1980 Laser Program Annual Report

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UCRL--50021-80-Vol.3

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Volume 3

1980 Laser Program Annual Report

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Publication Editor: John R. Strack
MS date: June 1981

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Contents

Volume 1

Section 1—Laser Program Overview

Introduction	1-1
Summary of Major Activities	1-4
Lead Laboratory Program for 1.06- μ m and Shorter-Wavelength Target Experiments	1-11
Program Resources	1-13
Program Facilities	1-14

Section 2—Solid State Laser Systems and Technology

Introduction	2-1
Operating Systems	2-2
Novette	2-38
Nova	2-53
Research and Development	2-184
Basic Research	2-304

Volume 2

Section 3—Target Design

Introduction	3-1
Ignition and Scientific Breakeven	3-3
A Calculational Comparison of D-T, Neon, and Argon Implosions	3-11
Charged-Particle Fusion	3-17
Plasma Theory	3-24
LASNEX and Atomic Theory	3-71

Section 4—Target Fabrication

Introduction	4-1
Glass-Sphere Development and Production	4-3
Coatings and Layers	4-9
Double-Shell Technology	4-29
Characterization and Analysis	4-34
Cryogenic Systems Development	4-50
Target Factory Studies	4-60
Fusion Target Development Facility (FTDF)	4-63

Section 5—Diagnostics

Introduction	5-1
Three-Channel Soft X-ray Streak Spectrometer	5-2
Pinhole Imaging Streak Camera	5-5
Multilayer X-ray Interference Mirrors for an Improved SXRSC Spectrometer	5-9
Target Spatial Definition for Low-Energy X-ray Spectral Measurements	5-10
Production and Characterization of Soft X Rays for Diagnostics Calibrations	5-13
Hyper-Filter-Fluorescer Spectrometer for X Rays Above 120 keV	5-15
Streak-Camera Recording of Simultaneous Optical and X-ray Signals	5-19
A 662-keV X-ray Calibration Facility	5-21
22 $\times$ Microscope Streak Camera	5-23
Microchannel Plate Reflector	5-31
Fuel Density Measurements Using Bromine L X-ray Line Imaging and Stark Profiles	5-32
Fabrication of High-Aspect-Ratio Microstructures for High-Energy X-ray and Particle-Imaging Applications	5-34
Point Imaging with X Rays (PIX)	5-36
Evaluation of X-ray Backlighting Diagnostics for Nova and Novette Laser Targets	5-38
Direct Fuel (ρR) Measurements by Neutron Activation	5-40

Debris Collection and Background Measurements with Bromine-Filled Targets	5-45
Automated Radiation-Chemistry Diagnostic System	5-49
Stimulated Brillouin and Raman Side-Scattering Diagnostics System	5-52
4 ω Optical Probing	5-53
Analysis of Bremsstrahlung Emission as an Electron-Density Diagnostic	5-56
CCD Applications in ICF Diagnostics	5-62
Automated Streak-Camera Calibration Using a CCD Readout	5-67
Optically Driven Semiconductor (Auston) Switch	5-68
A Fast Optically Triggered Silicon Switch for Framing-Camera Applications	5-70
Novette Target Alignment and Viewing Optics (TAV)	5-72
Section 6—Management and Analysis of Experimental Data	
Introduction	6-1
Fusion Experiments Analysis Facility	6-3
Beam Spatial Profile Analysis	6-7
Dante Temporal Unfolds	6-11
Pyroelectric Detector Data Processing	6-15
Raman Spectrograph Data Processing	6-18
Crystal Spectrograph Data Analysis	6-21
CCD Streak-Camera Processing	6-22
Novette Diagnostic-Configuration Data Base and Graphic Presentation	6-28
Target Shot Books and Archiving	6-32
Section 7—Laser Fusion Experiments and Analysis	
Introduction	7-1
Target-Interaction Experiments at 0.53 μm and 0.35 μm	7-2
Analysis and Summary of Wavelength-Scaling Experiments and Comparisons with Simulations	7-22
Stimulated Raman Scattering in Large Plasmas	7-35
Space-Time-Resolved X-ray Measurements from Accelerated Disk Targets	7-44
Backlighter Characterization Experiments	7-53
Angular Distribution of High-Energy X Rays from Shiva-Irradiated Targets	7-65
Ultra-High-Pressure Laser-Driven Shocks in Gold	7-69

Volume 3

Section 8—Advanced Lasers	
Introduction	8-1
Pulse Compressor Systems	8-3
Laser Theory and Design Introduction	8-32
Advanced Concepts	8-55
Advances in Atomic and Molecular Physics	8-70
Section 9—Energy and Military Applications	
Introduction	9-1
Accomplishments Associated with the HYLIFE Chamber Concept	9-4
A New Liquid Metal Chamber of High Repetition-Rate-Operation: the JADE Chamber	9-91
Ultra-High-Pressure Shock-Wave Experiments	9-99
Section 10—Advanced Isotope Separation	
Introduction	10-1
AIS Facilities	10-5
Laser Research and Development	10-17
Atomic-Vapor Physics	10-30
Computational Facilities	10-47
Alternative Applications of Technology	10-51
Planning Methodology	10-58

Section 8 Advanced Lasers

Advanced Lasers

Contents

Introduction	8-1
Pulse Compressor Systems	8-3
Beam Quality in Raman Amplifiers	8-4
System Experiments and Technology Developments in the RAPIER System	8-12
RAPIER B-Amplifier Pulsed-Power Development	8-17
Advanced Architectures for Raman Compression	8-20
A Method for Suppressing Second-Stokes Parasitics	8-23
Pulse-Stacker Laser Concepts	8-26
Summary of Stacker-Compressor Technical Issues	8-29
Laser Theory and Design Introduction	8-32
Rare-Gas-Halide Kinetics Experiments	8-32
Rare-Gas-Halide Kinetics Theoretical Studies	8-44
Dissociative Electron Attachment and Vibrational Excitation of F ₂	8-52
Advanced Concepts	8-55
Free-Electron Lasers	8-55
Energy-Storage Solid State Lasers—V:MgF ₂	8-60
Chemically Pumped Iodine Lasers	8-67
Advances in Atomic and Molecular Physics	8-70
Theoretical Studies of the Electronic States of CF ₃ I	8-71
Photon-Atom Interaction Modeling	8-72
Effects of Strong Laser Fields on Chemical Reactions	8-73
Collisionally Pumped X-ray Laser Studies	8-75
Quasiclassical and Quantum Wave-Packet Studies of Reactive Scattering	8-77
New Theoretical Methods in Electron Scattering	8-79
TAMP Remote Computational Facility	8-81
References	8-83

Advanced Lasers

Introduction

The objective of the LLNL advanced laser effort is to develop a short-wavelength laser driver system capable of achieving the ICF milestones of high pellet gain ($G > 50$) and power production. To be an effective laser driver, the selected system must have high efficiency and low installed cost. For efficient target coupling, the laser should operate at short wavelengths. The specific technical performance requirements placed upon the laser driver system have evolved from new information obtained in target irradiation experiments, computer simulations, and laser and reactor systems analysis. Our current understanding of laser/target coupling and implosions physics suggests an approximate set of performance parameters for fusion-reactor drivers that include several megajoules of energy, a pulse width of 5 to 20 ns (see Section 3), and an overall efficiency in excess of 5%.

Several advanced laser concepts we have been pursuing are illustrated in Fig. 8-1. A laser system based on the use of an advanced, flashlamp-pumped solid-state medium with high thermal conductivity would have the inherent advantages of an energy storage laser, and would build upon the highly developed technology of glass lasers. Two techniques, pulse stacking and Raman pulse compression, may allow efficient nonstorage laser media such as KrF to be used as high-power laser systems.

Based on extensive evaluation of various advanced laser candidates conducted over the past several years, we have concluded that the efficiency, scalability, and short wavelength of the rare-gas-halide (RGH) lasers combine to make them at present the most attractive advanced-laser candidate for ICF applications. System analyses and small-scale experiments indicate that an RGH laser system with 3 MJ of energy and 3 to 5% efficiency could be built at a current-dollar cost of \$300/J. Our programmatic activities during

Fig. 8-1

Introduction

1980 reflect this conclusion, in that our research efforts emphasized RGH lasers and the pulse-stacking and Raman-compression techniques required to make them suitable for fusion applications.

The second, complementary element of our program is the continuation of a vigorous search for alternative laser concepts that offer the possibility of higher efficiency, better performance, and lower cost than the pulse-compressed RGH laser system.

In 1980, DOE designated LLNL the Lead Role Laboratory for Advanced Lasers. In this assignment LLNL provides technical management of the Advanced Lasers Program sponsored by the DOE Office of Inertial Fusion (DOE/IF); LLNL's responsibilities will include program planning, technical direction of external contracts, and coordinating the advanced-laser research efforts of LLNL, Los Alamos, and Sandia. The Nevada Operations Office of DOE will administer that portion of the Advanced Lasers Program conducted outside the national laboratories.

Our main planning task is to generate and maintain a plan for the development and evaluation of advanced laser drivers, consistent with the resource envelopes and overall ICF program milestones set by DOE/IF. LLNL will provide technical management and direction within the framework of the approved program plan. An Advanced Lasers Program Management Office (PMO) has been established at LLNL, to perform the functions of the Lead Role Laboratory; required technical support will be drawn from programs within the nuclear-weapons laboratories and from industry and universities. A Program Coordinating Committee (PCC), composed of representatives from LLNL, Los Alamos, and Sandia, will assist the PMO in program planning and coordination of on-going advanced-lasers programs at the three laboratories.

The objectives of the Advanced Lasers Program are to identify, evaluate, and develop advanced laser drivers for future ICF applications. Specific objectives include

- Development and construction of a laser driver suitable for use in a high-gain test facility (HGTF), intended to demonstrate inertial fusion in the $G > 50$ regime.
- Development of commercially viable drivers for power reactors.

Based on budgetary and overall Program guidance provided by DOE/IF, the PMO prepared a five-year Advanced Lasers Program plan, adopting the pulse-compressed RGH laser as the baseline approach to achieving program objectives. The plan was approved by DOE/IF and presented to prospective private-sector participants in a meeting held in Las Vegas, Nevada in December 1980. The Program plan calls for the demonstration of a fully integrated short-pulse 10-kJ RGH laser system, and the development of a 100-kJ power-amplifier module suitable for use in a full-scale RGH laser system. The plan also provides for continued exploratory assessment of advanced laser concepts with superior cost/performance potential relative to the baseline approach.

In this section we summarize the results of advanced-lasers research conducted at LLNL in 1980. Our report begins with a description of results in RGH pulse-compressor research and development. We discuss basic experiments conducted with Raman amplifiers and integrated experiments on the RAPIER laser test bed. We also give a summary of activities involving large-scale RGH amplifier development, and describe some advanced concepts for Raman pulse compression and RGH pulse stacking. In the Laser Theory and Design subsection, we describe our supporting research on RGH kinetics.

This includes both small-scale experimental studies aimed at improving our current understanding of RGH laser performance, and theoretical efforts in the kinetic modeling of KrF, XeCl, and ArF lasers. We also discuss theoretical studies of several electron-F₂ collision processes that impact the kinetic modeling.

In the articles on Advanced Concepts, we present three advanced-laser candidates. We update our reporting on the free-electron laser (FEL) and the advanced solid-state V:MgF₂ laser, both of which were discussed in last year's annual report, and present the chemically pumped iodine system. Finally, in "Advances in Atomic and Molecular Physics," we describe several developments in theoretical and computational physics and quantum chemistry. These developments have emerged as a by-product of our programmatic activities, and broaden our ability to describe the microscopic physics of laser media.

