$ vacancy. Most XUV line emission from F-like ions is $2p-2s$ resonance emission which has been widely observed.^{120,129} However transitions of the form $2p^43p - 2p^43s$ exist that, like their Ne-like counterparts, are not expected to be seen unless amplified. These lines are predicted to have gain by the same models that predict Ne-like gains.^{119,130} Ne-like laser schemes have been extrapolated to other elements over a wide range of irradiance and target conditions.¹¹¹ Gain-length products of up to 16 have been produced on the brightest Ne-like lines however to date no observation has been made of F-like $2p^43p - 2p^43s$ emission.

The pumping scheme for a Ni-like laser is illustrated for Ni-like Yb in Fig. 48(a). The $4d$ levels are populated through a combination of direct collisional excitation from the ground state and cascading from upper levels. The population inversion is maintained between the $4d$ and $4p$ levels by fast radiative decay from the $4p$ levels to the ground state. Although several $4d-4p$ transitions are capable of being inverted, the largest gain is observed on the 50.23-Å $J = 0-1$ transition $3d_{3/2}4d_{3/2}0 - 3d_{3/2}4p_{1/2}1$. The upper state is pumped predominantly by collisional excitation from the $3d^{119}$ ground state. In Fig. 3-48(a) the collisional excitation rates¹³¹ are shown for an electron

temperature of 0.8 keV and an electron density of 10^{21} cm^{-3} , representative of the characteristics of Yb x-ray laser plasmas. Other transitions (not shown in the figure), predicted to have large gain^{117,118} include a second $J = 0-1$ transition (at 56.11-Å) and $J = 2-1$ and $J = 1-1$ $4d-4p$ transitions. In experiments, as discussed above, Ni-like lasers show significant gain only on the $J = 0-1$ transitions while the $J = 2-1$ and $1-1$ lines have little or no gain.

The cobalt-like ground state mostly consists of $3d_{5/2}$ and $3d_{3/2}$ holes in an otherwise closed $n = 3$ shell. The direct analog to the Ni-like $J = 0-1$ laser is illustrated in Fig. 3-48(b). The transition is $3d_{5/2}(3d_{3/2}4d_{3/2}0) (J = 5/2) - 3d_{5/2}(3d_{3/2}4p_{1/2}1) (J = 7/2)$. It is the same configuration as the brightest Ni-like line but with an extra $3d_{5/2}$ vacancy.

Figure 3-49 shows time resolved spectra taken with the SFFS streaked spectrometer looking along the axis of Yb and Ta x-ray laser foils. The upper Yb spectrum shown in Fig. 3-49(a) is from a 2.5 cm long Yb foil of thickness $75 \mu\text{g cm}^{-2}$ irradiated with $3.5 \times 10^{14} \text{ W cm}^{-2}$ of 2ω light in a 500 ps pulse. The main spectral feature is the Ni-like x-ray laser line at 50.23-Å; however a long wavelength "satellite" is also visible at 51.76-Å. The lower spectrum in Fig. 3-49(a) is from a 0.9-cm-long, $100\text{-}\mu\text{g cm}^{-2}$ Yb foil irradiated at $9.4 \times 10^{14} \text{ W cm}^{-2}$. The higher irradiance data shows an increase in emission from the long wavelength satellite. The lower spectrum in Fig. 3-49(b) is from a 1.7-cm-long, $127\text{-}\mu\text{g cm}^{-2}$ -thick Ta foil irradiated at $4.8 \times 10^{14} \text{ W cm}^{-2}$. The Ni-like x-ray laser line is evident at 44.83-Å, together with a similar long wavelength satellite at 46.07-Å. The upper spectrum is from a 2.5-cm-long, $95\text{-}\mu\text{g cm}^{-2}$ -Ta foil irradiated at a lower irradiance of $2.4 \times 10^{14} \text{ W cm}^{-2}$, although the Ni-like $J = 0-1$ is brighter the long wavelength line has disappeared. The satellite lines were also detected on the other on-axis spectrometer, the MCPIGS. The wavelengths quoted were obtained by calibrating the MCPIGS spectrometer with the known wavelengths of Cu-like $4-4$ resonance lines.^{123,124} None of the Ni-like or satellite lines were visible in the off-axis MCPIGS spectra.

The identification of the long wavelength satellite was made after reviewing RAC¹³² calculations of Ni-like $4-4$ lines of Yb and Ta and eliminating all reasonable possibilities for radiatively allowed transitions. The closest Ni-

Fig. 3-48. Simplified energy level diagram for: a) nickel-like ytterbium and b) the analogous levels in cobalt-like ytterbium showing the main levels communicating with the brighter 4d–4p transition in both ions. Collisional excitation rates are shown for an electron temperature of 0.8 keV and density of 10^{21} cm^{-3} . Spontaneous emission rates are shown in parentheses.

Fig. 3-49. Time resolved on-axis spectra from a) ytterbium foils; upper spectrum is from a 2.5-cm-long, 75- μgcm^{-2} -Yb foil irradiated at $3.5 \times 10^{14} \text{ Wcm}^{-2}$ lower spectrum is from a 0.9-cm-long, 100- μgcm^{-2} -Yb foil irradiated at $9.4 \times 10^{14} \text{ Wcm}^{-2}$; and, b) tantalum foils; upper spectrum is from a 2.5-cm-long, 95- μgcm^{-2} -Ta foil irradiated at $2.4 \times 10^{14} \text{ Wcm}^{-2}$ lower spectrum is from a 1.7-cm-long, 127- μgcm^{-2} -Ta foil irradiated at $4.8 \times 10^{14} \text{ Wcm}^{-2}$, showing the Co-like line on the long wavelength side of the brighter Ni-like $J = 0-1$ laser line. The Co-like line is stronger in the higher irradiance data.

like line calculated was in fact the $J = 0-1$ already identified. RAC calculations for Co-like Yb and Ta were compared with the measured line positions and the positions of the $3d5/2(3d3/24d3/2)0$ ($J = 5/2$) – $3d5/2(3d3/24p1/2)1$ ($J = 7/2$) analog to the Ni-like $J = 0-1$ was found to agree to within 5 eV. The RAC calculations tend to overestimate the energy of the $J = 0-1$ transitions by about 5 eV as they use an average level treatment of configuration interaction contributions to the energy of the upper $J = 0$ state. More detailed calculations using Grant's multiconfigurational Hartree Fock code¹²⁵ run with an optimized level scheme gave better agreement with both the Ni-like and Co-like observed wavelengths. The complete list of wavelengths, both calculated and measured, is given in Table 3-4. Also quoted in Table 3-4 are the disagreements (ΔE) between the calculated and measured transition energies. It can be seen that ΔE is similar for the Ni-like and Co-like lines, which is to be expected as they are affected by configuration interaction to a similar degree. It is unlikely that the lines are Cu-like or Zn-like as the spectra of 4–4 resonance lines is well known.^{123, 124} We have also seen that increasing the temperature of the plasma (through higher irradiance) increases the satellite emission [(Fig. 3-49(a)). Reducing the irradiance causes the line to disappear as shown in Fig. 3-49(b). This behavior is inconsistent with the line coming from a less ionized charge state such as Cu- or Zn-like.

The time history of the Co-like satellites are similar to that of the Ni-like $J = 0-1$. The emission time of both is of order 300 ps while background continuum and optically thick Cu-like lines (when seen) last for up to a nanosecond. The Co-like line emission reduces slightly earlier in time than the Ni-like emission in many shots, consistent with the temperature of the plasma decreasing after the absorption of laser light peaks.

Further evidence, for the identification of the satellites, is the observation of gain on these lines. The level diagrams shown in Fig. 3-49(a) and (b) both show transitions that can be easily inverted. A steady state calculation for the gain on the Ni-like line using a detailed atomic data file with rates from Ref. 131 predicts a gain of 11 cm^{-1} at 50.23 Å assuming a 25% Ni-like concentration. This calculation assumes a Doppler broadened line at an ion temperature of

400 eV to give a stimulated emission cross section of $8.5 \times 10^{-16} \text{ cm}^2$. (More complete calculations using the XRASER time dependent kinetics code coupled with the LASNEX hydrodynamics code show gains between 5 and 10 for the expected conditions within the Yb and Ta x-ray lasers.^{119, 120} In experiments gains of order 2.5 cm^{-1} have been observed on Yb and Ta $J = 0-1$ lines.¹³³

Applying a similar steady state treatment to the Co-like system in Fig. 3-45(b) (using a detailed atomic data file with 781 levels) we see that with a stimulated emission cross section of $3.9 \times 10^{-16} \text{ cm}^2$ and 25% Co-like ion concentration, we would expect a gain of 5.6 cm^{-1} for the plasma conditions in Fig 3-45. Hence for equal quantities of Ni- and Co-like ions we would expect gains in the ratio 2:1. If there is significant buildup of the lower laser level population due to self absorption (trapping) of the 4p–3d dump line, the ratio of the gains would be more complex. The Co-like dump line absorption cross section is smaller by a factor of 3 than that of the Ni-like dump line, hence the optical depth of the line should be smaller and the Co-like line should not be affected by trapping while the Ni-like line may be affected.

The Co- and Ni-like line intensities from several plasma lengths were fit to the scaling for the axial emission from a distributed source of amplified spontaneous emission given by Linford et al.¹²⁷ The $3.5 \times 10^{14} \text{ W cm}^{-2}$, 75- $\mu\text{g cm}^{-2}$ -thick Yb data illustrated by the upper curve in Fig. 3-49(a) showed a peak gain of $2.2 \pm 0.2 \text{ cm}^{-1}$ on the Ni-like $J = 0-1$, while the Co-like line had a gain of $0.7 \pm 0.3 \text{ cm}^{-1}$. The gains are approximately in the ratio 3:1, which may indicate that there is almost as many Co-like as Ni-like ions in the plasma. A similar analysis for the 127- $\mu\text{g cm}^{-2}$ -thick Ta data gave $2.3 \pm 0.2 \text{ cm}^{-1}$ as the peak gain at 44.83 Å and $2.1 \pm 0.4 \text{ cm}^{-1}$ as the gain in the Co-like Ta line which may indicate a large Co- to Ni-like ion ratio.

The observation of gain on the Co-like lines adds weight to the identification. Measurements of the gain and emissivity (the two parameters that are fit) may be a powerful tool to indicate the ratio of Ni-like to Co-like ions. The measured Co-like gain of 2.1 cm^{-1} in the Ta targets is a strong indication that the plasma may be overionized and that we could drop the irradiance and not necessarily see a reduction in Ni-like gain. This is an important point because in earlier experiments we had been concerned that the plasma was under-ionized. Indeed, the

Table 3-4. Calculated and observed wavelengths for the Ni-like J = 0 - 1 and the analogous Co-like line.

	Observed (Å)	Calculated (Å)	ΔE (eV)
Yb ⁴²⁺ $3d_{3/2}4d_{3/2}0 - 3d_{3/2}4p_{1/2}1$	50.23 ± 0.02	50.16	-0.34 ± 0.1
Yb ⁴³⁺ $3d_{5/2}(3d_{3/2}4d_{3/2}0) (J=5/2) -$ $3d_{5/2}(3d_{3/2}4p_{1/2}1) (J=7/2)$	51.76 ± 0.02	51.71	-0.23 ± 0.1
Ta ⁴⁵⁺ $3d_{3/2}4d_{3/2}0 - 3d_{3/2}4p_{1/2}1$	44.83 ± 0.02	44.76	-0.43 ± 0.1
Ta ⁴⁶⁺ $3d_{5/2}(3d_{3/2}4d_{3/2}0) (J=5/2) -$ $3d_{5/2}(3d_{3/2}4p_{1/2}1) (J=7/2)$	46.07 ± 0.02	45.97	-0.59 ± 0.1

data from the upper spectrum in Fig. 3-49(b) is for a lower irradiance, thinner target which shows significant gain on the Ni-like $J = 0-1$ but no Co-like line. This latter experiment may have a more optimized ionization balance for Ni-like and at the same time allow more efficient use of the pump laser energy.

Further experiments that could be interesting are to vary the target geometry and use the Co-like line to study the effect of trapping on the Ni-like kinetics. As mentioned above, the Co-like line should be free of trapping while the influence of trapping on the Ni-like line is of concern. Hence the observation in Ta of the high gain on the Co-like line despite the low intensity relative to that of the Ni-like line may be an indication that the source function of the Ni-like line is larger than the gain would indicate. This is due to a build up of population in the lower laser level.

One question that arises is where are other Co-like lines? Some are predicted to have gain but at least a factor of two less than the observed Co-like line. No other line emission is reproducibly observed on-axis which could be Co-like.

In conclusion, we have identified 4d-4p laser lines in Co-like ions that are analogs to Ni-like x-ray lasers. These lines appear with gains in qualitative agreement with what would be expected if there was a significant Co-like population. This is the first reported observation of analogous lasers from a neighboring charge state within a collisionally pumped x-ray laser. Also we have demonstrated a working x-ray laser operating at an optimum wavelength for producing holographic images of living cellular material. In addition, we observed 7 gain lengths at 43.185 Å in Ni-like tungsten in the first demonstration of an x-ray amplifier operating within the "water window." It remains to enhance output power and characterize and improve coherence in order to have an x-ray laser source at 44.83 Å suitable for holography. Increased efficiency, perhaps to the point that pump sources smaller than Nova may be usable, also remains a goal.

Author: B. J. MacGowan

3.5.3. Ne-Like X-ray Laser Kinetics Experiments on the Nova Laser

In 1984, an experimental team produced the first unambiguous demonstration of an extreme

ultraviolet amplifier at LLNL's Novette laser.¹³⁴ A lengthy series of investigations conducted from 1985 to 1987 on the two-beam facility at the Nova laser uncovered some questions relating to our basic understanding of the mechanisms governing charge state distributions, excited state populations (which create the inversions necessary for stimulated emission), and photon trajectories.¹³⁵ Following an extensive shutdown of the two-beam facility in FY 1989, we began a detailed experimental investigation of these so-called "kinetics" issues.

Although a variety of methods have been proposed for producing XUV and soft x-ray amplifiers and a number of these are being experimentally investigated by various groups around the world,¹³⁶ our interest lies in understanding the physical mechanisms at work in an XUV amplifier setup involving neon-like ions.

Neon-Like Extreme Ultraviolet Amplifiers.

A neon-like ion is an atom with enough electrons removed so that only 10 remain bound to the nucleus (as in neutral neon, hence the terminology). Because this electronic configuration represents a closed-shell structure, the neon-like species is especially stable and is present to a significant extent in plasmas of medium atomic number ($Z \sim 25-50$) over a large range of temperatures. In addition, the atomic physics of excited neon-like ions lends itself to the buildup of population inversions. The first (i.e., lowest-lying) excited state has the electron configuration $2p^33s$, which is strongly radiatively coupled to the $2p^6$ ground state. Population in the $2p^33s$ state radiatively decays to the ground state in a characteristic time of about 10^{-13} s. The next higher-energy excited state has configuration $2p^33p$ and is forbidden by angular momentum selection rules from decaying directly by dipole radiation to the ground state. Higher-order mechanisms take place in a characteristic time of about 10^{-10} s. Hence, even a randomly distributed ensemble of excited neon-like ions very quickly develops an inversion between the 3p and 3s levels and radiates amplified spontaneous emission (ASE) on different transitions ranging in energy from about 30 to over 120 eV, depending on atomic number.

In this study, we concentrated our effort on two elements, germanium ($Z = 32$) and silver ($Z = 47$). The specific energy levels of relevance to the XUV amplification in these elements are shown in Fig. 3-50. Our past experiments with neon-like

Figure 3-50. Top: excited state energy levels in neon-like germanium and silver involved in XUV amplification. Bottom: schematic of relative line positions and axial spectrometer coverage (shaded rectangles). Note that, in silver, the weak $J = 0$ line shifts in position relative to the other lines, and the excited state splittings are increased.

selenium unambiguously identified five lines that undergo ASE in these systems;¹³⁷ the number of states (hence, the number of lines) involved arises from fine-structure splitting in both the upper and lower laser levels produced by interactions between the angular momenta of the core and the excited electron. The transitions are conveniently labeled by the total angular momentum, J , of the upper laser level as shown in Fig. 3-50. Experimentally, what has been seen in the XUV for all elements tested to date is a dominance, in both gain and absolute intensity, of the two lines whose upper state is the $3p_{3/2} J = 2$ lines (as shown in Fig. 3-50), followed by the long-wavelength $J = 2$ and the $J = 1$. The weakest is usually the $J = 0$. Table 3-5 summarizes our observed gains for thin five lines in germanium, selenium, yttrium, molybdenum, and silver.

These observations are to some extent at variance with our understanding of the populating mechanisms in these systems. Originally, it was believed that the dominant means of populating the upper laser levels was direct excitation from the ground state through electron-ion collisions. Straightforward (but highly detailed) rate equation analysis of the collisionally modeled excitation kinetics predicted the highest gain on the $J = 0$ line, where, ironically, the lowest gain has actually been observed.¹³⁸ Subsequent modeling efforts, which included the effects of dielectronic recombination, reduced the disparity somewhat, but any reasonable parametrization still puts the $J = 0$ gain at least comparable to that of the bright $J = 2$ lines.¹³⁹

Outside researchers have offered two explanations as possible solutions for the " $J = 0$ problem." The first speculates that we are fundamentally mistaken about the dominant pumping mechanism in our amplifiers, that recombination cross sections are about an order of magnitude higher than we predict, and that recombination determines the bulk of the excitation in our systems.¹⁴⁰ Because recombination naturally favors higher angular momentum states, the $J = 2$ states automatically come out with the highest gain and the $J = 0$ states come out with the lowest. The second solution suggests that the spectral linewidths of the amplified lines are narrowed below their Doppler width by an effect called "Dicke narrowing." It is suggested that this effect, along with plasma turbulence, (which gives large local differential Doppler shifts), causes the

effective macroscopic gain to be reduced because of frequency mismatch between adjacent turbulence cells.¹⁴¹ The Dicke narrowing is most pronounced for the $J = 0$ level, and, hence, its gain would be most reduced.

From 1985 to 1987, experiments were done to test some aspects of these proposed explanations. An accurate, absolutely timed streak-camera recording of the $J = 2$ angular footprint versus time demonstrated that effectively all of the amplification was occurring in the main part of the driving laser pulse.¹⁴² By the time the Nova pulse had peaked, the XUV amplification had turned off. This is not the temporal signature expected to be characteristic of a recombinantly driven amplifier. In addition, the absence of significant fluorine-like gain in the XUV indicates a problem with the recombination-predominant explanation.¹⁴³ We are presently developing the diagnostic capability to measure XUV lines with sub-Doppler resolution, which will yield line-shape information that can be used to test the Dicke narrowing hypothesis.

Some of our past observations point to the need for additional research to better understand some of the processes involved. In long selenium amplifiers, where the gain-length product exceeds 12, pronounced asymmetry is apparent in the angular footprint of the beam.¹⁴² The time-integrated XUV intensity pattern is double-lobed, with each lobe about 11 mrad wide and centered about 8 mrad above and below the plane of the target foil. There is an intensity minimum at the physical on-axis location. The targets are composed of a layer of lasant deposited on a thin plastic substrate, producing a natural target asymmetry. This asymmetry carries over to the footprint pattern in that the lobe centered above the lasant side of the target is about five times brighter than the lobe centered above the plastic side. Also, spatially resolved images of the selenium amplifier lines indicate a source region of less than $70\text{ }\mu\text{m}$,¹⁴⁴ which is quite small compared to the few hundred micrometers predicted for an expanded plasma from a hydrodynamically thin foil target.¹³⁸ These observations, coupled with the early time history indicated from the absolutely timed streaked measurements mentioned above, strongly indicate that the real neon-like XUV amplifier is predominantly a high-density, nonuniform medium, evolving in time to the low-density-gradient amplifier previously suggested.

Table 3-5. LLNL-measured gains for neon-like XUV amplifiers.^a

	Germanium ^b Z = 32	Selenium ^c Z = 34	Gain (cm ⁻¹) Yttrium ^c Z = 37	Molybdenum ^c Z = 42	Silver ^b Z = 47
J = 0	1.0 (19.64 nm)	2.4 (18.243 nm)	(15.50 nm)	0.0 (14.16 nm)	d
J = 2	2.6 (23.22 nm)	4.0 (20.638 nm)	4 (15.50 nm)	4.1 (13.10 nm)	6.6 (9.93 nm)
J = 2	2.8 (23.63 nm)	3.8 (20.978 nm)	4 (15.71 nm)	4.2 (13.27 nm)	4.3 (10.03 nm)
J = 1	2.3 (24.74 nm)	2.2 (22.028 nm)	2.9 (13.94 nm)	d	
J = 2	1.8 (28.64 nm)	3.5 (26.294 nm)		d	

^a These gains are derived from measured relative intensities by fitting to the so-called Linford function.¹⁸ Also measured by the LLNL group but at the Phebus laser at Limeil, France, were gains for the J = 2 strontium lines:¹⁹ 4.4 cm⁻¹ (16.41 nm), 4.0 cm⁻¹ (16.65 nm).

^b Measured in the 1989-90 campaign.

^c Measured in the 1985-86 campaign.

^d These lines have been observed in our on-axis XUV spectra for silver, but more data are needed to determine gains, if any.

Since the two-beam facility was shut down in 1987, several other groups from around the world have succeeded in producing neon-like XUV amplifiers. A team from NRL demonstrated amplification from neon-like copper and germanium in 1987¹⁴⁵ and from neon-like zinc, arsenic, and selenium in 1988.¹⁴⁶ Also in 1988, a British group produced and characterized a germanium amplifier.¹⁴⁷ In 1989, a team from the University of Rochester developed neon-like germanium amplifiers.¹⁴⁸ In all these cases, the targets were hydrodynamically thick with single-sided irradiation at 1.06 μm . These observations were very similar: all gains were measured higher than those that we measured, and the $J = 0$ line gain stood out as being remarkably high, comparable to the $J = 2$ lines. The angular footprints in these cases were also similar, with a single observed lobe at approximately 10 mrad above the target plane. Given that a thick target inherently develops and maintains a strong electron density gradient, refractive photon trajectories can produce such a footprint.

Ne-Like X-ray Laser Experiments. In FY 1989, we were challenged to design and conduct experiments that would critically test our modeling. With a limited shot budget for that year (we had 27 shots in 5 days from late August into September), we concentrated for the most part on two specific paths of inquiry: the demonstration of XUV amplification in neon-like silver and the somewhat extensive characterization of neon-like germanium amplifiers under various ionization conditions. Silver has by far the highest atomic number ($Z = 47$) of any target material that we use for the production of a neon-like XUV amplifier; even the demonstration of lasing is an exercise in the extrapolation of hydrodynamics and atomic physics scaling relations. Importantly, the energy-level separation for silver is significantly larger than for selenium. This results in shorter-wavelength XUV ASE radiation, but, more importantly, a number of $n = 3$ to $n = 2$ electric quadrupole transitions can be seen in the 3- to 4-keV region which originate in the upper laser levels. Because of their multipolarity, the lines are optically thin and, hence, provide a direct measure of the excited state population density.

Germanium, in contrast, has a lower atomic number ($Z = 32$) than selenium and is, therefore, slightly easier to ionize. The twofold increase in

Nova beam energies that have been achieved since we conducted our first experiments in 1986 have given us a unique opportunity to study ionization kinetics and their effects in real amplifier-geometry targets. Specifically, we can achieve maximum irradiances in line focus geometry of about 10^{15} W/cm^2 in 1 cm of target length (with an obvious inverse relationship between target length and maximum target irradiance this irradiance is the total energy inverted in two opposite beams divided by the area of a focal spot and the direction of the Nova laser); the nominal selenium irradiance is $6 \times 10^{13} \text{ W/cm}^2$. Because the average electron energy scales approximately as the electric field strength of the Nova pulse (i.e., the square root of the intensity), we can achieve about a factor of 3 variation of the collisional excitation energy in either direction from nominal, provided we efficiently couple the pump laser energy into the plasma by appropriately varying target thickness.¹⁴⁹ This allows us to substantially under- and overionize germanium amplifiers. The overionized case is of particular interest because, once the L-shell ionization energy is breached, a large variety of charge states can be produced. Some of these, such as fluorine-like ions, are predicted to show gain on certain 3p-3s transitions. Germanium is also of interest as a point of comparison with the work done elsewhere.

The practical emphasis through most of FY 1989 was to assemble and field an array of diagnostic instrumentation in the two-beam chamber in order to generate high-resolution temporal- or, in some cases, angular- or spatial-resolving spectra in the XUV and x-ray. The design of the array had to be flexible enough to accommodate the varying energy coverages required by the widely different experiments performed in the chamber. The diagnostics included two spectrographs from a microchannel-plate intensified grazing-incidence spectrometer (MCPIGS), Rowland grating XUV survey with time-gated microchannel-plate detectors (one located along the amplifier axis and the other located well off-axis); a streaked, flat-field spectrograph, which is a varied-line-space grating instrument coupled to an x-ray streak camera for detection of the XUV emission in the on-axis location opposite to the on-axis MCPIGS; two streaked, crystal x-ray spectrographs (the Keanetech and XCSS); eight time-integrated crystal x-ray spectrographs (the Henways); a

time-integrated spatially imaging crystal x-ray spectrograph (the SIS); and an x-ray pinhole camera.

Some of our first data came from the silver amplifiers. We were limited to target lengths of approximately 1 cm because of the high irradiances required to strip the 37 electrons. We observed bright, directional ASE emission on the major $J = 2$ lines at 9.93 and 10.03 nm from a 9-mm target shot at maximum irradiance. A cursory survey of the effects of varying target thickness indicated that the inversions were enhanced at greater areal densities. Using samples from the thickest targets prepared for the series ($150 \mu\text{g}/\text{cm}^2$), we initiated a brief length study to determine the gains. Axial XUV spectra from this survey are shown in Fig. 3-51. The relatively abrupt onset of the neon-like emissions with increasing length indicates high gain coefficients (Table 3-5).

The numbers should be interpreted with some caution, however, because of the poor statistics. We could not discern any asymmetry or pointing of the amplified lines, which is as expected for such a low gain-length product amplifier. The interesting electric quadrupole (E2) lines were unambiguously identified in both the streaked and time-integrated x-ray spectra. The intensities of the E2 lines were quite low and suffered fairly high backgrounds from the nearby allowed 3s–2p neon-like lines. This pointed out the need for a higher-resolution focusing-crystal spectrograph, of the so-called “von Hamos” design. We are developing such an instrument for fielding in the two-beam chamber in 1990.

Bright lasing was observed in germanium amplifiers irradiated at $9 \times 10^{13} \text{ W}/\text{cm}^2$ (1.5 times “nominal” irradiance) on all five neon-like lines (Fig. 3-52). As recorded by the SFFS, all amplified lines displayed emission characteristics seen before in selenium: very bright, short temporal duration (~ 250 ps FWHM), and all positioned together very early in time (before noticeable spontaneous line or continuum emission). The axial MCP-IGS spectra gave no indication of any beam pointing or asymmetry, even for the longest-length targets (which is, to some extent, in contrast to our experience in selenium). Gains for the five lines were fairly easily measured (although the bright $J = 2$ lines required heavy filtering at the longer amplifier lengths) and are listed in Table

3-5. The x-ray spectra were dominated by neon-like $n = 3$ to $n = 2$ and $n = 4$ to $n = 2$ emission.

On examining the behavior of germanium amplifiers under varying irradiances, we maintained a scaling relation between the target areal density and the irradiance. Using a fixed reference value of $30 \mu\text{g}/\text{cm}^2$ at $6 \times 10^{13} \text{ W}/\text{cm}^2$, we varied the areal density as the square root of the irradiance. Thus, at $2.4 \times 10^{14} \text{ W}/\text{cm}^2$ (a so-called “4x” target), the areal density was $60 \mu\text{g}/\text{cm}^2$, and so on. This was done to keep the pump laser “burn-through” condition centered in time to the peak of the pulse intensity, in accordance with the hydrodynamic calculations of London and Rosen.¹⁴⁹ Table 3-6 illustrates the scope for this study, showing the 12 shot numbers distributed in irradiance-length space.

The first (and quite remarkable) result is that amplification in the XUV was observed for all shots in the matrix. The shots at 1.5 times nominal irradiance resulted in the brightest amplified lines. Both of these results are similar in character to some British data, in which thick targets were shot single-sided at $1.06 \mu\text{m}$.¹⁴⁷ Lasing similar to the nominal case but reduced in amplitude continued to be seen at two and four times nominal. However, at eight times nominal irradiance ($5 \times 10^{14} \text{ W}/\text{cm}^2$), we noted a dramatic and unexpected shift in the timing of the amplified lines. For this determination, the time-resolving character of the SFFS proved invaluable. Although there was no absolute timing fiducial on the streak camera (this observation has recently been validated on measurements which did include an absolute time fiducial), it was clear from both the onset of continuum and from the measured shot-to-shot jitter of the streaks that the $J = 2$ lines (the only lines visibly amplified) had shifted to later time and that, relative to the nominal-irradiance condition, the shift in peak ASE output was 320 ± 70 ps. This effect can be seen in the SFFS temporal lineouts in Fig. 3-53 and was clearly reproducible. Assuming that the nominal case was identical in absolute timing to that measured for selenium, the eight-times-nominal-case ASE peaks just as the pump laser turns off. This temporal behavior strongly suggests a recombination-dominated population kinetics.

It is also interesting and somewhat surprising that no fluorine-like lines have undergone ASE in the entire matrix, although some fluorine-like XUV lines which do not lase may be present in the twice-nominal and higher

Figure 3-51. XUV spectra from three silver amplifiers of different lengths. The extremely rapid rise of the amplified neon-like lines above background is indicative of a very high gain coefficient.

4 cm nominal Ge XUV amplifier

Figure 3-52. XUV spectrum from a long germanium amplifier shot at approximately $9 \times 10^{13} \text{ W/cm}^2$. These data are from the time-gated on-axis MCP-IGS spectrograph, with an instrumental resolution ($E/\Delta E$) of 1100. The gain has been deliberately set high to record the three weaker amplified lines. The two bright $J = 2$ lines at 23.2 and 23.6 nm are saturated (their intensity exceeds the instrumental dynamic range); on this same intensity scale, they would intrinsically be about four times brighter than recorded. The broad, low pedestal seen at the base of these bright lines is an instrumental artifact associated with detecting such intensities. The brightest spontaneous-emission line in the spectrum is the $2s_{1/2}$ - $2p_{3/2}$ Na-like line at 22.65 nm; it is barely discernible on the scale shown.

Table 3-6. Matrix of the germanium ionization study.

Irradiance	Length (cm)		
	2 cm	3 cm	4 cm
$(6 \times 10^{13} \text{ W/cm})$			
0.5		•	
1.5	•	•	••
2		•	••
4		••	
8	••		

Bullets indicate that Nova laser shots have been taken.

Figure 3-53. Temporal lineouts of the 23.6-nm $J = 2$ line from 2-cm neon-like germanium amplifiers, recorded by the streaked flat-field spectrograph. The black trace is from a shot of near nominal irradiance, about 9×10^{13} W/cm². In the blue trace, the irradiance was about 5×10^{14} W/cm², around six times higher. The peak of the XUV line intensity shifts by 320 ± 70 ps. Some background line and continuum emission is seen (especially evident in the higher irradiance shot as a long-lived pedestal on which the lasing line sits) because of the low spectral resolution of the instrument.

irradiance spectra. The analysis has been hampered somewhat because of $n = 2$ to $n = 2$ lines in L-shell ions detected in second order on the axial MCPIGS. Different gratings and filtering, which will be used in future experiments, are expected to alleviate this problem at least partially. The x-ray spectra have shown that under four times nominal irradiance, the ionization balance is running well into the L-shell, with identifiable lines from ion species as stripped as boron-like.

Experimental Program for Next Year.

Clearly, our experiments have just scratched the surface in addressing the experimental challenges associated with understanding these neon-like XUV amplifiers. We are planning an ambitious program for next year to more thoroughly probe the phenomena studied in 1989 and to expand the scope of the investigation to touch on some of the long-outstanding questions. We plan to measure the line width and line shape of the amplified XUV lines, in particular the $J = 0$. These measurements will test whether Dicke narrowing or other physics influencing line shape are present and whether they might be affecting macroscopically observed gains. To make the measurements, Lawrence Berkeley Laboratory (LBL) researchers are developing a new instrument. This high-resolution XUV spectrograph will, like the SFFS, use varied-line-space grating to focus to an x-ray streak camera for time-resolved detection and will be placed in the on-axis location used by the SFFS. Resolutions ($E/\Delta E$) of about 25,000 are expected.¹⁵⁰ We will also try to replicate the results of References 12–15, i.e., using a 1.06- μm laser for pumping and solid targets, as a starting point for investigating $J = 0$ gains under different laser-matter couplings. A second major new effort will be made to produce purely recombinative lasers by means of short pulse excitation. A new oscillator at Nova should provide 20-ps, 100-J pulses sufficient to drive mid-Z ions to or past the neon-like state. The dominant processes following this excitation would be recombinative at very high electron densities (since the plasma would not have had much time to expand), making them quite conducive for recombinatively fed XUV amplification. This type of experiment would effectively “force” the population kinetics into one predominant mode. An alternative experiment, also proposed for

next year, calls for long-pulse (>5 ns) excitation of very thick “slab” targets in an effort to set up steady-state hydrodynamic and radiative conditions. Each of the two extremes of pulse length should be simpler in principle to model than the present conditions (perhaps better optimized for amplifier behavior) of 0.5-ns pulses and exploding foil geometry. From an experimental standpoint, we must measure the entire angular footprint from these sources because in either case, the beam may not follow the mechanical geometry of the target. Therefore, we are modifying the on-axis MCPIGS to incorporate a 60-mrad imaging fore-optics assembly in front of the instrument (consisting of a cylindrical mirror and a turning mirror). With this modification (which is similar to a modification reported in 1987 by Whelan et al.¹⁴⁴), a measurement of the angular divergence of the beam can be made by collimating. Also, if it is of interest, we will be capable of forming a spectrally resolved, one-dimensional spatial image of the amplifier.

We will continue to study silver amplifiers. A major concern here is the development of a von Hamos high-resolution crystal x-ray spectrograph for distinct, uncompromised detection of the E2 lines for level population measurements. Also, we will study the interesting areal density scaling seen for silver. In germanium, we will continue to look for gain in fluorine-like lines. We have the pump power to shoot a 7.5-cm target at twice-nominal irradiance or a 4-cm target at four times nominal irradiance. Low gain lines may show up or they may not, resulting in yet another open question. We want to probe the irradiance conditions between the “4x” and “8x” cases to try to determine where and how the shift in amplifier timing occurs. To do this, we will set up absolute timing fiducials on all streaked instruments, in particular the SFFS, to accurately (to within a few tens of picoseconds) cross-check the XUV and x-ray emissions to the pump pulse. Finally, we will try to measure the electron temperature in the amplifier plasma by admixing in some K-shell emitter such as aluminum.

Authors: G. M. Shimkaveg, D. J. Fields, R. S. Walling, A. L. Osterheld, A. R. Fry, F. R. Kelly, R. E. Stewart, D. L. Matthews, W. H. Goldstein, and M. J. Eckart

3.5.4. Development of X-ray Holography for Imaging Biological Microstructures

Much of modern biology is directed toward the analysis of the function of macromolecular structures (e.g., chromatin-enzyme replication complexes, nuclear pore structures, structures of protein complexes, elements of the mitotic apparatus) smaller than the resolution limits of light microscopes (~ 2000 Å). This trend should steadily increase as the international genomics effort provides genetic information about the proteins that mediate cell function. The human genome effort itself requires the development of new techniques to reduce the time and cost of sequencing the entire human genome. We have been developing imaging techniques based on coherent scattering of x-rays as potential tools for providing high resolution images of biological microstructures and for providing more cost effective methods of DNA sequencing.

Microimaging with X-ray Laser

Holography. Analysis of biological structures ranging in size from a few hundred Å to $0.1\ \mu\text{m}$ is difficult to achieve using existing technologies. Visible light microscopy is limited to analysis of structures larger than $0.1\ \mu\text{m}$ by the diffraction limit of light. Electron microscopy and scanning tunneling microscopy can only be applied to analysis of structures that have been disassembled to facilitate probe beam access and stained to reveal selected structures. In addition electron microscopy requires fixation with electron dense crosslinking agents and dehydration. NMR spectroscopy can only be applied to analysis of homogeneous structures and diffraction approaches can only be applied to crystallized molecules. New approaches to intermediate resolution analysis are clearly needed. X-ray laser holography is an attractive tool for analysis in this resolution range. It can be applied to analysis of few hundred Å sized macromolecules inside more complex structures (e.g., living cells) as long as the target structures can be labeled. The short, sub-nanosecond image acquisition times that can be achieved using x-ray laser holography minimize the effect of x-ray damage on image quality.

Previous research efforts in development of x-ray laser holography for imaging of biological microstructures has concentrated on x-ray scattering physics, numerical simulations of holograms, and the development of numerical reconstruction procedures for x-ray holograms.

These were undertaken to evaluate the theoretical feasibility of x-ray holography for biological studies.

The optimal wavelength for producing an x-ray hologram of biological objects in water has been determined by a detailed investigation of the physics of coherent x-ray scattering from biological materials such as protein. The optimal wavelength is the x-ray wavelength that maximizes the coherent x-ray scattering and simultaneously minimizes the x-ray absorption. Our analysis of the coherent scattering cross section for a 300 Å diam protein sphere in water as function of wavelength for protein in water shows that the maximum cross section is at a wavelength just above the carbon K edge at 43.7 Å. Operating above the carbon edge at about 44 Å also minimizes the photo-absorption from the carbon thereby minimizing the x-ray heating to the sample. Thus, the optimum wavelength for holographic imaging of biological materials in water is about 44 Å. This wavelength is outside the "water window" between the carbon and oxygen K edges (43.7 - 23 Å) which was previously believed to be the optimal wavelength regime.

The damage caused by the x-rays absorbed by the object during the creation of the x-ray hologram also has been investigated. The production of an interpretable hologram requires that sufficient photons be scattered to allow detection of approximately 100 coherently scattered photons from each object pixel. The heating which the biological object undergoes in order to generate such scattering has been determined using a hydrodynamics model and the calculated cross sections for scattering and absorption. This analysis suggests that the intensity of scattering required to achieve a spatial resolution of 300 Å will result in sample heating to approximately 6000 K. This will cause a rapid hydrodynamic expansion and destruction of the object. This rapid expansion will compromise the spatial resolution of the hologram unless the exposure time is less than 50 ps. For times less than 50 ps, the object pixels do not have sufficient time to expand significantly beyond their initial size. Thus, the hologram must be produced before the x-ray heating induced explosion can destroy the object. These effects force the use of x-ray sources such as x-ray lasers which have the necessary high peak powers for producing holograms on these time scales.

The x-ray damage problem can be significantly reduced if the biological object has structures of interest within it labeled with high x-ray scattering cross section materials such as gold. Microspheres of gold as small as 10 Å in diam can be attached to structures within cells with techniques such as immunogold labeling. This allows specific structures to be labeled and the x-ray fluence required to image these structures can be significantly reduced. The use of 300 Å gold microspheres reduces the required x-ray fluence by a factor of 60. This results in the object reaching temperatures approximately at the boiling point. The resulting expansion is reduced allowing exposure times of a few hundred picoseconds, well within the capabilities of existing x-ray lasers. Using larger gold spheres will result in temperatures below the boiling point and in negligible expansion. For these lower resolution cases, the exposure time is limited by the time scales for biological motion and other x-ray sources such as high brightness synchrotrons could potentially be used.

Computer codes for the simulation of and reconstruction of x-ray holograms have been developed. These codes are used to evaluate different holography systems and to determine their realistic performance capabilities. These codes are validated by testing their predictions and reconstructed images against visible light holography experiments with known objects. These codes have been used to demonstrate for two-dimensional holography: (1) that holograms made with limited numbers of photons can be reconstructed to produce interpretable images, and (2) that gold labeling can significantly improve the contrast of biological objects allowing holograms of complex objects to be made with reduced x-ray fluences.

We have designed a proof-of-principle experiment to demonstrate x-ray laser holography of biological microstructures. The goal of this experiment is to create sub-1000 Å resolution holograms of gold labeled human chromosomes utilizing the existing Ta x-ray laser at 44.8 Å operating in a double pass configuration. Specific sites on a human x chromosome will be labeled with 750 Å diam gold microspheres. The relative positions of the labeled sites within the chromosome will be determined using x-ray laser holography.

The nickel-like Ta x-ray laser, whose 45 Å wavelength is optimal for biological objects, will be utilized in a single view Fourier

transform holography geometry. In this geometry, shown in Fig. 3-54, the x-ray laser illuminates both the distribution of gold spheres in the object and a reference scatterer. A narrow bandpass multi-layer x-ray mirror is positioned between the x-ray laser and the sample. The narrow bandwidth (~5%) of the mirror prevents the broadband soft x-ray continuum produced during the x-ray lasing process from illuminating the sample. A set of pinholes is used to collimate the x-ray laser beam and isolate a coherent region of the beam. X rays which are not scattered, but transmitted by the sample are blocked by a beam stop. Scattered x-rays from the gold spheres interfere with x-rays scattered from the reference scatterer producing the hologram fringe pattern at the detector. This geometry has been demonstrated in visible light tests.

The chromosome will have four or five specific sites labeled with gold microspheres. The gold labels will be attached using *in situ* hybridization with gold labeled, chromosome-specific DNA probes. A gold sphere reference scatterer will be attached using a similar procedure to the chromosome. The labeled chromosome will be placed on an x-ray transparent 500 Å thick boron nitride window. Preliminary experiments will utilize dry chromosomes. Subsequent experiments will use a chromosome in a physiologically equivalent environment. This will be accomplished by placing a boron nitride cover slip over the wet specimen.

Simulations of this experiment have been carried out and the resulting hologram reconstructed numerically to provide a two-dimensional image of the labels in the chromosome. These results show that a signal-to-noise ratio of greater than 500 can be achieved at 750 Å resolution using existing x-ray lasers.

These experiments, if successful, will be followed by multiview holography experiments that will generate true three dimensional images of the labeled chromosomes. These results set the stage for two, more biologically interesting studies.

One study will involve analysis of the organization of the 300-Å DNA fibers within the chromosome. These fibers will be made x-ray dense by incorporating iodine into selected bases along the DNA molecule. This will be accomplished during normal DNA replication so that the organization of the fibers should be

Fig. 3-54. The single view Fourier transform geometry for the initial experiments.

unperturbed by the labeling process. Other proteins will remain unlabeled so that the DNA fibers should have high contrast relative to other chromosomal proteins. The second study will involve analysis of the proximity of newly replicated DNA sequences to enzymes associated with the replication complex in interphase nuclei. The newly synthesized sequences will be labeled with iodine during the replication process so their locations relative to gold references scatterers attached to proteins thought to be part of the replication complex may be determined using x-ray holography. These two studies will provide biologically important information and will allow evaluation of the utility of x-ray laser holography for intermediate resolution analysis of chromatin and other macromolecular structures.

Development of High Resolution X-ray Imaging for DNA Sequencing Analysis. The goal of the Human Genome Project is to determine the sequence of bases in the DNA in the chromosome. Existing methods of accomplishing this are based on electrophoresis and are both slow and costly. Lower cost, faster methods are clearly desirable if the 10^{10} bases in the human genome are to be sequenced. We have been developing high resolution x-ray imaging as a potential lower cost, faster method for DNA sequence analysis.

This method is based on the application of x-ray diffraction techniques with 1.54-Å x-rays to determine the ensemble average of the sequence of labels along the individual elements of crystalline DNA. The crystalline DNA in this application is comprised of a large number (e.g., $\sim 10^{14}$) of identical DNA molecular fragment approximately 1000 Å in length. Each fragment is treated so that all bases of one of the four types of bases in DNA are labeled with an efficient x-ray scatterer. These identical, labeled copies are obtained either by clonal multiplication in a bacterial host, or by an in vitro amplification method such as the polymerase chain reaction (PCR). The labeled DNA fragments are made linear and oriented with their long axes perpendicular to a beam of coherent x-rays (see Fig. 3-55). X-ray diffraction patterns are made with each type of base labeled in turn. The Fourier transforms of the coherent x-ray scattering patterns from four

samples of the same target DNA sequence, one labeled with a DNA base label for adenine, one labeled with a DNA base label for thymine, one labeled with a DNA base label for guanine, and one labeled with a DNA base label for cytosine, gives the information needed for complete DNA sequence analysis.

The interpretation of the relative positions of the labels along the copies of the target DNA sequence is facilitated by labeling one end of each copy of the target DNA sequence with a distinctive, high atomic number label. This allows specification of the positions of the labels relative to the distinctive end label. In general, the end label will be separated from the target DNA sequence by a distance that is comparable to the length of the target DNA sequence. This eliminates the phase problem which makes conventional crystallographic analysis so difficult.

An experiment to test the concept of DNA sequencing by coherent x-ray scattering is currently under construction. Initial experiments will begin in July. This experiment will utilize a rotating anode CuK x-ray source. A carbon crystal monochromator is used to select out and collimate the 1.54 Å CuK line. The sample of oriented DNA sample is placed within this beam. Diffracted x-rays are recorded on a detector placed downstream from the x-ray source. Undiffracted x-rays which pass through the sample are blocked by a beam dump. Adequate spatial coherence is obtained by placing the sample in the far-field of the x-ray source. The target DNA fragment will be 60 Å long with a reference high atomic number atom at one end and two label atoms at the 40 and 60 Å positions. Numerical simulations of this technique have been carried out. Using a simplified model of DNA, the label positions can be clearly recovered from calculated diffraction patterns. These simulations have shown that the locations of 256 labeled bases can be recovered with high accuracy. A more detailed simulation that includes the exact positions of all the atoms in the DNA fragment to be used in the experiment has been developed. The results from this simulation include: calculated diffraction patterns from unlabeled DNA molecules which are comparable to measured patterns; calculated diffraction patterns from labeled DNA molecules which clearly show the effect of

Fig. 3-55. An x-ray diffraction approach to DNA sequence analysis.

the labels on the diffraction pattern; and successful recovery of the label sequence through Fourier transformation from the calculated diffraction pattern.

Author: Jim Trebes

3.5.5. Short Pulse Driven Recombination X-ray Lasers.

Abstract. The advent of compact high peak power short pulse lasers has made tabletop recombination x-ray lasers pumped by these systems a real possibility. In this paper we discuss some of the relevant issues for these lasers and highlight a number of issues that require consideration. Results of simulations designed to estimate potential gains in the quasi-steady-state region are described. The past year has also seen the development of the necessary spectroscopy diagnostic to experimentally examine the soft x-ray spectra from short pulse laser-produced plasmas; sample data from this instrument will be shown.

1. Introduction. Progress in x-ray laser development has been rapid since the first demonstrations of significant gain length products in the x-ray region in 1984.^{151,152} This progress has occurred in both the areas of x-ray laser source development and applications. Amplification has been observed at wavelengths as short as 50.3 Å,¹⁵³ and significant applications in the area of x-ray laser holography¹⁵⁴ and microscopy¹⁵⁵ have been demonstrated.

To date, many demonstrated x-ray lasers rely on medium to large scale inertial confinement fusion class lasers as their pump energy source.^{156–158} It is clear that reducing the size and cost of x-ray lasers will be an important step towards increasing their potential applications. Thus there is very significant interest in demonstrating high output power, short wavelength, tabletop x-ray lasers.

The recombination x-ray laser scheme¹⁵⁹ pumped by a short pulse laser is a natural candidate for such a system, and will be the focus of this paper. It should be noted that other x-ray laser schemes pumped by short pulse lasers have also been proposed.^{160,161} The use of short pulse lasers to produce recombination pumped inversions was first discussed in Refs. 162 and 163. In this work we provide an overview of small scale recombination x-ray lasers, and

discuss how systems pumped by short pulse lasers fit into this overall picture.

The remainder of this paper will be organized as follows. Section 2 provides an overview of small scale recombination lasers, while Section 3 describes in further detail systems pumped by optical field induced ionization. In Section 4 we will discuss our plans for experiments in this area in the coming year and show spectra from a spectrometer recently developed to carry out these experiments. Section 5 summarizes the results discussed in this paper and indicates directions for further work.

2. Small Scale Recombination X-ray Lasers.

Recombination x-ray lasers have been demonstrated in a variety of geometries,^{152,157,158} including those in which the plasma amplifying medium is cooled by expansion^{152,164,165} and/or radiation.^{157,168} While the repetition rate of these systems can be fairly high, significant reductions in size and cost of these recombination lasers are still necessary in order to increase their accessibility for applications.

There are essentially two paths for developing small scale recombination x-ray lasers. The first of these relies on scaling down in pump laser size the somewhat "traditional" approach of irradiating solid or foil targets with relatively long pulses ($t \geq 10$ ps) and allowing the plasma so created to cool via expansion or radiation. Schemes of this type were first used to demonstrate gains in recombination systems pumped by moderate sized solid state^{157, 158} and carbon dioxide¹⁵² lasers of the type used for inertial confinement fusion. Indeed, in recent experiments along these lines amplification has been reported in H-like C¹⁶⁶ at 182.2 Å and Li-like Al¹⁶⁷ at 105.7 Å in plasmas produced by relatively long pulse ($t \sim 3$ -5 ns) tabletop Nd-glass lasers. These results show progress towards the development of tabletop x-ray laser sources. However, very short x-ray laser wavelengths do not appear to be accessible with these scaled down "traditional" recombination laser schemes. This is due to the difficulty in inverting resonance transitions when the plasma ionization occurs through electron collisional processes.

The second approach to small scale recombination x-ray laser development relies on optical field induced ionization of a gas target to produce a completely fully stripped plasma on a very short time scale, while ideally leaving

the electrons at a very cold temperature.^{162,163} An attraction of this scheme is the possibility of inverting the resonance line, which may lead to very short x-ray laser wavelengths at achievable pump intensities. In addition, schemes such as this provide the ideal way of initially preparing the target medium so as to produce gain later in time on nonresonance transitions of the type more commonly seen in previous work^{152,157,164} on recombination lasers.

In the following section, we discuss some of the issues relevant to the second of these two approaches, and estimate the small signal gains that may be obtainable.

3. Recombination X-ray Lasers Pumped by Optical Field Induced Ionization. In the optical field ionized approach, XUV gain is produced on two time scales. The first is called the "transient" case,^{159, 162, 168} and refers to the above mentioned possibility of producing gain on resonance lines. This is possible due to the strongly nonlinear nature of field induced ionization, which provides for the possibility of completely eliminating any population in the lower level of the resonance line. A simple picture of such a system is that immediately after the pump laser pulse, inversion occurs in H-like ions on both the 2-1 and 3-2 transitions during the initial cascade of population through the excited levels of the ion in question. The small signal gain in the transient regime thus persists for on the order of a radiative lifetime of the lasing transition. The major leverage for short pulse laser pumped x-ray lasers is thus their ability to access short wavelength resonance line x-ray lasers at high efficiency by use of this "transient" regime.

The second time scale of interest is that of the "quasi-steady-state" (QSS) regime,¹⁶⁹ which can be as long as hundreds of picoseconds. This is the regime in which previous "traditional" recombination x-ray lasers pumped by larger drivers have operated. In this case, the gain increases as the plasma expands and cools, and peaks well after the pump laser pulse is over; inversions are produced between excited levels (such as the 3-2 transition in H-like ions). The time scale for the gain duration is thus of order a recombination time, as opposed to a radiative lifetime for transient inversions.

In order to produce gain using these ultra short pulse lasers, two conditions must be satisfied. (A detailed discussion of both of these issues can be found in Refs.162 and 163.) First of

all, the target media must be ionized to the fully stripped state (or to the He-like stage in the case of Li-like recombination schemes). For the experiments under consideration here, this occurs primarily through optical field induced ionization operative in the tunneling regime.¹⁷⁰ Secondly, the residual electron energy arising due to above threshold ionization (ATI)¹⁷¹ and other effects must be low enough so that the system is far from ionization equilibrium, and strong three-body recombination commences immediately. This second criterion is of special importance for resonance line x-ray lasers, although maintaining a cold electron distribution is also important to recombination laser performance in the QSS regime, as will be seen below.

Estimates of the required laser intensity to strip various XUV candidate laser ions of interest have been given in Ref. 162. As an example, this model indicates a laser intensity between $5 \times 10^{18} - 10^{19}$ W/cm² is necessary to fully ionize carbon, which is necessary for producing gain on the 33.74-Å 2-1 resonance line. For a given candidate x-ray laser wavelength, slightly less pump intensity is required for a Li-like ion, as compared to a H-like ion. Similar behavior is observed for longer pulse systems.

Estimates of the residual electron energy have been discussed by several authors. In Ref. 162 the electron energy distribution arising due to ATI has been discussed in some detail, and it is concluded that $E_{av} \sim 0.1 \cdot E_{pond}$, where E_{pond} is the average ponderomotive potential seen by the ions. More recent detailed PIC simulations including relativistic and space charge effects basically agree with this conclusion,¹⁷² provided that the relationship

$$1/t_{laser} < \omega_{pe} < \omega_{laser} \quad (1)$$

is satisfied.

Two other effects of potential importance for these systems are inverse bremsstrahlung absorption and stimulated Raman backscatter.¹⁶² With respect to the former, simple estimates suggest that subpicosecond lasers are important so as to avoid significant electron heating as the electron heating rate is typically a few tens of eV per picosecond. Stimulated Raman backscatter growth times appear to be of the order of $1/\omega_{pe}$, which is of the order of tens of femtoseconds. The resulting hot electrons may degrade the suitability of the gas media for

recombination lasers. All of these issues require further study and are currently under examination. It seems clear, however, that shorter wavelength light is preferred, as well as pulsewidths as short as possible. Both of these conditions will act to minimize the residual electron heating.

We now turn to a discussion of the expected gains for these systems. While the observation of gain at short wavelengths on resonance lines is of primary interest, demonstrating amplification at longer wavelengths in the QSS regime is a useful first step towards validating the feasibility of these short pulse laser recombination schemes. In addition, study of the plasma behavior in the QSS regime should be somewhat more realizable experimentally, especially in initial work.

In order to estimate the gains that may be expected in this case, we have written a computer program to simulate the QSS phase in these plasmas. (Extension of these calculations to the transient regime is underway.) The plasma is assumed to be formed at an initial electron temperature T_{e0} and radius R_0 ; the resulting plasma is allowed to cool by adiabatic expansion. The code combines a hydrogenic kinetics code with a similarity solution for the hydrodynamics. The initial temperature is chosen based on estimates for the magnitude of the various electron heating mechanisms discussed above. Trapping of resonance radiation is considered, including the effects of bulk Doppler shifts.¹⁶⁵ Results from this code have been compared to the published work of Pert¹⁶⁹ and the agreement is very good.

Figure 3-56 shows the results for a simple calculation for H-like Li. Li was chosen as an element for study because fully stripping Li should be quite feasible in initial experiments (the required intensity from Ref. 162 is approximately 10^{17} W/cm².) The $n = 3-2$ line is at 729 Å, while the 2-1 transient gain candidate line is at 135 Å. Note also that the candidate 4f-3d lasing line in Li-like N is at 748 Å, and requires nearly the identical intensity as Li; by extension it is thus a candidate lasing line also although we present detailed calculations only for H-like Li.

Note from Fig. 3-56 the relatively long duration of the gain compared to the hypothetical short pulse pump laser; this demonstrates well the different time scales between the QSS and "transient" regimes. Note

also the relatively high small signal gain. Figures 3-57a and 3-57b show the dependence of the peak gain on the initial plasma radius and temperature. In Fig. 3-57a, the gain is cut off at low T_e due to the increasing importance of radiative trapping of the 2-1 line (for simplicity $T_i \sim T_e$ is assumed, and the resonance line is taken to be Doppler broadened). Note that one can to some extent trade initial density for initial temperature.

Figure 3-57b shows the dependence of the peak gain on initial electron temperature. A crude estimate of the initial temperature for 0.5-μm irradiation in 0.5 ps yields $T_{e0} \sim 50-80$ eV, with inverse bremsstrahlung absorption and residual ATI energy each contributing roughly equally. It is thus apparent that electron heating will be an important issue; this also points out the importance of reducing the laser pulsewidth and wavelength to as small values as possible.

Similar calculations to those above have been carried out by Burnett and Enright¹⁶³ for both the QSS and transient case. These results indicate significant gains (of order 5-200 cm⁻¹) in both the transient and QSS regimes. In addition, in this work the importance of thermal conduction for high temperature, small initial radius plasmas is also pointed out. Such a situation may occur when the laser heated region is surrounded by a cold gas, as is typical of an experimental gas target geometry. Thermal conduction has the effect of reducing the electron cooling time and thus increasing the gain; this may thus act to reduce the severe effects associated with a high initial T_{e0} that are shown in Fig. 3-57a.

4. Future Experimental Plans and Diagnostic Development. Our first experiments in this area will be motivated by the considerations discussed above. For our initial work we will be considering the Li-like ions of N and Ne as well as H-like He as candidate lasing ions. Based on the criterion presented in Ref. 162, the required pump laser intensities to produce the necessary He-like N and Ne and H-like He ions are approximately 10^{16} W/cm², 10^{17} W/cm², and 3×10^{15} W/cm², respectively. Intensities of this order should be readily obtainable in the near future using the high power picosecond laser system under development at LLNL. The candidate x-ray laser ions and wavelengths are shown in Table 3-7.

Fig. 3-56. Time evolution of T_e , N_e , and 729-Å (3-2 transition) gain in H-like Li for the quasi-steady case. Initial conditions: $T_{e0} = 50$ eV, $N_{e0} = 5 \times 10^{19}$ cm⁻³, $R_0 = 10\mu\text{m}$.

Fig. 3-57a (Left). Dependence of the peak quasi-steady 3-2 gain in H-like Li on initial plasma radius. For this case, $T_{e0} = 50$ eV and $N_{e0} = 5 \times 10^{19} \text{ cm}^{-3}$. Fig. 3-57b (Right): Dependence of peak quasi-steady 3-2 gain in H-like Li on initial plasma temperature. Here $R_0 = 10 \mu\text{m}$.

Table 3-7. Candidate x-ray laser ions and associated wavelengths.

Ion	1_{OSS} (Å)	$1_{\text{Trans.}}$ (Å)
Li-like N	748 Å (4f-3d)	209 Å (2s-3p)
	512 Å (5f-3d)	
Li-like Ne	292 Å (4f-3d)	88 Å (2s-3p)
	200 Å (5f-3d)	
H-like He	1640 Å (3-2)	304 Å (2-1)

Our initial experiments will be aimed at looking at fluorescence emission from these ions in plasmas produced by the interaction of a high power picosecond pulse with a high density ($N_i \sim 10^{19}$ ions/cm³) gas plume provided by a supersonic gas jet. Examination of the resulting spectra will yield information regarding the plasma kinetics and state of ionization. In particular, we will search for evidence that the excited states are populated primarily by recombination processes. Should these experiments look promising, in the future we plan to attempt proof of principle soft x-ray amplification using an elongated gas jet geometry.

One of our accomplishments in the past year has been to develop and test the spectrograph necessary for these experiments. In general, for looking at spectra from these short pulse laser produced plasmas, both high sensitivity and high data acquisition rates are required. In addition, adequate spectral resolution (typically 1 Å) to identify the lines of interest and temporal resolution are necessary. The instrument chosen for this application is a variation of a streaked soft x-ray spectrograph (SFFS) used in recent LLNL x-ray laser research.¹⁷³ This instrument utilizes a variable line spaced grating to produce a flat field diffracted image of the incident x-rays which is suitable for coupling to a streak camera; details of the spectrograph have been presented elsewhere.¹⁷⁴ Operation of this instrument in a slitless mode provides for the necessary sensitivity. A 12-μm/pixel CCD camera coupled to image grabbing software on a Macintosh computer is used to record the image off the streak camera intensifier; this allows for rapid data acquisition rates.

Simple estimates suggest that the soft x-ray fluorescence yields arising from 1 J, 1 ps irradiation of a dense gas jet of the type to be used in future x-ray laser experiments is comparable to that obtained from the blow-off plasma produced by 1 mJ, 1 ps irradiation of solid targets. Hence, as a demonstration of the required spectrometer sensitivity, we have recorded time resolved spectra from a 1 mJ, 1 ps laser produced plasma of known emissivity.¹⁷⁵ Figure 3-58 shows a CCD image of a single-shot spectrum obtained from a preformed silicon plasma irradiated 200 ps later by a 10^{14} W/cm²,

583 nm picosecond pulse. As has been previously conjectured,^{176, 177} the Si V and Si VI lines observed are long lived and superimposed on a much shorter duration (<20 ps) continuum.

These results validate the feasibility of SFFS for obtaining soft x-ray spectra from short pulse laser-plasma interaction experiments. In the near term we plan to introduce a cylindrical focusing mirror which will increase the recorded image intensity by a factor of 50. This should provide sufficient sensitivity to see even weak lines from the candidate x-ray laser ions of interest.

5. Summary. In this paper we have given a brief overview of short pulse laser pumped recombination x-ray lasers. While the study of these systems is in the early stage, it is apparent that these systems appear quite feasible and offer the opportunity for extremely efficient scaling to short wavelength due to the possibility of inverting the resonance lines of H- and Li-like ions. Simple calculations also indicate the promise of these systems for producing high gain in the QSS regime due to the favorable initial conditions for simple plasma expansion arising from the very short heating pulse. A strong dependence on initial radius and temperature is predicted, although the latter may be modified due to thermal conduction effects. Much work remains, however, particularly in the area of quantifying further the electron energy distribution resulting from the pump laser pulse; this is especially important for the transient case. Initial experiments in the coming year will involve looking at fluorescence spectroscopy for the Li-like ions of N and Ne, as well as H-like He. A streaked soft x-ray spectrograph (SFFS) suitable for these experiments has been developed in the past year and shown to have sufficient sensitivity and data acquisition rates for use in these short pulse x-ray laser experiments.

Authors: C. J. Keane and O. L. Landen

Major Contributors: J. N. Bardsley, L. DaSilva, and D. L. Matthews

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Fig. 3-58. Sample streaked data from the SFFS instrument, taken from a preformed Si plasma irradiated 200 ps after preformation by a 10^{14} W/cm², 583 nm picosecond pulse.

3.5.6. Development of Focused X-ray Sources and Their Applications

In conventional laser matter interaction experiments x-ray emission is primarily from the hot-low density coronal plasma. This makes studying the higher density, cooler regions difficult. To alleviate this problem we have developed a facility capable of irradiating solids with x-rays. By using x-rays we can uniformly heat thin foils at near solid densities and avoid the steep density and temperature gradients which arise with laser absorption in underdense plasmas.

The approach we have chosen for achieving high x-ray fluxes is illustrated in Fig. 3-59. Broadband x-rays generated by irradiating a gold target with an optical laser are focused by an ellipsoidal mirror onto a remote target. The advantage of this approach over the more common direct proximity heating^{178, 179} is that we can achieve the same flux at 1.2 as they can at 1 to 2 mm. Because of this large separation between source and sample, we can eliminate debris and optical laser light, salvage the sample after irradiation and easily diagnose the sample. Our current focusing optics intercept a solid angle of $\sim 5 \times 10^{-3}$ sr and has a magnification of 1:1. We currently have 2 ellipsoidal mirrors, one nickel plated and one gold plated, allowing us to irradiate targets with different x-ray spectrum. In Fig. 3-60 we show the calculated x-ray spectrum on target for these two mirrors assuming a 200 eV blackbody as the x-ray source. The high energy cutoff is due to the reduction in mirror reflectivity whereas the soft x-rays are attenuated by the 2000 Å aluminum debris/optical shield.

Initial tests of these focusing optics indicated that x-ray focal spots of 100 ± 20 µm FWHM can be achieved. This measurement was made at the Janus laser facility by irradiating a gold target with 1–20 J of 0.53-µm laser light in a 1-ns pulse and placing x-ray film at the focal plane. The maximum x-ray flux obtained was 6 ± 2 J/cm². For certain applications it may be desirable to irradiate regions larger than 100 µm. With our current optic, it is possible to operate at larger spot sizes, but at reduced flux. By moving the sample away from the focal point an annular region can be irradiated.

One application of this facility is to study the change in reflectivity of Bragg crystals and multi-layered mirrors under high x-ray fluxes.

This is an important issue in laser plasma spectroscopy of intense sources (e.g., ICF targets). In a recent experiment, a tungsten/carbon multi-layer mirror ($2d = 53$ Å) was placed at the x-ray focus and the reflected x-rays were temporally resolved with an x-ray streak camera.

Figure 3-61 shows the time history of the reflected x-ray signal for two different x-ray fluxes. For x-ray fluxes of 0.02 J/cm² (20 MW/cm²), the reflected x-ray signal peaks with the laser intensity (within the 100 ps uncertainty). In contrast at flux levels of 6 J/cm², which is sufficient to vaporize the multi-layers, the reflected x-ray signal continues to increase after the peak of the x-ray irradiation. There are several factors that can account for this result. First, as the multi-layers are heated and the $2d$ spacing increases, the x-ray streak camera, which collects x-rays with Bragg angles of $46.1 \pm 1^\circ$, will see a different spectral region. Initially x-rays with wavelengths of 38 Å are recorded. A 10% change in the $2d$ spacing would shift the Bragg peak to 43 Å near the 44 Å absorption edge of carbon. The importance of this is that carbon contamination of the x-ray optic causes a factor of ~ 5 increase in x-ray intensity across the carbon absorption edge. The higher reflectivity of tungsten carbon mirrors for x-rays below the carbon edge ($\lambda > 44$ Å) would also enhance this effect. Another possible explanation for the increase in reflected x-rays is that ion diffusion in the multi-layer mirror during annealing leads to an increase in the first order reflectivity. An increase in multi-layer mirror reflectivity after laser annealing has been observed by Ziegler et al.¹⁸⁰ Spectrally resolved measurements of the reflected x-ray signal are planned to determine the relative importance of these effects.

Recently, this facility has been implemented on the Nova 2-beam laser facility. Using one beam, we have irradiated an Au target with 3.5 kJ of 0.53 µm in a 1 ns square pulse and achieved x-ray fluxes of $\sim 10^3$ J/cm². At these flux levels, we can produce near solid density plasmas with temperatures of ~ 100 eV. One of the experiments we are planning involves irradiating a thin (1000 Å) Au disk tamped by LiH (see Fig. 3-59) which maximizes the x-ray energy deposition in the Au while reducing the effects of hydrodynamic expansion. The emission and absorption of this tamped and lower z untamped targets can be easily studied with this facility.

Fig. 3-59. Experimental setup for broadband x-ray focusing experiments.

Fig. 3-60. Calculated x-ray spectra on target for gold and nickel plated ellipsoidal mirrors with a 2000-Å aluminum debris/optical shield and assuming a 200 eV blackbody.

Fig. 3-61. Time history of x rays reflected from a multilayer mirror of $2d = 56 \text{ \AA}$ for low and high x-ray fluxes. The absolute intensity levels cannot be directly compared due to differences in filtering and detector sensitivity.

In addition to focusing broadband x-ray sources, we have also investigated the possibility of focusing an x-ray laser to achieve higher fluxes. Our initial designs considered using a 100× demagnification Schwarzschild microscope.¹⁸¹ However, detailed ray tracing of these systems indicated that current multi-layer deposition techniques preclude such high demagnification systems because of the wavelength mismatch between the primary and secondary mirror. That is, the bandpass of each mirror varies across the surface due to the change in multi-layer spacing. This effect is enhanced for systems with small radius of curvature. As a result we have designed a 40× and 20× demagnification system which should allow us to achieve x-ray fluxes of $\sim 10^5$ J/cm² ($\sim 10^{15}$ W/cm²) in a 3-μm focal spot. These intensities are sufficient to study such non-linear optical phenomena as harmonic generation and four-wave mixing.¹⁸² Of particular interest is mixing an x-ray laser with an optical laser since it offers the possibility of producing an x-ray laser tunable over a 1–2 eV range.¹⁸³

As we have illustrated, focused x-ray sources offer a unique opportunity to produce high x-ray fluxes in a clean debris-free environment. Using the existing facility, we hope to begin producing and studying strongly coupled plasma which are of interest to the ICF, astrophysics and atomic physics community.

Authors: L. Da Silva, R. Falcone, D. Matthews and J. Trebes

3.6. Diagnostic Developments

3.6.1. X-ray Imaging with Very Short Aperture Times

During the last few years, several instruments utilizing electronic gating of microchannel plates (MCPs) have been developed and installed on Nova. While having application-specific front ends, these instruments have the common characteristic of taking a snapshot of a dynamic x-ray image with an aperture time of less than 100 ps. In 1989, six more gated imaging instruments were built, and qualitative improvements were made in the performance of these instruments.

Prior to 1989, several instruments¹⁸⁴⁻¹⁸⁶ were developed using gated MCPs for x-ray imaging. This year, we developed a standard gated

imager design suitable for most of our anticipated imaging needs. This design (Fig. 3-62) is intended for efficient fabrication and assembly, while allowing for flexibility of electrode configurations. We have installed five of these imagers onto the Nova Ten-Beam chamber and assembled one more for the University of Rochester for use on the Omega laser facility. At Nova, imagers were installed onto all four channels of the Quad 22x Wölter microscope¹⁸⁷ and a fifth behind a pinhole array forming the West Axial Pinhole Camera (WAXS). The Rochester instrument is similar to the WAXS, having a meander-line electrode and a multiple pinhole array.

On the Quad 22x microscope we record four images staggered in time from 100 ps to >1 ns. The active area of each imager is 15 × 35 mm, and the optical aperture is 90 μm wide. All electrical parameters are controlled by a computer-based system that provides local operator control as well as synchronization with the Nova countdown clock.¹⁸⁸ The WAXS is a manipulator-mounted instrument with an array of 12 pinholes in front of a 6-mm-wide serpentine strip. As a result, we obtain 12 equally spaced two-dimensional images in 650 ps.

We have performed analyses and experiments in order to better understand and improve the performance of gated MCP imagers. As a result, shorter gate times, higher spatial resolution, and a variety of image geometries are now obtainable.

Two models have been used to simulate the temporal history of an electron shower propagating along a channel, a discrete dynode model and a fully randomized Monte Carlo model.¹⁸⁹ For conventional MCPs having a nominal 40:1 length-to-diameter ratio, both give similar results which have been confirmed experimentally. For electrical pulse shapes, the highly nonlinear gain function of the MCP results in an optical gate time only about one-third as long as the electrical gate. As the gate pulse width narrows to the order of the electron transit time through the MCP (about 100 ps for conventional plates), the gain drops dramatically (Fig. 3-63). This is because the electrons are not under the influence of the peak electric field during much of their journey. We have measured gate times as short as 50 ps (Fig. 3-64) with conventional plates, but with a greatly reduced gain. As a result, we are now investigating thinner (L/D = 20) MCPs as a

Fig. 3-62. Sectional drawing of the standard gated imaging module, shown with its film holder.

Fig. 3-63. Gated imager gain and light output drops dramatically as the electrical gating pulse width approaches the electron transit time through the MCP.

Fig. 3-64. The effective aperture of a gated imaging module is measured to be 50 ps.

means for even shorter optical gates. The Monte Carlo model is more applicable here as there is greater variation in both the number and the impact energy of electrons striking the channel walls.

In an effort to make up for the gain lost at very short aperture times, we have investigated enhancing both X-ray quantum efficiency and MCP gain.¹⁹⁰ Cesium iodide (CsI) has been used on MCPs for some time as an efficient, low-energy, X-ray photocathode. We found that by evaporating CsI normal to both the front and back surfaces of an MCP, the gain is also enhanced. The CsI coats the channel walls and is a better secondary emitter than the lead-oxide surface achieved by the normal hydrogen-fired reduction of MCPs.

The spatial resolution limitations of gated imagers have also been investigated. The principal contributors to the resolution are electron energy spread at the MCP channel exit aperture and the combination of phosphor screen and film emulsion thicknesses along with the high numerical aperture optics used to capture the image. (Images are generally recorded on photographic film in contact with a fiber-optic window.) To minimize the effects of electron energy spread we have only a 0.5-mm gap between the MCP and the phosphor screen. The phosphor bias voltage is a minimum of 3 kV, but we run at higher voltages when resolution is important. We now use thinner phosphor screens, 0.6 mg/cm² of P-11 phosphor, and use T-MAX 3200 film (for its thinner emulsion) instead of Kodak 2484 when seeking maximum resolution. As a result of this work, the line spread function of a gated imager is typically 35 μ m FWHM, with a corresponding 50% contrast at a spatial frequency of 12 lp/mm.

In addition to improving gated imager performance by modifying physical parameters, we have also increased the utility of these imagers by developing meander-line strip geometries. The gateable portion of the MCP is defined by a conductive coating that serves as both a photocathode region and an electrical microstrip transmission line. In order to maintain short gate times, the transmission-line characteristics must be maintained for frequencies up to 10 GHz. Typically, this is a straight strip traversing the MCP, allowing up to 40 mm of active length. This year we produced two imagers having meander-line photocathodes, resulting in an active length of

100 mm.¹⁹¹ One of these is the WAXS; the other was delivered to the University of Rochester for use on the Omega laser. In the Rochester imager, 14 pinholes were arranged to fall on the meander line, producing 14 images spaced 50 ps apart (Fig. 3-65). Further work is being done to obtain a more uniform gain along the long meander line.

Gated imagers are an increasingly important diagnostic tool for laser target diagnostics. There are now seven gated x-ray imaging instruments on the Nova Ten-Beam chamber. By the end of 1990, we expect five more to be activated on the Ten-Beam chamber and two more on the Two-Beam chamber. Improvements in temporal and spatial resolution as well as dynamic range are being made incrementally as these instruments are fabricated.

Author: J. Wiedwald

Major contributors: P. M. Bell, R. Bonner, R. L. Hanks, J. D. Kilkenny, O. Landen, G. D. Power, and R. E. Turner

3.6.2. Fast Neutron-to-Light Converter for Burn History Measurements

In the course of experiments to measure the burn history of inertial confinement fusion targets using fusion gamma rays, we discovered fast, neutron-induced signals, but found no evidence of the desired gamma-ray-induced reactions.¹⁹² In this section we report the rationale and the details of our experimental setup for gamma-ray measurements and describe the recorded neutron-induced signals.

Rationale for Gamma-Ray Experiment. Compressed ICF targets filled with deuterium-tritium (D-T) fuel burn with the release of energetic neutrons and alpha particles through the reaction $T(d,n)^{193}\text{He}$. The burn history of such a target can provide valuable information with which to verify models of the implosion physics. To diagnose the burn time and burn duration, it is desirable to look directly at one of the fusion-reaction products. Most burn-time and burn-history diagnostic instruments are based on the detection of the 14-MeV neutrons.^{194,195} Because most fusion neutrons escape from the target without collision, temporal information about the burn is preserved at a detector outside the target; the recorded signal is delayed by the neutron flight time. Thermal motion of the plasma ions Doppler broadens the neutron energy spectrum (and thus the time-of-flight spectrum).

20-00-0989-2898A

Fig. 3-65. A 14-frame series of X-ray images using a meander-line electrode. The first image is at the bottom left, last at upper right. Images are spaced 50 ps apart.

To obtain signals from which the burn history can be deduced, neutron detectors must be close enough to the target to prevent excessive temporal broadening. Also, because the neutron speed is much less than the speed of light, detectors must be thin enough to prevent excessive temporal dispersion within the detector volume; this limits detector efficiency.

The 16.7-MeV gamma rays produced in the minor branch of the D-T reaction, $T(d,\gamma)^{196}\text{He}$, could provide an alternative measurement technique. Burn history would be deduced from gamma ray rather than neutron signals. A gamma-ray-based measurement is attractive because the gamma rays travel at the speed of light, are unaffected by the plasma temperature, and have a large interaction cross section in many materials. Thus, the temporal issues associated with neutron based measurements are not important. However, because of the small branching ratio¹⁹⁵ (5×10^{-5}), sensitivity becomes the dominant concern for ICF targets whose maximum neutron yields are just above 10^{13} . Now detector distance and volume must be selected to maximize detection solid angle and detector efficiency.

Details of Gamma-Ray Measurement.

Figures 3-66 and 3-67 show our experimental arrangement and illustrate the gamma-ray measurement concept. The heart of the system is a cone of Pb glass (Schott SF7) that acts as a gamma-ray-to-light converter. The burning fuel isotropically emits gamma rays, some of which enter the converter and interact with it via pair production. The resulting electron-positron pairs have a total kinetic energy of nearly 16 MeV. In the few picoseconds required for these relativistic particles to slow down to nonrelativistic speeds, they produce about 200 photons of Cerenkov light per MeV of initial particle energy. The Cerenkov light is preferentially emitted at angles peaked about 54° with respect to the particle trajectory. The cone shape was chosen to reflect a portion of the light towards the output surface of the converter. An $f/10$ Cassegrain telescope images the light to a high-speed streak camera for recording.

The vertex of the converter was placed close (8–26 mm) to the target to maintain a large detection solid angle. A 2-mm-thick Hevimet shield (90% tungsten alloy) to the front and side of the converter prevents target x-rays with energies below 200 keV from reaching the converter. The entire optical path from the

converter to the streak camera is enclosed to keep stray light from entering the detector.

The optical system was selected because the telescope already existed for another experiment. The telescope images light from the converter onto the slit of the streak camera. Light from the telescope's secondary mirror diverges slightly and travels 35 ft to a mirror that reflects the light to an aspheric lens for final focusing onto the slit. The beam passes through a hole in a 6-ft-thick concrete wall that acts as a shield to prevent fogging of film caused by neutron interactions with the camera on high-yield shots. The telescope collects less than 0.2% of the converter light; a larger collection solid angle would greatly increase detection efficiency.

Overall, system response should be less than 20 ps full width at half-maximum. The streak camera response is about 15 ps FWHM, and the 15-mm-long converter contributes about 10 ps to the temporal dispersion because of the difference in the speed of 16.7-MeV gamma rays and 350- to 450-nm Cerenkov light in Pb glass. At Nova, a fiducial system distributes energy from a single 100-ps optical pulse to various instruments for cross-timing their time bases. We pass a fiducial pulse through an etalon to produce a series of pulses 10 ps apart and inject the pulse train into the streak camera. This provides us with cross timing for other instruments and a calibration of our streak camera time base.

Results—Neutrons, Not Gamma Rays! Good streak camera images were recorded in experiments performed at the Nova Ten-Beam facility. The targets were 1-mm-diam spherical glass shells 2-mm-thick filled with 25 atm of equimolar D-T gas. To achieve an implosion, we directly irradiated each target with about 35 TW of 351-nm light in a 1-ns temporally square pulse. Signals were recorded for five targets, whose yields ranged from 5×10^{12} to 2×10^{13} D-T neutrons. Each target implosion produced a clean, fast signal like that shown in Fig. 3-68. The raw signal widths range from 175 to 225 ps FWHM. The signals rise cleanly from background, reach a peak, and drop to background with a somewhat longer fall time than rise time.

Two experiments show that the signals are induced by neutrons, not fusion gamma rays. In the first experiment, implosions were done with the vertex of the converter placed 8 and 22 mm from the target. For signals induced by fusion

Fig. 3-66. The conversion process and experimental setup of the converter and shielding. The converter cone has a vertex half angle of 26.5° and a base-to-vertex distance of 15 mm. Its surfaces are polished and at a slightly different angle than the inside of the shield. This allows total internal reflection of Cerenkov light from the air-glass interface.

Fig. 3-67. Experimental setup of the optical system.

Fig. 3-68. Typical burn-history signal from an image recorded with a high-speed streak camera for a target that produced 10^{13} neutrons.

gamma rays, the change in distance would produce no change in the time interval between the signal and the burn time. The signals recorded for the 22-mm distance were 235 ps later than those for the 8-mm distance relative to the burn time measured with our standard microchannel-plate neutron-emission time system¹⁹⁴; just the shift expected for a neutron-induced signal.

In another experiment, we determined the time interval between the signal and the incident laser power by cross-calibrating the streak camera time base to the incident laser power. For this target, we assumed a neutron-induced signal and determined the burn time to be 1050 ps after the leading edge of the laser pulse, in good agreement with the 920 ± 75 ps determined with the neutron-emission time system. The estimated burn time would be 1530 ps had we assumed a gamma-ray induced signal. The streak camera images show no evidence of signals at the time when the gamma rays arrive at the converter.

Analysis of our weakest signal (Fig. 3-69) provides additional insight into the system response and the light-coupling efficiency between converter and streak camera. This implosion produced the lowest yield (5×10^{12} neutrons) and had the largest target-to-converter distance (26 mm). The streak camera gain was set high enough so that it could detect individual photoelectrons emitted from its photocathode. We estimate the number of detected photoelectrons to be about 5000. With chi-squared fitting techniques, we determined the detector temporal resolution to be 18 ps, just the resolution of the streak camera, and the number of detected fusion events to be about 3000. Thus, only one or two photons per detected fusion event contribute to the streak camera signal.

We saw no hint of a signal induced by fusion gamma rays in the Pb glass. We did, however, observe a fast, neutron-induced signal. Why we did not observe any fusion gamma rays remains a question. The experimental setup is sensitive enough that only a few events are required to produce a detectable signal. Indeed, our weakest signal represents only 3000 neutron events. Since the cross section for pair production is ten times greater than for a neutron interaction, and the energy conversion to light is more efficient for gamma rays than for neutrons in Pb glass, we would expect to see some gamma-ray events with our system even with the low branching ratio,

especially with the detector 8 mm from a target yielding 2×10^{13} neutrons. Perhaps the neutron-to-light conversion is more efficient than expected, or we need a better collection geometry for the Cerenkov light. Additional work is needed to detect the fusion gamma rays.

We did observe clean, fast, neutron-induced signals from the Pb glass with widths ranging from 175 to 225 ps FWHM. Quadrature removal of an estimated 125-ps dispersion caused by neutron flight time in the converter suggests that the detector response to a single event convolved with the target burn duration is about 125 ps. Since most of this width is probably due to the burn duration,¹⁹⁶ the converter response to 14-MeV neutrons is probably much less than 100 ps FWHM. By improving the converter geometry and the coupling between the converter and the streak camera, we expect that a 100-ps detector system sensitive enough for neutron yields of 10^{11} can be built.

More work is necessary to determine the temporal response and sensitivity of the Pb glass to 14-MeV neutrons and the mechanism by which neutron energy is converted to light. Future experiments are planned to distinguish between two possible conversion mechanisms: (n, γ) reactions followed by the interaction of the resulting gamma rays with the Pb glass; and direct neutron collisions with the Pb glass that excite short-lived scintillation states. Should the Pb-glass response time be much less than 100 ps, the resolution for a Pb-glass based instrument will depend on converter thickness, and thus will be a compromise between temporal response and detection sensitivity.

Authors: R. A. Lerche, M. D. Cable, and D. W. Phillion

Major contributors: G. L. Tietbohl, T. D. Schwinn and G. Mant

3.6.3. Development of Line Broadening Diagnostics for High Density Implosion Experiments

I. Introduction. The broadening of spectral lines has proven to be a useful diagnostic of plasma density in a variety of plasmas produced by lasers and other devices. In particular, the high density plasmas produced in inertial confinement fusion experiments are especially amenable to diagnosis by this technique.¹⁹⁷⁻²⁰⁴ In ICF implosion experiments to date, a variety

Fig. 3-69. Burn-history signal from an image recorded with a high-speed streak camera for a target that produced 5×10^{12} neutrons. Statistical fluctuations cause the observed signal noise. The total signal represents about 3000 detected neutrons. The units of amplitude for Figs. 3-68 and 3-69 are similar.

of dopants (typically Ne, Kr, or Ar) have been added to the fuel in ICF capsules for purpose of fuel density diagnostics. The dopant is typically added in trace amounts so as not to perturb the overall capsule performance; the widths of these lines are generally easily resolved experimentally and serve as a diagnostic of fuel N_e . The typical range of electron densities surveyed is 10^{22} to 10^{23} cm^{-3} for Ne and 10^{23} to 10^{24} cm^{-3} for Ar. In the most recent experiments of this type, time resolved fuel densities in the range of 10^{24} cm^{-3} have been measured using the widths of the Ar He β and L β lines at 3.3 and 3.9 keV, respectively, in capsules indirectly imploded using the Nova laser.

Future ICF implosion experiments on Nova and elsewhere, however, are expected to reach both significantly higher fuel densities ($N_e \sim 10^{25}$ cm^{-3}) and pusher areal densities $\rho\Delta R$. In this regime, it is expected that the usefulness of Ar as a diagnostic may be impaired due to the combination of the widths of the Ar β lines becoming too broad (near 100 eV) to measure reliably and the increased attenuation of the Ar lines due to the higher pusher areal density. There is thus interest in developing broadening diagnostics operating at higher photon energy ($h\nu \sim 5\text{--}7$ keV) and with appropriate widths ($E/\Delta E \sim 1000$).

In this paper, we examine line broadening in Ne-like Xe as a candidate for this second generation density diagnostic. (Results of this work have also appeared in literature.²⁰⁵) Xe is an obvious choice to examine in that it is a noble gas and can be added quite easily to the fuel. Also, simple estimates indicate that $T_e \sim 1.5\text{--}2$ keV should be sufficient to produce Ne-like Xe. Electron temperatures in this range should be achievable in near term implosion experiments.

The remainder of this paper is organized as follows. Section II will describe the background theory and the code used to compute the lineshapes. In Sec. III sample calculations carried out with this code will be shown for H-like Ar, and detailed results presented for Ne-like Xe. Sec. IV will summarize the results and discuss plans for future work.

II. Background Theory and Code Description.

In the past several years, a variety of state of advanced line broadening codes suitable for the general user have become available.²⁰⁶

²⁰⁸ Among these codes are two generalized programs^{202, 203} capable of computing line profiles for complex multielectron emitters given

a set of atomic data for the radiator and basic knowledge (density, temperature) regarding the surrounding ion and electron perturbers. Both of these codes work in the standard quasi-static ion/impact electron approximation, which is described briefly below. (Validity criteria for these approximations have been considered elsewhere.²⁰⁹) Preliminary comparisons between the line profiles generated by these programs have been made,²¹⁰ and generally the two codes give very similar results.

The line profiles shown in this paper were calculated using one of these two programs, namely the TOTAL code. Figure 3-70 shows the schematic levels considered in the calculation. The radiative transition $j\text{--}k$ (typically although not necessarily a $\Delta n \geq 1$ transition) is the line whose profile is of interest. Level i is a nearby level with typically the same principal quantum number and opposite parity. The $j\text{--}k$ transition is broadened by both neighboring ions and electrons; for each, the broadening process is mediated by the strength of the dipole matrix element connecting levels i and j . The $i\text{--}j$ and $j\text{--}k$ lines are known as "Stark" and "radiative" transitions, respectively. In an actual calculation for a given radiative transition, many Stark levels i are typically considered.

We now move to a discussion of the broadening processes themselves, and consider the ion and electron broadening each in turn. The ion broadening is treated in the quasi-static approximation. In this limit, the perturbing ions near the radiator are assumed to be static during the lifetime of the radiative transition under consideration. The radiator sees a static electric field (ion microfield) arising from these ions, which leads to traditional Stark broadening of the upper level(s) of the radiative transition in question. This microfield results in levels of different parity (levels i and j in Fig. 3-70, for example) being mixed. The characteristic value of the microfield is given by

$$E_0 = 2.6eZ_p N_p^{2/3} \quad (1)$$

where Z_p and N_p are the perturber ion charge and density, respectively. The value of the field in Eq. (1) is known as the Holtsmark field, and is very nearly equal to the field produced by a point charge of magnitude $Z_p e$ acting at a distance $R_p = [(4\pi/3)N_p]^{-1/3}$ from the radiator.

Fig. 3-70. Schematic level drawing for line broadening calculations, showing "Stark" and "radiative" transitions i-j and j-k, respectively.

There is a distribution function $f(E)$ associated with the ion microfield, and this is computed using the APEX method.²¹¹ In general, the line profile $I(\omega)$ is given by an integral over the field distribution $f(E)$,

$$I(\omega) = \int P(E) J(\omega, E) dE \quad (2)$$

Here $J(\omega, E)$ is the line profile for a single value of the ion microfield. It should be noted that the advent of APEX represents a major time savings in these calculations, and makes the use of generalized line broadening codes such as TOTAL possible.

The ion broadening may be estimated using several simple formulas. For hydrogenic ions, the dominant broadening process is the linear Stark effect, and the full width half maximum of the broadened line may be estimated as follows:

$$\Delta E(\text{eV}) = 2.3 \times 10^{-16} \sqrt{g_i f_{ij}} \lambda_{ij} Z_p N_p^{2/3} \quad (3)$$

Here all quantities are as explained previously. For nonhydrogenic ions, the major process of importance is the quadratic Stark effect, which may be estimated by:

$$\Delta E(\text{eV}) = 5.7 \times 10^{-39} g_i f_{ij} \lambda_{ij}^2 Z_p N_p^{4/3} \quad (4)$$

In Eqs. (3) and (4), it is important to note that to estimate the broadening, one should choose the perturbing level i with the greatest product $f_{ij} \lambda_{ij}$. Note that this quantity is basically proportional to the electric dipole matrix element connecting levels i and j .

The electron broadening is treated in the impact approximation using a binary collision operator discussed by Griem.²¹² One can simplify this operator, and neglecting strong collision corrections and 3-J symbols the electron broadening may be estimated to first order by (referring to Fig. 3-70),

$$\Delta E(\text{eV}) = 10^{-24} \frac{N_e}{\sqrt{N_e}} f_{ij} \lambda_{ij} \quad (5)$$

where λ_{ij} is in Å and other units are as in Eqs. (3) and (4). Note that the electron broadening is basically proportional to the product of the collision operator and dipole matrix element connecting levels i and j . The relations given in Eqs. (3)–(5) are good to a factor of two or three, and should be used for rough estimation purposes only.

The code thus works as follows. A set of atomic data consisting of energy levels and oscillator strengths is led to the program. The appropriate Stark and radiative transitions are identified; the microfield distribution generated by APEX is used to create a mesh of microfield points. For each field value consistent, the radiator Hamiltonian is diagonalized to find the new Stark shifted energy levels. To each of these new resonances, the electron broadening operator is applied to determine the electron impact width. This process is repeated for each value of the ion microfield. The resulting set of resonances and associated electron widths is then used to generate the final line profile according to Eq. (1).

In the following section, we present results for both Ar and Xe obtained using the TOTAL code.

III. Results: Ar and Xe. In this section we present calculations of detailed line shapes for both K-shell Ar and Ne-like Xe, and refer the results back to the estimation formulas in Sec. II. For all results in this section, Doppler broadening has not been considered, as the Doppler widths (typically 1–2 eV) for both Ar and Xe are less than the pressure broadening effects discussed here.

Figure 3-71 shows the line profile for the Ar L_β line as computed by the TOTAL code as well as the line broadening codes of R. W. Lee.²⁰⁴ The plasma conditions chosen are typical of those obtained in recent Nova implosion experiments.²⁰¹ Note the good agreement between the two models. The underlying processes behind the broadening of the Ar L_β line are shown in Fig. 3-72, where the line profile is shown for a single value of the ion microfield ($E = 1.0$ in units of the Holtmark field). Note that there are a total of nine components due to the symmetry of level shifts in $|m_j|$. Three of the components are not visible in Fig. 3-72 and lie near line center; these have very small amplitude and hence the profile has a central minimum. From Fig. 3-71 it is apparent that the

Fig. 3-71. Computed line profile for the Ar L_{β} line as computed by the TOTAL code and the line broadening program of R. W. Lee.

Fig. 3-72. Profile of the $\text{Ar } L_{\beta}$ line for a, the single value $E = E_0 = 1.0$ of the normalized ion microfield [the normalized field strength E_0 is given by Eq. (1)].

Stark splitting of the $n = 3$ manifold is the primary broadening process.

The analytic estimates given by Eqs. (3)–(5) are in accord with Figs. 3-71 and 3-72. The major perturbing level in this case is the 3s level, and atomic structure calculations show that $f_{ij} = 1.5 \times 10^{-4}$ and $\lambda_{ij} = 3.31 \times 10^5 \text{ \AA}$. Inserting these values into Eq. (3) we find $\Delta E \sim 22 \text{ eV}$, which is in reasonable agreement with the actual width of 30 eV. From Eq. (5) the impact width is 1.6 eV, which also agrees well with the detailed calculation. The latter is smaller than the Stark width, as expected. The electron broadening and distribution of microfields act to fill in the dips in the wings of Fig. 3-72, leading to a complete profile of the type shown in Fig. 3-71. Figure 3-73 shows the full width half maximum of both the Ar He β and L β lines vs N_e , and indicates the density dependence of these linewidths for purposes of line broadening diagnostics.

We now consider results for Ne-like Xe.