Authors: E. V. George, L. D. Pleasance, and T. N. Rescigno

Pulse Compressor Systems

Rare-gas-halide (RGH) pulse compressor systems are currently the leading advanced laser concept for fusion drivers; as such their study comprises a major part of our advanced laser experimental and analytical program. In the LLNL program, we have emphasized systems which use backward-wave Raman scattering as a means of achieving a significant degree of pulse compression. Our major experimental facility for this purpose is the RAPIER (Raman Amplifier Pumped by Intensified Excimer Radiation) KrF laser system, on which we perform small-scale testing of the physics and engineering issues of stacker-compressor systems, and develop technology and components for UV electron-beam-pumped lasers.

In 1980 we used these facilities for a detailed study of pump-beam intensity-averaging and propagation issues in Raman amplifiers; the results of these studies are discussed in the next article. We integrated the RAPIER KrF discharge laser "front end" with the 25-J, 50-ns-pulse e-beam-pumped A-amplifier for a first series of full-system stacker-compressor experiments; these developments are described in "Systems Experiments and Technology Developments with the RAPIER System." We also continued development of pulsed-power components which could drive a

B-amplifier at the 500-J, 150-ns level; this would facilitate study of the technological issues attendant on e-beam-pumped KrF lasers with multiple pulse-forming lines and diodes, and with a laser volume large enough to approach the medium conditions expected in very large KrF lasers. These efforts are detailed in "RAPIER B-Amplifier Pulsed-Power Development."

Our initial system experiments and analyses have dealt with "linear" laser chains. Certain multipass configurations of greater conceptual complexity may offer significant savings in cost or improvements in extraction efficiency, however, and we have begun to study such concepts as applied to RGH stacker-compressor architectures. These analyses are presented in "Advanced Architectures for Raman Compression." Second-Stokes suppression in Raman amplifiers is also a significant issue; progress in that area is discussed in "A Method for Suppressing Second-Stokes Parasites."

Pulse stacking represents an alternative generic approach to RGH pulse compression. We have consequently carried out extensive studies of various pulse-stacking schemes, and summarize the key results of these analyses in "Pulse-Stacker Laser Concepts." The last article in this subsection, "Summary of Stacker-Compressor Technical Issues," discusses the trade-offs that must be addressed when comparing backward-wave Raman compression and pulse stacking.

Authors: J. R. Murray and R. A. Haas

Beam Quality in Raman Amplifiers

The Raman amplifier energy-extraction studies which we reported last year showed that energy extraction, gain, and second-Stokes limits are consistent with predictions.¹ We conducted these measurements using highly filtered uniform-plane-wave Stokes and pump beams in a geometry with a very low Fresnel number, thus eliminating the complications introduced by diffraction and by spatial variation in the beam intensity profiles. We extended these studies this year to amplifiers with large Fresnel numbers and beams with significant intensity and phase noise. These amplifiers have much more complex properties than do one-dimensional amplifiers using exactly counterpropagating uniform plane waves, but they are also more representative of the issues which must be addressed in the design of large devices.

Relay Imaging of the Stokes Beam in a Raman Amplifier Chain. In sophisticated neodymium laser designs, we use lens pairs to relay an aperture image through the amplifiers.² This gives a more uniform intensity profile in the amplifier, which improves performance.³ Relay imaging gives a more uniform intensity profile by preventing the growth of diffraction patterns as the laser pulse propagates through the chain. The reduction in the ratio of peak to average intensity in the beam profile gives lower peak values of the B integral,³ and less distortion due to nonlinear refraction. The lower ratio of peak to average fluence also allows the amplifiers to operate at higher average fluence before they suffer damage at points of peak fluence.

Raman amplifiers are subject to very similar limitations.⁴ The growth of a parasitic second-Stokes wave is proportional to the integral of first-Stokes beam intensity over the length of the amplifier. This is also true for the nonlinear phase distortion B integral, and optical damage in the UV is at least as much of a problem as for glass lasers. Therefore the performance of Raman amplifier chains should benefit significantly if the Stokes beam is transported by image relays.

Figure 8-2 shows an interstage relay chain specialized for the conditions present in Raman amplifiers. The black lines represent the intensity profile of the beam propagating through the amplifiers; note that the figure is not drawn to scale. The output end of an amplifier has the highest Stokes intensity, and therefore also has the highest contribution to second-Stokes gain and the highest potential for window damage. The most uniform intensity profile is at the image plane, so we position this plane inside the Raman amplifier near the output window. The interstage relay then images this plane onto a similar position in the following amplifier stage. The beam expansion in the relay is chosen to give the desired input fluence to the following stage. Interstage isolation against second-Stokes radiation is integrated with the relay lenses. We are presently using a "hard" input aperture as shown in this figure. This gives noticeable diffraction fringes in the low-intensity parts of the Raman amplifier. The sharp changes in intensity at the edge of the beam also enhance self-focusing effects. At present our intensities are too low to see any problems from hard apertures, but we expect that "soft" beam shaping will be required in large, high-fluence amplifiers as it is in glass lasers.

Fig. 8-2. A Raman amplifier linear chain with relay imaging.

Spatial Averaging of Pump Beam

Nonuniformities. A simple backward Raman amplifier with exactly counterpropagating Stokes and pump waves (Fig. 8-3) is very sensitive to nonuniformities in the pump intensity profile.⁵ The small signal gain for the Stokes wave is proportional to the integral of pump intensity seen by the Stokes wave as it propagates through the amplifier. This means that transverse intensity variations in the pump beam profile are exponentiated in a small-signal amplifier, and become much more severe in the Stokes beam than they were in the pump. The sharp Stokes intensity spikes which develop as a result will also diffract and disrupt an initially smooth Stokes-beam wave front, as shown in Fig. 8-4 (Refs. 4 and 5). The high local intensity in these spikes also promotes conversion to second-Stokes in the spikes at a much lower average Stokes intensity than for a smooth Stokes beam.

In theory the conversion of a few intensity spikes to second-Stokes parasitics is tolerable, but we observe in practice that small zones converted in this way diffract sufficiently to destroy a larger region of the Stokes beam than the size of the original hot spot. This means that there is a rather limited range over which we can operate an amplifier when second-Stokes production becomes significant in the hot spots of the beam. Saturated amplifiers show somewhat less severe perturbations, since the Stokes intensity growth becomes linear rather than exponential with regard to pump intensity. Nevertheless, they still require pump beams of very uniform intensity when using an exactly counterpropagating geometry. The practical realities of large excimer laser devices make it difficult to meet this condition without taking severe penalties in cost, efficiency, or both. This prompts us to study other Raman amplifier geometries which do not put such severe constraints on pump-beam quality, so that the KrF pump lasers can be designed for lowest cost and highest efficiency.

For a Raman amplifier to average out pump nonuniformities, the propagation directions of the Stokes and pump beams must be set at a slight angle to each other, so that a point on the Stokes beam sweeps across the pump wave front and averages out the pump intensity fluctuations. This does not affect the physics of the Raman

amplifier. The pump beam appears to the Stokes beam only as a stored energy density. Its exact direction of propagation is not important if the direction is generally backward with respect to the direction of propagation of the Stokes beam, although some parametric coupling effects must be considered for forward propagation.

The degree of averaging is improved if one uses multiple pump beams at slightly different angles, as shown in Fig. 8-5. A Stokes beam propagating through this array encounters large local variations in pump intensity from the nonuniformities and interference of the separate beams. The pattern of intensities varies rapidly with position down the cell as the Stokes beam propagates, however, and its average gain becomes smooth. Amplifiers in which the Stokes and pump beams are not exactly counterpropagating also allow the beams to

Fig. 8-3. A one-dimensional Raman amplifier with exactly counterpropagating pump and Stokes pulses.

Fig. 8-4. Schematic of the mechanism by which high local gain can destroy the Stokes phase front.

Fig. 8-5. A Raman amplifier pumped by multiple pump beams at an angle to the Stokes axis.

be separated without the use of dichroic optics. This is a significant practical advantage.

We have studied the practical performance of a Raman amplifier which is pumped by multiple pump beams at slight angles to the axis, using the KrF discharge lasers of the Rapier front end.^{1,3} The output of a KrF discharge laser typically shows severe modulation in both intensity and phase due to discharge nonuniformities,⁶ so this is a good test of the averaging of pump beams of poor spatial quality. We show the geometry of the amplifier in Fig. 8-6. Three independent pump beams each about 1 cm in diameter cross in a 2-m-long methane-filled Raman amplifier at angles of about 3 mrad. A Raman oscillator generates a Stokes probe pulse, and using a hard aperture we select from the pulse a small area of uniform intensity. We relay this area (as in Fig. 8-2) to the output end of the Raman amplifier cell. The Stokes beam diameter in the Raman cell is 7 mm. This image plane is then further relayed and expanded onto photographic film for a study of the uniformity of Raman gain in the amplifier. We were able to vary the pump beam intensities and the methane pressure in the cell to give a Stokes small-signal gain between e^1 and e^6 . We always adjusted the input Stokes intensity to keep the amplifier in small-signal gain conditions, since this is the most critical case for averaging pump intensity variations.

Fig. 8-6. Geometry of a three-beam-pumped Raman amplifier experiment on RAPIER.

Figure 8-7 is a photograph of the three overlapping pump beams as they transit the Raman amplifier. The severe nonuniformities in intensity distribution are obvious. Figures 8-8 and 8-9 show the profile of the Stokes beam at the image plane in the amplifier, with (Fig. 8-8) and without (Fig. 8-9) a small signal gain of 5 Np (i.e., e^5). The fringe patterns in Fig. 8-8 are from etalon effects in windows and filters; the beam diameter is 7 mm. There is some variation in gain between the center and edges of the amplifier. This is due to regions around the edges of the gain volume that are not illuminated by all three beams for the entire length of the amplifier; these regions see a smaller average pump intensity. It is clear, however, that the intensity profile of the amplified Stokes wave is not severely degraded by the gain in the amplifier, and is far more uniform than the pump intensity distribution. Note that the violent intensity modulations on the pump beams (Fig. 8-7) do not transfer to the Stokes wave.

We measured the far-field intensity distribution of the pump and Stokes beams to verify the absence of phase distortion in the amplifier. The pump beams exhibit a phase distortion perpendicular to the discharge axis, which is typical of excimer discharge lasers. Also, they are perhaps three to five times the diffraction limit of the pump laser aperture in this dimension. The Stokes beam, on the other hand, showed a far-field pattern that is equal to the

Fig. 8-7. Intensity profiles of the three pump beams as they propagate through the amplifier of Fig. 8-6.

diffraction limit of a 7-mm aperture to within the precision of measurement. The measurements of amplitude and phase profiles were repeated over a range of conditions, with similar results.

We also studied, for comparison, the performance of a Raman amplifier with a single pump beam propagating at an angle of about 1.5 mrad to the Stokes beam axis. This configuration gives some averaging of pump hot spots in one dimension, but gives much less overall averaging than the amplifier with three pump beams. Pump beam hot spots were smeared into lines in the plane of crossing, but the noise impressed on the Stokes beam was still unacceptably high for small-signal gains above about 3 Np. Above that gain we began to see diffraction of the small hot spots impressed on the Stokes beam by the mechanism of Fig. 8-4, with a corresponding loss of Stokes beam quality. This demonstrates the expected result: pump beams which are too noisy for acceptable performance in a single-beam-pumped amplifier at a small angle relative to the Stokes beam can nevertheless give acceptable performance in a multibeam-pumped amplifier where intensity averaging of the beams is more effective.