Figures 3-74 and 3-75 show Ne-like Xe 3–2 and 4–2 spectra as a function of N_e for $10^{24} \text{ cm}^{-3} < N_e < 10^{26} \text{ cm}^{-3}$. In the lowest density case the standard 2p–3d, 2s–3p, and 2p–3s lines are visible. (Note that no kinetics model is being used here; the relative line intensities are proportional to their gf values.) As the density increases, upper states are mixed through Stark broadening and electron collisional broadening also becomes significant.

In terms of focusing on a particular line for diagnostic purposes, it is evident from Figs. 3-74 and 3-75 that the 2–4 Ne-like Xe lines show more density sensitivity in the regime of interest ($N_e \sim 10^{25} \text{ cm}^{-3}$) than the 3–2 lines. We thus consider as a possible diagnostic the Ne-like Xe 4C (2p–4d) line at 6.237 keV, as experimentally this is expected to be one of the brightest lines in the spectrum. As shown in Fig. 3-75a, the FWHM of this line at $N_e = 10^{25} \text{ cm}^{-3}$ is approximately 10 eV. For density diagnostic purposes this is difficult but possible to measure (recall the Doppler width in this regime is of order 1–2 eV). The required instrumental resolution is 3000–6000, which is achievable. Simulations indicate that the photon energy of this line is sufficiently high so that the emission can escape from the fuel region.

The physics behind the broadening of this line is elucidated in Fig. 3-76. Here the profiles as well as the amplitudes are shown for both the full microfield and the $E = 1.0$ cases.

Examination of these results shows that the 4C line is primarily impact broadened, in contrast to

the Ar lines. The impact width of the 4C line is about 7.5 eV. A small additional amount of broadening arises due to perturbation of the 4C transition upper level due to the Stark effect. This effectively leads to a distribution of energies for the 4C upper state and thus adds approximately 2.5 eV to the FWHM.

This conclusion is also in keeping with the estimates in Eqs. (3)–(5). The major perturbing level in this case is the 2p4d state, with $f_{ij} = .2465$ and $\lambda_{ij} = 164 \text{ \AA}$. Substituting into Eqs. (4) and (5), we find that the quadratic Stark shift is small ($\Delta E < 1 \text{ eV}$) and the electron impact width is of order 9 eV. This is in good agreement with the full calculation shown in Fig. 3-76, and suggests that the small additional increment to the width arising from Stark broadening is due to large values (albeit at low probability) of the normalized ion microfield.

Figure 3-77 shows the computed width of the Ne-like Xe 3C and 4C lines versus N_e for $T_e = 2 \text{ keV}$, and indicates the typical density sensitivity of the 3–2 and 4–2 lines for diagnostic purposes. (The 4C curve is highlighted beyond $N_e = 5 \times 10^{25} \text{ cm}^{-3}$, as simple estimates²¹³ indicate that the continuum is lowered to near $n = 4$ at this point.) Finally, the TOTAL results in Figs. 3-71 through 3-77 do not include ion broadening and other effects. Simple estimates²¹⁴ indicate that ion broadening due to perturber/xenon elastic collisions is roughly of the same order as the electron impact broadening for electron densities greater than 10^{25} cm^{-3} . In addition, for both the ions and electrons, coulomb screening of the perturber during its interaction with the radiator is significant; as an example, it is this effect that results in the linewidths in Fig. 3-77 scaling slower than linearly in N_e . The screening for the electrons is included in TOTAL in the standard manner through use of a Debye cutoff in the collision operator.²⁰⁸ Detailed calculations of the ion broadening including screening effects are underway, and the results are expected to yield roughly a doubling of the widths shown in Fig. 3-77. In any case, observable broadening of the Ne-like Xe 4–2 lines will indicate densities of order 10^{25} cm^{-3} or greater, which is valuable in itself. Finally, from the experimental point of view, the potential increased linewidth also implies a possible significant reduction in required diagnostic spectral resolution.

Note again that the 4–2 line shows greater density sensitivity than the 3–2 line. The 5–2

Fig. 3-73. Inferred N_e vs observed full width at half maximum of the Ar beta lines ($T_e = 1$ keV).

Fig. 3-74. Spectra of 3-2 Ne-like Xe for four different electron densities at $T_e = 2$ keV.

Fig. 3-75. Spectra of 4-2 Ne-like Xe for four different electron densities at $T_e = 89.2$ keV.

Ne-like Xe 4C region: $N_e = 10^{25} \text{ cm}^{-3}$, $T_e = 2 \text{ keV}$

Fig. 3-76. Profiles of the Ne-like Xe 4C line. (Left) Profile of the line using the complete microfield. (Right) Profile of the 4C line for $E = E_0 = 1.0$ in units of the Holtzmark field.

Fig. 3-77. Inferred N_e vs observed full width at half maximum for the Ne-like Xe 3C and 4C lines. ($T_e = 2$ keV). For the 4C line, the region where the continuum is lowered to near $n = 4$ is indicated. Ion and electron impact broadening are considered; ion perturber/radiator collisions are not included (see text).

lines should be yet more sensitive; simple estimates indicate their width to be of order 25 eV. This is currently being examined in more detail. These lines are expected to be quite weak in intensity experimentally, however.

With calculations of the Ne-like Xe line profiles available, it is apparent that the next step in this work is to try and experimentally identify the 3-2 and 4-2 transitions of interest. In recent direct drive experiments²⁰⁵ the 3-2 and 4-2 lines were clearly evident, albeit at much lower fuel density. In the near future we expect to perform indirect drive implosions on Nova using Xe doped capsules to test whether Ne-like Xe can be produced, and to examine the spectra for evidence of broadening. Simulations indicate that it should be possible to produce Ne-like Xe, and thus there is good reason to expect that the measurement of the Ne-like Xe linewidths will allow measurement of the fuel density to be made.

Summary. In this paper, the first use of the TOTAL code to calculate spectra from complex ions has been reported. The code treats line broadening in the standard quasi-static ion/impact electron approximation. For Ar, good agreement was found between TOTAL and calculations carried out using the codes of R. W. Lee. The program was then used to compute line profiles for Ne-like Xe. It was found that the 4-2 lines represent the best candidate for plasma diagnostics in that they combine acceptable widths with reasonable intensity. The width of the 4C (2p-4d) Ne-like Xe line was found to be 10 eV at $N_e = 10^{25} \text{ cm}^{-3}$. The detailed line profile calculations for both Ar and Xe were found to be in good agreement with simple estimates for the ion and electron broadening. The Ne-like Xe lines of interest have been observed in recent Nova direct drive implosion experiments. In the near future, indirect drive experiments are planned to test the feasibility of Ne-like Xe as a fuel density diagnostic.

Author: C. J. Keane

Major Contributors: R. W. Lee, B. A. Hammel, L. J. Suter and A. L. Osterheld

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3.6.4. Soft X-ray Detection with Diamond Photoconductive Detectors

Introduction. High speed photoconductors (PCDs) have been used to measure the x-ray emission from a variety of high temperature plasma sources²¹⁵⁻²¹⁹ using GaAs or InP as the detector material. In the study of x-ray emission from laser plasma interactions, the optical sensitivity of these material requires the use of light blocking filters to eliminate scattered laser radiation which distort the flat spectral response of the detector.²¹⁷ Diamond has been used as a photoconductive material to eliminate optical sensitivity and recent measurements have shown that diamond has a flat response between 200 and 2200 eV. Diamond has an extremely small absorption coefficient for energies less than its 5.5 eV bandgap; and, therefore, it is insensitive to scattered laser radiation with a wavelength greater than 2200 Å. These properties make diamond a suitable detector material for total x-ray power measurements from high Z laser produced plasmas. As an example, we have used diamond PCDs to measure the x-ray yield from uranium disk targets as a function of laser intensity.

Theory of operation. A single carrier model has been shown to be accurate in modeling diamond PCDs.²²⁰ In this model the photogenerated carriers, electrons and holes, are treated as a single average carrier. The temporal evolution of the charge carrier density, n , is determined by a simple linear rate equation

$$\frac{dn}{dt} = \frac{P(t)}{\gamma V} - \frac{n}{t} \quad (1)$$

where P is the absorbed power, t is the carrier lifetime, V is the excited volume, and γ is the average energy to form an electron pair. The sensitivity, S , in A/W can be derived from this expression

$$S = \frac{e\mu\tau V_0}{L^2 \gamma} = \frac{e\mu(\epsilon)\psi\tau\epsilon}{L\gamma} \quad (2)$$

where V_0 is the applied bias voltage, μ is the carrier mobility, ϵ is the applied electric field, and L is the length of the PCD along ϵ . Typical

sensitivities are approximately 6×10^{-4} A/W with carrier lifetimes near 100 ps.

Experiment. The detector was a 1-mm-cube of insulating (IIa) natural diamond with Ohmic contacts on two opposite faces of the detector. The detector was mounted in a high bandwidth 50 Ω housing which terminated an impedance matched transmission line. The 1000 V bias voltage was supplied through a capacitive bias tee with a rise time of 60 ps. One of the uncontacted faces of the detector faced the radiation source with no filters to distort the measurement or the incident x-ray power. A small leakage current, less than 1 nA, flows through the detector without x-ray illumination. The pulsed x-ray signal was recorded on a high speed oscilloscope.

The detector acts as a radiation controlled resistor in the detection circuit where the decrease in resistance is proportional to the absorbed radiation power. If the output signal is small compared to the applied bias voltage, the signal is proportional to the absorbed power. If the signal level becomes a significant fraction of the applied bias then the output of the PCD must be corrected for the finite impedance of the recording system which is typically 50 Ω . This can be corrected analytically with the expression

$$I_L = S\Phi = \frac{I_{out}}{1 - I_{out} \frac{R_0}{V_0}} \quad (3)$$

where Φ is the x-ray power incident on the detector, I_L is the linear or corrected current, I_{out} is the measured detector current. This correction has been applied to all of the data in this paper.

The radiation source was generated from a laser plasma interaction. One arm of the Nova laser was used to irradiate uranium disk targets. The operating wavelength of the laser was 0.35 μm . Typically, 2000 J of laser light in a 1-ns flat topped pulse was focused upon a 1-mm-diam, 50- μm -thick disk of uranium located at the center of the Nova Ten Beam target chamber which is four meters from the diamond PCD. The focused laser intensity was varied from 5×10^{13} to 3×10^{15} W/cm².

The time integrated radiated spectrum and total radiation yield was measured with a photoemissive vacuum diode array (Dante).²²¹ A

low resolution emission spectrum is reconstructed from the output of 13 absolutely calibrated channels covering the energy range from 180 to 2000 eV. The diamond PCD and Dante had nearly identical views of the target, 47° from the target normal. The laser plasma x-ray spectrum is dominated by a broad band emission from 200 to 1500 eV with a mean energy of approximately 700 eV. A few percent of the radiated power appears in the niobium L lines and the uranium M-bands above 2 keV. The emission above 2 keV was monitored with an absolutely calibrated x-ray crystal spectrograph. A typical composite spectrum is shown in Fig. 3-78. The incident power equals the absorbed power because the 1-mm-thick detector stops all x-rays with an energy less than 7 keV.

A sequence of 19 uranium laser-plasma experiments were performed to compare the total emitted x-ray energy emitted by a target as measured with a PCD and the Dante array. The x-ray power incident on the PCD varied from 160 to 6700 W and the linear PCD output current varied from 0.12 to 6.14 A. A plot of the data is shown in Fig. 3-79. An excellent linear relationship exists between the incident x-ray power and the linear current output [from Eq. (2)] of the PCD. A power law fit yields $I_L \sim \Phi^{1.06}$ with $R = 0.99$. Using the absolute calibration of the PCD, the conversion efficiency of the uranium disk target as a function of laser intensity was measured. The conversion efficiency is defined as the total radiated x-ray energy, assuming a Lambertian radiator, divided by the incident laser energy. This is shown in Fig. 3-80. As expected for the x-ray conversion efficiency for a high Z radiator the conversion efficiency decreases as the laser intensity increases.²²² A least squares fit shows that the conversion efficiency scales as the incident laser intensity to the -0.26 power.

A few experiments were performed with shaped laser pulses on gold disk targets. For these experiments, the recording system for the Dante detector array was a set of 6 GHz oscilloscopes. The system response time was approximately 100 ps. An identical oscilloscope was attached to a single diamond PCD. The time resolved Dante measurement was analyzed at discrete times to yield the x-ray power as a function of time. This requires a complete spectral unfold at each time step and very accurate interchannel timing. The PCD data was linearized according to Eq. 2 and compared to the

Fig. 3-78. The emission spectrum from a uranium laser plasma. The spectrum below 2000 eV was measured with a multichannel filtered vacuum photoemissive diode array and above 2000 eV by a crystal spectrograph.

Fig. 3-79. A comparison of the linear current output of a diamond PCD to the x-ray power measured with the Dante photoemissive diode array. A linear relationship is measured. A power law fit to the data scales as $\Phi^{1.06}$, where Φ is the x-ray power incident on the detector.

Fig. 3-80. Uranium x-ray conversion efficiency as a function of laser intensity for 0.35- μ m light. Conversion efficiency is defined as total radiated x-ray energy (assuming a Lambertian source) divided by the incident laser energy.

x-ray power as measured by the Dante. The results are presented in Fig. 3-81. There is excellent agreement in the measured shape of the x-ray power. This is a further verification of the single carrier model (Eq. 1), i. e., linear operation of the detector for excitation pulses many times longer than the carrier lifetime.²¹⁷

Conclusion. The x-ray conversion efficiency of uranium has been measured with one detector, a diamond PCD. These measurements are in excellent agreement with an independent measurement using a photoemissive diode array. Also a measurement of the total radiated x-ray power made with a diamond PCD is in excellent agreement with a time resolved measurement using the array. Diamond PCDs promise to be simple and effective detectors for the measurement of x-ray power. The device is fast and insensitive to visible radiation and simple to operate.

Authors: D. R. Kania, L. Pan, H. Kornblum, P. Bell, O. N. Landen, and P. Pianetta

3.6.5. Multichannel Wölter X-ray Microscope

We have developed a multichannel imaging x-ray microscope for producing time and space resolved high resolution x-ray images of laser fusion experiments. The system utilizes five sectors of a grazing incidence Wölter x-ray microscope²²³ optic to produce images on five independent detectors. Typically four gated microchannel plate and an x-ray streak camera are used for detectors. The gated microchannel plate detectors allow four independent two-dimensional snapshots to be produced at four separate times during the experiment. The streak camera allows a continuously time resolved, one-dimensional image to be produced simultaneously with the gated images. The applications of this diagnostic include laser fusion implosion symmetry studies, hydrodynamic instability studies, and plasma physics studies.

The multichannel x-ray microscope was developed by modifying an existing Wölter x-ray microscope.²²⁴ A Wölter x-ray microscope consists of two rotational symmetric grazing incidence x-ray mirrors. The surface of the first mirror is a hyperboloid, while the surface of the second is an ellipsoid. The object to be imaged is placed at the focus of the hyperbola. The image created by the hyperbola acts as a virtual object for the ellipsoidal segment, which then produces

a real image of this virtual object. This arrangement allows for low aberration, high magnification imaging with grazing incidence optics. The Wölter optic used has a magnification of 22. The x-rays reflect from the nickel coated optic at a grazing angle of one degree allowing the microscope to be used with x-ray energies up to ~3 keV. Typically only a 15° sector of the optic is used instead of the full 360° available. The sector is defined by an x-ray opaque aperture at the front of the optic. The use of a small sector minimizes effect of figure errors on the resolution and results in a measured spatial resolution of ~8 μm .

Multiple images can be obtained by using multiple sectors if the images can be separated in space. A schematic arrangement for accomplishing this is shown in Fig. 3-82. The object can be a self-luminous laser fusion capsule or an x-ray backlit object. Five sectors are defined by an aperture at the front of the optic. Using five sectors simultaneously would result in the overlap of five images on the optic axis at the image plane. This overlap is avoided by deflecting four of the images off axis with grazing incidence plane mirrors inserted behind four symmetrically placed sectors on the Wölter optic. Each mirror was fabricated from fused silica, superpolished, and nickel coated. The measured surface figure is better than $\lambda/15$ and the surface roughness over 1- μm scalelengths is better than 3 Å rms. Significant x-ray scatter is observed from these mirrors, indicating that on 100-Å scalelengths the surface roughness is considerably worse. Two mirrors were required to deflect the images from each sector to allow sufficient space between images for an x-ray streak camera to be placed on the central axis. A single mirror at 2° would reduce the high energy response of the system to well below 3 keV.

The images deflected off-axis by the plane mirrors are detected by gated microchannel plates.²²⁵ The gated microchannel plates have CsI photocathode material coated inside the microchannel to increase the quantum efficiency for x-ray detection. The conducting strip line is coated on the microchannel plate surface. Gating of the plate is accomplished by launching a fast (~200 ps), high voltage (800 V) pulse down the stripline. Due to the nonlinear response of the plate as a function of voltage, gate times less than 100 ps can be obtained. Each of the four microchannel plates can be timed independently with a jitter of less than 50 ps. The spatial

Fig. 3-81. A comparison of the x-ray radiated power as measured by the Dante array with a high speed oscilloscope for each channel and a single PCD recorded on a single high speed oscilloscope. The error bars represent 10% error expected in the Dante measurement.

Fig. 3-82. A schematic arrangement of the multichannel x-ray microscope.

resolution of the channel plates is about 35 μm corresponding to a resolution limit in the object plane of 1.6 μm . The x-ray streak camera is also operated with a CsI photocathode and typically run with a time resolution of 100 ps and a spatial resolution of 80 μm . The spatial resolution in the object plane is 3.6 μm .

Author: Jim Trebes

3.6.6. Thick-Aperture Design for Neutron Penumbra Imaging

Introduction. Penumbra coded-aperture imaging of neutrons is used to determine the spatial extent of fusion reactions in the imploded core of inertial confinement fusion targets. This method has been used,²²⁶ to obtain neutron images of sources yielding $\geq 5 \times 10^{11}$ neutrons at a two-point resolution of 60 μm over a 400- μm field of view. We wish to increase the resolution to 10 μm over a field of view of at least 100 μm to image neutron sources producing $\leq 10^{11}$ neutrons.

High-resolution penumbra imaging of weak sources of penetrating radiation requires an aperture with a particular set of characteristics. First, adequate contrast requires use of an aperture several mean free paths thick. The resulting aperture must often be many centimeters thick. (The mean free path of 14-MeV neutrons in gold is 3.4 cm, for example.) Next, the aperture point-spread function (PSF) must have a sufficiently sharp cutoff. In addition, to obtain an image of a relatively weak source with an acceptable signal-to-noise ratio requires that the aperture be quite close to the source. Finally, penumbra imaging requires numerical reconstruction of a coded image to obtain an estimate of the source. The computational difficulty of the reconstruction is vastly reduced through use of an isoplanatic aperture and linear reconstruction techniques. Therefore, imaging system magnification must be reasonably uniform across a useful field of view.

Simple conical apertures^{227,228} have acceptably sharp PSFs if the desired field of view is much smaller than the aperture diameter. But the resolution achievable with such apertures suffers if (as is generally the case) the vertex of the cone must be moved away from the source plane to provide acceptable isoplanaticity. Better performance is available with toroidal-segment apertures having a parabolic taper,²²⁷ but only when the aperture can be placed sufficiently far from the source (see

below). Here I present a simple analytic approach to aperture design; the result is a taper in which the radius of curvature varies inversely with distance from the source plane, which I call an IRC aperture. Toroidal-segment and conical apertures can be obtained as special cases of the IRC design.

The appropriate IRC aperture to image Nova targets is 6-cm-thick with $\leq 10 \mu\text{m}$ resolution over a field of view of 150 μm at a source-plane-to-aperture-entrance separation of 2 cm. The characteristics of PSFs obtained with such a taper are examined in detail.

Images from point sources are simulated by detailed ray tracing through the aperture. The radiation within the aperture is assumed to be exponentially attenuated with the total mean free path of 14-MeV neutrons in gold. The detailed effects of scattering within the aperture can be neglected when the detector subtends a very small solid angle relative to the aperture, as is generally required when high resolution is desired. In this case, only the fraction of the incident radiation that scatters precisely forward will be incident on the detector, causing a uniform increase in the image background level. Furthermore, if the aperture also subtends a small solid angle relative to the source, the increase in background level becomes negligible. For example, for a 0.64-cm-diam gold aperture placed 5 cm from the neutron source and 20 m from a detector, the background level in the image will be increased by $<0.4\%$, mostly due to the forward peak of the elastic-scattering cross section.²²⁹ Noise and distortion due to quantization or other detector physics are not considered.

The idealized point-spread images are deconvolved with a Wiener filter to determine the effective resolutions and magnifications referred to the source plane. The resolution is $<3 \mu\text{m}$ at the center of the field of view, but it deteriorates to 10 μm at the edge of the desired field of view. Magnification is uniform within $\pm 6\%$ over the field of view. The aperture PSF is sufficiently isoplanatic and distortion free to allow the use of fast linear reconstruction methods in unfolding the coded images.

Thick-Aperture Design Technique. The cutoff of a penumbra-aperture PSF determines the resolution of the aperture. One way to make the aperture isoplanatic, then, is require that the aperture cutoffs corresponding to different points in the field of view be as similar as possible. To

carry this out, it is first necessary to develop an analytic approximation of the cutoff function. I begin with the specific case of a ray intersecting a toroid, because this geometry reduces to the simple problem of a line intersecting a circle in cylindrical coordinates. Solution of the resulting quadratic equation gives the path length of material intersected by the ray as a function of the angular displacement of the ray; this path length, divided by the appropriate mean free path, is the natural logarithm of the cutoff function. Next, I show how to make this cutoff function independent of the position of the source point within the field of view as desired. Finally, I generalize the result from the toroidal geometry by taking the minor radius of the toroid as the variable radius of curvature for a cylindrically symmetric aperture; thus we obtain a differential equation that can be integrated numerically to determine a particular aperture taper. Boundary conditions arising from the desired imaging geometry complete the specification of the aperture.

In Fig. 3-83, consider the intersection of a ray with a circle of radius R centered at the origin of the coordinate system; the ray leaves the source at angle θ . To determine the length Δz of the chord intercepted by the ray we locate the intersection points, which are determined by the equations

$$\left\{ \begin{array}{l} \rho - R = \theta(z - z_i) \\ \rho^2 + z^2 = R^2 \end{array} \right\}, \quad (1)$$

where z_i is the distance between the source and the aperture at the point of first intersection (tangency).

The system is easily solved to give the path length Δz through the circle as

$$\Delta z \cong 2R(1 - q) \sqrt{\theta^2 + \frac{1}{(1 - q)^2} - 1}, \quad (2)$$

where it is assumed that $\theta \ll 1$. In addition, $z_i / R \ll 1$, so that $q \equiv z_i \theta / R \ll 1$. To first order in q , then, we have

$$\Delta z \cong 2R \sqrt{\theta^2 + 2q}. \quad (3)$$

In this analysis, only the initial portion of the cutoff will be considered, $\Delta z \ll z_m$, the mean free path within the aperture material, so $\theta \ll z_m / 2R$. With this restriction, the first term under the radical can be neglected, so we have

$$\Delta z \cong 2\sqrt{2R z_i \theta}, \quad (4)$$

and it is now clear that Δz will be approximately independent of z_i if R varies as $1/z_i$.

Now take R as the (variable) radius of curvature of a tapered cylindrical aperture. Rays from different source points in the field of view will intersect the aperture with different values of $z_i \propto z$. If R_0 is the radius of curvature at the front face of the aperture, and if that face is separated from the source by a distance z_0 , we have

$$R = \frac{z_0}{z} R_0; \quad (5)$$

replacing $1/R$ with the usual expression for the curvature, we obtain a second-order, nonlinear, ordinary differential equation

$$\frac{d^2 \rho}{dz^2} = \frac{z}{z_0 R_0} \left[1 + \left(\frac{d\rho}{dz} \right)^2 \right]^{3/2}. \quad (6)$$

Equation (6) can be written as a pair of coupled first-order differential equations that can be solved by standard numerical methods to yield the IRC aperture taper.

It is desirable to specify the aperture in terms of its initial diameter d_0 , the source-plane-to-aperture-entrance separation z_0 , the aperture thickness z_t , and the field-of-view diameter d_{fov} . The geometric field of view (defined in Fig. 3-84) is chosen to maximize the sharpness of the aperture PSF. Together, these parameters specify the boundary conditions on the aperture taper. In practice, initial conditions (including a guess for the initial radius of curvature) must be specified to begin the calculation. The differential equation is then integrated along z to obtain the values of the taper function and its derivative at the aperture exit. This process is iterated, using a simple root-solving procedure to obtain successive guesses for R_0 , until the desired values are obtained at the exit.

It is interesting to note that the toroidal segment and conical aperture designs discussed in previous work are actually special cases of the IRC taper. In the toroidal aperture, the radius of curvature is constant. This circumstance is an approximation of the IRC taper when $z_t \ll z_0$, so that

$$R = \frac{z_0}{z} R_0 \cong R_0. \quad (7)$$

Fig. 3-83. Geometry of the theoretical development. A ray from a source point cuts a circle of radius R .

Fig. 3-84. Lines drawn tangent to the entrance and exit surfaces of the aperture define the geometric field of view.

In other words, when the aperture is much farther from the source plane than its thickness, an aperture with a constant radius of curvature is satisfactory.

The conical aperture can be viewed as a degenerate case of the IRC aperture in which the field of view is small and/or the aperture is particularly thick. Specifically, when

$$\frac{z_0}{z_0 + z_t} \frac{d_{fov}}{d_0} \ll 1, \quad (8)$$

the slope of the aperture does not vary significantly between the entrance and exit, R_0 goes to infinity, and a conical aperture is obtained.

IRC Aperture For Nova Laser Experiments.

An aperture for Nova laser target experiments must meet several requirements. We expect the targets of interest to yield roughly 10^{11} neutrons. This is a weak source, so to obtain maximum signal in the image we must minimize the source-aperture distance: we therefore take $z_0 = 2$ cm. The choice of aperture thickness is complicated in this case by the trade-off between image contrast and aperture solid angle. Description of the optimization procedure is beyond the scope of this paper: we designed a gold aperture 6 cm thick. Because neutron-emitting regions in a variety of shapes can arise in the proposed experiments, we take $150 \mu\text{m}$ as the field of view. Within the field of view, the desired resolution is $\leq 10 \mu\text{m}$.

These constraints suffice to determine the IRC aperture. (The resolution deteriorates towards the edge of the geometric field of view, so it was necessary to choose a geometric field of view of $240 \mu\text{m}$ to obtain the desired resolution over an effective field of view of $150 \mu\text{m}$.) Figure 3-85(a) shows the IRC taper function. The aperture initially tapers inward because the diameter of the geometric field of view is greater than the entrance diameter. (Use of an aperture with a much larger diameter would improve the image signal-to-noise ratio; the small aperture was used here to allow the high spatial resolution of the present analysis. Similar results for resolution and isoplanaticity have been obtained at coarser resolutions for apertures with diameters as large as 1 mm.) The aperture is relatively conical (that is, it has a large radius of curvature) near the entrance, and flares increasingly towards the exit; this is characteristic of IRC apertures.

The image formation is simulated by tracing rays from specific source points through the aperture. The transmission of each ray is recorded on a fine grid whose spacing corresponds to $1\text{-}\mu\text{m}$ pixel widths in the source plane. In calculating the attenuation in the aperture material, it is assumed that the radiation is continuous (as opposed to discrete particles) and that it is attenuated exponentially with an e -folding length of 3.4 cm. Figure 3-85(b) shows the cylindrically symmetric PSF of the on-axis source point. Figure 3-86(a) plots x-axis rows of the image matrix (x-axis lineouts) of the aperture cutoffs for an on-axis source point and for a source point displaced by $50 \mu\text{m}$; the x-axis scaling has been referred to the source plane by dividing by $1 + M$, where M is the system magnification (defined below), and the lineouts have been superimposed. Two lineouts are shown for the displaced point; they are cuts of the two-dimensional images taken parallel to the displacement. The (+) curve is shown in the correct orientation; the (–) curve has been reversed. The initial sharp portion of the cutoff is quite isoplanatic for all lineouts, but the displaced source point generates an asymmetric PSF, in which one edge has a larger-amplitude tail than the other.

To determine the resolution of the aperture, the coded image must be deconvolved. Here, where the images are generated analytically, it is desirable to use linear techniques that do not extrapolate the data in generating an unfold. Maximum-entropy techniques that have been suggested previously²³⁰ to unfold noisy experimental data are inappropriate. Therefore we use a parametric Wiener filter defined in the Fourier transform domain,

$$A_{\text{ref}}^{-1} \left[1 + \frac{k}{|A_{\text{ref}}|^2} \right]^{-1}. \quad (9)$$

In Eq. (9) A_{ref} is the transform of an idealized penumbral-aperture PSF, the rectangle function $\Pi(\rho/D)$, with a diameter D equal to that of the IRC on-axis PSF.²³¹ The parameter k adjusts the extent of the filtering, the precise nature of which can be exactly determined by inverse transforming the taper function in Eq. (9) (term in square brackets). By keeping k small (2×10^{-4}) in unfolding these analytic images, I avoid significant numerical errors that develop near the zeroes of A_{ref} , while preserving the high-frequency components

Fig. 3-85. IRC design for a neutron penumbral-aperture microscope: (a) taper function; (b) point-spread function.

Fig. 3-86. Resolution and isoplanaticity of the IRC aperture: (a) detail of the PSF cutoff; (b) deconvolved resolution functions (normalized to have a peak amplitude of unity) for various source-point offsets.

in the unfold. I call the deconvolved PSF the *resolution function*; it corresponds directly to the usual notion of PSF in uncoded, one-to-one-imaging systems. I define the resolution as the FWHM of the resolution function.

Figure 3-86(b) shows the resolution functions for on-axis and displaced point sources. For small offsets the resolution is excellent, and there is virtually no distortion; the FWHM for on-axis points and for points displaced $25\text{ }\mu\text{m}$ is $<2\text{ }\mu\text{m}$, the Nyquist limit for the simulated imaging system. Broadening and some distortion appear at an offset of $50\text{ }\mu\text{m}$: the FWHM is $5.3\text{ }\mu\text{m}$, and a double peak is visible. At an offset of $75\text{ }\mu\text{m}$ the FWHM has risen to $10.3\text{ }\mu\text{m}$. Thus the resolution is $10\text{ }\mu\text{m}$ over an effective field of view of $150\text{ }\mu\text{m}$. Note that the broadening is largely due to the change of shape in the central bright region of the off-axis PSFs. Figure 3-87 shows the variation of resolution with source point offset.

Like the resolution, it is easiest and most intuitive to define the magnification of the imaging system with respect to the reconstructed resolution functions. I therefore define the magnification as the ratio of the change in the position of the centroid of the resolution function to the corresponding change in the source-point position. For the Nova IRC aperture, the magnification (Fig. 3-87) is uniform to $\pm 6\%$ within the $150\text{-}\mu\text{m}$ useful field of view. The IRC aperture therefore produces little geometric distortion in the image in this example, at least with respect to its first-order moment. (Higher-order distortions are present, however.)

Summary. A technique for designing thick apertures for high-resolution penumbral imaging of penetrating radiation has been presented. The technique is based on a simple analysis of the intersection of rays emanating from a source region with a circle; the analysis is generalized through the use of a differential equation involving the radius of curvature of a cylindrically symmetric aperture. When the radius of curvature at any point along the aperture varies inversely with the corresponding separation from the source point, a particularly isoplanatic aperture taper is obtained.

The technique has immediate application in a neutron penumbral-aperture microscope planned for the Nova laser facility. A hypothetical aperture design for such an instrument achieves at least $10\text{-}\mu\text{m}$ resolution over a $150\text{-}\mu\text{m}$ field of view with no significant first-order distortion. The performance is

noticeably better than that of previous thick-aperture designs.

The IRC aperture design technique may be useful in other systems for imaging penetrating radiation. In particular, it could be optimal for use in thick-pinhole imaging systems in which a small-diameter aperture must view a relatively large field of view at a small source-aperture distance.

Author: D. Ress

3.6.7. Signal-to-Noise Ratio Calculations for Penumbral-Aperture Imaging

We recently demonstrated a new instrument, the neutron penumbral-aperture microscope, by producing images of the reacting deuterium-tritium fuel in laser-driven inertial-confinement-fusion targets.²³²⁻²³⁴ The 14-MeV fusion neutrons passed through an aperture and formed coded images of the active objects. These coded images were analyzed to produce reconstructed images of the burning fuel. An analytical expression for the signal-to-noise ratio (SNR) in a reconstructed image allows us to rapidly predict the performance of an imaging system with a particular target. In this section, we develop a general expression for the SNR expected in a reconstructed penumbral-aperture image.

Penumbral imaging is a coded-aperture imaging technique that consists of two distinct steps. The first step, which is illustrated in Fig. 3-88, is the formation of a coded image. Conceptually, this is similar to pinhole imaging but with the essential difference that the aperture is larger than the size of the source. A source of incoherent radiation casts a geometrical shadow through a circular aperture, thereby producing a coded image that consists of a bright, uniformly illuminated central region surrounded by a partially illuminated penumbra. All the information about the spatial distribution of the source is encoded in the penumbra. However, because the coded image is the superposition of the source distribution with the aperture transmission function, it does not represent the actual appearance of the source itself. Indeed, the image formed by each point in the source is a disk-shaped pattern whose size and position uniquely characterizes the position of its associated source point. When the aperture transmission function (point-spread function) is spatially invariant (isoplanatic), the

Fig. 3-87. Magnification and resolution of the example IRC aperture vs radial displacement in the source plane. Although the resolution deteriorates with increasing radius, the magnification remains roughly constant.

Fig. 3-88. Concept of penumbral imaging. A source casts a geometrical shadow through a circular aperture to produce a coded image. The penumbra of the coded image contains information about the shape of the source.

superposition becomes a convolution. In the second step, the source distribution is decoded (deconvolved) from the coded image using numerical techniques and a knowledge of the aperture shape.

Our development of an SNR equation for a reconstructed image pixel is based on simple counting statistics. We start by considering the image produced by a source element with a diameter Δs and an aperture whose diameter is r_{ap} and source-to-aperture distance is l_1 . The aperture to image plane distance is assumed to be much greater than l_1 . The image is a uniformly illuminated disk surrounded by a penumbral ring in which the illumination diminishes to a level determined by the transmission, τ , of the source radiation through the material forming the aperture. Information about this resolution element is coded entirely in this ring; no information is contained in the central region or in the uniform pedestal. The average signal (information) count density (neutrons/cm²) across the penumbral ring relative to the count density in the central region is $(1-\tau)/2$. The total number of image counts N_i containing useful information from the i th source element for a source strength N_{si} may be written as

$$N_i = \epsilon(1-\tau) \frac{r_{ap}\Delta s}{4l_1^2} N_{si} \quad (1)$$

where ϵ is the particle detection efficiency, and $r_{ap}\Delta s/4l_1^2$ is the fractional solid angle subtended by the penumbra.

Other source resolution elements also produce disk-shaped images. Where these images overlap the information bearing penumbra of the original image, they contribute to the background level and reduce the SNR. The total number of background counts N_b is the sum of the contributions from all other source resolution elements:

$$N_b = \epsilon \frac{r_{ap}\Delta s}{4l_1^2} \sum_{j=1}^{n_p} [\eta_{ij} + (1-\eta_{ij})\tau] N_{sj} \quad (2)$$

where n_p is the number of resolution elements in the source, η_{ij} is the relative overlap between the images of source elements i and j , and N_{sj} is the source strength of the j th source element. Setting $\eta_{ii} = 0$ accounts for the background counts within the penumbral ring caused by the original source element i .

Assume the detection follows Poisson statistics, then the variance in the information counts is given by:

$$\sigma_i^2 = \sigma_{si}^2 + \sigma_{bi}^2 = c_1^2 (N_i + N_{bi}) \quad (3)$$

where σ_{si} and σ_{bi} are the standard deviation in the signal and background counts, respectively; and c_1 accounts for the pulse-height characteristics²³⁵ of the detector. ($c_1 = 1$ for an ideal detector.) The SNR for source element i may be written as

$$SNR = \frac{N_i}{\sigma_i} = \frac{1}{c_1} \left[\frac{\epsilon(1-\tau)r_{ap}\Delta s N_{si}}{4l_1^2 \left\{ 1 + \frac{2}{(1-\tau)} \sum_{j=1}^{n_p} [\eta_{ij} + (1-\eta_{ij})\tau] \frac{N_{sj}}{N_{si}} \right\}} \right]^{1/2} \quad (4)$$

Some simplification of the SNR is possible. Exact numerical calculations have shown that $\eta_{ij} \leq 0.5$. As a worst-case estimate and a good approximation for targets much smaller than the aperture we can take $\eta_{ij} = 0.5$ for $j \neq i$. Now,

$$SNR = \frac{1}{c_1} \left[\epsilon \frac{(1-\tau)^2}{(1+\tau)} \frac{r_{ap}\Delta s}{4l_1^2} \frac{N_{si}^2}{N_s} \right]^{1/2} \quad (5)$$

where the total source strength N_s is given by

$$N_s = \sum_{j=1}^{n_p} N_{sj} \quad (6)$$

If the strengths of all source elements are identical, then $N_s = n_p N_{si}$ and we may write

$$\text{SNR} = \frac{1}{c_1} \left[\epsilon \frac{(1-\tau)^2}{(1+\tau)} \frac{r_{ap} \Delta s}{4l_1^2} \frac{N_s}{n_p^2} \right]^{1/2} \quad (7)$$

The SNR equations presented in this section may be applied to nonpenetrating and penetrating radiation. In practice, physical characteristics of the imaging system such as sharpness of the aperture PSF cutoff, magnification, and detector resolution set a minimum limit for Δs . For the thick apertures used in neutron imaging experiments, l_1 is the distance to a plane near the midplane of the aperture and r_{ap} is the radius in that plane.

Authors: R. A. Lerche and D. Ress

3.6.8. Ion Temperature Analysis of Implosions of DT-Filled Capsules

The demonstration of a high gain implosion through inertial confinement fusion will ultimately depend upon our ability to compress fuel to a high density. This in turn hinges upon our ability to keep the fuel as cool as possible during the initial phases of the implosion, since cool fuel compresses more easily than hot fuel. Hence, laser pulse shapes are now being

tailored²³⁶ to drive the implosion along the lowest possible adiabat. It is therefore desirable to be able to measure the fuel temperature, T_i , with a high degree of reliability. However, as we become more successful at imploding on a low adiabat, the task of determining T_i becomes more difficult. We have improved our measurement and analysis techniques²³⁷ and report here the precision to which T_i can be determined for DT-filled capsules with the neutron detectors currently in use at the Nova laser-fusion facility.

We begin with an analytic analysis of the relative error, σ_{T_i}/T_i , where σ_{T_i} is the standard deviation uncertainty in T_i . The technique currently used for deducing the fuel ion temperature is to measure the temporal broadening of the pulse of neutrons arriving at a detector a known distance from the target. The greater the broadening, the greater the fuel ion temperature at emission time,^{238,240} namely

$$T_i = a(\Delta t)^2, \quad (1)$$

where Δt is the FWHM of the thermally broadened neutron time-of-flight (TOF) signal, and a is a constant. Since

$$\frac{\sigma_{T_i}}{T_i} = \frac{2\sigma_{\Delta t}}{\Delta t} \quad (2)$$

then we need only derive a relation for $\sigma_{\Delta t}/\Delta t$. From propagation of errors, we have

$$\sigma_{\Delta t}^2 = \left(\frac{\partial \Delta t}{\partial \Delta t_m} \right)^2 \sigma_{\Delta t_m}^2 + \left(\frac{\partial \Delta t}{\partial \Delta t_{pr}} \right)^2 \sigma_{\Delta t_{pr}}^2 + \left(\frac{\partial \Delta t}{\partial \Delta t_b} \right)^2 \sigma_{\Delta t_b}^2 + \left(\frac{\partial \Delta t}{\partial N} \right)^2 \sigma_N^2. \quad (3)$$

Here Δt , Δt_m , Δt_{pr} , and Δt_b correspond to the FWHMs of the true (unfolded) TOF signal, the measured (raw) TOF signal, the detector impulse (prompt) response, and the neutron

emission burn width, respectively. In the last term, N refers to the number of detected neutrons. If distributions are assumed, then the uncertainty, $\sigma_{\Delta t}$, can be expressed as^{239,240}

$$\frac{\sigma_{\Delta t}^2}{(\Delta t)^2} = \frac{(\Delta t_m)^2 \sigma_{\Delta t_m}^2 + (\Delta t_{pr})^2 \sigma_{\Delta t_{pr}}^2 + (\Delta t_b)^2 \sigma_{\Delta t_b}^2}{(\Delta t)^4} + \frac{c_1^2}{2N} \quad (4)$$

where c_1 in the last term is a parameter that characterizes the pulse height distribution via

$$c_1 = \left[1 + \left(\frac{\sigma_q}{\langle q \rangle} \right)^2 \right]^{1/2}. \quad (5)$$

Here, $\langle q \rangle$ and σ_q are the mean and standard deviation of the detected charge distribution for 14-MeV neutrons. The first term of Eq. 4 represents a limit imposed by the finite temporal resolution of the detector; the second term represents a statistical limit. For small amounts of thermal broadening, the first term dominates, whereas for low numbers of detected neutrons, the statistical limit dominates.

We present in Fig. 3-89 the results of Eq. 4 as a function of detector distance and neutron yield for DT-plasma temperatures of 1 keV and 10 keV for a typical detector. Particularly for low yield shots ($Y_n^{DT} < 10^9$), we see that there exists a detector position that minimizes σ_T/T_i . Hence, it is beneficial to locate the detector somewhat closer to the target on low yield shots to increase N , thereby reducing the contribution from the second term in Eq. 4. We wish to test this hypothesis next with data taken on a recent series of implosions of DT-filled capsules at Nova.

The neutron TOF detectors used in the measurements consist of a 2.4-cm-thick by 4.6-cm-diam piece of Bicron BC418 fast ("quenched") plastic scintillator optically coupled to a microchannel plate photo multiplier tube. The signal from each detector is split and connected to two Tektronix 7912 transient digitizers set at different sensitivities. These units digitize and record the input voltage signal as a function of time. During the experiments, one detector was placed 18.33 m from the target, and another was 8.36 m, henceforth to be referred to as the 20 m and the 10 m detector, respectively. The neutrons emitted during an implosion arrive in a single burst, hence signals from individual neutron interactions in the detector "pile up," and the detector is run in "current mode." Before such a signal can be interpreted, the charge distribution and the impulse response from 14 MeV neutrons need to be calibrated.

The detector charge distributions were measured at our neutron generator laboratory, where deuterium is accelerated onto a tritium target, thus generating 14-MeV (DT) neutrons. A

typical charge distribution thus accumulated (in single-hit mode) with one of our detectors is shown in Fig. 3-90. The upper dotted histogram represents the raw charge distribution, the lower dotted histogram the background due to scattering, and the solid histogram the background subtracted distribution. After correcting for the effect of finite threshold, the mean charge deposited, $\langle q \rangle$, standard deviation, σ_q , and c_1 as defined by Eq. 5, were determined for each detector. We find for DT neutrons that $c_1 = 1.34$ and 1.40 for the 20-m and 10-m detectors, respectively. The resulting uncertainty introduced into T_i due to the $\pm 10\%$ uncertainty in c_1 is negligible.

The impulse response was measured on a low yield DT-neutron shot by placing each detector 50 cm from the target. At this close distance, thermal broadening is negligible, and all the neutrons in the pulse arrive simultaneously as an impulse. The resulting signal for the 20-m detector is shown in Fig. 3-91b, and has FWHM of 1.25 ns.

The first step in the data analysis is to deconvolve (unfold) the impulse response from the measured signal, as illustrated for the 20-m detector in Fig. 3-91. The solid curve in Fig. 3-91a represents the observed "raw" voltage signal, as recorded on a shot with DT-neutron yield of 7×10^9 , and has FWHM of 3.26 ns. We accomplish the unfold in an iterative fashion by guessing an initial "solution," folding it with the impulse response shown in Fig. 3-91b, then comparing with the measured signal. Based on the comparison, we improve our initial guess, and so on, until the solution, folded with the impulse response, is in good agreement with the observed signal. The unfolded solution thus arrived at is shown in Fig. 3-91c. The dotted curve in Fig. 3-91a represents the final comparison of solution folded with impulse response, based upon which we terminated the iteration.

The double peak structure in the unfolded TOF distribution results from the shoulder on the leading edge of the raw signal. Dividing the total charge of the signal, Q , by the average charge $\langle q \rangle$ per detected neutron (determined from the calibration run) indicates that $N \approx Q/\langle q \rangle \approx 170$ neutrons interacted in the detector. Therefore, the structure in the unfolded distribution is only a statistical fluctuation. This conclusion is further supported by: (1) Monte Carlo simulations of observed signal shape as a

Fig. 3-89. Temperature error versus detector distance from target for DT neutrons: (a) 1-keV plasma, (b) 10-keV plasma. Line pairs represent one standard deviation error limits in estimating the temperature for target yields of 10^8 (solid), 10^9 (dashed) and 10^{10} (dotted) DT neutrons. The curves are calculated for 1.5-cm thick, 2-cm diam plastic scintillator with a DT-neutron detection efficiency of 14% and $c_1 = 1.4$. In the calculations, temporal widths of $\Delta t_{pr} = 1064$ ps and $\Delta t_b = 100$ ps, and uncertainties of $\sigma_{\Delta t_{pr}} = \sigma_{\Delta t_m} = \sigma_{\Delta t_b} = 50$ ps are assumed.

Fig. 3-90. Detected charge distribution for 14-MeV neutrons for the 10-m neutron TOF detector, as measured in a separate calibration run at the neutron generator facility. The upper dotted histogram represents the raw distribution, the lower dotted histogram the background due to scattering, and the solid histogram the corrected (background subtracted) charge distribution.

Fig. 3-91. Deconvolution of the detector response from the observed voltage signal. (a) The solid curve on the top trace is the observed ("raw") neutron TOF signal, and the dotted curve is the result of convolving the impulse response with the unfolded "solution." The degree of agreement with the observed signal is a measure of the quality of the unfold. (b) The DT-neutron impulse response for this detector as measured on a separate shot with the detector located close to the target. (c) The unfolded ("true") neutron TOF signal.

function of shot yield and detector distance;²⁴⁰ and (2) the lack of such statistical fluctuation in the signal from the 10-m detector.

The unfolded voltage signal, $V(t)$, is next transformed to a differential neutron multiplicity spectrum. Since $dQ(t) = I(t)dt = [V(t)/50 \Omega]dt$, where 50Ω is the impedance of the circuit, $I(t)$ the current at time t , and $dQ(t)$ the differential charge in dt at t , then $dN(t) = dQ(t)/\langle q \rangle$ is the differential number of neutrons detected, or

$$\frac{dN(t)}{dt} = \frac{V(t)}{\langle q \rangle \cdot (50\Omega)}. \quad (6)$$

Defining $Y(t) \equiv \delta N(t) = [dN(t)/dt]\delta t$, where δt is the bin width, then based on Eq. 7 in Ref. 243, the error for each "data point" is $\delta Y(t) = c_1 \sqrt{Y(t)}$. The final TOF spectrum thus arrived at for the 20-m detector is shown in Fig. 3-92a, and that for the 10-m detector in Fig. 3-93a.

To determine T_i , a thermal TOF distribution, assumed to be, was fitted to the data by χ^2 minimization. The region fitted, shown by the two dotted vertical fiducial lines in Fig. 3-92a, is held fixed. The parameters varied are T_i , centroid (t_0), and overall amplitude (A_0), and the best fit is determined by mapping out the χ^2

vs T_i profile, as shown in Fig. 3-92b. The T_i corresponding to the overall minimum is equated with the best-fit T_i for the spectrum, and the (statistical) uncertainty is related to the curvature. Approximating the χ^2 profile as parabolic, then an increase of one above the minimum χ^2 corresponds to a change in T_i of one standard deviation.²⁴¹ This is illustrated by the dashed and the dotted lines in Fig. 3-92b, and gives $T_i = 0.95 \pm 0.14$ keV, where the error quoted is the statistical error.

For the same shot, a similar T_i analysis of the data from the 10-m detector is shown in Fig. 3-93, and results in $T_i = 0.90 \pm 0.05$ keV. The smaller relative statistical error for this detector results from the sharper curvature of the χ^2 space near its minimum, and reflects the higher number of neutron interactions (1500 in the 10-m detector vs 170 in the 20-m detector). We next estimate the total uncertainty in T_i for the 20-m and 10-m detectors, making use of Eqs. 2 and 4 and assuming distributions.

To decompose the uncertainty into the various components, we assume that $\Delta t_b = 0.10$ ns and $\sigma_{\Delta t_{pr}} = \sigma_{\Delta t_m} = \sigma_{\Delta t_b} = 50$ ps. For the 20-m detector, we have $\Delta t_m = 3.26$ ns, $\Delta t_{pr} = 1.25$ ns, $\Delta t = 2.13$ ns, $c_1 = 1.34$, and $N = 170$. Then combining Eqs. 2 and 4, we get

$$\begin{aligned} \frac{\sigma_{T_i}}{T_i} &= 2 \left[\frac{(\Delta t_m)^2 \sigma_{\Delta t_m}^2 + (\Delta t_{pr})^2 \sigma_{\Delta t_{pr}}^2 + (\Delta t_b)^2 \sigma_{\Delta t_b}^2}{(\Delta t)^4} + \frac{c_1^2}{2N} \right]^{1/2} \\ &= 2 \{ [(1.29 \times 10^{-3}) + (1.89 \times 10^{-4}) + (1.21 \times 10^{-6})] + (5.28 \times 10^{-3}) \}^{1/2} \\ &= 2 \{ [(1.48 \times 10^{-3})] + (5.28 \times 10^{-3}) \}^{1/2} \\ &= 0.16 \end{aligned} \quad (7)$$

for the relative error in T_i . We see that the dominant term in the uncertainty is the last term, namely, the statistical term. For the 10 m detector, we have $\Delta t_m = 1.94$ ns, $\Delta t_{pr} = 1.39$ ns, $\Delta t = 0.95$ ns, $c_1 = 1.40$, and $N = 1500$. Then

$$\begin{aligned} \frac{\sigma_{T_i}}{T_i} &= 2 \{ [(1.16 \times 10^{-2}) + (5.93 \times 10^{-3}) + (3.07 \times 10^{-5})] + (6.53 \times 10^{-4}) \}^{1/2} \\ &= 2 \{ [(1.76 \times 10^{-2})] + (6.53 \times 10^{-4}) \}^{1/2} \\ &= 0.27. \end{aligned}$$

PT-U-10520-2

Fig. 3-92a. Final neutron TOF spectrum for the 20-m detector corresponding to the unfold shown in Fig. 3-91c. The solid curve represents the best fit temperature of $T_i = 0.95$ keV, and the dashed and dot-dashed curves map out the range defined by the statistical error of $\delta T_i = \pm 0.14$ keV. The region included in the fits is marked by the two vertical dotted fiducial lines. (b) The χ^2 versus T_i profile from which the best fit T_i is determined. The statistical uncertainty, δT_i , is determined from the curvature in the neighborhood of the minimum, as illustrated by the dashed and dotted lines, namely, $\Delta\chi^2 = +1$ corresponds to $\Delta T_i = \pm \sigma_{T_i}$.

PT-U-10520-3

Fig. 3-93. Same as Fig. 3-92 except for the detector at 10 m from the target. The deduced temperature is $T_i = 0.90 \pm 0.05$ keV.

The dominant terms here are due to the uncertainties in the FWHMs of the raw TOF signal and the impulse response. Even though the statistical uncertainty is small, the total uncertainty in T_i is almost a factor of two larger for the 10-m detector than for the 20-m detector. The optimum detector location for this shot may be slightly less than 20 m, but 10 m is clearly too close.

We have also done a temperature analysis for a higher yield shot where $Y_n = 2.8 \times 10^{10}$. In this case, the MCP of the 10-m detector saturated, but the 20-m detector did not. The number of neutrons interacting in the 20-m detector was about 500, and the overall statistical quality of the signal was better than that shown in Fig. 3-92. The resulting temperature was $T_i = 0.95 \pm 0.09$ keV. Again using Eq. 7 to estimate the total error, where $\Delta t_m = 2.76$ ns and $\Delta t = 2.13$ ns, gives

$$\begin{aligned}\frac{\sigma_{T_i}}{T_i} &= 2\{ [(9.25 \times 10^{-4}) + (1.90 \times 10^{-4}) + (1.21 \times 10^{-6})] + (1.80 \times 10^{-3}) \}^{1/2} \\ &= 2\{ [(1.12 \times 10^{-3})] + (1.80 \times 10^{-3}) \}^{1/2} \\ &= 0.11.\end{aligned}$$

Since the contributions to the uncertainty from the response of the detector and from the statistical term are nearly the same, we suspect that the optimal detector position for this shot was near to the 20-m detector.

Other sources of error are now considered. (1) If scattering in the chamber wall and along the flight path leads to a slightly broadened signal, this broadening would be misinterpreted as thermal in origin. Based on shots where both detectors gave useful data, we estimate that this effect on T_i is 5% or less. (2) The finite width of the scintillator, ΔL , causes a slight broadening of the signal. Since the impulse response was measured at a location different from where the actual data were recorded, this broadening due to ΔL might not completely cancel when the impulse response is unfolded from the actual data. The uncertainty thus introduced into T_i should be small for DT neutrons, though less than 5%. (This effect could be accounted for during the analysis.) (3) The finite duration of neutron emission tends to broaden the measured TOF signal. Based on calculations of the implosions with LASNEX,²⁴² this broadening is predicted to be negligible. (4) A more serious concern arises, though, if emission occurs from regions of fuel undergoing substantial motion. The resulting kinematic distortion of the neutron TOF spectrum would then not be handled properly in our analysis, and our deduced T_i would be systematically too high. This point is still an open question.

In summary we have examined to what extent fuel T_i can be determined from implosions of DT-filled capsules using our current detectors. On low yield shots, the limiting factor in trying to minimize $\delta T_i/T_i$ typically is the number of neutrons interacting in the detector. Moving the detector somewhat closer to the target can be beneficial, provided that the impulse response is understood in detail. On higher yield shots, $\delta T_i/T_i$ is less sensitive to detector distance, and a location somewhat farther from the target is preferable. We conclude that we can measure 1 keV temperatures to about $\pm 20\%$ or better on shots with DT-neutron yield $Y_n > 5 \times 10^9$ with a detector at 18.3 m. Operating a second detector somewhat closer to the target is useful for improving the statistical quality of the data on low yield shots. For the yields measured in this series of shots, though, the detector located at 8.4 m was too close to allow a high quality T_i determination.

Authors: B.A. Remington, R.A. Lerche, M.D. Cable

3.6.9. A Technique for Shell Compression Measurements of Laser Fusion Targets by Neutron Activation of a Rubidium Tracer

Introduction. The measurement of the neutron averaged areal density $\langle \rho \Delta R \rangle$ of the compressed pusher of an ICF target by neutron activation has become a well established technique.²⁴³⁻²⁴⁷ Its implementation involves the activation of one or more elemental constituents of the pusher by 14

MeV neutrons produced in an implosion. Pusher debris containing these activated atoms is then collected on metal foils placed near the target. These foils are transferred to a radiation detector that measures the individual activity levels. The $\langle \rho \Delta R \rangle$ is determined from the expression

$$\langle \rho \Delta R \rangle = \frac{CA}{\sigma Y N_o f \eta_c \eta_D e^{-\lambda t_d} (1 - e^{-\lambda t_c})} \quad (1)$$

where C is the number of detected counts, A is the average atomic weight of a pusher atom, σ is the activation cross section, Y is the neutron yield, N_o is Avogadro's number, f is the fraction of atoms of the isotope to be activated compared to the total number of pusher atoms, η_c is the fraction of activated atoms collected, η_D is the detection efficiency, λ is the decay constant, t_d is the time to load the detector, and t_c is the counting time. The angle brackets indicate the measured value is an average weighted by the neutron production rate as well as an average over all neutron paths through the shell.

Previous $\langle \rho \Delta R \rangle$ activation measurements have mostly used the $^{28}\text{Si}(n,p)^{28}\text{Al}$ reaction on microshell targets made of either glass²⁴³⁻²⁴⁶ or silicon-doped plastic.²⁴⁷ The principal deficiency of this technique is that the fraction of Al collected (η_c) cannot be measured directly. One must therefore either measure the Al collection fraction by comparison with some other method of determining $\langle \rho \Delta R \rangle$ and then assume η_c is constant from shot-to-shot, or, as is more commonly done, the Al collection fraction is inferred by assuming that it is equal to the measure collection fraction of ^{24}Na , a common radioactive tracer in irradiated glass shells. Although these assumptions are reasonable, at Nova, under certain conditions, the amount of Na collected varies from shot-to-shot by as much as a factor of 6 and, furthermore, as shown below, the Al and Na collection fractions can differ from each other by as much as a factor of 2.

An extension of this activation method that does not suffer from these collection problems uses natural rubidium as the tracer. Both natural isotopes of Rb undergo (n,2n) reactions with 14.1 MeV fusion neutrons resulting in the activation products $^{84}, ^{84m}, ^{86}, ^{86m}\text{Rb}$. The isomers, ^{84m}Rb ($t_{1/2} = 20.5$ min) and ^{86m}Rb ($t_{1/2} = 1.02$ min), have half-lives short enough to allow reasonable count rates for making the shell $\langle \rho \Delta R \rangle$

measurement. By irradiating the target with reactor produced thermal neutrons prior to the implosion experiment, the radioactive tracer ^{86}Rb ($t_{1/2} = 18.8$ d) is formed throughout the shell. Comparing the ^{86}Rb activity in the target before the implosion to the activity on the collector foils after the implosion gives the fraction of rubidium that was recovered. Thus the same element is used for both the $\langle \rho \Delta R \rangle$ measurement as well as the collection fraction determination. As a further benefit, the ^{86}Rb shell activity can be compared to that of a RbCl flux monitor to accurately assay the Rb content in the shell.

Decay Scheme and Cross Section for ^{84m}Rb .

The most probable reactions (excluding elastic and inelastic scattering) of 14-MeV neutrons with Rb are (n,2n) reactions. Relevant properties of the (n,2n) reaction products from Refs. 252-254²⁴⁸⁻²⁵⁰ are listed in Table 3-8. The decay scheme for ^{84m}Rb , the reaction product we find most convenient to use, shows the three gamma rays emitted by the decay of the isomer (Fig. 3-94). Because internal conversion competes with gamma emission there are 6 separate decay modes as listed with their relative probabilities in Table 3-9.

There is a large variation among the measurements of the 14.1 MeV $^{85}\text{Rb}(n,2n)^{84m}\text{Rb}$ cross section.²⁵¹ We remeasured it using our neutron generator facility that produces a well characterized source of 14 MeV neutrons by directing a 200 keV deuteron beam onto a tritiated Ti-on-Cu target. The number of DT neutrons produced was monitored using an absolutely calibrated silicon charged-particle detector that counts the DT alpha particles emitted from the accelerator target. RbCl samples (0.5 g) were irradiated with a flux of $10^5 \text{ n/cm}^2\text{-s}$ for 20 minutes at positions 4-8 cm from the neutron source and at angles of 0, 30, 45, 75, and 90 degrees to the incident beam direction. Small corrections were made for the angular spread of neutrons across the sample, neutron scattering in the tritium target and in the RbCl sample, and attenuation of gamma rays leaving the sample. The ^{84m}Rb produced in the sample was then counted in the absolutely calibrated gamma detector described in the next section. The cross section obtained at 14.1 MeV was 0.514 ± 0.080 b which is close to the average of previous measurements compiled in Ref. 251.

Gamma-Ray Detection. A pair of coaxial high-purity intrinsic-germanium well detectors

Table 3-8. 14 MeV Neutron Activation Characteristics of Rubidium

	^{85}Rb		^{87}Rb	
Natural Abundance (%)	72.17		27.83	
Reaction	(n,2n)		(n,2n)	
Reaction Product	^{84}Rb	$^{84\text{m}}\text{Rb}$	^{86}Rb	$^{86\text{m}}\text{Rb}$
Threshold (MeV)	10.49	10.95	10.48	11.03
Cross-Section (mb)	1130	514	620	520
Half-life	32.77 d	20.5 min	18.82 d	1.02 min
Gamma-1 Energy (keV)	882	216	1077	556
Gamma-1 Intensity (%)	74	34	8.79	98.2
Gamma -2 Energy (keV)		248		
Gamma-2 Intensity (%)		63		
Gamma -3 Energy (keV)		465		
Gamma -3 Intensity (%)		31

Fig. 3-94. The decay scheme for $\text{Rb-}^{84\text{m}}$. The probability per decay for emission of each of the three gamma-rays is shown. Electron emission by the internal conversion process competes with the gamma transitions.

Table 3-9. Decay modes of $^{84\text{m}}\text{Rb}$ derived from γ intensities and IC coefficients in Ref. 8.

Decay Mode	465→Grnd	465→248	248→Grnd	Rel. Probability
1	γ			0.310
2	IC			0.050
3		γ	γ	0.332
4		γ	IC	0.008
5		IC	γ	0.293
6		IC	IC	0.007

γ = gamma emission; IC=internal conversion

are used to count the gamma rays emitted by the collected rubidium. These detectors have an active volume of 145 cm^3 each and are operated in a face-to-face geometry (Fig. 3-95) with the collector foils wound onto a Lucite spindle and placed entirely inside the 1.5-cm-diam, 4-cm-deep wells, thus providing nearly 4π detection solid angle. Resolution in the 100 to 500 keV energy range is approximately 2.2 keV FWHM. The detectors are operated independently, but with the gains matched so that the 4096 channel pulse height spectra from the two detectors can be summed for analysis after the experiment. A typical summed gamma spectrum is shown in Fig. 3-96.

When the 248-keV and 216-keV gamma rays are emitted in cascade, they cannot be resolved temporally and therefore produce summing effects which prevent the absolute detection efficiency from being determined in the usual way which is to use calibrated sources that produce photopeaks in the energy range of interest. Instead, we developed a 3-step calibration procedure. First a ^{152}Eu NBS standard source (1 mCi) was used to calibrate each detector at a position 20 cm away from their front faces in order to remove the effects of summing. The next step was to produce $^{84\text{m}}\text{Rb}$ activity (typically 5–10 nCi) in a 0.5 g cylindrical plug of compressed RbCl powder using our 14-MeV neutron generator facility. This sample was then counted in the 20-cm position to determine its absolute activity (by comparison to the Eu-152 source). The final step was to count the sample at various positions in the detectors. The efficiency for counting $^{84\text{m}}\text{Rb}$ at these positions for the three photopeaks was then determined from the peak area divided by the source strength as shown in Fig. 3-97. Accurate peak areas were obtained for all spectra using the peak fitting code MICROSAMPO^{252, 253} run on an IBM-AT personal computer.

Rubidium Assay. To accurately estimate $\langle p\Delta R \rangle$, the initial amount and distribution of Rb in the shell must be known. The distribution of Rb in glass shells is assumed to be uniform. This assumption is consistent with elemental analyses using characteristic x-rays produced by a tightly focused electron micrograph beam to interrogate different locations on the exposed wall of a hemispherical shell.

The absolute amount of Rb is determined by comparing the ^{86}Rb activity in a shell to the ^{86}Rb activity in a RbCl flux monitor that was

irradiated in close proximity to the shell. This is accomplished in the following way: 4 target shells are irradiated in a reactor-grade graphite holder for approximately 4 hours in a thermal neutron flux of $9 \times 10^{13} \text{ n/cm}^2\text{-s}$ at the Los Alamos Omega West Reactor. Each holder also contains four $1.5 \times 1.5 \text{ mm}$ hollow cylindrical flux monitors made of Epon 826 epoxy and Versamid 140 catalyst. (Irradiated empty epoxy cylinders produced insignificant amounts of radioactivity and showed no evidence of contamination from trace amounts of Rb in the epoxy.) Before sealing, each cylinder is filled with 1/2 microliter of a diluted (50:1) aqueous RbCl solution which is evaporated to dryness. This technique produces nearly identical ($\pm 2\%$) flux monitors containing an accurately known quantity of rubidium (7.1 μg) that is comparable to the amount contained in the target and consequently produces a level of activity nearly the same as the target thus allowing the two to be counted in the same physical location in the detector. We find, for instance, that in microshells used for direct-drive implosions,²⁵⁴ 1% of the atoms in the shell are Rb atoms.

Shell $\langle p\Delta R \rangle$ Measurements. The debris collector is a 25- μm -thick titanium foil formed into a cone (4 cm diam $\times$ 8 cm length) and is placed 5 cm from the target where typically 2 to 6% of the rubidium is collected. The foil is recovered from the target chamber, transferred to a counting lab, and loaded into the detector assembly. Transfer and loading takes 1–2 minutes. The counting data are collected from each detector in 10 consecutive 1-minute 4096 channel spectra followed by ten 10-minute spectra to accommodate the half-lives of the various activities contained in the target debris. [Typically the following activities can be observed following a direct-drive glass implosion: $^{84\text{m}}\text{Rb}$ (20.5 min), $^{86\text{m}}\text{Rb}$ (1.02 min), $^{86\text{m}}\text{Rb}$ (18.8 d), ^{28}Al (2.24 min), ^{29}Al (6.6 min), ^{24}Na (15.02 h), and ^{42}K (12.36 h).] The collector foil is then divided into sections (usually 5) which are counted individually using ^{86}Rb to determine the fraction of Rb collected η_D as well as the distribution of debris on the foil. This distribution is then used along with the position-dependent detection efficiency (Fig. 3-97) to calculate an average detection efficiency η_C . (Notice that the detection efficiency for the 216 keV and 248 keV gamma rays does not strongly depend on the detailed debris distribution.) Finally, the total $^{84\text{m}}\text{Rb}$ counts C in a 40-minute

Fig. 3-95. Schematic diagram of the detector geometry. Two coaxial germanium well detectors are positioned face-to-face thus forming a 4π detection solid-angle to the collector foil placed in the wells.

Fig. 3-96. High-resolution gamma spectrum from a glass sample containing 1 atom-% Rb that had first been irradiated by thermal neutrons and then by a short burst of 14-MeV neutrons. The activities produced by fast neutron interactions are Rb-84m, Al-28, Al-29 while Rb-86 was formed by thermal neutron capture.

Fig. 3-97. The measured probability for detection per decay of each of the 3 Rb-84m gamma rays as a function of source location (0 cm is the location midway between the detectors, 4 and -4 cm are at the bottom of the wells). Notice that the efficiency is relatively insensitive to location for the 216-keV and 248-keV gamma rays.

period starting 10 minutes after the shot (to allow time for background from ^{28}Al to decay) is used to compute the shell $\langle p\Delta R \rangle$ value using Eq. 1. The natural background (~ 1 cpm) is subtracted from the total counts as is a comparable contribution due to Compton events from ^{86}Rb and ^{24}Na gamma rays. The $\gamma \langle p\Delta R \rangle \eta_c$ product required to produce a statistically significant number of counts (about 30) is approximately 5×10^8 .

Authors: Stephen M. Lane and
Michael B. Nelson

3.6.10. The Generation of Ultra-Short Electrical Pulses

Introduction. The generation of picosecond high voltage pulses is essential for driving gated X-ray cameras and highly parallel transient digitizers. Until recently, 1-ns risetime 2.5-kV UHF vacuum planar triode pulsers have been used to drive pulse sharpening avalanche diodes²⁵⁵ generating sub-100 ps FWHM 1–2-kV pulses. However, the bulk, jitter, and expense of vacuum triode drivers has limited the application of the avalanche pulse sharpening technique.

Kilovolt-level sub-100 ps pulses can also be generated using the Auston²⁵⁶ photoconductive switch (PCS), but a complete PCS system is not competitive with the avalanche diode system in terms of overall cost or size. We have therefore concentrated on development of avalanche diode pulsers with the primary objective being to phase out vacuum triode drivers in favor of power MOSFET drivers. In addition, we have studied the avalanche shock wave process in some detail.

Progress—Power MOSFET Drivers.

Conventionally packaged MOSFET switching time is dominated by source lead inductance L_s ~ 8 nH and source dynamic resistance $r_s \sim 2 \Omega$, with an associated time constant $\tau_s = L_s/r_s \approx 40$ ns. Under near-linear conditions, MOSFET switching time is approximately $2.2\tau_s \approx 88$ ns, while under high overdrive levels switching time can be as low as 10 ns.

Microwave packaging techniques can lower L_s to < 3 nH, but this approach is expensive and not amenable to future hybrid circuit options. Instead, we have taken a circuit design approach to achieve faster switching with conventional packaging.

The voltage drop across the source lead during switching directly reduces the effective gate drive. This loss can be compensated through gate drive preshaping. We used the inductive current mirror circuit shown in Fig. 3-98 as a simple and symmetric means to generate the preshaped drive. Compensation occurs when $I_g/I_o \approx L_s/L_g \approx R_s/R_g$.

The SPICE plot in Fig. 3-99 shows the resulting 120V shaped gate drive as viewed from the gate terminal to ground. The internal gate-source voltage can be seen to be 10–20V, well above the 4V MOSFET threshold and yet within the specified 40V maximum rating. Clamp diode D1 prevents ringing and reverse overvoltage.

The inductive current mirror concept was extended to the driver stages preceding the final output stage. The complete circuit uses six transistors, operates from a logic level trigger, and delivers ~ 40 -kW, 2-ns risetime pulses while drawing one watt DC at a 1-kHz repetition rate.

Nonlinear Transmission Line.

Approximately 1-ns transients are needed to initiate shock waves in high voltage avalanche diodes, about $2\times$ faster than packaged MOSFET capabilities. To bridge this gap, a high-voltage diode-loaded nonlinear transmission line (NLTL)^{257–259} was developed for risetime compression. The circuit diagram is shown in Fig. 3-100.

The NLTL exhibits a voltage dependent propagation velocity, $V_p = \sqrt{LC(V)}$ where $C(V) = C(0)/(1-V/V_f)^m$ is the diode capacitance, with V_f being a materials-related constant and m being a function of dopant profile within the diode. Since $C(V)$ decreases with increasing voltage, the higher voltage part of an applied edge propagates faster than the lower part, resulting in edge steepening as it propagates down the line, as indicated in

Fig. 3-100. The risetime is limited, or is in a state of shock, when the dispersion of the media equals the compression arising from the nonlinearity.

Compression is limited by the cutoff frequency of the line, $\omega_c \approx 2/\sqrt{LC}$ and by diode cutoff frequency $\omega_{rc} = 1/(r_d(0)C(0))$ where r_d is the diode series resistance. In practice, the primary limitation is ω_{rc} , and for $\omega_{rc} < 5\omega_c$ we have from²⁵⁹

$$\frac{I_o}{I_g} \approx \frac{R_g}{R_s} \approx \frac{L_g}{L_s}$$

Fig. 3-98. MOSFET inductive current mirror for compensation of parasitic source lead inductance L_s .

Fig. 3-99. SPICE plot of inductive current mirror V_g , V_{gs} , and V_d .

$$\text{Compression} = T_L - T_h = N \left[(LC (V_L))^{\frac{1}{2}} - (LC (V_h))^{\frac{1}{2}} \right]$$

Fig. 3-100. Non-linear transmission line circuit diagram.

$$T_{\min} = \frac{8.8}{\omega_{rc}} \frac{1}{\sqrt{1 - V_{\max}/V_f - 1}} \approx \frac{8.8}{\omega_{rc}} \left(\frac{V_f}{V_{\max}} \right)^{1/2} \quad (1)$$

Since our NLTL is the first known diode-loaded NLTL to operate at the kilovolt level, it is useful to calculate $T_{\min}$ as a function of diode breakdown voltage, V_b . If we let $V_{\max} = V_b$ and use the relations²⁶⁰

$$r(0) = \frac{W_b}{Aq\mu N}, \quad C(0) = \left(\frac{q\epsilon N}{2V_f} \right)^{1/2} A$$

$$\text{and } V_b = 60 \left(\frac{E_g}{1.1} \right)^{3/2} \left(\frac{N}{10^{22}} \right)^{-3/4}.$$

We get, with Eq. (1),

$$T_{\min} = \left[\frac{3.75 \times 10^{-24}}{\left(\frac{E_g}{1.1} \right)^2} \right] \left(\frac{\epsilon V_b}{q\mu} \right)^{4/3}$$

where A = Area, ϵ = dielectric constant, $q = 1.6 \times 10^{-19}$, μ = mobility, N = impurity concentration in the bulk region, E_g = band gap, and $W_b = (2\epsilon V_b/qN)^{1/2}$ is the depletion width at breakdown. $T_{\min}$ is plotted in Fig. 3-101 for silicon and GaAs diodes.

In a uniformly periodic NLTL, shock formation does not occur until the edge has propagated some distance, so the LC sections near the input have a considerably higher ω_{lc} than needed. By reducing ω_{lc} of the input section to match the input risetime rather than the output risetime and then tapering successive LC sections upwards in ω_{lc} , we can keep the wave front in continuous shock. An experimentally optimized taper rate resulted in ~20% risetime compression per section and reduced the total number of sections by ~3x.

Using a 5-section NLTL, tapered in both ω_{lc} and impedance, we compressed a 1-kV, 2-ns MOSFET input to a 4-kV, 1-ns output as shown in the SPICE simulation in Fig. 3-102. Avalanche breakdown, however, is not modeled in Fig. 3-102. In actuality, the diode at the end of the unterminated NLTL is subjected to a double amplitude reflection, driving it into shock

avalanche breakdown at ~2.5 kV, as desired, and before 4 kV is reached.

Avalanche Diodes. Conventional impact avalanche breakdown spreads at a relatively slow rate limited by carrier saturation velocity, $V_{sat} \sim 10^7$ cm/s. In our diodes, an avalanche shock wave propagating at $V_x = 10 - 100 V_{sat}$ is initiated by a large and rapidly rising electric field which drives the diode to several times its static breakdown field strength, E_b .

As the avalanche shock wave transits the bulk region of the diode, it leaves a highly conducting plasma in its wake. During transit, a large displacement current flows causing the diode to switch from a high to low impedance state. Experimentally, we have observed 30 ps, 2.3 kV transitions in a 500 VDC rated diode. The diode terminal voltage resulting from this action is seen in the SPICE simulated breakdown waveform in Fig. 3-103 along with an edge-differentiated 100 kW output pulse.

Although this "single-shot" process was first discovered²⁶¹ in 1979, we have recently related it to microwave TRAPATT²⁶²⁻²⁶⁴ (Trapped Plasma Avalanche Transit Time) diode oscillators dating to 1967, and have thereby opened a large body of relevant literature to our application. We will now discuss TRAPATT theory and then present our theory on switching time.

The avalanche wave velocity can be solved for from the geometry shown in Fig. 3-104. The electric field gradient in a uniformly doped semiconductor is given by Poisson's equation

$$dE/dX = qN/\epsilon$$

where N = impurity concentration in the bulk region. With an applied voltage transient we get

$$dE/dT = J_{drv}/\epsilon$$

In summary we have where J_{drv} is the drive current density and ϵ is the dielectric constant. As can be seen from

Fig. 3-104, at breakdown we have

$$E(X) = E_b = E_b - (qN/\epsilon)X + (J_{drv}/\epsilon)T$$

from which we get wave velocity

$$V_x = X/T = J_{drv}/qN.$$

Fig. 3-101. Diode limited NLTL risetime versus breakdown voltage for abrupt junction diodes.

Fig. 3-102. 3:1 pulse compression on a five section tapered NLTL.

Fig. 3-103. SPICE simulated avalanche breakdown and 100 kW output pulse.

Fig. 3-104. Electric field gradient in a uniformly doped semiconductor.

The threshold for avalanche shock formation occurs at

$$V_x = V_{\text{sat}} = J_{\text{drv}}/qN. \quad (2)$$

J_{drv} can be related to drive voltage slew-rate, $\dot{V}_{\text{drv}}$, by

$$J_{\text{drv}} = \frac{I_{\text{drv}}}{A} = \frac{C V_{\text{drv}}}{A} = \frac{\epsilon}{W_b} \dot{V}_{\text{drv}}$$

and combining with Eq. (2),

$$\dot{V}_{\text{drv}}|_{\text{Thresh}} = \left\{ \frac{2q}{\epsilon} \left[\frac{60^{4/3} \left(\frac{E_g}{1.1} \right)^2 10^{22}}{V_b^{4/3}} \right] V_b \right\}^{1/2} \quad V_{\text{sat}} = 2.62 \left(\frac{1.3 \text{ kV}}{V_b} \right)^{1/6} \frac{\text{kV}}{\text{ns}} \text{ for silicon.}$$

The avalanche wave transit time is given by

$$T_{\text{sw}} = W_b/V_x = W_b q N / J$$

where $J = J(T)$ is the current density during transit. One key issue is the maximum theoretical switching speed. Previous theories^{263,265} have assumed that the driving force during switching is J_{drv} or $J_{\text{drv}} + J_{\text{load}}$, the current density due to the drive and load, but this results in a typical wave velocity of only $\sim 2V_{\text{sat}}$. In our proposed theory, the total current density during switching is

$$J_{\text{sw}} = J_{\text{drv}} + J_{\text{load}} + J_c$$

where J_c is the current resulting from the collapse of the depletion capacitance as the avalanche wave sweeps across the depleted base region, analogous to parallel plates closing together.

For our typical diodes, $W_b = 100 \mu$, $N = 2 \times 10^{14}/\text{cm}^3$, $J_{\text{drv}} = 326 \text{ A}/\text{cm}^2$, $J_{\text{load}} = 1300 \text{ A}/\text{cm}^2$ and by inference from an observed 30 ps switching time, we get an average $J_c = 10,650 \text{ A}/\text{cm}^2$. Clearly, once the avalanche wave is launched by J_{drv} the powerful regenerative mechanism of $J_c + J_{\text{load}}$ dominates and J_{drv} has only a secondary contribution. The threshold for regenerative action occurs when $(J_c + J_{\text{load}}) \geq J_{\text{drv}}$.

$$J_{\text{drv}} = qN V_{\text{sat}} = \frac{\epsilon}{W_b} \dot{V}_{\text{drv}}$$

or

$$\dot{V}_{\text{drv}}|_{\text{Thresh}} = \frac{qN}{\epsilon} V_{\text{sat}} W_b. \quad (3)$$

Substituting expressions for N and W_b from the previous section into Eq. (3), we arrive at

Future Directions. The MOSFET/NLTL

avalanche diode pulser resulting from this work is sufficiently compact and inexpensive that new gated X-ray camera structures and ultra-fast transient digitizer topologies can be considered. We are currently investigating impulse radar as a potential spin-off technology.

There is a broad range of possible enhancements to our pulser system. SPICE simulations show 600 ps MOSFET switching is possible using low inductance hybrid circuit packaging. This speed would eliminate the need for the NLTL, so an entire 100-kW pulse generator could easily be placed within a single hybrid circuit package.

NLTL speeds could be increased $>10\times$ with the development of high-voltage hyperabrupt junction GaAs diodes resulting in improved energy-conserving pulse-width compression on the tapered NLTL.

Optimization of avalanche diode doping profiles, the design of high punch through structures, and the use of other materials such as GaAs, SiC, and diamond could lead to higher switching speeds and reduced drive requirements. In addition, area scaling and Marx bank structures remain as potential avenues to higher power.

Author: Thomas E. McEwan

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4 Nova Laser Science and Technology

Section editors: J. R. Murray, H. T. Powell

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4. Nova Laser Science and Technology

4.1. Introduction

During 1989, the Nova Operations and Laser Science groups were involved in many activities in direct support of the Nova experiments program and basic laser performance improvements.

One of the major accomplishments was completing the installation of platinum-free glass in the 46-cm and 31-cm amplifiers along with the installation of fused silica lenses in the input and output of the 46-cm and 74-cm spatial filters. Upon completion of these modifications, the Nova laser was reactivated, characterized, and brought to full operational potential.

During system activation, a single beamline (BL 7) was chosen as a prototype for evaluation purposes. This prototype beamline allowed us to evaluate system damage issues along with crystal conversion efficiency based on crystal configurations, Type I/Type II versus Type II/Type II, and the effects of beam depolarization on conversion efficiency. During our prototype experiments, we also observed transverse stimulated Brillouin scattering (SBS) at 3ω in our 74-cm lenses and beam splitters which established 3ω operation limits for the Nova facility at pulsewidths of 2.0 ns and longer. The SBS problem disappears when $\approx 0.5 \text{ cm}^{-1}$ bandwidth is added.

Additional improvements to system performance were accomplished by carefully evaluating beam uniformity and initiating changes to beam apodizers, amplifier configurations, beam fill-factors and pinhole sizes.

Other areas of major effort were beam pointing and diagnostics. A program to improve and carefully characterize and maintain our streak cameras was initiated. We are also continuing to develop a reliable diagnostic to evaluate 3ω power delivered to target chamber center. Several experiments were conducted to quantify beam-pointing accuracy and determine areas for improvements.

Several engineering projects offering improvements to Nova operations and

performance were undertaken. A chamber lens extraction tool was developed to allow easy removal and replacement of 74-cm target chamber lenses. The control system is undergoing a complete upgrade to improve performance, reliability, and ease of modification. The flashlamps used in the 46-cm amplifiers were carefully evaluated to determine the cause of lamp spotting and seal failure. An effort in oscillator development was begun to provide us with short pulse (10–20 ps) and broadband ($\sim 12 \text{ \AA}$) capability.

Nova Operations continues to operate a 12-hr day from 7:00 a.m. to 7:00 p.m. We utilize two shifts with a 6-hr overlap. This amount of overlap allows significant time for operations personnel to support additional projects and eliminates system down time due to lunch and dinner breaks. We have set aside Monday as maintenance day and utilize Tuesday through Friday for system operation.

A summary of Nova experiments is presented in **Table 4-1**. This has been our most productive year to date. Note the distinction between the number of system shots and experiments. Nova can perform numerous experiments during a single system shot. For example, we may fire a target shot to the 10-beam chamber with 7 beamlines, propagate 2 beams to the 2-beam chamber, and perform a laser science experiment on the remaining beam. We would count this as one system shot and three experiments. Beamlines 6 and 7 have significantly more shots since they are used for the 2-beam chamber operation. Beamline 10 had a higher than average shot total due to an additional experimental station on this arm. This experimental station is located in the switchyard at the output of arm 10 and is often utilized for laser science experiments.

Detailed discussions on all these experiments and upgrades are provided in the following sections.

Authors: Glenn Hermes, Tim Weiland

Major Contributor: Demos Kyrazis

Table 4-1. Nova shots for calendar year 1988.

		System Shots	Ride-Along Shots	Total Target	Classified Target	Total Laser Science	Calibration 1ω	3ω	Total Miscellaneous Inc. Cals. Burns, Etc.
Jan.	88	28	0	25	4	0	0	3	3
Feb.	88	23	11	23	11	1	5	1	10
Mar.	88	59	18	21	8	5	6	2	11
Apr.	88	44	8	19	11	33	0	0	0
May	88	47	3	30	24	14	4	2	6
Jun	88	55	4	38	31	16	2	3	5
Jul	88	50	10	16	2	38	5	1	6
Aug	88	54	17	39	23	18	8	0	14
Sept	88	60	7	51	45	4	1	3	12
Oct	88	64	0	61	50	0	1	2	3
Nov	88	64	19	28	25	52	1	2	3
Dec	88	0	0	0	0	0	0	0	0
Total		548	97	351	234	181	33	19	73

- (1) System + Ride-along = Target + Laser + Miscellaneous.
- (2) During May and June we began using quick turnaround target alignment, beams were not realigned to similar targets.
- (3) Late August and September we began using quick turnaround laser and target alignment, no beam or pinhold realignment between shots, and no target alignment between similar shots.

4.2. Nova Performance

4.2.1. Introduction

Our goals for Nova are to provide 40–50 kJ of 0.35- μm energy to fusion targets in a precise temporally shaped pulse with good power balance and high alignment accuracy for all ten beams. This is part of an overall activity which we now call the Precision Nova project. The activities of the Nova laser science group during 1989 were directed toward these goals. In March of 1989, we completed the Nova Phase II upgrade. This included replacing the laser disks in the 31.5-cm and 46-cm amplifiers with disks free of platinum-particle inclusions, replacing the large spatial filter lenses with lenses made of platinum-free fused silica, and changing the laser staging to improve beam uniformity. After completing the upgrade, we demonstrated the increased energy possible from Nova at its fundamental wavelength of 1.05 μm and the increased energy available at the target at the third-harmonic wavelength of 0.35 μm . In separate series of shots, we demonstrated laser output energy at 1.05 μm of 125 kJ that exceeded Nova's original performance goals and we delivered 45 kJ of third-harmonic energy into calorimeters inside the ten-beam target chamber.

We continued to refine Nova's temporal pulseshaping capability demonstrated on the Phase II prototype beamline,¹ and during the year began target experiments with two-step continuously shaped pulses. We also began to improve the power balance among the beams at the target at the third-harmonic wavelength by working to reduce beam-to-beam differences in frequency conversion efficiency. An important part of this effort was to reduce the depolarization of the 1.05- μm beams incident on the frequency conversion arrays. On a demonstration beamline, we reduced the depolarized fraction of the energy from approximately 4% measured in 1986² to less than 1% measured in November, 1989.

We identified and studied other critical issues that might affect our ability to routinely deliver high third-harmonic energies to targets. These included: (1) demonstrating a practical method to suppress transverse stimulated Brillouin scattering in the fused silica focusing lenses; (2) studying the optical damage limits of the most damage-prone optics in Nova and

concluding our energy goals should be possible without significant damage to the laser and target chamber optics; and (3) finishing the data analysis of the prototype beam experiments which demonstrated methods to significantly improve Nova's third-harmonic conversion efficiency.

During the year, we formally organized a Precision Nova project team to make improvements to the laser necessary for future implosion experiments that comprise the hydrodynamically equivalent target physics (HEP) campaign. These experiments will require the laser to deliver to targets:

- Routine energy of 40 kJ at 3ω .
- Power imbalance of <5% rms throughout a shaped pulse.
- Flexible pulseshaping with both continuous and picket fence pulses.
- Beam pointing accuracy to less than 10 μrad rms.
- Beam simultaneity better than 15 ps.

We estimate that completion of this project will take two to three years.

Author: D. R. Speck

4.2.2. Prototype Beamline Results

To investigate the constraints on the energy and power available from the Nova laser and frequency conversion system after the completion of the Nova Phase II project, we conducted tests on a prototype beam. We chose beamline 7 as the prototype, the first to be retrofitted with platinum-particle-free laser glass and fused silica spatial filter lenses. This beam can be directed into the two-beam target bay, where space was available to construct a diagnostic station to study harmonic conversion at the high power available from a rebuilt beamline.

Output at 1.05- μm Wavelength. The goals of the prototype beam tests at 1.05 μm were to show that a Phase II beam can meet Nova's original performance goals, and to determine the energy limits at which Nova can safely operate. During 1988 and early 1989, the prototype arm fired more than 450 shots for laser and target experiments, many of which exceeded the original performance goals at 1.05- μm wavelength for a single beamline. High energy shots included 32 shots at more than 8 kJ in a 1-ns duration pulse and 11 shots at more than 11 kJ in a 2.5-ns pulse. During the series we minimized

the threat of optical damage by maintaining a spatially uniform beam (see section 4.2.4.), which we continuously monitored with photographs taken in the output sensor. The only damage observed in the laser chain was a minor damage spot on the input lens to the final spatial filter.

Prior to the high energy shots, we measured the passive inline transmissions of the optics and small signal gain of each of the amplifier sections in the beamline,³ in order to obtain the component parameters necessary for our laser propagation and energy extraction codes (Fig. 4-1). With this data we constructed Table 4-2 that lists the amplifier and beam parameters for each stage appropriate for a 12.0-kJ shot in a 2.5-ns duration pulse. We define fill factor as the energy in the beam divided by the product of the aperture area and the maximum locally averaged fluence. Square pulse distortion is defined as the ratio of the gain experienced by the first photon of the pulse to that of the last photon. B-integral is the phase retardation, expressed in units of radians, accumulated because of the n_2 , the nonlinear index of refraction of the glass. Figure 4-2 shows the measured output energy versus input energy of the amplifier chain and the fit to the data generated by our Frantz-Nodvik gain saturation model that uses the measured beamline component parameters as input.

Frequency Conversion Experiments. The third harmonic conversion efficiencies of Nova output beams are lower than those predicted by calculations of third harmonic conversion of ideal beams, with the deviation becoming larger at high intensities. A major cause of this reduced conversion efficiency is a small amount of elliptical polarization that is nonuniformly distributed across the beam (Fig. 4-3). The sources of this depolarization are primarily the vacuum-loaded spatial filter lenses (see section 4.2.6.) and the laser disk amplifiers. This depolarization causes nonoptimal polarization mix ratios for the Type II/Type II conversion process, which degrades the third harmonic conversion efficiency. To further investigate Type II/Type II third harmonic conversion at high intensities and to evaluate two possible corrections for the low efficiency, we used our diagnostic station⁴ to study third harmonic conversion on the prototype beamline. The diagnostic station (Fig. 4-4) consisted of a Nova frequency conversion array, two calorimeters for

measuring 3ω energy, and a streak camera, a near field camera and a far field camera for measuring the temporal and spatial profiles of the generated 3ω beam.

We measured the depolarized fraction of the 1.05- μm beam energy at the frequency conversion arrays at 5-kJ laser output both before and after the glass replacement.⁵ The depolarization decreased from 4.0 to 1.7% with the glass replacement.

In our harmonic conversion tests we measured the conversion efficiency with this lower level of depolarization. We define the intrinsic third harmonic conversion efficiency as the average third harmonic fluence exiting the crystal array divided by the average first harmonic fluence at the output of the laser. Therefore, the definition of intrinsic conversion efficiency does not include the transmission of beam transport optics or beam obscurations present in the operation of the ten-beam target chamber. The Type II/Type II array produced a peak third harmonic conversion efficiency of 55% at a 1ω drive intensity of 1.5 GW/cm². At 2.5 GW/cm², the conversion efficiency decreased to 45%, presumably from depolarization effects. The transmission at 0.351 μm of the Type II/Type II KDP array used in the prototype beam experiments was less than optimal, which resulted in conversion efficiencies that are 3–6% lower than would have been obtained with a pristine array.

We conducted full aperture tests of two possible modifications to the frequency conversion system that reduce the effects of depolarization. In small aperture experiments, we have previously demonstrated that greater than 70% conversion to the third harmonic can be attained by repolarizing the output beam before the Type II/Type II KDP crystals,⁶ or by replacing the first element of the conversion system with Type I KDP doubling crystals.⁷ Type I/Type II third harmonic conversion is relatively insensitive to depolarization because first harmonic light polarized orthogonal to the ordinary axis of the doubler does not interact in the conversion process in either the doubler or the tripler.

In one series of tests we repolarized the prototype beam immediately before the conversion crystals using a full aperture multi-layer TiO₂/SiO₂ polarizer manufactured by OCLI. This polarizer was designed to provide a polarization extinction ratio that is variable over a range of two orders of magnitude by

Cold Transmission of Amplifiers		Small Signal Gain of Amplifiers		Measured Transmissions	
9.2's	0.9465 / stage	9.2's	6.69 / stage (6.87 in 1984)	Input sensor to MidDisk	0.60
15's	0.9609 / stage	15's	3.74 / stage (3.72 in 1984)	MidDisk to MidChain	0.83
20.8's	0.9143 / stage	20.8's	8.53 / stage (8.87 in 1984)	MidChain to Output Sensor	0.70
31.5's	0.9130 / stage	31.5's	15.1 / stage = 1.72 / amp		
46's	0.9298 / stage	46's	8.23 / stage = 1.69 / amp		

40-00-0492-1203
27 NAS

Fig. 4-1. Schematic of Nova beamline 7, including the transmission (T) and area magnification of each stage.

Fig. 4-2. Frantz-Nodvik gain saturation model with beamline 7 data points.

Fig. 4-3. The depolarization, consisting of 1.7% of the beam energy, after the glass replacement, is distributed in four lobes.

Fig. 4-4. Schematic of the Phase II diagnostic station.

changing the incident beam angle.

Repolarization improved the 3ω conversion efficiency at 2.5 GW/cm^2 drive intensity from 45% to greater than 65% and produced a 35% increase in the 3ω energy generated from 8.1 kJ of $1.05\text{-}\mu\text{m}$ drive in a 1-ns pulse. The conversion efficiencies have not been corrected for polarizer losses, which are approximately 2% for the horizontal polarization. The third harmonic conversion efficiency was insensitive to the polarization extinction ratio over the measured range of 20:1 to greater than 500:1. While it overestimates the conversion efficiency of the beam with 1.7% depolarization, our frequency conversion model⁸ fits the repolarized beam data well (Fig. 4-5). The model assumes that the beam is fiftieth power super-Gaussian in time and eighth power super-Gaussian in space with a Rician noise distribution based upon measured near field beam profiles that image the plane of the conversion crystals. The depolarization model included in the computer code divides the beam into 17 regions with various magnitudes of ellipticity of the polarization that sum to a total beam depolarization of 1.7%. However, the shape of the conversion efficiency curve depends upon the orientation of the polarization ellipse at each point in the beam, which has not been measured. The model also assumes that the 1.0-cm-thick doubling and tripling crystals induce a loss of 1.5% per surface and are detuned by $100 \mu\text{rad}$ internal angle, which may overestimate the angular phase-matching errors introduced by divergence in the beam and crystal alignment. Near field photographs of the third harmonic beam show that the regions of the beam with the highest depolarization convert poorly. After repolarization, the spatial distribution in fluence of the converted beam was much more uniform (Fig. 4-6).