Light Channel Cells for Raman Amplifiers.

The Raman amplifier experiments we have described above demonstrate the power of averaging many noisy pump beams to produce a clean amplified Stokes beam. The crossing beam geometry we used, however, has certain limitations. The pump beams overlap well only over a fairly short region, unless their propagation angles with respect to the Stokes axis of propagation are very small. This requires very long propagation

Fig. 8-8. Intensity profile of an amplified Stokes pulse in the amplifier of Fig. 8-6, pumped to a small-signal gain of 5 Np (150 $\times$).

Fig. 8-9. Intensity profile of the Stokes pulse in the amplifier of Fig. 8-6, without gain.

distances, to separate the Stokes beam from the mirrors which direct pump beams into the cell. There is also a nonlinear four-wave mixing effect, which can become important for small crossing angles if the gain per unit length in the cell is high. We discuss this further below.

We can overcome these limitations by crossing the pump beams at a somewhat

larger angle, and redirecting them back through a cell with reflective walls when they separate, as shown in Fig. 8-10.

Figure 8-11 shows a view through such a cell used in RAPIER experiments. This cell consists of two uncoated glass plates facing each other. Rays propagating by multiple reflections from the plates give multiple images of the sword at the end of the cell. The reflectivity of even bare cell walls can be very high for large angles of incidence, so that pump propagation by reflection from the walls does not introduce a large loss. If the angle of incidence is not large enough to give high reflectivity on a bare surface, the wall reflectivity can be enhanced by metal or dielectric coatings. The theoretical reflectivity for s-polarized light incident on a glass plate at an angle of α rad with respect to the surface is about $1-4\alpha$. This agrees with our experimental measurements of pump propagation at 248 nm in the glass light-channel cell shown in Fig. 8-11.

Grazing-incidence light channels such as these have a long history in the laser literature,⁷⁻⁹ as well as in infrared and x-ray optics. Along with the present stacker-compressor application, they are also under consideration¹⁰ for coherent summation of poor-quality neodymium laser beams in a Raman amplifier. The earliest Raman light-channel amplifier⁷ actually used walls coated with an efficient Lambertian scattering surface rather than a reflecting surface. This gave a very uniform pump intensity, at some cost in transport efficiency.

The cell of Fig. 8-10 has another interesting and useful property for the design of large laser systems. The entire pump

beam can be extracted by a Stokes beam which is considerably smaller than the aperture of the cell, since the pump is reflected back through the center of the cell several times. This allows an excellent fill factor for extraction of energy from the pump into the Stokes beam. At the same time, the Stokes intensity profile can vary smoothly to zero at the edge, avoiding diffraction from abrupt intensity changes. This cannot be done in a conventional laser amplifier if efficient extraction is required, because extraction is poor at the perimeter of the beam where the intensity is low. In the Raman light-channel cell the pump energy which is not extracted in the low-intensity outer regions propagates back through the center of the Stokes pulse and is extracted there.

A further property of light-channel cells may prove useful in stabilizing the output Stokes pulse shape in the presence of pulse-to-pulse fluctuations in the pump and Stokes input signals. A large fraction of the pump energy in such a cell is outside the region that is extracted directly by the Stokes pulse. This has a stabilizing effect on the Stokes pulse width, since the reservoir of pump energy outside the directly extracted region takes a finite time to find its way into the extraction zone. The Stokes pulse will stabilize at a width given by the time required for a pump photon to propagate from the edge of the intense Stokes extraction zone to the wall, and then back to the edge of the extraction zone.

A potential problem that must be taken into account in the design of light-channel cells is conversion of individual pump beamlets to forward first-Stokes radiation

Fig. 8-10. A Raman amplifier pumped by multiple beams in a light-channel cell.

propagating exactly coincident with the pump. A cell with a cylindrical reflective wall, for example, tends to produce a line focus for each pump beamlet. The very high forward Raman gain for propagation through these foci can convert a significant part of the pump energy to forward first-Stokes radiation, which is wasted. This probably restricts practical light-channel cells to two-sided or four-sided parallel plate geometries. The cell walls need not be of particularly high optical quality, but concave warping which produces line foci can lead to optical damage or forward first-Stokes energy conversion, and must be controlled.

Detailed simulations are required for more precise studies of the properties of light-channel amplifiers. Their behavior is considerably more complex than amplifiers in which the extractable energy is stored in stationary excitations rather than in propagating photons. This is particularly true when they are integrated into regenerative designs such as the Raman accumulator concept discussed in "Advanced Architectures for Raman Compression" later in this section.

We studied the performance of a light-channel cell in the same geometry used to study the three-beam-pumped cell described earlier in this article. An image of the Stokes input aperture is relayed to the output end of a light-channel cell formed by two 3-m-long glass plates separated by 12.5-mm spacers. Two 10- by 10-mm KrF pump beams enter the cell at an angle of about 20 mrad with respect to the axis. They then propagate by grazing-incidence reflection between the glass plates. Each beam makes five transits through the Stokes beam before exiting the cell. Five wall reflections at this angle produce very good pump beam averaging, but give a fairly low pump transmission coefficient (about 0.66) through the cell, and an average pump intensity in the cell of about 0.82 of the input. This is acceptable for a preamplifier, but unacceptable for a saturated final amplifier in which efficiency is important. Thus these saturated stages must use either fewer reflections or more highly reflective walls.

Figure 8-12 shows the profile of the input-beam to this cell, while Fig. 8-13 shows the amplified beam output at a small-signal gain of 3 Np. The beam diameter is 8 mm, and the beam-limiting aperture is a

decagon, not the round aperture used for Figs. 8-8 and 8-9. The fringe patterns are from etalon effects in windows and filters; note the Airy patterns from three small damage spots. The averaging is similar to what we observe in the three-beam amplifier, although the light-channel cell has superior uniformity due to the more uniform integral of pump intensity down the light channel. The intensity profile of Fig. 8-12 is very well preserved.

Four Wave-Mixing Effects in Raman Amplifiers with Multiple Pump Beams. Each pair of pump waves which overlap in multiple-beam amplifiers can interfere to produce a three-dimensional array of interference fringes in the pump intensity distribution. These appear to the Stokes wave as a nonuniform gain, and impose an intensity modulation on the amplifier beam which diffracts to form secondary Stokes waves at the first-Stokes frequency, propagating at an angle to the primary first-Stokes wave. For certain specific phase-matched relations between the wave vectors of the two pump waves and the primary and secondary Stokes waves, the secondary waves can grow over long distances. In general, however, the secondary waves quickly get out of phase with the grating and soon average to zero if the pump waves are coherent.

This is merely a reexpression in different terminology of the pump intensity averaging discussed above. If the pump waves are not perfectly coherent with respect to each other, however, the cancellation of the non-phase-matched secondary Stokes waves is no longer exact, and they will be observed in the output from the amplifier. The energy in these waves is generally small; it can become significant, though, if the angles between the

Fig. 8-11. The view seen through a 5-cm-aperture light-channel cell used in RAPIER experiments.

pump and Stokes waves are very small and the gain per unit length in the amplifier is large. We have observed these waves in the three-beam-pumped amplifier shown in Fig. 8-6. We have also compared their intensity and propagation to theoretical models, in order to determine those conditions in a Raman amplifier under which secondary Stokes waves could be a significant loss mechanism.

We consider first the phase-matched case in which the secondary Stokes wave is expected to grow. This is a nondegenerate case of the degenerate four-wave mixing

Fig. 8-12. Intensity profile of the Stokes pulse in the amplifier of Fig. 8-10, without gain.

Fig. 8-13. Intensity profile of an amplified Stokes pulse in the amplifier of Fig. 8-10, pumped to a small-signal gain of 3 Np (20×).

Fig. 8-14. Phase-matched geometry for nondegenerate four-wave mixing in a backward Raman amplifier.

geometry first proposed by Hellwarth^{11,12} for phase conjugation. The secondary Stokes wave is also a "quasiconjugate" of one of the pump waves, although it cannot be an exact conjugate since its direction of propagation is not exactly back along the pump wave input direction. The phase-matching condition for four-wave mixing is simply that momentum must be conserved, i.e.,

$$\vec{k}_p + \vec{k}_p' + \vec{k}_s + \vec{k}_s' = 0$$

where $\vec{k}_p$ and $\vec{k}_p'$ are the two pump wave vectors, $\vec{k}_s$ is the primary Stokes wave vector, and $\vec{k}_s'$ is the secondary Stokes wave vector. It can be shown that this phase-matching condition is equivalent to the statement that the scattered secondary Stokes wave $\vec{k}_s'$ does not get out of phase with the grating formed by the interference of the two pumps.¹²

Figure 8-14 shows the vector addition of the four vectors $\vec{k}_p$, $\vec{k}_p'$, $\vec{k}_s$, and $\vec{k}_s'$ to give zero in a phase-matched geometry with counterpropagating pump and Stokes waves. (We ignore here the phase-matched geometry in which pump waves are allowed to propagate in both directions, although that geometry may be germane for regenerative systems.) The geometry of the phase-matching diagram shows that the angles α and β satisfy $\beta = (|\vec{k}_p|/|\vec{k}_s|)\alpha$, and the angle γ between $\vec{k}_p$ and $\vec{k}_s$ is $(\beta - \alpha)/2$. This gives $\beta = 1.07\alpha$ and $\gamma = 0.035\alpha$ for KrF-pumped methane, which is to say that for values of α of a few mrad, as in our experiments, the phase-matched case requires $\vec{k}_p$ and $\vec{k}_s$ counterpropagating to within about 100 μ rad.

Consider now a case which is not phase-matched, as shown in Fig. 8-15. We draw here a case in which the Stokes wave vector is parallel to the axis of the amplifier; the two pump waves cross it at approximately equal angles, as in the experiment of Fig. 8-6. Here there can be two waves formed, $\vec{k}_s'$ and $\vec{k}_s''$, which lie in the plane containing the two pump wave vectors and which propagate at angle β with respect to $\vec{k}_s$. For small angles, $\beta \approx (|\vec{k}_p|/|\vec{k}_s|)\alpha$, and $|\vec{K}| \approx (\beta^2/2)|\vec{k}_s|$, where $|\vec{K}|$ is the wave vector mismatch. As the waves $\vec{k}_s'$ and $\vec{k}_s''$ propagate, the contribution from one point in space will cancel the contribution from a point further on that is π rad out of phase,

that is, $|\vec{K}|L_D = \pi$, where L_D is a dephasing length for the grating. The geometry becomes slightly more complex in three dimensions, but for small angles the result is the same. Each pair of pump waves will interact separately, so for N pump waves we expect $N(N-1)$ secondary waves at angles of $\pm\beta_N$ with respect to the primary Stokes wave.

If the phase of a pump wave changes over the distance L_D , the cancellation of the secondary waves is no longer complete and they can be detected in the output from the amplifier. We might model this by defining a pump autocorrelation length L_C such that the pump electric field $\mathcal{E}(r, t)$ obeys

$$\int \mathcal{E}(r, t) \mathcal{E}^*(r, t') dt = \mathcal{E}_0^2 e^{-(|t-t'|/L_C)}$$

We can then find a value for F , that fraction of the output energy of a backward Raman amplifier appearing as secondary waves, assuming an amplifier gain per unit length g Np/cm:

$$F = \frac{2N(N-1)g^2L}{N^2L_C \left[\left(\frac{\pi}{L_D} \right)^2 + \left(\frac{4}{L_C} \right)^2 \right]}$$

Here L is the length of the amplifier and N is the number of pump beams (all assumed to be of equal intensity, dephasing length, and correlation length). This assumes that all of the waves are present simultaneously in the entire amplifier volume. F will be smaller if there are regions where some waves are not present. Note that this expression for F has the proper behavior in limiting cases. For $L_C \rightarrow \infty$ (perfectly coherent pump beams), $F \rightarrow 0$ unless $L_D \rightarrow \infty$ (perfect phase-matching), for which case F grows without limit. Also for $L_C \rightarrow 0$ (perfectly incoherent pump beams) we can form no recognizable gratings and $F \rightarrow 0$.