The Type I/Type II conversion scheme produced comparable results. Figure 4-7 summarizes the data, showing a maximum conversion efficiency of 70% at a 1ω drive intensity of 2.4 GW/cm^2 . Our third harmonic conversion model, with the doubling parameters modified for Type I conversion, fits the measured conversion efficiencies well. The model includes an estimate of the deviation in our experiment from the optimal detuning angle of $200 \mu\text{rad}$, averaging the effects of detuning a 1.2 cm thick doubling crystal by $155 \mu\text{rad}$ and by $245 \mu\text{rad}$, with the Type II tripler detuned by $55 \mu\text{rad}$ in both cases. The model also includes Fresnel losses

on the crystal surfaces that are inferred from measurements. The Type I/Type II conversion scheme generated over 8 kJ of the third harmonic from a 12.5-kJ, 2.5-ns duration pulse (Fig. 4-8).

We observed further improvements in the third harmonic frequency conversion efficiency after replacing 15 m of the 68 m of air in the beam tubes between the final spatial filter lens and the KDP crystals with argon. With the Type II/Type II array, the conversion efficiency did not decrease at high fundamental drive intensities with argon, remaining constant at 55%. A smaller improvement was observed in Type I/Type II third harmonic conversion. Replacing a portion of the air path with argon can improve the conversion efficiency by reducing the path length over which nonlinear processes in nitrogen can occur. Two effects important for third harmonic conversion are: the optical Kerr effect, that causes polarization ellipse rotation;⁹ and stimulated rotational Raman scattering (SRRS),¹⁰ which frequency-shifts some of the first harmonic energy out of the bandwidth range that the KDP crystals can convert efficiently to the third harmonic. A spectrometer that sampled a small portion of an 8-kJ, 1-ns duration pulse immediately in front of the KDP array showed a strong first Stokes line that contained several percent of the beam energy. With 15 m of the air path replaced by argon, the spectrometer showed no SRRS signal.

We also investigated the second harmonic generation capabilities of the two conversion schemes. The Type II/Type II array with the crystals tuned in quadrature¹¹ converted a 7.7-kJ, 1-ns duration $1.05\text{-}\mu\text{m}$ pulse to the second harmonic with 74% efficiency. Full aperture repolarization had little effect on the second harmonic conversion process, which is less polarization sensitive than third harmonic conversion. The Type I/Type II KDP array, which has a characteristically poorer second harmonic conversion capability, produced a peak second harmonic conversion efficiency of 60%. Although not as efficient as Type II/Type II in quadrature, 60% conversion efficiency may be sufficient for most second harmonic experiments envisioned for the Nova ten-beam target chamber.

Tests on the Nova Phase II prototype beam indicate that Nova can exceed its original performance goals at $1.05\text{-}\mu\text{m}$ wavelength and that this output can be efficiently converted to the third harmonic. Energy output exceeded

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24NAS

Fig. 4-5. Type I/Type II 3ω conversion efficiency improves with repolarization.

Fig. 4-6 . Repolarization produces more nearly uniform conversion among crystal segments.

40-00-0492-1201

27NAS

Fig. 4-7. Type I/Type II measured 3ω conversion efficiency.

40-00-0492-1202
27NAS

Fig. 4-8 . Type I/Type II third harmonic energy generation.

Nova's performance goals on many shots with only minor damage to one optical component. A Nova Type II/Type II conversion array provided up to 55% third harmonic conversion efficiency; two possible upgrades provided 65% to 70% conversion efficiency. The upgrades require the installation of damage resistant full aperture polarizers between the output of the laser and the KDP arrays or the substitution of Type I doubling crystals for the Type II doublers that are the first elements in the present conversion arrays.

Author: Bob Ehrlich

4.2.3. Transverse Stimulated Brillouin Scattering Experiments Using the Prototype Upgrade Beamline

Introduction. In 1988 we conducted a series of experiments to study the onset of transverse stimulated Brillouin scattering in fused silica at 350 nm. Transverse SBS was first observed more than twenty years ago,¹² and for some time it has been anticipated that SBS in optical components might become an important energy loss mechanism in high-energy pulsed lasers.^{13,14} Eggleston and Kushner have recently published¹⁵ an analysis of transient transverse SBS in fused silica windows used in large nanosecond-pulse lasers. Our experiments confirm these predictions. We have also demonstrated that an increase in laser bandwidth suppresses transverse SBS as predicted. These results were published recently.¹⁶ We summarize them briefly here.

The acoustic wave generated by transverse SBS in fused silica at 350 nm has a decay time of about one nanosecond, so it is most appropriate to treat pulses of length less than a few nanoseconds with a transient stimulated scattering model.^{15, 16} That model shows that one can define a threshold parameter H for transverse SBS which is

$$H = I t_a t_g = \phi t_g,$$

where I is the intensity of the laser in the component, t_a is the growth time of the SBS acoustic wave, and t_g is the growth time of the transversely scattered optical wave. If threshold occurs at the end of the pulse, then t_a is just the laser pulse length and $I t_a$ is the laser fluence ϕ . The optical growth time t_g is the

shorter of either the time it takes the scattered optical wave to propagate across the transverse aperture of the component or the laser pulse length. Literature values of the Brillouin parameters of fused silica predict that transverse SBS will become significant at 350 nm when H exceeds a value of about 2.3 J ns/cm^2 . H scales to other wavelengths approximately as λ^2 , so transverse SBS is most important (i.e., has the lowest threshold) for short wavelength, large components and pulse lengths, and high laser fluence. The Nova Prototype Upgrade beamline can generate 3ω outputs which exceed this threshold by more than a factor of two for pulses longer than about 2 ns, so it is important to understand how transverse SBS affects its performance.

We set up an experiment on the diagnostic station shown in Fig. 4-4 to study the onset of transverse SBS in the large fused-silica beamsplitter in that station during the long-pulse frequency conversion studies discussed in the preceding section. The Nova beam enters from the right side of the figure at a beam diameter of 69 cm and is tripled using either a Type I or a Type II doubler followed by a Type II mixer as previously described. The beam then enters the beamsplitter, which is an 80-cm-diam window of ultraviolet-grade fused natural silica (Heraeus Amersil). The window has an average thickness of 3.75 cm and a wedge angle between the faces of 0.36 degrees to separate the two reflections for diagnostic purposes. The entrance surface is uncoated and reflects a sample to a beam diagnostics package, and the exit surface has a sol-gel anti-reflection coating for 350 nm. The transmitted beam strikes a full-aperture absorbing glass calorimeter. We used a calibrated filter of UG-11 glass (Schott) placed before the calorimeter to absorb unconverted fundamental and second harmonic radiation for some of these data points, and for others absorbed the full beam energy in the calorimeter and subtracted the fundamental and second harmonic components using relative energies determined from calorimeters in the diagnostics package which observed the entrance-face reflection. In addition to energy measurements this package has a streak camera that gives the time dependence of the incident 350 nm radiation, and a near field camera which observed the intensity distribution of the 350 nm radiation exiting the frequency conversion array.

Fig. 4-9. Beam profile of the Nova laser and the expected pattern of transverse SBS radiation. The polarization of the 350-nm beam is perpendicular to the disk split for Type I-II frequency conversion, and tipped at an angle of 35° for Type II-II conversion (heavy arrows at window periphery). We expect transverse SBS radiation to grow perpendicular to the polarization. The laser pulse length is comparable to the transit time across the aperture, so only limited areas of the beam can go over SBS threshold at a fixed time into the pulse. The heavily shaded area can see radiation growing perpendicular to the pulse for the full 2.4-ns pulse length: this will be the area which goes over threshold when the parameter H just reaches its threshold value at that time. The medium-shaded area is the additional area that goes over threshold if the time to threshold is 1.6 ns. The obscuration in the center of the beam causes a significant reduction in these areas for Type II-II conversion.

Figure 4-9 is a face view of the fused-silica window showing the incident laser beam shape. There is a dark strip with a width of about 6 cm in the center of the beam due to the split disks in the 46-cm laser amplifiers. There is an additional obscuration of about 120 cm² composed of 0.5-cm-wide strips running across the aperture which is not shown in **Fig. 4-9**. These strips are shadows of the support web that holds the array of 27 × 27-cm frequency conversion crystals. For the purposes of this experiment we shall ignore the support web obscuration and assume that the intensity is uniformly distributed over the two shaded D-shaped profiles in a total area of 3300 cm². The scale of the strip obscurations (and of other minor nonuniformities in the intensity profile) is much smaller than the laser pulse length, so the transverse SBS radiation sees a gain which averages over them.

Figure 4-9 also illustrates some of the complexities we expect to see in a non-ideal beam having obscurations and a laser pulse length which is not short compared to the transit time across the aperture. Suppose that the SBS threshold occurs at the end of a 2.4-ns pulse. Only those parts of the beam which can see photons propagating perpendicular to the beam polarization for the full pulse length of 2.4 ns will go over threshold at this point, which means only those regions which are heavily shaded in **Fig. 4-9**. There are some areas of this sort for both I-II and II-II frequency conversion, but the different polarization directions for the two types of conversion (as shown by the large barbed arrow) cause the area over threshold to be much smaller for II-II conversion because the scattered light must cross the dark disk split where there is no SBS gain. The medium-shaded regions show the additional areas which go over threshold when threshold occurs at 1.6 ns into the pulse. Here a much larger fraction of the beam area is involved, but the area affected with II-II conversion is still limited by the disk split. If we allowed the pulse to continue for the full 2.4 ns at the same intensity, there would be time for transverse scattered radiation blocked by the disk split to reemerge into an illuminated region and see further gain, so an even larger fraction of the beam area would be over threshold at that point.

We diverted a sample of the transmitted beam from within the heavily shaded region of **Fig. 4-9a** to a streak camera on some of the I-II shots to measure the time dependence of the

transmitted 350-nm radiation in the region which should be the first to show SBS loss. We also installed several instruments to observe side-scattered light on various shots, not all of which were active on every shot. These diagnostics included photodiodes in the positions shown on **Fig. 4-9**, light-sensitive paper wrapped around the periphery of the window to make qualitative observations of the distribution of side-scattered radiation around the window, calorimeters to measure the fluence exiting the window over a small area along the periphery, and a photographic Fabry-Perot interferometer to observe the spectrum of the side-scattered light.

We also studied transverse SBS in a Nova target focus lens. This lens was also made of fused silica (Corning 7940) and had a center thickness of 8.2 cm, edge thickness of 4 cm, and diameter of 80 cm. Its profile was biconvex, but one surface was nearly plane: it had a focal length of 300 cm. The lens had sol-gel anti-reflection coatings on both sides. The beam entered the lens through the steeply curved side polarized perpendicular to the disk split (**Fig. 4-9a**). The reflected beam diagnostics could not be used with the lens, but the other diagnostics were the same as for the window. The lens had an additional 16-cm diam central obscuration: this blocked about an additional 100 cm² of the beam, which was negligible.

Transverse SBS in the Fused-Silica Window.

The most accurate signature of SBS loss is the decrease in transmission through the window within the heavily shaded region of **Fig. 4-9**, since this measurement is not affected by details of the beam profile or scattering within the window. **Figure 4-10** shows the time dependence of the incident and transmitted pulses within this region for a typical shot with 5750 J in 2.4 ns at 350 nm within the 3300-cm² beam area, with Type I-II frequency conversion. The transmission begins to decline at about 1.6 ns into the pulse, and by the end of the pulse has fallen to about 30%. **Figure 4-10** also shows the predicted pulse shape for the simple threshold model discussed above with a square laser pulse of the same pulse length and fluence and a range of values of threshold parameter H . The experimental behavior is consistent with the predicted value of SBS threshold, giving an onset of serious loss at H of about 2.3 J ns/cm². Loss develops more slowly in the real beam than in a simple threshold model, as one would expect.

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Fig. 4-10. Shape of the incident (reflected), transmitted, and side-scattered pulse within the heavily shaded region of Fig. 4-9a for a shot which reaches threshold at about 1.6 ns into the pulse (5750 J, 2.4 ns). The transmission begins to drop rapidly for a threshold parameter slightly greater than two. The high-frequency noise on the top of the pulse is typical of the streak cameras used to measure these pulse shapes. The recovery of transmission at the end of the pulse is discussed in the text.

The whole-beam transmitted energy falls more slowly than the intensity in the small regions which are the first to go over threshold, but the losses can still become quite large. **Figure 4-11** shows a graph of energy incident on the window versus energy transmitted. The calorimeters for the incident and transmitted beams were cross-calibrated for zero-SBS-loss conditions using data from the 1-ns frequency conversion studies at up to 5 kJ energy at 350 nm: for these conditions the parameter H is less than 1.5 and we expect (and observe) zero SBS loss. The scatter in the 1-ns points which are plotted in **Fig. 4-11** shows the typical experimental scatter which we observe in data from these calorimeters. The graph also shows a number of experimental points at 2.4-ns pulse length and up to a little more than 8 kJ incident energy. It is clear that there is significant energy loss for these pulses at energies above about 5 kJ, with losses rising to nearly 2 kJ at the highest energy shots.

Figure 4-11 also shows the predicted energy transmission through the beamsplitter using a simple transient transverse SBS model¹⁶ with a range of values of H . This equation will always overpredict energy loss, since it assumes that the laser pulse length is much shorter than the aperture transit time and therefore that the entire area goes over threshold simultaneously. Our experiment violates this condition, as shown in **Fig. 4-9**. The model does suggest that the threshold value of H is less than about 3 J ns/cm^2 . For this particular beam geometry, H must exceed about 3.5 J/cm^2 to give a serious energy loss.

The intensity of side-scattered light within the window is a useful relative measure of the total quantity of scattered light and the evolution of SBS, but the interpretation of these intensities is not simple. The scattered light signal at the edge is a complex summation of SBS and other scattered light from multiple points in the aperture at different transit times, and will not show a simple behavior. **Figure 4-10** shows a side-scattered light signal from that shot displaying some of these complications such as the long "tail" from transit-time effects. One can see that the scattered light signal will rise more slowly with increasing pump intensity than the scattered light signal from the highest-gain part of the aperture, but should rise more rapidly than ordinary linear scattering from defects, surfaces, and other features. Well above

threshold the SBS scattered intensity will saturate, but there will also be further evolution as other nonlinear processes set in.

Figure 4-12 shows the relative peak intensity of the side-scattered light within the window as observed by photodiodes PD 1 or PD 3 in **Fig. 4-9** for several different experimental conditions. For each set of data points the peak signals have been normalized to a value of unity at 2000 J and are connected with dashed lines. The solid line marked A shows the predicted exponential growth of peak SBS intensity using the transient gain model,¹⁶ and the line marked B shows a linear growth proportional to intensity, both normalized to unity at 2000 J. It is clear that the intensities observed are consistent with a sum of linear and exponential components and that there is a saturation at high intensities, as expected. Photodiodes PD 2 and PD 4 show lower intensities and an intensity dependence which is closer to linear. This is also expected, since these diodes should see a lower SBS contribution than the diodes which are in the direction of highest SBS gain.

Light-sensitive paper wrapped around the periphery of the lens shows that the side-scattered intensity peaks perpendicular to the beam polarization, as expected. Shots predicted to go over threshold at about 1.6 ns have high intensity over roughly that part of the periphery which one would expect from **Fig. 4-9**. The fluence exiting the edge of the lens can be as high as several J/cm^2 for input energies of 7–8 kJ. We estimate that the energy exiting the side of the window is somewhat less than that missing from the beam in transmission in shots well over threshold, indicating that there is a significant energy loss to other processes under these conditions.

The Fabry-Perot interferometer which was available at the time of the experiment was not of high quality, and the spectra of side-scattered light are not very useful. They are adequate to show that there is a line shifted by the expected Brillouin shift of about 35 GHz on some shots just over threshold. Shots well over threshold suggest that the spectrum is broadened, perhaps by multiple SBS or self-phase-modulation. These shots are also slightly over threshold for stimulated rotational Raman scattering in the air path,¹⁷ however, which further complicates the spectrum and its analysis.

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Fig. 4-11. Transmitted versus incident energy for the fused silica window. The points marked with a small dot are shots at 1-ns pulse length, for which SBS is always below threshold. The scatter in these points shows the typical scatter in data taken on this test station. The other data points are coded as follows: (open circle) window, Type I-II conversion, narrow band; (solid circle) window, Type I-II conversion, broad-band; (open triangle) window, Type II-II conversion, narrow band; (X) lens, Type I-II conversion, narrow band. The point indicated by the arrow is the point for which pulse shapes are shown in Fig. 4-10.

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Fig. 4-12. Relative peak intensity of side-scattered light versus incident energy. The data points are coded as described in the caption to Fig. 4-11. Curve A shows the expected exponential growth of transverse SBS intensity. Curve B shows a linear growth in intensity as expected for simple linear scattering. These curves and the data are normalized to a value of unity at 2 kJ. The data are consistent with a sum of contributions varying as A (with gain saturation) and B. Note that the broad-band side-scattered radiation shows only linear growth.

Suppression of Transverse SBS by an Increase in Laser Bandwidth. We installed a phase modulator on the Nova oscillator to demonstrate the effect of an increase in laser bandwidth on transverse SBS: this strategy was chosen because a well-controlled frequency modulation avoids intensity spikes due to mode interference which could lead to nonlinear self-focusing and damage to components in the laser chain. The modulator uses the transverse electro-optic effect in lithium niobate and imposes a frequency modulation on the 2.4-ns laser pulse as shown in Fig. 4-13. We used a streak camera display of the time-dependent fringes of a Fizeau interferometer to measure the frequency deviation and its envelope. One can define an "average bandwidth" for this modulation which is the average value of the frequency deviation envelope of the FM signal. This average bandwidth is either 5.5 or 2.75 GHz over the 2.4-ns pulse for the two conditions under which we took data. The instantaneous frequency of the fundamental is tripled, so the average bandwidth of the 350-nm output is three times these values. Models for the effect of bandwidth on transverse SBS^{15, 16} predict that a bandwidth of about 20–30 GHz is required to suppress the SBS gain to below threshold at an incident energy of about 7 kJ in 2.4 ns. Then the average bandwidth imposed on the laser in the experiments was somewhat less than the prediction. The polarization was perpendicular to the disk split (Type I-II conversion, see Fig. 4-9a).

Figure 4-12 shows side-scattered light intensities for narrow-band shots (open circles) and frequency-modulated shots (filled circles) with conditions otherwise held constant. These intensities may be compared directly since they are from the same experimental setup. At shots from 5 kJ up to the maximum energy of the experiment at 7.3 kJ the side-scattered intensity is down by a factor of ~30 from the expected intensity without modulation. The two shots near 7 kJ are for the two different bandwidths, and the side-scattered intensities are indistinguishable. The side-scattered intensity from the entire sequence of modulated shots is consistent with a linear dependence on input intensity (Curve B). There is no loss in energy transmission to within the precision of the measurement (see Fig. 4-11).

The result of the experiment is that transverse SBS is suppressed to undetectable

levels for both a 16.5 and 8.3 GHz bandwidth on the laser. Theory^{15, 16} predicts that SBS should be slightly over threshold with the 8.3 GHz bandwidth, so the effective gain may fall a bit faster than suggested by that theory. The definition of "bandwidth" for a complex modulation such as that of Fig. 4-13 is debatable, however, and needs further study. It would be desirable to reduce the bandwidth further to make a quantitative determination of the bandwidth required for suppression and its dependence on the modulation waveform. We plan to acquire such data in the future when the proper apparatus and experimental time on the Nova facility become available.

The conversion efficiency of the tripler is bandwidth dependent, but calculations predict a negligible change in conversion efficiency for 350 nm bandwidth less than about 50–100 GHz. We see no difference in frequency conversion efficiency between modulated and narrow-band shots to within the precision of the experiment.

The transmitted intensity in Fig. 4-10 also shows a feature due to a frequency shift in the laser. Note that the transmission suddenly recovers to a high value during the roughly 100-ps-long fall time of the laser intensity. The intensity-dependent nonlinear index of refraction causes a phase shift in the laser during the pulse, and this shift is about 1.2 rad at the end of a 2.4-ns square pulse. As the intensity falls to zero, this shift must also fall to zero. A linear shift of 1.2 rad over 100 ps corresponds to a frequency shift of about 2 GHz during that period, which is tripled to a shift of 6 GHz at 350 nm. This abrupt shift is enough to take the pump out of resonance with the SBS acoustic grating, and the transmission promptly recovers.

Transverse SBS in the Fused Silica Lens. The angle between the pump beam and scattered light moving transverse to the beam within the lens has a linear change of 9° between the opposite ends of a beam diameter. The effect of this change in angle is equivalent to a frequency chirp of 300 MHz per degree.¹⁶ If transverse SBS goes over threshold at 2.4 ns into the beam, the scattered light must propagate for 2.4 ns, which is 70% of the aperture. A photon moving along a diameter of the lens perpendicular to the SBS polarization will then see an effective chirp of about 2 GHz during the pulse at a rate of about 0.8 GHz/ns. If threshold occurs at 1.6 ns, the effective chirp will be proportionally smaller, or about 1.4 GHz, at the same chirp rate. Theory

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Fig. 4-13. Waveform of the frequency modulation which was used to study the effect of laser bandwidth on transverse SBS. The frequency deviation is a 2.2 GHz sinusoid damped with a 1-ns exponential, and commences about 0.5 ns after the beginning of the pulse. The average bandwidth over the 2.4 ns laser pulse is either 5.5 or 2.75 GHz: the peak-to-peak deviation of the exponential envelope is about three times these values.

predicts that bandwidths of this magnitude, although many times larger than the Brillouin line width of 170 MHz, are not sufficient to show any significant SBS gain suppression.

The results of the experiment confirm that prediction, as shown by the lens data points in Figs. 4-11 and 4-12. We are unable to see any difference in SBS threshold or the magnitude of SBS loss between the converging beam in the lens and the collimated beam in the window, to within the precision of the experiment.

Damage to Optical Components as a Result of Transverse SBS. High intensity stimulated scattering experiments are frequently associated with optical damage,^{12, 15} and we saw several types of serious optical damage in these experiments. It is important to note that the damage is not the result of the transverse SBS process itself, but of subsequent interactions. The energy deposited in the optical component by the phonons participating in SBS is only a few tens of millijoules, which is much too small to have any damaging effect.

When transverse SBS is well over threshold, the process evolves so that the side-scattered light comes from a very shallow depth beneath the entrance face.¹⁶ This light can be much more intense than the input beam and has a long propagation distance in the optical component, so there is a high risk of beam breakup due to self-focusing. After two high-energy shots with incident energy of about 8 kJ and side-scattered losses of about 1.5–2 kJ, we saw two types of severe damage to the window within this region. The first of these types was a structure consisting of long linear “tracks” of scattering sites about 3–5 mm below the entrance surface and oriented perpendicular to the incident beam polarization. These features have the same appearance as damage due to filamentation and small-scale self-focusing, and we attribute them to exactly that sort of beam breakup in the intense side-scattered SBS radiation. We also saw damage to the subsurface layer of the optical entrance surface. Optical grinding and polishing operations typically leave behind a fine characteristic pattern of sub-surface “microcracks” where chemical bonds are broken, but the “crack” is sufficiently narrow that it scatters a negligible amount of light. After these two high energy shots, the microcracks on the entrance surface of the window appeared to have opened to sufficient width to show significant light scattering in the visible region.

The interactions of the side-scattered light with materials around the periphery of the optical component caused a third type of damage. The components used here are circular and are mounted in holders using elastomeric supports (O-rings or pads) which also have circular symmetry. Intense side-scattered light deposits energy in these supports and generates a highly symmetric circular acoustic transient, as shown in Fig. 4-14. That transient then comes to an acoustic focus at the geometrical center of the component, and stresses in the focus are sufficient to fracture the fused silica.

We installed an array of strain gauge transducers on the Nova focus lens during our experiments to follow the propagation of acoustic impulses within the lens and verify that the damage was indeed due to acoustic transients as we had inferred. Figure 4-14 plots the propagation of an acoustic transient generated at the time the laser is fired as it converges on the center of the lens. It is clear from the propagation velocity of this transient that it is a Rayleigh wave on the surface of the lens, and that when the laser pulse strikes the lens at $t = 1.5 \mu\text{s}$ after the system timing pulse, that transient is located at a radius of 39.6 cm on the lens face. This is exactly the location of the O-ring that forms the vacuum seal on the back face of the lens. Figure 4-15 is a sketch of the edge of the lens showing how this happens. The O-ring absorbs scattered light, which generates a plasma between the O-ring and lens surface. This impulse launches both a Rayleigh (R) wave on the surface and a compression (P) wave into the bulk of the lens. When we raised the laser energy above about 5 kJ in 2.4 ns, at which the SBS-side-scattered light had an energy of about 150–200 J, the Rayleigh transient focused to an acoustic intensity large enough to fracture the fused silica in the center of the rear surface of the lens, as shown in Fig. 4-16. We estimate that the energy absorbed by the O-ring was perhaps 50–100 J at this point.

This experiment shows that transverse SBS can easily provide enough scattered light to initiate a fracture in Nova optics. Absorption of something less than 1% of the Nova beam energy is enough to initiate fracture, however, so it is clear that other sources of scattering might also put circular components at risk in high-energy beams. There have been several previously unexplained instances of optical components which have developed small central fractures

Fig. 4-14. Radius vs. time plot for an acoustic transient in the focus lens. The transient propagates at a velocity of 3450 m/s, which is an excellent match to the Rayleigh surface wave (R) velocity of fused silica (3430 m/s). Compression (P) and shear (S) waves propagate at 5920 and 3430 m/s, respectively.

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Fig. 4-15. Sketch of the edge of the Nova focus lens showing the source of the acoustic transients.

Fig. 4-16. Photograph of the center of the back surface of the focus lens after 50–100 J absorbed in the O-ring generated a focusing transient strong enough to fracture the lens surface at the focal point.

during routine Nova operations, and we now infer that acoustic foci from transients generated by absorption in mounting pads or O-rings around their periphery were probably the cause of these fractures.

We also observed several other types of acoustic transient in the lens and beamsplitter during this series of experiments, including the compression wave shown in Fig. 4-15, but we have no direct evidence that these transients caused damage as we have for the Rayleigh wave.

Conclusions. The experimental results show that transverse SBS in optical components sets in roughly where predicted by models such as that of Refs. 15 and 16. High scattered light intensity can result in damage to the optical component from subsequent events, particularly from the focus of acoustic transients in the optical component.

An increase in laser bandwidth reduces the growth rate of SBS. We have demonstrated that an average bandwidth of about 8 GHz is sufficient to reduce transverse SBS to undetectable levels for any conditions accessible under normal operation of Nova after the upgrade is completed. This bandwidth is slightly less than the requirement predicted by a simple theoretical model, although this point needs further study. Bandwidths of this magnitude have no effect on the efficiency of frequency conversion in typical frequency tripling schemes. We shall install a permanent phase modulator on the Nova oscillator so that it operates routinely with bandwidth large enough to give negligible transverse SBS.

Transversely-scattered light can interact with absorbers around the periphery of circular optical components and generate highly symmetric acoustic transients which then focus on the geometrical center of the component and cause fractures. This mechanism must be carefully considered in the mechanical and optical design of optics which handle high energy pulses. As an example, note that none of the high-energy optics in the Athena design for future fusion lasers are round.

Author: J. R. Murray

Major Contributors: J. R. Smith, A. E. Brown, R. B. Ehrlich, R. B. Engle, D. T. Kyrakis, C. E. Thompson, T. L. Weiland, and R. B. Wilcox

4.2.4. Improving the Near-Field Beam Quality of Nova

A primary constraint on operating Nova at high output power and energy is damage to large and expensive optical components. To minimize damage and other undesirable nonlinear optical effects, it is important to make the beams as uniform as possible throughout the laser. Near-field photographs and burn patterns taken during Nova's early years showed sources of intensity nonuniformities that could be corrected. First, gradients in fluence across the beam arose from asymmetries in the beam input from the preamplifier and from nonuniform gain across the aperture of the amplifiers. Second, high spatial frequency noise was caused by diffraction from contamination, damage, and imperfections in the optics. Third, regular circular-ring-type modulation arose from spatial filtering of the diffraction patterns resulting from the hard apertures used to fix the input beam diameter and from beam clipping caused by imperfect alignment of components.

To improve the beam uniformity, we made changes in the preamplifier and in the amplifier beamlines concurrent with replacing the disks in the amplifiers. These were first tested on the Phase II prototype beam and include: (1) improving the beam input from the preamplifier by refining alignment and cleaning or replacing dirty and damaged optics; (2) replacing the circular beam-defining apertures in the preamplifier and at the input to each beamline with soft apodizing apertures; (3) decreasing the beam size through the smaller diameter disk amplifiers and increasing the magnification of the 20.8-cm to 31.5-cm spatial filter; (4) reversing one 31.5-cm amplifier to compensate for the gain asymmetry in the disk amplifiers.

These changes resulted in substantially improved beam uniformity. Circular-ring modulation has been almost completely eliminated. Gradients in fluence across the beam have been reduced but some asymmetries remain that could possibly be identified and corrected. At 8-kJ output in a 1-ns pulse, the peak-to-average fluence, measured on near-field photographs taken in the output sensor of the prototype beamline, was reduced by approximately 40%.

In the 1987 annual report we described the beginning of this work on the prototype beamline.¹⁸ Here we discuss the apodizers

chosen, changes in staging that reduced the beam modulation midway through the amplifier chains, and measurements made on the improved prototype beamline.

Improved Apodization . In the 1987 annual report, we reported on MALAPROP propagation calculations that showed circular-ring modulation of the Nova beam could be almost completely eliminated by substituting soft apodizing apertures for the hard apertures being used to fix the beam size in the preamplifier and at the input of each amplifier chain. As a result, we began a search for a durable, easy to fabricate apodizer to use on Nova. The apodizer design we adopted was suggested several years ago by John Trenholme¹⁹ and more recently by Steve Patterson.²⁰ The smooth circular aperture is replaced by an aperture with a serrated edge.

An aperture with a serrated edge works as an apodizer when the high spatial frequencies induced are stripped or diffract out of the beam. **Figure 4-17** illustrates this process. It displays the serrated aperture used at the arm input, the fast Fourier transform (FFT) of the apertured beam (equivalent to intensity pattern at the focal plane of the following spatial filter), the portion of the spatial-frequency spectrum transmitted by the spatial filter, and the inverse FFT of the filtered spectrum. The spatial profile of the soft aperture function is superimposed on the final figure. For comparison, **Fig. 4-18** is the equivalent sequence for a regular circular aperture. Proper filtering of the serrated edge produced a soft profile and suppressed modulation.

The choice of functional form for the serrations was somewhat arbitrary. We investigated three serration functions: a linear tooth, an S-shaped tooth, and the function $r = r_0 - \delta r \exp(\alpha \cos 2\pi\theta / \phi)$. The cosine exponential gave somewhat better results in both calculations and experiments, thus we adopted it as the serration function. The numerical values used for the input aperture on the beamlines were $r_0 = 12.8$ mm, $\delta r = 1$ mm, $\alpha = .693$, and $\phi = 5^\circ$. Refinements in fill factor are possible by making small adjustments in r_0 .

To improve the apodization of the split disks in the center of the 46-cm amplifiers we replaced the original diffusing apodizer with a linear serrated apodizer as shown in **Fig. 4-19**. To minimize energy loading on spatial filter pinholes and to ease alignment, we placed this linear apodizer at the input to the 37 to 94 mm

spatial filter (the second spatial filter in each amplifier chain). This new apodizer reduced the diffraction peaks along the disk split observed in near-field photographs taken in the output sensors.

Beam Size Reduction in Amplifier Driver Section. The magnification of the spatial filters in the Nova beamlines were originally chosen to uniformly fill the various size disk amplifiers to take advantage of their full clear apertures. As a consequence the periphery of the beam passes through disks very close to the edge cladding. These regions of the small disks are under stress from the fused glass bonding of the edge claddings, resulting in severe aberration of the beam at its periphery. Much of the energy in these aberrated portions of the beam is removed by the spatial filters, leaving a strong ring modulation at the cutoff frequency of the spatial filters (see **Fig. 4-20a**). The new 31.5- and 46-cm disks with epoxy bonded edge cladding do not exhibit this edge aberration and have good wave front across their full aperture.

To eliminate the aberration build-up in the small disks, we reduced the beamline input aperture size from 26 to 23 mm, decreasing the beam diameter by about 10%. This change eliminated the ring modulation on the beam at the input to the 31.5-cm stage, approximately half-way through the amplifier chain. **Figure-20a** and **b** illustrate the improvement that this single change made to the beam at the input to the 31.5-cm stage.

To take advantage of the full aperture of the new amplifier disks we increased the magnification of the mid-chain 20.8–31.5-cm spatial filter to compensate for the reduced input aperture size. The result is a smoother beam input into the 31.5-cm amplifiers that fills their aperture. MALAPROP calculations of 15-kJ output from the beamline in a 3.0-ns pulse indicate a modest increase in B integral and pulse distortion from the smaller beam in the front half of the chain. We believe this is more than compensated for by the better beam quality.

Measurements of the Output Beam. To monitor the near-field beam during the Phase II prototype beam experiments, we took beam photographs in the beam 7 output sensor. We modified the near-field camera to photograph the beam after an effective propagation distance of 68 m beyond the final spatial filter lens, which is approximately the distance to the frequency conversion crystals on the Phase II

(a)

(b)

(c)

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Fig. 4-17. Calculated apodization produced by a serrated aperture and spatial filter showing: a) the beam uniformly illuminating the serrated aperture; b) the Fourier transform of (a) with the cut-off frequency of the spatial filter; and c) the inverse Fourier transform of the filtered spatial frequency spectrum.

(a)

(b)

(c)

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Fig. 4-18. Same as Fig. 4-17 but for a uniformly illuminated circular aperture.

Fig. 4-19. The serrated strip used to apodize the split in the 46-cm amplifier disks.

(a)

(b)

Fig. 4-20. Photographs of burn patterns taken at the input to the 31.5-cm stage showing the improved beam uniformity with a smaller beamline input aperture.

diagnostics table. We used our standard methods of calibrating and developing the film, which we subsequently scanned with a densitometer, and analyzed with computer codes to determine the fluence distribution and the size of the beam.

The film response was determined with the "dlogc-2spot" code that generated a density versus log exposure curve by comparing pairs of calibration spots.²¹ The response curves generated from different pairs of spots on the same piece of film were compared to check consistency.

The digitized densitometer data, film response curve, beam energy and pulse length were input to the "beam-analysis" routine to produce a near-field intensity profile, from which an intensity histogram, radial average plot, lineout, contour plot, and integrated intensity distribution were generated.

We analyzed photographs taken before and after decreasing the beam size in the driver section and replacing the input aperture with a serrated apodizer. The reproductions of photographs in Fig. 4-21 a and 4-21 b compare the output beam in the two cases, showing the reduction of the concentric rings that arise from spatial filtering of the diffraction patterns produced by the hard aperture and aberrations near the edges of the amplifier disks. Reversing one 31-cm amplifier to "flip" its horizontal variation in gain improved but did not eliminate the asymmetric illumination across the beam. Lineouts taken through photographs of 5-kJ shots show that the fluence varies by up to 35% along a diagonal diameter of the beam. This effect is reduced at higher energies by gain saturation and is much less evident in 12-kJ shots.

A good indication of the improvement in uniformity that our efforts produced is the reduction of the peak-to-average fluence in the output beam (Fig. 4-22). We measured the peak-to-average fluence to be 1.45 ± 0.08 on shots at 5 kJ in 1.0 ns and 1.25 ± 0.05 at 8 kJ. Before the changes, the comparable peak-to-average ratios were 1.65 ± 0.1 and 1.40 ± 0.08 . The average fluence is calculated by dividing the output energy by the beam area determined from the photographs. The peak fluence is chosen at the level below which 99.5% of the total energy is contained. The remaining 0.5% of the energy is contained in points that lie along the edges of the disk split. This modulation was subsequently reduced by the installation of a serrated disk

split apodizer. The peak fluences that we measured compare well with modulation calculated with the MALAPROP code.²²

We calculated the effective area of the beam from plots of radially averaged intensity versus radius. The effective area is the area of the flat-topped portion of the beam, plus the area of the decreasing intensity regions around the edge, weighted by intensity, minus the area of the 6-cm-wide disk split. Analysis of several shots at different energies (at both 1 and 2.5 ns) shows that the $r_0 = 12.5$ -mm serrated input aperture and the spatial filter modifications slightly reduce the effective output beam size, changing somewhat with output energy (Fig. 4-23). The effective beam area is about 7% smaller than with the previous configuration for 5-kJ output and 2% smaller at 12-kJ output. The effective area ranges from 3050 to 3300 cm². The increase in size as a function of energy is probably caused by gain saturation in the large amplifiers, results in the wings being amplified more than the higher-intensity circular flat-top region of the beam. The fill factor, defined as the effective area of the beam going through the amplifiers divided by their clear aperture, is now 71% for a 5-kJ shot and 77% at 12 kJ. We installed slightly larger input apertures ($r_0 = 12.8$ mm) to recover the small decrease in beam size. Photographs of the beam with the larger apodized aperture show no obvious changes in the quality of the beam although the analysis of these photographs has not been as detailed as with the 12.5-mm apertures.

Author: D. R. Speck

Major Contributors: R. B. Ehrlich, J. K. Lawson, R. L. Saunders, and T. L. Weiland

4.2.5. Assessment of Damage to 1-Meter Diameter Nova Turning Mirrors

I. Introduction/Background. A continuing problem with all large, high-peak-power laser systems such as the LLNL Nova laser is damage to key optical elements. The high reflectivity, dielectric coatings that are used on Nova to steer the output laser beams to the target chamber are perhaps the weakest elements (in terms of damage) on the system. In this section we summarize the results from an assessment of damage to the 1-m diam mirrors on Nova. Further details of this work can be found elsewhere.²³

(a)

(b)

Fig. 4-21. Photographs of the output from beam 7 taken in the output sensor
a) before and b) after modifications to improve the beam uniformity.

Fig. 4-22. Peak to average spatial modulation of beam 7 versus output energy for 1 ns pulses measured with old laser configuration •, measured with improved laser configuration ◊, calculated with MALAPROP +.

Fig. 4-23. Measured areas of beam 7 versus output energy with old laser configuration • and improved laser configuration ◊ with $r_0 = 12.5$ mm serrated input aperture.

The Nova HR coatings receiving the highest fluences are located at the mirror denoted by the letters "C" and "E" in Fig. 4-24.²⁴ For example, for a nominal 125-kJ, 1-ns shot, the average fluence on these mirrors is about 4 J/cm².

Assuming a modulation of the beam intensity of 1.5:1 (typical for Nova) this corresponds to peak fluences of roughly 6 J/cm².

The "C" and "E" mirrors have the same HR coating materials. However, the "C" mirrors serve a diagnostic function that allows some light transmittance and consequently they have fewer layers.

Recently an "E" mirror that had been badly damaged was removed from the system. The measured reflectivity across the mirror surface is shown in Fig. 4-25. In the most heavily damaged regions the reflectivity had dropped to roughly 66%. In addition, the mirror had a uniform "haze" layer over the area of the beam imprint. This layer could be easily removed by wiping, and contributed about 5% to the overall drop in reflection. Microscopic inspection of undamaged areas of the HR coating (that is, areas outside of the beam path) showed a large number of "particles" with diameters up to 30 μm .

Inspections of other mirrors on Nova showed similar results although the other mirrors do not seem to be as badly damaged. Furthermore, microscopic inspections of the other HR coatings show that "particles" are always present in the coatings. Consequently we assume that the assessment of the damage to one of the Nova mirrors given here is a fair representation of that expected for the other high fluence mirrors on Nova.

II. Analysis of Nova HR Coating Damage.

A. Density and Morphology of Coating Defects.

Microscopic inspection of undamaged HR coating regions shows a large number of visible defects (Fig. 4-26 a). Based on the photomicrographs of the coating, we estimate there are about 80 to 100 mm² in the 5–30 μm range. From the shape of the defects (i.e. circular) one naturally draws the conclusion that these are small particles embedded in the coating. However, we suggest that they are not particles simply because they are too large relative to the coating thickness and thus would not give the observed smooth coating surface. This point is illustrated in Fig. 4-27 where the total coating thickness (about 5 μm) is shown relative to an embedded 30- μm particle. Note that if many particles of the 30- μm size were

embedded in the coating one should be able to easily see them when looking at the coating surface from an oblique angle. Since the coating surface is very smooth (less than 25–50 Å rms surface roughness) then it is assumed that the defects are not particles (i.e. "spheres") but are more nearly two-dimensional (i.e. plate-like).

Of course the next obvious question is what are they if they are not particles? So far that remains unknown; all our attempts to quantitatively analyze such defects have failed.

B. Correlation of Damage with Observed Defects. Photomicrographs taken of damaged and undamaged regions of the Nova HR coating allow a one-to-one visual comparison of damaged and undamaged regions of the coating (Fig. 4-26 a and 4-26 b). We observed that the average damage "spot" was larger than the average defect and that in many instances the damage at one or more adjacent sites coalesced into a larger area. We also found that the density of damage spots was greater than the density of defects. For example, the density of damage spots in Fig. 4-26 b was about 300 mm² or about three times the original defect density.

The fact that we did not observe any regions remaining undamaged led us to conclude that visually observed defects do indeed cause damage. However, they don't seem to occur in sufficient numbers to account for all the damage. This latter observation has to be tempered with two facts: first the photomicrography of two different regions on the mirror surface and the defect density may have been slightly different at the two sites; second, a number of "barely visible" defects (<2 to 4 μm) were seen in the undamaged coating that were not included in the original (undamaged) defect-density estimate.

C. Source of the "Haze" on the Mirror Surface. Analysis of the particles that produced the haze on the surface of the mirror show they are ZrO₂ and SiO₂; the same materials as the coating. Thus the coating material blown off during damage created particles of ZrO₂/SiO₂ that litter the damaged surface. SEM images of the particles show that they are less than about 1 μm in diameter and have both jagged and smooth surface morphologies. The smooth particles probably nucleated and grew from a vapor phase whereas the jagged particles are simply ejected solid debris, i.e. not vaporized.

One possible scenario for the damage process is shown in Fig 4-28. Before damage, the coating

ACS-S42-U-5986-4

Fig. 4-24. Schematic diagram of a Nova beamline showing the positions of various mirrors. The mirrors in the high fluence positions are labeled as "C" and "E" and their key characteristics are summarized in the insert.

BE 533 U 5986 2

Fig. 4-25. Measured reflectivity at 1.06 mm at various locations on the mirror surface.²⁴ Wiping the haze from the surface increases the reflectivity about 5% above the values shown. The small diamonds indicate areas of heavy damage.

RC-U-9286-005

Fig 4-26. Photomicrographs of the (a) undamaged and (b) damaged region of the Nova mirror HR coating.

RC-U-9286-006

Fig. 4-27. Schematic diagrams showing the effects of (a) spherical or (b) plate-like defects on a nominally 5-mm-thick Nova HR coating. Note the large amount of surface roughening that would be easily visible for the case of the spherical (3-D) defects.

(a) Before

(b) During

(c) After

LW-S12-U-5986-5

Fig. 4-28. Simple physical picture of the proposed process for damage to Nova HR coatings (a) before, (b) during and (c) after damage.

RC-U-9286-004

Fig. 4-29. Predicted temperature rise at the incident surface of an impurity vs. the impurity absorption coefficient. The calculated temperature rise is per Joule/cm² of absorbed incident laser energy. Thus, for example, for 10 J/cm² of absorbed incident fluence, simply multiply the temperature rise by 10.

contains a number of defects that, when irradiated, heat to a temperature high enough to vaporize the surrounding $\text{ZrO}_2/\text{SiO}_2$ coating. Based on the normal boiling points of these materials, this temperature has to be greater than 3000 K. The injected material consists of a mixture of solid fragments and molten or vaporized SiO_2 and ZrO_2 . The vapor cools, nucleates and forms either single or agglomerated particles.

Based on the extent of damage (about 20–30% of the coating area) and the original coating thickness of about 5 μm , one can estimate the thickness of the debris layer. This is about 2 μm assuming the debris has a porosity of 50%.

The observation that the vaporized SiO_2 and ZrO_2 nucleates, grows and agglomerates is not surprising. Simple thermodynamic theory^{25–27} shows that growth will occur once a molecular cluster exceeds a critical radius, r_c , given by:

$$r_c = 2GV/(RT \ln(p/p_c)) \quad (1)$$

where G is the surface energy (erg/cm^2), V is the molar volume (cm^3/mole), R the molar gas constant ($\text{erg}/\text{mol K}$) and T the temperature (K). The driving force for nucleation and growth is the departure from equilibrium represented by the ratio of measured and equilibrium vapor pressures, i.e. p/p_c . Assuming SiO_2 vaporization at 3000 K and subsequent cooling to 1000 K, it is easy to show that the critical cluster radius is less than the SiO_2 molecular dimension. This simply implies that the system is so far from equilibrium under these conditions that nucleation occurs almost instantaneously. The point here is that simple theory predicts that the SiO_2 (or ZrO_2) vaporized during damage rapidly nucleates and grows forming many of the small particles that comprise the debris we see littered on the mirror surface. Moreover the debris particle size distribution ($<1.0 \mu\text{m}$) and morphology is consistent with this observation.

III. Possible Damage Mechanism in the Current Nova HR Coatings. It is clear from the data presented above that damage is caused by localized heating at defect sites in the coating. The question that always remains is: exactly how is the laser energy coupled into the material? One mechanism that is generally accepted for thin films is the so-called impurity model that simply assumes bulk absorption of the energy by some extrinsic defect, e.g. “dirt” buried in the coating. However, one question that

is rarely asked in relation to the impurity model is what could the impurity be? In other words, what materials have a bulk absorption high enough to give the temperatures required for vaporization (i.e. damage) and is it likely they would be in the coating?

To assist in our search for such materials, we used a simple 1-D heating model to help bracket the probable range of absorption coefficients. This model assumes an absorption profile through the defect given by

$$\frac{I}{I_0} = \exp(-\alpha x) \quad (2)$$

where I_0 (J/cm^2) is the incident laser fluence, I is the fluence at position x (cm), and α (cm^{-1}) is the absorption coefficient. Corresponding temperature rise through the absorbing region is:

$$\Delta T(x) = \frac{\alpha I}{\rho C_p} \exp(-\alpha x) \quad (3)$$

with a maximum at the front surface ($x = 0$):

$$\Delta T(0) = \frac{\alpha I}{\rho C_p} \quad (4)$$

where ρ is the material density (g/cm^3) and C_p is the specific heat ($\text{J}/\text{g} \cdot \text{K}$). Thus for a given laser fluence one can get a reasonable idea of the impurity absorption required to heat the sample to the boiling point of the coating materials (i.e., SiO_2 and ZrO_2) and then, in turn, bracket possible impurity materials that might be present. [For these “back-of-the-envelope” estimates we have neglected the enthalpy change associated with melting (i.e. heat of fusion).]

There is one thing to note about Eq. (4), that is, it is pulse length independent. This is valid only if the characteristic absorption depth (i.e. $1/\alpha$) is long compared to the characteristic thermal diffusion distance x :

$$x = 2(\kappa t)^{1/2} \quad (5)$$

where κ is the thermal diffusivity (cm^2/s) and t is the laser pulse width (s). This is a very important point because it is well known that damage thresholds generally seem to scale with

laser pulse width somewhere between $t^{0.25}$ – $t^{0.5}$. In Table 4-3 is a list of common dielectric coating materials with their associated thermal physical properties; these data are from Lange et al.²⁸ The calculated characteristic thermal transport distance (x) for a 1-ns pulse is given in the far right column; for these dielectrics the thermal transport is typically less than 0.1 to 0.2 μm . Consequently if the absorption coefficients are significantly less than $1/x$ (i.e. $<10^5 \text{ cm}^{-1}$), the damage threshold should scale only weakly with pulse length in the nanosecond-or-less region. As will be shown below, even relatively strong absorbing dielectrics have absorption coefficients less than this at 1.06 μm . Thus, one would expect a very small pulse-length-dependence for the damage threshold if this impurity-heating model is valid. (In the case of the original Nova HR coatings we have found a negligible pulse length scaling of the 1.06- μm damage threshold between 1 and 10 ns (Ref. 29.)

For comparison, we've also included a list of metals in Table 4-3 with their thermal-physical properties and characteristic thermal diffusion distance in 1-ns. Note that for metals the optical absorption depth (<10 -2 mm) is much less than the thermal diffusion distance ($\sim 1 \mu\text{m}$). Thus one would expect damage due to metal particles to scale as approximately $t^{0.5}$ in the ns region (for particles $>1 \mu\text{m}$). This is indeed observed, as in the case of the Pt-inclusion damage in glasses.^{30,31}

The question that remains is what dielectric coating materials have absorption coefficients at 1.06 μm high enough to cause thermal heating to ~ 2000 – 3000K ? Figure 4-29 is a plot of calculated maximum temperature rise (per J/cm^2 absorbed) vs. absorption coefficient. Since the original Nova mirrors damage at roughly 7 – $8 \text{ J}/\text{cm}^2$ then the impurity absorption would have to be $>1000 \text{ cm}^{-1}$ to give the observed thermal heating. So, what are some possibilities? One is SiO , a suboxide of silica that is often proposed to be present in e-beam SiO_2 coatings. SiO has an absorption coefficient of $<100 \text{ cm}^{-1}$ at $1 \mu\text{m}$ (Ref. 32) so it would not give the required temperature rise. Some other common dielectrics that are also strong absorbers at $1.0 \mu\text{m}$ include FeO , Fe_2O_3 , CuO_2 , CuO , and NiO . Based on absorptions at $1.0 \mu\text{m}$ for these materials when doped in glass, we have estimated bulk absorptions (Table 4-4). Note that almost all of these have absorptions greater than 1000 cm^{-1} . So one could propose that the coating is contaminated by numerous

inclusions of materials of this type. However, that is unlikely; the defect density is about $100/\text{mm}^2$ so there would have to be a fairly large contamination problem within the coating chamber. Moreover, the contaminants would have to be flat-plates (not particles) to agree with the observation that the defects are largely two-dimensional, producing negligible surface roughening.

This leads us to two main conclusions regarding the mechanism for damage to Nova HR coatings. First, a damage mechanism based solely on absorbing dielectric particles (i.e. dirt) is possible but unlikely. Such a mechanism would require plate-like (2-D) particles that have very high absorption coefficients. The very high absorption values required to heat the particles to the point of damage eliminate all but the most strongly absorbing dielectrics. Moreover, the concentration of defects in the Nova coatings is so high it is difficult to believe there is an outside contamination problem of this magnitude.

Second, it is clear from the data that visible defects do lead to damage, so they must be absorbing energy by some mechanism. We believe it may be due to electron donor (or acceptor) states lying close enough to the conduction (or valence) band to undergo excitation by linear absorption. High densities of these states may be associated either with the bulk or boundaries of the observed defects leading to sufficient absorption to produce the heating observed. Thus one would expect the damage to be associated with observed physical defects. Notice that the scale length of the region containing these absorbing states would need to be >0.1 to $0.2 \mu\text{m}$ to explain the negligible pulse length scaling that is observed. This is on the order of a single-layer coating thickness suggesting a plate-like defect region equal to the coating thickness and up to about $30 \mu\text{m}$ in diameter.

Author: J. Campbell, G. Edwards

Contributors: R. Wolfe, E. Lindsey

4.2.6. Nova Beam Depolarization

To experimentally determine the effects of reversing spatial filter lenses, we measured the level of depolarization of one Nova beam just before and just after reversing the final two output spatial filter lenses. To measure depolarization, which is expressed as the fraction of the total beam energy contained in

Table 4-3. Summary of thermal-physical properties reported by Lange et al.²⁸ for a number of crystalline dielectrics often used in optical thin films and, for comparison, several metals.

Material	Index	k (J/cm ² •s•K)	C _p (J/g•K)	ρ (g/cm ³)	T (melting) (K)	T (boiling) (K)	κ (cm ² /s)	— x (μm)
I. Dielectrics								
Al ₂ O ₃	1.72	0.346	0.774	4	2325	3250	0.112	0.21
HfO ₂	2.25	~0.017	0.168	8.68	3137	~5600	0.012	0.068
LaF ₃	1.59	0.0508	0.462	5.94	1763	~2600	0.019	0.086
LiF	1.37	~0.017	1.61	2.635	1121	1950	0.004	0.040
MgF ₂	1.43	0.116	0.394	3.176	1536	2512	0.093	0.193
MgO	1.83	0.6	0.92	3.77	3098	3870	0.173	0.263
SiO ₂	1.44	0.14	0.737	2.32	1696	2500	0.082	0.181
ThF ₄	1.59	0.1	0.359	6.32	1380	~2000	0.044	0.133
ThO ₂	1.90	0.142	0.234	10	3540	4670	0.061	0.156
Y ₂ O ₃	2.10	0.0246	0.451	5.01	2683	---	0.011	0.066
ZrO ₂	2.25	0.011	0.46	5.7	2950	4570	0.0042	0.041
II. Materials								
Au	--	2.97	0.131	19.3	1336	3243	1.17	0.69
Pt	--	0.69	0.131	21.45	2042	4803	0.24	0.31
Cu	--	3.94	0.380	8.96	1356	2868	1.15	0.68
Ag	--	4.18	0.234	10.49	1233	2483	1.70	0.83
Al	--	2.04	0.896	2.71	--	--	8.42	0.58

Table 4-4. Estimated absorption coefficients at 1.0 μm for several possible dielectric contaminants.

Material	Absorption Coefficient (cm^{-1})
SiO	< 100
FeO	5200
Fe ₂ O ₃	960
Cu ₂ O	2180
CuO	3490
Ni	2640

the s-polarized component, we employ the 2ω null technique. The amount of second harmonic light generated is measured with a 44-cm calorimeter after passing through a 4-mm aperture at target chamber center to block the residual 1ω . As described in the 1986 annual report,³³ the Type II/Type II frequency-conversion array on the target chamber is tuned to mix the s-polarized light efficiently with the nominally polarized component of the beam. Since the conversion efficiency of the depolarized component is near unity, the energy in the depolarized component is approximately equal to one-half of the 2ω energy generated. Our frequency conversion model³⁴ is used to more accurately estimate the conversion efficiency of the s-polarized component.

As shown in Fig. 4-30, approximately 2% of the energy in beamline 3 was s-polarized at the onset of this experiment. The measurement was made with 4-kJ temporally Gaussian pulses with a full width at half maximum of 600 ps, as required for a ride-along experiment in the two-beam chamber. After reversing the spatial filter lenses, the level of depolarization decreased to approximately 0.8%. The variations observed in the depolarization level on successive shots are consistent with those of previous Nova depolarization measurements; the level of depolarization is larger on the first shot of the day than on shots that follow. We estimate that these measurements are absolutely accurate to within about 20%, due to uncertainties in the calorimeter energy measurement, transmission of the target chamber optics and calculation of the depolarization from the frequency conversion model. To correct the measured 2ω energy for the transmission of the target chamber optics, the 2ω transmission was inferred from a measurement of the 1ω transmission measured in April 1989. The uncertainties in this correction contribute to the error in the calculation of the absolute magnitude of depolarization, but not the relative levels of depolarization before and after the lens reversal.

The level of depolarization with the lenses reversed compares well with the measured 0.4%–0.6% contribution of the 46-cm amplifiers³⁵ and the calculated 0.2% depolarization induced by the spatial filters. However, no information is available as to how these two contributions to the depolarization of the output beam add vectorially.

Although the depolarization substantially decreased, the lens reversal did not have a noticeable effect on the third harmonic conversion efficiency on BL-3 at the normal 1ω operating intensity of 1.6 GW/cm^2 . The 3ω system conversion efficiency, defined as the energy measured in the target chamber 44 cm calorimeter divided by the 1ω energy measured in the output sensor, remained at approximately 40%. Code calculations predict that the decrease in depolarization should result in a 7% increase in 3ω generation efficiency at normal operating intensities and a 15% increase at the maximum operating intensity of 2.5 GW/cm^2 . Although our calculations indicated that reversing the spatial filter lenses introduces a small amount of aberration on the output beam, we did not observe any change in the far field spatial profile measured in the output sensor.

Author: R. Ehrlich

4.2.7. Pulse Shaping Capability

Introduction . In this section, we describe the techniques developed to produce temporally shaped pulses on Nova—both the computer model which calculates the pulse shape required at the input of the amplifier system to produce the desired pulse on target and the engineering hardware that creates this master oscillator pulse shape. We will also review progress made in the synthesis of nonuniform transmission lines, which are used to generate shaped electrical pulses that in turn produce shaped optical pulses. Lastly, we discuss results of experiments which test our ability to calculate and produce a desired master oscillator pulse and we show pulse shapes successfully generated for Nova ICF experiments.

Optical Pulse Shape Modeling . We have constructed a computer model which we use to calculate the pulse distortion produced in the Nova amplifier chains and harmonic conversion system. The model is based on the system diagram shown in Fig. 4-31. We use the code to determine the pulse shape at the amplifier system input that produces the desired pulse at the target. Calculation of the input pulse shape is a two-step process starting at the target and ending at the input to the amplifier chain. First, we calculate the $1.05\text{-}\mu\text{m}$ pulse shape required at the frequency converters to obtain the desired $0.35\text{-}\mu\text{m}$ pulse shape on target. Second, we use the resulting $1.05\text{-}\mu\text{m}$ pulse as input for a calculation

Fig. 4-30. Depolarization measurements before and after reversing final two output spatial filter lenses on BL-3.

Fig. 4-31. Shown is a schematic drawing of a Nova laser amplifier beamline. The gray boxes represent different amplifier stages; the numbers labeling the amplifiers denotes the clear aperture in centimeters.

that accounts for the 1.05- μm pulse distortion caused by saturation of the glass gain media in the amplifiers. Although the third harmonic conversion process is well understood,³⁶⁻³⁸ imperfections in the beam and in the crystal alignment reduce and alter the power dependence. Sources of the degraded conversion include a small nonuniform depolarization of the 1.05- μm beam ($\sim 2-3\%$ of the total beam energy) and an associated intensity-dependent rotation of the polarization ellipse³⁹ (see also section 4.2.2.), a nonuniform intensity distribution in the beam, and deviations of the local k-vector of the beam from the phase matching angle of the crystal. Averaged over the beam, these angular deviations are believed to be 50-100 microradians. Detailed modeling of these effects requires a knowledge of the properties of each of the ten beams and is thus difficult to do in detail. Hence we have applied a statistical model of the beam intensity, an average angular error from phase matching, and a polarization distribution derived from measurements made on a single beamline to construct a model of the third harmonic conversion on an average Nova beam. The model fits measured data.⁴⁰

The third harmonic energy delivered to the target is degraded by losses incurred between the conversion array and the target. These losses arise from the obscuration caused by the 'eggcrate' crystal mounting assembly, a central obscuration that prevents unconverted light from hitting the target, and the surface losses of the KDP array, lens and debris shield. Total losses are nominally 12% but can be larger depending on the condition of the surfaces. Measured conversion efficiencies based on energy delivered to target chamber center are the ratio of third-harmonic energy measured by calorimeters in the target chamber to the incident 1.05- μm energy measured by the beamline output diagnostic modules. Since the pulses used during the measurement of the energy ratios are nominally temporally square, measuring this ratio at various drive energies allows us to determine the intensity dependence of the harmonic conversion. We then compensate for the obscuration and surface losses by scaling the model of the frequency conversion to agree with the measured ratios.

We used this empirically adjusted model to determine the 1.05- μm pulse shape required at the conversion crystals to produce the desired third harmonic pulse on target. Using the

transmissions of the individual beams, we calculate the 1.05- μm energy required from each beam to produce the specified 0.35- μm energy on target. The differences in optics transmissions and conversion efficiency between the ten beamlines require adjustment of the 1.05- μm drive to produce the same output 3ω energies. This results in differences in pulse shape because it produces different levels of amplifier saturation. This is presently the largest cause of instantaneous power imbalance among the beamlines.

To account for pulse distortion from saturation in the amplifier chain, we calculate the dependence of the amplifier gain on the amount of energy extracted from the laser media using the Frantz-Nodvik solution⁴¹ for a saturating amplifier, i.e.;

$$F_{\text{out}} = F_{\text{sat}} \ln \left[G_0 \left(e^{F_{\text{in}}/F_{\text{sat}}} + 1 \right) \right] \quad (1)$$

where

F_{out} = output fluence
 F_{in} = input fluence
 F_{sat} = saturation fluence
 G_0 = small signal gain.

This equation neglects passive losses and integrates the photon-transport equation in time. A two-level saturation fluence model was chosen to match measurements made on Nova laser glass with saturation fluences of 5.11 and 1.37 J/cm² partitioned with weights of 0.91 and 0.09 respectively.⁴²

The small signal gains in the model we obtain from measurements of the gain of a low-energy pulse amplified by a single amplifier stage in individual Nova beamlines. We observe differences in gain of the same stage in different beamlines of up to $\pm 10\%$. For our input pulse shape modeling, we use measured values of small signal gain and passive losses, with each arm having its own measured parameters.

To account for the distributed passive losses in the amplifier chain, we break each Nova amplifier stage into a series of lossless amplifiers separated by transmission losses. Some surface losses caused by birefringence in the laser amplifier components and scattering from the Brewster's angle surfaces do occur and to account for these, we take average surface losses

of 0.3% per surface before and after each laser disk. The 0.3% was chosen to match the measured average transmission of the Nova beamlines. Bulk losses present in the amplifier glass are assumed to be approximately 0.0011-cm^{-1} , consistent with off-line measurements made on Nova amplifier disks.⁴³ Bulk losses are approximated by breaking the amplification process into small steps and taking a loss between the steps. Steps are chosen small enough to be insensitive to additional steps. To account for transmission differences among the beams, we adjust the transmission of spatial filter lenses and other optics between amplifier stages (see section 4.2.2.) to match measured beamline transmissions.

To model the effect of saturation on the temporal pulse shape, we must introduce time dependence into our Frantz-Nodvik model. We find that a small modification to the time independent solution works extremely well. We write the fluence at an arbitrary time t in the pulse as,

$$F(t) = \int_0^t I(t') dt' \quad (2)$$

$F(t)$ is proportional to the stored energy extracted from the amplifier per square centimeter of output beam area up to a time t and provides the input to a Frantz-Nodvik calculation of the pulse distortion produced by the amplifier chain. In order to calculate the required input pulse shape, we first use Eq. (2) to calculate $F_{\text{out}}(t)$ from the desired $1.05\text{-}\mu\text{m}$ output intensity $I_{\text{out}}(t)$. $F_{\text{out}}(t)$ is used as input to our Frantz-Nodvik model of a Nova beamline that calculates the F_{in} required to produce F_{out} for each time interval in the pulse. The model solves equation (1) for each amplifier segment and compensates for the transmission losses described above to find $F_{\text{in}}(t)$. The derivative of $F_{\text{in}}(t)$ then yields the required input pulse shape.

Spatial variations in beam fluence, which we account for in the frequency conversion modeling, have little effect on these calculations. Because the Nova beam is highly uniform, we obtain the required accuracy using a spatially averaged fluence. The beam fluence distribution is nominally a "flat top" spatially, which can be approximated by a 50th-power supergaussian with a 10% rms deviation in the flat region. We are confident from our

propagation calculations and actual measurements of pinhole closure that under normal circumstances the pinholes on Nova show no change in transmission over the few nanosecond duration of our pulses. Similarly, at the pulse lengths at which we operate, other gain processes are not important and we achieve a good approximation with a simple Frantz-Nodvik calculation.

Pulse Shaping Apparatus . Two master oscillators, their associated electro-optic modulators and a system for compressing pulses combine to provide a wide range of pulsewidths and shapes available from the Master Oscillator Room (MOR) for input into the Nova amplifier system. Pulse durations routinely available are from 10 ps to 10 ns. The shortest pulses available are produced by compressing 80-ps Gaussian pulses from the mode-locked oscillator, and are approximately sech^2 in shape. The mode-locked oscillator is continuously tunable in pulse length to produce Gaussian pulses from 80 to 950 ps in duration. Pulses from 500 ps to 10 ns in duration are generated with arbitrary shape, constrained by the 150-ps response time of the electro-optic modulators and by our ability to generate large amounts of contrast with the modulators. Compressed pulses, Gaussian pulses and shaped pulses can be combined to make complex pulse shapes for target irradiation, or separated to propagate through different amplifier chains to provide x-ray backlighting sources or, alternatively, irradiation pulses for separate experiments. **Figure 4-32** is a schematic drawing of this oscillator system showing the oscillators, electro-optic modulators, pulse compressor and beam combining optics.

At the heart of the master oscillator system is an actively mode-locked and Q-switched (AMQ) Nd:YLF laser,⁴⁴ that generates a train of Gaussian pulses. This oscillator is continuously tunable in pulse duration from 80 to 950 ps by varying the depth of modulation of an acousto-optic modulator and by choice of mode-selecting intercavity etalons. Single pulses are selected from the train using Pockels cells. These pulses are either delivered to the main amplifiers, temporally compressed,⁴⁵ or used to drive fast photoconductive switches⁴⁶ in the electro-optic pulse shaper.

A single longitudinal mode (SLM), Q-switched Nd:YLF laser that emits a 300-ns pulse provides the input pulse to electro-optic

Fig. 4-32. Schematic of the optical pulse shaping apparatus used on Nova. A square 4-kV wave is reflected and shaped by the nonuniform transmission line, then routed to the pockels cell. The shaped optical pulse is windowed in time by the 'slicer' pockels cell, which improves the contrast while also slicing away unwanted residuals.

modulators in the pulse shaping system. The modulators consist of Pockels cells driven by a shaped high-voltage pulse. The shaped voltage pulse is produced by a novel high-voltage pulse generator^{47, 48} that uses silicon photoconductive switches illuminated by amplified pulses from the AMQ oscillator, and a nonuniform impedance transmission line to generate programmable pulse shapes. Programmed impedance versus distance in the transmission line determines the generated output voltage versus time which we will now describe.

Nonuniform Transmission Line Synthesis.

Until recently, the nonuniform line used on Nova was a conductive strip suspended within a conducting 'U' shaped channel by a series of adjustable dielectric micrometer extensions.^{47,49,50} Known as the VITL, we had problems with its hysteresis, tedious trial and error adjustment, and an inability to make pulses required for some recent experiments. We discovered during the past year that it was possible to make pulse shaping lines that worked better than the VITL, by hand cutting variable width microstrip transmission lines onto copper-clad fiberglass printed circuit board. We have done this to great success, and have repeatably generated temporally square pulses, pulses with a leading foot, and picket fence pulses. One key factor in making these lines is the rapid feedback available from our diagnostic equipment. A streak camera with a CCD camera was configured to capture a sample of the optical pulse from the shaping apparatus, and to display the result every 3–5 seconds. This has enabled the technician to cut away the copper from an initially uniform transmission line in a one direction process, while immediately observing the pulse. For most of the required pulse shapes, it has taken 1–2 hours to manually cut a pulse shaping line.

This year we also demonstrated the feasibility of a-priori generation of pulse shaping lines. We extended previous work on an inverse scattering technique known as layer peeling, which was applied to the synthesis of nonuniform initially charged or reflecting transmission lines.⁴⁸ Numerically, we change the impedance of the line at each point, starting at the beginning of the line and working in one direction, in order to reflect the required amount of the input wave to satisfy the output wave requirement. This is identical to what the technician does when he manually cuts the line,

but the difference is that the optimum solution can be found. On the computer, the problem is repetitively solved until the maximum amplitude voltage pulse of the desired shape is obtained, utilizing as much of the input pulse energy as possible. An advantage of manually cutting the line though, is that when the technician cuts the line, he directly observes the output optical pulse, taking into account the transfer functions of the Pockels cell, photoconductive switch, and cables. The layer peeling method developed so far has been shown to generate the precise electrical pulse asked for (Fig. 4-33), but the other system components have not yet been characterized. We developed several computer codes: one that performs the layer peeling/inverse scattering; a one-dimensional finite difference code⁵¹ for nonuniform transmission line analysis; and a code that generates the vectors required for manufacturing the line (Fig. 4-34). Work is currently underway to measure the necessary transfer functions, so that numerically synthesized pulse shaping lines can be used on Nova.

Experimental Results . The results to date validate our model of the Nova amplifier chain for temporally-square pulses, picket fence pulses, and high-contrast continuous pulses. **Figure 4-35** compares the pulse shapes calculated by our model with the measured shapes from a shot designed to extract a large amount of energy from an amplifier chain in a square temporal pulse. **Figure 4-35a** shows the required input pulse shape calculated by our model for a single-beam energy of 12.0 kJ at 1.05 μm incident on its frequency conversion array in a 2.4-ns temporally square pulse. To test the internal consistency of our model, we used this measured shape as input to our amplifier chain gain model and calculate the 1.05- μm pulse shape at the output of the beamline (**Fig. 4-35b**). The 0.35- μm pulse shape calculated also agrees well with the shape measured by the third harmonic streak camera (**Fig. 4-35c**). Thus, the modeling approach described at the beginning of this section agrees well with experiment.

The limits of the laser chain model and the pulse shaping apparatus were explored further. We successfully generated picket-fence pulses using two different methods, and generated pulses with a peak-to-foot contrast exceeding 100:1. **Figure 4-36a** shows a "picket fence" pulse which consists of successive impulses or "pickets"

QT-KC-U-7693-5

Fig. 4-33. Comparison between the desired electrical pulse, to which we designed a nonuniform pulse shaping transmission line, and the measured pulse from the line. The two overlay nearly exactly. The unwanted pulse after 3.5 ns is an artifact that is ultimately sliced away.

Fig. 4-34. Photograph of the nonuniform stripline prior to installing the cover ground plane. It was fabricated in .125 RT-Duroid microwave circuit board with type 'C' connectors.

TCF-AS22-U-11498-9

Fig. 4-35. Comparisons that illustrate the predictive capability of our pulse shaping model: a) The calculated input pulse shape predicted to create a square temporal output pulse compared to the shaped pulse measured in the master oscillator room. b) The output 1 ω pulse shape calculated from the measured input pulse compared to the 1 ω pulse measured at the output of the amplifier chain. c) The calculated third harmonic pulse compared to the measured third harmonic pulse.

Fig. 4-36. Typical shaped output pulses at the third harmonic ($0.35\ \mu\text{m}$) generated for ICF experiments on the Nova laser system include a) picket fence and b) complex temporal shape pulses.

before a longer shaped pulse. This pulse was generated by selecting a single pulse from the AMQ oscillator, and with passive optics, then splitting, delaying, and recombining two or more of the resulting pulses onto one beamline. Amplitude and timing of each picket can be independently adjusted. This series of pickets was then combined with the pulse from the pulse shaping system. Since the pulse shaping system is triggered by a short pulse from the same AMQ oscillator train, and the electrical circuitry exhibits negligible jitter with respect to its triggering pulse, the two components of the picket-fence pulse are precisely synchronized. Relative timing is set with optical delays.

A similar pulse has been generated by the electro-optic pulse shaper by using the nonuniform transmission line just described. This system simplifies alignment and increases reproducibility. Unlike pickets derived from a Gaussian pulse, the shape of individual pickets is variable, but their pulse length is limited to greater than 300 ps duration by the temporal response of the electro-optic modulators.

Figure 4-36b shows a stepped continuous pulse consisting of a low level "foot" followed by a programmed rise and a high power main pulse. The low power "foot" exhibits lower harmonic conversion efficiency because of its low intensity, but is similar to pickets in its effect on the target. As a demonstration, we generated a 3.1-kJ, four-step, third harmonic pulse of this type with maximum contrast of 125:1 (**Fig. 4-37**).

Summary. We have demonstrated the ability to temporally shape high energy pulses from a multibeam, frequency-converted, Nd:glass laser amplifier system. We achieve these results by using a flexible optical pulse generator that incorporates two master oscillators and an adjustable electro-optic pulse shaping system with a time response of 150 ps. The nonuniform line synthesis methods have been demonstrated and will soon be used to automate the electrical pulse shaping system. This apparatus is augmented by a computer model that includes the nonlinear effects of gain saturation in the laser amplifiers and the intensity dependence of the harmonic conversion and calculates the pulse required at the input to the amplifier system to obtain a desired frequency converted pulse on target. Using the computer model to aid in setting up the pulse generating hardware, we routinely deliver 20 to 40 kJ third harmonic pulses to ICF targets which are either

temporally square or complex temporal shapes designed to reduce shock heating in the target. We are currently using these complex shaped pulses in experiments designed to demonstrate the benefits of pulse shaping on target performance.

Authors: S. C. Burkhart and J. K. Lawson

Major Contributors: R. B. Wilcox, T. L. Weiland, and E. H. Padilla

4.2.8. High Energy Capabilities of the Nova Laser at 1.05 μm and 0.35 μm

During 1988 and early 1989 in parallel with testing the prototype beamline, we worked to rebuild the nine remaining Nova beamlines. The rebuild included installing Pt-particle-free laser glass in the 31.5-cm and 46.0-cm amplifiers,⁵² substituting lenses made from fused silica for the BK-7 lenses at the output of the last two spatial filters in each beamline, and upgrading all beams to improve their near-field uniformity.

After completing the rebuild of the amplifier beamlines in March 1989 and reviewing the results from the prototype beamline, we were confident that Nova could exceed its design goals of 80 to 120 kJ of output energy in a 3-ns pulse and 80 to 120 TW of output power in a 0.1-ns pulse. To demonstrate this, we designed a set of experiments to progressively increase the 1- μm laser output energy and power from previous operating levels (≈ 50 kJ/80 TW) to more than 120 kJ for pulse durations of 2.5 ns, to more than 80 kJ for 1-ns pulses, and to more than 120 TW for 0.1-ns pulses. These pulse durations are of interest because 2.5 ns is close to the duration at which Nova's high energy performance goal was specified, and output is limited by damage to optics: 1 ns is where energy and power output are of most interest for ICF experiments and 0.1 ns is where output power has reached its maximum and is limited by beam breakup from self-focusing instabilities.

During this series, we measured the 1.05- μm energy with our standard full-aperture (80 cm) calorimeters and the temporal pulse shape with 20-ps-resolution LLNL streak cameras that obtained their input signals from the beamline output sensors. The calorimeters have an absolute accuracy of $\pm 5\%$ and are repeatable to within approximately 1%. Their calibrations can be traced to a reference calorimeter (of 5-cm diam) in our calibration laboratory that was

Fig. 4-37. Measured four-level continuous pulse at $0.35\ \mu\text{m}$ produced on one beamline of Nova as a demonstration of Nova's pulse shaping capabilities. The energy in the pulse was 3.1 kJ.

Table 4-5. The average beamline energy at 1.05 μm , the standard deviation, and the spread in energy among Nova's ten beamlines during shots to demonstrate high energy and power.

Shot number	Ten-arm average (kJ)	σ (%)	Spread (%)
(2.5 ns)			
1940307	6.98	10.2	14.8
1940408	7.71	5.5	9.3
1940b414	8.08	1.9	3.4
1940418	9.81	2.5	4.1
1940503	9.78	1.5	2.7
1940508	12.47	1.7	2.7
1941105	11.96	1.2	2.0
(1 ns)			
1940605	5.86	2.6	5.0
1940709	7.07	2.1	3.5
1941205	8.20	3.5	6.9
1941213	7.77	4.0	6.6
1941307	8.20	1.0	1.8
(0.1 ns)			
1941310	1.04	6.9	12.4
1941313	0.97	5.7	9.7
1941405	1.04	3.8	6.2
1941411	1.14	2.1	3.4
1941803	1.33	2.1	3.6

optically calibrated at the National Institute of Standards and Technology with an absolute accuracy of $\pm 3\%$. For each shot in the series, **Table 4-5** lists the pulse duration, the average energy per beam, the standard deviation in energy among the beams in percent, and the maximum deviation from the average energy in percent.

During shots at each pulse duration, we continually adjusted the $\lambda/2$ waveplates that set the input energy to each beam to achieve the best possible beam energy balance. At the end of the shots at each pulse duration, the differences in energy among the beams were substantially reduced; the standard deviation approached 1% for the best shots.

Figure 4-38 compares the experimental data taken in this series with the projected maximum energies allowable from the rebuilt amplifier system.^{53,54} The output energies obtained confirmed that the projected energies are safe operating values and may be conservative. Maximum output energies and powers measured were 125 kJ in a 2.5-ns duration pulse, 82 kJ in a 1-ns pulse, and 125 TW in a 0.1-ns pulse. No observable damage to optics occurred on any shot in this series.

Figure 4-39 shows the instantaneous power histories of each of the ten beams from a 124-kJ shot, derived from the temporal shape of the output pulse measured with our streak cameras and the beam energy measured by the calorimeters. The indicated instantaneous power imbalance of the ten beams is less than $\pm 10\%$ peak to peak throughout most of the pulse. The accuracy of this measurement is limited by intensity fluctuations resulting from coupling the beams into the streak cameras and by noise from the streak cameras. Differences in pulse shape at this energy of up to $\pm 10\%$ are expected to occur because of differences in gain and transmission among the individual beamlines that influence their saturation characteristics.

Later in the year to demonstrate that Nova can reliably deliver increased energy on target, we did a series of shots whose goal was to provide more than 40 kJ of 0.35- μm light into the target chamber for target irradiation. In this series, the Nova laser provided 100 kJ of 1.05- μm output in a 2-ns square temporal pulse to the frequency conversion arrays on the ten-beam

target chamber. The laser pulse was frequency chirped to a bandwidth of approximately 15 GHz by an electro-optic modulator in the master oscillator room to suppress stimulated Brillouin scattering in the focusing lenses on the target chamber which could possibly damage them.⁵⁵ The frequency conversion system performed as expected and provided 43% of the 1.05- μm output as 0.35- μm radiation in the target chamber. **Table 4-6** summarizes the results of the five high energy shots done during the series, giving for each shot the 1.05- μm energy incident on the frequency conversion arrays, the standard deviation in 1.05- μm energy among the ten beams, the 0.35- μm energy in the target chamber, and the standard deviation in the 0.35- μm energy. Two shots were focused through a 4-mm aperture in the center of the target chamber and the 0.35- μm energy measured in calorimeters mounted on the opposing target chamber ports; the other three shots were onto targets and no direct measurement of the 0.35- μm energy was possible. It was estimated from the measured 1.05- μm laser output and the conversion efficiencies calculated from the shots into the calorimeters. On the highest energy shot, the calorimeters measured 45.3 kJ of 0.35- μm light.

In order to maximize the 0.35- μm output with minimum threat of damage to the laser-amplifier components, we attempted to balance the 1.05- μm energy output from Nova's ten beams. After making minor adjustments for energy imbalance observed on the first shot, we obtained a standard deviation (σ) from the average energy among the ten beams at 1.05 μm of 1% on three of the four following shots. The higher σ on the fourth shot resulted from anomalously low output from beam 1, probably caused by misalignment of a spatial filter pinhole. Careful inspection to the most threatened optics in the laser and the target chamber after the series revealed no damage, demonstrating that target campaigns at greater than 40 kJ of 0.35- μm energy are possible with Nova.

Author: D. R. Speck

Major Contributors: C. Bibeau, R. B Ehrlich, C. W. Laumann, D. T. Kyrazis, M. H. Henesian, J. K. Lawson, P. J. Wegner, and T. L. Weiland

MB-23-U-8884-03

Fig. 4-38. Estimates of the maximum 1.05-nm energy available from the Nova amplifier system with acceptable risk of damage to components. Data shown were taken after the amplifier-system upgrade. No damage to laser components occurred.

MB-23-U-8885-01

Fig. 4-39. Overlay of the measured instantaneous power from the ten beams on 124-kJ shot. The standard deviation from energy balance was 1.7%; power imbalance throughout most of the pulse was less than $\pm 10\%$ (Shot No. 1940508).

Table 4-6. Summary of shot series to provide >40 kJ of 0.35- μm energy on Nova targets.

Shot #	1.05- μm Energy (kJ)	$\sigma(1\omega)$ (%)	0.35- μm Energy (kJ)	$\sigma(3\omega)$ (%)
19103103	105.7	2.8	45.3	8.5
19103109	98.6	0.95	42.6	9.3
19103106 ¹	102.6	0.90	(44)	--
19110103 ¹	103.4	4.8 ²	(44)	--
19110107 ¹	100.2	1.1	(43)	--

¹Target shots, no 0.35 μm calorimetry

²Beam 1 anomalous

4.3. Precision Nova

4.3.1. Introduction

During 1989 we took the initial steps on the Precision Nova project to improve our ability to do more precise experiments. Although this is a multi-year, multifaceted project, the major thrust is to understand the causes of pointing errors and laser power imbalance on targets and to devise solutions. The following articles deal with quantifying the accuracy with which each of the ten Nova beams are pointed, beam steering by flashlamp pump induced aberrations, prototyping of an improved incident beam diagnostic system, and optical streak camera fidelity. The goals of this project are to provide positioning of each of the focused beams to better than 25 microns which is equivalent to 8 microradians uncertainty in pointing. The optimum use of shaped, third harmonic pulses to more efficiently compress targets, requires the power in each beam be within 5 percent of the average power. In the prototype diagnostic, the third harmonic power is assessed by whole beam calorimetry in combination with full beam aperture streak camera sampling of the intensity versus time of each beam.

Understanding the causes of the pointing and power balance errors will tell us where to put our resources to reduce these experimental variations that contribute to the range of performance that is observed in targets. More consistent target performance translates to shorter target campaigns to discern the influence of target design variables.

Author: R. Smith

4.3.2. Beam Targeting Accuracy

We have begun a multi-year effort to improve the targeting accuracy of Nova. As a first step, we conducted a series of experiments to determine the current pointing accuracy, to look for changes since an earlier series of pointing tests⁵⁶ in 1986, and to evaluate different diagnostics and beam configurations for the pointing tests. We also hoped to find any conspicuous errors in the alignment procedure which might quickly lead to an improved pointing accuracy.

The pointing tests consisted of aligning special targets with the normal target alignment procedure, irradiating with all ten beams,

recording the focal spots with x-ray cameras, and analyzing the data to determine the actual shot position of each beam compared to the position at which it was aimed. We used the same 50- μm -thick gold disk targets as used previously, with ten 100- μm -diam holes placed symmetrically on two circles. The target was placed perpendicular to the poles of the chamber, so that the five beams from the east were aimed at the outer holes and those from the west at the inner holes. Separate x-ray cameras imaged the two sides of the targets. **Figure 4-40** shows the target geometry and the general placement of diagnostics.

The purpose of the holes in the targets was to provide a spatial fiducial on the x-ray images which would allow the aim point to be calculated. If the beams were properly aligned, they would overlap the holes (typical x-ray image size was greater than 200 μm), and some of the x-rays would leak through to the opposite side to be imaged along with the beams spots on that side. Although the Nova beams are not aligned directly to the target, but rather to the alignment reticle which is subsequently replaced by the target, the diameters of the two circles of holes and their placement on the circles were designed to coincide with crossing points of lines on the alignment reticle.

The diagnostics included the two 8 $\times$ Kirkpatrick-Baez (KB) x-ray microscopes used previously, as well as a single axial pinhole camera (APC) imaging the east beams. The 5- μm pinhole used in the APC potentially provided 2 to 3 $\times$ better on-axis resolution than available from the 8 $\times$ cameras, and it produced no off-axis distortion. Both types of cameras were filtered to image photon energies between about 1.0 and 1.5 keV.

The earlier series used several laser pulse lengths to distinguish the effect of beam steering by plasma formation at the target. They found that their shortest pulse length of 200 ps gave the best results. We chose the 200-ps pulse length for this series to minimize this effect and to allow direct comparison to the best of the previous data. The nominal 3ω energy for each beam was 300 J.

We also investigated defocusing the beam somewhat to reduce hot-spot formation at the target and to make data reduction easier. Although the first useful shot of the series was at best focus, like all the earlier data, defocusing to 1 mm beyond best focus proved to be

Fig. 4-40. schematic of the pointing target geometry and its placement relative to the diagnostics

advantageous, and the remaining shots were taken that way. In all, the 8× KB cameras produced six new data sets for the east beams (even numbered beamlines) and four for the west beams. The APC produced two useful data sets for the east beams. Although the APC did produce the expected improvement in spatial resolution, its data was not significantly different than that from the corresponding 8× KB camera, and there were not enough data to draw major conclusions. Consequently, only the 8× data was used in the following analysis.

Data reduction required determining the actual shot positions on the film, and then translating these film coordinates to those of the central plane of the target chamber to compare to the known aim points in that coordinate system. The digitized film data was analyzed on a Lasers Experiments Analysis Facility (LEAF) VAX computer using a high resolution Ramtek graphics screen and a resident graphic package. The positions of the beams and the fiducials on the film were determined by visual estimation of their image centers using contours of equal film density.

The transformation from film coordinates to target chamber center coordinates is a linear transformation involving translations, scaling, and, in the case of the 8× data, rotations. To avoid the potential errors of determining this transformation directly from the known orientation, position, and magnification of the cameras, the positions of the fiducial holes were measured when the targets were fabricated, and the transformation coefficients were calculated from the measured fiducial positions on the film. The linear transformations are given by:

$$A_1 X_f + A_2 Y_f + A_3 = X_T \quad (1)$$

$$B_1 X_f + B_2 Y_f + B_3 = Y_T \quad (2)$$

where X_f and Y_f are film coordinates of one of the fiducials, and X_T and Y_T are target-chamber-center coordinates of the same fiducial obtained when the targets were fabricated. If all five of the fiducials were visible, there were five sets of (X_f, Y_f) and (X_T, Y_T) for each film image and five sets of equations like Eqs. 1 and 2, but only three unknowns per equation, so the transformation coefficients A_i and B_i were over specified by the data. Consequently, a least squares fit gave not

only the A_i and B_i values, but an estimate of the combined measurement errors of the fiducial positions on film and target. Given the A_i and B_i values, the film positions of the beams were translated to target-chamber-center coordinates and compared to their known aim positions. The estimated error in determining the centers of the film images of the beams is $\pm 30 \mu\text{m}$, and the estimates of the fiducial position errors from the least-squares fit is $\pm 15 \mu\text{m}$, giving a total estimated error for the data reduction of $34 \mu\text{m}$, RMS.

All the pointing shot data from the 8× KB cameras is shown in Fig. 4-41, including the two comparable shots from the previous series. This figure shows an east view of the target plane at the center of the chamber, as seen by the target plane imager used to align targets. The data from the previous shot series are indicated by the triangles.

The first conclusion to be drawn from the data of Fig. 4-41 is that there has been no significant overall change in the Nova pointing accuracy since the earlier pointing shots. Most of the earlier data points are within the groupings of the current data (all except beamlines 4, 6, and 9). This implies that the static quality of the laser disks in the 31.5-cm and 46-cm amplifiers has no influence on pointing accuracy, since all this glass has been changed since the earlier series.

The second conclusion is the overall beam targeting accuracy of Nova. The single number which characterizes targeting accuracy is $120 \mu\text{m}$ RMS, obtained from the deviations between aim point and shot position for all the data in Fig. 4-41. This number is substantially larger than the $75 \mu\text{m}$, RMS, obtained from the earlier pointing series, but we believe this to be the consequence of more data and not a real difference in pointing accuracy.

Two numbers provide better characterization of targeting accuracy, one to characterize the size of the clusters of data points and a second to characterize the offsets of these clusters from the aligned points. Figure 4-42 shows circles whose radii are the RMS offsets of the individual data points from the average position of the cluster for each beam. It also shows the aim points and a line indicating the size of the offset between cluster and aim point. The RMS offset of the clusters from the aim points is $93 \mu\text{m}$, and the RMS deviation of the shot points from the offset positions is $72 \mu\text{m}$.

02-03-1092-2743 pub

4-41. An east view of the target plane showing the aim points for each beam and the positions of all pointing shot data from the KB 8× cameras; the triangles indicate data from a previous shot series in 1986.

02-03-1092-3153 pub

4-42. Offsets of the average positions of the shot data from the aim point for each beamline and circles which indicate the RMS deviations of the shot data from the average offset positions; the dashed lines indicate the 1σ polarization direction for those beam where it differs from the radial direction in the figure.

Additional but less quantitative conclusions can be seen from the cluster shapes and offsets. **Figure 4-41** shows that the shot positions within the clusters are most often linearly rather than randomly distributed, and they are arrayed approximately radially with respect to the center of the target disk. Only five of the offsets are significant with respect to the cluster size, that is the circles do not include the aim points in **Fig. 4-42**, and these five offsets are oriented approximately in an outward radial direction. The 1ω polarization direction is different than the radial direction for only four of the beams (dashed lines in **Fig. 4-42** for beams 3, 4, 7, and 8), but three of these have significant offsets, and it appears that the offsets align better with the radial direction than with the 1ω polarization direction. The meaning of these latter conclusions, if any, is not yet clear and is the subject of ongoing investigations.

Author: J. E. Murray

Major Contributors: C. Bibeau, P. E. Young, K. S. Jancaitis.

4.3.3. Pump-Induced Aberrations in Nova Disk Amplifiers

One of the goals of the Precision Nova project is an improved targeting accuracy, which requires an understanding of the sources of distortion on the Nova output beams. As a first step toward this goal, we measured the flashlamp-induced steering in the Nova 31.5-cm and 46-cm amplifiers to determine its impact on the beam quality and pointing of the Nova output beams. We found two distinct types of beam steering with different causes: a large component near the edges of the disk and a smaller component throughout the disk aperture. This section describes our measurement technique, the experimental results and causes of the two steering components, the implications for Nova, and the implications for future solid-state laser systems.

Measurement Technique. We measured local flashlamp-induced beam steering as a function of time at many spatial locations rather than measuring the phase as a function of position (with interferometry) at many times during the flashlamp pulse. We chose the steering measurement primarily because it was much easier to implement, but also because it provided better time resolved data, which we expected to

be important for diagnosing the sources of induced steering. In fact, either type of measurement gives both phase aberration and steering, since steering equals the gradient of the phase. Integrating our steering measurements over position gives the flashlamp-induced phase front changes.

As shown in **Fig. 4-43**, we established a steering measurement capability by adding a lateral-effect photodiode to our gain measurement facility.⁵⁷ Placing the photodiode in the focal plane of the compound lens made it sensitive only to the angle changes of the beam at the amplifier. In principle, the four currents from the lateral-effect photodiode (+x,-x,+y,-y) gave all the information required to determine vertical and horizontal steering versus time. A change in beam angle at the amplifier in the x-direction $\Delta\alpha_x$, for example, is given by:

$$\Delta\alpha_x = C \frac{\Delta v_x}{\Sigma} \quad (1)$$

where Δv_x is the difference in voltages created by the two x-currents, Σ is the sum of the voltages (proportional to the total power), and C is a calibration constant, that includes the position sensitivity of the lateral-effect photodiode, the load resistors, and the focal length of the lens. In practice we wanted to digitize the data, and since the difference voltages were typically less than 0.1% of the individual voltages, we would have required more digitizing accuracy than was available (10-bit digitizers). Consequently we used analog subtraction to obtain the horizontal and vertical difference signals and digitized only the difference signals. The sum Σ was obtained from a second photodiode which monitored the total beam power, using an integrating sphere to minimize the effect of beam steering. Spatial filters and RG-1000 filters reduced fluorescence and flashlamp light into both detectors. We matched the four load resistors to a few parts in 10^5 at $\sim 10 \text{ K}\Omega$. The measured rise and fall times (10–90%) of the steering detector was $\sim 10 \mu\text{s}$.

A small-angle prism with a known deflection angle provided direct calibration of the steering detector. The prism was inserted into the probe beam and rotated about the beam direction. Plotting $\Delta\alpha_x$ versus $\Delta\alpha_y$ produced a circle for equal sensitivities in the x and y directions, with diameter $2\Delta v/\Sigma$. Knowing the

Fig. 4-43. Schematic of the measurement set-up.

deflection angle of the prism then gave the calibration constant C in Eq. 1. The error in determining C was $\pm 10\%$, and it resulted primarily from air turbulence in the beam path between the probe laser and detector, which typically caused the center of the circle to wander in time. The same problem limited the resolution of our measurements to about $\pm 2 \mu\text{rad}$ over 1.5 ms, although it was substantially better for smaller time intervals.

Experimental Results. Measurements on Nova 46 and 31.5-cm amplifiers, displayed in Figs. 4-44 and 4-45 respectively, showed two distinct types of flashlamp induced steering—a large effect near some of the disk edges and a substantially smaller effect throughout the disk aperture. Edge steering is confined to within a few centimeters of the disk edge, and it rises in amplitude very rapidly near the edge to values an order of magnitude larger than central steering. The largest measured edge steering was for steering in the horizontal direction near the central split of the 46-cm amplifier, as seen in Fig. 4-44a, but smaller edge steerings are also apparent in the vertical direction at the top and bottom of both amplifiers, Figs. 4-44b and 4-45b. Central steering, on the other hand, occurs in both amplifiers, throughout the clear apertures, and its amplitude doesn't change dramatically with position.

Edge Steering. The experimentally observed characteristics of edge steering implied a connection with the disk edge cladding.⁵⁸ First, it was strongly localized near the disk edges and increased rapidly as the edge was approached, and second, the direction of edge steering was always directly toward the edge. Both of these features can be seen in Figs. 4-44a, 4-44b, and 4-45a. These observations suggested that the amplified spontaneous emission (ASE) absorbed by the cladding was causing the edge steering.

Calculations confirmed that the ASE absorbed by the edge cladding would cause steering with the observed characteristics and with the proper magnitudes. Although there is not sufficient time for the disk to be heated significantly by conduction from the heated cladding, the thermal expansion of the cladding itself distorts the disk adjacent to the cladding. This distortion occurs at acoustical transit times, and is therefore quite fast. It is this surface distortion of the disk near the cladding that causes edge steering. Figure 4-46a shows the calculated distortion near the split of a 46-cm

disk using NIKE2D⁵⁹ and exaggerated by a factor of 5000.

The magnitude of the beam steering created by surface distortions at the input and output faces of the disks can be easily calculated by applying Snell's law at each surface. Figure 4-46b shows the geometry, with θ and ϕ the original angles of incidence outside and inside the undistorted disk surfaces, respectively; α and β the changes in surface normal directions caused by the distortion with positive angles as shown; and $\Delta\theta$ and $\Delta\phi$ the changes in transmission angles due to the distorted surfaces. By applying Snell's law at the input and output surfaces and using small angle approximations one finds the angular deviation of the beam:

$$\Delta\theta = (\beta - \alpha) \left(\frac{n \cos\phi}{\cos\theta} - 1 \right) \quad (2)$$

where n is the refractive index of the disk. At Brewster's angle, Eq. 2 becomes:

$$\Delta\theta = (\beta - \alpha)(n^2 - 1) \quad (3)$$

and at normal incidence:

$$\Delta\theta = (\beta - \alpha)(n - 1) \quad (4)$$

The predicted steering amplitudes based simply on these surface deflections agree very well with our measurements. Figures 4-47a and 4-47b compare predictions to measurements near the split of the 46-cm amplifier and at the top and bottom of both amplifiers, respectively. All calculations assumed an exponentially decreasing temperature into the cladding from an ASE fluence of 10 J/cm^2 on a disk cross section perpendicular to the disk surfaces.⁶⁰ The agreement is remarkably good, particularly if the top and bottom measurements for a given amplifier are averaged to account for possible mis-centering by a few millimeters of the disks in their holders.

Calculations also correctly predict no measurable steering at the right or left sides of both amplifiers, because the edge cladding is too far from the edge of the available clear aperture. Measurements at other azimuthal positions show edge steering comparable or less

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Fig. 4-44. Pump-induced steering in a 46-cm amplifier versus a) horizontal position at vertical = 0, and b) vertical position at horizontal = -0.1; positive steering means up for vertical steering or towards the right for horizontal steering (toward the high gain side of the amplifier). Measurements were made with one silver and one black end plate on the amplifier, at a 20 kV charge voltage, and at 1.5 ms after the main bank was triggered.

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Fig. 4.-45. Pump-induced steering in a 31.5-cm amplifier versus a) horizontal position at vertical = 0, and b) vertical position at horizontal = 0, with the same sign conventions and measurement conditions noted for Fig. 4-44.

a)

b)

Fig. 4-46.a) Calculations of the distortion caused by a heated edge cladding, expanded by a factor of 5000; b) the geometry used to calculate the effect of surface deformations on transmitted beam angle.

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Fig. 4-47. Comparison between measurements and calculations for the induced steering a) near the split in a 46-cm amplifier, and b) near the top and bottom of the clear aperture in 31.5-cm and 46-cm amplifiers.

than that at the top and bottom of the amplifiers.

As the ultimate check on this connection between edge steering and edge cladding, we tested a single pair of 46-cm half-disks which had no edge cladding attached at the disk split. The results showed no increase in steering near the split, indicating conclusively that the cladding caused edge steering.

Central Steering. The major experimentally observed features of the central steering are its well defined direction and its very complex time structure. Keeping in mind the $\pm 2 \mu\text{s}$ resolution of these measurements, Figs. 4-44 and 4-45 show that the central steering is almost entirely in the horizontal direction, i.e. in the planes of incidence of the disks. Figure 4-48 shows the time structure of horizontal steering signals at two aperture positions and for different flashlamp conditions for a 31.5-cm amplifier. The major features of the structure always reproduced well from shot to shot, but they changed substantially with both position and illumination condition. The first 3 rows of Fig. 4-48 show steering waveforms with only +side lamps fired, only -side lamps fired, and both sides fired (normal condition). The waveforms show opposite steering directions for the one-sided pumping cases, and substantially larger magnitudes than for "both." Figure 4-48 also shows that the algebraic sum of the +side and -side waveforms, shown in the last row, reproduces the "both" waveform remarkably well in both magnitude and time structure. This agreement, in fact, finally convinced us that most of the time structure in the waveforms was indeed real and not an artifact of the set-up.

We believe that the cause of the central aperture steering is disk bending due to differences in thermal input from the flashlamps on the two disk faces. Since the thermal input is fast relative to the mechanical response time of the disk, this produces mechanical oscillations about the new equilibrium configuration which cause the steering to oscillate in time. Disk bending explains the qualitative features of the steering, and we show analytically that disk bending from thermal stress can account for the observed steering amplitudes. Also, disk bending has previously been shown to cause oscillating, whole-beam aberrations in the 9.4-cm disk amplifier,⁶¹ although this bending resulted from

optically induced mechanical shock rather than thermal stress.

Differing pump fluences on the two faces of a disk cause bending as with a bimetallic strip. Although the surfaces remain parallel in the bent condition, this still causes beam steering because the disks are inclined at an angle to the beam direction, as shown in Fig. 4-49. For the case of parallel surfaces bent to the same radius of curvature R and oriented at Brewster's angle, the angles α and β in Eq. 3 become $\alpha = 0$ and $\beta = \tan^{-1}(\Delta x/R) \approx (d/R)\tan\phi \approx d/nR$, for a disk thickness d , and Eq. 3 becomes:

$$\Delta\theta = \frac{(n^2 - 1)d}{nR} \quad (5)$$

Figure 4-49 and Eq. 5 show that the direction of the steering is away from the center of curvature of the disk surfaces.

This disk bending model agrees with two experimentally observed characteristics. First, the bent disk steers only in the horizontal direction (the disk is perpendicular to the beam in the vertical direction) and second, the observed steering directions correspond to those from the expected bending directions of the disks in the amplifier. This is most easily understood for the one-sided pumping configuration. The disk faces closest to the lamps will be the hotter, causing the disk to bend toward the lamps with a center of curvature away from the lamps. Consequently the disk bending model predicts steering towards the lamps for one-sided pumping. Figure 4-50 shows one-sided pumping data from the -side for the 31.5-cm and 46-cm amplifiers, and it indicates negative steering (in the minus horizontal direction) for both. Pumping from the +side reversed the steering. Both results are as predicted.

The case with both sides pumped is somewhat more complicated because now the difference in pumping between the two faces creates the bending and therefore the steering. However, the measured steering in all cases was at least qualitatively consistent with this explanation based on disk bending.

We calculated the bending that would be expected from uniformly illuminating a disk from one side only, to determine if this bimetallic-like bending could produce the observed steering magnitudes. We used single-

Fig. 4-48. Steering versus time at two different aperture positions of a 31.5-cm amplifier and with various flashlamps energized: "+side" lamps only, "-side" lamps only, and "both" sides; the final row shows the algebraic sum of the "+side" and "-side" waveforms; the vertical calibration factor is the same as for Fig. 4-49.

Fig. 4-49. The geometry of a bent disk, showing the relationship steering direction and center of curvature.

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Fig. 4-50. Measured, time-averaged steering for side pumping of the 31.5-cm and 46-cm amplifiers.

sided pumping from the minus side for comparison with calculations, because it gave steering waveforms that were well enough behaved for a reasonable estimate of the average induced steering. (Too little of the temporal waveforms were available for pumping from the -side or pumping from both sides, as can be seen in Fig. 4-48.) The aperture-averaged steering values from the measured one-sided pumping data of Fig. 4-50 are 15 and 22 μrad , respectively for the 31.5-cm and 46-cm amplifiers.

To calculate the expected steering from the disk bending model, we used an analytic solution⁶² for high-aspect-ratio disks with temperature variation through the thickness only, a situation reasonably close to one-sided pumping in either amplifier. The solution gives the resultant radius of curvature R of the plate as:

$$\frac{1}{R} = -\frac{12\alpha}{d^3} \int_0^d T(z) T\left(z - \frac{d}{2}\right) dz \quad (6)$$

for a temperature distribution into the disk $T(z)$ with $z = 0$ at the surface, a disk thickness d , and an expansion coefficient α .

We estimated $T(z)$ in Eq. 6 for one-sided pumping using a variation of our flashlamp model.⁶³ This model assumes that all the pump photons reach the disks and that all pass through the disks only once, which means a 100% transport efficiency from flashlamps to disks and no multipass pumping. We used the actual lamp and circuit parameters of these tests, the measured transmission of Ce-doped flashlamps, and the measured absorption of 4.3 cm thick LG-750 laser glass⁶⁴ disks with a doping density of $2 \times 10^{20}/\text{cm}^3$.

The flashlamp model also required a nonradiative decay rate $1/\tau_{nr}$ to calculate $T(z)$, which we determined from

$$\frac{1}{\tau_{nr}} = \frac{1}{\tau_f} - \frac{1-f\beta_{RZ}}{\tau_r} \quad (7)$$

using the observed fluorescence lifetimes τ_f of the disks, the radiative lifetime τ_r and R to Z branching ratio β_{RZ} for LG-750, and the reabsorbed fraction f for the R-to-Z fluorescence at 880 nm. Using measured values of $\tau_f = 436$ and 370 ms for the 31.5 and 46-cm disks, respectively, $\tau_r = 409$ ms⁶⁵, $\beta_{RZ} = 0.39$ ⁶⁴, and assuming $f = .90$,

Eq. 7 gives τ_{nr} values of 1718 and 991 ms for our 31.5 and 46-cm disks.

Figure 4-51 shows the calculated $T(z)$ values for the two amplifiers. Those for the 31.5-cm amplifier are somewhat larger, because its pumping fluence is somewhat higher. Computing the integral in Eq. 6 then gives $R = 5.0$ and 3.9 km, respectively, for the 4.3-cm-thick 46- and 31.5-cm disks. Substituting into Eq. 5, gave predicted steerings of 15 and 19 μrad , respectively, compared to the measured values of 15 and 22 μrad .

These calculated and measured steering values are much closer than we expected, given the assumptions of the calculation. Nevertheless, the agreement demonstrated that disk bending can explain the observed steering magnitudes.

Implications for Nova—Edge Steering. The major influence of flashlamp induced steering on Nova is edge steering near the split in the 46-cm disks, where the clear aperture extends to within a few millimeters of the edge cladding. In fact, the central few centimeters of the Nova output beam are blocked intentionally because that portion of the beam does not harmonically convert. This loss of beam area near the split, that amounts to ~6% of the total (excluding the 13 mm-wide area lost to the disk holders), is caused by edge steering. Otherwise the effect of edge steering on Nova is negligible because there is adequate distance between the cladding and the beam edge. Our measurements show that ~25 mm constitutes an “adequate” distance for the Nova 43-mm-thick disks. Because of the octagonal shape of the Nova disks, the beam is close to the cladding at only a few places, and measurements at those points show that the edge steering is less than 10 μrad per amplifier at radial positions out to 95% of the clear aperture, i.e., 95% of 31.5 or 46 cm.

We also measured the gain for the pair of disk-halves without cladding at the BA split to determine how much, if any, of the 6% loss in beam area could be recovered by removing this cladding from all the Nova 46-cm disks. We placed loose absorbing plates at the split to block ASE paths between the disk-halves. We found that the average gain coefficient was 6% lower than expected from the same disks with cladding, but in reasonable agreement with predictions.⁶⁶ Without the absorbing plates, we found that the gain was lower still, as expected, but that it fluctuated substantially with bank

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Fig. 4-51. Calculated temperature distributions into the disk faces used to calculate the expected steering from one-sided pumping.

energy and from shot-to-shot in a manner we do not understand.

Removing the cladding at the BA split would therefore increase the Nova output beam area by ~6%, but it would also decrease the gain coefficient and stored energy in the 46-cm stage by ~6%. At long pulse lengths where Nova is currently limited by damage to optics beyond the 46-cm stage, this would mean 6% more energy output, assuming enough drive to make up the loss in 46-cm stage gain. However, at short pulse lengths, Nova is ΔB limited in the 46-cm stage, so the increase would be less, because ΔB accumulates in amplifiers inversely with the gain coefficient. If damage thresholds were improved enough that the 46-cm stage could be heavily saturated, there would be little or no improvement because the decrease in stored energy would be about the same as the increase in beam area.

Central Steering. Predicting the influence of central steering on the Nova output beam required measuring the steering by amplifiers in different positions within a stage. Intra-stage differences occur because the pumping on the outside faces of the disks can be significantly affected by adjacent amplifiers. Since induced steering comes from pumping nonuniformities, it follows that induced steering would also be influenced by intra-stage position.

We simulated different stage positions with black and silvered end-plates, as we did previously for the gain asymmetry measurements.⁶⁷ One silvered and one black end-plate simulated an end-stage position, and two black end-plates simulated an isolated amplifier position. For central-stage positions, we extrapolated the steering from the end-stage and isolated-amplifier results. This method was used because we anticipated destructive parasitic lasing from two silvered end-plates. We assumed, as previously, that flashlamp light from further-than-adjacent amplifiers was insignificant.

Figure 4-52a shows measured and extrapolated steering data for the 46-cm amplifier at the time of $G_{\max}$. Extrapolation to central stage steering, $St(\text{central})$, was done as follows for each horizontal position:

$$St(\text{central}) = St(\text{end}) + [St(\text{end}) - St(\text{isolated})] \quad (8)$$

Amplifier position evidently has very little effect on +side steering, since the +side half-disks are well away from the amplifier ends. The -side steering, on the other hand, goes from strongly negative for the isolated amplifier to essentially neutral for the central stage position.

Figure 4-52b shows the horizontal steering at $G_{\max}$ for an isolated 31.5-cm amplifier. We assume that its steering is relatively insensitive to intra-stage position because previous gain measurements⁶⁷ showed that intra-stage position had very little effect on its gain profiles.

We have not measured the smaller disk amplifiers on Nova, but we expect their contributions to the total aberration on the output beam to be substantially less than for the power stage, because of the larger amounts of beam expansion between them and the output.

Figure 4-53 shows the expected steering from the summed contributions of the 31.5 and 46-cm stages. The +horizontal direction and positive steering direction correspond to toward-the-frame at the 46-cm stage. (Recall that the individual amplifier measurements show no significant steering variations in the orthogonal direction, i.e. vertical at the 46-cm stage.) The contribution from the 46-cm stage includes two end-stage and two central-stage amplifiers all mounted with +sides toward the frame, and the spatial pattern was scaled by 46/74 to account for up-collimation. The contribution from the 31.5-cm stage assumes five amplifiers, four mounted with +sides toward the frame and one reversed; its pattern was scaled by 31.5/74 to account for up-collimation and inverted with respect to the axis to account for the 31.5-to-46 spatial filter. The right-hand vertical axis in **Fig. 4-53** also shows the transverse positions of rays in the focal plane of the Nova target lens (focal length = 3.2 m at 1053 nm).

Figure 4-53 shows that the power in the Nova output beam steers away from the central split, with the +side steering substantially more than the -side. The only exception is the small amount of beam power near the central split which is dominated by edge steering. The intensity-weighted average steering for the total beam is ~4 μrad (12 μm at the target plane), and for the two sides individually it is +12 and -4 μrad (38 and -13 μm), respectively. Since the diffraction limits at 1053 nm for the Nova output beams before and after the

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Fig. 4-52 Steering at G_{max} versus position for a) the 46-cm and b) the 31.5-cm amplifiers; positive steering is to the right or toward positive horizontal positions.

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Fig. 4-53. Steering at the Nova 74-cm output from the combined effects of the five 31.5-cm and four 46-cm amplifiers on Nova; positive steering refers to toward-the-frame at the position of the 46-cm stage.

harmonic generators are $\pm 1.8 \mu\text{rad}$ ($\pm$ half angle to the first null for a 70-cm round beam) and $3.9 \mu\text{rad}$ (for a 27-cm square beam), this pump-induced aberration is one to several times the diffraction limit in either case. Its most noticeable consequence would be an astigmatic focus, since the induced aberration is confined to the plane of incidence of the disks.

We found no data from Nova on the 1053-nm beam quality with which to directly compare this data, but measurements were made on Novette at 1053 nm,⁶⁸ and the induced aberrations would have been similar to Nova's, since the stagings of the two lasers were similar. The Novette measurements showed a beam spread (including 80% of the beam energy) of $\pm 17 \mu\text{rad}$ before the harmonic generators. However, there is no mention of an astigmatic focus.

More recent Nova data⁶⁹ on beamline 7 at 351 nm do show an astigmatic best focus of $\pm 23 \times \pm 37 \mu\text{rad}$, parallel and perpendicular to the disk split, respectively. The direction of the observed astigmatism is correct, but its magnitude is much too large to have been caused by these pump-induced aberrations.

Given these two pieces of system data, we concluded that the pump-induced aberrations from the Nova amplifiers are an insignificant part of the total aberration on the output beam. Nevertheless, the source and magnitude of this effect needs to be considered in any future laser design.

Implications for Future Glass Laser Systems. Laser system costing and performance calculations should include the effects of pump-induced central and edge steering. Edge steering reduces the useful clear aperture of a disk and increases the minimum spacing between beams, and central steering influences the output beam quality and the targeting accuracy of the system.

Pump-induced central steering is caused by non-equal pumping of the faces of amplifier disks, that occurs primarily for the end disks in an amplifier. This implies that future amplifiers should be built close-packed like the Nova 46-cm amplifiers for efficient flashlamp coupling between modules. This and possible compensation techniques require further investigation and are areas of on-going research.

Pump-induced edge steering is caused by thermal expansion of the edge cladding. Modeling indicates the following rule-of-thumb for the extra disk aperture required for stored energies similar to the 31.5-cm disks of Nova:

Holding edge-steering to less than $5 \mu\text{rad}/\text{disk}$ requires a distance of about one-half the disk thickness between the edge of the beam and the edge cladding. It's also worth noting that the development of a low expansion edge cladding would eliminate the edge-steering problem entirely.

Author: J. E. Murray

Major contributors: H. T. Powell, D. Nikkel, G. F. Ross, and J. D. Wintemute

4.3.4. 3ω Diagnostics

Overview. Past incident beam diagnostics recording both energies and pulse shapes at harmonic wavelengths which are incident on the target do not meet present objectives set forth by the "Precision Nova" project. Several designs have encountered a variety of technical problems—all of these problems stem from the fact that the diagnostic uses the highly-aberrated retro-reflection from the Nova focus lens. The feasibility and cost of a using full aperture beamsplitter and second focus lens to reduce the beam size was studied and determined to be unacceptable. Performance goals for an improved incident beam diagnostic are motivated by the objective of balancing Nova's beams 3ω power to better than 5%. Our initial design goals were: energy measurement's precision of $\pm 1\%$, and streak measurements sampling the beam's full aperture with less than 10% variation in sampling across the aperture.

During 1988 and most of 1989, we developed and implemented a design based on recollimating the retro-reflection from the focus lens and then reducing the diagnostic beam's intensity to a manageable level with an uncoated reflector. Initial tests were conducted in the two-beam target bay on the "Phase II" diagnostic platform. These tests were largely successful, showing good agreement with reference measurements of both pulse shape and energy. Prototype diagnostics were then installed on beamlines 5 and 6 in the ten-beam target chamber. Subsequent tests demonstrated that the incident beam diagnostics agreed well with the 44-cm reference calorimeter over a period of many weeks. Significant disagreement between the diagnostic energy and reference was seen, however, over a range of intensities. This disagreement was attributed to the polarization sensitivity of the design. We believe stress induced birefringence in

the focus lens causes polarization ellipse rotation at high intensities, via the nonlinear index of refraction.

Incident-beam-diagnostic recollimating design. A schematic of the first prototype incident-beam diagnostic is shown in Fig. 4-54. Major components of the diagnostic include the following: strut assembly, negative lens, uncoated beamsplitter, UG-11 filter, positive lens, 7-cm calorimeter, diffuser/fiber taper/fiber optic, and diode.

The 3-cm wide supporting struts were narrow enough to fit within the disk split, and allowed for ± 10 mm of transverse motion of the focus lens. The disk splits on the majority of the beamlines are not vertical in the diagnostic spool, however. The struts, therefore, were designed to be sufficiently rigid to prevent the package from bending more than 2 mm when installed on beamlines that are not horizontal.

The diagnostic reflection off the second surface of the focus lens contains a large amount of spherical aberration. Photos of the reflected beam were taken to document these aberrations and determine the distance away from the focus lens where caustics around the edge of the beam begin to form. To minimize fluence on the lens, the fused-silica negative lens was placed along the optical axis before caustics begin to form in the reflected beam. To reduce costs, this lens was designed with spherical surfaces. While aspheric surfaces could be used in this lens to correct virtually all of the spherical aberration, alignment tolerances would become very small. The negative lens recollimates the diagnostic reflection at a diameter of 160 mm, and collimation varies less than 1 mrad across the aperture, sufficient to propagate the short distance in this diagnostic. Intensity modulation across the aperture of the collimated diagnostic beam was roughly 3:1.

To reduce the total fluence in the diagnostic beam, an uncoated fused silica beamsplitter was used to reflect the beam out of the incident beam's aperture. At this reduced intensity a Schott UG-11 glass filter can safely separate the fundamental and third harmonic wavelengths without damaging or exhibiting non-linear transmission. Previous studies have shown that no significant nonlinearities exist at fluences below 100 mJ/cm^2 .⁷⁰ Peak fluence incident on this filter was 40 mJ/cm^2 at 3000 J 3 ω incident energy.

After reflecting from the surface of the 7-cm calorimeter (Hoya-glass B430) the f/2 beam

comes to a focus, with a spot diameter of 6 mm on the diffuser/fiber taper combination (see section 4.3.5.). This device is a tapered fused silica rod with an input aperture of 6 mm and an output aperture of 200 μm . The diffuser/fiber taper is used to increase coupling efficiency into the 50- μm fiber optic, while still sampling the entire beam aperture. Because the fiber taper accepts only nearly collimated light, it is necessary to use a diffuser to scatter a fraction of the light into the acceptance cone of the taper.

Significant distortion of the pulse shape can be caused by stimulated Brillouin scattering at high intensities in the fiber optic. Studies have been made to determine the peak intensity that can be transmitted through the fiber taper and fiber combination without pulse distortion.⁷¹ The fiber was utilized at an intensity below the SBS threshold.

The accuracy of this design depends on the stability of the chamber optics transmission and the reflectivity of the focus lens. A system to measure changes in these factors, the "calibration factor," is included in the diagnostic package. The 7-cm calorimeter was designed to translate out of the beam, allowing the light to focus onto a silicon photodiode. At target chamber center, an insertable positioner with ten silicon photodiodes provides the second half of the ratio measurement. The detection system is synchronous with a chopper at the alignment laser source. All photodiodes were selected based on uniformity of response across their surface to minimize alignment errors between measurements.

A few compromises are incorporated in the diagnostic design. A slight increase in the size of the central obscuration from 2.5 to 5.3% of the beam area was necessary in order to accommodate the negative lens and mounting hardware. Although 2ω measurements are not readily made with this configuration of the diagnostic, 2ω energy can be measured accurately with a 160 mm diameter calorimeter replacing the reflective beamsplitter.

Off-Line Full-Aperture Tests. The prototype diagnostic was tested off-line on a full-aperture test station (see section 4.2.2. for a description of the station). Reference diagnostics measuring the full aperture energy and pulse shape were available using this station.

Excellent agreement between the diagnostic and reference calorimeters was obtained. A variety of energies and pulse lengths were

Fig. 4-54. Prototype incident-beam diagnostic (recollimating design).

tested, ranging from 100 ps to 2.5 ns, and from 150 J to 3200 J at 3ω . RMS variation in energy measurement between the diagnostic and reference calorimeters was $\pm 1.9\%$ for 16 shots. No damage occurred to any of the diagnostic optics. The S/N ratio of the diagnostic calorimeter signal was such that reliable measurements of 3ω energies as low as 100 J were possible.

Good agreement between pulse shapes recorded with the incident beam diagnostic and a full-aperture reference was observed. A comparison of 100-ps, 1-ns, and 2.5-ns shaped pulses is shown in Figs. 4-55, 4-56, and 4-57. Signal levels through the tapered fiber appear adequate to record reliable 3ω pulse shapes above 1 TW. No pulse shape distortion due to SBS in the fiber optic was noted. A Bell Labs UV fiber (with no measurable dispersion at 351 nm) was installed for two 100-ps shots, and dispersion was reduced to within the error of this measurement.

On-Line Tests: Beamlines 5 and 6. Two sets of the incident beam diagnostic were installed onto beamlines five and six during the month of June. A comparison of incident-beam-diagnostic energy measurements with a reference calorimeter was made over a period of three months. Approximately 290 shots were fired to both beamlines; RMS variation between the diagnostic and the reference energy measurement was $\pm 3.0\%$.

Signal levels on the streak camera were marginal using the UV fiber and a Tektronix camera lens on the input end of the streak camera. Replacing the Tektronix camera with a lens system with higher UV transmission (see section 4.3.5.) increased signals to an acceptable level. Then good agreement between the 1ω and 3ω streaks was demonstrated.

During October, a series of eight shots was performed to measure the linearity of the incident beam diagnostic over a range of intensities. Ratios of energy measured in the incident beam diagnostic and the reference calorimeter are plotted in Fig. 4-58. Although there is appreciable scatter in the data set, a non-linearity of approximately $9\%/TW$ was observed. This nonlinearity is attributed to stress-induced birefringence in the focus lens inducing polarization ellipse rotation at high intensity levels.⁷² Since the reflectivity of the uncoated splitter varies significantly as a function of polarization shifts, there is a shift in apparent energy measured in the incident beam

diagnostic. Additional tests measuring the effect of relieving the stress (not at vacuum) on the focus lens confirmed this picture. Approximately 15% less energy reached the diagnostic calorimeter in this configuration, as compared to when the focus lens was stressed under vacuum.

Initial measurements of the calibration factor using photodiodes in the diagnostic and target chamber center were encouraging. However, additional data are required to validate the system. Results from a single measurement were very good. The ratio measured was stable to less than 1% over a 2-mm range of motion for the focus lens in x, y or z. The reproducibility of the insertion mechanism at chamber center was also acceptable. Nova scheduling conflicts precluded obtaining conclusive results in three other areas concerning long term performance. The first is whether the ratio is reproducible over a period of several days when no shots are fired and no chamber conditions are varied. The second concern is whether the cw measurement agrees with calibration shots on an ongoing basis. A final concern for the system is the impact on the Nova schedule in order to make the measurements. The time involved needs to be as short as possible so as not to slow the overall Nova shot rate. Measurements of the calibration factor will be made at a frequency in accordance with the rate at which the chamber optics change their transmission and reflection. We will determine this frequency once a fully activated system is installed.

Limited Tests of Diffuser Design. From the data in Fig. 4-58 it was deemed that a diagnostic which was insensitive to changes in the polarization of the beam was required. A different conceptual design was developed late in 1989, using a diffuse reflector to attenuate the diagnostic reflection. Fig. 4-59 diagrams a prototype diagnostic that was installed onto beamline six.

A 2-mm thick piece of Spectralon (a material based on polytetrafluoroethylene made by Labsphere) that scatters the diagnostic reflection in both reflection and transmission was placed adjacent to the output surface of the negative lens. This diffuser attenuates the diagnostic reflection by scattering nearly 100% of the light in a reflected Lambertian distribution. A pinhole is used to image the diffuser into the entrance aperture of an integrating sphere, thus limiting the field of view of the diagnostic. A

Fig. 4-55.

Fig. 4-57.

Fig. 4-56.

Fig. 4-55, 56, and 57. Comparison of pulses recorded through UV fiber-optic couple air-path coupling to LLNL streak camera.

Fig. 4-58. Non-linearity of energy measurement in prototype incident. Beam diagnostic is approximately 9% per terawatt.

Fig. 4-59. Refined incident-beam diagnostic was found to be polarization insensitive and linear in energy response.

charge amplifier was used to integrate and amplify the diode's signal. Three shots were taken using this diagnostic using 2-ns square pulses, with intensities ranging from 0.7 to 1.8 TW. Agreement of the diagnostic and reference energy measurements were within 3.5% RMS. Further development of this design to include pulse shape measurements, is planned for 1990.

Author: C. W. Laumann

Major contributors: J. R. Smith, J. B. Richards, J. A. Caird

4.3.5. Streak Camera Fidelity

Streak Camera Performance. The information base on streak camera performance has grown significantly. As a result, data quality has improved as well. The quantity of data remains high as these CCD coupled streak cameras are used throughout the system to record high speed pulse shapes ranging from several nanoseconds down to 15 ps. They are used at all three Nova wavelengths and can use fiber optic inputs or direct light input. The goal of these instruments is to record the pulse shapes to within an accuracy of $\pm 10\%$ at any point along the trace.

One of the key issues involving camera stability is whether or not the camera changes its response characteristics over time. The main technique used to determine this is to input a temporally flat pulse into the camera and record how flat the camera response is. The primary purpose of this flat field characterization is to correct raw data taken from the camera. Since the nonuniformities of the camera can be quantified, their effect can be compensated for and removed from the final CCD image. Many of the cameras that were characterized in the past have been recharacterized recently so we have some indication of their stability. We have found that this flat-field characteristic of the cameras is relatively stable over time. In general, changes measured over the course of one year have been less than $\pm 10\%$ and typically around $\pm 6\%$. This represents worst case deviations across the entire sweep.

In addition to flat-field response, the high end linearity of the dynamic range is now routinely being analyzed. The cameras vary among themselves with the 10% nonlinearity level occurring between 700 to 3300 counts, depending on the camera and sweep card used.

This saturation level information is included in the on-line data processing to warn operators of possibly saturated data.

The lower limit of the dynamic range was tested on-line during a shot series involving high contrast ratio pulses. During these "pulse podiatry" tests, signal levels of 30 to 40 counts in the foot of the pulse were determined to be accurate to within 10%. Combining this result with the characterization data gives the net result that the LLNL streak camera has a useful dynamic range that approaches 100:1.

A significant improvement has been made in the lens coupling between the slit at the front end of the streak camera and the photocathode. A new lens system has been designed to replace the old one which had very poor (5% at 3ω) UV transmission. The new version is estimated to transmit 68% at 3ω , 82% at 2ω and 77% at 1ω . In addition, it should transmit 55% at 263 nm. No new hardware is required as the lens mounts in the same C-27 housing as the old lens. The new system has 1:1 magnification and is made up of four identical silica-calcium fluoride doublets. Two doublets are mounted in a cell, with one cell on each side of a central shutter.

The primary benefit of the new lens is that it makes the streak camera more sensitive to UV light. This in turn allows for less light to be propagated in the fiber optic transmission lines for a given signal level at the streak camera. This is particularly important in order to avoid nonlinear interactions such as stimulated Brillouin scattering from occurring within the fibers, particularly at 3ω .⁷³ Such pulse deforming effects take place only if the intensity within the fiber exceeds a certain threshold. The new lens ultimately allows for operation at levels well below this critical threshold.

Reliable data was obtained with LLNL streak cameras during shot campaigns using a pulse width of 15 ps. Specific results depended upon which camera was used because of significant variations in their sweep speeds. In addition, care must be taken to keep signal levels low so as not to saturate the camera. The primary camera recorded a pulse width of 15 ps when the pulse was expected to be 15 ps. This implies that the camera's resolution is at least that good when the proper precautions are taken.

Data Quality. While efforts continue to develop a 3ω incident diagnostic package that adequately samples the incident beam, the

design of the temporal portion of an early prototype was quite successful. This is mainly due to the use of a large aperture fiber optic taper that can efficiently couple light from the diagnostic reflection off the focus lens to the recording instrument. The inherent problem with using the primary ghost reflection off the Nova focus lens is collecting the beam given its extreme aberrations. If the entire beam is to be captured, as was the case in the early prototype and is shown in Fig. 4-60, aberrations limit the beam's best focus to 4 to 6 millimeters in diameter. Such a beam is extremely difficult to efficiently couple to the multiplexing fiber optics of the streak camera network. While the taper has proven to be effective in this small aperture system, its effectiveness with large aperture diffusers is undetermined. The principal concern is low signal levels that result from the use of two diffusers. A first diffuser is needed to uniformly illuminate a region directly in front of the taper. A second diffuser is needed to redirect the light into the taper which, to a first approximation, only accepts collimated light. Signal levels can be increased by using weaker diffusers but only at the expense of uniform spatial sampling of the beam. Sampling concerns are especially important when high intensities and nonlinear effects can affect both the beam and the diagnostics. Low signal levels and uniform sampling are two of the reasons why photodiodes are being considered for 3ω temporal diagnostics. The obvious concern with photodiodes is their relatively slow response time. Whichever transducer is used, a final concern is the lack of a convenient reference measurement by which to determine the absolute quality of the temporal data.

Prototype tests have shown that streak camera data from the 1ω output sensors can be significantly improved by modifying the coupling configuration into the fiber relay system. Once again, as shown in Fig. 4-61, the use of a fiber optic taper is the key. The beam in the output sensor package is well controlled and thus makes the interface much cleaner than in the case of the 3ω sensor package. The old configuration relied on collection by a small 50 μm fiber which was susceptible to speckle variations that result from the scatter of the coherent laser beam off the diffuser as seen in the far field. The new configuration gives more consistent signal levels and more reliable pulse shape measurements because the taper provides

a larger aperture over which to sample scattered light from the diffuser. Regardless of the details of the streak camera configuration within the output sensor, the sensor itself is known to have a polarization dependence that affects the beam sampling. Again, this is particularly important at high intensities where nonlinear polarization effects produce variations across the beam.

Overall, LLNL streak camera data continues to improve with our understanding of hardware limitations and of system implications regarding how accurately the light that enters the camera is representative of the light that ultimately hits the target .

Author: J. Richards

Major Contributors: J. A. Hatch and C. E. Thompson

4.4. Nova Engineering

4.4.1. Introduction

By the end of 1989, the Nova laser had been operating for 5 years. During this period, many changes have occurred and much has been learned; such that Nova has become a reliable, versatile, and productive facility. In the process, Nova either met or exceeded all its original performance milestones, and the operational efficiency has increased dramatically.

The experience gained in laser science, target physics and operations have led to new requirements on the Nova facility. For example, the magnitude of plasma instabilities can be diminished by reducing the coherence of the focused beam on the target. This is accomplished, in part, by increasing the frequency bandwidth of the laser beam. We have also learned that the large, fused silica focus lenses on the target chamber are structurally damaged when the energy of the pulse at 0.351 μm exceeds certain limits. The damage mechanism is an intense acoustical wave initiated by transverse stimulated Brillouin scattering. Increasing the bandwidth of the laser beam also suppresses the SBS. But we now have the engineering requirement that the method for increasing bandwidth be fail-safe, otherwise we have the potential of simultaneously destroying all 10 target chamber focus lenses at a cost of over \$1,000,000. Two approaches for increasing the

Fig. 4-60. Output sensor interface to 1 ω streak cameras.

Fig. 4-61. Prototype spool configuration illustrating the fiber optic taper and UV optical fiber providing a 3ω temporal diagnostic.

bandwidth of the Nova master oscillator are reported (see Sec. 4.4.4.).

To better understand the laser/plasma interactions, more detailed diagnostics of the laser beam are required. We are currently constructing an advanced laser diagnostics station on beamline 26. The entire frequency converted laser beam will be sampled with a fused silica beam splitter, and the sampled beam will be projected to an equivalent focal plane for measurement (see Sec. 4.4.5.).

Any state-of-the-art system is bound to exhibit signs of aging after five years of vigorous service. Nova is no exception. As discussed in Sec. 4.4.3., the flashlamps in the large amplifiers have been a weak link in system reliability. Our efforts to correct that problem are described. Also many advancements in computers, networks, graphics and support software have been made since the Nova control system was originally built. Therefore we are upgrading our control system, with the expectation that this will both increase system reliability and reduce maintenance costs. In addition, the new control system will increase the effectiveness of operation personnel because we will also be able to combine functions once these changes are made, and exchange of data between various Nova subsystems can be made directly (see Sec. 4.4.2.).

Author: D. Kyrazis

4.4.2. Control-System Upgrade

The Nova computer control system is being upgraded to replace outdated systems with new computers and networks that follow current industry standards. This will allow expansion of capabilities and the use of computer systems from many vendors. The new system is based on workstation technology interconnected with the Ethernet network and the X Window System version 11 for graphics. During 1989 a preliminary design of the control system was completed and the project was approved. We expect to complete the project at the end of fiscal year 1991.

The Nova computer control system was designed in 1979 and 1980. The system was designed to control the four major subsystems of the Nova laser: power conditioning, laser alignment, laser diagnostics, and target diagnostics. The control system is a two-level computer hierarchy consisting of a top level of three VAX (Digital Equipment Corp.) 11/780

computers interconnected via shared memory and 67 front end processor (FEP) computers connected to the three VAX computers with a fiber optic network called Novanet. (See Fig. 4-62 for the control system network schematic.) The FEP computers are LSI-11 computers and are connected to the hardware that controls the laser. The three VAX computers provide the user interface and system integration functions. There are six Ramtek 9450 graphics controllers connected to the VAX computers that provide 17 color displays for the users. Ten of the graphics display screens have touch sensitive panels that are used for user input. The control room has four consoles, one for each subsystem. Each console has three display screens. Only the center screen in each console can accept touches, the other two screens are for display only. There are six other single screen control stations located throughout the facility for doing shot setup and test work.

The original computer hardware for the control system is reaching the end of its useful life. Expansion is not possible and it is beginning to get difficult to find service for the hardware. The maintenance costs for the hardware are high. Novanet system is also nearing the end of its useful life and more components are failing. Some integrated circuits are no longer manufactured. There is no immediate danger for catastrophic failure because of a large stock of spares for Novanet but the problem will eventually have to be faced. We have proposed to replace the VAX 11/780 computers, Ramteks, and Novanet systems to provide a longer life to the control system and allow expansion of the system.

The new control system will be comprised of a network of workstations connected by Ethernet. (See Fig. 4-63 for the new network schematic.) The new user interface systems will be VAX station model 3100-38 workstations. Each workstation will have 32 megabytes of memory and a 52 megabyte disk for virtual memory paging. There will be five workstations, one for each of the existing subsystems and one for a new subsystem, the shot director's console. The workstations will replace the center Ramtek screen in each console and a keyboard and mouse pointing device will replace the touch screen. The other two screens in the consoles will be X windows, terminals connected to the workstations via Ethernet. A special mouse/keyboard multiplexor device developed at LLNL will be used to allow a single mouse to

NOVA CONTROL SYSTEM ARCHITECTURE

All connections between computers, workstations, and X-terminals will use the Ethernet network (IEEE Standard 802.3). 21 X-terminals and workstations will be used in the new control system.

Fig. 4-63 New Nova control system computer network.

be connected to any of the three screens. This configuration allows any subsystem to be controlled from any screen. The existing system screens were dedicated to a particular subsystem and only screens with touch panels could be used for control inputs. The five workstations will be served by a dual host MicroVAX 3800 file server system that provides central file storage. A dual host system was chosen to improve reliability. If one server processor fails the other will automatically take over the file service tasks. The file server has a capacity for 4.8 gigabytes of disk space versus 2.13 gigabytes in the old system.

The graphics system software which was written for the Nova project will be replaced with the X-11 Window System standard software. The X-11 standard allows the use of remote displays across the network and different hardware platforms without changing programs. We have purchased a graphics editor called GMS from S-L Corp. which will be used to develop the user interfaces. This editor allows a developer to work directly with an operator to design the look and feel of a new screen. This rapid prototyping capability will help make better control screens faster. The capabilities of X-Windows and the S-L product help us design user interfaces that are similar to the Apple Macintosh computers. The Macintosh interface is known by many operators, is easy to learn, and is preferred over the old style user interfaces.

The Novanet network to the FEPs will be replaced with Ethernet. The capability for the replacement was developed in the Isotope Separation program. We are keeping most of the LSI-11 based FEPs, 36 for laser alignment, 11 for diagnostics energy processors, and 8 for streak cameras. These will all be interconnected with Ethernet. The three power conditioning FEPs will be replaced by two MicroVAX 3400 processors running the VAXELN real time operating system software, the nine target diagnostics system FEPs by a single MicroVAX 3400 system running the VMS operating system. All the new systems will use the Ethernet network for communications.

To use workstations and the X-Window system, all of the user interface software will be replaced. However, we will not rewrite any existing software that can be preserved. The ADA language is being used for all new software along with the VAX ADA compilation system and Code Management System product from DEC

to aid program development and keep track of different versions of programs to control software releases. Building the original Nova control system required 67 man-years of effort. The new control system will require about 15 man-years of effort to build. Effort is saved by reusing parts of the existing system, using industry standards, and using new software tools such as the graphics editor to produce new programs.

There will be a total of 21 control screens in the Nova laser facility when the new system is installed. Any of these screens can control any part of the laser. We are increasing the computing power by a factor of 9 to provide better and faster user interfaces. System management of the network will be easier with a centralized file server system. The cost of maintenance for the new system will be cut in half and upgrades to the hardware will be easier and cheaper. Thus far we have purchased the graphics editor and have ordered the first file server processor and three workstations. A demonstration of the new controls technology is planned for March 1990.

Author: W. G. Labiak

Major Contributors: Bob Demaret, Vicki Miller, Dave Kroepfl, Dave Johnson, Paul Van Arsdall, Pete Ludwigsen and Barry Cline

4.4.3. Flashlamp Failure Studies

Since its activation, the most serious and costly problem with the Nova power conditioning system has been flashlamp failures.⁷⁴⁻⁷⁶ The failures have either been due to cracks that develop in the glass-to-metal seal, or to flashlamp explosions. The seal failures result in air leaking into the flashlamps thus preventing them from triggering. We believe the explosions are caused by cracks in the lamp envelopes, which may result from brown contamination spots found on the outside of nearly all of the Nova lamps. This year, we identified a solution to the seal failure problem, which has been the most troublesome so far. Work is in progress to identify the source of, and solution to the contamination problem, as well as to find the magnitude of its potential impact on Nova in the future.

Seal Failures. The seal failures occur only on lamps from one of our two vendors and only on the model of lamp which is used in the 46-cm amplifiers. Of the 2000 lamps of this type

installed in Nova, about 400 have failed to date, while about 1000 remain in the laser. Since replacing this quantity of lamps is quite expensive, it is desirable to continue to have at least two qualified sources to minimize prices through competition. Hence we have worked to identify the source of the problem, shared that information with vendors, and developed improved prototype.

We identified the cause of the seal cracks by examining the failures microscopically. Opaque material possibly emanating from the tungsten rod becomes deposited on the air side of the glass-to-metal seal (see Fig. 4-64). This material absorbs flashlamp radiation resulting in thermally induced stress and ultimately cracking the quartz. The vendor proposed two solutions to this problem: (1) coating the tungsten wire with glass to shield the seal from any contamination that might originate on the wire, or (2) reversing the seal assembly to place the fragile portion inside the aluminum base cover, thus preventing the irradiation of any contamination on the seal. A new qualification test was devised to ensure that the tests of proposed solutions were statistically valid, based on failure rate data from Nova. The lamps were installed in a Nova 46-cm amplifier box, without disks, and fired for 1000 shots. Greater than one failure resulted in rejection of the design. The coated wire was the lower cost design and was therefore tested first. Two lamps had seal failures after a few hundred shots and the design was disqualified. The reversed seal assembly design passed the 1000-shot test without failures, thus requalifying this vendor. Our second vendor avoids the seal cracking problem completely by using a soldered glass-to-metal seal. A third vendor from Great Britain also failed to qualify due to seal cracks. We are now beginning to purchase lamps to replace the remaining lamps on Nova with the unreliable seal design.

Envelope Cracks. Nearly all of the lamps on Nova suffer from brown contamination spots on the surface of the quartz envelope. A typical flashlamp panel from a 46-cm amplifier might have several hundred spots, each over one cm in diameter. A micrograph of a typical spot is shown in Fig. 4-65. The spots result in small cracks that propagate near the surface of the quartz, presumably as the result of thermal stresses that occur when the spots absorb flashlamp radiation. The cracks are typically

only 10-20 mm in depth and do not extend past the area covered by the brown spot. We believe that under certain circumstances, these cracks may grow in length and depth, thus weakening the envelope and causing the flashlamp to explode during a shot. An effort is underway to identify the source of the spots, clean or replace the contaminated lamps, and prevent spots from forming on new lamps.

Author: D. W. Larson, A. C. Erlandson

Major Contributor: W. Sawvel

4.4.4. Oscillator Development—FM Operation of the Nova Master Oscillator

The transverse stimulated Brillouin scattering problem that potentially threatens the target focusing lenses on Nova can be controlled by sweeping the frequency of the initial pulse from the master oscillator room (See Section 4.2.3. of this section). The experimental demonstration of this effect used an Austin-switch-driven Pockels cell to create the frequency sweep on the 2.5-ns output pulse from the MOR's oscillator and electro-optic pulse chopper. This demonstration worked very well, but it would be difficult to implement in a fail-safe configuration for routine operation on Nova because the go/no-go decision would have to be made very quickly, in about 100 ns or so. Since SBS can lead to target lens fracture, which could potentially cause mechanical failure and subsequent lens implosion into the target chamber, failure to impress the frequency sweep on the oscillator pulse could lead to serious damage in the target chamber as well as catastrophic damage to the target lenses.

Another means for obtaining a frequency-swept oscillator pulse is the so-called FM mode of operation for the oscillator.⁷⁷ In principle this approach gives no amplitude modulation of the laser output which is desirable from the point of view of minimizing self-focusing of the beam caused by the nonlinear index of refraction. The basic idea is to couple the longitudinal modes of an oscillator in such a way that the instantaneous frequency sweeps sinusoidally approximately at the $c/2L$ mode-spacing frequency of the resonator. Mode coupling is accomplished with an intracavity phase modulator, detuned slightly from the $c/2L$ frequency. The magnitude of the frequency sweep, which depends directly on the

Fig. 4-64. Conventional and reverse re-entrant flashlamp seals.

Fig. 4-65. Micrograph of brown spot showing both developed and undeveloped cracks.

modulation depth of the modulator and inversely on the detuning, can be more than 100 times that of a single pass through the modulator. Since a frequency sweep comparable to the full-width half-maximum of the gain bandwidth ($\sim 4 \text{ cm}^{-1}$) has been demonstrated with Nd:YAG,⁷⁸ we felt the prospects of generating the required 0.5 cm^{-1} in Nd:YLF which has a substantially wider bandwidth were excellent.

A major attraction of the FM-laser approach was that the FM process would be allowed to reach steady-state before Q-switching the oscillator, as with the current mode-locked MOR oscillators. We assumed (and planned to experimentally verify) that Q-switching would not influence the FM character of the oscillator output, which would allow the fail/safe electronics the $\sim 1 \text{ ms}$ time period prior to Q-switching in which to confirm the frequency chirp. Another significant attraction was the potential simplicity of the modification required to the MOR oscillator. In principle, the Nova long pulse oscillator (LPO) could be converted for FM operation by the simple installation of a phase modulator—all the operational aspects of the oscillator could remain the same.

To test FM operation of an LPO, we installed a phase modulator in an off-line Nd:YAG LPO. We used Nd:YAG as a first step to avoid the potential birefringence effects with Nd:YLF, and because there were examples of successful FM operation of Nd:YAG lasers in the literature.

The phase modulator used a 1-cm-long, Brewster's-angle-cut LiNbO_3 crystal with a 4-mm square aperture and oriented with the c-axis in the plane of incidence. The voltage was applied along the c-axis, and the modulator was matched to a $50\text{-}\Omega$ cable. We measured a single pass modulation depth of 0.25 radians for an applied voltage of 20 V, RMS.

We used both temporal and spectral diagnostics to investigate the effects of the phase modulator. A photodiode with enough sensitivity to display the cw prelude signal also had a measured risetime (with an oscilloscope) of 0.6 ns, fast enough to observe longitudinal mode beating. For spectral information we used an optical spectrum analyzer with a free spectral range of 300 GHz and a measured finesse of $\sim 40,000$. To observe the spectrum of the LPO we drove the spectrum analyzer with a 0.5 ms voltage ramp which was synchronized to

display the last portion of the prelude. The spectrum analyzer did not respond linearly to the fast-ramp voltage, but the resulting distortion did not mask the essential features of the spectrum.

Figure 4-66 shows a time trace (a) and frequency spectrum (b) for the LPO with the modulator installed but not driven (i.e. free running). The figure shows a very complicated time behavior for the free running LPO, as expected, because of interference between the many longitudinal modes that were simultaneously oscillating. Observation of this time behavior with a slow temporal detector showed only a few percent amplitude modulation after the initial relaxation oscillations damped out, indicating that the total power in all the oscillating modes was constant to that level.

The spectrum of the free running LPO, shown in Fig. 4-66b, is reasonably well behaved, indicating a single set of 5 or 6 adjacent longitudinal modes. The spectrum changed with time, "rolling" back and forth randomly by one or two modes, but it did not jump cavity modes. The installation of a thin etalon (2 mm) in the LPO cavity allowed the spectrum to be tuned continuously over several tens of GHz. These observations indicate there were no significant intracavity etalons (with spacing greater than about 1 cm) and no significant feedback from extracavity components.

Figure 4-67 shows the time behavior and typical spectra when the phase modulator was driven near the $c/2L$ frequency of the LPO where it should have produced FM operation. The time behavior is substantially different than for the free-running case, but the amplitude modulation shows no sign of damping out. The spectra (for two values of detuning) show some characteristics of the desired FM spectra, namely the higher mode amplitudes near the ends of the spectrum and the general indication of symmetry, but the zeros are washed-out and the detailed symmetry is not present. (Note that the changes in the mode spacing apparent in the spectra results from the nonlinear sweep rate of the Fabry-Perot, as previously mentioned.) Increasing the drive power to the modulator or decreasing the detuning increased the overall spectral width, but it did not decrease the amplitude modulation.

We also set up a streaked Fizeau interferometer, as shown in Fig. 4-68, to directly

(a)

(b)

4-66. a) Temporal and b) spectral traces of the free-running LPO; the frequency spacing between modes in b) is 175 MHz.

(c)

(b)

4-67. a) Temporal and b and c) two spectral traces at different detunings for the LPO with the phase modulator active; the frequency spacing between modes for both spectral traces in is 175 Mhz.

Fig. 4-68. The experimental set-up for the streaked Fizeau interferometer.

observe the sinusoidal frequency variation on the beam. The Fizeau created an interference pattern of linear fringes which was imaged onto the input of a streak camera or the photocathode of an SIT camera. The fringes were separated in frequency horizontally and streaked vertically, resulting in a frequency-versus-time display on the streak camera. It was necessary to Q-switch the LPO to get enough intensity for the streak camera, and to pulse-slice a portion of the Q-switched envelope to prevent multiple traces by the streak camera.

Figure 4-69 shows streaked data obtained with different powers to the modulator and different detunings. The straight vertical fringes of **Fig. 4-69a** with the modulator off indicate no frequency variation during the pulse. Temporal modulation is evidenced by the nonuniformity of the streaks in time. (The apparent multiplicity of lines within a fringe is an artifact of the Fizeau setup, since it was also present with a single mode output.) The effect of the modulator is evident in **Figs. 4-69b and c**, which show the expected sinusoidal frequency swings caused by FM operation for large and small detunings, respectively. **Fig. 4-69c** shows a frequency swing of about 42 GHz, substantially more than the 15 GHz needed. However, the streaks consist of broken lines in the temporal direction, indicating the same type of temporal modulation seen in **Fig. 4-67a**. The same LPO could be easily and stably FM mode-locked by tuning the modulator to the $c/2L$ frequency of the cavity, as shown in **Fig. 4-69d**.

The partial FM operation indicated by the amplitude modulation in **Figs. 4-67 and 4-69** suggests that something was preventing the LPO from settling into a pure FM mode of operation. Two possible explanations which we investigated were cavity length fluctuations and a build-up time for the FM process which was substantially longer than the 4-ms prelude time of the LPO.

We measured the cavity length fluctuations interferometrically with a cw, single frequency Nd:YAG probe laser. The coherence length of the probe laser was sufficiently long that reflections from the front and rear mirrors created Fabry-Perot fringes with a finesse of 17, which were then monitored with a photodiode. With the LPO lamps and cooling water off and with the intracavity beam path about 99% enclosed to prevent index fluctuations in the air, we measured a cavity length change of about .015

μm over times as short as 1 ms. This represents a cold cavity length stability of about 1 part in 10^8 . Turning on the cooling water and lamps at levels well below the lasing threshold did not significantly change this stability. With the LPO above threshold, the heating of the YAG rod caused a continuous path length change in the rod throughout the 4-ms pump pulse. The total length change during the pump pulse was estimated to be about 0.15 m in the 4 ms. In addition to this approximately steady increase in path length, there was a five-fold increase in the random cavity length fluctuations to 5×10^{-8} at this pumping level. Nevertheless, this level of cavity stability should have been adequate to establish at least intermittent FM operation of the LPO.

Another possible problem for the FM mode of operation was the build-up time required, since the LPOs have a prelude period of only few milliseconds. To test this possibility, we measured the FM build-up time of a cw Nd:YAG laser using the same intracavity phase modulator as before. The particular cw Nd:YAG laser used was not sufficiently stable for careful build-up time measurements. In fact FM operation occurred only at a few discrete and very narrow (~ 10 Hz) frequency tunings. In contrast, stable FM operation should occur for detunings of $\sim 1\%$ or more of the drive frequency (~ 2 MHz) on either side of resonance with the $c/2L$ cavity frequency. Nevertheless, the laser was adequately stable for periods of a minute or so to demonstrate nearly pure FM laser operation, as evidenced by its spectrum and a reduction in temporal modulation.

Figure 4-70 shows the time development (a) and output spectrum (b) for FM operation of the cw YAG at one of the stable values of drive frequency tunings. The time trace was obtained by unblocking the laser cavity, and triggering the oscilloscope on the first relaxation spike. The spacing between modes of the spectrum is the drive frequency (~ 175 MHz).

For a pure FM output, the intensity would be constant in time and the mode amplitudes in the spectrum would be exactly symmetric about the center of the spectrum. The time trace in **Fig. 4-70a** shows a smooth decrease in amplitude modulation to a minimum of about 10%. We suspect that the residual modulation results from an intracavity etalon caused by the normal-cut YAG rod. The spectrum in **Fig. 4-70b** is nearly symmetric, but not exactly, which is probably

(a)

(b)

(c)

(d)

Fig. 4-69. Streaked Fizeau fringes showing frequency in the horizontal and time increasing downward, with a) the phase modulator off, b) the phase modulator on and a large frequency detuning, c) the same applied modulator voltage but small detuning, and d) the same applied voltage but zero detuning; the horizontal fringe spacing is 30 GHz.

(a)

(b)

Fig. 4-70. a) Temporal and b) spectral traces for the cw Nd:YAG oscillator operating in a nearly pure FM mode.

associated with the residual amplitude modulation. Consequently, this does not represent pure FM operation of the laser. Nevertheless it is substantially closer than what we obtained with the LPO, and its build-up time to nearly pure FM operation should indicate if the LPO has an insufficiently long prelude period.

We measured the desired build-up time by monitoring the decrease in relaxation oscillation amplitude and the development of an approximately modulation-free output after unblocking the laser cavity. Figure 4-70a shows that this time is about 0.5 ms, well within the prelude time of the LPO.

We have no explanation for why the LPO would not FM properly. However, this was to have been an interim fix at best for preventing transverse SBS in the focusing lenses, and it does not provide enough bandwidth for the planned beam smoothing experiments on Nova. Furthermore, part of its appeal was the possibility that it could have been implemented with a very simple modification to the existing LPOs. Since this turned out to be not the case, this effort was abandoned in favor of longer term solutions with substantially more bandwidth, like the crossed-phase modulation work described below.

Author: J. E. Murray

Major contributor: R. Wilcox

4.4.5. Oscillator Development—Generation of Temporally Smooth Broad Bandwidth Laser Pulses via Cross-phase Modulation in an Optical Fiber

Summary. In 1989 we developed and implemented a novel concept in broad bandwidth laser light generation which employed the phenomena of optical field cross-phase-modulation (XPM) in a single-transverse-mode optical fiber.^{79,80} The impetus for developing this technique arose from our desire to have a reliable, variable bandwidth oscillator source at 1.053 μm with a nominally “smooth” temporal envelope that could easily be pulse-shaped and propagated down the Nova chain with minimum damage risk at typical operating fluences. Such a bandwidth source is now envisioned to be used with other system components such as reflective or transmissive “echelons” and/or phase-scrambling optics (e.g., distributed phase

plates), so as to smooth the Nova irradiance distribution at target.

In addition, we constructed a stable, electro-optically Q-switched Nd:glass laser oscillator with 10–15 cm^{-1} bandwidth. This oscillator acts as the pump source for the statistical intensity fluctuations in the optical XPM interaction. The “noisy” pump source was co-propagated with a temporally smooth single-mode, single-frequency laser pulse, along orthogonal polarizations of a birefringent fiber. When the fiber length was kept short (<5 m), the single-mode pulse acquired a statistical phase noise and a spectral bandwidth up to 15 cm^{-1} , without incurring any discernible temporal noise. Stimulated Brillouin and Raman scattering, and group velocity dispersion (GVD) effects were avoided at the shorter fiber lengths tested. With longer fibers, up to 20 m, bandwidths as large as 40 cm^{-1} were generated, but GVD effects caused a measurable degree of temporal noise and an asymmetrical power spectrum. The optical XPM technique is not limited to statistical pump sources. With a narrowband, temporally shaped pump source or multiple sources, it should be possible to induce a power spectrum of almost any format on a single-mode laser pulse.⁸¹

The experimental layout for implementing statistical XPM is shown in Fig. 4-71. As mentioned, two laser sources are required: a single-mode, single-frequency Nd:YLF oscillator (so called long-pulse oscillator, linewidth $\Delta\nu \leq 100$ MHz); and a broadband Nd:phosphate glass oscillator (line width $\Delta\nu \sim 450$ GHz), both operating nominally at a 1.053- μm wavelength. The temporal structure of the glass oscillator is dominated by intensity spikes with time widths on the order of $1/\Delta\nu$, which was verified by analyzing time traces from a 2-ps resolution streak camera. A square pulse of ~ 15 -ns duration was electro-optically sliced out of the 100-ns Nd:glass laser pulse, and combined at the face of the optical fiber with a ~ 10 -ns square temporal slice from the single-frequency long-pulse oscillator. The Nd:glass laser pulse serves as the “pump” for the XPM interaction and is coupled into the fiber with its linear polarization axis carefully aligned along one of the fiber polarization axes. The single-frequency Nd:YLF laser pulse or “signal” is centered in time within the “pump” envelope, and is coupled into the fiber along the orthogonal polarization axis. XPM results when the intensity fluctuations on

Fig. 4-71. Experimental layout for cross-phase-modulation (XPM) of single-mode (LPO) and multi-mode (BBO) laser sources in a single-mode birefringent optical fiber for generation of temporally smooth wide-spectral-width pulses.

the “pump” which typically exceed a peak to average ratio of ~ 8 , produce an intensity-dependent refractive index modulation along both the “pump” and the “signal” polarization axes. The refractive index modulation produces a time-varying phase modulation of the “signal” pulse envelope, resulting in significant spectral broadening. The “pump” spectrum also broadens, by about $2 \times$ in width, via the process of self-phase-modulation (SPM). The pump and signal beam are spatially separated by a high rejection ratio polarizer located after the fiber optic output coupling optics.

Typical spectra are shown in Figs. 4-72a–d for a “pump” energy of $\sim 5 \mu\text{J}$, measured at the fiber output face, and a fiber length of 12 m. The average pump intensity was estimated at $\sim 1.4 \text{ GW/cm}^2$ in the $4 \times 6 \mu\text{m}$ core of our standard fiber, giving an average phase modulation depth of 2.2 rad/m . Figure 4-72a shows the single-mode spectrum ($<0.5 \text{ \AA}$) prior to XPM, limited by the resolution of the optical multichannel analyzer and grating spectrometer. With XPM from the pump pulse, Fig. 4-72b, the single-mode linewidth broadened symmetrically to $22.7 \text{ \AA}$ FWHM. The “pump” spectrum increased from $14.3 \text{ \AA}$ FWHM, shown in Fig. 4-72c, to $23.5 \text{ \AA}$ FWHM, shown in Fig. 4-72d, as a result of SPM in the fiber. We typically observed a central “spike” in the XPM spectrum (Fig. 4-72b) located at the wavelength of the unmodulated signal pulse. This “spike” was substantially reduced by adding a second output polarizer to the long-pulse oscillator pulse-slicing Pockels cell, and hence we surmised that the central frequency component was unmodulated light leaking through the slicer.

When the pump energy (as measured at fiber output) was doubled to $10 \mu\text{J}$, we observed that both the signal and pump spectra developed large asymmetries. Typical spectra are shown in Figs. 4-73a and b. The signal spectra, Fig. 4-73a, appears to develop a red wing, with an overall broadening towards the blue wavelength side. The pump spectra, Fig. 4-73b, develops a blue wing with an apparent increased broadening on the red side. In more typical operations, Fig. 4-74 shows the XPM induced bandwidth on the signal pulse as a function of the pump energy for several fiber lengths tested. These results apply to a fiber with a low value of polarization velocity mismatch (estimated from manufacturers specifications to be $\sim 0.2 \text{ ps/m}$). A second fiber, one commonly used for optical pulse compression

work, was also tested at a length of $\sim 17 \text{ m}$ and gave very poor results. An XPM spectral width of only $1.5 \text{ \AA}$ was obtained with $8 \mu\text{J}$ of pump energy. The polarization walk-off for this fiber was estimated at 1.1 ps/m , some $5 \times$ greater than our standard fiber.

Detailed numerical modeling of XPM and SPM in a birefringent optical fiber is now being performed. Our model includes the polarization velocity mismatch between the fiber birefringence axes, and includes group-velocity-dispersion effects on temporal reshaping. Preliminary simulations indicate that the spectral asymmetries shown in Figs 4-73a–b are at least partially caused by a temporal walk-off, or more precisely, a walking-ahead of the pump pulse relative to the signal pulse. An observation of stimulated Raman scattering (430 cm^{-1} shift) on the pump pulse was also made at the higher pump energies and longer fiber lengths, and could also account for asymmetrical spectra. With a large temporal walk-off typical of the pulse compression fiber mentioned above, simulations show that the XPM interaction is almost completely washed-out. Each time-point on the signal pulse envelope samples a wide variation in pump-induced phase shifts due to the relative travel of the pump envelope and hence, over the entire fiber length the net phase-shift for time points on the signal pulse average to near zero.

Numerical simulations also qualitatively agree with our observation of increased temporal noise on the single-mode pulse envelope, with the 20 m fiber length in particular at $>5 \text{ mJ}$ pump energy. With large pump-induced phase modulation per fiber length, the GVD of the fiber causes temporal reshaping or self-focusing of the signal envelope. Time points on the signal envelope that have a large phase derivative or, in other words, a large instantaneous wavelength shift will propagate at a different phase velocity down the fiber. Hence, phase modulation will be converted to amplitude modulation on the pulse envelope. Experimentally, no obvious spatial self-focusing was observed in the fiber even with large XPM spectral widths, as we observed that the single-transverse mode quality and diameter of the signal beam did not vary.

The XPM bandwidth system design is presently being finalized and plans are being made for installation into the Nova master oscillator system. We anticipate using the XPM

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Fig. 4-72. Typical wavelength spectra of the single-mode and broadband oscillator pulses: (a) single-frequency, single-mode pulse before entering optical fiber (linewidth $< 0.5 \text{ \AA}$; (b) single-mode spectrum after XPM in fiber (linewidth $\sim 22.7 \text{ \AA}$); (c) multi-mode pulse before entering fiber (linewidth $\sim 14.3 \text{ \AA}$); (d) multi-mode pulse after XPM in optical fiber (linewidth $\sim 23.5 \text{ \AA}$). Experimental conditions: multi-mode energy $\sim 5 \mu\text{J}$ pulse, power $\sim 330 \text{ W}$, peak intensity $\sim 1.5 \text{ GW/cm}^2$, fiber length = 12.2 m, fiber core = $4 \times 6 \mu\text{m}$.

Fig. 4-73. Typical wavelength spectra at high multi-mode pulse energy: (a) single-mode pulse spectra is shifted towards blue wavelengths (linewidth ~ 35 Å); (b) multi-mode pulse spectra is shifted towards red wavelengths (linewidth ~ 24 Å). Experimental conditions: multi-mode energy ~ 10 μ J, power ~ 660 W, peak intensity ~ 3 GW/cm², fiber length = 12.2 m, fiber core = 4×6 μ m.

Fig. 4-74. Measurements of XPM generated spectral width (full-width at half maximum) on single-mode pulse as a function of multi-mode (broadband oscillator) pulse energy coupled into fiber at fiber length of 2, 10, and 20 m. Same fiber type as used in Figs. 4-72 and 4-73.

system in the summer of 1990 as a variable bandwidth source for beam smoothing and broad bandwidth propagation experiments on Nova.

Author: M. Henesian

Major Contributors: P. Chaffee, F. Patterson, D. Browning, R. Wilcox, S. Burkhart, and H. Powell

4.4.6. Two-Beam Diagnostic System

We have designed, constructed and activated a full-aperture beam diagnostic on beamline 26 of the Nova two-beam facility. The activation of the diagnostic is an important milestone in our campaign to convert beamline 26 at the two-beam into a flexible facility for laser-plasma interaction studies relevant to LMF design. The two-beam laser diagnostic system (TBLDS) was developed specifically to provide detailed measurements of the full-power frequency-converted Nova beam at the target plan during experiments. On-line measurement of the space and time-resolved laser intensity at the target plan will provide target physicists with valuable information needed to understand absorption and scattering processes in the target and enable laser physicists to directly observe the effects of beam-smoothing techniques on the focal plane irradiance. Both groups can then directly correlate beam-conditioning with target performance.

The TBLDS was designed and built under a stringent set of guidelines dictated by experiment's needs and facility requirements (see Table 4-7). To provide accurate focal plane measurements during target shots, we needed to develop a precise equivalent plane camera. We divided the problem of equivalent plane imaging into two parts:

1. accurately reproducing the intensity pattern at chamber focus using the minimum number of 74-cm components, and

2. re-imaging the equivalent focus with a small-aperture, high-precision optical system.

Step 1 required a full-aperture beam sampler that would faithfully reproduce the spatial intensity profile and phase-front of the beam. Damage, wave front, and alignment issues led us to use a full-aperture beamsplitter located prior to the chamber focus lens for this purpose. The 82-cm diam, 6-cm thick bare fused silica splitter directs ~3.5% of the beam to a second Nova focus

lens. The focus formed by this lens becomes the object for our imaging system. The thickness of the beamsplitter was chosen to achieve the best compromise between reflected wave front aberrations due to optical figure and transmitted wave front aberrations resulting from B-integral. The beamsplitter blank was finished by Eastman Kodak Co. to produce reflected and transmitted wave fronts of better than $\lambda/4$ at 0.633 μm . By holding the splitter in a special two-point mount which simulated a V-block, we were able to maintain the manufactured wave front (see Fig. 4-75). The Nova focus lenses are manufactured to within $\lambda/6$ and we estimate that the equivalent focus is produced within an RMS wave front error $\sim\lambda/3$ at 0.528 μm .

A diagram of the diagnostic is shown in Fig. 4-76. A real equivalent focus, although bothersome in that it requires the use of a vacuum vessel, is important in that it physically separates the subsequent imaging system from the primary lens. The real focus allows us to evaluate the quality of our imaging system on-line by placing resolution targets at the focus, and provides wavelength separation, image relaying, field-of-view and other alignment benefits. In addition, vacuum loading the primary lens produces the induced birefringence of $\sim\lambda/2$ at 0.351 μm that is present in the target beam.^{81,82} The expanding f/4 beam is collected and re-collimated by a fused silica doublet consisting of an aplanatic/equi-convex lens pair, with the equi-convex element being used as a vacuum window to minimize both glass and birefringence. An uncoated mirror following this telescope directs the nominally 2-cm beam to a diagnostic module located at floor level while dumping unwanted beam energy. The bare mirror and beamsplitter are used at a 13 degree angle of incidence but are tilted in orthogonal plans to compensate for the ~10% difference in "s" and "p" reflectivities at this angle.

In addition to a focal-plane camera, the TBLDS module contains a complete set of fundamental beam diagnostics to measure energy, temporal pulse shape, near-field irradiance, phase, and frequency spectrum. The specifications for these measurements are listed in Table 4-8. Measurements can be made at the 1.054-, 0.528-, and 0.351- μm wavelengths, although not simultaneously. As an example of data we have obtained so far, Fig. 4-77 shows a high-resolution photograph of the full-aperture second harmonic beam, the first obtained on

Table 4-7. TBLDS Design Guidelines

- Operating wavelength; 1, 2 and 3 ω
- Dynamic range:
 - 1 to 100 mW CW (alignment)
 - 100 to 15000 J pulsed
- B-integral: minimize
- Energy measurement:
 - accuracy - 5% at target chamber center
 - calorimetry compatible with existing Laser Diagnostics electronics
- Pulse shape measurement
 - aperture averaged
 - spatial sampling uniform to 5% across aperture
 - temporal resolution: streak camera limited
 - data format compatible with existing camera networking and software
- Target plane imaging
 - resolution: 20 mm
 - field of view: ± 2 mm at target chamber center
 - Recording medium: film
 - exposure: 4 to 5 simultaneous, covering ~ 2 orders of magnitude
 - image size: comfortable for viewing, compatible with the PDS densitometer
- Near-field imaging:
 - resolution: mm at KDP plane
 - imaged plans: KDP plane ± 10 m
 - recording medium: film, CCD, CID
 - image size: see above
- Spectroscopy:
 - aperture averaged
 - resolution $< 0.5 \text{ \AA}$
 - recording medium: film, OMA
 - option for streaking
- Interferometry:
 - measure planes: KDP plane ± 10 m
 - resolution: 5 mm at KDP plane 1/10
 - recording medium: CCD, CID
 - calibrate
- Alignment:
 - active 1 ω alignment capability
 - tolerances: minimize
 - near-field viewing capability for centering
 - far-field viewing capability for pointing
 - 5 mR field - acquisition
 - 1 mR field - alignment
- Discriminate from back-scattered light
- Avoid laser-induced damage
- Avoid SRRS from atmospheric nitrogen in propagation paths
- Fit into existing facility with minimal structural change
- Locate downstream of phase-scrambling components (for beam-smoothing)
- Minimize number of full-aperture components to reduce cost and footprint
- Provide a back-scattered diagnostic beam

Fig. 4-75. The TBLDS 82-cm-diam beamsplitter mounted in its gimbal, (a) showing reflected and (b) transmitted wavefronts of better than $\lambda/4$.

Fig. 4-76. Optical layout of the TBLDS. A full-aperture beam-splitter provides both incident and back-scattered diagnostic beams. The equivalent focus formed by the duplicate primary lens is re-imaged in the diagnostic module located at floor-level beneath the telescope.

Table 4-8. TBLDS Design Specifications

- Measure the following laser parameters on a target shot:
 - energy
 - pulse shape
 - irradiance at the target plane
 - irradiance at selected near-field planes
 - frequency spectrum
- Measure 1, 2 and 3 ω (not simultaneously)
- Reliable calibration
- Polarization-insensitive sampling
- Minimize energy loss and aberrations to beam going to target
- Minimize glass to control non-linear propagation problems
- Easy alignment
- Discriminate from back-scattered light
- Avoid laser-induced damage
- Avoid SRRS from atmospheric nitrogen in propagation paths
- Fit into existing facility with minimal structural change
- Locate downstream of phase-scrambling components (for beam-smoothing)
- Minimize number of full-aperture components to reduce cost and footprint
- Provide a back-scattered diagnostic beam

Fig. 4-77. TBLDS photograph of a 1.5-kJ, 1-ns second harmonic beam at the plane of the primary focus lens. The fine x-hair in the photo is a temporary alignment aid located directly in front of the diagnostic primary.

Nova. The more sophisticated far-field and interferometry measurements are presently being activated.

The TBLDS is designed to be flexible and easily configured for different experiments. To configure the TBLDS for operation at a particular wavelength, the end of the vacuum telescope is translated along the beam propagation direction to position an aperture at the new focus, and the doublet translated an additional amount to recollimate the beam at the new wavelength. These adjustments are made using a standard Nova LSI-11 stepper-motor controller.⁸³ Additional positioning of the many components in the diagnostic module is made efficient through the use of multiple motor setpoints stored in the controller which define each sensor configuration. The TBLDS is fully integrated into the existing Nova alignment and diagnostic controls system.

The conceptual design for the TBLDS was completed and reviewed in December 1988. Construction started in October 1989 and the first shots into the diagnostic were conducted

successfully in early 1990. The data from the TBLDS will be used to characterize the phasefront, near-field uniformity, and focusability of the current Nova beam for comparison with existing data from opposed-port diagnostic⁸⁴ and Phase II experiments.^{85,86} In addition, the diagnostic will be used to characterize the frequency conversion of broad-bandwidth ($\sim 10\text{--}20\text{ cm}^{-1}$) pulses and the focusing properties of the full-aperture beam after implementation of smoothing by spectral dispersion (SSD)⁸⁷ with distributed random phase screen optical elements.^{88,89} The laser data provided by the TBLDS will be used to normalize the frequency conversion and diffraction codes that we have developed and thus will improve our ability to predict and tailor the laser irradiance presented to fusion targets.

Author: Paul Wegner

Major Contributors: M. Henesian and L. Seppala

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5 Target Science and Technology

Section editors: T. P. Bernat, R. C. Cook

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5. Target Science and Technology

5.1 Introduction

The charter of the Target Science and Technology (TS&T) Program is to provide targets for a wide variety of experimental program activities such as:

- Fusion experiments
- X-ray lasers
- Laser matter/plasma interaction studies
- Materials at extreme densities and/or temperatures
- Instabilities at high accelerations
- Laser driver characterization
- Diagnostic development

To satisfy this charter, the TS&T program maintains a multidisciplinary capability which includes all of the following:

- Fabrication-production
 - Micromachining
 - Microassembly
- Development
 - Coating processes
 - Engineering/environmental supports
 - Alignment systems
 - Transportation systems
- Characterization
 - Physical measurements
 - Chemical analysis
- Material research
 - Material properties at low

temperatures

- New materials
- Environmental stability

In the following sections we summarize a few of the development and fabrication activities pursued during this reporting period.

5.2. Measurements, Modeling, and Fabrication

5.2.1. Microchannel Plate Imager Coating

Coating techniques developed by Target Science and Technology to supply targets for laser fusion experiments have been used to coat x-ray imaging microchannel plates (MCP). By doing these coatings "in house" we are able to

improve quality control and performance through continued development work.¹⁻³

The majority of MCPs are utilized as multiple-frame gated x-ray imagers that image target implosions. MCP's are SiO₂ wafers 50 cm in diameter, 0.5 mm thick with a honeycomb of 10–20 μ m diam channels through the plates. The rear side has a metal coating over its entire surface forming a ground plane. The front has its coating in the form of a stripline. An electrical pulse (100 ps) acts as a fast shutter activating different areas of the plate over time as the pulse propagates along the stripline. As an incident x-ray photon strikes an active area of the MCP, electrons are emitted. These electrons, seeing a high potential difference relative to the ground plane, accelerate through the channel plate. The channels are set at an 8° bias angle allowing these accelerating electrons to strike the walls of the channel producing more electrons. From the output side of the MCP the amplified cloud of electrons is accelerated onto a phosphor and the resulting light recorded. (Fig. 5-1)

MCPs have a limited lifetime and when a new plate is needed, a mask of the desired stripline configuration is chosen and mounted over a blank plate inside an E-beam evaporator. To prevent deposition inside the channels the plates are mounted at a steep angle relative to the evaporation hearth. There are two mounts inside the coating chamber, one at a 15° angle and the other a 45° angle. Both mounts are located 40 cm above the hearth. The selected mask and blank MCP combination are placed in the 15° mount. A second plate is mounted without any masking in the 45° mount to form the ground plane.

In a typical coating sequence the plates are brought under vacuum and baked at 220°C for 2 hours to remove residual moisture inside the channels. The plates are allowed to cool, then 5000 Å of copper (Cu) are deposited immediately followed by 500 Å of gold (Au) to prevent oxidation of the Cu.

When this procedure is completed, the two plates are turned over and their mounting positions swapped. A different mask may be

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Figure 5-1. Schematic diagram of a microchannel plate.

selected as required, and the above procedure repeated resulting in two complete plates. Completed plates are then placed in a dry nitrogen atmosphere for storage until needed.

To overcome the gain loss of the Cu/Au coating, cesium iodide (CsI) is used to enhance the quantum efficiency of the MCP. Prior to installation in the imaging system, the stripline surface is coated at a near normal angle of incidence to the 8° bias angle with one μm of CsI in a thermal evaporation system. This allows coating farther down the channel wall with concurrent increase in gain caused by the higher secondary emission coefficient of the CsI compared to the glass of the MCP. (Fig. 5-2)

In addition to microchannel plates, the CsI coating capability has been expanded to include photocathodes for x-ray streak cameras.

The articles in this section describe some of the work we have been doing on Rayleigh-Taylor instability physics.

Authors: M. D. Saculla and R. J. Wallace

5.2.2. A Unified Integrated Model For Sputter Coating Uniformity

Introduction. The coating flux received by substrates in a sputter coating process can be thought of as consisting of two components: a diffusive flux, contributed by fully thermalized atoms, and a directed flux, contributed by atoms which have retained a major part of their original velocity (i.e., both speed and direction). Near the source, the directed flux prevails, whereas far away from the source, the diffusive flux dominates. The sputtered source atoms undergo collisions with the inert gas atoms and become thermalized. The fraction of thermalized atoms increases with distance from the source.

The two different flux components have often been modeled separately. Such models work only in the extremes, and are not useful in situations where both mechanisms contribute significantly. We have developed an integrated model of the process that assigns weights to the two mechanisms; the weighing factors are then related to the process and equipment parameters.

Modeling. The initial directional flux can be described by simple, line-of-sight models.⁴ As the sputtered atoms are thermalized by the sputtering gas (typically Ar) their net movement can be described by diffusion equations. The transition from the directed to the diffusional

flux is gradual. Thus, in a region where the contributions from the two flux mechanisms are of the same order of magnitude, neither of the two extreme models (i.e., directional and diffusional) will be adequate to explain the data and predict trends.

Our model takes into account the gradual shift from one mechanism to the other in a continuous manner. On the basis of the assumptions listed in Ref. 5, we can use the steady-state diffusion equation to describe the transport of the thermalized atomic flux:

$$\frac{d^2C}{dz^2} + \frac{1}{r} \cdot \frac{d}{dr} r \frac{(dC)}{(dr)} = 0 \quad (1)$$

where C represents the concentration of the diffusing species, i.e., metal atoms r represents radial distance from the center line and z represents distance along the center line with the boundary conditions of a specified diffusional flux, $f(r)$, at the chamber entrance, and zero concentrations at all exposed surfaces. We solved these equations using the Topaz 2-D finite element program.⁶ We have shown earlier that the coating uniformity on a small substrate is insensitive to the initial flux distribution. For the purpose of this investigation, we imposed a uniform flux at the chamber entrance.

The calculations of the directed flux distribution are straightforward: the top of the substrate facing the target receives the full directed flux (corrected for surface orientation), whereas the bottom part, being in the shadow, gets no flux.

Let ϕ be the fraction of thermalized atoms at any position z. Let m be the mean free path of the atoms at the chamber pressure and temperature, and n the expected number of collisions undergone by the metal atoms at the position z. Then, following the treatment given by Ecker et al.,⁷ and Townes,⁸ we can write:

$$\phi_z = 1 - e^{-k(z/m)} \quad (2)$$

where k is a phenomenological parameter describing the thermalizing efficiency of collisions.

Let f_0 be the initial directed flux, f_{dr} and f_{th} be the directed and thermalized flux contributions at position z, and f_n be the net flux

Figure 5-2. Completed microchannel plate with a serpentine stripline.

at position z . All are at $r = 0$. Then, following assumption 8, we can write

$$f_n = (1 - f) f_{dr} + f f_{th} \quad (3)$$

f_{th} is numerically calculated from Eq. 1.

Equations 1, 2 and 3 completely define the model. The mean free path can be estimated from kinetic theory.⁵ We define, for the purpose of this model, the quantity "flux uniformity" (or simply "uniformity") as:

$$\text{Uniformity} = \frac{f_{n,b}}{f_{n,t}} \quad (4)$$

where the subscripts b and t denote bottom and top, respectively.

Algebraic manipulations of Eqs. 3 and 4 yield the following expression for flux uniformity:

$$\text{Uniformity} = \frac{\phi_b f_{th,b}}{\phi_t f_{th,t} + (1 - \phi_t) f_0} \quad (4a)$$

$$\frac{(1 - e^{-jz_b/m}) f_{th,b}}{(1 - e^{-kz_b/m}) f_{th,t} + e^{-kz_t/mf_0}} \quad (4b)$$

Experiment. The design of our experiment was to determine the diffusive and directed fluxes under normal sputter deposition. These data are needed to confirm the validity of our model, and further refine our model for more accurate predictions. If we use a cylindrical object as the substrate, the directed flux will deposit the thickest coating on the top surface of the substrate facing the target. The shadowed surface should receive no coating if the coating flux consists of 100% directed flux. On the other hand, when the sputtering pressure is in the tens of millitorrs, the diffusive flux will become the dominant component as distance from the sputtering target increases. When the thermalization process is complete, the substrate sees no sputtering source. Under that situation, the coating thickness on a round substrate will be uniform. Our modeling work aims at conditions between these two extremes. By measuring the coating thicknesses at the top and bottom

surfaces of a round substrate, we can empirically determine the composition of the coating flux.

We used a 1/8 in. glass rod as substrate. A schematic of the experimental set-up is shown in Fig. 5-3. The sputtering chamber has a maximum sputtering distance of 12 inches. We used a 3-in. sputtering ring gun and a copper target. Copper can be scraped off easily from the glass rod and film thickness is determined by a profilometer. All depositions were made at 20 mTorr and room temperature. Where the mean-free path is 0.03 cm we varied the substrate to target distance (d) while holding the ground plane to target distance (L) at 12-in. constant. At each value of (d), we monitored film thicknesses at locations 4, 5 and 6 of the glass rod. Two other spots, 1 and 3, on the ground plane were used as coating reference points to monitor the consistency of the coating runs.

Results. As expected, the coating thickness ratio between locations 6 and 4 decreases as the substrate is moved down toward the ground plane. Both flux rates drop as d/L is increased. However, the ground plane has a stronger shadowing effect on location 6 than 4 because the latter depends exclusively on diffusional coating flux. Thus the flux rate ratio decreases also. As the substrate is moved closer to the source, thus decreasing the value of d/L below 0.5, the flux ratio decreases rapidly. The flux rate ratios are plotted in Fig. 5-4.

We have estimated the constant k , defined in Eq. 2, from our data to be 0.006. We adjusted the value of k until we got a satisfactory fit between the data and the model calculations.

The model predicts correct trends over its entire range. Also, the model predicts correct uniformity values for d/L between 0.125 and 0.75. As the flux decreases away from the target, and as the influence of the ground plane increases, the errors in the data as well as the diffusional model increase. This may explain the disagreement between the model and the data for d/L greater than 0.75.

This modeling helps us realize the effects of ground plane on sputter coating. The coating results would be different if a microsphere is supported on a stalk instead of a microscope slide during sputter deposition. A calibration run can be meaningful only if the coating configuration is similar to the actual coating condition.

Authors: R. S. Upadhye and E. J. Hsieh

Major Contributors: E. L. Pierce and M. K. Kong

*Numbered thickness monitoring locations

Figure 5-3. Experimental configuration.

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Figure 5-4. Model and data comparison.

5.2.3. Plasma-Assisted Coating Technology

CH coating is required in the construction of ablatively driven laser fusion targets. At LLNL we have excelled in providing ultrasMOOTH, thick coatings. However, uncertainties of exact chemical composition lead to problems in modeling the radiatively driven ablation of the coating. We are therefore working to better understand the details of our coating chemistry and deposition processes. The CH plasma coating project has operated with little modification or addition to the technology for several years. To improve the reproducibility of coater operation we focused our effort on: improving the mechanical systems of the coaters, analysis of the coatings, and characterization of the plasma generator.

Relatively little is known about the structure of plasma polymers since they do not have a regular repeat unit and are insoluble in most solvents, and hence not amenable to many traditional chemical analytical techniques. This year we started an experimental program to permanently define the composition and physical properties of the plasma polymers. Early experiments using infrared spectroscopy showed by the appearance of an OH stretch and a C = O stretch that the plasma polymers appeared to react in air. We next followed the change in mass of several films deposited on aluminum foil on exposure to air. Over a three-month period the coatings increased by 3% in weight. Heating the coatings in a vacuum oven later did not reduce the weight. We believe at this time that molecular fragments (free radicals) trapped in the coating react with either oxygen or water vapor incorporating oxygen chemically into the polymer structure. We have begun to develop other techniques for following the coating reaction including near-infrared spectroscopy and chemical analysis. Chemical analysis has shown that the stoichiometric formula for the plasma polymer is CH_xO_y , where x is from 1.2 to 1.4 and y is a function of time ranging from 0.03 at 1 month to 0.06 at 3 months. An exact analysis of the CH material is complicated by variations in the composition with position in the coater. The microshells typically are coated in a small region of the coater, but additional coating is required to perform a chemical analysis.

In the last year we have also characterized the rate of deposition of CH coating as a function

of gas flow rate using a laser reflectance technique. We found that the coating begins to deposit immediately after the discharge is started. We also characterized the distribution of coating rate as a function of radial position in the coater. The rate drops from a maximum at the center line to 0.5 of the maximum at a radius of 1.3 cm from the center line. This information is needed to determine the effect of shell bouncing conditions and position on the coating rate.

Perhaps the least understood component of the plasma coater is the helical resonator that generates the discharge. The resonator consists of a helix of a 1.98-mm copper wire of 28 turns grounded at both ends surrounding a quartz tube that carries the reactive gases. The entire apparatus is enclosed within a cylindrical metal container that is capacitively coupled to the coil. The coil is electrically driven at usually the fifth winding from the bottom. The drive electronics consist of an oscillator driving a broad band amplifier connected to a pi matching network. This electrical network has many variables that we have gradually and empirically learned to optimize for producing smooth coatings. This year we began measuring the impedance of the helical resonator as a function of frequency in an attempt to better characterize the discharge device. We have now documented the impedance characteristics of each resonator. The recorded impedance has been useful in resetting resonators after disassembly for cleaning or electrical repair. Previously, several runs lasting 3 to 4 days each were needed to regain normal operating conditions. We are now looking at the effect of drive frequency on the operation of the resonator. We have found that operating off resonance by $\pm 5\%$ results in a drop in coating rate of 50%. In the future we hope the impedance test will help us set the coaters immediately to optimum performance without the need for months of empirical system adjustment.

Author: S. A. Letts

5.2.4. Low-Density Hydrocarbon Foams

Introduction. We have been developing foams and fabrication processes for a direct-drive, high-gain ICF target. The design calls for the foam to hold the DT fuel in a spherical shell configuration in the target, which is then compressed to high density and ignited by the laser energy. The foam should have a density of

50 mg/cm³ or less and be composed of carbon, hydrogen, and oxygen. In addition, the foam should have a cell size of 4 μ m or less and a pore size of 1 μ m or less to stably hold the liquid DT. The foam must be dimensionally stable at cryogenic temperatures and be chemically stable in the radiation environment of the liquid DT. The foam must be machinable or moldable to the required dimensions and tolerances. Some designs also call for an overcoating of a few μ m of a low Z material.

In 1989, two foam materials were under development: polystyrene (PS) foams made from an inverse emulsion system; and carbonized resorcinol-formaldehyde (CRF) foams made from condensation polymerization gels of resorcinol and formaldehyde, which are dried to form resorcinol-formaldehyde (RF) aerogels and then carbonized to produce the final CRF foam. Because of shifting priorities our efforts on these materials were scaled back toward the end of 1989, and experimental work terminated early in 1990. We successfully completed our development of the PS foam, producing materials with densities as low as 25 mg/cm³ with 2- to 3- μ m cell sizes. Our development of CRF foams was almost completed, though some concerns remain with respect to density uniformity and reproducibility. Background information and the results of previous development work can be found in Ref. 9-18.

Polystyrene Foams. Our PS foam development was completed in 1989 with the perfection of techniques to lower the foam density to 25 mg/cm³ while maintaining the cell size at 2-3 μ m. This was accomplished by adding a non-reactive filler, ethylbenzene to the oil phase, thus reducing the relative amounts of styrene and divinylbenzene. Because of the structural similarity between styrene and ethylbenzene, we expected that the formation of the inverse emulsion would be unaffected by this substitution. This appeared to be the case, as we were able to produce small-celled emulsions. Polymerization of these emulsions was accomplished in the standard way, but then prior to drying, a 50/50 isopropanol/water mixture was forced through the wet foam. The mixture washed out not only the inorganic residue from the initiator and most of the surfactant, but also the inert ethylbenzene.

The foam could then be dried in a drying oven. The resulting foam was not only a chemically pure CH material with 2- to 3- μ m

cell size, but also could be produced with densities as low as 25 mg/cm³, depending upon the level of ethylbenzene used. As expected, as the level of ethylbenzene was increased the morphology of the resulting foam became more lacy until finally the material lost structural integrity at densities less than 25 mg/cm³.

A U. S. Patent application detailing this significant modification of the existing technology for producing low-density inverse emulsion polystyrene foams has been submitted.¹⁹

Carbonized Resorcinol-Formaldehyde Foams. During 1989, most of our effort was focused on understanding and improving the carbonization process with a goal of optimizing foam properties. Our primary interest has been the dependence of foam properties on the carbonization temperature. Figure 5-5 shows the results of a thermogravimetric analysis of the RF foam. The technique measures mass loss as a function of temperature at predetermined temperature ramp rates. In this experiment samples of RF foam were heated in an inert atmosphere at a ramp rate of 5°C/min and then held at specific temperatures (shown on Fig. 5-5). Two features are worth noting. First, the carbonization process follows the temperature closely at this ramp rate since very little additional mass loss is noted when the temperature is held constant. Secondly and more important, is that the carbonization process, at least as measured by mass loss, is essentially complete by 700°C. At this temperature the mass loss was about 50%. Samples of foam were carbonized at temperatures of 400 to 1000°C in 100°C increments with resulting densities roughly the same for all carbonization temperatures. The density dropped from a value of 32 ± 1 mg/cm³ for the uncarbonized RF foam to 25 ± 3 mg/cm³ for the foam carbonized at 500°C or higher. There was a steady increase in the modulus up to about 700°C. The value plateaued here at about 0.4 MPa. Ramp rate and inert gas pressure in the carbonizing furnace had no apparent effect on the resulting foams.

Studies showed that the foams carbonized at 500°C were difficult to machine since the dust from the machining process tended to agglomerate on the grinding tool, making it difficult to maintain and/or repeat a small tolerance contour. The 1000°C CRF foams were found to be easier to handle, machine, and inspect.

Figure 5-5. TGA scans of eight RF samples ramped to the indicated temperatures at 5°C/min and held at that temperature.

One remaining difficulty with the CRF foams centers on the issue of density homogeneity. There are often significant differences (up to 20%) in the densities of CRF foam pieces taken from different batches, from within the same batch, and often from within a single piece. Part of the source of these variations may be the difficulty in measuring the volume of sometimes irregularly shaped pieces, but some of the apparent variation is real. Should this work be continued, the source of these variations will have to be understood and controlled.