We now compare these predictions to experimental results. Figure 8-16 shows the far-field pattern of the three pump beams incident on the amplifier of Fig. 8-6. Figure 8-17 shows the far-field pattern of the Stokes output. Note that both film images were severely overexposed to ensure that all scattered radiation would be detected. The size of the far-field spots in these

photographs is therefore not representative of the true size of the far-field images. In particular, the central primary Stokes beam is about 100 times as intense as any of the secondary beams in this particular image. We observe in Fig. 8-17 the six secondary Stokes beams predicted by this analysis, each at the proper angle as given by the three pump beams.

Figure 8-18 shows the fraction F of the primary Stokes beam which appears in this case in secondary beams as a function of the square of the small signal gain in Np/m. The value of F is proportional to g^2 , as predicted, and is not affected by parameters such as methane pressure. There is a considerable scatter in the data, but we do not find this surprising; the pump autocorrelation length is affected by discharge instabilities that do not reproduce very well from shot to shot.

Fig. 8-15. Non-phase-matched geometry for four-wave mixing with pump waves propagating at equal angles to the Stokes axis in a backward Raman amplifier.

Fig. 8-16. Far-field pattern produced by the three pump beams of the amplifier shown in Fig. 8-6.

The value of F is consistent with estimates of $L_D \approx 7$ cm (from the geometry of the experiment) and $L_C \approx 5$ cm (from the measured pump laser line width).

We have also looked at F for a two-beam pumped amplifier aligned to the phase-matched condition of Fig. 8-14, which implies $L_D \rightarrow \infty$ in the expression for F . The fraction scattered into the secondary wave was up by about a factor of 2 from a non-phase-matched geometry with $\alpha \approx 2$ mrad. This, again, is consistent with estimates of L_C . Four-wave mixing experiments in the past have typically been conducted using media length $L \ll L_C$, so the L_C limit on conversion has not been noted. For long interaction lengths, however, it must be taken into account.

We have demonstrated here that non-phase-matched four-wave mixing effects can be detected in Raman amplifiers, and that they can be significant in small preamplifiers when the gain per unit length is high and

multiple pump beams cross at small angles. Mixing effects will not be important in large amplifiers with larger pump crossing angles and low gain per unit length. Consider, for example, an amplifier which converts a 50-ns input pulse to a 10-ns output pulse, has a compression of 5, and with pump waves having L_C of a few cm as in the present experiment. The small-signal gain will be limited to 5 in such a device, in order to achieve reasonable extraction efficiencies. This corresponds to $g = 0.2$ Np/m, so even at the very small crossing angles of Fig. 8-18, the fraction lost to secondary waves is well below 1%.

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Fig. 8-17. Far-field pattern of the primary and six secondary Stokes beams produced by the pump beams of Fig. 8-16.

Fig. 8-18. Prediction and measurement of the fraction of the far-field Stokes energy detected in secondary beams.

System Experiments and Technology Developments in the RAPIER System

This year we operated the RAPIER front-end discharge laser and A-amplifier as an integrated system, for a first series of KrF stacker-compressor experiments. In this article we will first describe the sequence of events when the system is fired, and then discuss several subsystems in more detail.

Figure 8-19 shows the layout of the RAPIER stacker-compressor experiment. A shot begins when a trigger generator erects two Marx banks on the A-amplifier. The Marx banks then begin to charge the two water Blumlein pulse-forming lines, which produce the e-beams used to excite the laser gas. The lines are fully charged after $\sim 1 \mu\text{s}$. In the meantime a pulse from the trigger generator that has been delayed by $\sim 0.5 \mu\text{s}$ fires a trigger to charge a fast intermediate energy-storage capacitor on the β_0 -amplifier of the front end. About 200 ns later, in quick succession, this trigger pulse fires the two small lasers of the KrF oscillator and the command trigger of the β_0 -amplifier. A narrow-line-width (~ 1 GHz) output pulse from the KrF oscillator arrives at the β_0 -amplifier just as its gain begins to rise, and

▲ Fig. 8-19. RAPIER system stacker-compressor layout.

injection-locks an unstable resonator regenerative amplifier cavity on it. After this point all timing in the system is by optical delays and laser-triggered switches, rather than by electronic circuits.

The β_0 -amplifier emits a 20-ns narrow-band pulse containing about 750 mJ of energy. About 400 mJ leaves the front-end room at this time, to provide input beams to the A-amplifier stacker and four trigger pulses for the A-amplifier output switches. While the Blumlein voltage pulse is propagating towards the electron guns in the A-amplifier, part of the β_0 beam is split into three sequential 20-ns pulses of about 30 mJ each. These pulses are optically delayed so that they begin to enter the A-amplifier gain volume just as the electron guns begin the 60-ns excitation pulse. The three pulses follow each other in sequence through the gain volume, in a triple-pass geometry at slightly different angles (Fig. 8-20). Each pulse emerges with an energy of 5 to 8 J.

Figure 8-21 is a photograph showing the optical layout around the A-amplifier in operation. The output beams are 8 by 10 cm at the output of the amplifier, and are set to be slightly focusing so that they are 4 by 5 cm after a propagation distance of about 25 m. The beams reflect from a folding mirror about 12 m from the A-amplifier and then reflect from three 5- by 5-cm stacker mirrors, so that they arrive simultaneously at a 3-m-

Fig. 8-20. Optical configuration of a three-pass, three-beam optical cavity on the A-amplifier.

long Raman light-channel cell at the right of the picture.

Back in the front-end room, meanwhile, the rest of the β_0 -amplifier output fires the two discharge lasers, β_1 and β_2 , which pump the Raman part of the front end. These lasers are excited by pulsers charged to 60 kV dc

Fig. 8-21. Photography of the A-amplifier ▲ as configured for system stacker-compressor experiments.

Fig. 8-22. Laser-triggered output switch on the A-amplifier water Blumlein. ▼

and switched with very low jitter by laser-triggered spark gaps. The β_0 output is also used as a source to injection-lock unstable resonators on these two amplifiers. Part of the β_1 output is focused through a methane cell to produce a forward superfluorescent Stokes pulse, which is then amplified in two preamplifiers pumped by β_1 and β_2 . The Stokes output from the front end at this point is typically 100 to 150 mJ in a 2- to 3-ns pulse. The Stokes output propagates to the A-amplifier-pumped Raman cell with a delay chosen so that it meets the three pump pulses from the A-amplifier in the Raman cell at the right of the figure.

In last year's annual report we described a low-jitter pulser triggered by KrF-laser breakdown in high-pressure SF_6 gas.^{13,14} We showed then that the low UV breakdown threshold of SF_6 allowed us to form a conducting channel between the electrodes of a high-voltage spark gap using only a few mJ of KrF radiation. We also showed that we could switch both dc and pulse-charged high-voltage spark gaps over a very wide

range of conditions with subnanosecond jitter. As shown in Fig. 8-19, this year we have replaced many electrically triggered elements of the RAPIER system with simpler and more reliable laser-triggered equivalents.

Figure 8-22 shows a laser-triggered output switch that we have installed on the A-amplifier. We have designed the switch so that two conducting channels are formed in parallel on each of the Blumlein pulse-forming lines. This reduces the inductance of the switch, and demonstrates that laser triggering can be applied to multichannel switches. The laser-triggered switch is a much simpler device than its electrically triggered counterpart, described in last year's annual report.¹³

We find that the laser-triggered switch has fewer misfires than the electrically triggered version, but that it must be cleaned after every hundred shots or so to maintain its performance. The discharge channels are always in the same position in the laser-triggered switch; this increases the local electrode erosion and the deposits on the switch insulator near the channels. Eventually these deposits begin to cause surface breakdown along the insulator, leading to switch misfires. We suspect that this problem can be alleviated by changes in electrode material and insulator design.

We have also installed laser-triggered pulsers on the β_1 and β_2 discharge lasers in the front end. This allows us to drive these lasers with simple dc-charged pulsers and still maintain nanosecond jitter. The previous configuration used an electrically triggered fast intermediate storage system, which required more care and maintenance and was somewhat inferior in jitter and reliability. As in the A-amplifier output switch, we find that erosion of stainless steel electrodes limits their life. At present we find a life of about 5×10^4 shots, with sulfur deposits cleaned out of the switch about every 10^4 shots. Even as it stands, this is a fairly long lifetime for a spark gap switching 150 J at 60 kV, but we feel it can be improved further using tungsten electrodes and a gas mixture containing O_2 , which burns away sulfur deposits during operation. These modifications are now being tested.

We have begun to use the MALAPROP beam-propagation code,^{13,15} developed for Nova to simulate propagation in KrF

stacker systems, in order to quantify issues associated with channel cross talk, focusability, and peak fluences for optical damage. We are testing these predictions by calculating the propagation of beams from the A-amplifier and comparing the result with experimental data (Figs. 8-23 to 8-26). Our comparisons this year were qualitative, but they demonstrate the potential for this kind of interactive analysis in the design and operation of KrF laser systems.

KrF lasers are very high-gain devices. Large KrF amplifiers will therefore require a uniformly saturating input field in order to drive the loaded gain low enough to prevent

Fig. 8-23. MALAPROP calculation of the intensity distribution in a slice of an A-amplifier output beam.

Fig. 8-24. Intensity distribution (burned film) of an actual A-amplifier beam.

Fig. 8-25. MALAPROP calculation of intensity distribution in an A-amplifier output beam, through an amplifier output window with a 1% loss due to randomly distributed dust particles.

Fig. 8-26. Densitometer scan of an actual A-amplifier beam with a mildly etched and dirty output window, for comparison with Fig. 8-25.

parasitic oscillations. The usual result is that the output beam from the laser is "clipped" by the structures at the edge of the amplifier aperture. Diffraction effects become quite important in the propagation of these sharp edges, even though only a small fraction of the beam energy may be

involved. On stacker mirrors several tens of meters downstream from the amplifier, local intensity spikes climb considerably above the average intensity and force a significant reduction in the safe mirror fluence in the system. The RAPIER A-amplifier is, in fact, so limited: although our stacker mirrors have a 20-ns damage threshold of somewhat over 1 J/cm^2 , we find ourselves forced to limit the average fluence to between 0.3 and 0.5 J/cm^2 to avoid mirror damage.

Figure 8-23 shows, as an example, a MALAPROP simulation of an initially square and uniform 10-cm-wide A-amplifier beam which slowly focuses to 4 cm wide at a distance of 34 m (the entrance to the Raman cell). The peak intensity exceeds the average intensity at this point by a factor of about 1.5. This pattern is qualitatively similar to the patterns actually observed (Fig. 8-24), although the propagation distance for this particular experimental shot was somewhat shorter than in the calculation. Note the Airy patterns from damage spots.