As part of our investigation of the factors that control the microstructure of RF and CRF materials, we used viscometry to measure the effect of temperature on the rate of gelation of the resorcinol-formaldehyde reaction mixture.²⁰ The viscosity of the reaction solution rises rapidly in the first 30 minutes to a plateau value of 20 to 40 poise where it holds for as much as 24 hours until gelation occurs, which is marked by a rapid increase to several thousand poise. As expected, raising the cure temperature accelerates the gelation process.

Work on the carbonized composite polystyrene/resorcinol-formaldehyde (CPR) foam, described in the 1988 Laser Program Annual Report, has also been terminated with our change in priorities. However an application for a U. S. Patent on the material has been filed.²¹

Authors: R. C. Cook, S. R. Buckley, B. L. Haendler, F. M. Kong, S. A. Letts, and G. E. Overturf III.

Major Contributors: C. L. Giles and J. D. Aldridge.

5.2.5. Production of Spherical Shells for Silica

Introduction. Single shell ICF designs, under certain conditions, require that the fuel be levitated as a cryogenic shell a few hundred microns thick.²² Maintaining such thick liquid DT layers is quite difficult, since the liquid fuel slumps under the influence of gravitational forces. A number of approaches to this problem have been suggested.²³⁻²⁵ Sacks and Darling²² have suggested the use of a small pore size, open cell, rigid foam structure as a sponge to precisely define the layer contours and to stabilize the liquid against gravitational slumping. Both inorganic (e.g., silica, alumina, borosilicates)

and organic (e.g., polystyrene, carbonized resorcinol-formaldehyde) materials are potential candidates for the foam matrix for ICF targets. This report describes the recently concluded project to make the foam targets from silica aerogel, conducted jointly by the Laboratory and Airglass, a commercial enterprise in Sweden.

Although the basics of making low-density silica aerogels have been well-known for many years,²⁶⁻²⁸ the technology to fabricate ICF targets from this material was not developed, and constituted the subject of this investigation. Several approaches can be taken to make hollow microspheres. One of them is to use the droplet technology developed at the Laboratory.²⁹ This approach was not directly applicable to silica aerogels, since the material is quite different from organic polymers, and since any vapor-liquid interfaces generated at the alcogel boundary tend to shrink and damage the resulting aerogel.

It was, however, possible to modify the droplet tower technology for this project. The essence of the implemented technique is to use specially designed nozzles to blow microshells from a partially gelled solution of TMOS (tetramethoxy-silane), methanol, water, and a small quantity of ammonia (as a catalyst for the hydrolysis and gelling reactions). The microspheres are separated from the nozzle by a gas jet, and carried through the gas stream containing ammonia, which serves as a catalyst. These spheres are then collected over a foam bed, washed and cleaned, transported to the oven, and supercritically dried.

Process Chemistry. The chemistry of aerogel formation is quite complicated, and not fully understood. Excellent reviews of the subject are given by Iler²⁶ and Schmitt.²⁷ The following brief treatment is presented for the sake of continuity and completeness.

Silica aerogel can be made by a variety of techniques, described by Iler.²⁶ The technique we used in this work is described by Henning;²⁸ An orthosilicate is hydrolysed according to the following overall reaction:

where R represents an alkyl group, such as methyl or ethyl. (The present process uses tetramethoxysilane, or TMOS).

The silicic acid undergoes the following overall reaction:

which is followed by gelation:

The actual reaction mechanism is a function of many variables, such as the alkyl group, the nature and concentration of the catalyst, and the solution pH. Also, the hydrolysis and gelation reactions can take place in parallel, rather than in series, as implied above. The above reactions are meant to provide just stoichiometric information; neither the kinetics nor mechanistic information are being conveyed here.

The hydrolysis reaction can be catalyzed by either acids or bases. The mechanism and the kinetics of the reaction are a strong function of the nature and concentration of the catalyst. Airglas has found that ammonia is a desirable catalyst,²⁸ for TMOS hydrolysis, whereas Schmitt found nitric acid to be an acceptable catalyst for TEOS (tetraethoxysilane) hydrolysis. Similarly, either a base or an acid will catalyze the gelation reaction. For both the TMOS and the TEOS based processes, ammonia appears to be the preferred catalyst.^{27,28} The recipe for making the gelling solution is given by Henning.²⁸

Droplet Generation. The process of stable, collectible droplet formation imposes two conflicting requirements on the design of the equipment: The time spent in free fall should be as long as possible. This allows time for the hydrolysis to proceed as far as possible, thus making the shells stronger and better able to withstand the final impact upon collection.

The time spent in free fall should be as small as possible. This keeps the droplet velocity low, and allows the shells to land more softly. The key to the resolution of this dilemma is to recognize that an extra degree of freedom is available to us through the control of the concentration of ammonia in the gelling solution, and in the gas entering the inner nozzle and the reaction chamber. This allows us to control the reaction rates for hydrolysis and gelation, and thus to select the time of free fall through the reaction chamber which strikes a reasonable

balance between the conflicting requirements cited above.

The droplet generation system consists of three major components: the nozzle assembly, together with the flow controllers for the gelling solution, nitrogen, and ammonia; the reaction chamber, together with the ammonia flow controller; and, the collection mechanism. **Figure 5-6** shows a schematic of the droplet generation system, the reaction chamber, and the collection mechanism.

The nozzle assembly used for the droplet generation comprises two nozzles. The gelling solution flows through the outer nozzle, and a mixture of nitrogen and ammonia gas flows through the inner nozzle. The purpose of the inner nozzle is to facilitate the creation of a gas bubble which is trapped by the gelling solution flowing in the annulus between the inner and outer nozzles. Ammonia is added to the inner gas bubble to act as a catalyst to speed up the gelling process. The ammonia helps form the inner skin and stabilizes the inner bubble.

The inner nozzle can be moved up and down during the run to control the cross sectional area available for the flow of the gelling solution. In addition, a variety of tips can be selected for both the outer and the inner nozzles, which gives us additional degrees of freedom for adjusting and controlling the liquid and gas flows.

A transverse jet of nitrogen is used to strip the droplet from the tip of the nozzle, that minimizes the formation of a tail. It also helps deflect the droplet from the nozzle, thus reducing the collision probability of the droplets. It also imparts a tumbling and vibrating motion to the droplet which helps center the inner bubble. The position of the transverse jet nozzle, shown in **Fig. 5-6**, can be independently adjusted in the horizontal and vertical directions.

The droplet separated from the nozzle falls through a plexiglass reaction chamber about 1 m in height, and about 40 by 30 cm in cross section. The droplets are observed in this chamber with the help of a strobe light, which is the primary diagnostic aid during a run.

The flow rates of liquids and gases and the geometric dimensions, such as the positioning of the inner nozzle and the transverse jet, are adjusted on the basis of the information obtained through this observation. A high concentration of ammonia is maintained in the reaction chamber to aid the hydrolysis and gelling

Figure 5-6. The droplet generation system for silica aerogel microspheres.

reactions, and to stabilize the droplet by helping form a skin on the outside of the droplet. Collection of the soft, hollow alcogel spheres is a difficult problem.

If not properly done, the shells break or deform. Initial attempts to collect them over liquids such as methanol, ethanol, water or ammoniated water resulted in shell damage. The present method is to collect the shells over foam, which slows down the falling shell, and softens the impact. The original foams were made from a solution of water, methanol and soap. It was later observed that addition of a small amount of ammonia to the soap solution yields better results. The collected shells are moved to a ventilated bench, and a small quantity of methanol is poured over them to break the foam. It is important to remove the shells as soon as possible from the solution, and wash them several times with methanol to remove the soap and other contaminants on the surface. The alcogel shells cannot be dried by ordinary means (e.g., in air at ambient temperature and pressure), because the surface tension of the solvents (water, methanol) in the alcogel matrix tends to damage the pores progressively as the vapor-liquid interface recedes from the surface. Instead, the solvent is first taken to supercritical conditions, and then removed gradually from the matrix. A detailed description of the theory and practice of this operation is given by Schmitt.²⁷

Authors: Ravindra S. Upadhye.

Major Contributors: Charlotte King, LLNL; Anna-Greta Johansson, University of Lund, Sweden; and, Sten A. Henning, Airglass, Sweden.

5.2.6. Non-Destructive Fill Pressure and Permeability Measurement

Introduction. The fill pressure of diffusion filled plastic direct drive implosion targets has generally not been known accurately.³⁰ There are two major reasons for this; First, there is no practical nondestructive method of measuring the internal pressure after filling. This means that we have to assume that all the shells reach the external pressure (the fill pressure) within the allocated amount of time. Second, even if we were to know the initial fill very accurately, we do not know what the fill is at the time of the experiment, since up to several hours can elapse between target depressurization and implosion,

and the leakage rates are not known very accurately. The reported permeabilities of polyvinyl alcohol (PVA), which forms a permeation barrier for the shells, vary widely, and are very sensitive to trace quantities of water,³⁰ so that estimates of leak rates based on the accepted permeability values contain corresponding uncertainties.

Our recent work indicates that it is now possible to measure nondestructively both the fill pressure and leakage rate for any given shell accurately. We shall outline here our experiments, and describe the development work on a test station for routine characterization of shells. A more thorough description of the experimental work and an analysis of the data will be found in "A Non-Destructive Method for Measuring D₂/DT Fill Pressure and Permeability for Direct Drive Plastic Shells".³¹

Experimental Work. A schematic of the apparatus used in this experiment is shown in Fig. 5-7. The filled plastic shell is mounted in a sample holder and placed inside a cryostat, where it is cooled below the critical point of D₂ or DT. Once the temperature stabilizes, a heating laser is fired into the shell to heat it momentarily. A video image of the shell shows a distinct phase change if the D₂ is in the condensed state when the heating laser is fired. On the other hand, if the set temperature is too high, so that the D₂ inside the shell is not condensed at that temperature, no change in the video image is apparent upon firing the laser. We can successively approximate the temperature at which the D₂ liquifies by repetitious heating and cooling, until the dew point temperature is narrowed down to 0.1 K. We have found the final temperature to be consistently reproducible within 0.1 K, which represents the accuracy limit of this apparatus. The dew point temperature is then used to calculate the saturated vapor density of the appropriate isotope using a fit to data from Souers.³² The room temperature target pressure calculated from this density is then the target fill pressure.

Figure 5-8 shows the measured fill pressures of two different shells as a function of the elapsed time at room temperature. After making a measurement of pressure (as described in the previous section), we raised the shell to room temperature, held it there for a known amount of time, and measured the pressure again by freezing it. These repeated measurements gave us

Figure 5-7. Dew point test system schematic.

Figure 5-8. Pressure vs time history of shell 1 & 2.

the data shown in Fig. 5-8. The values of $t_{1/2}$ were estimated from the two data sets, and were found to be approximately 50 hours for both; this is indicated by the solid line in Fig. 5-8. That both the shells appear to leak at the same rate is not surprising, since the shell parameters (such as the inside diameter, and the wall thickness) are carefully controlled. The leakage rate coefficient (or, equivalently half-life) may depend on many variables, such as the thickness and uniformity of the PVA layer, local defects in the PVA layer, the thickness and quality of the outer coating, and ambient conditions, such as humidity and temperature. Changes in these variables may lead to large variations in the leakage rate measurements even on shells from the same batch.

The work presented here demonstrates that we will be able to use this new technique to obtain accurate fill pressure and leakage rate data, and as a quality control tool. As Fig. 5-8 shows, the leakage rate coefficient (or half-life) is constant over a period of time, and gas leakage from the shells is a first order process. Therefore, two or three measurements over a period of a few hours are sufficient to establish useful bounds on the leakage rate as a by-product of the quality control measurements. More data points over a longer time period will, of course, be necessary to obtain more accurate estimates of the leak rates. Furthermore, we have established that

repeated freezing and heating do not have a significant impact on the D₂ permeability of the shells (if this were not true, the curves of Fig. 5-8 would have shown significant deviation from linearity). This indicates that the cryogenic measurements are truly nonintrusive.

We are currently designing a system (comprising a cryostat, sample handler, temperature controller, vacuum pumps, a laser source and diagnostic optics) to implement this technique as a routine characterization and quality control tool. We would like the new system to be considerably simpler than the one on which this work was performed. The new system will have a faster cool-down rate, and will hold several shells at a time. The present accuracy in pressure measurement is limited mainly by the accuracy with which we can measure the dew point temperature. If we can measure the dew point to within 0.02° K, our pressure estimates will be within 2% of the true value over the entire pressure range.

Author: J. Sanchez.

Major Contributors: The shells used in this experiment were made by E. M. Lilley. The PVA and CH coatings were done by K. King and L. A. Witt respectively. The shells were characterized by C. King and C. Adam. They are all LLNL employees.

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6 Advanced Drivers

Section editors: W. H. Lowdermilk, H. T. Powell

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6. Advanced Drivers

6.1. Introduction

The objectives of the Advanced Drivers Program are to: (1) upgrade the performance capabilities of Nova; and (2) develop future drivers for the ICF Program. In 1989 we began planning the Precision Nova Project to improve Nova's capabilities for irradiating ICF targets. The principle objectives of this project are to improve energy and power balance among Nova's ten beams; improve pointing accuracy and develop an improved diagnostic package to measure beam energies at fundamental and harmonic wavelengths over the full beam aperture. Another project was begun to improve the uniformity of irradiation intensity delivered to the targets on Nova. The principle of such beam "smoothing" is to intentionally reduce the spatial and temporal coherence of the laser beam. The main activities within this project are: development of large aperture phase plates to provide spatial incoherence; and development of broadband, multifrequency laser oscillators to provide temporal incoherence.

In 1989, the effort on future ICF drivers was focused on two laser concepts intended for application to the Laboratory Microfusion Facility (LMF). These concepts are "Loki," a laser-pumped solid state laser and "Athena," a flashlamp-pumped Nd:glass laser. This year saw completion of our studies on Loki. Our studies showed that the Loki concept was unlikely to have a cost advantage over the more conventional Athena design, and it was therefore decided not to pursue this concept further. A major accomplishment in support of Loki was the invention and development of a new laser material--chromium doped lithium calcium aluminum fluoride (LiCaF). LiCaF has other potential applications, including ultra-high peak power lasers, for which a modest development activity will continue. The rapid progress made in improving the size and optical quality of this material was enabled by an ongoing program to develop an advanced crystal growth process. This program is supported by the Director's Initiative Fund.

The major development effort was devoted to the Athena laser. A multipass laser design was chosen for Athena because it has improved performance and lower cost than the master-oscillator-power-amplifier (MOPA) design of Nova and earlier glass lasers at LLNL. Another important characteristic of Athena is that the optical aperture is subdivided (segmented) into a 4×4 array. The aperture is segmented to prevent amplified spontaneous emission (ASE) losses in the amplifier and transverse stimulated Raman scattering (SRS) in the frequency converter crystals. Segmentation also allows optical components to be fabricated with reasonable aspect ratios (aperture to thickness ratio).

Athena can meet all technical requirements for the LMF, and was designed conservatively to provide a high level of confidence in its performance. All the components and materials required for Athena have been demonstrated subscale.

Athena development activities were focused on reducing costs significantly below those of previous Nd:glass laser systems used for ICF research. The principal ways to reduce cost are:

1. improve system efficiency,
2. increase energy storage density,
3. increase laser-induced damage thresholds,
4. improve engineering design and;
5. reduce manufacturing costs.

During this year our activities have been primarily in the first three categories. We continued developing the multisegment amplifier. This amplifier has higher efficiency and higher energy storage density than previous glass laser amplifiers. It is also a compact and modular design that further reduces its manufacturing costs. We continued developing optical materials and thin film coatings with improved damage resistance. We also continued a systematic search for frequency conversion materials with higher conversion efficiency. The development of a lower cost pulsed power system was continued, and we began development of an advanced alignment system to efficiently align the multiple beamlines required for a multimegajoule LMF laser.

The on-set of stimulated Raman or Brillouin scattering in optical components is a potential limitation in laser performance that impacts the laser design and must be quantitatively characterized. This year we made measurements of the stimulated gain coefficients in relevant materials, and made theoretical and experimental studies of the suppression of these effects by increasing the bandwidth of the laser pulse.

This year we completed the first level of design codes necessary to optimize the laser design for both performance and cost. These codes model all aspects of system performance and have been validated by comparison to the extensive database from Nova and other laboratory experiments.

We continued the Optical Materials Research Program, which is funded by the DOE Office of Basic Energy Sciences. The topics investigated this year include measurements of the nonlinear refractive index of optical materials, investigation of laser induced damage mechanisms and the studies of spectroscopy and energy transport dynamics of laser materials.

Author: W. H. Lowdermilk

6.2. Multi-Pass Laser Architecture

6.2.1. Introduction

To achieve our cost and performance goals for a megajoule-class ICF laser driver to follow Nova, we need to make substantial changes in the laser design and component manufacturing methods, as well as to develop a better understanding of trade-offs in both areas. During 1989 we had a small but active effort in megajoule laser technology described here. This effort included multisegment amplifier research and modeling, understanding the sources of optical loss in Nd:glass laser systems, conceptualizing new methods for energy extraction from laser amplifier chains, testing advanced pulsed power components to drive our amplifiers, and exploring new methods for aligning very large multibeam laser systems.

Our amplifier research centered around the construction and testing of a 2×2 aperture multisegment amplifier (MSA) where each aperture was roughly 30-cm square. Our MSA design departs significantly from past amplifier designs. Unlike previous amplifiers, it contained

one-half of its flashlamps in a region between the disks where they provided pumping in both directions. This increases the electrical efficiency by preventing loss of pump radiation. However, the most dramatic design departure of the MSA was its construction in a single-disk-long module. This construction method is readily extendable to the 4×4 amplifier size which we envision for megajoule systems. However, this design puts a large premium on joining these modules closely together to allow cross-pumping between modules in order to maintain high electrical efficiency. We have developed a very general model to describe flashlamp-pumped disk amplifiers which does a very good job of fitting and, to some degree, predicting the observed amplifier gains in both the single and multisegment amplifier cases.

A cost-effective system must provide an efficient energy extraction from these large amplifiers. A fundamental limit on energy extraction is imposed by optical losses at 1- μm in the laser disks. Hence we have studied the sources of such losses which include bulk absorption, surface scatter, as well as depolarization (since the polarization must be maintained in a Brewster disk system). We now understand the major contributors to the bulk loss, and with our glass manufacturers, have made significant strides to improve that figure. In terms of multipass architectures for energy extraction, we have explored two new concepts that utilize a small angular change to separate the input and output beam. Both concepts are passive and do not require an actively pulsed switch. Also both are compatible with a Brewster disk system.

We are also developing pulse power components for the large energy storage banks which will be necessary in a megajoule-class laser. Specifically, for economy we are trying to handle energy in very large units. For this we are developing, with the help of outside companies, large high-energy-density capacitor cans as well as very high-current switches. Our tests of prototypes in both areas are described.

Finally we are also considering new alignment concepts. Since the size of an individual amplifier subelement is limited to roughly 30-cm because of amplified spontaneous emission inside the disk, we recognize that a megajoule-class laser will have hundreds of apertures compared to Nova's ten (or 20 if one counts the 46-cm split disks as two). We are

evaluating new alignment techniques that can potentially align hundreds of beams in an acceptable time as well as accurately focus them on a target in the required positions.

Author: Howard T. Powell

6.2.2. Disk Amplifier Design and Performance

Introduction. In 1989, we made significant progress towards our goal of developing Nd:glass disk amplifiers for megajoule (MJ) fusion systems. We directed our efforts towards testing two different amplifiers: 1) the single-segment amplifier (SSA), which has been described in previous annual reports;¹⁻³ and 2) a new multisegment amplifier, which we designed and built with many features we have proposed for an MJ upgrade of Nova. As result of this work, we have greatly expanded the database we have used to develop a disk amplifier performance model (described in Article 6.2.3.). This model will be essential for optimizing the design of amplifiers for an MJ system.

All our previous SSA data were taken with disks made of Schott LG-660 silicate glass. This year, we measured the performance of the SSA with disks made of Schott LG-750 phosphate glass. Phosphate glasses generally have larger cross sections, higher gain, and greater ASE losses than silicate glasses. Our amplifier model predicts performance equally well for phosphate glass disks as for silicate glass disks.

We also measured, in absolute units, the spectra of pump radiation in the SSA. These measurements were made to improve our knowledge of pump cavity physics, and to test our amplifier model. Specifically, we used the measured spectra to calculate cavity transfer efficiencies, which are used in our amplifier model to characterize the transfer of radiation from the flashlamps to the disks. We found that the cavity transfer efficiencies determined from gain measurements using our model were 25% smaller (with uncertainty of $\pm 12\%$) than those calculated from the spectra. The reasons for this discrepancy are unknown, but the agreement seems remarkably good considering the differences in the two approaches.

In 1989, we also made the first measurements of the performance of a multisegment amplifier. MSAs have been proposed as a way of scaling Nd:glass disk amplifiers to very large size.⁴⁻⁶ The primary purpose of the measurements was to test the MSA pump-cavity geometry in

which central flashlamp arrays pump disks on both sides.

The test bed MSA was a one-disk-long module holding four Brewster-angle disks. It was designed to amplify a beam divided into four segments arranged in a 2×2 array, thus simulating one-fourth of the 4×4 MSA we proposed for a megajoule upgrade of the Nova laser. We used the amplifier model to design the 2×2 MSA to operate at high values of stored fluence per disk (stored energy per unit beam area per disk) at some expense in storage efficiency.

In the laboratory, the 2×2 MSA achieved 2.9% storage efficiency at 2 J/cm^2 stored fluence per disk with phosphate glass. This is the highest efficiency that we have achieved for an amplifier operating at this large stored fluence per disk.

The cavity transfer efficiency of the MSA was 13% higher than that of the SSA, when the two amplifiers were compared at similar pump cavity conditions. Measurements show that a major factor causing the improved MSA performance is the higher efficiency of central flashlamp arrays compared to side flashlamp arrays.

Single Segment Amplifier: Gain Measurements. The SSA was designed to amplify a square beam (with 32-cm sides) using four rectangular disks. Thus, it tests a pumping geometry similar to that of multisegment amplifiers that have been proposed for MJ lasers. The SSA was built (1) to generate a large database of amplifier performance so that an empirical amplifier model could be developed; and (2) to improve storage efficiencies over those achieved in previous amplifiers.

We have reported, in previous years, extensive gain measurements made on the SSA using disks made of LG-660 silicate glass.¹⁻³ This year, we report the results of similar SSA measurements using disks made of LG-750 phosphate glass. These measurements were used to develop the amplifier model, described in Article 6.2.3. that predicts amplifier performance for disks made of either phosphate or silicate glass.

Experimental. The SSA was loaded with four LG-750 phosphate glass disks that were 4.05 cm thick and doped with Nd at a concentration of $4 \times 10^{20} \text{ ions/cm}^3$. The measured fluorescence decay rate of these disks was 345- μs .

The disks used the epoxy-bonded glass edge claddings developed for Nova.⁷

SSA gain measurements were made using the same experimental apparatus described in past years. However, in this case, the gain coefficients inferred from the measurements were multiplied by a factor of 1.03 to correct for the slight wavelength mismatch between the probe laser (a cw Nd:YAG laser tuned to 1.052- μm) and the peak of the lineshape function at 1.053- μm .⁸ Stored energy densities were calculated using a recently measured value of the stimulated emission cross section for Nd in LG-750, $3.5 \times 10^{-20} \text{ cm}^2$.⁹ In calculating the energies stored in the disks, we took into account the measured variation in stored energy density as a function of position in the disks. Storage efficiencies were calculated by dividing the total energy stored in the disks by the total electrical energy delivered to the flashlamps.

Gain measurements were made using three of the combinations (number of lamps per side and reflector types) we used in our previous SSA tests: 3 lamps per side - Rabl reflectors (3R); 5 lamps per side - Winston reflectors (5W); and 8 lamps per side cylindrical reflectors (8c). The shapes of these reflectors are shown in Fig. 7-4 of Ref. 2. The Rabl and Winston reflector shapes were designed to minimize the amount of light reflected back through the lamps. For each reflector type, data were taken over a range of pulse lengths and pulse energies.

The flashlamps had 1.5-cm-bore diam and Xe fill pressures of 300 Torr. The flashlamps were preionized to improve storage efficiencies.

Measured Storage Efficiencies. The SSA data taken with the LG-750 disks and 5W reflectors are presented in Fig. 6-1a, in which storage efficiencies are plotted versus stored fluence per disk. Data taken previously with the LG-660 disks and 5W reflectors are presented in Fig 6-1b for comparison. Results are given for several different $3(LC)^{1/2}$ pulse lengths. (L is the inductance and C is the capacitance of the circuits driving the flashlamps.)

The curves drawn through the data are predictions made by our amplifier model using a single cavity transfer efficiency function that accounts for the transfer of radiant energy from the lamps to the disks. It is apparent that the model predicts the SSA performance equally well for either silicate or phosphate laser glasses.

Storage efficiencies measured using the 3R and 8c reflectors are presented in the article describing the amplifier model. Generally, we found that storage efficiencies measured with the LG-750 phosphate glass disks were comparable to, or greater than those measured with the LG-660 silicate glass, especially at low stored fluence per disk. This is due to the higher LG-750 absorption coefficients in the Nd-pump bands. However, as stored fluence per disk was increased, the storage efficiency fell more rapidly with the LG-750 glass than with the LG-660 glass. This was the result of the higher stimulated emission cross section, higher gain, and greater ASE losses in the LG-750 glass. These effects cause the differences in shapes of the LG-660 and LG-750 efficiency curves shown in Fig. 6-1a and Fig. 6-1b.

Single Segment Amplifier: Measurements of the Pump Radiation Spectrum. We measured the spectrum of radiation incident upon a disk in the SSA pump cavity in absolute units. These measurements were undertaken to improve our understanding of the physics of pump cavities, and to determine the effect of 1.05- μm fluorescence on depumping the disks. In the future, the data we generated can be used to validate ray-trace codes that model pump cavity performance.

We analyzed the pump cavity radiation using the same apparatus we previously used to analyze flashlamp radiation,¹⁰⁻¹² with some slight modifications. As shown in Fig. 6-2, a scatter plate sampled the radiation incident upon one of the interior disks in the SSA. Scattered light passed through a single hole in the side of the amplifier and was detected by two instruments, which were used sequentially. These instruments, apertured to detect only radiation from the scatter plate, were: 1) a Tektronix J-20 spectrum analyzer (that uses a silicon Vidicon detector at the output plane of a 0.25-m spectrometer) to measure the spectral distribution; and 2) a Molelectron P1-35 pyroelectric detector with flat spectral response to determine the time dependence of the intensity. The J-20 was gated (25- μs FWHM) with a rotating mechanical shutter to obtain time-resolved spectra. The spectra were corrected for the nonlinearity of the Vidicon response and for the wavelength dependent sensitivity of the instrument and its filters. The scatter plate was a $0.9 \times 4.8 \times 4.8$ -cm piece of

(a)

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(b)

05 00 0893 796 / A p.u

Figure 6-1. SSA storage efficiency plotted versus stored fluence per disk for: (a) 4.05-cm thick, 4%-doped LG-750 phosphate glass disks, and (b) 4.3-cm thick, 4%-doped LG-660 silicate glass disks.

Figure 6-2. Schematic diagram showing relative positions of the SSA disks and the apparatus used to measure pump cavity spectra.

Soviet MC-20, a white ceramic with flat diffuse reflectance throughout the visible and near IR.

We calibrated the pump cavity spectra absolutely using a calorimeter. As shown in Fig. 6-2, the calorimeter was placed directly above the scatter plate in a position which, by symmetry, receives the same radiation exposure. In these positions, neither the scatter plate nor the calorimeter measurably affected the radiation detected by the other.

The spectra were calibrated in three steps: (1) the pyroelectric detector signal was converted to intensity by comparing the time-integrated pyroelectric detector signal with the fluence measured by the calorimeter; (2) the time-dependent intensity in the 0.4 to 1.0- μm band was determined from pyroelectric detector signals recorded with two longpass filters. These filters began to transmit at 0.4- μm (Hoya L-42 filter) and 1.0- μm (Corning 7-56 filter); and 3) the J-20 spectrum was converted to spectral intensity by comparing the spectrum, integrated from 0.4 to 1.0- μm , with the intensity of radiation in the 0.4–1.0- μm band at the time the J-20 shutter was opened.

We made the calorimeter from a piece of copper (with dimensions $1/8 \times 2 \times 2$ in.) bead blasted on one side and coated with black chrome. Measurements show that the diffuse reflectance of black chrome for normal-incidence radiation is $3 \pm 1\%$ over the wavelength range 0.35 to 1.5- μm that covers the range of significant output by the Ce-doped quartz flashlamps. Black chrome has a high damage threshold for flashlamp radiation. Temperatures were measured by four thermocouples that were epoxied onto the back surface and shielded from the pump radiation by an aluminum plate.

Fluences were calculated as the product

$$(1-R_d)^{-1} r c t (\Delta T), \quad (1)$$

where R_d is the diffuse reflectance of the calorimeter surface (estimated to be 3%), r is the density of copper, c is the specific heat of copper, t is the thickness of the calorimeter plate, and ΔT is the temperature rise resulting from firing the flashlamps. To test the calorimeter calibration, we exposed the calorimeter to known fluences using a 1.5-cm-bore Ce-doped flashlamp operated in the open. The tests indicated that the calibration was accurate to within a few percent for light incident at near normal incidence. However, when the calorimeter was

placed near the lamp so that radiation was incident over a broad range of angles, the temperature rise of the calorimeter was approximately 8% lower than expected. This was probably the result of increased specular and diffuse reflections for light incident at large angles. Thus, it is reasonable to regard the measured fluences as lower limits which are probably within 5–10% of the actual fluences.

The spectra were measured under conditions in which the flashlamps closely approached the quasi-steady state in which transient effects, such as arc expansion, had a minimum effect. Thus, most spectra were measured at the time of peak flashlamp radiant, with the flashlamps operated at relatively long pulse lengths [$3(LC)^{1/2} = 678 \mu\text{s}$].

The spectra were measured with both 8c and 5W reflectors in the SSA. For some of the measurements made with the 8c reflectors, the LG-750 disks were removed.

Calorimetry Results. A measure of the efficiency of the pump cavity is the fluence ratio, $R_f = E_{\text{inc}}/E_{\text{rad}}$, where E_{inc} is the fluence measured by the calorimeter, and $E_{\text{rad}} = \eta_r E_{\text{del}}$ is the fluence radiated at the surface of the lamp. E_{del} is the measured electrical energy per unit of bore area delivered to the flashlamp. Values of the radiant efficiency of the flashlamps, η_r , were measured previously for an open 1.5-cm bore Ce-doped lamp operated under similar circuit conditions, using the techniques described in Ref. 12.

Figure 6-3 shows fluence ratio R_f plotted versus flashlamp explosion fraction (f_x) for the 5W and 8c reflectors, the latter both with and without disks. Explosion fraction is defined as the electrical energy delivered to a flashlamp divided by the energy required to explode the flashlamp on a single shot.¹⁰ Comparing the two lower curves (for data taken with disks), it is apparent that the measured fluences do not increase linearly as the number of lamps is increased, just as was found for the SSA gain tests. This is because the flashlamps themselves absorb a larger fraction of the pump cavity radiation as their number is increased. That R_f is so high at low flashlamp loadings indicates that there are multiple reflections of light in the pump cavity, especially when the SSA was run without disks. As f_x is increased, R_f decreases due to the increased reabsorption of light by the lamps as they become more opaque. Without disks R_f is obviously higher because of

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Figure 6-3. Fluence ratio versus explosion fraction for (a) 5W and 8c reflectors with disks, and for (b) 8c reflectors without disks.

the absence of Nd-absorption. The uncertainties for the R_f values are estimated to be $\pm 11\%$, due mostly to systematic calibration errors.

Measured Pump Cavity Spectra. Figure 6-4 shows a measured pump cavity spectrum (for the 8c reflectors, with disks) overlayed with a measured open lamp spectrum. Both spectra were measured at a peak flashlamp radiant output power per unit area of 32 kW/cm^2 . Here, the peak output powers for the SSA lamps are those of otherwise equivalent clear-fused quartz flashlamps, calculated from peak electrical input powers using data from our study of open lamps.¹²

The pump cavity spectrum has many features not present in the open-lamp spectrum. By comparing the pump cavity spectrum with the absorption spectrum of LG-750 glass, it is apparent that most of these features are caused by the absorption of light by the disks. However, as noted on the plot there are some obvious features caused by the lamps that show up on line reversal, absorption of radiation by Xe line transitions that usually appear in emission.

Line reversal occurs when light emitted by the hot plasma near the center of the lamp is preferentially absorbed, at the line transitions, by a relatively cool layer of plasma near the lamp wall. At the line transitions, the optical absorption depth is generally less than the thickness of the cool layer. In the continuum, the absorption depth is usually much greater. Therefore, the line emission radiated by the lamp originates in the cool layer near the wall, while the continuum radiation originates in the hot plasma near the lamp center giving line reversal.

While some line reversal is evident in the open lamp spectrum in Fig. 6-4, the line reversal is much more pronounced in the pump cavity spectrum. This suggests that the cool layer of plasma is made thicker by operating the lamps in the pump cavity, rather than in the open.

Figure 6-5 shows more evidence for the thickening of the cool layer in the pump cavity. In this plot, the blackbody temperature required to match the measured intensity of the radiation, for wavelengths greater than $1.0\text{-}\mu\text{m}$, is plotted versus instantaneous electrical input power per unit area. The blackbody temperatures were calculated using the filtered pyroelectric detector data, calibrated absolutely by the calorimetry data. Radiation in the infrared is sensitive to the condition of the cool layer of the

plasma because this portion of the lamp emission spectrum is dominated by strong line transitions. Data for the three different pump cavity conditions, and for a 1.5-cm bore open lamp, are presented. The blackbody temperatures for the pump cavity radiation all fall within a narrow range, and, above 20 kW/cm^2 , are lower than the blackbody temperatures for the open lamp. The differences in temperature between the pump cavity and open lamp spectra increase with input power, consistent with our observation that line reversal in the pump cavity spectra becomes more pronounced as input power is increased.

The mechanism causing the thickening of the cool plasma layer is not understood. It is plausible, however, that the infrared radiation returned to the lamp by the pump cavity and was absorbed by the cool layer causing the cool layer to expand toward the lamp wall.

An implication of our observations is that the reabsorption of infrared radiation by the lamp does not contribute efficiently to increased output power. Infrared reflecting lamp coatings are probably not useful at high lamp loadings. In fact, it appears that the reabsorption of infrared radiation is, to some degree, detrimental since lamp emission at some wavelengths is decreased due to the thickening of the cool layer.

Figure 6-6 compares spectra measured with the two different reflector types. These spectra were measured at peak flashlamp output powers per area of 36 kW/cm^2 . In the near infrared, where the flashlamp plasma is most opaque, the 8c and 5W spectra are in agreement. According to thermodynamics theory, this is the expected result when the flashlamps themselves are the dominant absorbers in the pump cavity. At the shorter wavelengths, where the plasma is less opaque, the 8c spectra exceed the 5W spectra. This is expected when there are objects in the pump cavity other than the flashlamps that are important absorbers (such as the disks).

Figure 6-7 shows the near-IR portion of a pump cavity spectrum measured at the time of peak gain, with the 8c reflectors. The lamps were fired at an explosion fraction of 0.29. The spectrum clearly shows several reversed Xe lines. The $1.05\text{-}\mu\text{m}$ Nd fluorescence, that depumps the disks, is less intense than the flashlamp radiation. From this data, we see that Nd-fluorescence makes only a small contribution to the cavity power at

Figure 6-4. A pump cavity spectrum (8c reflectors, with disks) and an open lamp spectrum, measured at 32 kW/cm^2 peak radiant output power per unit area.

Figure 6-5. Blackbody temperatures, corresponding to the intensity of radiation at wavelengths greater than $1.0\text{-}\mu\text{m}$, plotted versus electrical input power per unit area for a) the three pump cavity conditions and b) a 1.5-cm bore open lamp.

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Figure 6-6. Comparison of pump cavity spectra measured with the 5W and 8c reflectors.

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Figure 6-7. Near-IR spectrum showing 1.05-μm fluorescence in the pump cavity.

1.05- μm leading to upper-state depumping. (See Ref. 13 for a brief discussion of flashlamp depumping rates.)

Cavity Transfer Efficiencies determined from Pump Cavity Spectra. In our amplifier model, we use the cavity transfer efficiency η_t to describe the transfer of radiation from the flashlamps to the disks. It is defined as:

$$\eta_t = \frac{A_d \cdot S t}{A_f \int d\lambda I(\lambda, P_e) (1 - e^{-\alpha_g(\lambda)t}) \lambda / \lambda_0} \quad (2)$$

where S is the average Nd-pump rate in the disks (in units of power per unit volume), $I(\lambda, P_e)$ is the angle-averaged lamp emission spectrum (spectral intensity per unit bore area), α_g is the absorption coefficient for the disk glass, λ_0 is the laser wavelength, t is the disk thickness, A_f is the total bore area of the flashlamps, and A_d is the total surface area of the disks. When the amplifier model is used to determine cavity transfer efficiency, S is taken from the gain measurements (taking ASE losses into account), and η_t is inferred using our modified Trenholme-Emmett flashlamp model.¹⁴ Values of S and η_t are affected by errors in the stimulated emission cross section and the fluorescence decay rates.

We have developed a more direct way of determining cavity transfer efficiency using the measured pump cavity spectra to calculate S . We assume that the light incident on the slab is isotropic, and take into account the angular dependence of both the Fresnel reflections and the path length through the slab. In this case, S is given by

$$S = \frac{2\pi}{t} \int d\theta \int d\lambda \frac{I_{\text{inc}}(\lambda)}{\pi} \times \cos\theta \sin\theta (1 - e^{-\alpha_g(\lambda)t/\cos\theta}) \times \frac{\lambda}{\lambda_0} T_F(\theta) F \quad (3)$$

where θ is the angle of incidence, $I_{\text{inc}}(\lambda)$ is the measured pump cavity spectrum, $T_F(\theta)$ is the polarization-averaged transmittance taking Fresnel reflections into account, and F is the end-loss correction factor. S is changed by only a few percent by changing the angular distribution of light from isotropic to Lambertian.

The factor F is necessary since less light is incident upon the end-looking disk surfaces than the interior-looking disk surfaces, and S represents an average pump rate for all disks in the amplifier. In calculating F , we assumed that the pump rates for the end-looking disk surfaces are reduced from those of the interior disk surfaces by the ratio $E_{\text{inc,e}}/E_{\text{inc,i}}$, where $E_{\text{inc,e}}$ is the calorimeter fluence measured in front of an end-looking disk surface, and $E_{\text{inc,i}}$ is the fluence measured in front of an interior disk surface. Thus, some error is introduced in F by ignoring the differences in the spectral distributions of light incident upon the interior and end-looking disk surfaces. We estimate this error to be less than 2%.

For the SSA, F is given by

$$F = \frac{6 \cdot E_{\text{inc,i}} + 2 \cdot E_{\text{inc,e}}}{8 \cdot E_{\text{inc,i}}} \quad (4)$$

where the numbers 6 and 2 are the number of interior and end-looking disk surfaces. Fluences measured at the end disk were less than those measured at the interior disk by approximately 40%, regardless of f_x . Thus, F was approximately 90% for both the 8c and 5W data.

Figure 6-8 shows that the cavity transfer efficiencies determined from gain measurements are typically 25% smaller than those determined from the spectra. The cause of the differences between the η_t inferred from gain measurements and spectral measurements are not known, but the differences do not seem large given the differences in the techniques. Notice, however, that the second shape of the cavity transfer curves and relative magnitudes for the 8-lamp and 5-lamp cases are in good agreement. We estimate the absolute calibration errors for the η_t determined from spectra to be $\pm 12\%$.

Multisegment Amplifier. We have proposed to use multisegment (previously called "logpile") amplifiers (MSA) to scale Nd:glass disk amplifiers to megajoule-class size.⁴⁻⁶ Unlike the amplifiers used in previous Nd:glass fusion lasers, the MSAs would amplify beams divided into many segments. Segmentation reduces disk size, and ASE losses, while keeping beam transport relatively simple since several beam segments can be handled as a single beam. MSAs are very compact and have a smaller parts count compared with other amplifiers. They are

Figure 6-8. SSA cavity transfer efficiencies determined from gain measurements (solid curves) and from spectra, for both the 5W and 8c reflector types.

probably the most economical approach to building large fusion drivers.

We have proposed a full-scale 4×4 MSA suitable for a megajoule laser with four rows of disks pumped by three rows of central flashlamp arrays that pump disks on both sides, and by two rows of conventional flashlamp arrays that pump disks on one side. To simplify assembly and maintenance, we would build these 4×4 MSAs in one disk-long interchangeable module, with each module holding a 4×4 matrix of disks. We built and tested the single 2×2 MSA module shown in Fig. 6-9 which represents one-quarter of a full-size 4×4 module. The primary purpose for building the 2×2 MSA was to test the performance of the MSA pump cavity, with its unique central array, and to prove the one-disk-long module concept. Also, in designing and building a working amplifier, we addressed many of the important engineering and design issues generic to MSAs.

Design. In order to reduce the number of disks required to store a given amount of energy in an amplifier, it is desirable to operate at high values of stored fluence per disk (defined as the total energy stored in a disk, divided by the area of the largest beam aperture that could be amplified by that disk). However, by increasing stored fluence per disk, storage efficiency is generally decreased, due to accompanying increases in ASE losses and reabsorption of pump radiation by the flashlamps. The result is increased size and cost of the pulse power system driving the flashlamps. Consequently, there is a cost tradeoff between stored fluence per disk and storage efficiency.

We designed the 2×2 MSA according to performance predictions made with our empirical amplifier model, which was developed from SSA data. (See Article 5.2.3 for more discussion of these simulations.) Specifically, we selected the flashlamp pump area ratio (1.03), pulse length, and pulse energies to optimize storage efficiency for a high value of stored fluence per disk. Pump area ratio is a measure of the flashlamp density and is defined as $p \cdot d/s$, where d is the flashlamp bore diameter and s is the distance between lamp centers in the side array. As pump area ratio is increased, stored energy density increases and storage efficiency falls because of increased reabsorption of light by the flashlamps. In a cost-optimized system, depending on costs of disks and the pulse

power system, a different pump area ratio may be chosen.

The four disks in the MSA were made of LG-750 phosphate glass doped at 3×10^{20} Nd ions/cm³. The disk dimensions were $56.8 \times 30.0 \times 3.75$ -cm, excluding the 0.6-cm-thick edge claddings used to prevent parasitics. The disks were oriented at Brewster's angle with respect to the amplified beam segments. The segments measured 29-cm on each side.

To approximately balance the pumping between the central and side arrays, twice as many flashlamps (12) were used in the central array as in each of the two side arrays (6). The side arrays were similar to the flashlamp arrays used in all our previous amplifiers with cylindrical reflectors placed behind each lamp. However, the central flashlamp array was fundamentally different. It illuminated disks on both sides and used only small diamond-shaped reflectors between the lamps to limit optical transfer from lamp to lamp. Figure 6-10 shows the positions of the flashlamps and the reflectors in the central array.

The flashlamp dimensions were 2.5×61 -cm (bore diameter $\times$ arc length), and the Xe fill pressure was 200 Torr. The lamp envelopes were made of Ce-doped quartz. The flashlamp circuit capacitance and inductance were 188- μ F and 147- μ H, respectively. Each circuit drove three flashlamps in series. The flashlamps were preionized, with prepulses initiated 200- μ s before the main pulses and with energies per unit bore area of 0.1 J/cm².

For safety, to keep high-voltage transients away from diagnostic equipment and personnel, the flashlamp arrays were electrically isolated from the outside case and grounded back at the pulse power system. To simplify assembly and maintenance, the flashlamp arrays and disk holders were designed for rapid installation or removal from the module.

Simulation of Adjacent Modules. Since we tested only a single module, its performance was affected by the loss of pump radiation through its large open ends. For some of the tests, to reduce end losses, a flat silver reflector was placed over one end of the module. Because the reflector formed a virtual image of the module, it simulated two modules coupled together. As shown in Fig. 6-11, the two modules were simulated with their disks oriented in either "X" or "diamond" configurations, depending on

Figure 6-9. The 2 × 2 MSA.

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Figure 6-10. Positions of the flashlamps and diamond-shaped reflectors in the central array.

Figure 6-11. Schematic diagram of the simulated "diamond," "X," and "V" configurations (see text for details).

which end of the module was covered with the flat reflector. In these terms, the one-disk long module by itself formed a "V" configuration.

The ends of the module that were not covered with the reflector were covered with black-anodized boxes to prevent flashlamp radiation from leaking into the laboratory. Small holes were cut in the silver reflectors and in the black anodized boxes so that the probe beam could be scanned vertically and horizontally from the disk center.

Gain Measurements. The apparatus used to measure gain in the 2×2 MSA is the same apparatus used in our SSA experiments. We measured gain as a function of flashlamp explosion fraction while holding the probe beam fixed at the center of one of the disks. We also measured gain as a function of horizontal and vertical displacement from the disk center while holding the explosion fraction fixed at 0.2. Like the SSA measurements, these measured gains were converted to stored energy densities using the newly measured value of the cross section for LG-750 glass.

To find the average stored energy density of the entire disk, the stored energy density at the center of the disk was multiplied by a factor that corrected for the shape of the horizontal energy density profile. (The vertical profile was flat.) We assumed that the stored energy density profile did not change with the energy delivered to the lamps.

To calculate the total stored energy in the disk, the average energy density was multiplied by the extractable volume of the disk, $V_{\text{ext}} = (l - t/n)wt$, where l is the pumped length of the disk, w is the pumped width of the disk, t is the disk thickness, and n is the index of refraction.

Storage Efficiencies. Storage efficiencies were calculated by dividing the peak extractable energy stored in the disks by the total electrical energy delivered to the flashlamps. We assumed that the energy stored in the other three disks was the same as the energy stored in the probed disk. By symmetry, this should be the case. This assumption is supported by two laboratory measurements: 1) within experimental uncertainties ($\pm 1\%$), equal amounts of electrical energy were delivered to each of the two side arrays; and 2) our calorimetry measurements described below, which show that radiant energy is divided equally between the two sides of the central array.

In Fig. 6-12, storage efficiencies measured for the three MSA configurations are plotted versus flashlamp explosion fraction. Of the three configurations, the "V" configuration (the module by itself with open ends) had the lowest storage efficiency because it had the greatest end losses. The simulated "diamond" and "X" configurations had higher storage efficiencies because the flat silver reflector reduced the end losses by approximately one-half.

Extrapolated Performance of Longer Amplifiers. Figure 6-12 also plots the extrapolated storage efficiency of an interior module sandwiched between other neighboring modules. The performance of such a module would be almost completely unaffected by end losses. The efficiency of the interior module was calculated by adding to the average storage efficiency of the "diamond" and "X" configurations the difference between this average and the storage efficiency of the "V" configuration.

The storage efficiencies plotted in Fig. 6-12 can be used to predict the storage efficiencies of 2×2 MSAs of arbitrary length. In the extrapolation technique, the storage efficiency of the module on each end of the amplifier is approximated with the storage efficiency of either the simulated "diamond" or the simulated "X" configuration, whichever is appropriate. Then, the storage efficiency of an entire MSA (assembled from a number of MSA modules) is found by averaging the storage efficiencies of the end modules with the storage efficiency of the interior modules, with the appropriate weighing.

Performance Comparisons. Figure 6-13 plots the extrapolated storage efficiency of a 2×2 MSA which is seven disks long (the approximate length proposed for the Nova Upgrade) versus stored fluence per disk. Storage efficiencies of three other amplifiers are also plotted. These amplifiers are: (1) the single segment amplifier, operated at a lower value of pump area ratio (0.79) than was used in the MSA; (2) the Nova 46-cm amplifier; and (3) the Nova 31.5-cm amplifier.

At high values of stored fluence per disk, the MSA reached higher storage efficiency than has previously been achieved in Nd:glass disk amplifiers, including the Nova amplifiers. The Nova amplifiers are less efficient because of the use of: (1) higher pump area ratios; (2) elliptical disks that allow some pump light to pass from

Figure 6-12. Measured storage efficiencies of the simulated "diamond," "X," and "V" configurations, and the extrapolated storage efficiency of an interior module sandwiched between other adjacent modules.

Figure 6-13. Storage efficiencies for: (1) the SSA; (2) the MSA (extrapolated to an amplifier 7 disks long); (3) a Nova 46-cm amplifier; and (4) a Nova 31.5-cm amplifier.

one side of the amplifier to the other without passing through a disk; and (3) lower Nd-doping $\times$ disk thickness products. Nonetheless, the Nova amplifiers reached higher stored fluences per disk because of their higher pump area ratios.

At low values of stored fluence per disk, however, the SSA was more efficient than the MSA. This was largely due to its smaller pump area ratio. As a result, and because of limits on the amount of energy that could be safely delivered to the flashlamps, high values of stored fluence per disk were not accessible by the SSA.

Since other differences between the SSA and MSA affect storage efficiency, simple comparisons of storage efficiencies do not accurately indicate the performance of the MSA pump cavity geometry relative to that of the SSA. Such a comparison requires comparisons of MSA and SSA cavity transfer efficiencies at equal pump area ratios.

Cavity Transfer Efficiencies. To determine the relative performance of the MSA and SSA pump cavities, we analyzed the MSA and SSA data with our amplifier model. Except for one process (the transfer of optical power from the flashlamps to the disks in the pump cavity), the model calculates rates for all the energy transfer processes that occur in the amplifier, either by direct measurements or by empirical modeling. Once these rates were known, we simulated the gain measurements by calculating a cavity transfer efficiency that completely accounts for the transfer of power from the flashlamps to the disks.

Cavity transfer efficiency is a function of the pump area ratio. As pump area ratio is increased, cavity transfer efficiency is decreased since a larger fraction of the pump radiation is reabsorbed by the flashlamps themselves rather than by the disks. However, the stored fluence is increased, due to increased pump radiation. Cavity transfer efficiency is also affected by the pump cavity configuration. Here, configuration refers to the physical arrangement of the flashlamps, reflectors, disks, and blast shields. The MSA configuration is radically different from the SSA configuration because of the use of the central flashlamp arrays.

To obtain a measure of the relative efficiency of the MSA and SSA pump cavity configurations, with the effect of different pump area ratios removed, the cavity transfer

efficiency of the SSA was linearly interpolated to the same pump area ratio as was used in the MSA. Also, the cavity transfer efficiency of the MSA was extrapolated to that of an amplifier 4 disks long, the same length as the SSA, so that end losses would be weighted equally.

Figure 6-14 plots cavity transfer efficiencies of the MSA and the SSA (both for amplifiers 4 disks long, at a pump area ratio of 1.03) versus flashlamp output power per unit of bore area. For a given pump cavity, the cavity transfer efficiency is determined by flashlamp output power per unit area since the spectral properties of the flashlamps, in both absorption and emission, are functions of flashlamp output power per unit area.¹¹ The cavity transfer efficiency of the MSA was approximately 13% higher than that of the SSA. This is a measure of the relative pumping performance of the 2×2 MSA and SSA pump cavity configurations. We attribute the improved performance of the MSA mainly to the more efficient central flashlamp array.

Central versus Side Arrays. To determine the relative performance of central and side arrays, we measured the gain profile of an MSA disk while firing only the single central flashlamp array by itself, and while firing only the two side flashlamp arrays by themselves. In order to remove the effects of end losses from the measurements, the central array was tested in the simulated "diamond" configuration, while the side arrays were tested in the simulated "X" configuration. The flashlamps in both the central and side arrays were driven at an explosion fraction of 0.2.

Figure 6-15 shows the resulting measured gain profiles. By integrating these two curves and comparing their ratio, we find that 14% more energy was stored when only the central array was fired than when the two side arrays were fired. Thus, under the conditions of this experiment, it appears that the Nd-pumping efficiency of the central flashlamp arrays is 14% greater than that of side flashlamp arrays. We conclude that the higher Nd-pumping efficiency of the central flashlamp array is a major factor contributing to the higher cavity transfer efficiency of the 2×2 MSA relative to that of the SSA. (We believe that another important contributing factor is the positioning of the vertically-oriented lamps in the MSA so that the edges of the disk holders do not occlude the flashlamp radiation, as is the case in the SSA.)

Figure 6-14. Cavity transfer efficiencies of the SSA (interpolated to the pump area ratio used in the MSA) and the MSA (extrapolated to the 4-disk length used in the SSA).

Figure 6-15. Gain profiles measured while firing either: 1) only the central array, in the "diamond" configuration, and 2) only the two side arrays in the "X" configuration.

Nonuniform Gain Profiles. One consequence of the more efficient pumping efficiency of the central flashlamp array is nonuniform pumping of the MSA disks: the side of each disk that faces the central flashlamp array is pumped more strongly than the side that faces a side flashlamp array. Some gain nonuniformity is also caused by the separate contributions to the gain profiles by the central and side arrays that do not result in a constant value when added together. The loss of pump radiation through the ends of the amplifier also causes gain nonuniformity. End losses cause the gain in the portion of the disks nearest the end of the amplifier to be lower than that of the opposite side. However, the end loss contribution to nonuniformity can be reduced for the amplifier as a whole by designing the amplifier to have an odd number of disks.

Figure 6-16 shows gain profiles measured for the "V," "diamond," and "X" configurations with the flashlamps fired at an explosion fraction of 0.2. Figure 6-16 also shows the extrapolated average gain profile for an amplifier 7 disks long. (Gain coefficients can be extrapolated in a way similar to the way that storage efficiencies can be extrapolated.) Since the extrapolated gain coefficient for the 7-disk (odd number) long amplifier is much more uniform than the measured gain coefficients, it is apparent that end losses are a major contributor to the nonuniformity seen in the measured configurations.

We know that gain nonuniformity contributes to beam steering due to nonuniform heating of the disks. Beam steering and gain nonuniformity could probably be reduced by more evenly balancing the pumping by the central and side arrays.

Radiant Efficiency of the MSA Flashlamp Arrays. We measured the time-integrated radiant (electrical-to-optical conversion) efficiencies of the flashlamp arrays used in the 2×2 multisegment amplifier. The measurements were made to compare the performance of the central and side flashlamp arrays, and to provide data by which modelers can validate combined flashlamp model and ray-trace codes.

We built a cylindrical calorimeter large enough to enclose the flashlamp arrays and to measure their radiated optical energies. The calorimeter dimensions were 88×109 -cm, height $\times$ diameter, and its walls were made of 3/16-in. aluminum. The calorimeter was

assembled in two halves which were placed on either side of the arrays. The inside surfaces of the calorimeter were coated with black chrome. Twenty-four thermocouples were connected in series and distributed over one half of the outside of the calorimeter.

The response time of the calorimeter was several seconds, equal to the time required for thermal equilibration from one side of the calorimeter wall to the other. The slight cooling of the calorimeter during the equilibration period was taken into account by linear extrapolation of the temperature curve back to the shot time. Radiant efficiency was calculated from the extrapolated temperature rise at shot time using the measured mass of the calorimeter and the specific heat of aluminum at room temperature.

The central flashlamp array was tested both with and without the diamond-shaped reflectors placed between the flashlamps. The flashlamps in the arrays were driven with $3(LC)^{1/2}$ pulse lengths of 500-ms, the same pulse lengths used in the MSA tests. Also, as in the MSA tests, the flashlamps were prepulsed with sufficient energy (about 0.1 J/cm^2) to attain the maximum radiant efficiency.

Radiant Efficiencies and Transfer Efficiencies. Figure 6-17 shows both radiant efficiencies and transfer efficiencies plotted versus explosion fraction. Transfer efficiency is defined as the radiant efficiency of the array divided by the radiant efficiency of a bare lamp. It is a measure of the effect of the flashlamp array geometry on radiant efficiency. Transfer efficiencies were calculated using previously measured radiant efficiencies for a 1.5-cm bore lamp, even though the lamps in the array had bore diameters of 2.5-cm. Using radiant efficiencies for lamps with bore diameters other than 2.5-cm introduces only small errors in transfer efficiencies since the output power of a flashlamp depends primarily on its instantaneous electrical input power per unit of bore area, regardless of bore diameter or of pulse length.¹² We estimate the errors in the transfer efficiencies to be $\pm 7\%$, and the errors in the radiant efficiencies to be $\pm 5\%$.

Notice that the transfer efficiency of the central array (with diamond-shaped reflectors) approaches unity while the others are significantly below.

Two loss factors cause the radiant efficiencies of the arrays to be lower than those

Figure 6-16. Gain profiles measured for the "diamond," "X," and "V" configurations, and an extrapolated average gain profile for disks in an amplifier seven disks long.

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Figure 6-17. (a) Radiant and (b) transfer efficiencies of the side flashlamp array, and the central flashlamp array, both with and without diamond shaped reflectors between the lamps.

of the bare flashlamp. These factors are: (1) absorption of radiation by the lamp plasma; and (2) absorption by the reflectors. The lower radiant efficiency of the side arrays compared to the central array is probably the result of greater absorption by both the plasma and the reflectors. We believe that both factors are more important in the side arrays, since: (1) the cylindrical reflectors in the side arrays reflect much of the radiation back through the lamps; and (2) the cylindrical reflectors in the side arrays are much larger than the diamond-shaped reflectors in the central array. Increased absorption by both the plasma and the reflectors cause the transfer efficiencies of the arrays to fall as explosion fraction is increased. The reasons for these conditions are: (1) the flashlamp plasma becomes more opaque; and (2) the lamp emission spectrum blue-shifts, so that a larger fraction of the optical power emitted by the lamps is absorbed by the silver reflectors that absorb most strongly below 450-nm. The radiant transfer efficiencies of the central array were reduced by removing the diamond-shaped reflectors. This demonstrates that the reflectors are effective in reducing the transfer of radiation from lamp to lamp, as they were designed to do.

The single-sided gain measurements reported above show that, at an explosion fraction of 0.2, the Nd-pumping efficiency of the central array is 14% higher than that of the side array. We find here a comparable improvement in the radiant efficiency of the central array (with diamond-shaped reflectors) compared to the side array, 18% (also at an explosion fraction of 0.2). However, some differences in the ratios of Nd-pumping and radiant efficiencies should be expected since Nd-pumping is sensitive to the spectral distribution of output power while radiant efficiency is not. Specifically, plasma absorption removes radiation from the major Nd-pump bands while absorption in the near UV, by the silver reflectors, does not.

Author: A. C. Erlandson

Major Contributors: K. S. Jancaitis, R. W. McCracken, J. E. Murray, and H. T. Powell

6.2.3. Disk Amplifier Performance Model

We have combined our earlier work that treated many of the individual pieces of the physics related to amplifier performance¹⁵⁻¹⁷ into a semi-empirical model which can be used to

analyze the results of our experimental measurements, and can be used to predict the effects of variations in amplifier design parameters on the expected stored energy density. The major new element introduced in the development of this model is an empirically determined radiant transfer efficiency that determines how well the light emitted by the flashlamps is transmitted to the laser glass. We find that for a given amplifier configuration, this transfer efficiency depends only on the instantaneous output power per unit bore area P_{out}/A from the flashlamps. Once this dependence is determined for a given amplifier geometry the model predicts both the magnitude and the time dependence of the gain observed in the Single-Segment Amplifier and Multi-Segment Amplifier module over the range of configurations tested.

The treatment of the energy transfer, storage, and decay in this model is broken down into several distinct steps. In all our analyses we begin with the experimentally measured electrical energy input to the flashlamp arrays. The conversion of the electrical input $P_{in}(t)$ into the optical output $P_{out}(t)$ of the lamps and the shape of the time-dependent output spectrum are determined using our standard flashlamp model.^{15,16} We use an empirically determined transfer efficiency function $\eta_{xfer}(P_{out}/A)$ (described in detail below) to model the transport of the flashlamp light to the laser disks. This transfer efficiency is the ratio of the instantaneous pumping efficiency of the inversion in the disks relative to that which would be obtained if all the flashlamp light passed through the disks once in a direction normal to the surface [$\eta_{pump}(P_{out}/A)$]. The single-pass pumping efficiency η_{pump} is calculated by convolving the predicted lamp spectrum with the glass absorption. The time dependence of the stored energy in the laser disks is then obtained by solving the differential equation,

$$E_{stored}(t) = \eta_{xfer}(P_{out}/A)\eta_{pump}(P_{out}/A)P_{out}(t) - k_{decay}(E_{stored})E_{stored}(t) \quad (1)$$

The decay rate includes both spontaneous and stimulated contributions. The latter component is estimated using the Monte Carlo ray-tracing method described in an earlier report, enhanced by the factor of 50% which we

find necessary to obtain agreement with our experimental observations.¹⁷

For a given amplifier configuration, the cavity transfer efficiency curve is generated by fitting the predictions of the amplifier model to the peak values of the gain coefficient observed experimentally. We have found that a functional dependence of the form

$$\eta(P_{\text{out}}/A) = A + B \exp(-C P_{\text{out}}/A) \quad (2)$$

where the coefficients A, B and C are varied to obtain the best fit and works well for both the SSA and MSA configurations. This functional form agrees with our expectations based on observations of the lamp radiative characteristics. As the lamp loading increases they become more opaque to their own output and present an increasing optical loss due to the imperfect optical recycling that takes place in the lamps. At the highest loading the lamps eventually become optically black and the absorption and loss saturates. In addition, as the power input to the lamps is increased their output shifts to the ultraviolet where the silver reflectors used in the SSA and MSA are less efficient.

The cavity transfer efficiencies inferred for the 3-, 5-, and 8-lamp per side SSA geometries are presented in Fig. 6-18. In each case the curve shown was obtained by averaging the fit coefficients obtained in separate analyses of the experimental data taken at one input pulse length (and/or glass type) for the amplifier configuration of interest. As one might expect due to feedback considerations the cavity transfer efficiency is a decreasing function of the number of lamps packed into each side of the amplifier.

To test our assumption that the transfer efficiency for a given configuration depends only on the lamp loading, we used the transfer efficiency curves presented in Fig. 6-18 in our amplifier model to calculate the expected peak storage efficiency for all of the data taken at each lamp packing fraction. The results of this analysis for the 5-lamp per side case using the Winston reflector shape are presented in Figure 7-4 of Ref. 2 of the previous section. In Figs. 6-19 and 6-20 we show the comparison of the experiment and model for some additional data taken with the 3-, 5-, and 8-lamp geometries. The fit between the predicted and observed results is a good one for the data except that taken at 8-lamps per side using LG-660 glass. We

believe that this disagreement is probably due to a degradation in the quality of the reflectors in the period between the experiments with LG-660 and LG-750 glass.

In addition to yielding a good fit to the measured peak gains, the model also does a good job in predicting the time dependence of the energy stored in the amplifiers. In Fig. 6-21 we show the measured and predicted temporal gain profiles for two cases with a short and long input pulse, respectively. Since the fitting process involved only the peak gain coefficients, this agreement for the time dependence of the gain indicates that the model treats most of the physics involved in the performance of the amplifiers.

We have also used our model to analyze the data taken for the MSA amplifier module. The inferred cavity transfer efficiency curves are shown in Fig. 6-22 and the comparison between the model and experimental values for the peak energy storage are displayed in Fig. 6-23. The derived cavity efficiencies for the SSA and MSA geometries are somewhat similar, but a direct comparison is somewhat misleading because the lamp packing fractions in the two cases are different. We can, however, use the data in Fig. 6-18 to obtain an interpolated value for the transfer efficiency of the SSA at a lamp density equivalent to that in the MSA, and (following the procedure described in the previous section) extrapolate the transfer efficiencies in Fig. 6-22 for the MSA module to get an estimate for the transfer efficiency in a full MSA amplifier.

Author: K. Jancaitis

Major Contributors: H. Powell, J. E. Murray and A. Erlandson

6.2.4. Optical Loss in Brewster-Disk Systems

Introduction. Optical losses are extremely important in fusion laser systems, especially when energy is extracted in a gain-saturated regime such as in the Nova upgrade design. Optical losses adversely affect both the energy extraction efficiency and the square pulse distortion. A loss measurements facility was established at Livermore in 1986, and our ability to measure losses using a cavity lossmeter¹⁸ and a laser calorimeter¹⁹ have been described in previous annual reports. This article discusses the addition of an integrating-sphere-based scatterometer to the loss measurements facility,

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Fig. 6-18. Inferred cavity transfer efficiencies for the single-segment amplifier for the 3-lamp configurations examined experimentally.

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Fig. 6-19. Comparison of the calculated (solid lines) and experimentally observed (data points) storage efficiencies in the single-segment amplifier using LG-660 laser glass.

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Fig. 6-20. Comparison of the calculated (solid lines) and experimentally observed (data points) storage efficiencies in the single-segment amplifier using LG-750 laser glass in the 8-lamp per side geometry.

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Fig. 6-21. (a) Comparison of the calculated and predicted time dependence of the gain in the single-segment amplifier at short ($\tau_{\text{input}} = 152 \mu\text{s}$) input pulse lengths. (b) Comparison of the calculated and predicted time dependence of the gain in the single-segment amplifier at long ($\tau_{\text{input}} = 927 \mu\text{s}$) input pulse lengths.

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Fig. 6-22. Inferred cavity transfer efficiencies for the multisegment amplifier module for the three amplifier configurations examined experimentally.

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Fig. 6-23. Comparison of the calculated (solid lines) and experimentally observed (data points) storage efficiencies in the multisegment amplifier module in its different experimental configurations.

as well as new results obtained using the instruments in this facility. Finally, we discuss a recently demonstrated ability to use the ringdown lossmeter with samples oriented at Brewster's angle to measure birefringence losses.

The Integrating Sphere Scatterometer.

Previous instruments in our loss measurement facility measured absorption loss (the laser calorimeter), or total loss (the ringdown lossmeter), which is the sum of absorption and scattering losses. In fact, the ringdown lossmeter is a rather sensitive instrument that can give varying results due to variations in operator alignment techniques. Furthermore, the ringdown lossmeter measurement is very sensitive to the parallelism of the surfaces of a sample (if used at normal incidence). A wedge in the sample leads to an apparent increased level of loss that is linear in the wedge angle.²⁰

To provide a more direct and accurate method to quantify scattering losses in optical materials, an integrating-sphere-based scatterometer was built²¹ and is shown schematically in Fig. 6-24. The measurements made with this system are not sensitive to wedges in the sample, unless Fresnel reflections from the sample surfaces do not exit the sphere (a much less stringent requirement than demanded by the ringdown lossmeter). Samples are held on a pedestal at the center of the sphere, while a lens, L_2 , focuses a Nd:YAG laser beam through the sample, and onto the exit aperture. The lens is antireflection coated and has very little scattering of its own so that it introduces little background signal. The exit aperture can be as small as 1-mm, so that scattered light is integrated and detected for all scattering angles greater than 5-mrad.

The scatterometer was used to measure scattering losses in a series of pilot-melt laser glass samples^{21,22} that had been previously measured using the laser calorimeter and ringdown lossmeter.¹⁹ The results of these measurements are shown in Table 6-1.²³ For these samples the results confirm that scattering losses are a small fraction of total passive loss, and that the total loss is very close to the measured absorption loss. This observation leads us to believe that cheaper polishing methods could be used to finish laser amplifier disk surfaces while still achieving adequately low loss provided that the laser damage performance of these surfaces remained adequate. In highly

pure fused quartz window material, however, we found that absorption losses were extremely low, and that scattering losses were comparable to, or greater than the absorption losses.

The scatterometer also proved to be useful in the quantification of losses in samples of crystalline laser materials grown at LLNL and elsewhere.^{20,24-26} In some Nd,Cr:GSGG crystals, the scatterometer was an important tool that demonstrates the necessity of sample parallelism.²⁰ In LiCaAlF crystals, the scatterometer quantified the scattering losses²³ in crystals produced by early experimental crystal growth runs.²⁴⁻²⁶ Advances in subsequent crystal production techniques reduced scattering losses substantially.^{25,26}

Separation of Nd^{3+} Absorption from Background Losses in Laser Glass. One of the most useful aspects of the ringdown lossmeter results from its capability to measure losses as a function of wavelength.¹⁹ We used its spectral scanning capability to confirm the likely contaminants responsible for background absorption losses in laser glass. As seen in Table 6-1, the magnitude of these losses is typically in the range of $0.1\% \text{ cm}^{-1}$. For advanced fusion laser designs, we would like these background losses to be less than $0.05\% \text{ cm}^{-1}$.

Figure 6-25 shows a typical spectral scan of losses in a phosphate laser glass sample at wavelengths near 1.06 microns. The most obvious absorption feature in the scan is due to the small thermal population of the $^4I_{11/2}$ lower laser level of Nd^{3+} that makes transitions to the $^4F_{3/2}$ upper laser level. We used the measured emission lineshape function for the $^4F_{3/2}$ to $^4I_{11/2}$ transition along with reciprocity theory to predict the lineshape function for the corresponding absorption transition²⁷ (which is not easily measured due to the extremely small population of the lower laser level). We fit the predicted absorption lineshape function to the measured loss curve in Fig. 6-25 to obtain a measure of the linearly sloping background absorption.²⁷ The slope of the background absorption was compared to the known absorption spectra of common transition metal impurity ions, and a decomposition of the loss spectrum into a sum of Nd^{3+} , Cu^{2+} , and Fe^{2+} was obtained, as shown in Fig. 6-26.

Clearly, the decomposition of the absorption spectrum into these components is not unique. Ion microprobe analysis of some of the laser glass

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Figure 6-24. Schematic of the integrating-sphere scatterometer which has been used to measure scattering losses in laser glasses and crystals.

Table 6-1. Comparison of loss measurements of Schott and Hoya pilot-melt laser glass samples (all loss coefficients are in units of 10^{-6} cm^{-1}).

Vendor	Melt #	γ^*_{calor}	γ_{scat}	$\gamma^*_{\text{calor}} + \gamma_{\text{scat}}$	γ_{ringdown}
Schott (LG-750)	714	960	40	1000	930
	716	910	18	928	880
	723	1390	43	1433	1290
	735	1240	12	1252	1120
	736	1270	42	1312	1180
	738	1070	65	1135	1020
	742	970	39	1009	1010
Hoya (LHG-8)	20-01	1120	23	1143	1050
	25-01	860	61	921	830
	28-01	820	40	860	810

* Corrected for Nd absorption and fluorescence.

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Figure 6-25. Nd^{3+} absorption is separated from background absorption using the spectral scan capability of the ringdown lossmeter.

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Figure 6-26. Illustration of how the ringdown lossmeter data for a 31.5-cm disk of LG-750 can be decomposed into absorption by Nd³⁺, Cu²⁺, and Fe²⁺.

samples, however, have revealed concentrations of Cu and Fe impurities that are generally consistent with this interpretation.²⁸

Depolarization Loss Measurements with the Ringdown Lossmeter. We modified the ringdown lossmeter so that it could be used to measure depolarization losses in laser disks mounted at Brewster's angle, as shown in Fig. 6-27.²⁹ We have shown that subtracting the loss per unit length measured at normal incidence from values measured at Brewster's angle yields depolarization losses as small as 10^{-4} per pass.³⁰ Results obtained for one quadrant of a 9-cm phosphate laser disk are shown in Fig. 6-28.³⁰ It is seen that the loss levels generally increase toward the edge of the disk, where increased effects of residual stress birefringence are expected. Actual losses due to depolarization in an amplifier chain can approach values measured in this experiment only if a consistent pattern of birefringence is encountered repeatedly from disk to disk, so that the phase retardations add coherently. Since it is likely that there will be some variation in birefringence patterns from disk to disk, some cancellation of retardation is expected, and the actual depolarization loss per disk may be less than what is measured using this technique.

The fact that the ringdown lossmeter measures losses by multipassing the laser disk increases the sensitivity of this technique substantially relative to classical methods of measuring residual stress-birefringence. While phase retardation increases linearly with distance, the intensity in the orthogonal polarization as well as depolarization losses increase quadratically. Current specifications for Nova laser glass place an upper limit of $\Delta n < 5\text{-nm/cm}$ for residual stress-birefringence in laser disks, while some future designs call for $\Delta n < 3\text{-nm/cm}$. We have shown that the ringdown lossmeter can readily measure birefringence levels down to $< 1\text{-nm/cm}$, which should be useful in quality control for future laser disk procurement.

Author: John Caird

Major contributor: Norm Nielsen

6.2.5. Passive Multipass Designs

The off-angle multipass (OAMP) geometry under consideration for the output stage of a Nova upgrade, makes possible an entirely

passive optical switch, based on the slight angular difference between successive passes. This section describes two angle-sensitive optical components, a multiorder wave plate and a kilolayer mirror, which could be made into passive switches.

The basic idea of an angle-sensitive switch is illustrated by the 2-pass amplifier geometry shown in Fig. 6-29. Here the switch consists of a wave plate and polarizer. The wave plate is oriented such that the relative phase retardation it introduces on the first pass is a multiple of 2π , and the beam remains vertically polarized. For the second pass, the mirror introduces a small offset angle $\Delta\theta$, such that the retardation introduced by the wave plate increases or decreases by π , changing the polarization to horizontal and causing reflection from the polarizer.

Three and four pass OAMP geometries that take advantage of angle-sensitive switches require a relatively high angular sensitivity in one direction and a minimal sensitivity in the other. Figure 6-30a shows a 4-pass OAMP geometry to help illustrate these requirements on angular sensitivity. The two lenses image-relay mirror M1 onto M2, and vice-versa, so that the beam returns to the same position on the mirrors after a round-trip in the cavity. Note that only the centerlines of the beam on each pass are shown to reduce the complexity of the drawing. The beams are actually collimated between the mirrors and lenses, and they come to a focus at the pinhole plane between the lenses. The entire region between the lenses must be evacuated, as for spatial filters.

The pulse is introduced into the cavity by reflection from the small diagonal mirror near the pinhole plane of the relay ("far-field injection") where the beams come to a focus. Figure 6-30b shows the angular orientation of the beam at the switch for each of the passes, where θ is the angle in the plane of Fig. 6-30a, and ϕ is the angle out of the plane. A pulse is injected below the plane of Fig. 6-30a (below is indicated by the dashed beamline) and toward M1. Mirror M1 is oriented orthogonal to the cavity centerline that reverses the sign of both angular offsets shown in Fig. 6-30b, so that pass 2 is above the plane of the figure. The switch is oriented such that its sensitive angular direction is along θ and its insensitive direction is along ϕ , and such that it transmits for the angle of pass 2. Mirror M2 changes ϕ but not θ , so that pass 3 also

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Figure 6-27. Schematic for the ringdown lossmeter as modified to make measurements of birefringence losses for beam propagation at Brewster's angle.

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Figure 6-28. Birefringence losses measured over one quadrant of a 9-cm Nova amplifier disk as a function of position.

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Figure 6-29. A two-pass geometry using an angle-sensitive waveplate to separate input and output beams.

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Figure 6-30. a) A 4-pass off-axis multipass geometry that utilizes an angle sensitive switch as a passive switchout; b) the angular orientation of the beam at the switch on each pass.

transmits through the switch. At mirror M1 again, both angles change sign so that pass 4 is at the angle that switches out of the cavity.

The utility of an optical switch for this application requires that the switching angle $\Delta\theta$ be in an appropriate range. It must be relatively small, because it determines the extra aperture required to accommodate losses from vignetting in OAMP geometries. On the other hand, it must be at least a few times the diffraction angle of the beam (which is limited by spatial filter pinholes), to provide adequate isolation between passes in the multipass cavity. Also $\Delta\theta$ determines the maximum diameter of the far-field injection mirror that influences the damage threat to the mirror. Consequently, its optimum value is a trade-off. Current estimates for $\Delta\theta$ fall in the range from about .05 to 0.3° .

We have identified two types of angularly sensitive optics which might be useful for this switching application, the combination of a multiorder wave plate and polarizer and the kilolayer mirror. Both potentially exhibit adequate angular sensitivity in one direction and little or no sensitivity in the other, and both could be scaled to 30-cm apertures with adequate quality and damage threshold.

Multiorder Waveplate Switches. The angularly sensitive switching action of a multiorder wave plate results from the change in extraordinary index with angle relative to the optic axis of the wave plate. The phase retardation Φ between ordinary and extraordinary waves after propagating a distance d at the angle θ to the optic axis (measured internal to the crystal) is given by:

$$\Phi(\theta) = [n_o - n_e(\theta)] \times 2\pi \frac{d}{\lambda} \quad (1)$$

for ordinary and extraordinary indices n_o and n_e , respectively, and wavelength λ . For full-wave and half-wave plates,

$$\Phi(\theta) = 2m\pi \quad (2)$$

$$\Phi(\theta + \Delta\theta) = (2m + 1)\pi \quad (3)$$

respectively, where m is an integer corresponding to the order of the wave plate. For a uniaxial crystal,

$$1/n_e^2(\theta) = \cos^2 \theta / n_o^2 + \sin^2 \theta / n_e^2 \quad (4)$$

Solving for θ and $\Delta\theta$ gives:

$$\sin^2 \theta = \frac{[(1 - \lambda m / n_o d)^{-2} - 1]}{(n_o^2 / n_e^2 - 1)} \quad (5)$$

$$\sin^2(\theta + \Delta\theta) = \frac{[n_e(\theta)/n_o - \lambda/2n_o d]^{-2} - 1}{(n_o^2 / n_e^2 - 1)} \quad (6)$$

For each integral value of m , there exists a θ_m given by Eq. 5 that gives a full wave retardation on transmission through the crystal and therefore no change in polarization. There is also a corresponding $\Delta\theta_m$ given by Eq. 6 that gives an additional half-wave of retardation that changes the polarization by 90° .

Figure 6-31a shows $\Delta\theta_m$ versus θ_m from Eqs. 5 and 6 for KDP.³¹ (Note that only every fifth solution is shown from 15° to 75° .) The $\Delta\theta$ magnitudes would be exactly symmetric with respect to $\theta = 45^\circ$ if the crystal thickness were an integral multiple of $\lambda/(n_o - n_e)$. Note also that any one of the half-wave solutions $\theta_m + \Delta\theta_m$ could be used with any one of the full-wave solutions θ_m .

The following approximation to Eq. 6 gives the shape of the $\Delta\theta$ versus θ curve in Fig. 31a, and it is useful for quickly estimating the usefulness of crystals:

$$\Delta\theta \approx \frac{\lambda}{2d(n_o - n_e)} \times \frac{(n_o + n_e)}{2 \sin^2 \theta} \quad (7)$$

At $\theta = 0$, the solution becomes:

$$\Delta\theta(0) \approx \sqrt{\frac{\lambda}{2d(n_o - n_e)}} \times \frac{(n_o + n_e)}{2} \quad (8)$$

For example, Eq. 7 gives $\Delta\theta = 0.130^\circ$ and 0.127° for KDP and KD*P, respectively, for 1-cm-thick wave plates at 1053 nm and at an orientation of $\theta = 45^\circ$.

Figure 6-31b also shows solutions for KDP, but only for $\theta = 25$ to 35° and for four wave plate

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Figure 6-31. Discrete solutions for the switching angle $\Delta\theta$ of a KDP waveplate versus θ for a) the full range of θ and a 1.0-cm-thick waveplate (only every fifth solution is shown between 15° and 75°), and b) a smaller range and four thicknesses.

thicknesses. Clearly, a switching angle in the desired range of .05 to .3° can easily be obtained for reasonable wave plate thicknesses.

To confirm these calculations, we measured and calculated $\Delta\theta$ at 633 nm for one of the old $5 \times 5 \times 1$ cm³ KDP type II SHG crystals used in the first 3×3 crystal array to demonstrate diamond turning and precision orientation of KDP crystals. Assuming that the c-axis was properly oriented at 58.7° to the surface normal (for type II SHG from 1053 nm), Eq. 6 gives $\Delta\theta = 0.0767^\circ$, using the dispersion data of Ref. 31. With a 15-mm collimated HeNe beam, we measured $\Delta\theta = 0.076^\circ$ near the surface normal, in excellent agreement with calculations.

We have identified four potential problems for the use of this wave plate in a multipass geometry. All give rise to depolarization that would decrease the effective transmission through the switch, but all are manageable with KDP or KD*P.

Anisotropic loss coefficients, as in KDP and KD*P, result in a slight change in the polarization direction away from 0 or 90° on transmission through the wave plate. The direction of polarization rotation is toward the axis of weaker absorption, e.g. toward the extraordinary axis in KDP, and it occurs for both half-wave and full-wave orientations of the wave plate. Further, the polarization rotation increases with the number of passes through the wave plate until the polarization is reset by a polarizer. The change in polarization angle $\Delta\psi$ is given by:

$$\Delta\psi = \tan^{-1} \left\{ \exp[\alpha_o - \alpha_e(\theta)d] \right\} - 45^\circ \quad (9)$$

which simplifies for small absorptions and uniaxial crystals to:

$$\Delta\psi = \frac{d}{2} (\alpha_o - \alpha_e) \sin^2 \theta \quad (10)$$

Since the solutions for $\Delta\theta$ are symmetric about $\theta = 45^\circ$, Eq. 8 shows that solutions less than 45° should be used to minimize $\Delta\psi$. For example, a 2-cm-thick KDP wave plate, for which $\alpha_o = .058/\text{cm}$ and $\alpha_e = .006/\text{cm}$, gives $\Delta\psi = 2.24^\circ$ for $\theta = 60^\circ$ and 0.74° for $\theta = 30^\circ$.

The seriousness of this polarization rotation is determined by the loss due to repolarization. Since a polarizer is required for switching, the

number of passes between repolarization would be two, unless a second polarizer were added to reduce this loss. Loss L from polarization rotation is given by $L = 1 - \cos^2 \Delta\psi \approx \Delta\psi^2$, and the above 2-cm KDP wave plate at $\theta = 30^\circ$ would give $\Delta\psi = 1.49^\circ$ after 2 passes for a negligible repolarization loss of .07%. Since the absorption for 90% deuterated KD*P is an order of magnitude less than for KDP, its loss would be even smaller.

A second source of depolarization is the half-angle diffraction spread of the beam $\delta\theta$, which results in a phase retardation $\delta\Phi = \pi\delta\theta/\Delta\theta$ for those parts of the beam at this maximum beam spread. However, $\delta\Phi$ in an OAMP is determined by the OAMP geometry, since $\delta\theta = (\text{pinhole diameter})/2(\text{focal length})$, and the switching angle $\Delta\theta$ is set equal to $(\text{pinhole spacing})/(\text{focal length})$. Consequently, $\delta\Phi = \pi/R$, where R is the ratio of the pinhole spacing to the pinhole diameter. The loss L due to repolarization goes as $1 - \cos^2(\delta\Phi/2) \approx \delta\Phi^2/4$, so the loss for those rays at the maximum beam spread would be $L \approx (\pi/4R)^2$. At worst, this loss amounts to a "soft" spatial filtering, with the loss increasing quadratically with ray angle to this maximum. In fact, R will probably be in the range of 8 to 10 for a practical OAMP to ensure isolation between passes, and to allow a reasonable-sized injection mirror. Therefore, this loss mechanism will also be negligible.

A third source of depolarization comes from the angular change $\Delta\phi$ which is at right angles to the switching direction $\Delta\theta$. Figure 6-32 shows the geometry with $k(\theta)$ indicating the direction for a full-wave of retardation, $k(\theta+\Delta\theta)$ indicating the half-wave retardation direction, and $k(\theta+\delta\theta)$ indicating the change from $k(\theta)$ by $\Delta\phi$ at 90° to $\Delta\theta$. Simple trigonometry gives $\delta\theta \approx \Delta\phi^2/2\tan\theta$. As before, $\delta\theta$ causes a phase retardation $\pi(\delta\theta)/\Delta\theta$, that in turn causes a depolarization loss $L \approx (\pi\Delta\phi^2/4\Delta\theta\tan\theta)^2$. Generally, the angular change in the two directions will be the same, so that $\Delta\theta = \Delta\phi$, and this loss will be negligible for acceptably small switching angles $\Delta\theta$. For example, even the largest estimated value of $\Delta\theta = 0.3^\circ$ gives $L = .005\%/\text{pass}$.

A fourth potential source of depolarization results from the temperature sensitivity of the wave plate. The consequences of a change in temperature can be obtained by differentiating Eq. 1:

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Figure 6-32. k-vector geometry for calculating the effect of an angular change $\Delta\Phi$ at right angles to the switching angle $\Delta\theta$.

$$d \frac{\Phi(\theta)}{dT} = \gamma(\theta) \Phi(\theta) + \left[\frac{dn_o}{dT} - \frac{dn_e(\theta)}{dT} \right] \times 2\pi \frac{d}{\lambda} \quad (11)$$

where $\gamma(\theta) = \gamma_o \cos^2 \theta + \gamma_e \sin^2 \theta$ is the linear expansion coefficient and

$$\frac{dn_e(\theta)}{dT} = n_e(\theta)^3 \times \left[\left(\frac{\cos^2 \theta}{n_o^3} \right) \frac{dn_o}{dT} + \left(\frac{\sin^2 \theta}{n_e^3} \right) \frac{dn_e}{dT} \right] \quad (12)$$

Table 6-2 gives $d\Phi/dT$ for two orientations of KDP and KD*P, and single-pass depolarization loss for $\Delta T = 1^\circ\text{C}$ for three wave plate thicknesses, using γ and dn/dT data from Refs. 3 and 1, respectively.

The values in the table show a fast increase in loss with θ , which makes smaller values of θ advantageous, as before. Also, KD*P has substantially lower losses because of its smaller $d\Phi/dT$.

In summary, KD*P is preferable over KDP, primarily because of its much smaller transmission loss at 1053 nm, but also because its temperature sensitivity is substantially less. A 1-cm-thick KD*P wave plate oriented at $\theta = 30^\circ$ would have a switching angle $\Delta\theta = 0.15^\circ$, a transmission loss per pass of 0.5%, and a depolarization loss/pass of 0.16% for $\Delta T = 1^\circ\text{C}$, all of which are quite acceptable. Furthermore, the technology for finishing KD*P already exists and the facilities for scaling its size to 30 cm exist, as demonstrated by the existing KDP harmonic generators on Nova. If its damage threshold can be increased to that of KDP, KD*P would be a nearly optimum choice for this application.

Author: James E. Murray

6.2.6. Advanced Pulsed Power Components

Our current approach to meeting the aggressive cost goals for a future ICF driver pulsed power system requires some component development. This approach involves switching large amounts of capacitive energy storage into a single large transmission line. The line would terminate at the amplifier, fanning out to the many flashlamps. This design reduces costs by

reducing fabrication and assembly labor associated with pulling and terminating cables, and building individual pulse forming networks. However, this approach results in very high peak currents, and therefore large mechanical forces, on bank components during faults due to arcing in the bank or capacitor failures.

Therefore, component development is required in several areas to accommodate these stresses within our cost goals. In particular, development is necessary to produce low cost, high current capacitors, switches and transmission lines.

In addition to component development, a 2 MJ prototype capacitor bank module is being built under contract at American Controls Engineering. This bank will be used to verify simulation codes and simulate fault modes. It will also be used to address system issues and to develop and test transmission lines and other pulsed power components.

Capacitors. The cost of capacitors is driven by the materials costs. Therefore, the cost is reduced if less material is used to store the same electrical energy. This is accomplished by increasing the stress on the dielectric in the capacitor. There are various approaches which might be successful, such as improving dielectric uniformity and quality; reducing the number of airborne particles in the manufacturing area; and optimizing the thickness and number of layers of dielectric in the capacitor. There are also other approaches that show some promise but are proprietary to the vendors.

We have initiated a development and testing program with three vendors (GE, Maxwell Laboratories, and NWL Capacitors) to identify and improve the state of the art in producing low cost capacitors for the LMF. Each vendor was asked to provide units for testing which would meet the list of specifications in **Table 6-3**. This list was chosen to give the manufacturers as much flexibility as possible to produce innovative designs, while still allowing us to compare the test results of units from various vendors. The testing is in progress in our capacitor test facilities which have been modified for this purpose. Three tests will be performed on units from each vendor. A dc life test will address capacitor bank charge time issues that affects the cost of power supplies for the LMF bank. In this test, capacitors are charged to their rated voltage, and left charged continuously until failure. The number of hours to failure is then measured. The bank shot life will

Table 6-2. Values of $d\Phi/dT$ and single-pass depolarization loss for $DT = 1^\circ\text{C}$ for three waveplate thicknesses shown at two orientations of KDP and KD*P.

	θ	$d\Phi/dT$ (rad/ $^\circ\text{C}$)	Loss/pass for $\Delta T = 1^\circ\text{C}$		
			d = 0.5 cm	1.0 cm	2.0 cm
KDP	30°	-0.10	.06%	.025%	1.0%
	58°	-0.27	.46%	1.8	7.3%
KD*P	30°	-0.057	.02%	.08%	.32%
	60°	-.151	.14%	>57%	2.3%

Table 6-3. Preliminary LMF Capacitor Specifications.

Life:	10,000 shots, 90% confidence
Cost:	2-3 ¢/Joule
Reversal:	<10%
Current:	.3A/J
Voltage:	22KV
Rep. Rate:	One shot per 30 minutes
Faults:	Survive 5 shots at 80% reversal
Charge time:	2 minutes
Configuration:	Single Bushing
Capacitance & Energy	Can vary as necessary to meet other requirements

be addressed by a critically damped shot life test that simulates the electrical conditions of an LMF module. Resistor banks are used to approximate the flashlamp load impedance and the units are charged and fired once per minute until failure. Finally, a reversal, or ringing test will determine the level of tolerance of the units to faults in the bank. A fault is simulated by replacing the load from the shot life test with a low impedance load made from a stainless steel sheet. This results in an underdamped circuit and higher peak current. The percentage of reversal can be adjusted by varying the dimensions of the stainless steel sheet. This value will be chosen based on fault impedance data from experiments on the prototype module.

Partial shipments of capacitors have now been delivered by each vendor and some dc and shot life testing is complete. These initial results are shown in Table 6-4. The mean shot life is used as a convenient way to compare the performance of the vendors' units to date. Additional data and Weibull statistical analysis are required before we can accurately assess progress toward our 10,000 shot, 90% confidence level goals. However, the preliminary data indicate that, while these designs would meet our cost goals, more work is needed to improve the shot lifetime of these units. Funding is presently not available to address these issues with another design iteration.

Switching. A cost effective high voltage current and charge transfer switch is the present limitation to the size of the baseline energy storage module for the LMF. The present switch of choice is the ignitron, a vacuum tube with a liquid mercury cathode. This type of switch is used on Nova and has the advantages of low cost, high reliability and simple triggering requirements compared to other switch candidates. The Nova ignitrons, however, will not handle the peak currents required by the LMF module. A joint ignitron switch development effort was initiated 3 years ago with Y-Program, the LLNL Pulsed Power Thrust Area, Richardson Electronics and Texas Tech University. Prototypes of a commercial ignitron switch resulting from this work were tested this year.³² The results indicate that the new device will switch 500 kA, 200 C at 10 kV for several thousand shots. This is greater than a factor of two increase from previously existing tubes. The baseline requirements for the LMF module are

500 kA, 200 C at 22 kV for 10,000 shots. Thus, two switches in series would be required to meet the LMF module requirements unless the voltage hold-off capabilities are improved. Additional testing is also required to verify the lifetime projections. Further development in this area will resolve these remaining issues.

Author: Douglas W. Larson

Major Contributor: Phillip M. Test

6.2.7. Laboratory Microfusion (LMF) Alignment System Concepts

Efficient and accurate alignment of the laser beam from oscillator to target is essential for productive operation of an LMF laser system. The LMF laser could have 32 to 56 beams, each consisting of 16 individually pointed beamlets.

Requirements and Design Guidelines. To establish a functional alignment system for the Athena laser, a study was conducted to identify the requirements and desired features. Several existing systems were reviewed, the operators of these systems and, the designer and potential users of the LMF were consulted. The resulting system requirements and desired features are presented below.

Beamlet pointing. The beam must be properly pointed through spatial filter pinholes, or the beam will be asymmetrically filtered or blocked completely. To correct the pointing of a beam with respect to the pinhole, either the pointing of the beam can be adjusted with optics, or the pinhole position can be adjusted until the beam passes through the hole. For LMF, as with Nova, the beam placement accuracy requirement is a few millimeters separation from the beam edge to any obstacle.

Positioning the surrogate target. Since the real target can not be used for beam alignment, a surrogate target is used. Before an alignment operation begins, the surrogate target must be moved to the target chamber center with some type of positioning mechanism. The surrogate target should be placed at the target chamber center for two reasons: (1) only near the chamber center do all the beams hit a target symmetrically, and (2) the limited field-of-view of the diagnostic equipment is aimed at the target chamber center.

Beamlet pointing at the surrogate target. The alignment system must accurately place the beams to the desired point on the target

Table 6-4. Preliminary Capacitor Test Data.