MALAPROP also helps us explore certain other problems which arise in excimer lasers, as shown in Fig. 8-25. In this case we have calculated the intensity distribution for the same optical propagation path as in Fig. 8-23, but in this simulation we have assumed that there is a 1% loss in the output window of the amplifier, due to a haze formed by randomly spaced dust particles and fluorine etching of the window. Here we see a much less regular diffraction pattern, with a peak-to-average intensity ratio of about 2. Figure 8-26 shows a densitometer scan across a film image at a similar position in the actual experiment, with a window which had noticeable etching and dust (note the logarithmic scale). The peak-to-average ratio and general appearance are very similar to the predictions of the model. This illustrates that even fairly small amounts of the dust which accumulates in KrF amplifiers may have a significant impact on the optical design of the system. We plan to continue more detailed studies of these effects in the coming year.

We conducted a series of small-signal Raman-gain experiments in a 5-cm-aperture, 3-m-long light-channel cell pumped by the three beams from the A-amplifier stacker. These were similar to the small-scale

experiments reported in the previous article, and gave similar results. The maximum energy extracted was 1.5 J from a 7.8-J pump energy at a small-signal gain of 5 Np and an input Stokes fluence of 1.5 mJ/cm² in a 3-ns pulse. We were unable to study higher extraction conditions at the time, because the Stokes beam quality from the front end was rather poor before we installed light-channel cells and optical relays. We plan to repeat this experimental series early next year. We did verify, however, that we were able to keep the system alignment and timing well within the necessary tolerances required for full-system experiments.

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RAPIER B-Amplifier Pulsed-Power Development

The B-amplifier of the RAPIER system is intended to produce approximately 500 J of optical energy from KrF molecules excited by the interaction of an electron beam with a mixture of noble and fluorine gases. Two electron-beam machines will inject beams 25 cm high by 125 cm long into opposite

sides of a laser cell 40 cm high by 30 cm wide (Fig. 8-27). Each electron-beam machine consists of three magnetically isolated diodes driven by three cylindrical 5-Ω water-dielectric Blumlein pulse-forming lines operating at 450 kV and producing 150-ns pulses. The lines are pulse charged from separate Marx generators located in an oil-filled tank underneath the pulse-forming lines. This extremely compact system will deliver 36 kJ of electron-beam energy to the anodes of the diodes and will occupy less than 15 m² of floor space.

The B-amplifier incorporates an extension of the technology we used to construct the RAPIER system's A-amplifier. Last year's annual report¹⁶ contains a description of both amplifiers, and details the technology changes we are making to satisfy our B-amplifier performance goals. During the past year, we have tested individual components required for the B-amplifier, and we have constructed a prototype B-amplifier line that will enable us to test all of the components simultaneously at full system performance. Figure 8-28 shows the prototype line early in its assembly.

The electrical design of the Marx generators is quite conventional; they are five-stage generators with two 0.3-μF capacitors connected in parallel in each stage. The five pairs of capacitors are charged in parallel to a maximum of 100 kV

Fig. 8-27. RAPIER B-amplifier.

Pulse Compressor Systems

and then switched into series by five triggered spark gaps. We have constructed and tested two full-scale Marx generators at a charge voltage of 100 kV. One generator is shown in Fig. 8-29. When installed in the B-amplifier, these generators will be operated at a charge voltage of 85 to 90 kV. We tested the two generators separately and together to ensure that they can be operated in close proximity to each other and to tank walls without high-voltage breakdown problems. The resulting mechanical assembly is capable of storing 45 kJ in a 2.1-m³ tank.

The balanced charging system for transferring energy from the Marx generators to the pulse-forming lines has presented the greatest design difficulties. The system requires electrical connections to all three of the nested cylindrical conductors of the pulse-forming line, with sufficient isolation between the connections to withstand voltages comparable to the 450-kV output of

the machine. We have selected an approach that permits all of these connections to be made through the end of the pulse-forming line, so that the lines may be stacked vertically as well as horizontally.

Figure 8-30 shows the physical layout of the interconnection region. A portion of the charging inductance for the inner line is located inside of the inner conductor to minimize the voltage applied to the connecting cable when the machine is fired. We have tested this system in a 60-ns pulse-forming line to voltages of 500 kV, 10% above the design voltage.

We have routinely operated the output switch (described in last year's annual report¹⁶) at 500 kV, also in a 60-ns pulse-forming line. We suggested last year that using a copper-tungsten material for the electrodes in the output switch would give better performance. In fact, we found it necessary to use this material, because the damage to stainless steel electrodes during a

Fig. 8-28. Prototype B-amplifier pulse-forming line.

normal firing is sufficient to substantially reduce the hold-off voltage of the switch. The erosion of the copper-tungsten electrodes is not nearly as severe.

We will electrically trigger the output switches of the B-amplifier pulse-forming lines by applying a pulse to a midplane electrode in each switch. The pulses for triggering the three output switches will be generated by a single trigger switch that is charged by the Marx generators. In normal operation, this switch will itself be triggered by a laser pulse generated elsewhere in the RAPIER system.

To prevent possibly serious damage to parts of the B-amplifier system, we must ensure that the output switch of a pulse-forming line is triggered whenever that line is charged by its Marx generator—regardless of whether this triggering occurred in response to a fire command or not. Since the output-switch trigger system is common to three lines, and since other electromagnetic coupling may also exist between the lines, we prefer that, if one Marx fires, its two neighbors also fire. We have demonstrated that this multiple firing can be accomplished by coupling the Marx generator triggers.

To ensure that the trigger switch fires even in the absence of the laser trigger pulse, we have added a midplane electrode to this switch, using the circuit shown in Fig. 8-31. In normal operation, this midplane electrode serves no function; in the absence of a laser pulse, however, the midplane voltage will continue to increase until the trigger switch fires. This occurs even if the Marx generators fire before they are fully charged. We have tested this system on a single line and verified its effectiveness.

In operation, the electron beams are subjected to large self-generated magnetic fields in the diodes where the beams are formed. These magnetic fields deflect electrons inwards toward the beam axis. To reduce these fields to a tolerable level, we subdivided each diode so that a magnetically separate 25- by 41-cm beam is formed for each pulse-forming line. The magnetic fields still deflect the beam electrons, producing a "pin-cushion" distortion of the beam emitted by a rectangular cathode. (The transverse electric fields are larger near the corners of the cathode and more effectively oppose the magnetic fields.)

We have produced a rectangular beam of the desired dimensions, voltage, current, and pulse length on the modified PI422 machine by using a cathode shaped to correct the distortion, as shown in Fig. 8-32(a).

Although the beam is rectangular in shape, beam electrons are still deflected through significant angles, leading to losses in the foil support structure. We have measured 70% transmission of the desired beam through a

Fig. 8-29. B-amplifier Marx generator.

Fig. 8-30. Interconnection region of balanced charging system.

structure that has a 90% geometric transmission. Although this loss is tolerable, there is certainly room for improvement. Preliminary experiments indicate that we can reduce these deflection losses by further subdividing each diode horizontally, as shown in Fig. 8-32(b), and by masking the cathode to prevent emission of electrons that would be lost to the ribs of the foil support plate.

We have tested each of the major components of the B-amplifier system to the point that they can now be integrated into a single pulse-forming system. We may be able to improve the performance of some of the components with further work. Although we have dealt with some issues related to the

operation of several parallel lines, other issues can be addressed only when multiple lines are available.

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Advanced Architectures for Raman Conversion

The techniques of pulse stacking and Raman compression are the most developed generic schemes for exploiting efficient but non-storing laser media such as the rare-gas halides. Pulse stacking appears to offer the highest overall efficiency in the case where laser fluences at each point of the amplifier chain are chosen to optimize the system's extraction efficiency. Systems analysis indicates, however, that the optical complexity of a multimegajoule-class stacker system may be prohibitive, with ~1000 individual optical channels required to achieve the desired total output energy and net compression.¹⁷ Backward-wave Raman systems offer a much simpler optical design, but have the liability of an extra energy conversion step (in the Raman medium) which degrades the overall system efficiency. The individual loss processes which result in lower extraction efficiency in both the RGH power amplifier and the Raman amplifier are difficult to compensate for with the amplifiers de-coupled.

Efficient extraction from Raman amplifiers requires that the compression factor be small, in the range from 1 to 3 (Ref. 18). Compression factors much higher than this result in the growth of unwanted second-Stokes radiation. On the other hand, compression factors in the range from 1 to 3 are not large enough for our use in power amplifiers operating with excitation pulse lengths from 100 to 1000 ns.

Efficient extraction from RGH laser amplifiers requires that the gain of the system be optimized to avoid the effects of amplified spontaneous emission (ASE), parasitic growth, and nonsaturable

Fig. 8-31. Output-switch trigger circuit.

Fig. 8-32. B-amplifier cathode geometries.

absorption. The extraction efficiency of the KrF amplifier generally exhibits a pronounced maximum in the gain regime from 3 to 5. Control of system front-end cost and complexity usually dictates somewhat higher gains than this, with an attendant reduction in amplifier extraction efficiency.

We have studied laser architectures in which the RGH pump medium is contained in an optical cavity, coupled to the Raman medium which is contained in a second optical cavity. Our preliminary analysis indicates that we can maintain nearly optimum extraction efficiency in both media during long-pulse operation. We have given this concept the name "Raman accumulator" since two high-Q regenerative cavities (the RGH cavity and the Raman cavity) interact in the Raman medium to cause a steady accumulation of energy in the first-Stokes pulse at the expense of the RGH pump pulse.

Figure 8-33 shows the elements of the system. The RGH medium is shown at the left, enclosed in a high-Q cavity called the RGH cavity. This cavity extends through a converter channel that houses the Raman medium at the right. As the RGH medium undergoes excitation, the RGH cavity fills with radiation at the pump wavelength. A Raman accumulator cavity is established which passes first-Stokes radiation down the bore of the converter channel, extracting

energy from the obliquely propagating pump radiation. This extraction technique leads to a high fill factor and reduces the constraints on pump beam quality. The radiation in the RGH cavity can be divided into two components, one propagating in each direction. The cavity lengths are selected so that the short first-Stokes pulse in the accumulator cavity interacts with half of the radiation in the RGH cavity on each round trip through the accumulator. Thus, the effective pump pulse width, τ_t , for Raman extraction is given by one transit time in the RGH cavity,

$$\tau_t = \frac{L}{c}, \quad (1)$$

where L is the RGH cavity length and c is the speed of light. The corresponding accumulator cavity length is $L/2$ and the length of the Raman medium is $L/4$. Successive passes of the pump and Stokes beams are permitted in their respective cavities, while excitation of the RGH medium proceeds for as long as acceptable extraction efficiency in the RGH amplifier can be maintained.

An important feature of this concept is that the portion of the pump radiation in the converter channel which is not converted to Stokes radiation is sent back to the RGH medium, where it works to control ASE and

Fig. 8-33. Schematic of the Raman accumulator architecture.

Pulse Compressor Systems

parasitics and keep the gain of the RGH amplifier in a regime where efficient extraction can occur. Extraction by the Stokes pulse serves to keep the gain in the RGH amplifier from becoming too low, where nonsaturable losses begin to reduce the extraction efficiency. This symbiotic interaction of the two coupled cavities is the principal advantage of the accumulator concept. It can be shown that overall efficiency of the accumulator system, η_{AC} , approaches the overall efficiency of the RGH medium, η_{RGH} , and hence that of the stacker, except for the small quantum defect between the pump and first Stokes frequencies:

$$\eta_{AC} = \left(\frac{\nu_s}{\nu_p} \right) \eta_{RGH} \quad (2)$$

Here ν_s and ν_p are the Stokes and pump frequencies, respectively.

We have used analytical models of the accumulator to calculate the extraction efficiencies and fluence levels in the device over a series of optical passes through the system. We have estimated the effects of optical losses and finite rise and fall times of the pulsed power on the RGH amplifier. Approximately 10 passes in the RGH cavity resulted in optimum transfer of energy to the accumulator.