Vendor	Energy Density (J/cm ³)	Energy per Unit (KJ)	Test Voltage (KV)	Shot Life		DC Life	
				# of Units	Mean Life (Shots)	# of Units	Mean Life (Hours)
A	0.49	50	22	6	18,800	2	630
B	0.90	40	20	7	3,800	2	64
C	0.55	25	25	3	330	*	*

* not tested

(surrogate target) by adjusting the focusing lens or final turning mirrors. The desired pointing accuracy for LMF has not been specified.

However, it is presumed to be about five times the precision Nova specification of 30 μm RMS (i.e. 150 μm RMS from the aim point) based on simple scaling of target size with energy.

Desired Features. Low cost is a primary constraint on the alignment system design. First, implementing multiple-use hardware whenever possible, using the same equipment on some or all of the beamlets will significantly reduce cost. In general, the best opportunity for multiple use of equipment is in a laser front-end or the target chamber center where all the beams are together. LMF will conceivably have two orders of magnitude more components for alignment than Nova. Therefore, these components should be designed to minimize operating and maintenance costs.

High accuracy in placing the laser beams on the target is a significant factor in improving target performance. There are several features of an alignment system that promote higher accuracy:

- A minimal number of operations is better since each alignment operation can contribute to alignment errors.

- Direct sensing of events eliminates the errors associated with remotely observing an event. For example, sensing beam position with a camera at the desired focal spot in the target chamber rather than observing beam scatter off a reticle with a remote imager.

- No compensation of the alignment beam path avoids the alignment errors it can induce. For example, some concepts use lenses or mirrors to account for chromatic differences of an alignment beam and the main beam.

- Use of the laser front-end to generate the alignment beams eliminates the need for alignment lasers. If no alignment lasers are used then no co-alignment of beams with the main beam or other alignment beams is necessary. Consequently, accuracy should be improved since co-alignments can be a source of alignment errors.

- Position verification of alignment parts, such as a lens, mirror,

etc., injected into the main beam insures that these parts do not cause alignment errors. Exceptions would be alignment insensitive parts such as image screens where content, not position, are important.

A short alignment time is another area of consideration. A short alignment time would have two benefits. If it were sufficiently short, alignment time would not affect the time between shots. If the alignment can be performed before other in-between shot tasks, e.g., amplifier cooling which typically takes over one hour in Nova, then the alignment time will not affect shot to shot time. Secondly, a short time between the start of alignment and the shot minimizes the opportunity for alignment drift due to vibrations, creep, temperature drift, etc. Ideally, an entire alignment would be performed just prior to a shot. Several features could shorten initial alignment, such as simultaneous alignment of all beams (32 to 56) or beamlets (512 to 896). Also, the automation of the entire beam and target alignment process would reduce time considerably since it will make possible the performance of many operations in parallel. For example, simultaneously sensing all the beam pointing errors on the pseudo target and moving all the final turning mirrors to correct these errors at one time.

Finally, a visual display of the operations is very desirable, as proven with Nova. Even when the operation is automated, visual monitoring is helpful in verifying the proper functioning of the system. If an error in the automation system occurs, it can be overridden and controlled manually with the assistance of the visual monitors.

Conceptual Design. During this past year conceptual design approaches have been investigated, in order to develop a system capable of minimizing the number of sensor packages required, therefore, reducing the initial capital cost and the operation and maintenance cost.

For the purpose of evaluating alignment system concepts, the basic layout of the Athena laser design was assumed to be as shown in Fig. 6-33.

This Athena multipass laser has a master oscillator that provides the input to a single preamplifier; the preamplifier is isolated at both input and output by spatial filters. The filtered output of the preamplifier is divided

Figure 6-33. Schematic layout of a single beamlet of the Athena multipass laser design.

and routed to a set of multipass amplifiers by means of a beam multiplexer. For the LMF design, the number of multipass amplifiers is between 32 and 56.

Each multipass amplifier encompasses two gain regions located at the opposite extremities of a 30-m cavity. The central 15 m of the cavity comprise a multipass spatial filter. Tilt of one end mirror of the cavity causes successive passages of the beam to be displaced one from another by an angle on the order of a milliradian. This serves two purposes. First, it permits the use of multiple aperture stops in the focal plane of the spatial filters; each stop being traversed only once by a multipassing high-power pulse. This is necessary because the plasma generated within a stop by the passage of a pulse exhibits a relaxation time which is long compared to the round-trip time within the cavity. Second, it provides a convenient geometry for injecting a pulse from the multiplexer. Kickout from the amplifiers is via a plasma-electroded Pockels cell and polarizing beamsplitter. A final spatial filter leads to a frequency tripling system and to the optics that focus the third harmonic beam upon the target.

One concept to be discussed below was proposed by Laser Power Research, Inc. (LPR) in a study funded by LLNL.³³ The LPR approach uses a double-pass interrogation scheme that employs laser range-finder technology to provide time-separated error signals from multiple locations in an Athena beamline. Using this technology, the proposed concept permits the alignment of the multiple beamlines proposed for an LMF Athena design by using a single alignment subsystem capable of characterizing the pointing, focus and centration errors of all key elements of the laser system, in less than 200 seconds per iterative step.

For this concept, an alignment beam is inserted via a beam splitter immediately beyond the output of the master oscillator. A schematic of the proposed alignment path is presented in Fig. 6-34.

By using a two-axis optomechanical beam segmentation device, the alignment beam can be translated from segment to segment of the 4×4 array that comprises each Athena beamline, while maintaining perfect collinearity of the optical path.

The collinear master oscillator and alignment beam is passed through the front end preamplifier and spatial filters, and is then sent

to beam splitters where beamline shutters direct the alignment beam sequentially through each beamline. The alignment beam is diverged with respect to the main beam during centering alignment and converged during pointing alignment. Retroreflective reference devices are located at key locations in the beam path, i.e. pinholes, spatial filter lenses, and the target chamber center. Certain reflectors reflect the alignment beam back to the front-end of the laser when the alignment beam is diverging and other reflectors when the beam is converging. The return time of the signals distinguishes the location of the reflector in the beam path. The reflected beam is evaluated by a front-end camera to determine the centering, pointing, and focusing of the beam at those locations.

To center the beams through the system, the alignment beam is converged slightly compared to the main beam. The convergence results in an expanded beam as the beam exits the second lens of a spatial filter (see Fig. 6-35). The expanded beam is sampled at several points where the beam exceeds the normal diameter. The samples are taken with small scraper mirrors. The sampled light is reflected off retroreflectives and back tracks its path to the laser front-end. The camera at the front end compares the magnitude and symmetry of the return signals with the expected return to determine the centricity of the beam. Likewise, a diverging alignment beam is injected to center the beam at the first lens in a spatial filter. All the reflected samples from one beam are sensed simultaneously and distinguished temporally. Because the beam edge samplings are polar offset, the samples from the various locations on one beamline can be simultaneously observed.

For pointing beams through the pinholes, a pinhole plate is moved into the spatial filter focus plane for each laser amplifier spatial filter. The plate will have a main beam pinhole, an open hole, and a meniscus lens (see Fig. 6-36). For pinhole pointing alignment, the meniscus lenses are positioned at the expected focus of the main beam. The alignment beam is slightly converged at the front-end to match the radius of the meniscus lens curvature. If the alignment beam is correctly pointed, the beam reflects from the lens precisely upon itself. The returned beam is imaged on the camera. An improper focusing or pointing of the beam at the pinhole is sensed and corrected by moving the end cavity mirrors. Most of the alignment beam

Figure 6-34. Schematic for the Laser Power Research alignment concept.

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Figure 6-35. Centering from reflected samples of the beam's edge.

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Figure 6-36. The multiholed plate inserted into each pinhole location.

energy passes through the meniscus lens to the next pinhole where it also provides a return signature. All the meniscus lens reflections are distinguished temporally. Therefore, as with centering returns, all the returns from the pinholes on a beamline can be observed simultaneously. The open hole in the plate is to pass a converging beam to the target chamber.

After completion of the beam alignment through the multipass amplifier cavity, the beam is passed through the optical switch and then guided by a turning mirror to the output spatial filter. The same sensing techniques described for the cavity spatial filter is used for this component and alignment is achieved by controlling the turning mirrors and possibly the movement of the spatial filter pinholes. Following the alignment of the output spatial filter, the beam is guided by the three or four turning mirrors to the frequency convertors and final focusing lens that align the individual beams to the target assembly.

In a separate study conducted by LPR³⁴, a target position sensing technique was proposed to position a surrogate and real target assembly at the target chamber center which can then be used to perform final pointing and alignment of the multiple laser beams to the target.

The technique proposed to position the target assembly used a wide-band interferometer, taking advantage of the matched path length requirement of a white-light Michelson interferometer. For this application, the white light takes the form of a wavelength scanned laser beam. By matching the reference arm of the interferometer to the nominal target range, the motion of fringes in the measurement arm provides both the direction and magnitude of target displacement errors. When a bandwidth of 10 nm is employed for such a sensor, the target assembly can be positioned within the micron regime.

Alignment of the laser beams to the target has, in the past, been accomplished by using sensing signals from reflecting spheres. A novel surrogate target concept has been proposed by LPR that involves the use of index-matched dielectric spheres; the properties of which are suited both for position sensing and for beamlet alignment. By virtue of their retroreflective nature, their use permits initial beamlet misalignments to be far greater than is permissible with more conventional targets.

The use of a surrogate target for direct alignment and position sensing while in close proximity to the target itself has led to the desirability for a reference element of minimal mass. The concept proposed uses an optical reference element for the surrogate target which can be separated by 50–100 mm relative to the real target; while, at the same time, exhibiting a mass of only 25 mg. This, in turn, permits the reference element to be mounted within the slender structure of the target support.

The target positioner and alignment subsystem, as shown in Fig. 6-37, combines the use of a side-band interferometric distance sensor, a HeNe differential range finder, and a two-axis centroid displacement sensor with a gravity vector mirror and a pair of servo-driven Resley prisms for the purpose of measuring target offsets. The interferometric distance sensor uses a tunable laser source in conjunction with a Lissajous-type reference cavity. Sensing of both surrogate targets and real targets is direct. The resulting positional accuracy for both surrogate and real targets is 7 μm RMS in each of three orthogonal axes. For surrogate targets, this accuracy is maintained throughout a 20-mm spherical volume which is centered within the chamber.

This system requires that a single 100-mm-diam perforation be made at the uppermost or lowermost point of the Athena target chamber; thereby allowing a single sensor package to be collocated with the target positioning mechanism. This arrangement is ideal from the standpoint of providing tight closed-loop control of the positioning mechanism. It also permits position sensing to be performed via an optical path that lies wholly within the target support arm. The number of first-order structural constraints imposed upon the system is three; these being associated with the degree of translational freedom available to the sensor support pier.

Author: R. J. Foley and C. S. Vann

Contributors: G. Flint, Laser Power Research Corp., San Diego, CA

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Figure 6-37. Schematic of a Monostatic Position Sensing Subsystem which combines a Wide-Band Laser Rangefinder with a Two-Axis Centroid Displacement Sensor.

6.3. Laser Optical Damage Studies

6.3.1. Introduction

The ability of optical materials to resist laser damage continues to be a major factor in the development of high-peak-power lasers used in ICF research. The reason for this is largely economic; the higher the damage threshold, the smaller and hence less costly the optical component.

For years one of the weakest links in the laser chain (in terms of damage) has been multi-layer coatings. During this past year we have carried out a three part program to improve the damage resistance of optical coatings. The key results from each of the three parts of the program are discussed in detail later in this section. In this introduction we will simply highlight these results.

One part of our effort has been to improve the existing multi-layer coatings that are prepared by e-beam evaporation. During this past year we have shown that the damage thresholds of e-beam deposited multi-layer $\text{HfO}_2/\text{SiO}_2$ thin films can be permanently increased by a factor of 2 to 3 by illumination with subthreshold fluences of laser light. This subthreshold illumination procedure is referred to as "laser conditioning" (see Sec. 5.3.3). The films used in our study were prepared by three different physical-vapor-deposition techniques: ion-beam sputtering, plasma plating and e-beam evaporation. Only the e-beam deposited films showed consistent and significant improvement with laser conditioning. Of the material pairs examined ($\text{HfO}_2/\text{SiO}_2$, $\text{ZrO}_2/\text{SiO}_2$ and $\text{TiO}_2/\text{SiO}_2$), $\text{HfO}_2/\text{SiO}_2$ gave the greatest and most consistent damage improvement with conditioning. The number of layers and the reflective or transmissive characteristics of the $\text{HfO}_2/\text{SiO}_2$ films were found to have little impact on laser conditioning of the film. Furthermore, the results show that the damage thresholds of a wide range of e-beam deposited coatings (HRs, polarizers, etc.) can be improved by laser conditioning. Based on this research, we have developed a technique for conditioning large mirrors and we are currently using it on Nova.

The second part of our coating effort has focused on the development of high quality, damage resistant, sol-gel coatings. Our study of highly reflective (HR) dielectric coatings

prepared by the sol-gel process has indicated that hydrated alumina and silica are the materials of choice for the high and low index components respectively. We can now prepare samples with up to 36 quarter-wave-layers by spin coating at room temperature colloidal suspensions of the relevant oxides. These have reflectivities of about 99% and damage thresholds in the range $30\text{--}40\text{ J/cm}^2$ at 1064 nm with 16-ns pulses at 30 Hz . This threefold improvement in threshold over that reported earlier for this system was attributed to an improvement in the preparative technique particularly in regard to cleanliness.

The third part of our research program has focused on the development of a novel, damage resistant rugate coating. These rugate multi-layer dielectric optical coatings, nominally consisting of 1000 or more optical-quarterwave layers, are prepared by reacting SiCl_4 and a halogenated dopant (e.g. GeCl_4) with O_2 in a microwave or RF-driven plasma. The dopant concentration is such that the refractive index different between adjacent quarterwave layers is about 0.02 or less. The deposition is carried out at high substrate temperatures ($850\text{ to }1100^\circ\text{C}$) producing a fully dense, fused silica coating. Surface damage thresholds of high reflectivity coatings prepared by this plasma process are comparable to those for fused silica. For example, at $1.06\text{ }\mu\text{m}$ and 16 ns we measure surface damage thresholds greater than 45 J/cm^2 compared to about $45\text{--}60\text{ J/cm}^2$ for super-polished optical fused silica.

During this past year we have designed and built a high-temperature, plasma assisted, chemical vapor deposition system to deposit the multi-layer rugate coatings of SiO_2 and doped- SiO_2 on flat substrates. The coater is designed to deposit up to several thousand alternating quarterwave layers of SiO_2 and doped SiO_2 on SiO_2 substrates at deposition rates up to several microns per minute. The substrate is resistively heated to about 1000°C during the deposition phase of the process. The plasma is driven by a 13.56 MHz RF unit capable of producing power densities of up to 140 W/cm^3 in the reaction zone. The coater is designed to be adaptable to microwave generated plasmas, as well as RF. Reactant gas flow rates of up to 10 slm can be achieved at a 10-Torr operating pressure. Reactants consist of O_2 , SiCl_4 and a volatile halogenated dopant (e.g. GeCl_4 or CCl_2F_2).

These gases react in the plasma volume producing SiO_2 with dopant concentrations (e.g. GeO_2) of up to a few percent. A variable dopant concentration is used to produce index differences (nominally about 0.01) between adjacent optical layers.

Some of the major technical issues associated with the design and construction of this coater include: flow distribution and its effect on deposition uniformity, substrate heater design and operation, choice of RF versus microwave driven plasma and diagnostics design and placement. To address many of these issues, we have relied heavily on numerical and experimental modeling of the fluid flow, reaction kinetics, and heat transfer within the coater.

In addition to our three-part program on optical coatings we have also carried out a study of composition effects in phosphate laser glasses. This composition study was of 41 phosphate glasses which are derivatives of a commercially available Nd-doped laser glass. Systematic variations were made in the quantities of alkali and alkaline earth components in the glass. Correlations were developed relating the compositionally averaged electric field strength and certain thermal, mechanical, optical and laser properties of these glasses. We found that the compositionally averaged field strength is a good parameter for predicting the effects of complex mixtures of alkali and alkaline earth components in the glass. However, it is inadequate for explaining the effects of glass components which are predominantly covalently bonded with oxygen.

Author: J. H. Campbell

6.3.2. Sol-gel Coating Development

Our major effort on sol-gel coatings has continued to be directed towards the preparation of highly reflective dielectric coatings for use at 1064 nm. For nearly six years we have prepared single layer antireflective coatings by a sol-gel process but high laser damage threshold highly reflective coatings had not yet been achieved. Our investigation throughout has been based on the use of colloidal suspensions as the coating media. These have several advantages over other materials. Suspensions are liquid and can be rapidly applied at room temperature by

conventional means. They require no cure or heat and are therefore well suited for large optics and the capital equipment cost is low. Finally, coatings can be easily removed and replaced with no damage to the substrate.

Our HR coatings consist of alternating layers of high and low refractive index oxides. Our investigation has indicated that AlOOH and SiO_2 are the best materials to use for this system. Colloidal suspensions of SiO_2 are readily prepared by the base catalyzed hydrolysis of $\text{Si}(\text{OC}_2\text{H}_5)_4$ in ethanol.³⁵ We continue to prepare single layer AR coats from this material on substrates up to 1 m in diameter. The laser damage threshold is very high³⁶ and this therefore becomes a natural selection for the low index component of our HR system. Although AlOOH has only a moderately high index (1.62) it was selected over several other higher index candidates such as HfO_2 and TiO_2 mainly because it had the highest laser damage threshold when deposited as a single layer. Colloidal suspensions of AlOOH are readily prepared by the hydrolysis of $\text{Al}(\text{OC}_4\text{H}_9)_3$ in acidified water.³⁷

We prepared our HR coatings on fused silica or BK-7 substrates by depositing alternating layers of AlOOH and SiO_2 from the relevant colloidal suspensions using a spin coater. **Figure 6-38** shows the reflectances obtained for a variety of samples containing from 26 to 40 layers. It can be seen that at least 34 layers are required for 99% reflection. **Figure 6-39** shows the transmission curve of a fused silica substrate coated with 36 layers. The reflectance of this sample is greater than 99% at 1064 nm. The narrow band third harmonic reflectance of about 95%, which is characteristic for this type of quarterwave stack, is also present.

The laser damage threshold measurements at 1064 nm (16 ns pulse, 30 Hz) range from 21–50 J/cm^2 with an average of 37 J/cm^2 and have no apparent dependence on the number of layers. We had previously reported damage thresholds in the range 5–21 J/cm^2 with an average of 12 J/cm^2 .³⁸ While the reason for this is not known for sure, we have recently taken more rigorous precautions to exclude contamination by careful filtration of the suspensions and the use of a clean room for coating fabrication.

Author: Ian M. Thomas

<u>Layers</u>	<u>Reflection at 1.06 μm</u>	<u>Damage 1064 nm/16 ns</u>
26	95.0 %	34 J/cm ²
30	98.0	45
30	98.5	28
32	98.0	38
32	98.5	50
34	99.0	21
36	99.0	37
40	99.5	40

Figure 6-38. Reflectance and damage thresholds of several $[\text{SiO}_2\text{-AlOOH}]^n$ HR coatings.

Figure 6-39. Transmission spectrum of [SiO₂-AlOOH]₁₈ coating.

6.3.3. Plasma Chemical-Vapor-Deposited Coatings

High temperature plasma assisted chemical vapor deposition opens the possibility to fabricate unique optical components for lasers³⁹ with high laser damage threshold.⁴⁰ The challenge is to transform this inside-tube technology to coat flat substrates. Our group has approached this problem in two parallel efforts. Workers at Schott Glaswerke, Mainz, FRG (all under contract to LLNL) are designing and building a new high temperature flat coater to use microwave (2450 MHz) energy to excite the plasma. Here at LLNL, we have designed and constructed a high temperature coater using RF (13.56 MHz) to excite the plasma. We chose to build a coater based on 13.56 MHz because it would in principle be scalable to a process suitable for coating large optics without major design considerations for plasma applicators.

The coater is designed to deposit up to several thousand alternating quarterwave layers of SiO₂ and doped-SiO₂ on SiO₂ substrates at deposition rates up to several microns per minute. The substrate is resistively heated from below and the plasma is driven by a 13.56 MHz RF generator capable of producing power densities of up to 140 W cm⁻³ in the reaction zone. Reactant gases (O₂, SiCl₄, GeCl₄ and CCl₂F₂) with a total gas flow rate of 10 slm can be processed at 10 Torr operating pressure.

Preliminary deposition experiments were conducted at Arizona State University using a radiantly heated tube furnace (Fig. 6-40). Due to coupling to the plasma, we were unable to determine the substrate temperature accurately using thermocouple sensors. Furnace power and gas flow control were sometimes affected by the electric fields of the radio transmitter. Nonetheless, single layer coatings up to 1.5-μm thick were obtained using the experimental parameters listed in Table 6-5. We found several important issues that the coater design must address including the need for good confinement of the plasma power for higher power densities, a hot enclosure surrounding the substrate to suppress high stress cold wall coatings and particulates, and good diagnostics of system parameters such as temperature, furnace power, plasma power and voltage.

Prototech Research Inc., in Chandler, AZ, was awarded a contract to design and build a research PCVD coater to operate within LLNL

specified goals coating a 15-cm-diam, 1-cm-thick silica substrate at 1000°C substrate temperature. This design addressed all the issues raised in the preliminary coating effort described above and is flexible enough to switch to a microwave plasma driver in the future, if necessary. The coating chamber is shown schematically in Fig. 6-41 and consists of a parallel plate RF electrode assembly for uniform plasma generation. The driven electrode is also the enclosure for the platinum/rhodium wire wound furnace. Precursor gases are admitted through an array of holes above the substrate and exhausted out the periphery. Temperatures are monitored by optical pyrometers in the furnace as well as at the substrate surface. The substrate is surrounded with opaque silica which is heated by absorbing infrared radiation from the furnace.

We successfully modeled the mass transport of the gas inlet design⁴¹ by simulating the deposition of a film on a substrate using the evaporation of naphthalene, a volatile solid, in a scale model of the PCVD reactor chamber. These modeling results demonstrate a uniform mass transport across the surface to ± 5% for a given set of flow and pressure parameters suitable for coating operations. This experimental mass transport behavior was compared with stagnation flow theory assuming uniform gas injection through the distribution plate above the substrate. As shown in Fig. 6-42, the experimental mass transfer rate agrees with theory below a Reynolds number of approximately 3. Enhanced mass transfer above this value is the result of localized jet impingement on the substrate. The Reynolds number is defined as,

$$Re = \frac{HV}{\eta} \quad (1)$$

where H is the distance separating the gas injection plate and substrate (cm), V is the superficial axial injection velocity (cm/sec) and η is the kinematic gas viscosity (cm²/sec).

The deposition chemistry as well as the importance of the plasma excitation frequency on the resulting thin films is unknown at this time. To begin addressing these problems, we have constructed a coating plasma test bed with the flexibility to accept antennas from a 1-kW RF generator at 13.56 MHz and from a 2.5-kW microwave generator operating at

Figure 6-40. Schematic of radiantly heated horizontal plasma CVD reactor.

Table 6-5. Summary of typical operating conditions used in our preliminary radiant heating plasma CVD experiments.

Pressure	0.3–3.0 Torr
Flow rate oxygen	100–500 sccm
Flow rate SiCl ₄	20–90 sccm
Flow rate GeCl ₄	0–10 sccm
RF power (CW exp.)	300 W
RF power (peak)	57–94 W
RF power (average)	8.5–15 W
Pulse width	10–15 × 10 ⁻³ s
Pulse Repetition Frequency	10–30 Hz
Temperature	1050–1250°C
Reaction time	60 min
Number of experimental runs	16

Figure 6-41. Schematic of high temperature RF-PCVD coater reaction chamber.

Figure 6-42. Dimensionless mass transfer rate versus Reynolds number.

2450 MHz in either pulsed or continuous wave operation. The test apparatus has been built to accept a Langmuir probe, an optical multichannel analyzer with both temporal and spatial resolving capability and a residual gas analyzer.

Author: Raymond M. Brusasco

Major Contributors: Jerald A. Britten, John H. Campbell, Donald W. Hughes, Mark W. Scrivener, Charles B. Thorsness, and Walter G. Unites

6.3.4a. High Temperature, Plasma Chemical-Vapor-Deposition of Optical Coatings

Introduction (Refs. 42-54). In this section, we discuss the results of experiments using high temperature ($T > 850^\circ\text{C}$) plasma-assisted chemical vapor deposition to produce damage resistant optical coatings. The coatings typically consist of several thousand quarter-wave layers of fused SiO_2 lightly doped with a glass network modifier or network former (e.g. GeO_2 , F) to give an interlayer index variation.

Currently, most multi-layer dielectric optical coatings used in high power laser systems are prepared using physical vapor deposition techniques (e.g., e-beam evaporation) and typically consist of a stack of alternating layers of metal oxides and/or fluids. Although the exact mechanism for laser induced damage in these coatings is still widely debated, there are a number of inherent weaknesses in the deposited structure that have been proposed to lead to optical damage.

It is well known that optical fused-silica produced by a high temperature chemical vapor deposition (CVD) process, such as that used by Corning, has high thresholds for both surface and bulk laser damage⁵¹⁻⁵³. In the Corning CVD process, SiO_2 is formed by vapor-phase hydrolysis of SiCl_4 and then deposited layer-upon-layer until a boule of SiO_2 is formed.⁵⁴ Thus, in general sense, the SiO_2 is a monolithic CVD "coating."

We believe that the reason for the excellent damage resistance of optical fused silica is the lack of potential damage causing weaknesses associated with dielectric coatings prepared by physical vapor depositions. First, the material is compositionally uniform. Consequently there are no distinct material interfaces with

thermal-mechanical or chemical incompatibility. Second, the chemical stoichiometry is uniformly controlled throughout the material by use of high-temperature, oxidizing deposition conditions. Third, any dust, dirt, or other particulate contamination is fully oxidized and either dissolved directly in the glass or incorporated as an insignificant, low absorption, scattering center. Finally, the fused silica is a fully dense, amorphous body without the high density of microstructural defects that are often produced by conventional physical vapor deposition processes.

Figure 6-43 illustrates our concept for preparing optical coatings by high temperature CVD. In brief, various low concentrations of either network modifiers or network formers are added to alternating layers of the silica glass matrix. Doping the SiO_2 with these materials produces the index difference needed to achieve the desired optical characteristics, yet the dopant concentration is low enough that the bulk SiO_2 glass remains largely unaltered. Thus there are no distinct material interfaces. The coating is produced under the same high temperature ($>1000^\circ\text{C}$), oxidizing conditions used for preparing bulk SiO_2 and is therefore amorphous, fully dense, with very low optical absorption and essentially defect-free. Since the dopant concentration is low, a large number of layers (several thousand) are required to produce the desired optical effect.

Selection of a CVD Coating Process. The CVD processes used to prepare the preforms for silica optical fabrication can, in principle, also be used to prepare optical coatings such as those described above. These processes can generally be broken down into two types. One is the so-called "soot" process in which the silica precursor (SiCl_4) is reacted thermally with O_2 to produce fine (0.1 to 1.0 μm) particles of SiO_2 ; the particles then deposit on the substrate wall and are subsequently sintered to a fully dense, amorphous SiO_2 monolith.⁵⁵⁻⁵⁷ The second general deposition process is plasma driven. In this case the reaction of the precursor and O_2 are driven by a microwave or RF initiated plasma.⁵⁸⁻⁶⁰ The plasma processes have the advantage that the thickness of a single deposition layer can be controlled to the angstrom level. In contrast, the layer thickness resolution of the "soot" process is at best several SiO_2 particle-diameters (a few tenths of a

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Figure 6-43. Conceptual view of multilayer optical coatings prepared by chemical vapor deposition of alternating layers of doped and undoped fused silica.

micron) and thus impractical for deposition of typical quarter-wave thick layers used in optical coatings.

Of the current plasma process, the one having the best thickness control at high deposition rate is the Schott plasma-impulse chemical vapor deposition process (PICVD).⁶⁰⁻⁶² This is the process we have selected for demonstration of the optical coating concept outlined here. The coatings prepared during the course of this study were deposited on the inside of a 1.7-cm SiO₂ tube since that is the configuration of the presently available plasma CVD systems.

In brief, the coating system consists of three major components: a gas supply system, the deposition unit, and vacuum pump (Fig. 6-44). All three units are under computer control. Reactant gases are continuously supplied to a 1.7-cm-i.d. × 65-cm-long quartz glass tube that serves as both the reaction vessel and deposition substrate. Immediately surrounding the quartz tube is coaxially positioned microwave cavity. This entire assembly is located in a tubular furnace that typically heats the substrate to 800–1100° C. Further details of the process and system are described elsewhere.⁶⁰⁻⁶²

Typical deposition conditions used for these experiments are given in Table 6-6. The plasma was driven in a pulse mode. Deposition rates were generally about 2 μm/min over the length of the tube (~50 cm); this corresponds to single-pulse layer thickness of about 30 Å at the nominal 100 Hz pulse repetition rate. The substrate temperature was usually kept at about 1000° C to reduce inclusion of product gases in the film and enhance the film quality.

The multi-layer optical coating was prepared by periodically mixing SiO₂ precursors gas (i.e. SiCl₄) with a halogenated dopant gas. Dopant gases were chosen because they are known to give stable silicate glasses, and they produce the desired index change without producing unacceptable thermomechanical stresses.⁶³ In most cases, co-dopants of Ge and F were to produce the desired index change.

Optical Characterization of Coatings Prepared by PICVD. The measured transmission of two typical PICVD coatings is shown in Fig. 6-45. The peaks in the curves (i.e. regions of low transmission) indicate reflection bands. The first coating (Fig. 6-45a) consists of approximately 1000 optical quarterwave layers with a

sinusoidal index variation and a maximum Δn of about 0.01 between adjacent layers. The layer thickness was adjusted to give a single maximum reflectance at 1064 nm. The shoulder on the short-wavelength side is due to small process variabilities. In general the measured bandwidth of the PICVD coatings is about 2–3 times the theoretical value. This suggests that with future improvement in process control, narrower bandwidth coatings with fewer layers should be possible.

The coating shown in Fig. 6-45b consists of about 5000 quarterwave layers with roughly every 1000 layers designed to reflect a different wavelength (hence 5 reflectance maxima). In this case, the tail on the short-wavelength side is caused by bulk absorption by the glass matrix and is not due to process variability. Since the deposition process is fully automated, it is possible to create coatings having any number of peaks either fully or partially reflecting. Moreover, by closely grouping the peaks such that they overlap, reflective coatings of any arbitrary bandwidth can be achieved.

The overall absorption loss⁶⁰ in optical thin films produced by the PICVD process is near the theoretical limit of that expected for the pure bulk material. The reason for this is clear; these coatings were made using the same deposition process used to prepare the full-density optical preforms used to fabricate optical fibers. The reported loss spectrum for the GeO₂-doped SiO₂ material used to prepare these coatings is approximately 1 dB/km at 1.0 μm corresponding to $\sim 10^{-6} \text{ cm}^{-1}$. This absorption loss is about 5 orders of magnitude lower than that for high quality dielectric thin-films produced by current e-beam evaporation processes.^{48,64}

Compositional and Structural Characterization. We have analyzed the dopant concentration profile through several of the PICVD layers using secondary-ion-mass-spectrometry (SIMS). Figure 6-47 shows one example of SIMS data from the top approximately 12–13 layers of a 2000 layer Ge-F co-doped SiO₂ coating. The dopant concentration variation gave a peak-to-valley refractive index differences of 0.016. The SIMS depth resolution is about 4 nm near the coating surface and becomes gradually poorer with increased sputter depth. Nevertheless, the individual coating layers are clearly resolved. The lobes seen on the leading and trailing edges of the first few Ge-doped layers are probably real and a

Figure 6-44. Schematic of the PICVD deposition system.

Table 6-6: Typical deposition conditions for preparing multilayer doped SiO₂ coating using Schott's PICVD process.

Substrate:	SiO₂ Heralux WG
Substrate geometry	
diameter:	1.7 cm i.d., 2.0 cm o.d.
length:	>50 cm
Substrate temperature:	850–1100° C
Pressure:	3 mbar
O₂ mass flow:	200 sccm
SiCl₄ mass flow:	50 sccm
Dopant flow:	<10 sccm
Dopants:	GeO ₂ , F
Dopant source:	GeCl ₄ , CCl ₂ F ₂
Pulsed plasma repetition rate:	100 Hz
Microwave frequency:	2.45 GHz
Microwave power (avg.):	1.5 kW
Deposition rate:	~2 μm/min.
Incremental thickness:	~30 Å
Index profile:	sinusoidal or square wave
Quarterwave thickness:	1800 Å
(at 1.06 μm)

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Figure 6-45. Measured transmission versus wavelength for (a) 1000 and (b) 5000 quarter-wavelength optical coatings prepared by the PICVD process. The dopant profile was sinusoidal in both cases with the amplitude of the index variation about 0.01.

Figure 6-47. Results from SIMS analysis of the top 12–13 layers of a F, Ge co-doped SiO_2 coating.

result of variability in the dopant mass flow controllers during this particular run.

Damage Threshold Measurements. Most damage threshold measurements of the PICVD coatings were carried out at 1064 nm, with a 16-ns pulse width and 30 Hz repetition rate. The details of the damage test system are described elsewhere.⁶⁵

Recall that the PICVD coatings were deposited on the inside of a 17.2-mm fused silica tube as described above. In order to properly test the coating, the tube was cut lengthwise into four pieces. The laser beam was directed onto the inner coated surface of the sample at about 10° angle of incidence. Not all of the coatings that were damage tested were designed to reflect at 1064 nm leading to significant transmission through the sample. This often caused damage at the output surface and significant retroreflections. Thus, we feel our measurements represent a lower bound or conservative value of the damage thresholds that can be expected for the PICVD coatings.

Figure 6-48 compares the measured damage thresholds of the PICVD coatings with the surface damage thresholds of super-polished (Zygo) fused silica and Ge-doped fused silica prepared using very high temperature (~1800° C) chemical vapor deposition. The undoped fused silica is the same high quality material used to fabricate lenses, debris shields, and other transmissive optical elements on the current 120 kJ LLNL Nova laser. Fused silica has the highest surface damage threshold of any material or coating on the Nova laser and consequently we use it as a "bench mark" to compare with new optical materials. As expected, the kilolayer PICVD coatings are nearly equivalent in damage threshold to the surface threshold of super-polished fused silica. The pulse length scaling for damage to fused SiO₂ is approximately:

$$D_T = 19 t_p^{0.4}$$

where D_T is the damage threshold (J/cm²) and t_p is the laser pulse width (ns).

Unfortunately, we only tested one PICVD coating at 1 ns where it gave a damage threshold of about 13 J/cm². This is somewhat lower than fused silica but is probably inconclusive without further tests (i.e. more sample statistics).

Conclusions. High reflectivity coatings, nominally consisting of 1000 to 2000 optical-

quarterwave layers have been prepared by reacting SiCl₄ and a halogenated dopant (e.g. GeCl₄) with O₂ in a microwave-driven plasma. The dopant concentration is such that the index difference (Δn) between adjacent quarterwave layers is about 0.02 or less. The deposition is carried out at high substrate temperatures (about 900 to 1100° C) producing a fully dense, fused silica coating. The motivation for using the high temperature conditions is the observation that commercially available CVD deposited fused silica (used for making focusing lenses and other high fluence laser optical components) has very high surface and bulk damage thresholds.

Measured damage thresholds of HR coatings prepared by this plasma CVD process are comparable to those for super-polished fused silica surfaces. For example, at 1.064 nm and 16 ns we measure PICVD coating damage thresholds of about 45 J/cm² compared to about 45–60 J/cm² for the surface of super-polished Corning optical fused silica. The plasma CVD coatings have the highest N-on-1 damage threshold of any optical coating tested to date by the LLNL Laser Program.

Author: J. H. Campbell

Major Contributors: J. L. Emmett, R. M. Brusasco, and F. Rainer (LLNL); R. Th. Kersten, V. Paquet, and H. -W. Etzkorn (Schott Glaswerke).

6.3.4b. Fluid Dynamics And Mass Transfer Characteristics In Radial-Flow, Multijet PCVD Reactors

The plasma-assisted chemical vapor deposition (PCVD) process for making multi-layer reflective coatings on optical substrates involves a complex interaction of plasma physics, gas-phase and surface chemical reactions, and flow and transport phenomena in the reactor. Briefly, the PCVD process seeks to build a laser damage-resistant reflective surface, tuned to a particular wavelength, by depositing alternating layers of material of slightly differing refractive index [e.g. layers of pure silica (SiO₂) alternating layers of material of slightly differing refractive index or fluorine] on an optical quality flat silica substrate. The dopant amount is small to reduce the stress in the films. The layer thickness must be precisely controlled, and the number of layers sufficient (on the order of a thousand) to create a reflector

Figure 6-48. Measured damage threshold at 1064 nm versus pulse width for super-polished, bare fused silica and coatings prepared by either PICVD or very high temperature (~ 1800 C) CVD. The data for the bare fused silica are from Refs. 51, 52 and 53. The PICVD coating samples consisted of 1000 or more quarter-wavelengths of doped SiO₂.

capable of >99% reflection of the desired wavelength of light.

The goal of the present effort is to use a high-temperature, radio-frequency (RF) coater to produce such a mirror with a damage fluence threshold of $>40 \text{ J/cm}^2$, with damage and reflective properties uniform over a 15 cm scale. The RF plasma initiates chemical reactions that led to SiO_2 formation in a reaction volume continuously swept by SiCl_4 and O_2 at a relatively low pressure (about 1 mbar).

A schematic of the reaction chamber of this coater is shown in Fig. 6-49. Reactant gas is introduced through ≈ 150 equally spaced small holes drilled into an upper injector plate. The exhaust gas flows radially out through holes drilled into the cylinder defining the reactor sidewall. The solid reaction product deposits on all reactor surfaces including the substrate. The number and size of the gas injection holes is determined such that the pressure drop through these holes is much larger than the radial pressure drop in the reaction volume, insuring an equal flow of gas through each hole.

It is clear that, regardless of complexities of plasma-gas coupling and reaction chemistry, radially uniform deposition on the substrate is not possible without uniform mass transfer, which is governed solely by the fluid dynamics in the reactor. Therefore, we have studied the mass transfer characteristics of a similar flow system both experimentally and theoretically⁶⁶ to determine conditions in which radial uniformity can be expected in actual deposition experiments.

Experiment. A flow system identical to that of Fig. 6-49 was constructed and mass transfer was studied by sublimation of naphthalene substrate into an air stream at ambient temperatures, under flow and pressure conditions similar to those to be used in the actual PCVD reactor. The naphthalene substrate was formed by hot-pressing naphthalene crystals into disks. Overall sublimation rates and radial uniformity were quantified by weighting and profiling the surface before and after an experiment.

Comparison with Stagnation Flow Models. The experiments have shown that for relatively low flow rates, radially uniform sublimation of the naphthalene substrate occurs, and from this we conclude that uniform deposition of SiO_2 will occur under similar conditions in the PCVD coater, as long as radially uniform plasma conditions and reaction kinetics are realized. A

stagnation flow model⁶⁷ that assumes uniform injection of gas into the reactor (i.e. injection through a uniformly porous surface), and that neglects edge effects, has been applied to this data. Model predictions for mass transfer rates agree very well with measured values for low flow rates, as shown in Fig. 6-50. However, as the flow rate becomes larger, experimental mass loss rates begin to deviate strongly from predicted values. This condition is caused by locally enhanced mass transfer as a result of jet impingement on the substrate. Figure 6-51 shows a naphthalene disk after a high mass-loss experiment. The dimpling pattern on the surface reflects the spacing of the gas injection holes above: troughs are directly under the holes, and ridges are formed between adjacent holes. In the actual coater, deposition to the surface, and not removal from the surface, occurs and this pattern would be reversed with troughs between adjacent holes.

A Jet Impingement Mass Transfer Model.

Heat and mass transport enhancement due to jet impingement has been the subject of many studies in the past. However, these studies have considered high jet velocities in systems for which increased transport rates are desirable. In our case, mass transport enhancement due to jet impingement is definitely not desirable.

We have developed a model for jet impingement for low velocity viscous flows in the hopes of quantitatively determining operating conditions wherein the local mass transfer irregularities in Fig. 6-51 will not occur. In the interest of tractability, a single jet in an axisymmetric geometry is analyzed. The influence of adjacent jets is implicitly accounted for by imposing a zero radial flow boundary at a distance midway between the jets. The geometry of our computational domain is shown in Fig. 6-52. Also shown on this figure are the boundary conditions imposed on our flow. The upper confining surface present in the experimental geometry (the upper surface of the reactor through which the injection holes are drilled) has been removed in our computational geometry to allow gas to escape. This condition does not affect the flow patterns in the region of interest near the substrate provided that the aspect ratio (height/width) of our computational cell is large.

The air flow is described by the Navier-Stokes equations of motion for steady, laminar,

Figure 6-49. Schematic of PCVD Reactor.

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Figure 6-50. Experimentally determined naphthalene mass loss rates compared with predictions of stagnation flow model.⁶⁷ (Stagnation flow Reynolds Number $Re = Q_t H / A \mu$, where Q_t is the total gas flowrate, H is the injection plate-to-substrate spacing, A is the injector surface area and μ is the gas viscosity.

Figure 6-51. Effect of jet impingement-enhanced mass transfer on naphthalene sublimation.

Figure 6-52. Single jet impingement model flow geometry, showing boundary conditions for velocities and naphthalene mole fraction.

axisymmetric flow of an incompressible fluid, written in cylindrical (r, z) coordinates as:

$$v_r \frac{\partial v_r}{\partial r} + v_z \frac{\partial v_r}{\partial z} = \frac{\partial P}{\partial r} + \frac{1}{Re_j} \left[\frac{\partial}{\partial r} \left(\frac{1}{r} \frac{\partial}{\partial r} (r v_r) \right) + \frac{\partial^2 v_r}{\partial z^2} \right] \quad (1)$$

$$v_r \frac{\partial v_z}{\partial r} + v_z \frac{\partial v_z}{\partial z} = \frac{\partial P}{\partial z} + \frac{1}{Re_j} \left[\frac{1}{r} \frac{\partial}{\partial r} \left(\frac{\partial v_z}{\partial r} \right) + \frac{\partial^2 v_z}{\partial z^2} \right] \quad (2)$$

where v_r and v_z are the r and z velocity components and $Re_j = 4Q/\pi d_j \mu$ is the jet Reynolds number (where Q is the gas flow rate through each hole, d_j is the jet diameter and μ is the gas viscosity). The assumption of incompressibility is valid since the air velocities are much less than the speed of sound in air.

Mass transfer is described in terms of the gas flow by:

$$v_r \frac{\partial X}{\partial r} + v_z \frac{\partial X}{\partial z} = \frac{1}{Re_j Sc} \left[\frac{1}{r} \frac{\partial}{\partial r} \left(\frac{\partial X}{\partial r} \right) + \frac{\partial^2 X}{\partial z^2} \right] \quad (3)$$

where X is the mole fraction of naphthalene in the air stream normalized with its value at saturation, and $Sc = \mu/pD$ is the Schmidt number (where p is the gas density and D is the diffusivity of naphthalene in air).

Equations 1–3 along with boundary conditions shown in Fig. 6-52 were solved by using a commercially available finite-element flow code FIDAP.⁶⁸ The solution was obtained in two steps: first, the flow field was calculated; and second, Eq. 3 was solved for X given this flow field. From this solution, the mass transfer rate at the substrate surface $J_g = \partial X / \partial z|_{z=0}$ was evaluated at different radial positions. Parameters influencing this solution are Re_j , the jet-to-substrate spacing H and the jet-to-jet spacing L .

The model as well as the experiments show that mass transfer nonuniformities are very sensitive to experimental conditions in the range of interest to our application. Representative model results are shown in Fig. 6-53. This figure shows the marked nonlinearity in the radial mass transfer uniformity with respect to Re_j .

Model results suggest that the axial distance wherein the jet centerline velocity decays to small values is dependent only on Re_j and L . Increasing Re_j increases the momentum strength, while increasing L decreases the effectiveness of shear stresses in dissipating this momentum, and leads to recirculation zones that insulate the central core of the jet. As long as the jet-to-substrate spacing H is greater than the effective jet dissipation length determined by Re_j and L , jet impingement mass transfer irregularities are not appreciable.

A series of experiments were performed in which Re_j and H were varied, and model predictions for the onset of jet impingement transport nonuniformities agreed well with experimental results.⁶⁶ This modeling and experimental effort has allowed us to define upper operating flow rates for a given reactor geometry in the PCVD coater, and has led to a design modification that allows changing the diameter of the gas injection ports.

Author: J. A. Britten, Major Contributor: C. B. Thorsness

6.3.5. Effect of Composition on the Properties of Phosphate Laser Glasses

Introduction. One of the major problems glass chemists face in developing new laser glass is tailoring the composition to achieve the desired properties for a given laser design. The approach often used in the past has been for the laser designer to specify the desired thermal, mechanical, optical and laser properties and have the glass chemist carry out a "composition study" aimed at optimizing these properties. The general glass composition space to be studied is usually selected based on past experience. Although this approach is often successful, it generally requires a new compositional study for each new laser design.

In the past, a number of attempts have been made to correlate various laser glass properties with corresponding composition changes.⁶⁹⁻⁸⁴ In general, these studies show the influence of various additives on the glass properties. However, what has remained illusive is a means of quantitatively predicting the effects of composition on the properties of complex, multicomponent laser glasses. To develop such predictive correlations from first principles is a formidable task because the details of the influence of glass composition on structure is not

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Figure 6-53. Model predictions of jet-induced mass transfer rate variation from stagnation point to flow separation point for various Re_j , with $H = 2$ cm, $s = 1$ cm.

well understood. Nevertheless, a large number of laser glasses have been melted and their properties measured (see for example Refs. 72, 80-84). Consequently, a database exists from which empirically derived composition versus property correlations could be developed.

During this past year we have taken a first step in trying to develop predictive correlations from known compositional databases.

Specifically, we have completed a compositional study of derivatives of Schott LG-750 phosphate laser glass and have attempted to correlate the compositionally averaged electric field strength with some thermal, mechanical, and optical properties of these laser glasses.⁸⁵ This work suggests that good predictive tools (in a quantitative sense) could possibly be developed that relate glass composition to laser glass properties. Such tools could greatly reduce the amount of compositional research required to develop new glasses for future laser designs.

Electric Field Strength as a Correlation Variable. Most laser glasses used in laser systems are complex, multicomponent compositions. Thus, one is faced with trying to relate glass properties to glass compositions containing a multitude of different network formers or modifiers of varying concentrations. A number of previous studies have shown that glass properties can often be correlated with the so-called "electric field strength" of the modifier cations. This concept was developed many years ago by Dietzel⁸⁶ and is derived from the Coulombic attractive force, F (dynes), between oxygen anions and the corresponding cations in the glass:

$$F = \frac{(q_c q_o e^2)}{(r_o + r_c)^2} \quad (1)$$

where r_o and r_c are the oxygen and cation radii (cm), q_o and q_c are the valence states of the anion (O^{2-}) and cation respectively, and e is the fundamental charge. Values of the field strength calculated from published cation radii can be found in the literature.⁸⁷ The term $[q_c e / (r_o + r_c)^2]$ is defined by Dietzel as the "cation field strength" relative to O^{2-} and has units of esu/cm^2 .

For the purposes of this work, we have attempted to correlate a compositionally averaged cation field strength E_{eff} with various laser glass properties:

$$E_{\text{eff}} = \sum_{i=1}^n E_i C_i S_i \quad (2)$$

where C_i is the mole% of the i th oxide, E_i is the field strength for the corresponding cation, and S_i is the cation/oxygen stoichiometric ratio (e.g., 2 for Na_2O). In this way, we can compare glasses having different or mixed modifier compositions. The field strength parameter has the advantage of being easily determined from the glass composition and does not require costly analytical measurements.

Unfortunately, this simple field strength picture is inadequate for describing variations in network formers on glass properties. This is understandable; the field strength model is based on an electrostatic picture and is expected to work well only for additives that have widely different electronegativities (i.e., high degree ionic character in the bonds). The network formers are largely covalently bonded and thus would not be expected to conform to the field strength model. To help illustrate this, the percent ionic character for various cation-oxygen bonds of several common glass formers or modifiers is plotted in Fig. 6-54. These were calculated from the electronegativity differences, using Pauling's suggested relationship⁸⁸:

$$\% \text{ ionic bond} = \left(1 - \exp \left[-(x_c - x_o)^2 / 4 \right] \right) \quad (3)$$

where $(x_c - x_o)$ is the difference in electronegativity between the cation and oxygen. Note that, as expected, the strong oxide-glass network formers (i.e., Si, B, P, etc.) have less than about 50% ionic character, whereas most strong modifiers generally have greater than 70% ionic character in the bond. In this study, we have held the network former components constant, therefore, only the variation of E_{eff} produced by the modifiers is considered. The compositionally averaged field strength values are expressed in units of $10^{-14} \text{ mol\% esu}/\text{cm}^2$, as per Eq. (2).

Figure 6-54. Estimated percent ionic character in a given chemical bond based on Pauling's electronegativities differences.⁸⁸ Typical oxide glass network formers and strong modifiers fall in the regions shown.

Glass Preparation and Property Evaluation.

The results presented here are based on analyses of 41 small-scale melts which are compositional derivatives of Schott's commercially available LG-750. The modifier compositions and measured properties of these glasses are summarized elsewhere.⁸⁵ Each composition was melted in a quartz crucible at temperatures near 1300°C, followed by refining and homogenizing for a few hours. The glass was cast, annealed and then cut into sections to allow fabrication of samples required for the evaluation of the properties of interest. Further details of the compositions, melting procedures and methods used to characterize the sample are given in Ref. 85.

The base phosphate glass used in this study (LG-750) is a proprietary composition so we report⁸⁵ only the alkali and alkaline earth components in the glass; in all compositions, they comprised a total of 31 mol%. Each glass also contained 1.6 mol% Nd₂O₃ (i.e., about 4 wt%). To make the data easier to assimilate, we have arranged the compositions into seven composition groups, as summarized in Table 6-7.

The effective field strength for the mixture was calculated from the composition and the individual cation field strengths using Eq. (2). In the following sections, we summarize some of the effects of various cation additions on measured glass properties. Further correlations and data are given in Ref. 85.

Thermal and Mechanical Properties. In general, the coefficient of linear thermal expansion (CTE), thermal conductivity, and Young's modulus all correlate with the effective field strength for the multicomponent alkali and alkaline earth cation mixtures.

Intuitively, one might expect the CTE to decrease with increasing cation field strength, as is observed in the data (Fig. 6-55a). This is based on the simple physical picture of the glass as an ensemble of anharmonic oscillators. For higher field strength cations, the binding energy between the cation and oxide is greater corresponding to a deeper interatomic potential well. Thus, increases in temperature would tend to produce less of a departure from harmonic oscillation around some mean interatomic distance. In contrast, for low field strength cations the oscillators lie in a shallower energy well and have greater anharmonic character. Therefore, the lower the field strength of the cation, the greater the increase in mean internuclear separation with temperature.

The thermal conductivity is observed to increase with increasing field strength for the glasses in this study (Fig. 6-55b). This good correlation is probably not an effect of field strength, as much as it is simply the impact of various cation additives on phonon mean-free path. It is generally well known that the greater the atomic weight, the greater the phonon scattering and the lower the thermal conductivity. Since, within a given group, the field strength is greatest for the lowest atomic weight cation (e.g., Li⁺ > Na⁺ > K⁺, etc.), then one might expect that the thermal conductivity would increase with increasing field strength.

We also found that in general, Young's modulus correlates with field strength although there is significant scatter in the data.⁸⁵ One of the reasons for the large spread in the data is that the alkaline earths have less effect on the modulus than the alkali cations (for comparable changes in field strength).

For some laser applications, the thermal shock resistance of the glass is important. In such cases, glass with a maximum value for the parameter $k(1-\mu)/E\alpha$, often referred to as a thermal-mechanical figure of merit, are favored. Here, k is the thermal conductivity, μ is Poisson's ratio, E is Young's modulus, and α is the thermal expansion. We found that of the glasses in this study, the ones with the highest compositionally averaged field strength have the greatest thermal-mechanical figure of merit.

Optical and Laser Properties. A plot of the refractive index (n_d) versus field strength is shown in Fig. 6-56a for the first five groups of glasses shown in Table 6-7. Note that proceeding from K to Li increases n_d , whereas Ba to Mg causes n_d to decrease. Also values of n_d for glasses containing mixtures of various alkali components fall smoothly on these curves. The refractive index (n) of a dielectric material in the visible and near visible region is directly related to its electronic polarizability by the well-known Lorentz-Lorenz relationship:

$$n = \left(\frac{1 + 2k\alpha_e/V_m}{1 - k\alpha_e/V_m} \right)^{\frac{1}{2}} \quad (4)$$

where α_e (cm³) is the electronic component of the polarizability, V_m (cm³/mol) is the molar volume of the glass and k is a constant equal to $4\pi N/3$, where N is Avogadro's number (mol⁻¹).

Table 6-7. Summary of glass composition groupings.

Composition "Group"	Melt Numbers	Key Composition Feature
1	A-6, 1, 4, 7, 14, 16, 9	Constant BaO (13.8%) and variable alkali oxide
2	A-2, 8, 11, 13, 15, 17	Constant MgO (13.8%) and variable oxide
3	A-22, 18, 19	Constant CaO (13.8%) and variable alkali
4	A-20, 30, 31, 27, 38, 40, 41, 44	Constant MgO (18.8%) and variable alkali oxide
5	A-24, 33	Constant MgO (22.8%) and Li ₂ O vs. K ₂ O
6	A-3, 5, 37, 39, 28, 29, 51, 34, 36, 32, 47, 50, 52	Mixed alkali metal and alkaline earth
7	A-21, 23, 8	Constant Li ₂ O and MgO; variable ZnO

Figure 6-55. Correlation of (a) coefficient of linear thermal expansion and (b) thermal conductivity with compositionally averaged field strength defined by Eq. (2) for all 41 glasses in this study. The solid line represents a linear regression to the data.

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Figure 6-56. (a) Refractive index (n_d) and (b) molar volume versus compositionally averaged field strength for glasses containing MgO, CaO, or BaO with both single and mixed alkali metal components.

Note that the index increases as the term α_e/Vm increases. Since the polarizability varies as the size of the ion, then, within a given group of the periodic table (e.g., Group IA or IIA) the index should increase with increasing atomic number. This trend is observed for the alkaline earths (with fixed alkali metal - see Fig. 6-57). However, within the alkali metal series (with fixed alkaline earth) just the opposite trend is observed, i.e., the index increases going from K to Li. Thus, as shown in Fig. 6-56b, the molar volumes of these particular glasses increase more rapidly than the polarizabilities going from Li to K with a fixed alkaline earth component.

Emission Cross-Section. A plot of emission cross-section versus field strength is given in Fig. 6-58. Note that for this entire range of compositions, the cross-section changes by less than 10%.

From the Judd-Ofelt treatment^{89,90} of rare earth spectra, the emission cross-section is related to the line of strength for a particular transition between two J states.

$$\sigma(\lambda_p) = k_s \left(\frac{\lambda_p}{\Delta \lambda_{eff}} \right) \left(\frac{(n^2 + 2)^2}{n} \right) S(J:J') \quad (5)$$

where $S(J:J')$ is the line strength, λ_p is the peak fluorescence wavelength, $\Delta \lambda_{eff}$ is the effective line width given by

$$\Delta \lambda_{eff} = \int \frac{I(\lambda)d\lambda}{I(\lambda_p)} \quad (6)$$

and n is the refractive index at the emission wavelength. The proportionality constant, k_s , in Eq. (5) is given by $[8 \pi^3 e^2 / 27hc (2J + 1)]$. Thus, the cross section reflects not only the effect of the glass composition on line strength but also the effects of composition on refractive index and effective line width.

Our results also show that for the entire range of modifier compositions used in this study, there is little change in line strength of the $^4F_{3/2} - ^4I_{11/2}$ transition. Changes in transition strengths reflect the changes in the local field around the neodymium ion. Specifically, increases in the noncentrosymmetric character of the local field cause a corresponding increase in transition strength. Thus, the data show that relatively modest variations in modifier

composition for these particular glasses seem to produce little change in the local Nd^{3+} environment.

A correlation of the emission band width ($\Delta \lambda_{eff}$) of the glasses in their study with the modifier field strength showed that doubling the field strength of the cation additives produces only about a 6% increase in bandwidth.

The small observed changes (<10%) in cross section, line strength and bandwidth with large variations in the modifier content leads to an important practical conclusion for the glass chemist. That is, for this base glass composition (LG-750), one can vary the alkali and alkaline modifier composition to optimize the thermal, mechanical, or glass-forming properties without adversely affecting these laser properties.

Fluorescence Lifetime. The fluorescence lifetimes of the glasses in this study are plotted versus the compositionally averaged electric field strength in Fig. 6-57. Fig. 6-57a is the data for all glasses, whereas Fig. 6-57b shows the individual effects of various composition mixtures of one alkali metal plus one alkaline earth oxide. In general, the greater the modifier field strength the lifetime will be shorter.

The radiative lifetime τ_r represents the maximum possible fluorescence lifetime and is calculated from the sum of spontaneous radiative emission probabilities for all transitions from the $^4F_{3/2}$ states to all terminal states. In contrast, the measured fluorescence lifetime represents the sum of all radiative and non-radiative processes. There are four major non-radiative processes that affect lifetime:⁷¹ (a) self-quenching by Nd (so-called "concentration quenching"), (b) quenching caused by impurity ions, particularly transition metals, (c) energy transfer from Nd to excite vibrational overtones of hydroxyl groups in the glass, and (d) multiphonon relaxation between Nd and the glass matrix.

For the glasses used in this study (see Fig. 6-58), we have held the Nd concentration constant, thus we tentatively assume that the concentration quenching effect remains approximately constant of this set of glasses. In addition, high purity starting materials have been used to eliminate quenching induced by non-OH impurities. The water content does vary somewhat from glass to glass. However, we have corrected for the water effect on the lifetime, using the correlation developed by Caird.⁹¹ Thus, the fluorescence lifetime data plotted in

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Figure 6-57. Measured fluorescence lifetime versus compositionally averaged field strength for (a) all glasses in this study, and (b) glasses containing only one alkali metal and alkaline earth component. The solid line in (a) represents a linear regression analysis of the data.

QT-J-U-9173-06

Figure 6-58. Emission cross section for the $^4F_{3/2} - ^4I_{11/2}$ transition versus compositionally averaged field strength for all glasses in this study.

Fig. 6-57 represents the combined effects of the glass matrix on both the radiative and multiphonon relaxation.

Author: John H. Campbell

Contributors: Joseph S. Hayden, Yuiko T. Hayden (Schott Glass Technologies).

6.3.6. High-Laser-Damage-Threshold Potassium Dihydrogen Phosphate Crystals

Introduction. The progressive development of inertial confinement technology will eventually require drivers of even higher fluence. One such driver has been in conceptual design at LLNL for a number of years. Current design parameters⁹² are for 10-ns pulses with fluences of 40 J/cm² at 1 ω and 25 J/cm² at 3 ω . Because these fluences are considerably higher than Nova, we initiated research into understanding and improving the damage resistance of KDP single crystals. We describe here recent success we have had in growing KDP single crystals with damage resistance at and above our requirements.

Experiment. Purification. Analyses have shown that reagent-grade KDP contains a variety of contaminant compounds.⁹³ The contamination is in the form of dissolved atomic species and a variety of organic and inorganic compounds, both particulate and dissolved. Because of the varied types of contamination, a multifaceted approach is necessary for the eventual ultrapurification of raw starting material.

As our operational purification process, we chose an initial recrystallization followed by ozonation to remove oxidizable material, ultrafiltration to remove nonoxidized particles, and UV light exposure to suppress the growth of bacteria. This purification scheme was eventually realized in the apparatus shown in Fig. 6-59. It consists of a 150-L stainless steel tank with a heating/cooling jacket for processing the raw starting material. Attached to the tank by means of heated tubing is a 10-gal/min magnetically coupled centrifugal circulation pump (March Manufacturing, Glenview, IL), a standard 10-in. column filter housing (Nuclepore sanitary housing, Nuclepore Corp., Pleasanton, CA) and a UV exposure lamp (UV water sterilizer, Aquafine, Valencia, CA). This provides the means to filter and UV illuminate the KDP continuously until the desired purity is achieved. The lid of the tank is removable for

ease of loading; however, the tank remains sealed during processing and the finished raw material is passed through an outlet valve and into the growth chamber by modest pressurization with ultrafiltered nitrogen gas. The entire apparatus is portable to facilitate transport from the preparation laboratory to the growth facility.

In typical operation, the raw starting material, nominally 36 kg of KDP, is dissolved at 50°C in 150 l of steam-distilled and standard-resin-exchange treated water. The solution is then exposed to 0.3 l/min of 3% ozone in air generated by a Polymetrics Model T-816 ozone generator. The oxidizable carbon content of the solution is periodically monitored by analyzing extracted samples with an OI Analytical Model 700 TOC Analyzer. Ozonation is discontinued when the level of oxidizable carbon falls below the detection limit (~5 ppb). The solution is then degassed and transferred to the growth system with ultrafiltered dry nitrogen.

Growth. Most single-crystal KDP is grown in processes that either rotate or revolve one or more seed crystals in a saturated KDP solution. The process usually involves a load-bearing surface within the growth environment, and can generate a considerable amount of particulate contamination and ultimately spoil a previously purified growth solution. To mitigate this effect, we developed the continuous filtration, continuous flow growth system shown in Fig. 6-60. The growth chamber consists of a 5-l spherical, double-wall reaction vessel. The vessel is open at the bottom where a 60-cm Graham condenser is attached. The crystal seed is positioned rigidly in the center of the vessel. Fresh growth solution is pumped through the condenser at 1 l/min with a magnetically coupled gear pump (Micropump, Cole-Parmer, Chicago, IL). The flow is directed onto the growing crystal faces and then returned to the pump from the top of the reaction vessel by means of a heated tube. To maintain purity throughout the growing period, a 10-in. cylindrical filter with 0.05- μ m pore size is placed between the pump and the growth vessel.

The degree of saturation in the growth system is maintained by means of a Lauda model RK-20 refrigerator/circulator. The reaction vessel and condenser are heated independently with circulator water at a flow rate of ~3.5 gal/min (the returning solution is heated to ~3°C above saturation to avoid crystallization in the

MX-018-U-9329/F 1

Figure 6-59. Schematic of the bulk KDP purification apparatus.

MX-018-U-9329/F 2

Figure 6-60. Schematic of the continuous flow and filtration crystal growth apparatus.

pump and filter). All crystals grown to date with this system have been produced by the temperature-drop method. Typical growth rates were from 1 to 2 mm per day, and the crystals were grown to $5 \times 5 \times 5$ cm in size.

Damage Testing. Because the test methods and facilities have been described previously,⁹⁴ we will only briefly describe them here.

Most of the data were obtained from experiments performed on the LLNL Reptile facility. Reptile uses three lasers to generate pulses at the four YAG harmonics (1.06, 0.55, 0.35, and 0.26 μm). Nominal operating conditions are 10 or 16-ns pulse duration, 1 to 30-Hz repetition rate, and peak fluence of up to 65 J/cm² at 1064 nm. Samples to be tested are mounted in a Nomarski microscope stage. The illuminated volume, nominally a cylinder 1 mm in diameter and 8 mm in length, is examined for defects prior to testing by an operator using back-lighted microscopy at 100 $\times$ magnification. After testing, the crystal volume is again observed at 100 $\times$ magnification and new (damage) defects are counted and characterized. The process is performed either by illuminating the test volume with a single pulse at preset fluences (1-on-1) or by illuminating the test volume with a series of pulses of increasing fluences beginning well below the damage-threshold (laser-conditioned or n-on-1).

Results and Discussion. Table 6-8 is a compilation of recent progress made at LLNL, Cleveland Crystals, Inc. (CCI), and the Fujian Institute in increasing the laser damage threshold of single-crystal KDP. The pre-1987 data point is the average of a selection of damage thresholds from an LLNL database that contains damage thresholds of over 600 individual KDP single crystals. The selection criteria for the pre-1987 data was a minimum of purification to include no filtering of the growth solution or minimal filtering at pore sizes greater than 2 μm .

The remaining data in the table have been obtained on crystals grown since 1987. These crystals were all grown in solutions of increasing levels of purification. The purification techniques used were ultrafiltration, filtration of the growth solution at pore sizes of less than 0.1 μm prior to growing, continuous filtration, filtration of the growth solution during growing at pore sizes equal to or greater than 0.1 μm , and continuous ultrafiltration for continuous filtration of the growth solution at pore sizes

less than 0.1 μm . Also, crystals grown at LLNL were by the continuous flow and continuous ultrafiltration method described in this paper.

Due to the time-consuming nature of the damage testing we performed on these crystals, it was not practical to measure all samples at all wavelengths. However, the majority of crystals were tested with laser-conditioning (n:1) at 1064 nm. Although the sample database is small, it is still possible to draw some significant conclusions.

The samples from LLNL were obtained from crystals grown under identical methods with the exception of filter pore size. Here, there was a decrease in pore size of approximately two orders of magnitude among the samples. This decrease in pore size resulted in a marked difference in damage threshold. Indeed, the sample grown at 0.05- μm pore size exceeded the fluence limits of the testing laser. This is fairly strong evidence in corroboration of a foreign-particle-derived damage mechanism.

Two of the samples from CCI were obtained from a crystal grown in a solution that was continuously filtered at 0.05 μm . However, the damage threshold of these crystals was slightly lower than the crystal grown in the LLNL continuous ultrafiltration and flow apparatus. This may be attributed to a difference in growth apparatus. The CCI method used a bearing surface within the growth chamber that can be a source of foreign particulate matter. This can in turn reduce the effectiveness of the filtering scheme thereby resulting in the potential incorporation of contamination into the growing crystal.

The Fujian crystals had damage thresholds lower than the CCI ultrafiltered (0.05- μm pore size) samples. Although we have limited information on the Fujian growth process, we do have some assurance that the crystals are grown from a solution that has been ultrafiltered prior to growth and that the growth mechanism does not incorporate continuous filtration.⁹⁵ This method exposes the crystal to the possibility of increased particulate contamination during growth and may explain the lower damage thresholds.

The above results, although from a limited number of samples, give considerable credence to the theory that particulate contamination in crystal growth solutions can be responsible for low damage thresholds in KDP single crystals. Furthermore, mitigation of the problem requires

Table 6-8. Comparison of the damage thresholds at 355 nm and 1064 nm and 10 ns for various KDP single-crystal samples.

Source	Filtration Method	Pore size	Damage threshold 10 ns (J/cm ²)			
			355 nm		1064 nm	
			1:1	n:1	1:1	n:1
Various	See text	-	-	-	-	16.0
Fujian	Ultrafiltration	-	-	-	10.8	22.2
Fujian	Ultrafiltration	-	-	-	-	>30.9
LLNL	Continuous flow&filtration	>0.2μm	-	-	11.3	18.1
LLNL	Continuous flow&filtration	0.2μm		22.3	-	37.4
CCI	Continuous filtration	0.2μm	-	-	12.3	31.0
CCI	Continuous filtration	0.2μm	-	26.2	17.4	43.4
CCI	Continuous filtration	0.2μm	-	-	20.0	37.6
CCI	Continuous filtration	0.2μm	-	-	10.7	36.3
CCI	Continuous ultrafiltration	0.05μm	17.0	25.7	44.4	-
CCI	Continuous ultrafiltration	0.05μm	17.8	24.2	32.1	-
LLNL	Continuous flow&ultrafiltration	0.05μm	-	28.6	-	>63.9

both ultrapurification of growth solutions as well as maintenance of purity during crystal growth.

We have demonstrated a correlation between growth solution purity and laser damage threshold of single crystal KDP. Controlled experiments implicate particulate contamination as a main source of laser damage. Furthermore, the highest damage thresholds are achieved when particulate contamination is actively controlled for the duration of the crystal growth. These crystal growth conditions have resulted in test samples that exceed the damage threshold requirement of the next-generation LLNL laser fusion driver. Research in this area is now being directed toward developing a practical application of the purification and growth techniques described herein for the production of large-scale KDP crystals.

Authors: Ken E. Montgomery and Fred P. Milanovich

6.4. Advanced Materials for Frequency Conversion

6.4.1. Introduction

Nonlinear crystals that harmonically convert infrared light to the visible and ultraviolet are essential components of high-power glass laser systems for inertial confinement fusion. On NOVA, the present ICF laser, potassium dihydrogen phosphate performs this function. Since the short wavelength light couples most effectively to the target, the efficiency of the frequency conversion process is a key criterion for judging the performance of these crystals. With uniform beams having optimal temporal and spectral characteristics KDP is a very good conversion crystal. However, large laser systems generally do not provide such "perfect" beams. Therefore we have continued our search for harmonic generating crystals that could be capable of much higher conversion efficiency than KDP for the shaped pulses, sub-optimal beam quality and broad spectral bandwidth expected in the next generation of fusion lasers. Of course, efficiency is not the only issue that determines the relative merit of one nonlinear crystal compared to another for an ICF driver. Processes that limit

the drive intensity (especially optical damage), and the relative cost of crystal growth and fabrication are also important. Table 6-9 lists the various properties that can influence the choice of a material for the frequency converter on a multimegajoule-scale laser facility. Ultimately the choice of material must be based on the cost/performance tradeoff.

We have made significant progress in our search for harmonic generating materials with enhanced efficiencies. We are now focusing on more accurate characterization of a small group of candidate materials. We also continued our investigation of the properties of deuterated L-arginine phosphate (DLAP). DLAP has remained of interest both as an alternative to KDP, and as a representative member of a large class of organic salts we are investigating. For high average power conversion, DLAP is attractive because it has a set of temperature insensitive phase matching directions which we report. For more immediate uses, we continued to investigate improved growth processes for KDP itself, which remains the only harmonic generating crystal available in the size and quality required for large fusion drivers. Improved growth methods that raise the damage threshold of KDP to our goal of 40 J/cm² were demonstrated.

Author: S. Velsko

6.4.2. New Materials Evaluation

A. Crystal Synthesis and Powder Survey.

We have completed the synthesis and preliminary screening of over 200 new crystals of chiral organic salts and related materials. A chiral organic salt is composed of (small) organic molecules that contain at least one small acentric unit with delocalized bonding attached to a chiral carbon center.⁹⁶ The chiral carbon ensures that the crystal containing the molecule will be noncentrosymmetric, while the acentric delocalized unit gives the molecule a large second order hyperpolarizability. This group of materials was chosen for study because it contains a large fraction of crystals which are more efficient harmonic generators than KDP for doubling and tripling near infrared light.⁹⁶ Since all of these crystals are water soluble, in principle it is possible to grow them to the large sizes needed for ICF using the same general method as is currently employed for KDP.

Table 6-9. Properties that determine the relative value of ICF harmonic generating crystals.

Intrinsic efficiency	Intensity limits	Scaling + fabrication	Other
Nonlinearity	Damage	Material cost	Thermal sensitivity
Angular sensitivity	Stimulated brillouin	Diamond turning	Moisture sensitivity
Chromatic sensitivity	Stimulated Raman	Handling % Waste	
Absorption	Two-Photon absorption	A.R coating	

The largest group of compounds we studied were the salts of optically active amino acids and their derivatives because these molecules are readily available commercially. Table 6-10 gives the amino acids and counterions from which we attempted to make salts. Of the approximately 400 possible compounds, 180 new materials with a characterizable morphology were obtained. We also synthesized a large group of nitrate salts,⁹⁷ and made a literature survey of available formate salts.⁹⁸ Figure 6-61 shows the range of nitrate salts containing one or two different cations which were investigated. A significant number of new noncentrosymmetric crystals containing nitrate ions were discovered in this work.

Of the compounds we produced, 22 had powder second harmonic generation (SHG) signals comparable to or larger than KDP.⁹⁹ This implies that the nonlinear coefficients of these 22 crystals are similar to or larger than those of KDP. The white-light refractive indices of these 22 compounds were subsequently measured so that their potential for angularly insensitive phase-matched doubling or tripling could be evaluated.¹⁰⁰ Low angular sensitivity, in addition to a large nonlinear coupling, is crucial for high-efficiency frequency conversion. Table 6-11 is a list of the 22 candidate materials that indicate the particular frequency conversion process (i.e., Type I or II doubling or tripling) for which they have low angular sensitivities. The phasematching properties of these materials are now being investigated in more detail using the single crystal measurement techniques previously developed in our laboratory. From these measurements, we will determine if the phasematching orientations with low angular sensitivity are also the ones that have the required large nonlinear coupling.

B. Direct Phasematching Measurements in Small Crystals. We have continued to characterize the phasematching properties of small single crystals of candidate harmonic generators using the NLOG technique.¹⁰¹ Here, small crystals are ground into a spherical shape and mounted on an apparatus that monitors the intensity of second or third harmonic light as the orientation of the sample is changed relative to the laser beam focused into the crystal. During the measurement, the crystal is immersed in a refractive index matching fluid to reduce light scattering from the sample surface. From these

measurements, the phasematching orientations and their associated nonlinear coefficients and angular sensitivities can be obtained.

The accuracy of the nonlinear coefficients determined by this technique is a major concern to us, since the intensity of the converted signal can be reduced by light scattering at the crystal-fluid interface due to the imperfect match between the refractive index of the crystal and the fluid in which it is immersed.¹⁰¹ For the mixing process $\omega_1 + \omega_2 = \omega_3$, the effective nonlinear coefficient d_{eff} is related to the measured intensity of the output beam I_3 by:

$$I_3 \propto \frac{d_{\text{eff}}^2}{n_1 n_2 n_3} l^2 I_1 I_2$$

where I_1 and I_2 are the intensities of the input laser beams with frequencies ω_1 and ω_2 ; n_1 , n_2 and n_3 are the crystal's refractive indices at ω_1 , ω_2 and ω_3 respectively; and l is the crystal thickness.

We have evaluated the accuracy of our measurements by comparing the d_{eff} values determined in spheres with those measured in polished plates for materials available in sizes sufficient to cut properly oriented plates.¹⁰² The results of these studies are shown in Table 6-12. Agreement between the sphere and plate measurements is excellent for low and moderate refractive index crystals. For crystals whose refractive indices are larger than 1.6, such as barium borate (BBO) and lithium borate (LBO), the measurements are poorer since immersion fluids with refractive indices greater than 1.6 have poor optical transmission in the infrared and ultraviolet. In those cases we are forced to use lower index fluids and accept the consequences of a poor match. The scattering problem is most severe for the 3ω measurements. Fortunately, most of the materials we have characterized have refractive indices less than 1.6, for which adequate fluids are available.

Table 6-13 lists the crystals whose phasematching properties were studied this year. This list includes both crystals which were discovered in the survey of the previous year and crystals which were reported in the literature.¹⁰³ The literature materials were chosen because the reported properties implied that they would be high efficiency doublers or triplers. Of the literature materials, only

Table 6-10. Amino acids and counterions used to produce chiral organic salts.

Amino Acids	Anions	Cations
Alanine	Fluoride	Lithium
Arginine	Chloride	Sodium
Asparagine	Bromide	Potassium
Aspartic acid	Nitrate	Rubidium
Cysteine	Sulfate	Cesium
Cystine	Chlorate	Ammonium
Glutamic acid	Phosphate	Thallium
Glutamine	Arsenate	Guanadinium
Histadine	Acetate	
Hydroxyproline	Trifluoro-	
Isoleucine	Acetate	
Leucine	Trichloro-	
Lysine	Acetate	
Ornithine	Methyl-	
Methionine	sulfonate	
Proline	Trifluoromethyl-	
Serine	sulfonate	
Threonine	Chloroacetate	
Valine	Formate	
	Lactate	
	Malonate	
	Oxalate	

	H	Li	Na	K	NH	Tl	Rb	Cs	Be	Mg	Zn	Cd	Ca	Sr	Pb	Ba	Al	Ga	Yb	Ir	Y	Tb	Gd	In	Sm	Nd	Pr	Ce	La	Ur
H	A		X	A	C	U	X	X	X	X	X	X	X	X	X		X		X		X		X	X				X		
Li	C	X	X	X	X	X	X	X	X	U		X	X	X	X					X		X						A	X	C
Na	C	A	X	I	A	X			X	X		X	X	X	X	X	X		C		C		C	A		C	C	C	C	X
K	C	A	I	X	A		A	X	X	A	A	X	A	A		A	A		U	A	A	A	C		C	A	A	A	A	X
NH	X	X	A	U		U	X	U	A	A	X	X	X	X		X		A		A		C	C					A	A	X
Tl	C		X			U			X	A	I	I	X				X				C							A	A	
Rb	A	S	X			U	U						X	X			X				X			X		U	U	A		
Cs	A	X			U	X	X	X	X	X			X	A			X				A		X	A				C	A	X
Be	X		X		X								X			X														
Mg	C				X								X			X	U											U	U	C
Zn	A	I													X	X												U	U	
Cd	A														X													X		C
Ca	C	S	S																									U		A
Sr	C		S																		X							X		C
Pb	C	S	X																		X							X	X	C
Ba	C	X																			X							X	X	C
Al	C																											X	X	
Ga	C																											U		
Yb	A																													
Ir	C																													
Y	C	U	S																									X	X	
Tb	C																													
Gd	C																													
In	A																													
Sm	C																													
Nd	C																													
Pr	C																													
Ce	C																													
La	C																													
Ur	A																													

Figure 6-61. Chart showing the cation combinations used in the synthesis of nitrate salts for nonlinear optics. A - acentric phase formed (bold indicates large powder SHG signal); C - centric ; X - mixed salt does not exist at room temperature; I - indeterminate phase; blank - no crystalline phase produced.

Table 6-11. New ICF frequency conversion candidates.

Compound	Doubling*		Tripling		
	I	II	I	II	III**
N-Acetyl methionine	$\gamma(1.00)$	$\alpha(1.15)$	β	---	---
Cd lactate	---	$\omega(1.03)$	ω	---	---
histidine CF ₃ COOH	---	$\alpha(1.14)$	---	γ	γ
histidine acetate	---	---	---	---	4
histidine fluoroborate	---	$\alpha(1.01)$	α	γ	γ
histidine CH ₃ SO ₃ H	---	$\alpha(1.13)$	---	γ	γ
histidine nitrate	---	---	---	---	4
histidine perchlorate	---	$\alpha(1.12)$	---	---	4
alanine CH ₃ SO ₃ H	---	---	γ	---	---
l-arginine oxalate	---	---	---	β	β
l-arginine CF ₃ COOH	---	---	---	α	4
TI acid glutamate	---	---	β, γ	---	---
methionine nitrate	---	---	α	---	4
methionine sulfone	---	$\gamma(1.12)$	---	---	β
l-valine bromide	---	---	---	---	---
leucine nitrate	$\alpha(1.08)$	---	---	γ	γ
leucine chloride	$\alpha(1.00)$	$\gamma(0.99)$	γ	---	---
CaNO ₃ urea		$\gamma(1.06)$	γ	---	β
KN ₃ + LaNO ₃ ·6H ₂ O	---	---	---	β	β
K ₂ Ba (NO ₃) ₄	---	$\gamma(1.06)$	---	---	---
RB ₂ NA (NO ₃) ₃	---	---	---	α	α

* Noncritical wavelength reported in microns.

** "4" indicates the type III is phasematchable for 1.06 μ m, but noncritical.

Table 6-12. Comparison of nonlinear coefficients deduced from sphere and plate measurements.

Material/ process		d_{eff} (pm/V)	
		Sphere	Plate
KDP	2 ω II	0.33 ± 0.01	(0.35)
	3 ω II	0.36 ± 0.01	(0.35)
LAP	2 ω II	0.70 ± 0.14	0.74 ± 0.10
	3 ω II	0.80 ± 0.16	0.73 ± 0.07
LAAc	2 ω I	0.28 ± 0.07	0.26 ± 0.02
	3 ω I	0.68 ± 0.14	0.67 ± 0.06
BBO	2 ω I	1.09 ± 0.15	2.02 ± 0.21
	3 ω I	0.22 ± 0.04	1.55 ± 0.12
LBO	2 ω I	0.83 ± 0.17	0.82 ± 0.05
	2 ω II	0.70 ± 0.14	0.68 ± 0.05
	3 ω I	0.38 ± 0.08	0.67 ± 0.04
	3 ω II	0.32 ± 0.08	0.55 ± 0.03

Table 6-13. Materials studied by the NLOG technique.

New Materials:	Process	
L-Arginine Acetate	(LAAc)	2ωI,2ωII,3ωI
L-Arginine Bromide	(LABr)	3I,3ωII
N-Acetyl Histadine	(NAHis)	3ωI,3ωII
N-Acetyl Tyrosine	(NATyr)	3ωI,3ωII
N-Acetyl Hydroxy- proline	(NAHP)	3ωI,3ωII
Thallium Tartrate	(DTT)	2ωI,2ωII
Literature Materials:		
Lithium Borate	(LBO)	2ωI,2ωII,3ωI,3ωII
Lithium Formate	(LFM)	2ωI,2ωII,3ωI,3ωII
Dipotassium Malate	(DKM)	3ωI,3ωII
Dipotassium tartrate	(DKT)	3ωI,3ωII
Diammonium tartrate	(DAT)	3ωI,3ωII
Strontium Formate dihydrate	(SFD)	2ωI,2ωII,3ωI,3ωII
Yttrium Formate dihydrate	(YFD)	2ωI,2ωII,3ωI,3ωII

lithium borate¹⁰⁴ and lithium formate (LFM)⁹⁸ were found to be more efficient than KDP for ICF doubling or tripling. Of the new crystals, only L-arginine acetate (LAAc)^{102,105} was found to be more efficient than KDP (as a Type I tripler).

C Alternative Candidates for ICF Frequency Conversion. We have now identified five materials as possible high efficiency alternatives to KDP for frequency converting large aperture, solid state fusion lasers.

Table 6-14 lists these materials along with their nonlinear coefficients, and angular sensitivity values relative to KDP for selected frequency conversion processes.

L-arginine fluoride is a more efficient Type I doubler,¹⁰⁶ and L-arginine acetate is a more efficient Type I tripler^{102,105} than DLAP or KDP. Preliminary investigations of the crystal growth of these materials shows that LAAc grows as easily as LAP and has a similar crystal shape.^{106,107} On the other hand, L-arginine fluoride grows more slowly than LAAc and has a more plate-like habit.¹⁰⁶ Both of these crystals have much higher melting points than LAP (230 °C vs. 120 °C) because the crystal structures do not contain included water molecules (waters of hydration). As a result, they are likely to behave much like KDP with respect to the sol-gel coating process.

LBO and LFM are commercially available crystals and are currently available in sizes of a few centimeters. Unlike the other crystals in this list, LBO is a high temperature, inorganic crystal grown from a flux.¹⁰⁴ We are currently evaluating in more detail the crystal growth, laser damage and absorption properties of these crystals.

Author: S. Velsko

Major Contributors: L. Davis, M. Webb, F. Wang, R. Reibold.

6.4.3. Development of Deuterated L-Arginine Phosphate (DLAP)

We have studied the crystal growth of DLAP (under contract to Cleveland Crystals, Inc.) Ideally, DLAP would grow with equal rates in the a and c crystallographic directions to maximize the yield of properly oriented frequency conversion plates from a single DLAP crystal. In general, the shape of DLAP crystals are plate-like, with the a dimension much

shorter than the b or c dimensions.¹⁰⁸ It is known from previous studies that DLAP grown from solutions with excess arginine do have more equal a b and c dimensions than those grown from stoichiometric solutions or with an excess of phosphoric acid.¹⁰⁸ However, crystals grown with excess arginine are much more prone to cracking, as well. Therefore, one aim of our investigation is to find other impurities that will inhibit the growth of the c and b crystal faces with respect to that of the a face when they are added to the growth solution in low concentrations. In an initial study, small amounts of transition metal ions such as chromium and copper were added to the growth solution since it is well known that these have a dramatic effect on the morphology of KDP crystals.¹⁰⁹ However, no such effect was observed in DLAP, even at impurity concentrations as high as 1%. In a second set of experiments, DLAP crystals were grown from solutions that contained 1–2% of other common amino acids besides L-arginine (See **Table 6-11**, Sec. 6.4.2.). Only threonine had an observable inhibiting effect on the b and c growth directions. In no case was the effect of impurities on crystal shape as large as that caused by varying the amount of arginine.

We also performed several moderate scale (10 cm) DLAP growth runs, but only one yielded a large volume of clear material. Other runs were terminated prematurely because of cracking or other problems. Attempts to grow DLAP at constant temperature were consistently plagued by the formation of large numbers of small, unwanted crystallites in the growth solution (self-nucleation). Initial attempts to grow with constant filtration (by analogy with the work done on KDP—see “High-Damage-Threshold Potassium Dihydrogen Phosphates,” by J. Atherton, Sec. 6.3.6.) did not yield any reproducible improvement in internal crystal quality, because the filtration rate was probably too low to eliminate the growth of molds and bacteria in the growth solution.

Four 32 × 32 mm² aperture, 5-mm thick frequency conversion plates supplied by Cleveland Crystals were sol-gel AR coated for 1, 2 and 3 ω. Although no problems were encountered during the application of the coatings or for several days afterwards, considerable fogging was noticed after approximately 1 month under ambient conditions. While this phenomenon is not

Table 6-14. Five materials with higher efficiency than KDP for doubling or tripling 1.06 μm .

Crystal	Process	$d/d(\text{KDP})$	$\beta_\theta/\beta_\theta (\text{KDP})$
LBO	2 ω I	2.8	1/3
LBO	2 ω II	1.9	1/5
LBO	3 ω I	2.4	2/3
LBO	3 ω II	1.6	1/3
LFM	3 ω II	3.0	2/3
LAP	2 ω I	2.4	1.5
LAP	2 ω II	2.3	1.7
LAP	3 ω I	2.4	1.9
LAP	3 ω II	2.3	1.6
LAAC	3 ω I	1.2	1/2
LAAC			
LAAC		2.4	3/4
LAF	2 ω I	3.4	1

understood, it will have to be eliminated for DLAP crystals to be viable frequency convertors for high power lasers.

Author: S. Velsko

Contributors: Gretchen Kennedy and Ian Thomas, Cleveland Crystals, Inc.

6.4.4. Thermal Behavior of Phase-Matched Harmonic Generation in DLAP

Introduction. Deuterated L-arginine phosphate (dLAP) is a very promising new crystal for use in frequency conversion devices for Nd^{3+} lasers.¹¹⁰ In addition to possessing larger nonlinear couplings and lower threshold powers than KDP, dLAP demonstrates a very unique property: temperature-insensitive phase-matching directions for second-harmonic conversion of 1- μm light. Such behavior is extremely attractive for high-average power applications where it greatly eases the performance requirements for thermal management techniques. Based upon the measured values of the thermal coefficients of refraction,¹¹⁰ temperature-insensitive phase-matching directions in dLAP have been predicted for both Type I and Type II second harmonic conversion of 1.064- μm light.¹¹¹ We have performed a set of experiments to verify these predictions.

Thermal Behavior of the Phase-Matching Loci. We have experimentally determined the thermal behavior of the dLAP second-harmonic generation phase-matching loci by measuring the thermal rates of change of the phase-matching angles in the principal dielectric planes.¹¹² These measurements were performed with an extremely high quality spherical sample of dLAP, fabricated by diamond-turning techniques, immersed in a temperature-controlled index-matching bath installed in the nonlinear optical goniometer apparatus.¹¹³

Figure 6-62 shows this apparatus consisting of a modelocked and Q-switched Nd^{3+} :YAG laser whose 1.064- μm wavelength, 100-ps duration pulses are incident upon the spherical sample mounted in a computer-controlled goniometer. The intensity of the second-harmonic light emerging from the sample is measured with a photomultiplier tube, while the remaining light at the fundamental is directed to an alignment diagnostic. **Figure 6-63** is a photograph of this

dLAP sphere generating a beam of light at the second-harmonic wavelength.

The measured thermal rates of change of the Type I and Type II second-harmonic generation phase-matching angles in the principal dielectric planes are given in **Table 6-15**. The thermal behavior of the Type I second-harmonic generation phase-matching locus is displayed schematically in **Fig. 6-64**. Since the thermal rates of change of the phase-matching angles have opposite signs in the principal planes, there must be phase-matching directions that have a zero-valued thermal rate of change. These are the points about which the locus distorts as the temperature changes. However, the thermal behavior of the type II second-harmonic generation phase-matching locus shown in **Fig. 6-65**, indicates that no temperature-insensitive phase-matching directions exist in this locus. The entire locus simply expands outward as the temperature increases.

Using the measured values of the thermal rates of change of the phase-matching angles in the principal dielectric planes, and a simple model that describes the phase-matching loci of dLAP, we have been able to estimate the locations and relative nonlinear couplings of the temperature-insensitive phase-matching directions in the Type I second-harmonic generation locus. These are given in **Table 6-16**; the azimuthal and polar angles are measured in a coordinate frame where the α axis is chosen as the polar axis, and the azimuthal angle is measured from the α - γ plane.

The discrepancies between the predicted and measured thermal behavior of these phase-matching loci are attributed to errors in the reported values of the thermal coefficients of refraction of dLAP reproduced in **Table 6-17**.¹¹⁰ Because the thermal rates of change of the phase-matching angles depend so sensitively on the thermal coefficients of refraction, even small changes in these coefficients can produce significant changes in the thermal derivatives of the phase-matching angles. The thermal coefficients of refraction were measured by the minimum deviation method with triangular prism-shaped dLAP samples. These measurements did not account for anisotropic thermal expansion, which caused thermal changes in the prism angles of the samples. By using the thermal expansion tensor of dLAP and

Figure 6-62. Schematic diagram of the nonlinear optical goniometer apparatus; see text for details.

Figure 6-63. Photograph of the dLAP sphere, 9.4 mm in diameter, immersed in the temperature-controlled index-matching fluid. A beam at the fundamental frequency, invisible in this photo, is incident on the sphere from the left, and the second-harmonic beam can be seen emerging from the right edge of the sphere. The lack of scattered green light from within the sphere is an indication of the high quality of the crystal.

Figure 6-64. Stereonet plot of the Type I second-harmonic generation phase-matching locus for $1.064\ \mu\text{m}$. The arrows indicate the measured directions of motion of the extremal points of the locus as the temperature increases; the best fit values of the rates at which the phase-matching angles change in the principal planes are given in Table 6-15.

Figure 6-65. Stereonet plot of the Type II second-harmonic generation phase-matching locus for $1.064\ \mu\text{m}$. The arrows indicate the measured directions of motion of the extremal points of the locus as the temperature increases; the best fit values of the rates at which the phase-matching angles change in the principal planes are given in Table 6-15.

Table 6-15. Measured behavior of the extremal points of dLAP phase-matching locus.

	θ (α - β plane) (at 24 °C)	θ (α - γ plane) (at 24 °C)	$d\theta/dT$ (α - β plane) ($\mu\text{rad}/^\circ\text{C}$)	$d\theta/dT$ (α - γ plane) ($\mu\text{rad}/^\circ\text{C}$)
Type I	$22.149^\circ \pm 0.011^\circ$	$35.304^\circ \pm 0.005^\circ$	$+248 \pm 4$	-38.3 ± 1.9
Type II	$37.463^\circ \pm 0.009^\circ$	$47.160^\circ \pm 0.005^\circ$	$+237 \pm 4$	$+30.9 \pm 2.0$

Table 6-16. Approximate temperature insensitive Type I SHG phase-matching directions.

azimuthal angle ϕ	polar angle θ	$\frac{d_{\text{eff}}}{d_{\text{eff}}^{\text{peak}}}$ ^a
$\pm 21.5^\circ \pm 0.5^\circ$	$33.3^\circ \pm 0.1^\circ$	0.97 ± 0.02
$\pm 158.5^\circ \pm 0.5^\circ$	$33.3^\circ \pm 0.1^\circ$	0.32 ± 0.02

^a $d_{\text{eff}}^{\text{peak}}$ is the maximum nonlinear coupling in the Type I phase-matching locus of dLAP.

Table 6-17. Reported values of the thermal coefficients of refraction for dLAP.¹¹⁰

Wavelength (nm)	dn_{α}/dT ($10^{-6}/^{\circ}\text{C}$)	dn_{β}/dT ($10^{-6}/^{\circ}\text{C}$)	dn_{γ}/dT ($10^{-6}/^{\circ}\text{C}$)
1014	-42.0 ± 0.6	-24.2 ± 0.8	-70.5 ± 0.5
579.07	-42.1 ± 0.4	-25.1 ± 0.6	-72.2 ± 0.5

an approximate knowledge of the orientations of the original sample prisms, we have been able to calculate the thermal rate of change of the prism angle in the original samples and thereby correct the original values for this effect. Because the resulting values of the thermal coefficients of refraction are only approximate, we then adjusted the values by a least squares procedure to agree with the measured thermal rates of change of phase-matching angles in the principal dielectric planes. Our best estimates for the thermal coefficients of refraction are given in Table 6-18.

Author: Charles E. Barker

Major Contributors: David Eimerl, Stephan P. Velsko, and Mark S. Webb

6.5. Advanced Materials for Laser Amplification

6.5.1. Introduction

The objective of our research and development program in advanced solid state laser materials is to identify and characterize new materials with the properties required for use as active laser media in a multimegajoule research laser or in an ICF power-plant reactor laser. In 1989 we continued our search for active media for a laser-pumped multimegajoule research laser for a Laboratory Microfusion Facility (LMF). The basic concept is for an efficient flashlamp-pumped laser operated in a long-pulse mode that pumps a gain medium having properties optimized for energy storage and extraction. This approach might lead to a lower cost LMF than Nd:glass laser technology. Since higher pumping efficiencies can be achieved, higher energy storage densities can be obtained, thus reducing the size of the laser, the length of gain medium in the optical path, and the nonlinear phase perturbation caused by the gain medium.

We are developing a new tunable Cr³⁺ activated laser, LiCaAlF₆ (Cr:LiCAF), which is a promising candidate for use as the pump laser. Cr:LiCAF has potential for high flashlamp pumped laser efficiency at about 800 nm, a wavelength ideal for pumping a Nd³⁺ activated gain medium. The laser performance of Cr:LiCAF is reviewed in the article "Laser Properties of

Cr:LiCAF." Our measurements demonstrated that absorbed pump energy can be very efficiently converted into laser output energy in Cr:LiCAF, and calculations show that the flashlamp pumping efficiency of Cr:LiCAF is about a factor of two higher than for Nd:glass. Cr:LiCAF was found to have very low thermal lensing compared with oxide materials. We also measured many of the important optical, thermal, and mechanical properties of Cr:LiCAF. Scattering losses caused by micron-size defects were identified as a serious problem, and improvements in crystal growth and post-growth annealing were developed to reduce these losses to suitably low levels ($\sim 10^{-3} \text{ cm}^{-1}$).

We identified several possible gain media that could be pumped by Cr:LiCAF. The most promising of these is SrF₂:Nd, that offers a 1300 μs lifetime, adequate cross sections for pump absorption and stimulated emission, and a very low nonlinear refractive index. The spectroscopic and laser properties of this medium are presented in the article "Spectroscopy and Gain Measurements of Nd³⁺ in SrF₂ and other Fluoride Hosts."

We made considerable progress in the development of an improved process for crystal growth of large slabs by the Bridgman method. Because of its importance to the ICF effort, Cr:LiCAF was chosen as the material to be grown in this development program. The installation of a large induction heated furnace has been completed, and growth of Cr:LiCAF slabs will be done in 1990.

Author: L. L. Chase

6.5.2. Laser Properties of Cr:LiCAF and CrLiSAF

Introduction. We have evaluated the laser performance of the two new materials: Cr³⁺-doped LiCaAlF₆, and LiSrAlF₆, (LiCAF and LiSAF). Very little was known about these materials prior to our investigations¹¹⁴ since we grew the first single crystals and recognized the potential for laser action.¹¹⁵ Our studies have included a complete characterization of the optical, thermal, and mechanical properties, as well as the actual laser performance. Most of this effort has concentrated on the Cr:LiCAF crystal, although Cr:LiSAF has been characterized to a limited extent as well.

Laser Performance. The absorption and emission spectra of Cr:LiCAF are pictured in Fig. 6-66. As is the case for most Cr-doped crystals,

Table 6-18. Corrected and adjusted thermal coefficients of refraction for dLAP.

Wavelength (nm)	dn_{α}/dT (10 ⁻⁶ /°C)	dn_{β}/dT (10 ⁻⁶ /°C)	dn_{γ}/dT (10 ⁻⁶ /°C)
1064	-37.3 ± 1.7	-53.0 ± 1.7	-63.0 ± 1.7
532	-36.4 ± 1.7	-53.4 ± 1.7	-66.9 ± 1.7

Figure 6-66. Absorption and emission spectra of Cr:LiCAF plotted on absolute cross section scales. The $E \parallel c$ and $E \perp c$ spectra are reported since the material is uniaxial.

the absorption spectrum consists of three bands. This is particularly useful for flashlamp-pumping, since approximately twice as much flashlamp light can be absorbed by the Cr^{3+} -doped crystals, compared to Nd-doped glass.¹¹⁶ The emission is peaked near 760 nm, and, in principle, may be tuned throughout the observed band, although the demonstrated tuning range thus far only includes the 720–840 nm region.¹¹⁵ Note that since LiCAF is uniaxial, spectra with the electric field parallel and perpendicular to the optic axis ($E \parallel c$ and $E \perp c$), are both reported in Fig. 6-66. The $E \parallel c$ emission spectrum has a higher cross section, resulting in a lower threshold than would otherwise be found for an isotropic system.

We have analyzed the flashlamp-pumped laser performance of a 0.25×3.15 in. laser rod of Cr:LiCAF in a conventional laser head.^{117,118} The measured slope efficiencies, η , are plotted in Fig. 6-67 as a function of the transmission of the output coupler, T . The data have been fitted to the equation:

$$\eta = \eta_0 \frac{\ln(1-T)}{\ln(1-T) + \ln(1-L_d)} \quad (1)$$

where L_d is the double-pass loss and η_0 is the slope efficiency that would be obtained if the material and cavity had zero optical loss. The value of $L_d = 47\%$ corresponds to a loss of 4.0%/cm. This loss level is corroborated by a Findlay-Clay analysis, and by direct measurements of the scattering. While this particular laser rod had high loss, the extrapolated efficiency of $\eta_0 = 4.14\%$ provides a strong indication that the efficiency of Cr:LiCAF will be high when the crystal growth technique is perfected and low-loss material becomes available.

Having observed the detrimental effects of passive loss in Cr:LiCAF, we next conducted a systematic investigation of the magnitude and nature of these losses for material grown by the Czochralski and gradient freeze methods.^{118,119} Both absorption and scattering were found to contribute to the overall loss. The parasitic absorption is caused by the existence of a small amount of Cr^{2+} on the Ca^{2+} site. This loss is typically small at the laser wavelength $\leq 0.1\%/cm$. The scattering losses, on the other hand, are significant, and are in the range of 1–4%/cm for the Czochralski-grown samples, and

0.1–0.8%/cm for the gradient freeze material. Interestingly, the scattering is about 50% greater for π , versus σ , polarization of light, suggesting that the defects giving rise to the scattering have a specific crystallographic orientation. Perhaps the most significant result is that Cr:LiCAF crystals exhibiting loss as low as 0.1%/cm have been generated by Atherton and co-workers at LLNL. Although these crystals are small (1–2 cm), they demonstrate that it is feasible to grow high-quality material.

We have also demonstrated laser action in the strontium analog of Cr:LiCAF, $\text{LiSrAlF}_6:\text{Cr}^{3+}$ or Cr:LiSAF. This material has interesting properties, in that it possesses a significantly larger emission cross section than Cr:LiCAF ($5 \times 10^{-20} \text{ cm}^2$ versus $1.3 \times 10^{-20} \text{ cm}^2$), and a wider tuning range.¹²⁰ The experimental tuning range is depicted in Fig. 6-68, where a Kr-laser was used as the pump source. Cr:LiSAF may ultimately prove to be a viable material for commercial purposes.

Thermo-optical and thermomechanical properties. The thermal, mechanical, and optical properties of LiCAF relevant to laser performance are listed in Table 6-19, where data for $\text{LiYF}_4:\text{Nd}^{3+}$ (Nd:YLF) is included for comparison.¹²¹ It is apparent from Table 6-19 that LiCAF and YLF have similar properties. The fracture toughness, K_{IC} (a measure of the tendency to fracture), is nearly the same, within experimental uncertainty. The change of refractive index with temperature, dn/dT , is negative for both materials, and, as we shall see, this gives rise to very low thermal lensing. The quantity, $(\rho C_p)^{-1}$, characterizes the temperature rise associated with a given deposition of heat. The expansion coefficient, α , and the thermal conductivity, κ , are also listed. Since Nd:YLF has been judged to possess adequate physical properties for low and medium power laser applications, we can suggest, on the basis of the data in Table 6-19, that Cr:LiCAF will be adequate as well for many applications.

Since thermal lensing can adversely impact the beam quality of a laser system, we have directly measured the magnitude of the thermal lens with the aid of Moiré deflectometry, for the flashlamp-pumped rod geometry. The Cr:LiCAF results are shown in Fig. 6-69, where they are contrasted with those of alexandrite, (Cr^{3+} -doped BeAl_2O_4). The prompt negative lens is due to the uneven heating of the laser rod, such that

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Figure 6-67. Plot of the flashlamp-pumped slope efficiency of Cr:LiCAF as a function of the output coupler transmission. The extrapolated slope efficiency, η_o , and the double-pass loss, L_d , were calculated with Eq. 1.

02-03-0293-0496

Figure 6-68. Tuning range of Cr:LiSAF obtained with Kr-laser pumping.

Table 6-19. Comparison of the properties for two laser hosts (see text for definition of quantities).

Laser host	K_{IC} (MPa $\sqrt{m}$)	dn / dT ($10^{-6} / ^\circ C$)	$(\rho C_p)^{-1}$ ($^\circ C / J / cm^3$)	α ($10^{-6} / ^\circ C$)	κ (W/(m $^\circ C$))
LiCaAlF ₆	0.36	-4.6, c -4.2, $\perp$ c	0.36	3.6, c 22, $\perp$ c	5.1, c 4.6, $\perp$ c
LiYF ₄	0.27	-4.3, c -2.0, $\perp$ c	0.33	8, c 13, $\perp$ c	4.3

TB-Payne-8438-2A

Figure 6-69. Measurement of the magnitude of the pump-induced thermal lens for Cr:LiCAF and alexandrite as a function of time.

the edges are hotter than the center. As the rod is cooled by the flowing water the sign of the lens reverses. It is clear from the data in Fig. 6-69 that the thermal lens of Cr:LiCAF is much less than that of alexandrite. In fact, the lensing along the *c* axis is virtually athermal. The low thermal lensing power is typical of fluorides, where the value of dn/dT is negative and tends to cancel the stress-induced electro-optic contribution to the lens.¹²²

Physics of the Cr³⁺ Impurity. The favorable laser performance of Cr:LiCAF and Cr:LiSAF owes to the adequate mechanical, thermal, and optical host properties; the feasibility of crystal growth, and the nature of the host site where the Cr³⁺ ion is located. We now consider how the crystal site impacts the transition strengths and energy levels of Cr³⁺.

The crystalline lattice of LiCaAlF₆ is depicted in Fig. 6-70,¹¹⁴ and Cr³⁺ is known to substitute for the Al³⁺ ion. The key element that needs to be recognized is the nature of the odd-parity field present at the Al³⁺ site, since this crystal field is responsible for inducing oscillator strength into the nominally forbidden d-d transitions of Cr³⁺. It is describable as having t_{2u} symmetry within the O_h point group, and may be imagined as resulting from the distortion where the three fluorines above the Al³⁺ are rotated clockwise, and the three below are rotated counter-clockwise, (see Fig. 6-70). This "twisting" of the fluorines gives rise to the polarization dependence of the absorption spectra.¹²³ The emission is also polarized along the *c* axis, thereby concentrating the emission in particular directions and allowing for a reduced laser threshold. We have also shown that the detrimental effects of excited state absorption (ESA) are mitigated in these systems, since the ESA band does not gain oscillator strength from the t_{2u} distortion, while the absorption and emission bands do. As a result, the emission cross section has a substantially higher cross section than does the ESA band in the lasing region.¹²⁴

Summary. We have comprehensively evaluated the laser performance, physics, and the thermal, mechanical and optical properties of Cr:LiCAF. While the crystal growth needs to be refined and perfected, we are convinced that Cr:LiCAF and Cr:LiSAF have significant potential as new laser materials.

Authors: S. A. Payne and L. L. Chase.

Major Contributors: L. K. Smith, J. A. Caird, J. E. Marion, H. W. H. Lee, W. L. Kway, B. W. Woods, R. S. Hughes, L. E. Davis, G. T. Ullery, E. R. Prochnow, G. D. Wilke, N. D. Nielsen, and B. H. T. Chai (Univ. of Central Florida).

6.5.3. Development of CrLiCAF and CrLiSAF

Introduction. The fluoride materials LiCaAlF₆:Cr³⁺ (Cr:LiCAF) and Cr:LiSrAlF₆ (Cr:LiSAF) have shown to be promising new solid-state laser materials.¹²⁵⁻¹²⁸ We have been developing them for possible use in a multimegajoule class laser-pumped-laser for a Laboratory Microfusion Facility. To date we have both laser-pumped, and flashlamp-pumped these materials. Using Kr-laser pumping, the laser efficiency of Cr:LiCAF was comparable to alexandrite, and that of Cr:LiSAF was somewhat lower. The flashlamp-pumped performance of Cr:LiCAF was also found to be encouraging, although there have been high losses in the rods used in the measurements. Given these favorable results, we have worked to further improve these materials and optimize their performance. In this article we discuss the growth and spectroscopic properties of Cr:LiCAF crystals with high Cr³⁺ concentrations and review our current understanding of the origin of the passive losses in crystals grown by Czochralski (CZ), gradient freeze (GF), and Bridgman methods. In this section we discuss the CZ and GF crystals. In Article 6.5.5. we discuss the Bridgman results.

High Cr³⁺ Concentrations. Crystals were grown by the Czochralski method with starting concentrations of 5, 10, and 15 mole %, and the corresponding measured concentrations in the crystals were 2.6, 5.2, and 8.0 mole %, respectively. The quality of the crystals was not noticeably affected by the Cr³⁺ doping level. The crystals were very uniformly doped, and a distribution coefficient for Cr³⁺ of 1.0 was deduced from the Cr³⁺ optical absorption coefficient at various positions along the length of the boules. It is possible that the lower concentration in the crystals compared with the starting values is caused by the presence of chromium oxides in the starting material, that is caused by the introduction of oxides during the growth process. In fact, we found that it was difficult to completely eliminate oxides from the CrF₃ starting material by hydrofluorination before growth, and a black layer of chromium

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Figure 6-70. Crystal structure of the LiCaAlF_6 lattice, reproduced from the work of Viebahn.¹¹⁴

oxide was found at the bottom of the residual charge in the crucible.

Identical absorption spectra and Cr^{3+} emission lifetimes were measured for samples with all of these Cr^{3+} concentrations. The absorption spectra are illustrated in Fig. 6-72, which shows the absorption cross section of samples with 0.3 and 8.0 percent Cr^{3+} . These spectra are identical to within experimental error. Thus, the optical properties of Cr^{3+} are not affected by doping level, at least up to 8 mole %, where a considerable fraction of nearest neighbor pairs is expected. This behavior results from the favorable crystal structure of Cr:LiCAF. Cr:LiSAF-neighboring Al sites are not coordinated by common fluorine ions, so interactions between Cr^{3+} nearest-neighbor pairs are very weak.¹²⁹

The absorption spectrum in the long wavelength part of Fig. 6-72 is expanded by a factor of 10 to highlight the presence of a broad absorption band peaked at 1 μm . This band is probably caused by the presence of Cr^{2+} ions on the Ca^{2+} site. Similar features have been observed and attributed to Cr^{2+} in other halide materials doped with Cr.¹³⁰⁻¹³³ Furthermore, in a Cr:LiCAF sample doped with high-purity Cr metal powder, the 1- μm absorption feature increased severalfold, which suggests that the absorption is enhanced in a reducing environment, as would be expected if it were caused by Cr^{2+} . This band does not contribute significantly to the loss in the 730–900 nm lasing region of Cr:LiCAF, and to be reduced by improving the stoichiometry of the CrF_3 starting material, and by controlling the crystal growth conditions. Fig. 6-73 shows the absorption near 1 μm in Cr:LiSAF doped with 2 mole % Cr. The Cr^{2+} absorption peak is reduced and shifted to longer wavelengths than in Cr:LiCAF. It will not be a significant source of loss in the 780–950 nm lasing region of Cr:LiSAF.

Since Cr:LiCAF and, to a lesser extent, Cr:LiSAF, have relatively small absorption cross sections at the peaks of their absorption bands, relatively high doping densities are required in order to efficiently absorb flashlamp light. The doping levels we have obtained in Cr:LiCAF are more than adequate for most applications. For example, Fig. 6-74 shows a calculation of the single-pass flashlamp-pumping efficiency of Cr:LiCAF as a function of Cr^{3+} doping density-thickness product in units of cm^{-2} .

For comparison, an identical calculation for LHG-8:Nd³⁺ laser glass is also shown. The calculated pumping efficiency is defined as the ratio of the available output energy in the laser transition divided by the flashlamp electrical input, neglecting radiative decay losses. These calculations are based on the assumption that every flashlamp photon passes through the medium only once, i.e., every area of the flashlamp is mapped onto the surface of the slab-shaped medium.¹³⁴ The flashlamp model assumes a xenon lamp operated at 20% of its explosion energy with a 300-ms pulse width at 10% power points.

The calculated pumping efficiency of Cr:LiCAF becomes nearly independent of doping-thickness. The calculated pumping efficiency of Cr:LiCAF becomes nearly independent of doping density at about $2 \times 10^{20} \text{ cm}^{-2}$. This corresponds to a Cr concentration of 4 mole % for a 5-mm-thick slab or for a rod with similar cross section. This doping level is easily achieved in Cr:LiCAF, and probably also in Cr:LiSAF. The pumping efficiency calculations used to obtain the plots in Fig. 6-74 are known to overestimate the measured values for laser disks in close-coupled cavities by about 30%.¹³⁴ Nevertheless, the calculations suggest that pumping efficiencies in the range 10–15% should be attainable.

Loss Measurements. For efficient extraction, optical losses at the laser wave length must be below $0.5\% \text{ cm}^{-1}$. We measured losses in Cr:LiCAF samples using several methods: (1) transmission of a semiconductor diode laser at 780 nm; (2) scattering losses at 1064 nm in an integrating sphere; (3) absorption loss at 1064 nm in a calorimeter; (4) Findlay-Clay analysis of flashlamp pumped laser data for a Cr:LiCAF rod. The diode laser data was corrected for ground state absorption of Cr^{3+} using the absorption spectrum in Fig. 6-73.

The integrating sphere and calorimeter data are shown for several samples grown by Czochralski and gradient freeze in Table 6-20. The scattering loss is $(1-2) \times 10^{-2} \text{ cm}^{-1}$ in the CZ material and $(1-8) \times 10^{-3} \text{ cm}^{-1}$ in the GF samples. In samples grown by both methods the scattering loss was anisotropic, about 50% larger for π than for σ polarization. The losses were lower in the 1-cm-diam GF samples (GF#1,3) than in the 2-cm-diam sample (GF#2). The losses are believed to be caused by the effects of strains in the samples that nucleate scattering defects along shear bands during cooling of the boules.

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Figure 6-72. Absorption cross-section spectra of Cr:LiCAF crystals. The spectrum of the 8% sample is displaced upward for clarity. Note that the cross-section scale at the right corresponds to the data for $\lambda > 725 \text{ nm}$.

70-00-0290-1327 pub

Figure 6-73. Absorption cross-section spectrum of Cr: LiSAF.

70-00-1089-2993pub

Figure 6-74. Calculated single-pass flashlamp-pumping efficiency for Cr: LiCAF and Nd ³⁺: LHG-8 glass.

Table 6-20. Comparison of intergrating sphere and calorimeter data for samples grown by Czochralski and gradient freeze.

Sample	Cr+3 conc. (mol%)	Polar- izaton	Loss at 1064 nm (%/cm)	
			Scattering	Absorpton
Czochralski:				
undoped	0	π	1.4	0.06
		σ	0.8	0.05
#1	~2	π	2.2	0.20
		σ	1.2	0.23
#2	3.7	π	2.4	0.67
		σ	2.2	0.90
#3	6.8	π	1.0	0.60
		σ	0.6	0.73
#4	8.8	π	1.5	0.61
		σ	0.9	0.67
Gradient freeze				
#1	~2	π	0.31	0.31
		σ	0.22	0.38
#2	~2	π	0.80	0.13
		σ	0.54	0.15
#3	~2	π	0.09	0.33
		σ	0.07	0.29

The absorption losses at 1064 nm are similar in the CZ and GF crystals. The very low absorption in the undoped CZ sample suggests that Cr doping is the source of the absorption in the other boules. The 1- μm absorption band in Fig. 6-72 is the cause of this absorption. The data in Table 6-20 show that this absorption does not increase linearly with Cr concentration. The fraction of Cr^{2+} centers apparently decreases at higher Cr doping levels. Note that the GF samples have absorption losses comparable to the CZ sample with a similar doping level (2 mole %).

Defect Characterization. Examination of several Cr:LiCAF samples by optical microscopy shows that there are micron-size scattering centers that form either clouds of isolated defects (smoke) or arrays of defects. The arrays are often oriented along lines in crystallographic directions corresponding to dislocations perpendicular to the c-axis. The arrays are sometimes also curved and suggest nucleation along the growth direction. Measurement of the spectral distribution of the scattering performed on both undoped and Cr-doped samples with a spectrophotometer shows that the scattering is only very weakly dependent on wavelength in the range from 0.2 to 2 μm . This shows that the defects that actually cause the bulk of the loss at the laser wavelength have dimensions $>1 \mu\text{m}$, since smaller defects would display the well-known Rayleigh dependence on the wavelength (inverse fourth power of the wavelength) for the integrated scattering loss. Thus, the defects observed in the microscope are the important ones.

Some evidence of the nature of the defects can be inferred from the magnitude of the losses by using a formula appropriate to the limit of large particle radius, $a > \lambda$. The scattering cross section for such large particles is twice the geometric cross section,¹³⁵ and the scattering loss, β , for spherical particles can be shown

$$\beta = \frac{3}{2} \frac{NV}{a} \left[\frac{m^2 - 1}{m^2 + 2} \right]^2 \quad (1)$$

where N is the number of scatterers per unit volume, $V = 4\pi a^3/3$ is their volume, and m is the ratio of the refractive index of the scatterers to that of the Cr:LiCAF matrix. If we assume that the scatterers are voids ($n = 1$) or oxide particles

($n \approx 1.8$), with radii $a \approx 1 \mu\text{m}$, the volume fraction NV occupied by the scatterers that is required to account for the measured values of $\beta \approx 10^{-2} \text{ cm}^{-1}$ is about 10 ppm. This estimate is roughly consistent with the number and size of the defects observed microscopically. For fluoride particles, CaF_2 , for example, $n = 1.43$; for such low dielectric discontinuities Eq. (1) predicts that volume fractions in excess of 1% are required to produce the losses. This suggests that the particles are oxides or voids, although they might also consist of some low-density fluoride material with an effective refractive index substantially lower than that of LiCAF. Electron microscopy, electron microprobe, and SIMS measurements are in progress, and the results of those investigations will be presented elsewhere.

Summary. We have shown that Cr:LiCAF is capable of very efficient flashlamp pumping and high-extraction efficiency at modest fluences. Cr^{3+} doping levels of at least 8 mole % are possible without noticeable degradation of crystal quality.

Author: L. L. Chase

Major Contributors: M. S. Shinn, S. A. Payne, W. L. Kway, and J. A. Caird

6.5.4. Evaluation of the $\text{SrF}_2\text{Nd}^{3+}$ Laser Material

Introduction. We have explored the possibility of utilizing a laser-pumped laser architecture (Loki) for a multimegajoule fusion laser system. The schematic description appears in Fig. 6-75, where the system is shown to be divided into two components: the pump laser and the storage medium. The pump laser requirement might be fulfilled by Cr:LiCAF, or by laser diodes. The identification of a suitable storage medium forms the subject of this section. Numerous issues impact the viability of the storage medium, including: (a) the absorption coefficient and cross section; (b) the emission lifetime, wavelength, and cross section; and (c) other issues such as crystal growth, fabrication, optical quality and the damage threshold. A much more detailed account of the work is reported in Ref. 136.

Spectroscopy and Gain Measurements.

Trivalent rare-earths, such as Nd^{3+} , may be substitutionally incorporated into the fluorite structure hosts: CaF_2 , SrF_2 and BaF_2 . The Nd^{3+}

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Figure 6-75. Schematic representation of a fusion laser concept (Loki), involving a laser-pumped storage medium. Some of the important properties of the storage medium are noted.

ion therefore requires an additional charge compensator to maintain the electrical neutrality of the lattice. The various possible types of charge compensation are represented in Fig. 6-76. We have found that there are two major types of charge compensator appearing in $\text{SrF}_2:\text{Nd}^{3+}$, involving nearest-neighbor (nn) interstitial fluorines (F_i^-), and clusters. The nn interstitial fluorine occupies the normally empty site next to the Nd^{3+} ion. The clusters can be quite complicated in their structure, but, in any case, lead to the formation of unfavorable centers where the Nd^{3+} luminescence is quenched. The relative amount of clusters and nn sites is described in Fig. 6-77, where it is seen that at Nd^{3+} concentrations below 0.2 at.%, and the crystal becomes practically free of clusters.

Nd^{3+} -doped CaF_2 did not prove to be a useful storage medium because excessive clustering occurs at Nd^{3+} concentrations an order of magnitude less than that characterizing SrF_2 . On the other hand, $\text{BaF}_2:\text{Nd}^{3+}$ is not useful because the next nearest-neighbor type of compensation predominates (rather than nearest-neighbor); the spectroscopic properties of this center are unfavorable because the emission cross section turns out to be too small. This situation occurs because the interstitial fluorine is responsible for inducing oscillator strength into the 4f-4f transitions of Nd^{3+} and, therefore, the further it is from the Nd^{3+} ion, the weaker the transitions are expected to be. Fortunately, the nearest neighbor compensation found for $\text{SrF}_2:\text{Nd}^{3+}$ does provide adequate absorption and emission cross sections.

The absorption spectra of $\text{SrF}_2:\text{Nd}^{3+}$ are shown in Fig. 6-78; the final states of the transitions are designated on the figure. The data in Fig. 6-78 may be used in conjunction with the Judd-Ofelt method of analysis to calculate that the radiative emission lifetime of the $^4\text{F}_{3/2}$ state is 1410 μs . For comparison we have measured the emission lifetime of $\tau_{\text{em}} = 1280 \mu\text{s}$ and the quantum efficiency of $\eta = 0.87$, and therefore can independently determine the radiative lifetime to be $\tau_{\text{rad}} = \tau_{\text{em}}/\eta = 1470 \mu\text{s}$, in agreement with the Judd-Ofelt calculation.

The emission spectrum associated with the $^4\text{F}_{3/2} \rightarrow ^4\text{I}_{11/2}$ transition appears in the upper right-hand side of Fig. 6-79. The peak emission cross section is calculated to be $1.68 \times 10^{-20} \text{ cm}^2$, from the radiative lifetime, refractive index, and the spectrum shown in the figure. The absorption spectrum from 720 to 750 nm is

reproduced in the upper left-hand side. The lower frames contain the results of a pump-probe experiment, in which an alexandrite laser operating at 796 nm served to pump the Nd^{3+} ions to the $^4\text{F}_{3/2}$ excited state. Both bleaching of the absorption and gain within the emission region have been observed, as indicated on the figure. The pump-probe experiment gave an independent measure of the emission cross section of $1.29 \times 10^{-20} \text{ cm}^2$, which, within experimental error, agrees with the determination based on the emission spectrum. Interestingly, the $^4\text{F}_{3/2} \rightarrow ^2\text{G}_{9/2}$ excited state absorption (ESA) transitions can be observed within the bandwidth of the $^4\text{F}_{3/2} \rightarrow ^4\text{I}_{11/2}$ gain region, as noted on the figure. Both Judd-Ofelt calculations and the energy level positions provide support for this assignment of the ESA lines. The impact of the ESA on the gain cross section is to reduce its magnitude by about 4%.

Several important physical and optical properties of $\text{SrF}_2:\text{Nd}^{3+}$ are summarized in Table 6-21. This material may potentially serve as a laser-pumped storage medium, in part because of the adequate absorption, emission, and gain properties, but also because of the low nonlinear refractive index. In addition, SrF_2 can be grown in large size with excellent optical quality. The nonunity value of the distribution coefficient of $k = 1.5$, however, suggests that it will be difficult to obtain uniform doping of the medium, although this problem is not inherently unsolvable. The maximum absorption coefficient of 0.30 cm^{-1} at 796 nm that is attainable without clustering is somewhat low, although the optimization of the Nd concentration and the system architecture has not yet been carefully considered.

Authors: S. A. Payne, J. A. Caird, L. L. Chase, L. K. Smith, N. D. Nielsen, and W. F. Krupke

Major Contributors: S. Townsend, R. C. Powell (Oklahoma State University), R. Sparrow (Optovac, Inc.)

6.5.5. Advanced Large-Dimension Crystal Growth

Introduction. The goal of our research is to develop new technologies for the growth of slab-shaped optical crystals with the large dimensions and high quality required for

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Figure 6-76. Illustrations of the perfect fluorite lattice, an uncompensated Nd^{3+} ion, Nd^{3+} compensated by nearest-neighbor and next nearest-neighbor interstitial fluorines (F_i^-), and an example of a cluster.

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Figure 6-77. Plot of the concentration of Nd³⁺ ions found in clusters as a function of the concentration of nearest-neighbor compensated centers (C_{4v}-sites). The symbol, α , denotes the absorption coefficient.

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Figure 6-78. Absorption spectrum of $\text{SrF}_2:\text{Nd}^{3+}$ plotted on an absolute cross section scale. The final states of the transitions giving rise to the absorption bands are designated.

MG-L-U-7812-12

Figure 6-79 Absorption and emission spectra shown in the upper frames; bleaching and gain as determined over the same spectral range with a pump-probe experiment are shown in the lower frames.

Table 6-21. Physical and optical properties of SrF₂:Nd³⁺.

Property	Value
Peak emission	
cross section	$\sigma_{\text{em}} (1036.5 \text{ nm}) = (1.7 \pm 0.2) \times 10^{-20} \text{ cm}^2$
Emission lifetime	$\tau_{\text{em}} = 1280 \pm 20 \text{ } \mu\text{s}$
Radiative lifetime	$\tau_{\text{rad}} = 1470 \pm 100 \text{ } \mu\text{s}$
Pump-band absorption	
cross section	$\sigma_{\text{abs}} (796 \text{ nm}) = (0.79 \pm 0.1) \times 10^{-20} \text{ cm}^2$
Maximum pump-band	
absorption coefficient	$\alpha (796) \text{ nm} = 0.30 \text{ cm}^{-1}$
Nonlinear refractive index	$n_2 = 0.50 \times 10^{-13} \text{ esu}$
Nd ³⁺ distribution coefficient	$k = 1.5$

advanced solid-state lasers. Limitations in both optical quality and scale-up of the Czochralski technique, the principal process currently used to grow most laser crystals, continue to inhibit the utilization of several important laser crystals. The process we are developing is flexible, and allows growth of many different materials, including oxides, fluorides, and semiconductors. The application includes optics, electro-optics, and electronics. The project is jointly supported by the ICF program and by LLNL Director's Initiative funds.

In the past year, we received and installed a state-of-the-art Bridgman furnace from Crystallox Ltd. (Wantage, U.K.), and used it to grow several slab-shaped crystals. We verified its operation up to 2100°C, and measured a temperature stability of 0.1°C, which we had predicted on the basis of an in-depth thermal analysis prior to delivery of the furnace.¹³⁷ We continued development of high-temperature sensing and seeding diagnostics. In particular, we demonstrated that sapphire optical fibers could be used to monitor and control temperature up to 1900°C. We also developed a novel seeding technique for high temperature crystals based on a large difference in electrical conductivity between the solid and melt phases. Early this year, we made a transition from a focus on the growth of YAG:Nd, a high-temperature oxide, to an emphasis on a lower melting fluoride crystal, LiCaAlF₆:Cr (LiCAF). LiCAF shows exceptional promise as an efficient pump laser for a laser-pumped-laser laboratory microfusion driver.¹³⁸ We grew high-quality, small-scale LiCAF in our existing gradient freeze furnace, and have produced several LiCAF slabs in the Crystallox furnace.¹³⁹ We characterized the Crystallox furnace both by modeling it with the induction heating model developed last year, and by measuring temperature profiles in the heaters and melt with translating thermocouples.¹⁴⁰ Finally, we made significant progress in determining the nature and cause of optical defects in LiCAF.

Furnace Development, Installation and Characterization. Our advanced process for growing high-quality slab-shaped crystals is based on the vertical Bridgman process, in which a shaped ampoule is lowered through a three zone furnace to grow the crystal (Fig. 6-80). The hot and cold zones are independently controlled to produce the melting point temperature in the middle of the adiabatic zone, with an

appropriate temperature gradient at that point. The adiabatic zone is designed to reduce lateral heat transfer, thereby flattening the melt/crystal interface. To grow a crystal by this technique, a seed crystal of the desired orientation is placed in a socket at the bottom of the ampoule, and feedstock of the appropriate composition is added to the ampoule. The ampoule is positioned in the furnace such that a portion of the seed is near the middle of the adiabatic zone. The feedstock is entirely in the upper part of the furnace, and is therefore completely melted down to the seed. Crystal growth is initiated by translating the ampoule downward in the furnace at the desired growth rate, and the slab-shaped crystal grows with the same orientation as the seed. To realize high-quality, slab-shaped crystals, we devised specifications for a large Bridgman furnace based on heat transfer and fluid dynamics calculations previously described.^{141,137}

The Bridgman furnace designed and built to our specifications by Crystallox Ltd. was delivered to the Laboratory in early 1989 (Fig. 6-81a). It was installed, and we performed acceptance tests verifying overall performance and temperature capabilities during the month of February. The Crystallox design is based on inductively heated graphite disks (called "susceptors" because they couple with the electromagnetic field produced by the coils) that conduct heat to a rectangular channel that accommodates the ampoule (Fig. 6-81b). In this furnace, two independently controlled coils, connected to two 40-kW, 25-kHz generators couple to graphite susceptor stacks separated by a fibrous graphite insulator that forms the adiabatic zone. A water-cooled horizontal copper baffle separates the two coils to prevent interaction between the upper and lower coils.

We specified that the furnace should operate up to 2100°C, to allow growth of high-temperature materials such as YAG:Nd that melts at 1970°C. In our induction heating calculations, we predicted that this temperature would be reached using about 25 kW from the two generators. Because the capacity of these generators is 80 kW (2 × 40 kW), we expected to easily reach the maximum operating temperature of 2100°C.¹⁴¹ The acceptance test results verified this analysis, as approximately 24 kW was required to reach 2100°C. We also specified that the furnace should have a temperature gradient of at least 10°C/cm at its

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Figure 6-80. In the Bridgman process, a prismatic ampoule is lowered through a three-zone furnace to grow the crystal. For slab-shaped optical crystals, this technique produces a melt with stable, laminar fluid flow, and a crystal free of hydrodynamically induced striations. The adiabatic zone flattens the crystal/melt interface.

Figure 6-81. (a) Photograph of advanced Bridgman furnace developed with Crystallox. (b) Schematic representation of Crystallox furnace internals, including susceptor, insulation, and coils.

maximum temperature to allow growth of high-quality crystals. In tests at Crystalox and at the Laboratory, this gradient specification was also met.

Temperature stability is also a critical determinant of crystal quality. In the Crystalox design, power is controlled rather than temperature, primarily because power can be measured at all levels more accurately than temperature. While changes in the power produce accompanying variations in temperature, the large thermal mass of the furnace attenuates these temperature fluctuations at the crystal interface to acceptably small levels.

The time constants of the furnace at 2000°C and 800°C are approximately 2 and 5 hours, respectively. These values were determined by turning off the power and plotting the logarithm of temperature versus time; the time constant is the negative of the initial reciprocal slope of this line. We measured the temperature stability of the Crystalox furnace with an Accufiber (Beaverton, OR) sapphire optical pyrometer. Over a two-day test period, we observed fluctuations in temperature of less than 0.1°C over a several-hour time period, but with a diurnal variation of 0.25°C. We attributed this long term oscillation to variations in the cooling water temperature. This long term fluctuation will be reduced by improving temperature stability of the closed-loop chiller, and may also be compensated for by using advanced computer control software in the coming year.

High Temperature Sensing and Seeding Diagnostics. Precision temperature sensors are required to measure both long-term drift (time scales of a few hours) and short-term fluctuations (1–100 sec.). We must detect and correct for long-term temperature changes to prevent the melt/crystal interface from migrating out of the adiabatic zone. The interface must remain within the adiabatic zone to retain the large axial temperature gradient and slightly convex isotherms necessary for growth of high-quality crystals. Short-term fluctuations caused by power instabilities may introduce striations or banding in high-temperature crystals, such as Nd:YAG, if the amplitude exceeds about 0.1°C. These short-term temperature excursions are dampened in the Crystalox furnace by the large thermal inertia of the susceptor stack, giving the large time constant discussed above. However, for some crystals it is important to detect short-

term fluctuations in temperature associated with unstable or periodic melt flow. In our experience growing LiCAF, we have not observed short-term fluctuations in either the susceptor or the melt.

For this effort, we are using optical fiber thermometry, a form of pyrometry that uses a narrow sapphire fiber to transmit ambient thermal radiation from the hot interior of the furnace to an external, filtered photodetector whose signal is used to determine the radiation temperature using Planck's law. This technique is particularly suited to precision temperature measurements in high-temperature induction furnaces because of the absence of sensor electrical circuits, and the material stability of sapphire at high temperatures.

While the technique has been verified at 1000°C, its use near 2000°C has not. To apply sapphire optical thermometry to this temperature regime, the thermal absorption (or emission) of sapphire must be negligible compared with that of the furnace interior at the same temperature.¹⁴² To determine the effect of thermal emission on temperature reading, we used an image furnace to heat a 1-cm length of the fiber to approximately 1900°C, and compared the thermal radiation signal with that from a calibrated light source at the same temperature. The ratio of the thermal emission (per unit length) to the calibrated transmission indicates a negligible error (see Table 6-22). Using Planck's law, the absorptivity of sapphire at this temperature can be deduced as 0.002 cm⁻¹, a negligible reduction in transmission. Thus sapphire optic thermometry at 2000°C appears feasible.

Seeding Diagnostic for YAG:Nd. The classical Bridgman technique is limited by its inability to directly detect the conditions necessary to begin crystal growth. In the Czochralski process, a seed is touched to the melt and slowly withdrawn while rotating; this entire procedure can be viewed by an operator. In the Bridgman process, the equivalent seeding procedure is "blind," and requires an indirect indication that the seed has only been partially melted (rather than completely or not at all). Thermocouples wrapped about the exterior of the crucible are an adequate seeding diagnostic for materials that melt below about 1400°C, because their accuracy relative to the temperature gradient at the seed is sufficient to guarantee partial melting of the seed. For

Table 6-22. Absolute error in fiber temperature measurement resulting from fiber emission (per unit length) at furnace temperatures.

Optical fiber temperature (°C)	Specific emission/ transmission (cm ⁻¹)	Absolute error (°C)
1800	0.0025	0.6
1875	0.0007	0.2
1950	0.0014	0.4

higher-melting materials such as YAG:Nd, a new technique was needed.

We have developed a self-calibrating technique for detecting seed meltback based on the 100- to 1000-fold change in electrical conductivity of ionic solids upon melting. We core-drill the ampoule base and seed crystal to accommodate a refractory-wire electrode (Fig. 6-82). When the seed end melts, the wire forms a closed circuit with the crucible through the highly conductive ionic melt, and the resultant abrupt drop in circuit resistance is detected electronically. Growth of the large crystal can then commence.

The contact angle between the crystal/melt interface at the wire is acute, as is the angle at the ampoule wall. Therefore, defects that propagate normal to the crystal/melt interface are focused to extinction above the wire and the crucible walls, and the dislocations created by the wire electrode are self-healing. We confirmed this effect on small-scale experiments growing YAG:Nd.

Bridgman Process for Growth of LiCAF. We made a transition from the growth of YAG:Nd to LiCAF in early 1989. We were originally growing YAG:Nd because of general interest in it as a laser material, and because it possesses a range of crystal growth challenges that made it an appropriate choice for process development. However, in the last year, LiCAF has emerged as a potential candidate for use in a laser-pumped-laser LMF driver to succeed the Nova laser in the ICF program. Fig. 6-83, is a schematic of the laser-pumped-laser that illustrates the use of LiCAF. In this concept, a LiCAF slab approximately $1 \times 10 \times 20 \text{ cm}^3$ absorbs flashlamp energy, and emits it at about 800 nm (long pulse) to a Nd-doped storage material. Light is extracted from the storage medium at approximately 1000 nm, as in the current Nova laser. The primary advantage of this approach compared to a more conventional flashlamp-pumped Nd-glass laser (such as Nova) is that the chromium in LiCAF is a much more efficient broad-band absorber than neodymium, so that coupling to the flashlamps is increased. Because of this increased efficiency, the volume of material required for an LMF driver could be reduced, thereby conceivably reducing the overall cost of the LMF facility.

One of the most important issues that must be addressed to explore this approach is the size and quality of LiCAF crystals that can be grown.

LiCAF crystals have been grown by the Czochralski process both at the Laboratory and at Allied-Signal under contract to the Laboratory since late 1988. However, these crystals are limited in their usefulness by their size (less than 5 cm diameter) and high scattering losses (typically $2 \times 10^{-2} \text{ cm}^{-1}$).¹⁴³ The LMF driver requires slabs of about 10 cm width and scattering losses $< 1 \times 10^{-3} \text{ cm}^{-1}$.

There are several aspects of the Czochralski process that may ultimately limit its ability to produce large-dimension, high-quality LiCAF which are not present in the Bridgman technique. First, Czochralski is inherently an open system that allows evaporation of volatile components unless a suitable encapsulant is used. Because LiCAF is a congruently melting material (the growing crystal has the same composition as the melt from which it is grown), changes in the melt composition produce a change in the composition of the crystal. Evaporation of AlF_3 is a significant problem in Czochralski growth of LiCAF; approximately 2-3% of the AlF_3 evaporates per day. This problem does not exist in the Bridgman process because the ampoule can be sealed to retard evaporation.

A second important limitation is that the Czochralski technique produces cylinders rather than slabs and it introduces a large degree of material waste, and adversely affects the economics of large-scale production. In the Bridgman process, a slab of near final dimension is produced directly.

Third, the Czochralski process tends to produce large thermal gradients in the crystal near the interface which are necessary for diameter and interface shape control. The large and nonlinear gradients in turn create large thermal stresses in the crystal that create dislocations which serve as nucleation sites for scattering defects, as well as stress-induced birefringence that distorts the beam. In the Bridgman process, smaller, more linear gradients can be employed that produce smaller thermal stresses. In addition, stress induced by a rigid ampoule wall can be reduced by the use of a soft ampoule liner.

Finally, scale-up of the Czochralski process to produce 10-cm diam crystals has a host of potential difficulties, including control of the melt/crystal interface and much larger thermal stresses which scale as the square of the radius. At some crystal diameter, these thermal stresses will exceed the fracture strength of the crystal.

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Figure 6-82. (a) Diagnostic for locating melt/seed interface based on difference in electrical conductivity between melt and solid. (b) The crystal/melt interface forms an acute angle with the more thermally conductive ampoule wall and wire meltback detector, causing defects to annihilate at the wall, or to each other.

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Figure 6-83. Schematic of the laser-pumped-laser microfusion facility using LiCAF as a pump laser for a neodymium-doped energy storage host material.

LiCAF Crystal Growth. We modeled the Crystalox furnace at the lower temperatures required for LiCAF growth (812°C), and characterized it with translating thermocouples. We also had to develop a different ampoule for LiCAF than we were using for YAG:Nd. This work is described in the next section. While this effort was underway, we grew small-scale LiCAF in our smaller gradient freeze furnace to examine the effects of temperature gradients, growth rates, crucible material, and diameter on the quality of the crystal.

We modified the radiation shields used for growth of YAG:Nd and added an alumina muffle to create nominal axial temperature gradients of between 8 and 16°C/cm to examine the effect of this gradient. We found that the lower gradients produced higher quality LiCAF with fewer planar dislocations. For a given nominal temperature gradient, the gradient at the seed always exceeded this value because of the design of the gradient freeze furnace. Temperature is ramped down linearly with time to grow the crystal. Hence, the larger gradients at the seed produce lower growth rates, since the growth rate is the temperature ramp rate divided by the temperature gradient. Growth rates ranged from 0.3 to 1.0 mm/hr, with the best material produced at the lower growth rates adjacent to the seed.

Temperature stability in the gradient freeze furnace does not approach that of the Crystalox furnace; the temperature generally is controlled to within 1°C as opposed to 0.1°C in the Crystalox furnace. The primary cause of this temperature variation is fluctuations in the low-conductivity water used to cool the furnace. We are procuring a closed-loop chiller this year with a temperature stability of 0.1°C to improve temperature control in the gradient freeze furnace. We observed that the crystal interface consistently broke down in the gradient freeze furnace when only a portion of the melt had crystallized, typically about 50%. This is almost certainly due to constitutional supercooling produced by a combination of temperature fluctuations and slight deviations from the congruent composition.

We also found that crystals grown at 1.25-cm-diam had consistently smaller scattering losses than 2.0-2.5-diam crystals. Losses were measured by LiCAF by insertion in a diode laser, in a hemispherical scatterometer,

and by calorimetry.¹⁴⁴ At 800 nm, the losses are primarily passive, and result from approximately 1 μm spherical particles or voids which can be easily observed by optical microscopy. In the gradient freeze crystals, the losses ranged from $8 \times 10^{-4} \text{ cm}^{-1}$ in the best crystal to about $1.5 \times 10^{-2} \text{ cm}^{-1}$ in the larger crystals. These losses are consistently lower than those measured in Czochralski-grown crystals. The origin and nature of these scattering sites are discussed further in the next section.

We applied the induction heating model developed in 1987 and refined in 1988 to the Crystalox furnace in order to examine the temperature field (gradients and interface shape) at the much lower temperatures required for LiCAF (relative to YAG:Nd). We examined the effects of graphite and crystal/melt thermal conductivity on the predicted temperature profile. A representative set of results is shown in Fig. 6-84. Figure 6-84a shows the in-phase magnetic field lines predicted by the model, from which the power deposition in conducting regions is calculated via Ohm's law. These power contour lines are shown in Fig. 6-84b. The logarithmic falloff in power dissipation from the exterior of the susceptor toward the interior is consistent with simple one-dimensional skin-depth considerations of induction heating. Finally, the temperature field produced by the power deposition profile of Fig. 6-84b is shown in Fig. 6-84c. Note that the left boundary of Fig. 6-84c is the furnace centerline, and that the wide face of the ampoule is displayed. This result predicts a flat melt/crystal interface, and a gradient of 7.5°C/cm. We have measured a melt/crystal temperature gradient of 8°C/cm using a thermocouple that translates through the melt in the Crystalox furnace, so that these results are representative of the conditions during growth.

LiCAF is a trigonal crystal, with highly anisotropic thermal expansion coefficients. LiCAF is grown along the "a" axis, normal to the "c" axis; the thermal expansion coefficients are 20×10^{-6} and $3.5 \times 10^{-6} \text{ }^\circ\text{C}^{-1}$, respectively. The graphite used for LiCAF ampoules has an expansion coefficient of $8 \times 10^{-6} \text{ }^\circ\text{C}^{-1}$, so that as the crystal cools the ampoule puts it under compression along the c-axis. In our small-scale gradient freeze crystals, we generally had to cut the ampoule away from the crystal, even though the crystal had not adhered to the ampoule wall. For the Crystalox furnace, we developed a

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Figure 6-84. Representative results of the induction heating model for the Crystalox furnace. (a) In-phase magnetic field lines. (b) Resulting power contours that show little power deposition in the magnetic field baffle plate and furnace walls, $a=10^4 \text{ W/m}^3$, $b=10^5$, and $c=10^6$. (c) Temperature field resulting from power deposition profile in (b). Isotherms spaced at 10°C .

split ampoule, held together by molybdenum nuts and bolts. The use of molybdenum, which has a lower expansion coefficient than graphite, causes the joint between the two ampoule halves to tighten as the ampoule heats up.

We've grown several crystals in this ampoule in the Crystallox furnace, with promising results. The scattering losses are comparable to the large dimension gradient freeze crystals, even though they are 10 cm wide as opposed to 2 cm in diameter. The defects are also comparable in appearance to those in the gradient freeze crystals. Once the defect composition is known, and its precursor composition is decreased, we will be able to obtain low scattering level LiCAF in the Crystallox furnace by growing the crystals from a high-purity melt of the correct composition.

The slab-shaped crystals grown in the Crystallox furnace show a tendency to crack during cool-down. Cracking may be due to traction forces induced by the crucible or second-phase material, or to internal tensile stresses created by nonlinear temperature gradients within the crystal. We have started thermal stress modeling to guide us in modifying the crucible and/or the furnace temperature field during next year.

LiCAF Defect Characterization.

Characterization of the defects in LiCAF is in its preliminary stage, but we've already made important progress. The predominance of the scattering defects in LiCAF are approximately 1 μm in diameter, and lie both in crystallographic axes and in randomly dispersed "smoky" regions. There are also occasional macroscopic inclusions ranging in size from 50–200 μm . These large inclusions are usually associated with large bubbles, and have been identified with electron microprobe as being Ca, Al, and CaCr-rich, and are presumably CaF_2 , AlF_3 , and CaCrF_5 . These compounds result from slight errors in the starting composition (noncongruent), and can be rectified by growth from the correct initial composition. We will use differential thermal analysis, and density measurements along the length of several crystals to determine the congruent composition in the next year to eliminate the large inclusions.

The much more important, smaller scattering sites have not yet been identified, but work is in progress on several fronts to determine their makeup. As was stated above, these defects are about 1 μm in diameter, and in the most heavily

concentrated areas, are present at about 100 ppm by volume. Based on their size and concentration in the crystals, a geometric scattering analysis indicates that the scattering sites are most likely either an oxide or a void (bubble), rather than a fluoride resulting from an incongruent starting composition.¹⁴³ This conclusion results from the relatively large change in the index of refraction necessary to produce the scattering levels observed. LiCAF has an index of refraction of 1.38, while a typical oxide (such as alumina) and a gas have indices of refraction of 1.8 and 1.0, respectively. By comparison, the index of refraction of calcium fluoride is 1.43.

Using optical microscopy, we have observed these defects located along the crystallographic axes and in randomly dispersed, smoky areas. The region around the crystallographically oriented defects is completely clear, suggesting migration of a species to these dislocation sites. We have also observed that the periphery of the crystals are clear of scattering sites, as are the Czochralski-grown crystals, first noted by R.C. Morris of Allied-Signal. This observation can be explained by volatilization of a gas, by high-temperature quenching of a species before it nucleates and grows, or by diffusion of a species out of the crystal during cool-down. We hope to determine the mechanism of defect formation, in addition to the identity of the defects next year.

We have been able to image a line of these defects using SEM, but have not yet been able to identify them with an electron microprobe. We've used SIMS to probe the LiCAF matrix and regions near the defects, and haven't found any differences between the two regions. However, the relatively small concentration of these defects makes it extremely difficult to be confident that the SIMS beam is sampling a defect. We are continuing our efforts using both microprobe and SIMS. We are also planning to use micro-Raman for probing regions of LiCAF with heavily concentrated scattering sites. We will also use proton NMR on similarly decorated LiCAF samples to determine if the scattering sites might contain H_2O or HF. Infrared spectroscopy indicates the presence of an H-O or H-F bond,¹⁴⁵ but we haven't been able to determine if this signal is caused by a surface layer of absorbed water or a bulk concentration of HF or H_2O . We have also initiated a series of hot isostatic press experiments to investigate the possibility that the scattering sites might be

condensed vacancies that can be annealed out of the material.

Future Plans. In FY90, we will continue to grow slabs of LiCAF in the Crystallox furnace, as well as small-scale cylinders of LiCAF in the gradient freeze furnace for materials characterization and determination of the congruent composition. We will determine the composition of the scattering sites in LiCAF, probably using one or more of the techniques described above. Once the defect composition is known, we will take the necessary steps to reduce its concentration. Because oxides are a classical contaminant of fluoride crystals, and because oxides are a likely candidate based on a geometrical scattering analysis, we are working to develop a reactive atmosphere process for LiCAF, based on the combination of HF and a fluorinating agent such as CF_4 . We will model thermal stress in the growing crystal slab, and reduce these stresses by a combination of modifying the furnace temperature distribution and the use of a soft crucible. A relatively new commercial control software package (LabView) will be interfaced to the Crystallox furnace to allow advanced control algorithms, such as adjusting furnace power based on cooling water temperature oscillations, to be implemented. Finally, we have begun discussions with crystal growth and laser glass companies regarding transfer of this technology in the coming year. Several of these companies have expressed interest in our crystal growth technology, and we expect to successfully transfer the technology to at least one or two of these companies next year. By transferring this technology to the private sector, we will help to ensure the availability of large, high-quality crystal materials for future Laboratory programs.

Authors: L. J. Atherton and J. F. Cooper

Major Contributors: D. H. Roberts, J. J. DeYoreo, D. R. Burns, G. G. Pollock, and A. A. Tesar

6.6. Non-Linear Optical Processes

6.6.1. Introduction

In 1988 we conducted an experimental series on the Nova laser to study the onset of transverse stimulated Brillouin scattering (TSBS) in large optical components, as described in Sec. 4.2.3 of this Report. TSBS is most important for high

fluence, large transverse beam dimension, long laser pulse length, and short wavelength. Nova is large enough that TSBS can become an important energy loss for pulses at 350 nm longer than about 2 ns. The Nova experiments demonstrate that TSBS losses can be reduced to a negligible value for SBS gains typical of fused silica by a small increase in laser bandwidth, and that this increase in bandwidth is too small to affect the efficiency of frequency conversion to 350 nm. A detailed description of these experiments and analyses was published this year.¹⁴⁶

Future larger glass fusion lasers will operate with longer pulses and higher fluences which put them well into the region where TSBS losses must be controlled. Reliable values are important for the peak Brillouin gain and acoustic decay time for the optical materials which we might use in these lasers. In 1989 we established a collaboration with SRI International to modify a stimulated scattering gain apparatus of theirs to measure backward and transverse SBS gains directly. At LLNL, we also studied backward SBS growth from noise in a few materials as an independent check on these results.

6.6.2. Gain for Stimulated Brillouin Scattering in Optical Components

The experiments at SRI International use a CW argon laser-pumped single-mode ring dye laser as a probe to measure the gain induced by pumping with a doubled single-mode Nd:YAG SS laser at 532 nm. The pulse length of the pump laser can be varied from about 10–40 ns, so some features of transiency can be explored in the low-gain regime. The initial results of these experiments will be presented as an invited paper at the CLEO/IQEC-90 conference and will also be published shortly.¹⁴⁷

Table 6-23 has a summary of preliminary results from the SRI measurements for several optical materials of interest, with backward SBS gains expressed relative to that of fused silica. The peak Brillouin gain γ is the relevant gain scaling parameter for pulses longer compared to the acoustic decay time τ , and the integrated line strength γ/τ is the relevant scaling parameter for pulses shorter than τ .¹⁴⁸ One can see that the acoustic decay time is shorter in the amorphous materials (glass and liquid) than in crystals, as expected. The data also illustrate the failure of the Lorentz-Lorenz

Table 6-23. Preliminary results on the relative SBS gain and acoustic lifetime of optical materials at 532 nm.

	steady-state gain γ (relative to fused silica)		acoustic lifetime τ in ns at 532 nm, backward scattering		transient gain γ/τ (relative to fused silica)	
	SRI	LLNL	SRI	LLNL	SRI	LLNL
fused silica	(1)	(1)	1.0	0.95	(1)	(1)
CaF ₂	0.3		3.5		0.08	
KDP**	1		1.5		0.67	<0.85*
LAP**	10	3.8	1.7	1.4	6	2.7
GGG	3		8		0.38	
methanol	4.3		0.46		9.3	
toluene	3		0.14		22	

Notes:

*Damaged at focus with this fluence, below SBS threshold

**These crystals are highly anisotropic, so SBS gain may have a significant dependence on the direction of propagation and polarization. These parameters were arbitrary in these preliminary measurements.

KDP = potassium dihydrogen phosphate frequency converter

LAP = levo-arginine phosphate frequency converter

GGG = gadolinium gallium garnet laser crystal

approximation to the calculation of polarizability in solids, though this is not a new observation.¹⁴⁹ As an example, the Lorentz-Lorenz model predicts an integrated line strength for gadolinium gallium garnet (GGG) about 6× larger than fused silica because of its high index of refraction (1.94). The measurements in the table show that it is less than half that of fused silica. It is for this reason that we must have direct Brillouin gain or Pockels tensor measurements such as those we are conducting here to evaluate Brillouin gains rather than estimating them using simple models.

Table 6-23 also includes relative gains and acoustic relaxation times inferred from backward SBS threshold measurements conducted at 1053 nm at LLNL over a similar range of pulse lengths. The acoustic decay times are inferred from the dependence of threshold on pulse length. They have been scaled from their measured values to those appropriate for 532 nm assuming classical acoustic damping,¹⁴⁸ which gives $\tau \sim \lambda$.¹⁴⁷ These threshold gain measurements are consistent with the SRI results.

The SRI apparatus can be calibrated to give absolute gain values if parameters such as the laser intensity profile are well-known. The absolute gain values derived in this way are lower than the values predicted from Pockels tensor measurements in the literature for those materials (fused silica, toluene, methanol) for which such measurements are available. Our values range from about 73% of the value predicted for fused silica to 85% of that predicted for toluene. The experiment should be able to give higher accuracy. We shall study this further in 1990 to determine whether the discrepancy is real or is due to difficulties with intensity measurements, corrections for transient scattering effects, or other experimental artifacts.

The SRI apparatus can be used to measure transverse SBS gain as well as backward gain, and we have demonstrated this using fused silica. The materials for which an actual transverse measurement is most important are the highly anisotropic frequency converter crystals such as KDP and LAP, since the dependence of SBS gain on propagation direction and polarization may be rather complex and require a detailed knowledge of the Pockels tensor. In 1990 we shall make transverse measurements in these materials.

We conclude from this year's data that the SBS gain in KDP is comparable to that in fused silica, so it will set similar limits to large fusion lasers. The gain in LAP is significantly larger than fused silica, so SBS in that crystal will be a more serious limit if it is used in such lasers. The gain in materials such as calcium fluoride is so low that SBS will be less significant than other limits, such as optical damage at high fluence.

Authors: G. W. Faris, and A. P. Hickman (SRI); J. R. Murray, and D. Milam (LLNL)

6.6.3. Phase Conjugation of Short Pulses

Phase conjugation is a promising technique for correcting phase error, locking together the phases of beams from different apertures, and automatically pointing beams. Conjugation by stimulated Brillouin scattering appears to be the technique best suited for phase reversal of a pulse with energy of about 1 J. SBS arises from the interaction of acoustic waves in a transparent medium with an intense laser beam. This interaction creates a density gradient capable of efficiently reflecting a beam with phase fronts that are reversed relative to those in the input beam. The simplest SBS conjugator consists of a volume of the scattering medium and a lens to focus the input beam into this medium. Because an intense beam is required to excite SBS, laser-induced breakdown and self-focusing may interfere with conjugation. The possibility of such interference is increased, and other problems related to the transiency of the SBS process are encountered, when the input pulse is of short duration. Therefore, selection for a given application of the optimum medium requires identification of the mechanisms that limit the fidelity of conjugation.

We performed several experiments aimed at separating the deleterious effects of self-focusing from those caused by transiency. One intent was to test the validity of a model¹⁵⁰ that allows calculation of the phase retardation as a function of pulse duration that is nonlinearly induced when a pulse with rectangular waveform and energy equal to the SLS threshold, E_t , is focused through a conjugator. In this calculation the input beam was assumed to have phase fronts distorted only by smooth aberrations. For an input having wavelength $\lambda = 1053$ nm and energy equal to the SBS threshold, Fig. 6-86 gives calculated values for the rapid (electronic and Kerr) contributions to

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Figure 6-86. Nonlinearly induced phase shift on an input pulse with energy equal to the SBS threshold as a function of pulse duration.

the nonlinearly induced retardation in a liquid CS₂, a solid SiO₂, and a compressed gas N₂ at a pressure of 60 atm. The predicted threshold-level retardation is for a 6-ns pulse that ranges from 0.01λ for compressed N₂ to 0.5λ for CS₂. Note that a conjugator is usually used at input energies well above threshold, and that the actual value of the induced phase shift will be that corresponding to the highest intensity reached prior to creation of a reflective grating.

A model that fully accounts for self-focusing in the conjugating medium is not presently available, although there have been several studies of the use of conjugation to reverse phase shifts nonlinearly induced in components external to a conjugator.¹⁵¹⁻¹⁵³ Lacking a model, we can still identify some instances in which the induced retardation is likely to be innocuous. Experience with fusion lasers suggest that induced phase shifts of magnitude less than 0.5λ are not likely to cause catastrophic alteration of the shape of a laser beam as it passes through the conjugating medium.¹⁵⁴ Such small phase shifts should be incorporated into the grating and reversed, particularly when the waveform of the pulse is rectangular and the magnitude of the induced phase shift is constant during grating formation. However, once reflectance is initiated, the simultaneous presence of both the input and the conjugate wave will cause an increase in the induced phase shift of the input wave, so the allowable value of initial input phase shift might be less than 0.5λ . We may additionally anticipate that a phase conjugation will be significantly impaired by the presence during grating formation of an induced phase shift of magnitude λ which is adequately large to cause beam trapping due to either small-scale or large-scale self-focusing.

Using the value 0.5λ as a measure of allowable phase shift, and the data in Fig. 6-86, perturbation of the input beam is expected for a threshold-level pulse of short duration in a CS₂ conjugator, whereas for a pulse of long duration such perturbation is expected only at input intensities well above the SBS threshold. To test this prediction, pulses with rectangular waveform and duration of 6, 13 or 70 ns, in a beam with Gaussian spatial intensity distribution and less than 0.5 waves of phase nonuniformity, were focused through a conjugator consisting of a lens with focal length of 50 cm and a 20-cm thickness of CS₂ that was centered on the beam waist. As the energy in the 6-ns pulse was increased above

$0.35 E_t$, the shape of the beam transmitted through the conjugator exhibited a power-dependent evolution from its Gaussian shape into a beam with two intensity maxima.

Interestingly, for input energies just above the SBS threshold, $1.0-1.3E_t$, the beam reflected by stimulated scattering had Gaussian shape, which suggests that a reversible phase perturbation was responsible for the intensity modulation observable on the transmitted beam at these input levels. At input energies greater than $1.3 E_t$, both the reflected beam and the transmitted beam contained two or more intensity maxima. Observations made with a 6-ns pulse at a particular input power could be reproduced using a pulse of equal input power and duration of either 13 or 70 ns. These results are consistent with the assumption that self-focusing in the CS₂ was responsible for the beam perturbation. An estimate of the maximum allowable phase shift was obtained by noting that for pulses with duration of either 6 or 70 ns, perturbation of the reflected beam occurred when the calculated phase shift of the input beam was $0.6-0.7 \lambda$. Finally, even the limited ability of CS₂ to serve as a conjugator for a 6-ns pulse may be regarded as surprising, since the critical power for self-focusing, 29 kW,¹⁵⁵ is comparable to the input power at SBS threshold, 33 kW (0.2 mJ).

If, as indicated by the experiments in CS₂, an induced phase shift of $0.6-0.7 \lambda$ is tolerable, data in Fig. 6-86 suggest that self-focusing should not interfere with conjugation in a compressed gas of pulses with duration greater than 6 ns and energy less than $60-70E_t$. However, distortion of the reflected beam was observed at input energies of $10-20E_t$ when either N₂ or CH₄ at pressures of 60-120 atm was used to conjugate pulses with durations of 6-13 ns in a beam with less than 0.5 waves of phase nonuniformity. When 3 waves of astigmatic phase error were added to the input beam, allowable input energies for a CH₄ conjugator at 60 atm were reduced to $2E_t$ for 6 ns, and to about $5E_t$ for 70 ns. Therefore, while the quality of phase conjugation in a compressed gas had an observable dependence on either the intensity or the duration of the input pulse, the results are not explained by calculated values for the induced phase shift on the input beam.

A second problem, transiency, is identified in the literature as a possible complication in the conjugation of a pulse with duration, t_p ,

comparable to the acoustic lifetime, t_b , of the scattering medium.¹⁵⁶ Stimulated scattering is in the steady state when t_p is greater than $15t_b$ and very transient when t_p is less than $5t_b$, although the solutions for the transient case hold for pulse durations intermediate to these limits.

Improvement in the quality of phase reversal for a short pulse has been obtained by using a priming pulse with duration t_p much greater than t_a and with relatively low intensity to create the grating prior to arrival of an intense short pulse.^{157,158} In these experiments, the quality of phase conjugation was judged by comparing the far field distribution of the conjugate beam to the diffraction limit.

We have investigated priming as a technique for improving the spatial uniformity in the near field of the phase conjugate beam. Using CH_4 at 60 atm for which $t_b = 14$ ns, we compared conjugation of a 6-ns rectangular pulse to that for a composite pulse consisting of a Gaussian waveform with duration of 250 ns on which an intense 6-ns segment positioned about 50 ns was superimposed after the temporal peak of the Gaussian. With the 6-ns pulse by itself, and with 3 waves of added astigmatic input aberration, the spatial distribution of the conjugate beam was distorted when the input energy exceeded 60 mJ, which is about twice threshold. With the Gaussian by itself, the SBS threshold was 80 mJ, and at threshold the induced scattering occurred during the interval of maximum power, indicating that the scattering induced by this pulse was nearly in the steady state. The leading edge of this 250-ns, 80-mJ Gaussian was used to prime the conjugator to a reflectance of less than 1% prior to arrival of the 6-ns segment containing an additional 130-mJ. With priming, the reflectance for the 6-ns, 130-mJ pulse increased from 25 to 75%, accurate reversal of the astigmatic phase error was obtained, and the reflected beam was free of amplitude modulation.

This result suggests that transiency is the mechanism that complicates conjugation of a short pulse in a compressed gas. Since the primed reflectance was small, the intense 6-ns segment of the pulse penetrated to the beam focus and incurred essentially the same nonlinear phase shift as in the nonprimed experiment, and additional phase shifts should have been induced by counterpropagation of the intense conjugate wave.

Therefore, conjugation of a short pulse is complicated if the medium has a large nonlinear refractive index, or if the scattering is highly transient. This raises the possibility that for conjugation of short pulses the optimum medium might be a liquid with short acoustic lifetime and only moderately large nonlinear refractive index, CCl_4 for example, that has been rigorously cleaned to deter laser-induced breakdown.

Author: David Milam

6.7. Optical Material Properties

6.7.1. Introduction

The Optical Materials Research Program is supported by the Office of Basic Energy Sciences (OBES) of the Department of Energy through its Division of Materials Science. This program involves basic research on topics and materials of importance to the Laser Program. Three research areas provide the framework for our research effort:

- Spectroscopy of laser ions in crystals and glasses.
- Measurement and modeling of the magnitude and dispersion of the nonlinear refractive indices of optical materials.
- Fundamental mechanisms of optical surface damage in optical materials.

Spectroscopy of Laser Ions. The research on spectroscopy of laser ions emphasized the change in polarizability of Cr^{3+} and Nd^{3+} ions when they are excited into their metastable upper laser level. This polarizability change, $\Delta\alpha$ was measured using four-wave mixing. We found that $\Delta\alpha$ was strongly dependent on the identity of the anion neighbors of both the Nd^{3+} and Cr^{3+} ions. This result can be understood as resulting from the influence of the neighboring anions on the energy and character of the opposite-parity excited states that determine the polarizability of these ions. These changes in polarizability produce an "electronic lens" in a pumped laser material. We find that the magnitude of this lens can be comparable with the thermal lens in single-shot laser applications.

We investigated several aspects of the behavior of two Nd^{3+} doped energy storage

media, SrF_2 and LaF_3 , pumped to high inversion levels by a long-pulse pump laser. This work was motivated by our interest in exploring laser-pumped lasers for laser fusion applications. Cross relaxation of excited Nd^{3+} ions (also known as Auger decay) and excited state absorption of laser pump light are two potentially important loss mechanisms. We found that excited state absorption of pump light causes insignificant loss off laser pump energy. On the other hand, we found that the Auger decay of excited Nd^{3+} ions can deplete about 10 percent of the pumped population in LaF_3 at inversion densities of about 10^{19} cm^{-3} (equivalent to energy storage densities of several joules per cubic centimeter).

In order to understand the laser pumping processes of the Nd^{3+} ion in various hosts it is important to know the lifetimes of the excited states involved in pumping, excited state absorption, and Auger decay. We measured the lifetimes of the important spin-orbit multiplets in several oxide and halide crystals. These lifetimes, which ranged from subnanosecond to many microseconds in duration, were much shorter (by up to two orders of magnitude) than the lifetimes predicted by the commonly used "energy gap law." In contrast with these transitions between different spin-orbit multiplets, nonradiative transitions among the sublevels of crystal field split spin-orbit multiplets of rare earth ions occurs on a time scale of a few picoseconds or less. We measured such decay rates for the $^4\text{G}_{5/2} - ^2\text{G}_{7/2}$ multiplet of Nd^{3+} in several glasses using an accumulated photon echo experiment. The transition rates increased with increasing energy of a crystal field state relative to the lowest state in the multiplet because of the increasing density of vibrational states in the host glass. We observed more subtle variations of the lifetimes within the inhomogeneously broadened envelopes of the crystal field levels. These variations are believed to result from a correlation between crystal field splitting and electron-phonon coupling strengths for different ion sites in the crystal.

Nonlinear Refractive Index. We measured the nonlinear refractive indexes of many optical crystals at the second and third harmonic wavelengths, 532 nm and 355 nm, of the Nd^{3+} laser. For wavelengths longer than the two-photon interband absorption threshold, we found that the wavelength dependence of the nonlinear index could be modeled by a simple

dispersion formula similar to the Sellmeier formula for the linear refractive index. The only information required to estimate the nonlinear index of a material at these wavelengths is the nonlinear index at one wavelength and the Sellmeier parameter for the linear index. At wavelengths below the two photon absorption threshold, the nonlinear indexes of several materials are larger by about a factor of two or so compared with their values at 1.06 μm .

Optical Surface Damage. We developed a new method, using pairs of picosecond laser pulses, for investigating the mechanisms for laser-induced surface modification and optical surface damage. This experiment can be used to distinguish between linear and nonlinear laser-surface interactions and to measure the lifetimes of the excitations (electronic or thermal) that result from these interactions. We found that the laser-surface interaction is caused by linear absorption for crystalline ZnS , and by a nonlinear process for BK-7 glass. A decay time of about 5 ns was observed for the excitation produced on ZnS surfaces; the "excitation" is most likely to be localized heating at the ZnS surface.

Author: Lloyd Chase

6.7.2. Dispersion of the Nonlinear Refractive Index of Optical Materials

The nonlinear refractive index n_2 strongly influences the propagation of intense laser beams through optical materials. It causes self focusing and amplification of spatial intensity fluctuations of an incident beam. The nonlinear index is defined by

$$n = n_0 + n_2 \langle E^2 \rangle, \quad (1)$$

where n is the refractive index in the presence of an optical electric field E and n_0 is the linear refractive index. The brackets imply a time average over the optical period. Our research is aimed at gaining an understanding of the magnitude and wavelength dependence (dispersion) of n_2 for optical materials. We previously measured values of n_2 for many glasses and crystals at the 1.06- μm fundamental wavelength of the Nd laser and investigated the dependence of n_2 on the composition and structure of optical materials.^{159,160} In 1989 we measured n_2 at the second and third harmonic

wavelengths, 532 nm and 355 nm, for many of the same materials. The measurements have enabled us to investigate the suitability of empirical models for predicting the dispersion of n_2 .

The dispersion of the linear refractive index is often modeled by using a Sellmeier dispersion formula, which, in its simplest form, predicts a wavelength dependence at wavelengths in the region of transparency given by

$$n_0^2 \sim 1 + \frac{A}{1 - \left(\frac{\lambda_0}{\lambda}\right)^2} \quad (2)$$

where hc/λ_0 is the average energy of excited states of the material that are important in determining its polarizability, and the parameter $A \sim 1$ for most materials. The form of this Sellmeier formula is suggested by the perturbation theory expressions for the linear index. In a similar manner, it is possible to deduce (from the perturbation theory expression for the third-order polarizability of a medium) an approximate formula for the wavelength dependence of n_2 when $\lambda \gg \lambda_0$. This formula is

$$n^2(\nu) = n_2(0) \left\{ 1 + \left(5 + \frac{5}{6k} \right) \nu^2 + \frac{A}{2(1+A)} \nu^2 \right\} \quad (3)$$

where $\nu = \lambda/\lambda_0$, λ_0 is the Sellmeier characteristic wavelength, A is the Sellmeier constant [(both are assumed to be the same as those used in Eq. (2)], and k is a parameter (~ 0.5) that depends on the detailed form of the electronic wave functions of the medium. Equation (3) predicts that the wavelength dependent correction to n_2 increases with decreasing wavelength as λ^{-2} . The first term (proportional to ν^2) on the right of Eq. (3) is the wavelength dependence of the third order polarizability of the medium. The second term is the wavelength dependence of the Lorentz local field correction; $f = (n^2 + 2)/3$. This is much smaller than the first term for most low-index materials. n_2 is proportional to f^4 .

We measured the ratio of the nonlinear refractive index of many crystals relative to that of fused silica at the 532- and 355-nm second and third harmonic wavelengths of the Nd:YAG laser. The measurements are done using nearly

degenerate mixing, as described in previous publications.^{159,160} The points plotted as squares in Fig. 6-87 are the logarithms of these n_2 ratios for some of the crystals investigated at the three Nd laser wavelengths. The dashed lines in the figure connect the squares in order to identify the materials at each wavelength. The solid lines are the predicted wavelength dependencies obtained using Eq. (2) with the published Sellmeier parameters, A and λ_0 , for each material and for the fused silica reference material. In most cases, the agreement is quite good. Some of the materials for which the dispersion is much larger relative to SiO_2 than predicted by Eq. (3) have strong two-photon absorption at either 355 or 532 nm, or both. This is true of the titanates, ZnO, and CdS. Eq. 3 is not expected to be a good approximation for these materials at these wavelengths. It is clear from Fig. 6-87 that the dispersion of n_2 in many low index, wide bandgap insulators is nearly identical to that of SiO_2 .

In order to obtain the absolute value of n_2 at 355 and 532 nm for all the crystals measured relative to fused silica, we used the measured Raman scattering cross section of liquid benzene as a nonlinear reference standard. This cross section has been measured reproducibly by several investigators at many wavelengths from 350 to 700 nm.^{161,162} The n_2 value for several of the crystals was determined by fitting the frequency dependence of the coherent anti-Stokes Raman scattering (CARS) signal of composite samples, measured as a function of the frequency difference between a Nd:YAG harmonic and a dye laser. These composite samples were cells composed of the material to be calibrated and containing a thin layer (~ 10 – $20 \mu\text{m}$) of benzene. The setup is shown in Fig. 6-88 and a typical CARS spectrum of a LiF cell containing benzene is shown in Fig. 6-89 for a laser wavelength of 532 nm. The solid line in Fig. 6-89 is a fit to the spectrum using the Raman cross section for the 992 cm^{-1} vibration of benzene and a fitted value of the third order susceptibility, $\chi^{(3)}$, of the LiF window, which produces the nonresonant background scattering level (the "0" level in Fig. 6-89). A nonlinear index of $0.3 \times 10^{-13} \text{ esu}$ is obtained from the fitted value of $\chi^{(3)}$. We obtained similar calibration data for the three fluoride materials LiF, NaF, and BaF_2 at both 532 and 355 nm.

The n_2 values for LiF at the three Nd laser wavelengths are plotted in Fig. 6-90. The 1064-

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Figure 6-87. Ratios of nonlinear refractive indices of some crystals relative to fused silica at Nd:YAG laser harmonics.

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Figure 6-88. Cell used to calibrate the n_2 for LiF using the C_6H_6 Raman cross section.

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Figure 6-89. CARS spectrum of the cell in Fig. 6-88. The smooth, solid curve is a fit to the data yielding $n_2 = 0.3 \times 10^{-13}$ esu for LiF at 532 nm.

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Figure 6-90. Nonlinear refractive index of LiF at three Nd laser wavelengths. Solid and dashed lines are plots of Eq. (3) with different values of $n_2^{(0)}$.

nm point was calibrated by time-resolved interferometry measurements,¹⁶⁰ and the points at 532 and 355 nm are calibrated from the benzene Raman cross section. The solid curve is a fit to the functional form of Eq. (3). Only $n_2^{(0)}$ is varied to fit the data. The fitted curve falls barely within the error bars for the three points. Although this could be due to errors in our measurements, it is also quite possible that the two different nonlinear reference standards used to calibrate the measurements are themselves uncertain. In fact, the two points calibrated from the benzene standard are mutually self consistent. The dashed curve, which is the solid curve scaled up by about 10%, passes through the 532- and 355-nm points, but it predicts an n_2 value that is 30% higher than the measured point at 1064 nm. An uncertainty of 30% in a nonlinear susceptibility measurement is not unlikely. In order to resolve this uncertainty, and calibrate the relative measurements in Fig. 6-87 with greater certainty, it is important to get absolute calibration from a third type of measurement at two or more of the Nd harmonics.

In summary, our measurements show that the nonlinear refractive indices of many optical crystals follow a simple dispersion law at wavelengths much longer than one and two photon absorption bands. A good estimate of the wavelength dependence of n_2 at long wavelengths can be obtained for each material from the Sellmeier parameters fitted to the dispersion of its linear refractive index.

Author: L. L. Chase

Major Contributors: R. Adair and S. A. Payne

6.7.3. Investigation of Optical Damage Mechanisms Using Picosecond Pulses

Laser-induced damage to optical elements is a technologically critical area for high-power laser research. The interaction of intense laser beams with the surfaces of transparent optical materials can involve linear and nonlinear absorption caused by surface states, defects, and impurities. Direct measurement of these interactions and their effects on the surface is very difficult because of the small volume of material affected and lack of reliable methods for probing the surfaces of insulators. It is known, however, that laser excitation of many transparent materials results in the ejection of charged and neutral surface constituents and

impurities.¹⁶³ In the present work we have used these emitted particles to detect accumulated optical surface damage on ZnS crystals and BK-7 glass caused by excitation of the surface by pairs of picosecond pulses.¹⁶⁴ From the ratio of the damage thresholds obtained with double-pulse and single-pulse excitation, it is possible to determine whether the laser-surface interaction is linear or nonlinear in the laser intensity and to measure the lifetimes of excitations (electronic, thermal, etc.) produced by the laser irradiation.

In our experiments, samples were irradiated in a vacuum system at a pressure of 10^{-9} Torr using a 580-nm dye laser operating at 10 Hz with a pulse width of about 1.5 ps and a maximum pulse energy of 300 μ J. The laser was focused to a 150- μ m-spot diam, yielding fluences up to 1.3 J/cm². A quadrupole mass spectrometer (QMS) detected neutrals emitted following the laser excitation. We monitored the Zn signal (64–68 amu) for ZnS and the Na signal (23 amu) for BK-7. The output of the QMS was sampled by a gated integrator for about 1 ms after each laser pulse or pulse pair, and the signal was plotted continuously on a chart recorder.

The neutral emission yield data obtained for a mechanically polished ZnS crystal using "single-pulse" excitation at 10 Hz is shown in Fig. 6-91. ("Single-pulse" means that one pulse arrives every 100 ms.) A different site and laser fluence were used for each trace that records emission as a function of the number of laser shots. At low fluences, as in Fig. 6-91(a), the neutral emission from the surface is very weak and is not detectable at the QMS sensitivities used in this work. This is consistent with the behavior of neutral emission yields observed in previous experiments with single shot excitation, where it was found that the neutral emission yield decreases rapidly with successive laser shots.¹⁶³ At a well defined threshold fluence, 0.4 J/cm² in Fig. 6-91(b), a large flux of neutrals is observed after several hundred laser pulses, and the number of pulses required to initiate the emission decreases rapidly with increasing fluence, as is evident from traces Fig. 6-91(b)–(e). This data sequence was repeated several times, and it was observed that the threshold fluence was repeatable to within 20–30%. Thus, for this case *cf.* excitation by a single pulse every 100 ms, we determined what we defined as the threshold fluence $F_t(1)$ by gradually increasing the fluence and noting the level at which neutral emission is observed

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Figure 6-91. Neutral emission yield for picosecond pulse excitation of ZnS at a 10-Hz repetition rate. Approximate fluences (J/cm^2) are: (a) 0.38, (b) 0.4, (c) 0.43, (d) 0.56, (e) 1.35. The inset shows the number of pulses required to cause damage as a function of fluence.

within the first 1000 laser pulses. The inset in Fig. 6-91 shows this to be a reasonable definition, since the damage fluence becomes relatively independent of the number of pulses near this threshold. We found that visible damage, in the form of a large roughened area comparable in size to the illuminated region of the sample, occurred in the first pulse or two after the neutral emission signal appeared. The most important point to emphasize is that, for any fluence above $F_t(1)$, there is apparently a cumulative modification of the surface on each laser pulse; this modification causes a catastrophic failure after a number of pulses. The requisite number of pulses decreases rapidly with increasing fluence until, at some fluence above 1 J/cm^2 , the damage occurs on the first shot.

In order to investigate the cause of the cumulative surface modification, we repeated the experiment of Fig. 6-91 using pairs of picosecond pulses at a repetition rate of 10 Hz. The pulses in a pair had equal fluences and were separated by a variable time delay, τ , of up to 8 ns. From data similar to that in Fig. 6-91 we determined a threshold fluence $F_t(2, \tau)$ above which cumulative optical damage occurs after at most 500 pulse pairs (the total number of pulses was the same as for the single-pulse case). We define $F_t(2, \tau)$ as the sum of the fluences of the two pulses at the accumulated damage threshold. For each value of τ we calculated the ratio $R(\tau) = F_t(2, \tau)/F_t(1)$, which is plotted in Fig. 6-92 as a function of τ . There are three distinct time intervals in Fig. 6-92.

Firstly, for $\tau < 2 \text{ ps}$, the pulses overlap temporally and form interference fringes that double the fluence in the maxima. Since the peak fluence in these maxima is twice as large as for a single pulse of the same average fluence, we expect a value of $R = 0.5$, which is close to the observed value. This result gives no new information since it results merely from a redistribution of the laser energy on the surface. Secondly, for $2 \text{ ps} < \tau < 1 \text{ ns}$, the measured value of R is 1 for over two orders of magnitude in τ . Thirdly, for $\tau > 1 \text{ ns}$, R increases toward the value of 2 that would be expected if each pulse of the pair interacts independently with the surface; each pulse must then exceed the single pulse accumulation threshold $F_t(1)$.

In order to qualitatively interpret the results in Fig. 6-92, and to demonstrate the general utility of these measurements, we present a

simple model based on these reasonable assumptions: 1) the cumulative surface modification results from an absorption mechanism proportional to I^n , where I is the total laser flux on the surface; 2) this absorption produces some excitation that decays exponentially in time with a time constant τ_c ; 3) the peak instantaneous value of the excitation must exceed a threshold value in order to cause surface modification. This is the simplest model that allows an interpretation of the data based on absorption varying as a power law in the intensity. Since the threshold value of the excitation must be the same with either single- or double-pulse excitation, we can write

$$\left(\frac{F_t(2, \tau)^n}{2} \right) (1 + e^{-\tau/\tau_c}) = [F_t(1)]^n, \quad (1)$$

where the two terms on the left are proportional to the value of the excitation just after the second pulse, i.e., when the excitation produced by the first pulse has decreased by the factor $\exp(-\tau/\tau_c)$.

The ratio R obtained from Eq. (1) is

$$R(\tau) = 2 \left(1 + e^{-\tau/\tau_c} \right)^{-1/n} \quad (2)$$

Note that this model does not deal with the physically uninteresting redistribution of intensity caused by interference at very short time delays where pulse overlap occurs. Although a fringe-visibility factor could be added to Eq. (2), neither the accuracy of the data, nor the utility of such a procedure justifies the added complexity. At short time delays ($\tau \ll \tau_c$), Eq. (2) predicts, $R = 2^{1-1/n}$. The value of $R = 1$ observed for ZnS, for values of τ ranging from the $\sim 2 \text{ ps}$ "interference limit" up to $\sim 1 \text{ ns}$, implies that linear absorption is the active mechanism; this leads to the conclusion that only the total energy of the pulse pair determines the threshold for surface modification. A best fit of Eq. (2) to the data of Fig. 6-92 is given by the dotted curve, which is calculated for $\tau_c = 5 \text{ ns}$. This value of τ_c implies that the optical damage threshold fluence for long pulses would be independent of laser pulse duration up to 5 ns.

The photon energy of the laser, 2.14 eV, is far below the bandgap energy of ZnS (3.54 eV), but it is in the region of very strong two-photon

Figure 6-92. Ratio of threshold fluences for pulse-pair and single-pulse excitation as a function of pulse-pair separation. The dotted line is calculated from Eq. 2 with $\tau_e = 5$ ns.

absorption. In fact, the 1-mm-thick samples strongly absorb the picosecond pulses, which have peak irradiances of about 400 GW/cm² at the accumulated damage threshold. Nevertheless, our results show that the laser interaction with the surface proceeds through a linear absorption mechanism that must involve defects or surface states. The "excitation" with a lifetime of 5 ns could be either local heating or free carriers. Temperature excursions would not, of course, decay exponentially, but the accuracy of the data does not justify an attempt to fit it to a thermal diffusion model. Further discussion of the damage mechanism can be found elsewhere.^{163,165}

Experiments were also performed on well polished BK-7 surfaces, and the ratio R is plotted versus time delay in Fig. 6-93. Although there is a much larger statistical fluctuation of the data points the value of R consistently falls between 1.5 and 2 with an average value of 1.75 for short time delays. This result shows that linear absorption is not the cause of the cumulative surface modification for BK-7. The data could be interpreted as a five-photon absorption process that creates an excitation with a lifetime of about 2 ns, as shown by the dotted curve in Fig. 6-93. Considering the scatter of the data, however, we can conclude only that the laser-surface interaction is highly nonlinear. It is conceivable that avalanche ionization may be the cause of surface modification, and it is unlikely that all of the assumptions of the simple excitation model discussed above would apply in this case.

In conclusion, these experiments have shown that pulse-pair excitation of neutral surface emission can be used to investigate laser-surface interaction mechanisms and to measure the lifetimes of the excitations caused by these interactions. Similar experiments at other pulse durations would also be of interest to investigate the dependence of the interactions on peak irradiance and fluence. These experiments are in progress.

Author: L. L. Chase

Major Contributors: H. W. H. Lee and R. S. Hughes

6.7.4. Dynamics of Laser-Pumped Nd³⁺ Laser Media at High Excitation Density

Laser-pumped solid state lasers offer the potential for improved beam quality, efficiency, and energy storage density. These performance advantages may be important for inertial confinement fusion laser drivers. Two concepts for large fusion drivers use either a tunable solid state laser or semiconductor diode laser arrays to pump long-lifetime rare earth energy storage media to high-energy storage density (~1 J/cm³). For high efficiency, losses in the pumping process and depletion of the pumped population must be avoided. Two potential loss mechanisms that can occur in laser-pumped rare earth doped storage media are excited state absorption (ESA) of pump photons and energy transfer upconversion of the excited ions. We investigated the importance of these loss processes in SrF₂:Nd³⁺ and LaF₃:Nd³⁺, which are two energy storage media selected for possible use in a fusion driver pumped by Cr:LiCaAlF₆ (Cr:LiCAF).

The energy level diagram in Fig. 6-94 illustrates the ESA of pump radiation for Nd³⁺. The principal laser pumping bands are the $^4I_{9/2} \rightarrow ^4F_{7/2}$, $^4I_{9/2} \rightarrow ^4F_{5/2}$, and $^4I_{9/2} \rightarrow ^4F_{3/2}$. As shown in Fig. 6-94, there is a resonance at ~740 nm that has been observed to cause excited state absorption of pump radiation of this wavelength in Nd³⁺ (Ref. 166). The loss due to ESA manifests itself by increasing the pump flux necessary to achieve a given excited state density. The important ESA transition at ~740 nm is the $^4F_{5/2} + ^2H_{9/2} \rightarrow ^2P_{3/2}$. The critical parameters that determine the magnitude of the loss due to ESA are the $^2H_{9/2} \rightarrow ^2P_{3/2}$ cross section and the lifetime of the $^4F_{5/2} + ^2H_{9/2}$ intermediate state. The optical rate must exceed the decay rate of the intermediate state for ESA to be significant,

$$\frac{\sigma_{\text{ESA}} I}{h\nu} \geq 1/\tau \quad (1)$$

where σ_{ESA} is the cross section for the $^2H_{9/2} \rightarrow ^2P_{3/2}$ transition, I is the pump beam intensity, and $1/\tau$ is the lifetime of the upper pump level. Using the Judd-Ofelt theory¹⁶⁷ to estimate oscillator strength for $^2H_{9/2} \rightarrow ^2P_{3/2}$, the ratio of the ESA cross section to the pump cross section is ~0.3. Chase and Smith¹⁶⁸ found that the lifetime of $^2H_{9/2} + ^4F_{5/2}$ was less than 5 ns for

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Figure 6-93. The ratio $R(\tau)$ for borosilicate glass.

Figure 6-94. The energy levels of Nd^{3+} show an excited state absorption for red pumping at $\approx 740 \text{ nm}$.

Nd:LaF₃ and Nd:SrF₂. For a free-running Cr³⁺:LiCaF laser, the pump laser pulse is a series of relaxation oscillations. Assuming pumping at 10 J/cm² with 300-ns wide spikes, the peak intensity is ~500 kW/cm². For these conditions the optical rate for ESA is

$$\sigma_{\text{ESA}} I / h\nu \approx 2 \times 10^4 \text{ sec}^{-1}$$

whereas $1/\tau \geq 2 \times 10^8 \text{ sec}^{-1}$. Because of the short lifetime of the intermediate state, ESA should be a negligible effect for laser pumping of Nd³⁺:LaF₃ and Nd³⁺:SrF₂. This conclusion is supported by our measurements which showed that the fractional ESA of pump light is about 0.1% of the total pump fluence, and it is caused by the thermal equilibrium population in ²H_{9/2} + ⁴F_{5/2}.

The second intrinsic loss mechanism we are considering, the upconversion process, occurs via the dipole-dipole interaction of two excited Nd³⁺ ⁴F_{3/2} ions.¹⁶⁹ As shown in Fig. 6-95, the upconversion process causes one ion to be promoted to a higher excited state while the other ion drops to a lower energy state. The rate of decay due to the dipole-dipole interaction increases rapidly with increasing excited state density. The magnitude of the upconversion rate is determined by the overlap of the absorption spectrum for the upward transition and emission spectrum of the downward transition and the distance between excited centers. For a given doping density of the rare earth impurity, the upconversion rate will increase with pump fluence as the excited state density increases. For Nd³⁺ upconversion of ⁴F_{3/2} ions should be strongest to the ⁴G_{7/2} and ²G_{9/2} states. An approximate rate equation for the decay of the excited state population due to upconversion is:

$$\frac{d[N]}{dt} = -\frac{[N]}{\tau} - 2\alpha[N]^2 \quad (2)$$

where [N] is the density of excited Nd³⁺ ions and τ is the lifetime of the ⁴F_{3/2} state at low excitation. The upconversion rate α has units of cm³/s and is a function of the number density and spatial distribution of the ⁴F_{3/2} ions. The decay is a highly nonlinear process with the upconversion being enhanced by energy migration at some excited state densities. The experimental procedure for measuring α and typical values have been reported in the

literature for several Nd¹⁷⁰⁻¹⁷² and Er^{173,174} doped materials.

We investigated ESA of pump light and upconversion in Nd:LaF₃ during pumping at the saturation fluence with an Alexandrite laser (730–800 nm). As shown in Fig. 6-96, exciting Nd:LaF₃ in the red yields blue and green emission. The emission below ~510 nm arises from ESA of pump light which populates ²P_{3/2} and ⁴D_{3/2}. The emission above ~510 nm arises from the ⁴G_{7/2} state and is indicative of upconversion of two ⁴F_{3/2} ions.

From the time-integrated emission intensities and quantum yields of the ²P_{3/2} and ⁴G_{7/2} states, the fraction of ⁴F_{3/2} population lost to upconversion and ESA was calculated.

⁴F_{3/2} ions lost to upconversion

$$\propto \frac{\langle \text{Intensity } ^4G_{7/2} \rightarrow ^4I_{9/2} \rangle}{\eta h\nu} \quad (3)$$

For 0.2% Nd:LaF₃, approximately 10% of the excited ions were lost to upconversion during pumping and about 0.1% of the excited ions are lost to ESA when exciting at the saturation fluence. From these data, we estimate the value of α in Eq. (2) to be about $2 \times 10^{-17} \text{ cm}^3/\text{s}$.