We can achieve switch-out of the accumulated energy as shown in Fig. 8-33. Here an optical pulse at the second-Stokes frequency is injected into a Raman switch at the end of the accumulator; it builds up and extracts the energy from the accumulator cavity in two passes through the Raman switch.

We have made preliminary scaling models of the Raman accumulator at the multi-megajoule level; these models allow us to compare the accumulator performance to

Fig. 8-34. Conceptualization of a reactor-scale Raman accumulator 2-MJ laser system.

that of other generic fusion laser systems. Our analytical models indicate the possibility of an overall system efficiency approaching 7%, at a temporal pulse compression factor of 50.

Figure 8-34 shows one possible architecture for a reactor-scale 2-MJ laser system. The RGH power-amplifier array is in the form of an annular stripline with a pulse length of 1 μ s. The amplifier array is segmented into 16 parts; each segment of the laser medium has an aperture of ~ 2 m, and is excited from the outer boundary by electron beams. The accumulators and front-end drivers for each segment are located directly above. The anticipated compactness of the Raman systems is readily apparent in this design, and with the potentially higher efficiency of the accumulator geometry compared to conventional Raman systems, this system architecture may represent the most practical fusion laser concept based on nonstoring media. Experimental verification and complete numerical simulation of the accumulator concept will be addressed as the next steps in concept evaluation.

Authors: E. V. George and J. C. Swingle

A Method for Suppressing Second-Stokes Parasitics

It is now well established, both theoretically and experimentally, that the performance of a Raman compressor is seriously constrained by second-Stokes parasitics. In this article we describe a spectral filter for use in KrF/CH₄ Raman compressor laser systems. We also describe laser-system architectures in which the spectral filter may be used to suppress second-Stokes build-up.

In the absence of absorption, the growth of the intensity I_{2S} of second-Stokes radiation within a Raman amplifier is given by

$$I_{2S} = (I_{2S}(0) + I_{sp}) e^{G_{2S}} \quad (3)$$

where

$$G_{2S} = \left(\frac{c\gamma^f I_S}{n_p + n_S} \right) L \quad (4)$$

is the second-Stokes gain, γ^f is the forward intensity gain coefficient, L is the amplifier length, I_S is the Stokes intensity, and n_j is the appropriate index of refraction for the pump ($j = p$) and Stokes ($j = S$) waves. The quantity $I_{2S}(0)$ is the intensity of second-Stokes radiation at the input of the Raman amplifier, due, for example, to earlier stages of amplification. The second-Stokes spontaneous-scattering source term is represented by an effective intensity

$$I_{sp} \approx \left(\frac{\pi \hbar \nu_S \Delta \nu_r}{\lambda_S^2} \right) d\Omega \quad (5)$$

Here, as an approximation, $d\Omega$ is the effective solid angle subtended by the output aperture of the amplifier as seen from the input, namely

$$d\Omega \sim \frac{\pi a^2}{L^2} \quad (6)$$

where a is the radius of the amplifier. For typical KrF/CH₄ Raman amplifiers $a \sim 1$ m, $L \sim 15$ m, $\Delta \nu_r \sim 0.1$ cm⁻¹, and the effective second-Stokes spontaneous noise intensity is $I_{sp} \sim 0.2$ W/cm².

When the second-Stokes intensity grows in magnitude to $\sim 10^9$ W/cm², serious depletion of the backward first-Stokes pulse occurs. Avoidance of this condition requires that the second-Stokes gain be restricted so that

$$G_{2S} < \ln \left[\frac{I_S}{I_{sp} + I_{2S}(0)} \right] \\ \sim \ln \left[\frac{10^9 \text{ W/cm}^2}{I_{sp} + I_{2S}(0)} \right] \quad (7)$$

If the second-Stokes noise input $I_{2S}(0)$ is much less than $I_{sp} \sim 0.2$ W/cm² (a nontrivial condition to achieve), then Eq. (7) requires that $G_{2S} < 22$. This is the condition that is generally employed to estimate second-Stokes parasitics limitations.¹⁹ On the basis of this simple analysis, it is clear that

exponential attenuation [$\exp(\alpha_{2S}L)$] between amplification stages, with $\alpha_{2S}L > 20$, will be required to effectively suppress second-Stokes build-up. Last year's annual report¹⁹ gives a more detailed treatment of second-Stokes limitation and suppression techniques.

To date, we have found no satisfactory in situ absorbers for KrF/CH₄ compressors which both pass the pump and Stokes waves with very low loss and efficiently attenuate the second-Stokes radiation. We have, however, developed certain architectures for staging KrF/CH₄ pulse compressors in which the pump and Stokes beams are decoupled.

As shown in Fig. 8-35, these architectures include (a) multipass extraction, (b) single-

pass multibounce, and (c) regenerative configurations such as the accumulator architecture discussed in the previous article. In all cases, the second-Stokes absorber must be positioned inside the cavity, so that second-Stokes radiation reflected from its surfaces is separated from the backward first-Stokes pulse. This approach, which simplifies the second-Stokes suppression problem, requires a filter which transmits first-Stokes radiation at $\lambda_S = 268$ nm and strongly absorbs second-Stokes radiation at $\lambda_{2S} = 291$ nm. From the previous analysis it is desirable to have a loss-length product, $\alpha_{2S} \mathcal{L}$, greater than 20 (here α_{2S} is the second-Stokes absorption coefficient and $\mathcal{L}$ is the thickness of the filter). To maintain a high first-Stokes efficiency, it is desirable to have the first-Stokes loss-length product, $\alpha_S \mathcal{L}$, as small as possible, but certainly not larger than 0.1. A value of $\alpha_S \leq 0.01$ would be considered ideal. Thus we require $\alpha_{2S}/\alpha_S = 200$, and desire $\alpha_{2S}/\alpha_S > 2000$.

Figure 8-36 shows the spectral absorbance of cerium-doped lithium yttrium fluoride ($\text{Ce}^{3+}:\text{LiYF}_4$, or $\text{Ce}:\text{YLF}$ for short) reported by Ehrlich et al.²⁰ The absorption measurement was for unpolarized light propagating parallel to the *c* axis of a 275- μm -thick plate of 1%-doped material; it is not corrected for reflection losses. The fluorescence spectrum is not corrected for self-absorption around 300 nm, and thus actual intensity in this region may be somewhat larger than indicated. Spectra taken at excitation wavelengths of 193, 248, and 254 nm were essentially the same. The inset shows a schematic energy-level diagram for $\text{Ce}:\text{YLF}$. The spectrum exhibits a very strong $4f \rightarrow 5d$ absorption band centered at ≈ 291 nm, and appears to be highly transparent at 268 nm.

The peak cross section for the 291-nm transition can be estimated from the data provided in Fig. 8-36. The sample was 275 μm thick and contained an estimated Ce^{3+} ion density of 1.6×10^{20} ions/ cm^3 , or 1 at.%. Note that nominal doping percentages are usually quoted in terms of the percentage impurity placed in the melt from which the crystal is grown. Since it is anticipated that the distribution coefficient for Ce^{3+} in YLF is less than unity, 1.6×10^{20} ions/ cm^3 is likely to be an upper bound; the derived cross section will therefore be a lower bound.

Fig. 8-35. Generic KrF/CH₄ Raman compressor amplifier stages with Ce:YLF second-Stokes absorbers: (a) multipass extraction, (b) single-pass multibounce, and (c) regenerative configuration.

The peak optical density of the 291-nm absorption band is 2, so the cross section is

$$\sigma(291) = \frac{2.3 \text{ (OD)}}{\rho(\text{Ce}^{3+})\mathcal{L}} \geq \frac{(2.3) (2)}{1.6 \times 10^{20} (0.0275)} = 10^{-18} \text{ cm}^2 \quad (8)$$

We can get a second estimate on this cross section using the Einstein relationship

$$\sigma = \frac{\lambda^2}{8\pi c n^2 \Delta\nu \tau_{\text{rad}}} \quad (9)$$

where λ is the wavelength of the inverse fluorescence transition $5d \rightarrow 4f$, $\lambda = 312 \text{ nm}$, $\Delta\nu$ is the fluorescence line width of the inverse transition, n is the refractive index, and τ_{rad} is the radiative lifetime. The fluorescence lifetime is 40 ns; since about 40% of the fluorescence emission occurs in the inverse transition at 312 nm, we can estimate $\tau_{\text{rad}} = 100 \text{ ns}$. Using these data, together with $\Delta\nu = 2000 \text{ cm}^{-1}$ and Eq. (9), we have

$$\sigma = \frac{(0.312)^2 \times 10^{-8}}{8\pi c (1.48)^2 (2000) (10^{-7})} = 3 \times 10^{-18} \text{ cm}^2 \quad (10)$$

The factor of 3 between this value and the value given in Eq. (8) may be due to an over-estimation of the Ce^{3+} ion density used to evaluate Eq. (8). As a compromise, pending direct measurement of the Ce^{3+} ion concentration in a spectroscopic sample of this material, we assume the cross section value of $2 \times 10^{-18} \text{ cm}^2$. This is a very characteristic value for quantitatively measured $4f \rightarrow 5d$ transitions in other rare-earth-doped materials. Using this cross-section value and a cerium ion doping of $1 \times 10^{20} \text{ ions/cm}^3$, we have $\alpha_{2S} = 200 \text{ cm}^{-1}$. Thus if we want $\alpha_{2S}\mathcal{L} = 20$ we require a 0.1-cm-thick plate of Ce:YLF.

In Fig. 8-36 it is not possible to determine the value of α_S , because the loss is so small. Figure 8-37 shows another spectrum of a nominal 1% Ce:YLF crystal. This particular sample may suffer from minor scatter and refraction effects. The nominal optical path length was 0.43 in. (1.1 cm). At a wavelength of 320 nm there is no absorption, so the zero level is set at this wavelength to remove Fresnel losses. The peak of the 291-nm

transition is off the scale, as it exceeds an optical density of 2. At 268 nm we observe a total loss density of 0.03. This loss includes any slight increase in Fresnel loss relative to that at 320 nm, due to dispersion of the index of refraction, scattering loss, and true Ce^{3+} ion absorption.

As a worse case we assume that all of the observed loss is due to Ce^{3+} ion absorption. Then $\alpha_S = 2.303(0.03)/1.1 \text{ cm} = 0.063 \text{ cm}^{-1}$. Using this value, a 0.1-cm filter plate would introduce an absorption loss of 0.006, and

$$\frac{\alpha_{2S}}{\alpha_S} = \frac{20}{0.006} \approx 3300 \quad (11)$$

Thus Ce:YLF not only satisfies the minimum requirement, but also exceeds the desired performance ratio. Note that Ce:YLF absorbs at the KrF pump wavelength (248.5 nm) with a cross section $\sigma(248) \sim 1.5 \times 10^{-20} \text{ cm}^2$. If a Ce:YLF filter plate is placed in the path of an intense KrF pump beam, considerable absorption and

Fig. 8-36. Absorption and fluorescence spectra of Ce:YLF (source: Ref. 20).

Fig. 8-37. Absorption spectrum of a nominal 1% Ce-doped YLF crystal.

heating of the filter will occur. Therefore, the Ce:YLF second-Stokes suppressor must be employed, as shown in Fig. 8-35.

The material itself and its practical use seem promising. It is fairly hard and stable. From an optical point of view, YLF is transparent²¹ to 150-nm radiation (in pure material), and is nominally free of two-photon absorption²² at $\lambda = 268$ nm. It has a relatively low nonlinear refractive index at $\lambda = 1.06$ μm of 0.6×10^{-13} esu (Ref. 23). It has a relatively high bare-surface damage threshold²⁴ at $\lambda = 1.06$ μm of 87 J/cm^2 for a 4.6-ns pulse and at $\lambda = 0.35$ μm of 300 J/cm^2 for a 40-ns pulse.²¹

The major (and perhaps only) drawback of YLF is that it is a crystalline material and therefore subject to an issue of growth. Sanders Associates has grown Tm:YLF crystals up to 10 cm in diameter, and believes that YLF can be scaled to much larger dimensions.²¹ It would be very beneficial if the filter thickness could be reduced sufficiently to permit vapor deposition or epitaxial growth on an inert substrate, to achieve large contiguous areas. Of course to obtain adequate absorption at 291 nm, one would have to greatly increase the Ce^{3+} doping level. Although Ce^{3+} tends to not go into YLF in conventional crystal growth, the possibility of a thin Ce:YLF film cannot be ruled out at present, especially if an attractive Raman-compressor laser system design can be worked out which requires a filter function appropriate for Ce:YLF.

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Pulse-Stacker Laser Concepts

Perhaps the simplest and most general approach to pulse compression is pulse stacking.^{25,26} In a pulse-stacking scheme, energy from an amplifier medium is extracted by a sequenced train of pulses that have been time-coded in some manner. Pulses in the train are separated after amplification, according to the predetermined code. After separation, the pulses are recombined and stacked, with an appropriate time delay added to each pulse. This assures simultaneous arrival at a point or plane of all members of the sequence.

The simplest coding is by angle. In principle, wavelength coding or polarization coding could also be used; we are less certain, however, of the aperture scalability of the dispersive elements which translate a time-dependent wavelength or polarization into the spatial separation required for pulse stacking. There are also other coding options, for a pump medium able to sustain rapid space/time variations in its population inversion density. We restrict the discussion here to pulse compression achieved by stacking angle-coded pulses in a medium of either high or low gain coefficient.

Efficient extraction of energy from a medium with a high gain coefficient (such as rare-gas halides) requires the sequence of pulses to be contiguous in time, in order to prevent loss of inversion energy due to the occurrence of superfluorescence or parasitic oscillation between extracting pulses. For a total amplifier gain duration, τ_p , each of a sequence of n contiguous pulses will have a pulse length τ_p/n . In the case of sequenced extraction of energy from a medium with a low gain coefficient, or from an energy-storing medium (a backward-wave Raman amplifier), each of the n pulses forming the sequence need not be contiguous in time. The pulses can have a shorter pulse length, $\tau_p/\kappa n$, where $\kappa > 1$, so long as the medium can store energy without significant radiative or collisional loss for a time τ_p/n , the interval between extracting pulses.

The power multiplication attainable by pulse stacking is related simply to the instantaneous power which can be drawn from the medium. For an amplifier producing an output intensity I_0 over an area A , and producing net power $P_0 = I_0 A$, the total power gain achieved at the plane where n beams are stacked together is simply $P/P_0 = \eta_0 n$, where η_0 is the optical train efficiency, including medium fill factors, scattering, and absorption losses by optical surfaces. Pulse stacking is attractive because η_0 can be made close to unity. Stacking is also efficient, although it is optically cumbersome for very large compression ratios.

Two generic approaches to pulse stacking with high-gain laser media are illustrated in Fig. 8-38. In one approach [Fig. 8-38(a)] each amplifier stage in the laser system

contains a pulse-stacking (short-pulse angle-encoding) and unstacking (short-pulse angle-decoding) optical system. Between amplifier stages the pulses are transported as a single short pulse (~ 10 ns). Optical relaying and spatial filtering of this short pulse between amplifier stages is then accomplished in the usual manner.

In the second approach [Fig. 8-38(b)], pulse stacking and unstacking are used only at the beginning and end of a train of amplifiers. Successively larger amplifiers are used in sequence to amplify the initial train of short low-energy pulses to the energy

level required for target irradiation. After amplification, the train of pulses is unstacked to produce the short high-energy pulse for target irradiation. In this scenario, laser-beam control may also be accomplished by optical relaying and spatial filtering between amplifier stages. An example of an amplifier stage which may be used in the first approach is discussed in Ref. 27; in this article we will discuss the characteristics of the second approach.

Figure 8-39 shows a pulse-stacking laser system in which the angle encoding and angle decoding are performed only on the input and

Fig. 8-38. Two generic approaches to pulse-stacking or angle-multiplexing systems for fusion lasers.

(a) Chain of separately angle-multiplexed amplifiers

(b) Angle-multiplexed chain of amplifiers

Fig. 8-39. Pulse-stacker laser system in which a train of contiguous short pulses is amplified by a chain of amplifiers; for simplicity only two beam lines are shown.

Pulse Compressor Systems

output laser pulses, as in Fig. 8-38(b). By properly arranging the amplifiers, each pulse requires only a single mirror between stages of amplification. This mirror redirects the pulse toward the next amplifier, and also provides control of the laser beam diameter. Depending on the amplifier sizes, these intermediate mirrors may be either plane or curved. This design is perhaps the most efficient pulse stacker concept in terms of the number and size of mirrors.

It may be difficult to propagate a laser pulse through the long distances required by an optical pulse stacker, due to diffraction effects and optical distortion of the laser pulse. These effects are in turn due to optical figure errors, amplifier fill factors, or optical medium inhomogeneity. The average pulse energy must be lowered significantly, to prevent optical surface damage if the laser pulse picks up a significant amount of intensity modulation. One method of reducing the effects of diffraction, once the beam picks up a phase modulation (due to figure error or medium inhomogeneity), is to provide optical relaying between stages of amplification, using optical imaging elements placed in each beam path between amplifier stages in such a way as to image the amplifiers onto each other (Fig. 8-40).

The intensity and phase profiles of the beam at one amplifier are thus reproduced at the succeeding amplifier. As long as the intermediate optical elements are sufficiently large, phase disturbances produced at one amplifier are not connected to intensity disturbances resulting from diffraction propagation to succeeding amplifiers.

Optical relaying introduces its own problems, however:

- The number of optical elements required increases significantly.
- Since the pulse-stacking telescopes magnify the beam diameter, the encoding angles are demagnified. Therefore, if the pulses are to be angle-decoded within a relatively short distance, the input pulses must be separated by very large angles, which spreads out the optical elements over a very large volume.
- Large input angles make it difficult to design a stacker/relay system with good optical beam quality without using exotic optical elements.

Figure 8-40 shows a relayed pulse stacker which reduces the angle demagnification problem. In addition to an optical relay between the amplifiers, a set of mirrors is provided to reset the encoding angle to a desirable value. These mirrors also provide

Fig. 8-40. Pulse-stacker laser system with optical-relay/filter-transport system for each angle-multiplexed beam line.

far more compact arrangement of the laser stages, by allowing more optimum packing of the various components. Nevertheless, the angles on the intermediate mirrors in the configurations of Figs. 8-39 and 8-40 are large enough to create significant optical-system problems.

In summary, we have identified two generic approaches to pulse stacking. These novel systems utilize telescopic optical-transfer techniques to simultaneously maximize beam fluence and amplifier extraction efficiency. This, in turn, minimizes mirror surface area and cost. Other pulse-stacker optical-train designs have also been identified; these use regenerative extraction techniques to greatly simplify the input optical train. Pulse stackers can achieve large and efficient compressions that yield high laser efficiencies ($\sim 5\%$) when used with efficient laser media such as KrF. Due to the pulse time delays required to work at large compressions, stacker architectures tend to encompass a large area. This contributes significantly to the generally high cost of pulse-stacker systems. Large-pulse compressions also require the multiplexing of many beam lines, which presents a serious problem for alignment, control, pointing, and tracking of the output laser beam. Clever innovation will be required to reduce the scale size and complexity of these systems, in order to realize their potential for fusion applications.

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Summary of Stacker-Compressor Technical Issues

We have outlined the physical and technological advantages and disadvantages of pulse stacking and backward-wave Raman scattering as distinct means of compressing pulse outputs from UV excimer lasers. The overall efficiency and cost of a pulse-compressor system depends on the pump efficiency, the architecture chosen, and the pulse-compression ratio desired. Figure 8-41 compares the efficiency of the compression subsystems for various combinations of stacking and Raman compression. The optical subsystem efficiency is estimated to be in the range from 85 to 95%, and is dependent on system fill-factor and mirror reflectivities. The pure backward-wave Raman compression efficiency is strongly dependent on compression ratio, as previously shown.

The region of interest shown in Fig. 8-41 is for compression ratios in the range of approximately 2 to 20. Smaller compression ratios are not considered, because the fluence gain at the Stokes frequency is too small. The compression ratio for a pure stacker is arbitrarily cut off at approximately 50 for this graph; the zone plot for pure Raman compression is arbitrarily terminated at a net compression of 20. The arrow indicates the direction of improved performance if a second Stokes absorber is used. Regenerative Raman techniques, such as the accumulator

Fig. 8-41. Pulse compression efficiency plotted as a function of total system compression for three compression techniques.

Pulse Compressor Systems

discussed in "Advanced Architectures for Raman Compression" earlier in this section, are capable of achieving efficiencies approaching those of a pure stacker (85 to 95%).

The efficiencies of various hybrid combinations are derived by multiplication of the appropriate zones of efficiency. Stacker-compressors occupy the performance region in Fig. 8-41, which begins at a net compression of 50 due to the integral nature of stacking and the minimum Raman compression for backward Raman amplifiers. Performance curves for stacker-compressors, with smaller total compressions, can be generated by displacing the appropriate stacker-compressor curves in Fig. 8-41 to the left. Stacker-compressors are more efficient than pure Raman compressors for compressions of 10 or more, and offer efficiencies in excess

of 75% for compressions as large as 75 (such as a stacking of 25 and a compression of 5). For compressions of larger pulses, one can choose either to stack in both the pump and Raman cells, or to utilize the Raman effect twice (the efficiency zones for these cases are not shown in Fig. 8-41). Compressions in excess of 200 can be achieved at fair efficiency with hybrid combinations.

The overall efficiencies which might be achieved with stacker-compressor systems are compared in Table 8-1, assuming typical values of the net pulse compression available with each architecture. Electrical efficiencies of about 5 to 8% are projected for pure pulse stacking of KrF amplifiers. Efficiencies of about 4 to 7% are projected for the stacker-compressor systems which feature operation at comparable net pulse compressions. These efficiencies are consistent with the baseline requirements projected for a fusion power-plant driver.

A qualitative comparison of the major features of these generic stacker-compressor systems is presented in Table 8-2. In this table a minus sign denotes performance that is less than desirable or relatively poor; the plus sign denotes relatively better performance characteristics.

Pure angle-coded pulse stacking is conceptually simple, and we anticipate its capacity for large compressions (>25) with relatively high efficiency. Because UV excimer-laser amplifiers are power-limited devices, however, due to the presence of nonsaturable photoabsorption losses, the staging of arrays of pulse-stacked KrF amplifiers into the multi-hundred-terawatt regime appears technologically complex and expensive. Such systems call for large numbers (≤ 1000) of electromechanically controlled beam lines and large numbers of

Table 8-1. Comparison of overall efficiencies of typical pulse compressor systems.

Architecture	Pure stacker	Pure Raman	Stacker-compressor
Overall compression			
Subsystem efficiency (%)	25 to 50 ^a	10 ^b	20 ^c
E-beam generation, transport, and deposition	← 60 to 80 ^d →		
Intrinsic pump medium	← 10 to 12 →		
Pump electrical	← 6 to 9.6 ^d →		
Pump optical transport/coupling	90	90	90
Raman extraction (methane)		55 to 85 ^e	75 to 85
Laser system electrical	5.4 to 8.6	3 to 7.3	4 to 7.3

^aLarger compressions are possible with a greater number of beams.