Blue and green emission, similar to that of Fig. 6-96, was also observed during pumping Nd:YAG, Nd:SrF₂, and Nd:YOS. The strength of the ESA and of the upconversion will both change as the host is varied. For hosts with a short lifetime for the ⁴F_{5/2} + ²H_{9/2} state ESA will be weak. For most fluoride and oxide host crystals ESA will be negligible, predominantly due to absorption of pump photons by the thermal population in ⁴F_{5/2} + ²H_{9/2}. For hosts that have good overlap of absorption and emission spectra for the pairs of transitions shown in Fig. 6-95, upconversion can be an important loss process for high pumping densities. For host crystals that have sites where Nd³⁺ ions are paired or clustered, the upconversion rate can also be large.

Author: J. H. Schloss

Major Contributors: L. L. Chase, S. A. Payne

Figure 6-95. Possible upconversion paths for two $\text{Nd}^{3+} 4F_{3/2}$ ions.

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Figure 6-96. ESA of pump light and upconversion in Nd:LaF₃. Exciting Nd:LaF₃ in the red yields blue and green emission.

6.7.5. Ultrafast Nonradiative Relaxation Dynamics of Nd^{3+} in Glass

Introduction. Nonradiative relaxation of excited electronic states is one of the most important processes affecting the operation of laser media. This relaxation directly affects the quantum efficiency of luminescence and the dynamics of the upper and lower lasing levels. An empirical relationship commonly used to describe nonradiative relaxation is the "energy gap law."¹⁷⁵ Briefly, this law states that the nonradiative relaxation rate decreases as $\exp(-\alpha \Delta E)$, where ΔE is the energy gap between the initial level and the terminal level for the relaxation and α is a constant for a given host media. Though widely and sometimes indiscriminately invoked, this empirical relationship is only appropriate for describing nonradiative relaxation involving large energy gaps (multiphonon processes) and weak electron-phonon interactions. Nonradiative relaxation has been extensively studied in these situations and occurs in milliseconds to below 1 ns.¹⁷⁶ The case of small energy gaps (permitting rapid single phonon processes) and strong electron-phonon interaction has received little attention because of the difficulty in observing the ultrafast relaxation. Our work addresses this problem. We have studied the extremely fast nonradiative relaxation (picoseconds to femtoseconds) involving the small energy gaps between crystal field states of an excited state Stark manifold of Nd^{3+} in glassy media.¹⁷⁷

Experimental Procedure. Our experiments consisted of measuring the optical dephasing rate ($1/T_2$) of various transitions in Nd^{3+} . These inhomogeneously broadened transitions were from the $^4I_{9/2}$ ground state to the $^4G_{5/2}$, $^2G_{7/2}$ excited state Stark manifold of Nd^{3+} in silicate, phosphate, and fluoride glasses. Femtosecond accumulated photon echo techniques¹⁷⁸ were employed at low temperatures (13 K) to measure the dephasing rates. Since the major contribution to the dephasing rate comes from the nonradiative relaxation, we were able to determine the nonradiative relaxation rates from our measurements. The femtosecond pulses were generated by a tunable, hybridly mode-locked dye laser synchronously pumped by the frequency doubled output of a mode-locked Nd:YAG laser operating at a 76-MHz repetition rate.

Results and Discussion. We found the dephasing rate to be highly wavelength dependent. The variation of the dephasing rate as a function of excitation wavelength in a silicate glass at 13 K is shown in Fig. 6-97. The inhomogeneously broadened absorption profile is also shown in the figure. Similar results were observed for the phosphate and fluoride glasses. Both single and biexponential echo decays were observed. The slow decay rate of the biexponential is designated by a cross and the fast decay rate by a circle. Single exponential decays are also designated by a circle. Figure 6-97 shows that the data can be grouped into spectral regions where the dephasing rate changes smoothly with the excitation wavelength. There are seven such regions encircled in the figure. We have shown that these regions correspond to the location of the seven inhomogeneously broadened Kramers degenerate crystal field components in this Stark manifold.¹⁷⁷ Thus we were able to spectrally resolve the individual crystal field states from under the broad inhomogeneous profile of the Stark manifold by using this novel T2-selective dynamic excitation spectroscopy.

There are four possible sources of the dephasing: population relaxation arising from phonon emission, population relaxation arising from Orbach processes, Raman pure dephasing, and dephasing arising from the dynamics of the tunneling states of glassy media.¹⁷⁷ Our data indicate that except for the lowest energy crystal field state, population relaxation arising from phonon emission contributes most to the dephasing. The order of the phonon emission process can be determined from the phonon density of states (DOS). A phonon DOS which spans all the relevant energy gaps would indicate that single phonon emission is dominant over multiphonon emission. We have calculated the phonon DOS from Raman scattering data.¹⁷⁷ The calculation shows the relaxation to proceed via single acoustic phonon emission and is extremely fast. At low temperature (13 K) we observed population relaxation occurring in a few hundred femtoseconds to tens of picoseconds. Since acoustic phonons, and not optic phonons, are involved, the nonradiative relaxation behavior is not expected to change much from host to host. This is consistent with the similar relaxation behavior we observed in silicate, phosphate, and fluoride glasses.

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Figure 6-97. Variation of the dephasing rate of the $^4I_{9/2}$ - $^4G_{5/2}$, $^2G_{7/2}$ transition of Nd^{3+} in a silicate glass as a function of excitation wavelength at 13 K. See text for explanation.

Generally, the dephasing rate increased as higher lying crystal field states were excited. This reflects the larger number of available terminal levels for the relaxation and the larger density of acoustic phonon states. However, the detailed plot of the dephasing rate vs wavelength in Fig. 6-97 revealed considerable unexpected structure. As higher lying states were excited, the relaxation rate did not increase in the expected monotonic fashion. Surprisingly, the relaxation rate decreased as the excitation energy was increased within each of the seven inhomogeneously broadened crystal field states. We have analyzed these results with a perturbation treatment of single electron-phonon coupling. Our analysis indicates a site-dependent electron-phonon interaction and a strong correlation between the inhomogeneous broadening of the crystal field states.¹⁷⁷

In summary, we have directly measured the rates of phonon-induced population relaxation within the $^4G_{5/2}$, $^2G_{7/2}$ Stark manifold of Nd^{3+} in various glassy hosts. At 13 K these processes occur in the femtosecond to picosecond time regime, which is consistent with previous calculations.^{179,180} We have also determined that the mechanism of the relaxation is single acoustic phonon emission. Since acoustic phonons are involved, the nonradiative relaxation behavior is predicted to change little with the host material. This agrees well with our observations. We have also disproved the common misconception that the relaxation rate should increase monotonically as higher-lying crystal field states are excited. To date, this is the only detailed study of such a rapidly relaxing system.

Author: H. W. H. Lee

Major Contributor: R. S. Hughes, Jr.

6.7.6. Polarizability Change of Optically-Pumped Laser Ions

Introduction. Cr^{3+} and Nd^{3+} -doped crystals and glasses serve as the active medium for many types of solid-state lasers. The understanding of the electronic structure of these ions, and the subtle manner in which the host affects them, are key aspects for developing effective, new laser materials. We have used the four-wave mixing (FMW) method to investigate the electronic properties of the Cr^{3+} and Nd^{3+}

impurities, and, in particular, we have measured the difference in the polarizability of the ground and excited states. This pump-induced polarizability change also gives rise to a change in the index-of-refraction, which can potentially cause a new type of lensing phenomenon in the laser medium.

The essential features of the FWM experiments are illustrated in Fig. 6-98, where three laser beams are arranged to overlap in the sample volume. It is imagined that a "pair" of beams, (i.e. E_1, E_3 or E_2, E_3), initially interferes such that a modulated intensity profile appears within the sample. The high-field regions of the interference pattern are expected to generate excited states of the impurity ions, thereby giving rise to a population grating. In essence, the interference pattern becomes imprinted on the sample, via the excited state population. This grating serves to scatter the third input beam shown on the right-hand side of Fig. 6-98, so as to produce the signal wave, E_s .

Cr^{3+} -Doped Oxides. The energy levels of the Cr^{3+} impurity are illustrated in Fig. 6-99. The $3d^3$ electronic structure gives rise to the 4A_2 , 2E , 4T_2 and 4T_1 states designated on the figure. For the case of oxide crystals, there is a much stronger transition above these $3d^3$ states, known as the charge-transfer (CT) band. The CT transitions are the crucial features which cause the excited state to have a higher polarizability than the ground state. The FWM experiment is sensitive to this polarizability difference, $\Delta\alpha_p = \alpha_e - \alpha_g$. By carefully measuring the scattering efficiencies, laser powers, beam waists, impurity metastable state lifetimes and the sample absorption coefficient, absolute values of $\Delta\alpha_p$ may be determined.¹⁸¹

First, we set out to definitively prove that the scattering is due to the polarizability change. To accomplish this, we measured the dispersion of $\Delta\alpha_p$, as shown in Fig. 6-100.¹⁸² Here the grating was produced with an Ar-laser operating at 488 nm, and the scattering was observed for three different probe laser outputs at 442, 633, and 1064 nm. The rise in $\Delta\alpha_p$ can be fit with the equation:

$$\Delta\alpha_p = (fc^2/m) \left\{ \left[(\omega_0 - \Delta)^2 - \omega_p^2 \right]^{-1} - \left[\omega_0^2 - \omega_p^2 \right]^{-1} \right\} \quad (1)$$

Figure 6-98. Illustration of grating formation by the (a) E_1 , E_3 (upper left) and the (b) E_2 , E_3 pairs of beams, (lower left), and the subsequent scattering of the (a) E_2 (upper right), and the (b) E_1 (lower right) beams from this grating to generate the E_s signal wave.

LW-025-u-1368-6

Figure 6-99. Energy levels of Cr^{3+} in an oxide host. The $3d^3$ electronic states are shown, as well as the higher lying charge transfer transition. Either the 4T_2 or the 2E state may serve as the metastable level, depending on the host. The ground and excited state polarizabilities are denoted as α_g and α_e .

CW-1018-U-5732-1A

Figure 6-100 Dispersion of the polarizability change measured at three probe laser wavelengths. The datum represented with an "X" is used to establish the absolute scaling.

where ω_p is the probe beam frequency, Δ is the energy difference between the ground and metastable states, and ω_0 and f are the frequency and oscillator strength of the transition making the dominant contribution to the polarizability, (CT band in this case). The data for the Cr^{3+} -doped oxides shown in Fig. 6-100 can be accurately fit with this equation, and the derived values of ω_0 and f noted on the figure are, in fact, in close agreement with spectroscopic measurements of the position and strength of the CT band. This conclusively proves that the CT band of the Cr^{3+} -doped oxides is responsible for the polarizability change.

The values of $\Delta\alpha_p$ obtained for numerous Cr^{3+} -doped oxides are shown in Fig. 6-101.¹⁸¹ The ordinate indicates that the values of $\Delta\alpha_p$ are typically in the range of $0.01\text{--}0.1\text{Å}^3$. The abscissa of Fig. 6-101 is the energy separation between the ^2E and $^4\text{T}_2$ metastable states of Cr^{3+} . Either of these two states may serve as the metastable level, depending on their relative energies. The theoretical line plotted in Fig. 6-101 is calculated by assuming that the overall polarizability change is the sum of that due to the populations of the ^2E and $^4\text{T}_2$ states in thermal equilibrium. Systems such as $\text{BeAl}_2\text{O}_4\text{:Cr}^{3+}$ and $\text{LaMgAl}_{11}\text{O}_{19}\text{:Cr}^{3+}$, where the ^2E is much lower than the $^4\text{T}_2$ state, give the lowest $\Delta\alpha_p$ values, while crystals for which the energy levels are nearby, such as $\text{Gd}_3\text{Sc}_2\text{Ga}_3\text{O}_{12}\text{:Cr}^{3+}$, have the highest values. It is therefore apparent that the $^4\text{T}_2$ state has a much larger polarizability change relative to the $^4\text{A}_2$ ground state, than does the ^2E .

Nd^{3+} -Doped Crystals and Glasses. We have also performed FWM measurements on Nd^{3+} -doped materials. The main parameter characterizing the FWM results is again $\Delta\alpha_p$ the difference in the polarizability between the ground and metastable states. The Nd^{3+} energy level diagram in Fig. 6-102 shows that the metastable and ground states are designated as $^4\text{F}_{3/2}$ and $^4\text{I}_{9/2}$. The transition from either of these $4f^3$ states, to the $4f^25d$ manifold, is rather strong as a result of the allowed nature of the $4f \rightarrow 5d$ transmission. The strength and position of the $4f^25d$ state determines the value of $\Delta\alpha_p$ (as was the analogous situation for the CT band of the Cr^{3+} -doped oxides).

The measured polarizability changes are shown in Fig. 6-103 for the Nd^{3+} impurity doped into a variety of different host materials, including fluoride and oxide crystals, and

fluorophosphate, phosphate, and silicate glasses.¹⁸³ The $\Delta\alpha_p$ values are seen to be quite similar for each class of host material, since the Nd^{3+} site is expected to be very similar. On the other hand, we see that there is a great deal of host dependence for the different types of materials. For example, the oxides yield $\Delta\alpha_p$ values that are four times larger than those of the fluorides. An important question regarding the data in Fig. 6-103 is related to the nature and origin of the host dependence. Insight into this area can be gained by using Judd-Ofelt theory to express the magnitude of $\Delta\alpha_p$ in terms of the basic parameters describing the Nd^{3+} electronic structure. The result is:¹⁸³

$$\Delta\alpha_p = \frac{3\tilde{\nu}_{5d}\alpha_{\text{FS}}}{7\pi} \left| \langle 4f | r | 5d \rangle \right|^2 \left[\frac{1}{(\tilde{\nu}_{5d} - \tilde{\nu}_{\text{ex}})^2 - \tilde{\nu}^2} - \frac{1}{\tilde{\nu}_{5d}^2 - \tilde{\nu}^2} \right] \quad (2)$$

where α_{FS} is the fine-structure constant, and $\tilde{\nu}_{5d}$, $\tilde{\nu}_{\text{ex}}$, and $\tilde{\nu}$ are the energies in wave numbers of the $4f^25d$ state, the $^4\text{F}_{3/2}$ metastable state, and the photon, respectively. The $4f\text{--}5d$ radial integral also appears in this formulation. The two unspecified parameters are the $4f^25d$ energy and the $4f\text{--}5d$ radial integral. Using the free-ion values of these quantities we calculate that $\Delta\alpha_p$ (free-ion) = 0.080Å^3 . This magnitude is similar to that which we have measured for the oxides, 0.05Å^3 , although it is much larger than the fluoride results, $0.01\text{--}0.02\text{Å}^3$. The sensitivity of the $\Delta\alpha_p$ values can be traced to the effect of the host on the $5d$ orbital, and therefore, on the $4f\text{--}5d$ radial integral. This idea has been previously noted in several other contexts, including the $4f\text{--}5d$ transition moment, the $4f\text{--}4f$ transition strengths, Raman spectroscopy, and two-photon studies.

Population lens. Since the excited state of the Nd^{3+} and Cr^{3+} laser ions in the active medium have a different refractive index than that of the ground state, the polarizability change can potentially generate a "population lens" in the laser material. This gives rise to a change in the focus of an optical element when it is pumped nonuniformly. We can roughly compare the magnitude of the ordinary thermal lens and this population lens for the case of a

MO-S70-U-4714 2043 (a)

Figure 6-101. Measured polarizability changes for a variety of Cr^{3+} -doped oxides plotted as a function of the energy separation between the 4T_2 and $2E$ metastable states.

TCF-S36-U-6092-11

Figure 6-102. Energy level diagram of Nd^{3+} illustrating the $4F_{3/2}$ and $4I_{9/2}$ metastable and ground states, respectively. The associated polarizabilities, α_e and α_g , are determined by the position and strength of the transitions to the $4f^25d$ states.

JAJ-S8-U-6092-10

Figure 6-103. Measured polarizability changes for a wide variety of Nd^{3+} -doped materials. The hosts are grouped into five types, as indicated at the bottom of the figure. The sample identifications refer to the glass designation from the *Laser Glass Handbook* (M-095, Rev. 2).

single-shot amplifier, by calculating the relative pump-induced refractive index change, $\Delta\eta_{\text{pop}}/\Delta\eta_{\text{th}}$.¹⁸⁴ By noting that

$$\Delta\eta_{\text{pop}} = (2\pi f_L^2 N_{\text{ex}}/n) \Delta\alpha_p \quad (3)$$

where f_L is the Lorentz field correction and N_{ex} is the excited state density, and also noting that $\Delta\eta_{\text{th}} = (dn/dT) \Delta T$, the ratio of the population lens to the thermal lens can be estimated to be 0.59, -0.47, and -0.60 for Nd:Y₃Al₅O₁₂, Nd:LiYF₄, and LG-750, respectively. It is important to emphasize that these results apply to a single-shot amplifier, and not to a multipulsed system.

Authors: S. A. Payne and L. L. Chase (LLNL); R. C. Powell (Oklahoma State University); and S. C. Weaver (Princeton University).

Major contributors: G. Wilke, E. Prochnow

6.7.7. Nonradiative Lifetimes of Nd³⁺ Pump Levels in Halide and Oxide Crystals

Introduction. There has been much recent interest in pumping solid state lasers with semiconductor diode laser arrays and other solid state and gas laser pump sources. Laser pump sources can provide very high pump fluxes in a narrow band of wavelengths, and this property can introduce both opportunities and problems that are not encountered with flashlamp pumping. One can, for example, envision rapid pumping of rare earth ions like Nd³⁺ to excited states that have a moderately high nonradiative decay rate, and circumvent the nonradiative losses by extracting the pumped population rapidly. On the other hand, sequential absorption of pump light, which becomes more likely at high pump intensities, might significantly reduce pumping efficiency. The realization of the promising features of laser pumping, and the avoidance of the problems, requires a good understanding of the nonradiative decay rates and mechanisms of rare earth ions.

The Nd³⁺ ion, whose energy levels are depicted in Fig. 6-104, is currently the most popular activator ion for pumping with laser diode arrays and solid state lasers. When Nd³⁺ is pumped with flashlamp pulses tens or hundreds of microseconds in duration, it is a good approximation to assume that the pump levels

feed the ⁴F_{3/2} level with a negligible time delay and that there is negligible non-thermal-equilibrium population in the levels above ⁴F_{3/2}. This may not be the case with laser pumping. In recent reports of the laser performance of Nd:YAG and other crystals pumped by an alexandrite laser, it was suggested that the efficiency was reduced at some pump wavelengths by excited state absorption of pump light originating from the transient population in the ⁴F_{5/2} state.¹⁸⁵ In order to assess this possibility it is important to know the lifetimes of the excited states above ⁴F_{3/2} in various crystals. In the present work we have measured the lifetimes of some of the potentially important pump levels of Nd³⁺, ⁴G_{7/2}, ⁴G_{5/2}, ⁴F_{7/2}, and ⁴F_{5/2} in some fluoride, chlorides, bromide and oxide crystals. These levels are identified by arrows in Fig. 6-104. Our measurements confirm and extend the recent work of Basiev, et al., who did lifetime measurements of some Nd³⁺ multiplets in LaF₃, SrF₂, and YAlO₃.^{186,187}

It will be useful to compare our measurements with those estimated from the empirical energy gap law, given by the formula¹⁸⁸

$$W = Ce^{-\alpha\Delta E} \left[1 - \exp\left(-\frac{E_{\text{phonon}}}{kT}\right) \right]^{-n} \quad (1)$$

where ΔE is the minimum energy gap to the next lowest multiplet, and C and α are empirical constants which are assumed to depend mainly on the host medium and only weakly on the ion. The parameter α is proportional to the number of phonons emitted, n , which is the smallest integer larger than $\Delta E/E_{\text{phonon}}$.

Experimental Results. A Nd:YAG-pumped dye laser was used to resonantly pump each state, and the luminescence spectra and lifetimes were measured using a 1-m spectrometer, a photomultiplier tube with a GaAs photocathode, and a transient digitizer. The digitizer output was plotted on a logarithmic scale for data analysis. The dye laser pulses had a complex temporal shape with a width of 10 ns, and the photomultiplier tube and electronics had a response time of about 6 ns. Lifetimes were measured either from the tail of the transient signal following termination of the pump pulse or, in the case of the ⁴F_{5/2} level, from the risetime of the ⁴F_{3/2} luminescence. When the

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Figure 6-104. Energy levels of Nd^{3+} . The arrows at the right identify the multiplets for which lifetimes are reported in Table 6-24.

decay or rise time was only slightly longer than the response time of the apparatus, we estimated lifetimes using the approximate formula $\tau^2 = \tau_m^2 - \tau_r^2$, where τ_m is the measured decay time or rise time and τ_r is the response time of the apparatus. We avoided photomultiplier tube saturation by maintaining a suitable combination of low-light fluxes and high tube voltages. The transient decays observed when pumping the ${}^4G_{7/2}$ and ${}^4G_{5/2}$ levels of Nd^{3+} in the chlorides and in $CsCdBr_3$ were nonexponential, possibly because of nonradiative decay within the Stark levels of the multiplets. For these cases, we have quoted the lifetime observed after the initial rapid decay when the decay curve became nearly a single-exponential time dependence.

The results of our measurements are summarized in Table 6-24. The first column gives the host and Nd^{3+} concentration. Under the heading of each multiplet, the first column gives the values of ΔE and n (in parentheses). The second column gives the measured lifetime in ns. The label "F" means that the lifetime was "fast," less than the minimum measurable with our setup (~ 5 ns). We were not able to measure the lifetime of the ${}^4F_{7/2}$ level since its emission could not be detected, and pumping into the ${}^4F_{7/2}$ level yielded the same risetime of the ${}^4F_{3/2}$ emission as for pumping into ${}^4F_{5/2}$. We conclude that the lifetime of ${}^4F_{7/2}$ is considerably shorter than that of ${}^4F_{5/2}$, as expected from the smaller energy gap below ${}^4F_{7/2}$ in Table 6-24. Our measurements for LaF_3 and SrF_2 agree well with the few 300 K measurements reported by Basiev, et al.,^{186,187} except for the ${}^4G_{5/2}$ level in SrF_2 , for which they report a lifetime of 16.8 ns and we get 55 ns. The reason for this difference is unknown.

Discussion. We first consider excited state absorption (ESA) of pump light originating from one of the pump levels above the ${}^4F_{3/2}$ metastable state. This will be a significant loss mechanism only if the pump flux is an appreciable fraction of the saturation flux $\phi_s = hv/\sigma_{ESA}\tau$ for the ESA transition, where σ_{ESA} is the ESA cross section and τ is the lifetime of the initial state from which the ESA originates.¹⁸⁹ The most likely initial multiplet for ESA is ${}^4F_{5/2}$, which has a lifetime of several microseconds in the chlorides and bromide and less than 5 ns in the fluorides and YAG. This is the state identified by Kliever and Powell as the initial state for ESA in YAG.¹⁹⁰ Assuming a

reasonable estimate for $\sigma_{ESA} = 10^{-19} \text{ cm}^2$, the estimated saturation flux is $5 \times 10^5 \text{ W/cm}^2$, in the chlorides and bromide, and it is in excess of $5 \times 10^8 \text{ W/cm}^2$ for the fluorides and YAG. It is possible that significant ESA will occur for the chlorides and bromide if they are pumped by the spiked output of a long-pulse solid state laser, which can produce peak fluxes of well over 1 MW/cm^2 . The transient ESA is clearly negligible, however, for fluorides and oxides. We conclude that any ESA of laser pump light for Nd^{3+} in oxide and fluoride hosts must be caused by the thermal population of the ${}^4F_{5/2}$ level or by some other effect that does not result from the nonequilibrium population of the laser-pumped multiplets.

Our measurements for Nd:YAG yielded lifetimes that were less than 5 ns for all of the states investigated. These results are in disagreement with the measurements of Zverev et al.,¹⁹¹ who measured a rise time of about 0.5 μs for the ${}^4F_{3/2}$ emission when pumping directly into ${}^4G_{7/2}$. They interpreted this risetime as the lifetime of ${}^4G_{7/2}$. Very recently, Cruz et al.¹⁸⁵ pumped directly into the ${}^4G_{7/2}$ multiplet in Nd:YAG with a picosecond laser and observed a rise time of 0.47 ns for the 1.06- μm gain. They also interpreted this rise time as the lifetime of ${}^4G_{7/2}$. This measurement is consistent with our upper limit of 5 ns. Table 6-24 shows that the energy gap below the ${}^4G_{7/2}$ level is almost identical to the ${}^4I_{11/2}$ - ${}^4I_{9/2}$ gap in all the crystals. This suggests that the lifetime of the ${}^4I_{11/2}$ terminal state for 1- μm laser operation is likely to be ≤ 1 ns in YAG and oxide glasses and about 10 ns in fluorides.

Basiev et al.^{186,187} noted that the nonradiative decay rates they measured for several multiplets of Nd^{3+} in some fluorides and $YAlO_3$ are over two orders of magnitude faster than those predicted from the energy gap law. This breakdown of the gap law occurs for energy gaps in LaF_3 as large as 1500 cm^{-1} , corresponding to the emission of five phonons. For example, the lifetime of ${}^4G_{7/2}$ at 300 K calculated from Eq. (1) is 4.7 μs , whereas the measured value in Table 6-24 is 56 ns, a factor of 84 lower. In $LaCl_3$, the calculated lifetime for the four-phonon decay of ${}^4F_{5/2}$ is 27 μs , and the measured value is 3.3 μs , only a factor of six lower. Apparently, the gap law is unreliable for the excited quartet multiplets of Nd^{3+} .

Several years ago, Perlin et al. calculated nonradiative decay rates for some excited states

Table 6-24. Energy gaps and measured lifetimes for some multiplets of Nd^{3+} in crystals. "*" indicates data was obtained from the risetime of the emission from $^4\text{F}_{3/2}$. "+" indicates additional information: about 90% of the Nd^{3+} ions in CaCdBr_3 are paired in neighboring Cd^{2+} sites.

	$^4\text{G}_{7/2}$		$^4\text{G}_{5/2}$		$^4\text{F}_{5/2}$		$^4\text{I}_{11/2}$
	$\Delta\text{E}(\text{n})$	τ	$\Delta\text{E}(\text{n})$	τ	$\Delta\text{E}(\text{n})$	τ	ΔE
	(cm^{-1})	(ns)	(cm^{-1})	(ns)	(cm^{-1})	(ns)	(cm^{-1})
SrF_2 (0.1%)	1530 (5)	18	1135 (4)	F	900 (3)	F	1400
LaF_3 (0.2%)	1478 (5)	56	1168 (4)	14	959 (3)	F	1478
LaF_3 (3%)	"	42.5	"	14	"	F	1478
$\text{Na}_2\text{Y}_3\text{F}_{11}$ (0.1%)	-	22.0	-	7	-	F	-
LiYF_4 (~1%)	1381 (7)	9	1024 (3)	F	930 (2)	F	1475
LaCl_3 (0.5%)	1697 (7)	24000	1134 (5)	7000	1004 (4)	3300*	1725
PbCl_2 (0.5%)	-	20000	-	6800	-	4200*	-
CsCdBr_3 (0.5%)+	1466 (9)	40000	700 (4)	1100	957 (6)	16500*	-
YAG (1%)	1152 (2)	F	736 (1)	F	860 (2)	F	1150

of Nd^{3+} in rare earth garnets.¹⁹² Their calculations yielded subnanosecond lifetimes for all four of the states in YAG which have lifetimes shorter than our time resolution of 5 ns, and they are also close to the $^4\text{G}_{7/2}$ lifetime measured in Ref. 185

Conclusions. The lifetimes of the $^4\text{G}_{7/2}$, $^4\text{G}_{5/2}$, and $^4\text{F}_{5/2}$ states of Nd^{3+} in several oxide and halide hosts are much shorter than predicted by the energy gap formula. The lifetimes of $^4\text{F}_{5/2}$ level in YAG and all of the fluoride crystals that we investigated are too short to significantly enhance the excited state absorption of laser pump light above the level caused by thermal population of the state.

Author: L. L. Chase

Major collaborator: L. K. Smith

6.7.8. Energy Level Assignments for MgO:Ni^{2+}

Introduction. The MgO:Ni^{2+} crystal has been regarded as a prototypical system among impurity-doped hosts for many decades because the Ni^{2+} ion is surrounded by a perfect octahedron of oxygens. Laser action has been demonstrated, although with poor efficiency.¹⁹³ One prominent loss mechanism is excited state absorption (ESA), in which a photon is absorbed from the excited state to a higher lying energy level, rather than contributing to laser action. Our study¹⁹⁴ of the ESA and unrelaxed luminescence spectra of MgO:Ni^{2+} serve to rationalize the poor laser performance of MgO:Ni^{2+} that has been observed, and also to provide new insight into the nature of the energy levels of MgO:Ni^{2+} .

The transitions to be investigated are schematically illustrated in Fig. 6-105. We have observed the $^1\text{T}_2 \rightarrow ^3\text{T}_{1a}$, ^1E luminescence and the $^3\text{T}_2 \rightarrow ^3\text{T}_{1a}$, ^1E ESA spectrum for the first time. These spectra are displayed in Fig. 6-106 along with the $^3\text{A}_2 \rightarrow ^3\text{T}_{1a}$, ^1E ground state absorption spectrum. All of these spectra are plotted relative to the energy of the $^3\text{A}_2$ ground state, so that they may be directly compared. The stronger unrelaxed luminescence band at $13,000\text{ cm}^{-1}$ is due to the $^1\text{T}_2 \rightarrow ^1\text{E}$ transition, while the $^1\text{T}_2 \rightarrow ^3\text{T}_{1a}$ transitions in the $13,500\text{--}16,500\text{ cm}^{-1}$ region are much weaker, as may be expected due to the influence of the spin selection rules. Similarly, the transition from the $^3\text{A}_2$ ground state and the $^3\text{T}_2$ excited state favor the $^3\text{T}_{1a}$ state, rather than the ^1E (see Fig. 6-106). These

spectra may also be analyzed in a more detailed manner to definitively assign several of the spin-orbit components of the $^3\text{T}_{1a}$ state (see Ref. 194 for further explanations). It has also been determined that the $^3\text{T}_2 \rightarrow ^3\text{T}_{1a}$ ESA band is likely to overlap the $^3\text{T}_2 \rightarrow ^3\text{A}_2$ laser transition, and that the two transitions have similar cross sections. Therefore, ESA to the $^3\text{T}_{1a}$ state appears to be responsible for the low laser efficiency of MgO:Ni^{2+} that has been observed.

Author: S. A. Payne

Major Contributors: G. T. Ullery, G. D. Wilke

6.7.9. Compilation of Rare-Earth Laser Transitions

Introduction. We have recently completed a detailed compilation¹⁹⁵ of all reports of laser action for crystalline solid-state lasers activated by paramagnetic ions (transition metals and rare-earths), that have appeared in the 1980s. While this review involves over 500 references and extensive tables of laser parameters, its details are clearly far beyond the scope of this article. We indicate the nature of our work by including two figures that describe the known laser transitions of rare-earth ions. Fig. 6-107 illustrates all of the known laser transitions of trivalent rare-earth ions, and the typical wavelengths at which they occur. The main requirements for the realization of laser action include the availability of a pump source, the existence of an adequate stimulation emission cross section and storage time, the absence of various competitive processes (such as cross relaxation, nonradiative decay, and excited state absorption), and the availability of a crystalline sample of adequate optical quality. There are probably many more transitions among the energy levels shown in Fig. 6-107 which satisfy these requirements in addition to the ones reported in the figure. Fig. 6-108 contains information similar to that expressed in Fig. 6-107, although it is plotted as a function of wavelength. The sparse number of long wavelength lasers is primarily due to the influence of the energy-gap, since the close spacing of the energy levels gives rise to an enhanced nonradiative decay rate and thus to a short energy storage time. The small number of short wavelength lasers is due to the absence of effective pump sources, and to reduced emission cross sections at shorter wavelengths.

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Figure 6-105. Energy levels of MgO:Ni^{2+} indicating the transitions investigated in this work.

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Figure 6-106. Actual spectra corresponding to the transitions indicated in Fig. 6-105. The energies have all been plotted relative to the 3A_2 ground state, so that the data may be directly compared.

Figure 6-107. Transitions of trivalent rare-earth impurities that were lased in crystal-line hosts, (dashed arrows refer to those transitions only lased in glassy materials). The wavelengths are indicated in μm .

Most of the laser transitions reported in Figs. 6-107 and 6-108 are not well studied, and therefore their potential is unknown. The data in our compilation provide an indication as to the great diversity of solid-state laser materials that exists. It is interesting to note, however, that laser action was demonstrated on eight of the transitions in Fig. 6-107, after they had been predicted by Caird in 1975.¹⁹⁶ Furthermore, we may expect many new lasers to be discovered, evaluated, and engineered in the future.

Authors: S. A. Payne and J. A. Caird

6.8. Athena/LMF Studies

6.8.1. Introduction

The National Academy of Sciences has established a committee which is in the process of conducting a major review of the national inertial fusion program for the Department of Energy. During the course of that review we have presented our concept for Athena, a future Laser Microfusion Facility using a glass laser driver which could achieve high fusion gain in the laboratory. This section will give an overview of that design and of the modeling and cost optimization studies which we are using to refine it.

Target designs for such a facility are discussed elsewhere in the Laser Program Annual Report. In summary, we judge that a facility having a laser driver energy and pulse length of about 4-5 MJ in 8 ns at 350 nm is an appropriate compromise between the cost, performance risk, and applications of a facility designed to demonstrate high fusion gain. If it is necessary to reach fusion yields as high as about a gigajoule, then the size of the facility rises to perhaps 8-10 MJ of driver energy in about 10 ns. A facility designed to demonstrate ignition and some gain would need to provide 1.5-2.5 MJ in 3-5 ns.

6.8.2. Overview of the Athena Concept for a Laboratory Microfusion Facility

The Athena driver is a modular, frequency-converted glass laser design concept that we have been developing over the last several years for use in a future Laboratory Microfusion Facility. The basic beamline design for this driver will be the same for any of the options discussed above, though details such as the

number of disks in each final amplifier will change somewhat depending on the pulse length for which we optimize performance. The Athena design allows us to size and cost these various possibilities for an LMF with good confidence that the performance of the facility will meet or exceed the nominal driver specifications. Our analysis—and this description—will focus on the final amplifier of the laser beamline and the beam transport optics between that amplifier and the target. These two categories of components represent about two-thirds of the cost of an LMF-scale facility, as shown in Fig. 6-109, and are therefore the most important part of the LMF facility design. The cost of smaller facilities tends to be dominated by other costs, such as small-scale “front end” hardware, buildings, and design labor. Final amplifiers and beam transport were only about a third of the cost of the Nova laser facility, and they would be an even smaller fraction of the cost of a smaller laser system.

The basic design philosophy for Athena is to require that all performance parameters or components must already have been demonstrated at some scale in the laboratory, though perhaps not all simultaneously. In some cases (such as size of frequency conversion crystals with high-damage threshold) that scale may be smaller than we require in the final system, but the path for scaling to the size required is clear. We believe that further development will give a somewhat better performance than the baseline Athena design which we present here, up to perhaps 30% greater energy on target for roughly the same quantity of laser hardware.

The Athena Beamline. Athena is a laser design based on a beamline with a laser energy output of 230 kJ in 8 ns at 1053 nm with frequency conversion in nonlinear crystals to the third harmonic at 350 nm. (Operation at 530 nm or 260 nm is possible, but not included in the baseline design.) After beam transport and frequency conversion, the energy on target from a single beamline is 145 kJ in a nominal 8-ns main drive pulse at 350 nm. The various options for an LMF facility use different numbers of these beamlines, as shown in Table 6-25 and 6-26. We shall describe the mid-range option (32 beamlines for an energy on target of about 4.5 MJ) in detail here.

The Athena beamline uses neodymium-doped phosphate glass of the same type as used

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Figure 6-109. Cost distribution for a megajoule-scale LMF facility compared to the 0.1 megajoule-scale Nova laser facility. This figure includes only the costs to purchase and install the facility as shown: development, auxiliary buildings, etc., are not included. The final amplifiers and beam transport account for about two-thirds of the cost of the LMF. For Nova these components are about one-third of the cost.

Table 6-25. Options for a Laboratory Microfusion Facility.

Option	Fusion yield	Driver required
High-end	~1000 MJ	~8–11 MJ/10 ns
Mid-range	~100–200 MJ	~4.5–6 MJ/ 8 ns
Ignition	some	~2–2.5 MJ/ 5 ns
		~1.5–2 MJ/ 3 ns

Table 6-26. Athena Beamlines for a Laboratory Microfusion Facility.

Option	Number of beamlines	Disks per beamline	350 nm output energy per beamline (kJ)	Design pulse length (ns)
High-end	56	14–16	165	10
Mid-range	32	12–14	145	8
Ignition	16–18	10–12	125	5
			95	3

in the Nova laser. Amplified spontaneous emission in the laser glass requires that we limit the aperture of an individual subaperture of the amplifier to about 30 cm. In the Athena design we have chosen to segment the entire optical system at a subaperture size of 30×30 cm. The amplifier uses a 4×4 array of these subapertures, as shown in Fig. 6-110, so that a single beamline contains 16 parallel subaperture beams in an overall size of about 1.5×1.5 m. It is essential for direct-drive fusion targets, and desirable for indirect-drive, to smooth the effective intensity profile on the target. The Athena design achieves this smoothing by operating each of the 16 subapertures at a different frequency within the gain bandwidth of the glass, and then overlaying the focal spots of all 16 on the target. This gives an effective laser bandwidth of $50\text{--}100\text{ cm}^{-1}$ (0.15–0.3%) at 350 nm.

Figure 6-111 compares the Athena amplifier design to the amplifiers of the Shiva and Nova lasers at LLNL. One can see from this figure that the Athena amplifier is a natural evolution from these earlier amplifiers and that the scale-up from Nova to Athena is of roughly the same magnitude as the scale-up from Shiva to Nova.

We are considering two major options for the architecture of the laser cavity that extracts energy from the final amplifier. There are significant advantages to multipass extraction of amplifiers when the design permits highly efficient use of stored energy. The Athena design falls into this category so both options are multipass architectures, though with different numbers of passes.

It is not difficult to arrange for double-pass extraction in an amplifier, and we are pursuing that as one option. Figure 6-112 shows a block diagram of the layout of a 4.5-MJ, 32-beamline Athena using the double-pass geometry. The stage gain (energy out/energy in) of the final amplifier in double-pass operation is about 100, so each final amplifier beamline requires an input energy of about 2.4 kJ. This implies a fairly large preamplifier between the small-scale hardware of the oscillator and “front end,” and the 32 final amplifier beamlines. Since we wish to have a minimum variety of different parts for mass production, the simplest concept for this preamplifier is to install two additional final amplifier beamlines to be used as lightly saturated preamplifiers having a stage gain of about 800. (The nominal 80 kJ output from each of

these two beamlines is about twice the input required for the 16 final amplifier beamlines which it drives, but it is necessary to have a design margin of this sort to accommodate beam balancing, propagation losses, and apodization at the final amplifier.) The double-pass option, then, requires that we purchase 34 beamlines of final amplifier hardware in order to put 32 beamlines on target.

The other option that we are pursuing is a multipass amplifier which contains an active switch to extract the output beam from the final amplifier cavity, as shown in the block diagram of Fig. 6-113. This option has 3-1/2 passes through the final amplifier rather than only two, so the stage gain is about 1.4×10^4 and the input energy to the final amplifier stages is negligible. For this option we must have an optical switch to direct the beam out of the laser resonator, and the switch must operate at the full output fluence of the laser and with low loss.

We demonstrated such a switch at the Athena single subaperture scale in experiments conducted several years ago, as shown in Fig. 6-114. The switch is a Pockels cell using a thin plate of electro-optic crystal (potassium dihydrogen or dideuterium phosphate). The switch rotates the polarization of the beam when a voltage is applied between two transparent, conducting plasma electrodes formed by striking an electrical discharge in a low-pressure gas on each face of the crystal. Our experiments show that it is not difficult to operate a switch of this size with an opening time of about 20–30 ns, which is adequate for the Athena design.

We believe that the 3-pass concept will give a less expensive and more compact LMF than the double-pass option, but the differences are perhaps a few tens of millions of dollars in cost and 10–15% in energy on target for a mid-range LMF such as a 32-beamline Athena. These differences are large enough to be worth understanding in some detail, but they do not have a major impact on the overall cost of the LMF. Also, since the switched multipass option is more complex than the double pass, we shall describe it in more detail here.

The Switched Multipass Option.

Figure 6-115 shows a diagram of the layout of the final amplifier optical cavity of the switched multipass. The cavity contains two amplifier assemblies between two plane mirrors, and an active optical switch assembly consisting

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Figure 6-110. The Athena amplifier assembly contains neodymium-doped phosphate glass slabs mounted at Brewster's angle in a 4×4 array. Each individual slab is about $32 \times 50 \times 4.5$ cm, and is roughly the same size as the 31.5-cm disks used in the Nova laser. Xenon flashlamps located in the spaces between the slabs and along the outside walls pump the amplifier.

Shiva

- 20 cm aperture
- 500 J output

Nova

- 46 cm, 2-segment aperture
- 10 kJ output
(20 × Shiva)

Athena

- 1.5 m, 16-segment aperture
- 100–200 kJ output
(10–20 × Nova)

Figure 6-111. The scaling from Nova to Athena is comparable to the scaling from Shiva to Nova. The Nova amplifier introduced segmented apertures with two segments each having a size of about 46×35 cm. The Athena amplifier merely increases the number of similarly sized segments to 16.

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Figure 6-112. Block diagram of a 4.5-MJ, 32-beamline Athena using double-pass final amplifiers. Each group of 16 beamlines uses one additional beamline as a preamplifier.

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Figure 6-113. Block diagram of a 4.5-MJ, 32-beamline Athena using Pockels-cell-switched, 3-1/2-pass final amplifiers. This option requires less hardware than the double-pass of Fig. 6-112, but requires a high-fluence, low-loss switch.

Figure 6-114. Prototype of a plasma electrode Pockels cell at the scale of the Athena subaperture (26×26 cm). Two low-pressure gas discharge chambers serve as transparent electrodes that apply a very uniform electric field to the crystal along the direction of propagation of the laser.

Figure 6-115. Diagram of the final amplifier cavity of the switched multipass option for Athena. The final amplifier contains two-disk amplifier assemblies separated by a spatial filter or relay telescope.

of the Pockels cell just described and a polarizer. There is also a relay telescope, consisting of two identical lenses separated by twice their focal length, that occupies half of the length of the amplifier cavity. This telescope serves several functions. The simplest of these is to limit the solid angle which can propagate through the amplifier by placing an aperture or "pinhole" in the far-field focal plane in vacuum at the center of the telescope. Since this plane is the Fourier transform plane of the beam in the amplifier, the pinhole clips off high-spatial frequency intensity or phase modulations growing from flaws in the optics or from nonlinear effects and thereby forces any intensity and phase noise on the beam to be of a scale size larger than the spatial frequency cutoff of the pinhole.

The telescope has another less obvious function. When the telescope is half the length of the cavity, the lens system forms a real image of every point within the left side of the cavity on the right side, and vice versa. (This is easiest to see when the cavity is symmetric, with a length $L/4$ on each side of the telescope, but this also holds for asymmetric arrangements.) When a real image forms in an optical system, the diffractive optical propagation distance from the original object goes to zero. A small intensity or phase distortion, such as a dust speck or a small optical flaw shows no diffractive ripple growth near the image since the effective propagation distance is near zero. Also intensity noise on the beam converts to phase noise by way of nonlinearities in the index of refraction at high intensities, and this process amplifies any intensity noise on the beam. Image-relayed propagation minimizes this noise growth.

Image-relayed propagation therefore gives a beam intensity profile, both within the amplifier and after switchout, that is much more uniform than it would be with long unrelayed paths over which modulation can grow by linear and nonlinear diffraction. Laser-induced optical damage to components has a threshold at some value of local beam fluence, so the average beam fluence can be much higher if there is a uniform beam intensity profile. The uniformity and high fluence makes it easier to extract energy efficiently from the amplifiers, and also reduces the area and cost of the beam transport optics in the laser system.

Figure 6-116 shows the beam path through the switched multipass. The beam enters near the center of the relay telescope just as the

Pockels cell is switched to the "on" position which leaves the polarization of the beam unchanged and allows it to pass through the thin-film polarizer and circulate in the cavity. The injected pulse bounces three times from the cavity mirrors, which gives $3\frac{1}{2}$ passes through the two amplifier assemblies. Each of these has a gain of 5.8 (1.75 Nepers). When 70% of the energy stored in the two amplifier assemblies is extracted into the output pulse, the stage gain (energy out/energy in) through the $3\frac{1}{2}$ passes is about 1.4×10^4 . As the pulse makes its last transit through the cavity, the Pockels cell is switched to the "off" position which rotates the beam polarization perpendicular to the plane of the figure. The thin-film polarizer then reflects the beam out of the amplifier cavity.

Figure 6-117 shows a close-up of the region near the focal plane in the center of the relay telescope. The two end mirrors are not exactly parallel, so there is a small angular separation between passes that separates the focal spots of the passes in the focal plane. The angular separation gives a simple way to inject the input pulse at small diameter using an input mirror near the focal plane of the spatial filter. When a high-energy beam passes through the pinhole of a spatial filter, the energy in the high-spatial-frequency noise components of the beam strikes the edge or face of the pinhole and generates a plasma that expands to fill the pinhole. If we try to put another optical pulse through the pinhole a short time later, the plasma may absorb or distort that pulse. Our design avoids that possibility. There are no low-loss closed optical paths as there would be in a design with a single pinhole, so parasitic oscillations are less important. Also the switch may not be perfect, in which case some of the output energy will remain in the polarization which stays trapped within the cavity. Our design protects the amplifier against damage due to any such radiation which might continue to circulate in the cavity.

Figures 6-115 to 117 show the concept behind the switched multipass design, but the actual layout is somewhat more complex. The amplifier and other optics are fully segmented at an aperture of 30 cm, shown in **Fig. 6-110**. Then a top view of the beamline corresponding to **Figures 6-115 to 117** actually looks like **Fig. 6-118**. There are four vertical stacks of four parallel subaperture beams each having its own relay telescope, switch aperture, and polarizer.

Figure 6-116. Beam path through the cavity of Fig. 6-115. The input beam enters in the center of the spatial filter and has 3-1/2 passes (or three mirror reflections) before the switch fires to extract it from the final amplifier.

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Figure 6-117. Detail of the beam paths in the far-field or at the center of sign has several useful properties: (1) Simple injection of the input pulse at small beam diameter; (2) No distortion of the beam by plasma generated by earlier passes; (3) No closed paths for parasitic oscillations; (4) Trap for any radiation not switched out of the cavity when the switch fires.

Figure 6-118. A top view of the switched multipass option for Athena showing the fully-segmented amplifier array. There are four vertical stacks of the four parallel subaperture beams shown in the diagram.

Each amplifier assembly contains seven disks at Brewster's angle in a 4×4 array. The cavity end mirror and polarizers are arrays of 16 subaperture mirrors or polarizers mounted on a common rigid plate, shown in Fig. 6-119. The aperture of the output beam from the beamline is about 1.5×1.5 m. The length occupied by the amplifier assemblies, switch, and polarizer is 15 m, so the length of the cavity must be twice that, or 30 m. A full beamline has an input energy of about 16 J for an output energy of 230 kJ at 1 nm in an 8-ns pulse, and with an average fluence of 20 J/cm^2 in the output beam.

The input beam injection uses a beam that is also an array of 4×4 subapertures, each with a nominal size of about 4×4 cm, as shown in Fig. 6-120. This beam enters through an injection relay telescope and strikes a 4×4 array of small splitter mirrors located in the center of the relay telescope array near the pinhole plane. Each of these mirrors, only one of which is shown in Fig. 6-120, picks up one of the 16 input beams and directs it to a position between two subapertures and about 1 m away from the central pinhole plane of the relay telescope array. At this point each input beam strikes a mirror that directs the beam towards the injection mirror for its particular subaperture. An apodizer trims the beam to the 3×3 cm size required for injection, and a polarizer and rotatable half-wave plate act as a variable attenuator to balance the input to each of the subapertures in the beamline. The beam then strikes a 75-cm focal length input lens which brings it to a focus at the pinhole plane in the center of the relay telescope. On its way to that plane, the beam strikes an injection mirror located about 15 cm from the focal plane, shown in Fig. 6-117, and that mirror directs it into the beam path for its particular subaperture.

The injection system is fully image-relayed through this path. The injection telescope forms an image of the input aperture of that telescope at a point just before the variable attenuator assembly, and the input lens images that plane onto the input aperture of the amplifier assembly.

Figure 6-121 is a perspective drawing of a full beamline to scale showing all of the features discussed above. The beam paths of the injection system and the 3-1/2 amplifier passes are shaded within the relay telescope, which is enclosed in a vacuum chamber not shown in this drawing. There is one 7-disk amplifier array and mirror at each end of the cavity. At the far end

are the Pockels cell assembly and a polarizer array that direct the output beam into the beam transport system which takes it on its way to the target chamber.

Figure 6-122 is a perspective drawing of the beam transport system that conveys the output beam from a single beamline to the frequency conversion crystal array, and then on to the target. The transport system contains three turning mirrors in the layout presented here, though we have considered other options. There is also an output relay telescope that images the output aperture of the beamline onto the frequency conversion array to maintain a uniform intensity profile on that array.

Figure 6-123 shows a more detailed layout of the final two turning mirrors, the frequency conversion array, and the lens array that illuminates the fusion target. The final turning mirrors are adjusted so that the optical axes of the 16 subapertures of each beamline coincide at the plane of the target. The individual subapertures within a beam are separated from each other by 5–10 cm due to obscurations in the amplifier assembly and mounting flanges for lenses and other components. The final turning mirrors can be configured as shown in Fig. 6-123 so that this separation is larger at the frequency conversion array, if that is convenient for the final mechanical design.

The frequency conversion array consists of two nonlinear crystals in series for each subaperture of the array. The first of these crystals is designed to convert two-thirds of the input 1053-nm energy to the second harmonic at 527 nm. The second crystal then mixes that green light with the remaining 1053-nm light to give the third harmonic at 351 nm. These crystal arrays are very similar to arrays which we now have on the Nova laser, as shown in Fig. 6-124. We have demonstrated that we can convert 70% of the input energy to the third harmonic in a Nova array of this sort, though we do not yet operate routinely at that level. For Athena we have assumed routine operation at 70%, with some room for further improvement. Conversion efficiencies above 85% have been achieved under well controlled laboratory conditions.

Efficient frequency conversion requires that the crystals be aligned to precisely that angle which gives equal propagation velocities (index of refraction) for the fundamental and harmonic beams in the doubling and mixing crystals. The index changes rapidly with wavelength in the

Athena
(Same scale as photo)

Figure 6-119. A sketch of a 4×4 subaperture array of mirrors such as those used for Athena cavity end mirrors or beam transport to the target. The individual mirrors in this array are somewhat smaller than the largest turning mirrors used in the Nova laser.

Relay/Spatial filter

Figure 6-120. The input beam injection array for the Athena switched multipass. The input beam enters the final amplifier as an array of small beams from a preamplifier stage. These strike an array of splitter mirrors that direct each of them to its appropriate position in the pinhole plane. Each input beam has a variable attenuator, and is fully relayed through the amplifier. Only one such beam is shown in this figure.

Figure 6-121. Sketch of an Athena beamline showing all the features discussed in Figs. 6-115 to 120. The beam paths of the injection system and the 3-1/2 amplifier passes are shaded in the relay telescope; this region is enclosed in a vacuum chamber not shown in this drawing. There are disk arrays and mirrors on either side of the relay telescope assembly. At the far end are the Pockels cell and polarizer arrays that switch the beam out of the cavity and direct it into the beam transport system.

Figure 6-122. Sketch of the beam transport system for Athena. The turning mirrors and output relay telescopes arrange the beams in the desired geometry around the target chamber.

Figure 6-123. Detail of the final section of the beam transport system. The final two turning mirrors repoint each subaperture so that all of the subapertures from a single beamline illuminate the same spot on the target. The beams are converted to the third harmonic in nonlinear crystals just before they are focused on the target.

Nova

**Angular alignment
tolerance achieved $\pm 50 \mu\text{ rad}$**

Athena (Same scale as photo)

**Angular alignment
tolerance spec $\pm 50 \mu\text{ rad}$**

Figure 6-124. An Athena frequency conversion array compared to the similar conversion arrays now in service on Nova. The Athena arrays are only slightly larger than Nova arrays, though higher damage threshold allows them to handle much larger energy.

ultraviolet, so the bandwidth over which this phase-matching condition has negligible effect on conversion efficiency is only about $2\text{--}3\text{ cm}^{-1}$. We wish to have different frequencies over a range of $50\text{--}100\text{ cm}^{-1}$ at 350 nm for intensity averaging in the beam. We must, therefore, cut each of the 16 doubling and 16 mixing crystals in the array at slightly different angles to the crystal axes so that the crystals in a particular subaperture are aligned properly for the frequency of the beam in that subaperture and the entire array can be pointed as a unit. This is no more difficult than cutting and mounting all the crystals with the same tolerances to precisely the same angle, as was done for the Nova array, though crystals are no longer interchangeable between subapertures.

After frequency conversion to 350 nm , the beam proceeds through a focusing lens into the target chamber, and on to the fusion target. There is a blast shield between the focusing lens and the target to protect the expensive optical components from target debris. There may be a random phase plate to control the beam's transverse spatial coherence to improve the uniformity of the intensity profile on the target. We estimate that the standoff distance from final focusing lens to target should be greater than about 15 m for a target yield of $100\text{--}200\text{ MJ}$ as appropriate for the mid-range facility with a 32-beam, 4.5--MJ driver. The high-end LMF would, of course, require a larger target chamber and standoff distance to avoid neutron, shrapnel, and x-ray damage to walls and optics. The design issues of high-yield chambers are discussed in more detail in another section of this report.

The Double-Pass Option. Figure 6-125 shows the layout of the Athena double-pass option, for comparison to Fig. 6-115. The double-pass option is really just the last two passes of the switched multipass option and uses the same split amplifier with internal relay. The input and output to the amplifier are separated in angle within the output relay telescope which images the amplifier output on target. We must inject a much larger energy here than for the switched multipass, so the injection mirror must be larger to avoid damage. The design must reach some compromise between poor fill-factor in the final amplifier due to a large angular separation between the two passes, and length of the output spatial filter. The design shown in Fig. 6-125 places the mirror at a position where the fluence

perpendicular to the beam path is about 25 J/cm^2 in the injected beam and the fill-factor penalty is similar to the switched multipass design. The fluence on the mirror itself is much smaller, because of the steep angle of incidence; however, damage to defects in optical coatings (which is usually the lowest threshold process) has only a weak dependence on angle of incidence. The output relay telescope is then about 42 m long, which is somewhat longer than other considerations would dictate. The amplifier cavity is somewhat shorter than the multipass since it is no longer necessary to leave space within the cavity for the switch and polarizer. The full beamline is again a 4×4 array similar to that shown in Fig. 6-118. The injection system is similar to Fig. 6-120, but larger to handle the higher energy. We must inject about 150 J into each subaperture, or 2400 J into the full beamline. The output of the beamline will be slightly greater for the double-pass because the double-pass has no switch losses.

The 32-beamline LMF Facility. Figures 6-126 and -127 show side and end views of a 32-beam LMF facility using the Athena switched multipass beamline. There are two laser halls at opposite ends of the building, each containing 16 beamlines which are stacked four high in four groups. In the center of the facility is a 7-m -radius aluminum target chamber designed to handle fusion yields up to $100\text{--}200\text{ MJ}$. Extensions from that chamber locate the final focusing lenses at a radius of 15 m from the target. This layout shows an illumination geometry designed for indirect drive targets that distributes the 16 beamlines at each end into three cones with apices at the target. There are six beams, each in cones with half-angles of 55° and 65° , and four beams in a cone at 35° . The central core of the LMF building around the target chamber is a cube about 50 m on a side. From this extend two laser halls about 20 m tall, 50 m wide, and 40 m long, which gives an overall building footprint of about $50 \times 125\text{ m}$.

Figure 6-128 shows the size of the existing Shiva-Nova facility at LLNL for comparison to the 32-beam, 4.5-MJ LMF of Figures 6-126 and -127. The LMF building is somewhat shorter than Shiva-Nova, and has about twice the floor area.

Summary. We believe that the Athena concept is the best current option for a Laboratory Microfusion Facility. The Athena driver technology uses similar concepts and components

Figure 6-125. Diagram of the final amplifier cavity for the Athena double-pass option. The double-pass option is just the final two passes of the multipass shown in Figs. 6-115 to 121. It requires no switch, but higher input energy. The input beam injection array is similar to that of Fig. 6-120, but is located in the output relay telescope between the final amplifier and the target.

Figure 6-126 side view and Fig. 6-127 End view of a 32-beamline, 4.5-MJ Athena Laser Microfusion Facility. There are two laser halls with a target area between them. Each laser hall contains 16 beamlines.

Figure 6-128. Sketch of the existing Nova Laser Facility for comparison with Figs. 126 and 129. The Athena building is a bit shorter than the Nova building and has about twice the floor area.

of similar size and performance to those which we have already demonstrated on the Nova laser, so the risk that it may not meet the baseline performance is low. Also there have been major production runs of similar components for Nova and other large glass lasers, so development and production costs are much easier to project than for less-developed laser candidates. We feel that the time has come to proceed with an LMF based on this concept.

Authors: D. N. Frank, J. T. Hunt, J. R. Murray, E. Storm, J. B. Trenholme

6.8.3. Athena LMF Cost Study

Introduction. The LLNL approach to the construction of a Laboratory Microfusion Facility for the National ICF and Weapons Physics Programs is to use the Athena solid-state multipass laser concept,¹⁹⁷ which is a very modular and flexible system. Because of its modularity and commonality of components, the individual parts count is reduced and much of the system can employ mass production techniques, thus reducing design, manufacturing and assembly cost.

Solid state systems, like Athena, can be modularly built up from a single prototype beamline, to 18 beamline (1.5–2.0 MJ) target facility and then to a larger 32–56 beamline (4.5–11 MJ) Athena-type driver as a national user facility. The smaller megajoule class experimental facilities can be used by design laboratories to conduct smaller yield preliminary experiments with a higher experiment and data rate.

Cost Analysis. The Laboratory Microfusion Capability Study Phase I Summary¹⁹⁸ (DOE/DP-0069) established a standard work breakdown structure (WBS), shown in Fig. 6-129, to be used by all laboratories for categorizing and comparing costs.

LLNL has used this WBS methodology to develop component work packages for the Athena LMF design. To date, we have prepared conceptual designs for the complete system, with detailed component designs for approximately 35% of the laser system. A preliminary design report for the experimental area, including environmental systems, optics protection, and system diagnostics, has also been prepared.

We have used three costing methodologies to establish our LMF cost estimates.

- Top-down costing: establishes the project cost scope and compares conceptual designs and system trade-offs.
- Modified bottom-up costing: establishes preliminary subsystem costs and identifies areas where significant performance and cost savings can be achieved.
- Bottom-up costing: establishes the final detailed system cost.

The “top down” costing is based on physics conceptual designs and uses scaling from similar existing laser systems. A “modified bottom-up” costing includes engineering conceptual drawing and component cost data from industrial suppliers. A “bottom-up” costing is performed on subsystems with detailed engineering component drawings, so that vendor quotes can be obtained for the type, quantity and delivery schedule required to meet the project criteria. To date, we have determined approximately 20% of the subsystem costs using “bottom-up,” 35% using “top-down” and 45% using “modified bottom-up.”

A cost breakdown is given in Table 6-27 for the baseline 32-beam, 4.5- to 6-MJ Athena system. These costs are based on industrial manufacturing estimates for the power amplifier assemblies, on manufacturers’ quotes for optical materials, finishing, and coating, and on designs and preliminary cost estimates by industry for alignment and power conditioning systems. Costs for the laser and experimental buildings were developed by architects at LLNL. As shown in Table 6-27, the total construction cost for a 4.5- to 6- MJ driver and a 200-MJ-yield experimental facility is \$585 million (FY89 dollars).

Table 6-27 also gives costs for a 56-beamline driver that could deliver 8.5–11 MJ at 0.35 μm and an experimental area capable of containing a 1000-MJ target yield, for weapons vulnerability, lethality, and effects experiments.

LMF Laser/Experimental Facility Options. The LLNL Athena LMF design has the flexibility to allow alternative laser energies, high-yield experiment facilities, and target irradiation configurations, either during the original facility construction stage or after construction of the baseline system. The incremental project costs for adding these LMF options are estimated in Table 6-28.

Athena Prototype Beamline. The LLNL ICF development plan describes the component characterization and manufacturing process

Figure 6-129. LMF work breakdown structure (WBS) for the LLNL Athena Driver.

Table 6-27. LMF Athena construction costs for two energy options (FY89 dollars).

WBS	Description	(4.5-6MJ) 32 beams	(8.5-11 MJ) 56 beams
		\$M	\$M
1.1	Experimental facilities (laser, target, and support buildings)	36.0	45.5
1.2	Other special equipment (optic shop and computer equipment)	4.5	6.5
1.3.1	Laser alignment and diagnostics	25.9	37.6
1.3.2	Front end (oscillators, preamplifiers, spatial and temporal coherence control)	30.3	35.7
1.3.3.1	Amplifiers with glass and flashlamps	88.3	127.9
1.3.3.2	Cavity (spatial filters, switch, and end mirrors)	58.1	83.1
1.3.4	Beam transport (coherence control and frequency conversion)	65.4	92.6
1.3.5	Power conditioning	39.3	68.1
1.4	Target area - 200-MJ Yield - 1000-MJ Yield	48.0	101.0
1.5	Project Office	14.1	21.9
1.6	Data acquisition and controls	22.0	27.9
	Subtotal	431.9	647.8
	Contingency (labor & hardware)	(17.7) --	(26.6) --
		(75.1) 92.8	(112.7) 139.3
	Industry technology development (line item)	60.0	60.0
	Total cost (line item)	\$584.7M	\$847.1M

Table 28. Project costs for LMF options.

	32 beamlines	56 beamlines
Laser energy and pulse length	4.5–6 MJ 8 ns	8.5–11 MJ 10 ns
Experimental area		
Initial high gain and yields to 200 MJ	\$585 M	----
1000-MJ full weapons effects chamber	----	\$847 M
Second direct drive experiment area (200 MJ) with beam support	\$200–220 M	\$205–225 M
Second experiment area with beam transport and 1000-MJ full weapons effects	\$200–220 M	\$260–280 M

research needed to achieve final cost validation. We can integrate these developments into an Athena prototype beamline and characterize the overall performance of an LMF system. Such a prototype would enable LMF scientists and engineers to begin the design, integration, and performance testing of the beamline at an accelerated schedule. The prototype also would provide confirmation of the system performance and would reduce the cost to construct the full LMF.

The prototype would consist of:

- Oscillator and beam injection system.
- Full scale 4 x 4 amplifier.
- Spatial and temporal coherence control.
- Multipass alignment and control system.
- Optical switch for a single beamlet (if required in final design).
- Beam transport for one beamlet.
- Frequency conversion for one beamlet.
- Beam and system diagnostics for one beamlet.

The prototype multipass amplifier would produce 230 kJ at 1.05 μm in an 8-ns pulse. We would extract, transport, and frequency-convert the energy from one beamlet, delivering 9 kJ at 0.35 μm , with spatial and temporal coherence control. All beamlet alignment, control, and laser diagnostics systems would be fully operational so that the overall system performance and beam quality could be verified. However, since the Athena design is fully segmented, we would not propose to complete a full-scale prototype, i.e., extract and convert all 16 beamlets to 0.35 μm to produce the full 145 kJ of 0.35 μm energy of an Athena beamline.

The prototype, if built, would be installed in our existing Bldg. 381 high-bay laboratory (Fig. 6-130). This laboratory was used for the Nova prototype beamlines (called Novette) in 1981. This facility has the required thermal and vibration stability, computer and master oscillator rooms, and adequate space for power conditioning in the basement.

Prototype System Cost. Because we would not be constructing a full LMF, economies of scale would generally not be achievable for the prototype. The design cost would be a greater fraction of the total cost than for the LMF, and facilities costs would be different than for the

LMF. Furthermore, the prototype would be constructed before the cost reduction manufacturing R&D was complete which requires older and more costly manufacturing processes. Therefore, the overall cost for the prototype could not be linearly extrapolated to determine the overall Athena LMF cost. However, a large fraction of the design engineering and some of the fabrication development costs incurred in building the prototype would reduce the Athena LMF design and construction costs and would allow a more accurate cost estimate for the full LMF to be made. The prototype cost summary in Table 6-27 used the same WBS as the LMF system.

Since the early 1970s, LLNL has worked to improve the laser and optical technology base; to transfer laboratory technology to industry; and to encourage development of lower-unit-cost manufacturing processes. In the absence of competition, high-technology industries have little incentive to develop low-cost manufacturing processes on their own beyond those minimally necessary. We have found that properly placed manufacturing R&D contracts oriented specifically to lowering manufacturing costs, have high leverage in developing and maintaining a competitive, qualified, multiple vendor network. This philosophy was responsible for the large difference (in some cases more than a factor of 10) in cost actually paid for Nova optics and the typical industry price for the same optics before Nova. We will continue this process during the conceptual design and system construction for Athena.

The cost estimates presented in this paper were developed by obtaining estimates from industrial suppliers based on engineering drawings of amplifiers, optical components, target chamber designs, and power conditioning equipment. We are preparing a full cost study report that will provide a detailed laser and experimental area detailed "bottom-up" costs.

Author: R. J. Foley

6.8.4. The NTES Laser Facility for Physics Experiments—A Study

The Nuclear Test and Experimental Sciences Program (NTES) requested that the Laser Program design a laser system capable of delivering a few hundred joules to a kilojoule of energy in a range of pulse forms to two experiment areas for a variety of laser physics

Figure 6-130. Athena prototype beamline installation in Bldg. 381 laser high bay.

and science experiments. Two options were considered in the study. In one option a new laser would be constructed in the Bldg. 381 high bay. In the other option the necessary modifications would be done to the Nova facility to meet the needs of the NTES directorate. The results of this study are reported in their entirety.^{199,200} What follows is a summary of the study.

The new facility was to have the characteristics shown in Table 6-29.

Experimental Capability. The NTES Physics Facility would provide experimenters with the capability to perform most of the laser physics and other scientific experiments discussed in Appendix E of Ref. 199. The system design would also accommodate future enhancements in the driver and experimental capability, as shown in Tables 6-30 and 6-31.

Building 381 High Bay Option. The proposed facility would be located in the high bay of Bldg. 381. An additional 7500 ft² of laboratory space to the north of the Bldg. 381 high bay would be constructed. A block diagram of the entire facility is shown in Fig. 6-131. The main components of the facility are the experiment area, the high-energy laser, auxiliary lasers, controls and data acquisition, and the building and modifications (facility).

The experiment area has two experiment chambers and some diagnostic hardware brought directly from Janus.

The twin-beam, high-energy laser subsystem uses a double-pass amplifier laser architecture and includes beam transport, laser alignment, and the diagnostics system. An ultra short-pulse and probe interferometry beam were also included in the study. Additionally, a low-energy beam that bypasses the high-energy laser subsystem is available for "set-up" shots. The energy pulse, which would have a Gaussian temporal shape is produced by one of three master oscillator/preamplifier (MOP) chains in the front end. These drive either of two high-energy laser beamlines or an ultra short-pulse and probe beamline. To increase the versatility of the NTES facility for the Laboratory, one of the two high-energy beamlines would be configured to serve both as a reliable, low-cost, high-energy laser for NTES, and as a test bed to allow the ICF Program to develop multipass laser technology for the proposed Laboratory Microfusion Facility. This dual-laser architecture uses both a double-pass laser cavity and a multipass scheme that employs a Pockels

cell switch. When performing NTES experiments, the amplifiers would be double-passed using a Faraday-rotator-polarizer combination as a switch. This configuration isolates the laser cavity components from the 1- μ m light back-reflected from targets.

When doing research for LMF development, this beamline would be conveniently converted to the multipass architecture by inserting a mirror near the end of the cavity. The beam will then bypass the rotator and will be routed to a test stand that contains the Pockels cell switch. Note that this test bed would never irradiate targets because there would be no isolation while in this configuration; this would prevent damage from back-reflection.

When operating at 10 ns, this beamline would produce between 5–8 kJ at 1.053 μ m.

Building 391 High Bay Option. For this option we proposed to use the existing Nova laser facility located in the high bay of Bldg. 391. An additional 5000 ft² of laboratory space to the east of the Bldg. 391 high bay would be constructed to house a new experiment area. The facility requirements are the same as the 381 high bay option. The experiment area has two experiment chambers and some diagnostic hardware brought directly from Janus. Two beams of the Nova laser through the 20.8-cm disk amplifiers would be used as the high-energy beams, and the additional prerequisite beam transport, laser alignment, and the diagnostics system would be purchased and installed. An ultra short-pulse and probe interferometry beam are also included in this proposal. Additionally, a low-energy laser subsystem is available for "setup" shots.

The energy pulse, which has a Gaussian temporal shape, is produced by one of three master oscillator/preamplifier chains in the front end. These drive either of two high-energy laser cavities or an ultra short-pulse and probe beamline.

Authors: D. J. Christie and J. T. Hunt

Major Contributors: D. N. Frank, M. A. Hennesian, D. W. Larson, J. K. Lawson, S. E. Mayo, H. G. Patton, J. B. Richards, and G. L. Tietbohl

Table 6-29. Facility Characteristics.

- Two experiment chambers.
- Two high-energy beams, each with 250–500 J per beam at a wavelength of $\sim 1.053\ \mu\text{m}$, and 200–400 J/1 ns per beam at $0.526\ \mu\text{m}$.
- A setup beamline that bypasses the laser cavity to be used for calibration of experiment area instruments.
- One probe beam (0.1 J rated at 30 ps) to be used for interferometry.
- An ultrashort-pulse beam (0.1 J rated at 1 ps).
- A minimum of three oscillators.
- Approximately $7500\ \text{ft}^2$ (Ref. 200) of additional laboratory space.

Table 6-30. General Facility Capabilities.

- Several experiments to be setup in the same time frame and conducted alternately, with the completion of possibly eight tests per day (approximately one shot per hour per beamline at full energy) and numerous lower-energy tests.
- Beams to be transported from one experiment area to another on a shot-by-shot schedule.
- Energies ranging from millijoules to a kilojoule to be delivered for different experiments.
- Experiments using from one to three beams (two 20.8-cm-diam beams and one 5-cm-diam beam) at durations of from 1 ps to tens of nanoseconds, using wavelengths of 1.053 and 0.53 μm .
- Alignment beams with wavelengths of 1.053 and 0.53 μm are available to both experiment areas.
- All of the laser beams to be delivered to either of the experiment areas.

Table 31. Future enhancements in driver and experimental capability.

Master Oscillator /Preamplifier	<ul style="list-style-type: none"> • Two general purpose beams capable of pulse durations as short as 30 ps and as long as 10 ns, with chirp and pulse-stretching capabilities, delivering energies of 0.1 J in 30 ps to 1–5 J in 1 ns, with cavity bypass capability (setup beam). • One ultrashort-pulse/probe beam capable of delivering 0.1-J pulses with durations of 30 ps or 1 ps; the probe would be frequency converted to either 3, 4 ,or 5 ω. • A repetition rate of more than one shot every ten minutes. • A pulse-to-pulse energy variation of less than 5% (or as quoted by vendor if reasonably close to 5%). • A prepulse ratio for the fundamental wavelength of less than 10^{-5}. • A beam divergence of less than ten times the diffraction limit. • Protection against back reflections from the main laser chain or the experiment areas. • Output and midchain calorimetry. • Output pulse shape and timing. • Prepulse beam energy. • Output beam profile.
High-energy, laser-cavity beamlines	<ul style="list-style-type: none"> • A pulse-to-pulse energy variation of less than 10%. • A prepulse ratio for the fundamental wavelength of $\sim 10^{-5}$. • A beam divergence of less than ten times the diffraction limit. • Input, output, prepulse, and midchain energy diagnostics. • Diagnostics for output temporal and spatial pulse shape, and timing.
Experimental area	<ul style="list-style-type: none"> • Two general purpose target chambers • Provisions will be made for the installation of Janus target diagnostics. • Provisions will be made for two more experiment chambers, which may be installed at a later date.
Computer control and data acquisition capability	<ul style="list-style-type: none"> • Computer-assisted control systems for pulsed power, beam transport, laser alignment, and laser diagnostics. • An oscillator-control system that will allow for the setting of the time of arrival of a pulse at the experiment area to within approximately 50 ps and with no more than approximately 20 ps jitter. • An early (300-ns) trigger system, delivering at least five synchronized (<20 ps), 100 V, and less than 1-ns-rise-time triggers to each experiment area.

Figure 6-131. NTES laser facility for physics experiments.

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7 ICF Applications

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7. ICF Applications

7.1. Introduction

Our ICF Applications group revised the HYLIFE-I power plant design. With HYLIFE-II, we aim to avoid the hazard of burning lithium in water and air, while still achieving desired target and beam performance. The CASCADE power plant design underwent revisions to eliminate vaporization; yet to be developed are ways to clear beam paths, protect beam ports, and reliably and economically insert the target into the reactor. We developed a vacuum system for both the HYLIFE and CASCADE reactor concepts.

We investigated stimulated Raman scattering for limiting laser propagation through an ICF reactor target chamber.

After evaluating the Laboratory Microfusion Facility (LMF), we expect it to meet all safety requirements. We also performed a thorough systems analysis for the LMF, identifying the facility's objectives, its essential systems, the systems requirements, the added features that would be desirable, and troublesome interfaces.

7.2. Reactor Concepts and Technology

7.2.1. HYLIFE-II

The HYLIFE-I power plant design carried out between 1977 and 1985 used falling liquid lithium jets between the target and the vessel walls to reduce the stress in the wall due to the explosive energy release and to protect the wall from neutrons and x-rays.

During 1989 we began a study¹ to revise the HYLIFE-I design, which we are calling HYLIFE-II. We are trying to avoid the hazard of burning Li in water and air by substituting the molten salt called Flibe (Li_2BeF_4). At the same time we are designing for beam and target performance that is now thought more likely to be achieved. HYLIFE-I assumed laser light or heavy ions at about 4.5 MJ in two narrow clusters of beams, giving a gain of 400.

Lasers are not considered practical for a liquid fall reactor because many beams ($\gg 8$),

each widely spread from one another in angle, are needed to illuminate the indirect-drive target, and these are difficult to get through the jet array. In addition, the low efficiency of making laser energy, the low predicted gains, and high cost must be overcome to make lasers practical power-plant drivers.

Heavy ions are thought to be a good driver candidate, but the predicted construction costs of accelerator systems must be reduced by about a factor of 3 before the cost of heavy-ion drivers will not significantly affect the cost of electricity for a 1-GWe plant. The required driver arrangement is now thought to be 5 to 10 (or more) narrow (± 0.01 to ± 0.02 radians) beams on each of two sides delivering 5 MJ for a gain of 70. This is 5 times lower gain than in HYLIFE-I. To get the same power, the repetition rate must increase from 1.5 Hz to about 8 Hz. The driver energy can be increased to reduce the repetition rate by a factor of 2 or so. After the target explosion occurs, some of the liquid jets will fragment, causing splash. Clearing the beamlines of splash droplets and globules so the next beam of heavy ions can propagate to the target is a major issue being addressed. At the low repetition rate of HYLIFE-I of 1.5 Hz, splash will fall from rest about 2 m between shots. This was not enough to clear the beamlines in HYLIFE-I, so this problem was never solved. A small amount of splash might be tolerable; e.g., if 10% of the beam path is blocked, then a 10% increase in driver energy might be a solution. In HYLIFE-II we are shortening the liquid fall from 5 m used in HYLIFE-I to about 2 m and injecting pulsed liquid jets that will pass through the beam path after the shot to clear splash liquid in preparation for the next shot.

Safety studies of HYLIFE-II show the tritium inventory to be so small that in an accident tritium is not the major issue. But ^{18}F , with its 2-hr half life, could contribute significantly to hazards in an accident, so release of only a small fraction of the ^{18}F inventory can be tolerated. Because fluorine is so active, we expect it to be tied up in the nonvolatile forms LiF and BeF_2 , so that in an

accident only a small fraction of the total would be expected to be released. In a liquid Li system with 1 kg of tritium dissolved, there would be a large release of tritium. Tritium will permeate easily through pipes, requiring special attention to prevent more than 40 Ci/d tritium release to the environment. We think it is possible to keep tritium release low enough by an appropriately sized tritium separator using a liquid sprayed into a pumped chamber.

Author: R. W. Moir

7.2.2. CASCADE

CASCADE is a power plant design carried out between 1982 and 1986 using flowing grains of sand consisting of layers of granules of pyrolytic (ceramic) graphite, BeO if needed and Li₂O or LiAlO₂. The sand flows "down" the sides of rotating cones, picking up heat from neutrons, x rays, and debris. The hot sand is thrown in an upward arc into hoppers that direct it to flow down through SiC heat exchanger tubes carrying helium to a turbine. The hot sand transfers its heat to the tubes by radiant heat transfer. The x rays will vaporize ~1 kg of the exposed layer of graphite following each fusion pulse and might cause the granules to fragment. Vaporized graphite will most likely condense in a different form, which would require daily remanufacturing of the graphite granules. Some fraction of the vaporized graphite will leave the rotating conical reactor chamber and condense onto surfaces colder than ~3900 K. This graphite must be recovered in some way. It must not interfere with the beam propagation, and in case of an air-ingress accident it could contribute to an explosive energy release if too much were allowed to accumulate. The lead from the target would fill this reactor chamber at a density too high for the heavy ion beams to propagate through, unless there were a self-pinched propagation mode.

During 1989 CASCADE underwent important revisions. It has been proposed to surround the target with a layer of frozen lithium of about 1 to 2 kg which will intercept the x rays and debris and shield the granules from the intense heating enough that no vaporization is predicted. In this case the added tritium breeding in the lithium avoids the need for BeO granules, and LiAlO₂ can be used. As it rushes out of the chamber in viscous flow, the Li entrains the 0.4 g of Pb, reducing its density enough to

permit propagation of heavy ion beams. Clearing the beam paths and protecting the beam ports from x rays, debris, and neutrons must still be investigated. Economic and reliable ways in which the target might be inserted in any reactor are yet to be developed, but it will probably be easier to inject a cryogenic target with a lithium shield than to inject a target by itself.

Author: J. H. Pitts

7.2.3. Starlight

The Starlight concept²⁻⁴ for an ICF reactor uses a softball-sized solid-lithium x-ray and debris shield that surrounds each fuel target as it is injected into the reactor. The shield vaporizes as it absorbs x rays and ion-debris energy emanating from the fusion reactions in the target. Use of the shield increases the energy deposition times at the surface of the first wall by four orders of magnitude (to >100 μs), which allows the energy to be conducted into the wall fast enough to prevent vaporization. Starlight operates at 5 Hz with 300-MJ-yield targets. It features a stationary, nonvaporizing first wall that eliminates erosion and shock waves that can destroy the wall; it may also allow arbitrary target illumination geometries, so that efficient coupling of either laser or heavy ion beam driver energy to the target can be achieved.

When neutrons penetrate the shield, the wall experiences neutron damage that limits its lifetime. We must therefore choose wall materials that have an economic lifetime. A specific design for laser drivers uses a 6-m-radius, 2 1/4 Cr-1 Mo steel first wall with a 1-m-thick Li₂O granule blanket (50% density). An economical lifetime (6.5 yr) and a favorable tritium breeding ratio are predicted. This is a low-temperature option for Starlight. Either water or helium is the likely coolant. Water is not compatible with Li₂O, and water may be less economical than He operating at higher temperatures. The use of He at low temperatures, while removing the concern of water and Li₂O reacting, is still not as economical as operating at high temperatures.

A high-temperature option requires the use of a ceramic first wall and breeding material. All the ceramics we examined are severely limited for first wall applications by He production. Further, a favorable tritium breeding

ratio with the high-temperature LiAlO_2 solid breeding material with a ceramic first wall is not possible unless Be is a major constituent of the ceramic. A thin layer of Be (in an inexpensive pressed powder form) behind the first wall will be effective in increasing tritium breeding only if the ceramic wall thickness is set by less stringent criteria than that from resisting buckling from high-pressure He. Then, the ceramic must have a much higher damage threshold from He accumulation (appm level), or the ceramic must have a He release rate at high temperatures that gives the material an economical wall lifetime.

Author: J. H. Pitts

7.2.4. Hydro*Star

We introduced a new ICF reactor concept⁵ called Hydro*Star that uses DT-ignited DD capsules and a water blanket. The concept uses an advanced heavy-ion driver concept and a target-chamber blanket that is 1 to 2 m of water. Each fusion pulse vaporizes nearly 0.5 cm of the inside surface of the water, thereby creating hot steam, which is vented directly from the target chamber into steam turbines. Thus, Hydro*Star is merely a large steam engine. Advantages of the new concept include a thermal efficiency near 50%, reduced plant radioactivity, longer wall component lifetimes, reduced risk of catastrophic accidents, and potentially lower cost of electricity. Disadvantages include the required advanced target designs, advanced driver concept, and difficulties in interfacing and propagating the driver beams through the high-pressure target-chamber environment.

Author: C. D. Orth

7.2.5. Vacuum Pumping for ICF Reactors

We developed a vacuum system⁶ for the HYLIFE and CASCADE ICF reactor concepts that consists of an axial-flow compressor followed by mechanical booster pumps in series having a combined total of four stages. Vacuum components are placed outside the reactors and operate at a temperature of 300 K. Pumping speeds necessary to remove hydrogen isotopes and helium from the reactors were conservatively calculated to be a maximum of 24 to 36 m^3/s at 1.3 Pa (10^{-2} Torr) when laser drivers are used and 310 to 620 m^3/s at 0.13 Pa (10^{-3} Torr) when heavy-ion-beam drivers are used.

Openings in the reactors, required for illumination of the fuel pellets and for establishing the flowing blankets inside the reactors, have high enough conductance's for removal of the gases from inside the reactors.

Author: J. H. Pitts

7.2.6. Constraint on Residual Gas Density in Laser Fusion Target Chambers Because of Stimulated Raman Scatter, UCID-21816

Laser propagation through an ICF reactor target chamber establishes a maximum on the allowable gas density, and hence on the chamber shot repetition rate. Small-angle forward stimulated Raman scattering is an important process in limiting propagation and setting this constraint. A conservative model—propagation through a uniform plasma with saturated plasma wave amplitude—leads to a fractional energy loss which scales as the square of wavelength and is independent of intensity. For a 10-MJ target in an Ar atmosphere, we expect substantial energy loss to onset very abruptly as the gas density increases beyond about 10^{16} cm^{-3} ; for a Flibe atmosphere, the acceptable density of LiF molecules may be about a factor of 3 smaller ($3 \times 10^{15} \text{ cm}^{-3}$). If the gas ionization process is dominated by x-ray photoionization, then density gradient stabilization will substantially increase the allowable density.

Author: R. A. Sacks

7.3. Laboratory Microfusion Facility

7.3.1. Neutronics and Radiological Safety Analysis

One goal of the Laboratory Microfusion Facility (LMF) is to conduct some 1400 target experiments over a 5-year period to demonstrate high gain in the laboratory. Projected nominal yields for these experiments range from <10 MJ to 1000 MJ. High gain experiments are expected to be conducted once per week. After high gain ($>10\times$ input energy) has been achieved, target yields above 100 MJ will be used to conduct weapons physics and weapons effects experiments. In addition, experiments will be conducted to assess the viability of inertial fusion for electric power production, space propulsion, and space power. This will extend the facility lifetime to as much as 30 yr.

All aspects of the facility will meet the U.S. Department of Energy (DOE) radiological siting requirements, the International Commission on Radiological Protection (ICRP) (and DOE) public dose limits, the DOE occupational dose limits, and the more stringent LLNL "as low as reasonably achievable" (ALARA) guidelines. It will also meet the EPA-proposed emission standards for radionuclides.

We have set two design goals that minimize the effect of radiation hazards on the operation of the facility. The first is that LMF workers will have access to the diagnostic platforms outside the target chamber within 24 hr after a high-yield shot. This is necessary for retrieving data and preparing for the next experiment. Although it is plausible that routine interior chamber operations may be done remotely, workers may occasionally require access to the inside of the chamber for special tasks.

Therefore, the second design goal is that workers can safely enter the chamber for a short time 7 days after a high-yield experiment. As quantitative measures of these two design goals, we have established dose rate limits of 2 mrem/hr outside the chamber 24 hr after a shot (considering all sources of decay radiation) and 100 mrem/hr inside the chamber 7 days after a shot.

We have previously reported⁷ many aspects of radiological safety concerning the LMF. Here we report further work covering beam tube activation, final optics activation, neutron heating of a diagnostic, electromagnetic pulse (EMP) generation, and radiation shielding and tritium recovery/disposal issues.

Aluminum and carbon steel are both acceptable beam tube materials if workers can wait for 1 to 3 days after a high-yield shot before close contact with the beam tube material. For the first few days, the dose rate from a steel tube is lower than that from an Al tube, but after about a week, the Al tube has a significantly lower dose rate.

Although all acceptable harmonic conversion materials generate a ^{16}N hazard (over 1000 rem/hr in some cases) immediately after a high-yield shot, the dose rate is negligible after a few minutes.

The BK-7 final mirror presents the most significant problem. One day after a shot the dose rate is over 20 mrem/hr, primarily due to ^{24}Na . For times longer than a week after a shot, the dose rate is dominated by ^{24}Na , which

decays rather slowly (half life ≈ 2.6 yr). This leads to a plateau in the dose rate vs time curve. The dose rate from ^{24}Na continues to build up over the life of the facility, and after 1000 shots, the plateau is slightly greater than 2 mrem/hr.

The EMP generated in the inner chamber will probably be quite large. Its late-time arrival will not impede data collection, and circuit hardening techniques can be employed to ensure the survivability of diagnostics.

A proposed neutron diagnostic placed 5 cm from a target will be destroyed by neutron heating and will generate a significant shrapnel hazard. A 2.6-m-thick concrete shield provides public protection from prompt radiation well within Federal Safety Regulations. The routine tritium release rate (<100 Ci/yr) will produce a dose to the local population that will be less than 0.1% of the 25 mrem/yr regulatory limit.

Author: M. T. Tobin

7.3.2. Systems Analysis

By the middle of 1989, great progress had been achieved in performing a thorough systems analysis⁸ for the proposed Laboratory Microfusion Facility (LMF). This analysis identified the facility's objectives, its essential systems, the systems requirements, the added features that would be desirable, and troublesome interfaces. Key design variables were identified; risk analysis was the most influential. Preliminary analyses were conducted to choose the optimum target-chamber first-wall material as frost, to select the optimum size of the target chamber and the optimum distance to the final optics, and to identify features of systems to protect the final optics from target emissions.

Author: C. D. Orth

7.3.3. Spherical Target—Chamber Gas Dynamics and Wall Stresses

We began to assemble a compendium⁹ of the physical processes and mathematical equations governing first-wall x-ray ablation, gas dynamics, and the history of first-wall stresses for spherical ICF target chambers. Such physical processes as dissociation, ionization, and gas reverberations were considered, as well as the effectiveness of ports to vent the chambers and thereby reduce wall stresses. We discovered that up to about 6 times more wall mass can be

vaporized by the cooling of the vapor originally ablated by the x rays. The first-wall stress, previously thought to be caused by an ablation impulse and a long-term stable interior pressure, is actually dependent on these and on the sub-millisecond interior pressure-time integral, all in relation to the period for radial wall oscillations. We found that the post-ablation pressure impulse dominates the wall stresses for chamber radii smaller than about 0.5 m.

Author: C. D. Orth

7.4. Defense Applications of ICF

7.4.1. Overview of Defense Applications of ICF¹⁰

We have examined the extent to which the LMF may contribute to the weapons physics and weapons effects communities' requirements in the first decade of the next century and beyond.¹⁰ Current Nova weapons experiments were considered and then extrapolated to LMF energy and spatial scales to give some measure of LMF utility for weapons physics

research. Three types of weapons effects testing that would be done in the LMF if it were available today were examined in detail. While future needs are currently unknown, these current potential contributions lend confidence to the notion that the LMF could play a major role in weapons effects testing when it does become available. The impact of new underground test limits, either to a few kilotons, and/or severe restrictions on the number of tests, is to increase the benefit of an LMF in this scenario. The limitations of direct drive fusion to provide the entire range of defense applications was also considered. Finally, the production of tritium and special nuclear materials and the advantages fusion has over fission in performing this mission were explored in detail. While providing an important defense capability, the LMF will address issues that need to be resolved to establish the feasibility of ICF for energy production.

Authors: W. J. Hogan and M. T. Tobin

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Editing/Proofreading

Jason Carpenter
Al Miguel
Pete Murphy
Joy Perez
Leslie Wilder

Art/Production Staff

Chuck Chapman
Galen Hazelhofer
Kim M. Lanto
C. Sue Lawrence
Daniel Moore
Don Ryder

Layout

Pamela D. Dickinson

Print Coordination

Ron Natali

Authors

Section 1: E. M. Campbell, D. L. Correll

Section 2: W.E. Alley, H.A. Baldis, R.O. Bangerter, J. Barnard, B.I. Bennett, R.A. Bosch, S.T. Brandon, D.A. Callahan, G. Caporaso, L.B. DaSilva, S. Dixit, D.C. Eder, G.D. Enright, K.G. Estabrook, A. Friedman, T. Godlove, D.P. Grote, J.A. Harte, M. Hermann, D.H. Hewett, C.J. Keane, D.S. Kershaw, J.S. Kollin, W.L. Kruer, A.B. Langdon, Y.T. Lee, D.A. Liberman, R.A. London, B.J. MacGowan, D.L. Matthews, S. Maxon, K.O. Mikaelian, J.K. Nash, T.W. Phillips, M.K. Prasad, M.D. Rosen, H.L. Rutkowski, R.A. Sacks, H.A. Scott, A.I. Shestakov, B. Shore, A. Szoke, J. Trebes, S.V. Weber, B.L. Whitten, S. Yu, G.B. Zimmerman

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Section 6: L.J. Atherton, C.E. Barker, J.A. Britten, R.M. Bruscaro, J. Caird, J.H. Campbell, L.L. Chase, D.J. Christie, J.F. Cooper, A.C. Erlandson, R.J. Foley, D.N. Franck, J.T. Hunt, K.S. Jancaitas, F.W. Krupke, D.W. Larson, H.W.H. Lee, W.H. Lowdermilk, D. Milam, F.P. Milanovich, K.E. Montgomery, J.E. Murray, J.R. Murray, N. Nielsen, S.A. Payne, H.T. Powell, J.H. Schloss, L.K. Smith, E. Storm, I.M. Thomas, J.B. Trenholme, C.S. Vann, S. Velko

Section 7: W.J. Hogan, R.W. Moir, C.D. Orth, J.H. Pitts, R.A. Sacks, M.T. Tobin

Major Contributors

R. Adair, C. Adam, J.N. Bardsley, C. Bibeau, A.E. Brown, D. Browning, D.R. Burns, J.A. Caird, P. Chaffe, B. Cline, L. Davis, B. Demaret, J.J. DeYorco, D. Eimerl, J.L. Emmett, R.B. Engle, B.A. Hammel, J.A. Hatch, J.S. Hayden, M.A. Henesian, D.W. Hughes, R.S. Hughes, K.S. Jancaitas, D. Johnson, C. King, K. King, M.K. Kong, D. Kroepfl, W.L. Kway, J.K. Lawson, R.W. Lee, E.M. Lilley, E. Lindsey, P. Ludwigsen, G. Mant, J.E. Marion, S.E. Mayo, R.W. McCracken, V. Miller, D. Nikkel, E.H. Padilla, F. Patterson, H.G. Patton, E.L. Pierce, G.G. Pollock, H.T. Powell, E.R. Prochnow, F. Rainer, R. Riebold, J.B. Richards, D.H. Roberts, G.F. Ross, R.L. Saunders, M. Sawvel, D. Schwinn, M.W. Scrivener, L. Seppala, M.S. Shinn, J.R. Smith, A.A. Tesar, P.M. Test, C.E. Thompson, C.B. Thorsness, G.L. Tietbohl, S. Townsend, G.T. Ullery, W.G. Unites, P. VonArsdall, F. Wang, M. Webb, R. Wilcox, G.D. Wilke, J.D. Wintemute, L.A. Witt, R. Wolfe, B.W. Woods, P.E. Young

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