^bOverall efficiency will be high if a lower compression ratio is chosen. Large, high-efficiency compressions require the use of a second-Stokes parasitics suppressor.

^cStacking of 4, backward Raman compression of 5.

^dRequires a major breakthrough in pulsed-power technology.

^eRequires a second-Stokes absorber or regenerative Raman technique, such as an accumulator system.

Table 8-2. Comparison of features of generic stacker-compressor systems.

Property	Raman compressor	Stacker	Stacker-compressor
Efficiency	-	++	+
Compression	-	+	+
KrF pump technology	-	+	+
Cost sensitivity to KrF energy/aperture	+	-	+
E/O mechanical complexity	-	-	-
Optical damage constraints	-	+	-
Pulse length trade-off	-	-	-
Pellet preheat	+	-	+

mirrors of relatively high optical quality. The cost, reliability, and pointing and tracking capability of such electrooptical mechanical systems is a serious issue that has not been resolved. One possible solution would be reducing the number of beam lines, by using large-aperture high-energy KrF amplifiers. Another issue for pure pulse stackers is channel cross talk, which leads to target preheat. If saturable absorbers are required to control the problem, then laser efficiency will be reduced. Pure pulse stackers also generally require a large area or volume, leading to high building costs. They also have a limited ability to generate variable-length pulses, which is a problem in systems that must be designed to produce the shortest pulse length required.

The performance characteristics of pure Raman compressors are currently limited by the second-Stokes parasitics problem (see "A Method for Suppressing Second-Stokes Parasitics" earlier in this section). Although the quantum efficiency of KrF laser-light Raman scattering in CH_4 is high ($\sim 93\%$), parasitics limit conversion efficiencies to $\sim 55\%$ for compressions on the order of 10. Except for this problem, the pure Raman approach is the most attractive, because of its inherent simplicity and high potential efficiency. Two candidates for solving the second-Stokes problem are a second-Stokes absorber, and the novel architecture of the Raman accumulator, which exploits unique operating regimes to circumvent second-Stokes parasitics. At present, the Raman approach is the only one that appears amenable to compact system architectures. Due to the current complexity of pulse-stacking scenarios, appropriate efforts in Raman system development are highly desirable.

Stacker-compressor systems combine the attractive features of both angle multiplexing and backward-wave Raman compression:

- Relatively high efficiency is achieved by using small Raman compressions ($\sim 5\times$).
- The target-preheat problem due to stacker-channel cross talk is eliminated by the Raman conversion process.
- The number of laser beam lines compared to a pure stacker is substantially reduced ($\sim 5\times$).

Unfortunately, optical alignment and control issues are still significant compared

to the ideal storage laser. All stacker-compressor laser system concepts, as presently conceived, rate a "-" in E/O mechanical complexity and pulse length flexibility (see Table 8-2). Many such systems are multicolor lasers, and thus require separate laser systems to provide pulses at appropriate frequencies. Major innovations in system architectures will probably be required if these systems are to satisfy cost and reactor efficiency requirements.

Another major technological issue that must be addressed for all UV excimer-laser stacker-compressor systems is the development of high-efficiency ($\gtrsim 98\%$), high-damage-threshold ($> 5 \text{ J/cm}^2$) optical elements. It is recognized that presently available transmitting and reflecting optics in the $0.25\text{-}\mu\text{m}$ spectral region will have to be substantially improved. Present conceptual designs of large systems call for optics capable of handling beam fluences of 5 to 10 J/cm^2 . Selected individual mirrors have shown damage thresholds in this range, but at present the average commercial sample suffers damage at only 1 to 2 J/cm^2 . Optical elements and materials suitable for service at $\sim 0.25 \mu\text{m}$ are considerably fewer in number than those suitable for service at $1 \mu\text{m}$, although it appears that adequate optics are likely to be developed using the methodology proven to be successful in developing $1\text{-}\mu\text{m}$ optics for Nd:glass systems.

Finally it is important to note that the stacker-compressor laser efficiencies discussed in this section are referenced to the laser output outside the reactor system and not to the target. The latter efficiency, η_D , is the one required to satisfy the driver requirements. This efficiency is the product of the laser output efficiency (see Table 8-1) and the optical efficiency, η_c , for coupling the laser output to the target through the reactor-laser interface. The value of η_c depends on the details and realities of reactor-vessel and containment-building design. To date, reactor designs, which include laser-coupling considerations, have required several (~ 4) reflections to couple the output of the laser to the target. Several "windows" are also required to isolate the reactor environment. For optical systems of this type, maximal values of η_c in the 80 to 90% range are likely. Clearly, critical issues such as final optics configuration, lifetime,

and target stand-off distance need further clarification. These considerations, together with target irradiation requirements (target size, pointing and tracking, etc.), must be accounted for in laser designs of the future.

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Laser Theory and Design Introduction

In the next three articles we present the results of analytical and experimental programs that have addressed the kinetics of rare-gas-halide (RGH) laser media. These programs include

- Kinetic experiments with electron-beam-pumped RGH laser media.
- Theoretical studies of KrF vibrational relaxation.
- Ab initio calculations of dissociative electron attachment and vibrational-excitation cross sections for F₂.

Additional contributions from the laser theory and design program to pulse-compressor systems and advanced concepts studies are included in articles before and after this subsection, respectively.

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Rare-Gas-Halide Kinetics Experiments

There is a superficial similarity among the kinetic mechanisms of all RGH media, especially in terms of energy-channeling and spectral-emission properties. According to our present understanding, for example, the production of one rare-gas ion or one rare-gas excited state typically leads to the production of one RGH upper-state molecule. Energy channeling then occurs by

- Ion-ion recombination (of a positive rare-gas ion with a negative-halogen ion).
- The neutral-channel reaction of rare-gas metastables with halogen-bearing molecules.
- Subsequent displacement reactions, such as $\text{ArF} + \text{Kr} \rightarrow \text{KrF} + \text{Ar}$.

If we could fully convert this efficient excited-state production to output-laser photons, the intrinsic-medium efficiency of electron-beam-pumped RGH media would be a very attractive 25%.

According to presently available experimental data, KrF comes nearest to realizing this potential, having an intrinsic efficiency of about 10% (see Ref. 28). According to conventional wisdom, KrF does not reach its full efficiency potential because of the presence of nonsaturable absorbers, which are formed during excitation and are precursors to the formation of the upper laser level. The presence of these absorbers results in a typical gain-to-loss ratio of about 10:1, which in turn limits the output fluence to about 10 MW/cm² (roughly 10 times the saturation fluence) and the extraction efficiency to about 50%. Although experimental evidence supports this view,^{29,30} specific questions about KrF remain, such as the identification of the absorbing species and the rate of cross-relaxation within the upper laser level (discussed in the next article). Gaining understanding of these issues may provide some hope for improving the efficiency of the KrF laser.

Despite the similarities in their kinetic mechanisms, there are substantial differences in RGH laser performance. As discussed in the last year's annual report,³¹ XeF is perhaps the best-understood case for having a substandard laser-medium performance. In XeF the C state lies substantially below ($\approx 650 \text{ cm}^{-1}$) the B state which is the upper level for the 350-nm B-X laser.³² The accumulation of population in the C state, and its finite coupling rate to the B state, substantially reduces the XeF gain and laser-extraction efficiency. This problem has been countered in some XeF experiments by heating the laser medium.³³ We believe this problem is much less important in XeCl, KrF, and ArF lasers, however, because their B-C splittings are quite small.³⁴

Although the main effort in our Advanced Lasers Program is based on pulse compression of e-beam-pumped KrF lasers, we are also interested in the use of alternative RGH media. To provide a basis for comparison, we are presently measuring the intrinsic-medium efficiencies and

maximum-output power limits of XeCl, KrF, and ArF lasers. According to previous studies, these lasers appear to be the most promising RGH candidates in terms of efficiency. Their respective output wavelengths of 308, 249, and 193 nm cover a range of bandwidths within which there may be differences in fusion target performance as well as variation in the damage resistance of thin-film dielectric-coated optics. Thus the basic data on the performance characteristics of these lasers may be combined with several other factors in choosing an advanced laser candidate for further development.

XeCl Laser. During 1980 we studied experimentally the XeCl laser. The history of XeCl differs markedly from KrF. The XeCl laser was first demonstrated in 1975 by Ewing and Brau, who employed a gas mixture of Ar, Xe, and Cl₂, and e-beam pumping at a current density of 150 A/cm² (Ref. 35). They found the XeCl output energy to be orders of magnitude smaller than that of KrF, however. The future for XeCl brightened considerably in 1978, when Champagne observed that the XeCl laser "compared favorably with the best results obtained for KrF" when he used a mixture of Ne, Xe, and HCl and an e-beam pumping current of 8 A/cm² (Ref. 36). He reported an optimum XeCl laser efficiency of 5% at an operating pressure of 4 atm. Champagne attributed the improvement in laser performance to the reduction in medium absorption caused by using Ne and HCl rather than Ar and Cl₂, respectively. Using excitation currents on the order of 40 A/cm², Rothe et al. echoed these results in 1979, claiming that "XeCl may be as efficient as KrF."³⁷

From a kinetic modeling standpoint, the XeCl laser results were at first difficult to understand. Setser's measurements had indicated that Xe metastables react rapidly with HCl, but do not produce XeCl excited molecules.³⁸ It was therefore presumed that with Ne-Xe-HCl mixtures, the XeCl laser depended solely on the ion-ion recombination mechanism. This was rather different than the multiple pathways postulated for KrF. An examination of measured kinetic rates available at that time, however, made the ion-recombination seem unlikely. Specifically, the bimolecular

charge-transfer rate from both Ne⁺ and Ne₂⁺ to Xe had been measured to be quite slow,³⁹ so there did not appear to be an effective channel for making Xe⁺.

Significant progress toward understanding the XeCl system was made at LLNL and elsewhere during 1980. Work by Collins and Lee showed that the rate for termolecular charge transfer, Ne₂⁺ + Ne + Xe, was fast enough that at atmospheric pressure and above there was no significant bottleneck in the formation of Xe⁺ (Ref. 40). Our experimental studies of XeCl fluorescence, however, indicated that the attachment rate of electrons to halogen-bearing molecules (to form Cl⁻) was usually the rate-limiting step in the formation of XeCl. Simply using higher halogen-donor pressures to increase the Cl⁻ formation rate was not beneficial, because the Cl donors quenched the RGH excited states at a gas-kinetic rate.⁴¹

We found that the primary advantage of using HCl rather than Cl₂ was the greater rate of attachment, and ultimately the larger formation rate of XeCl excited states. The results of Allan and Wong⁴² are crucial to understanding this mechanism. They showed that the cross section for attachment to HCl at the threshold electron energy is 38 times greater for HCl in $v = 1$ compared to HCl in $v = 0$. The cross section for HCl in $v = 2$ is 880 times greater (see Fig. 8-62 in the following article). Hence we now presume that the formation of Cl⁻ proceeds by vibrational excitation of HCl (which occurs with a large cross section), followed by attachment. We observed evidence for this two-step mechanism in the dependence of XeCl fluorescence on excitation power (current). We found a considerably different current-dependence for Cl₂. These results were investigated with a kinetic model (described in the next article) which supports this mechanism. We also studied other Cl donors, such as CCl₄ and CCl₃F, which we found to produce XeCl fluorescence greater than when using Cl₂ mixtures but less than HCl mixtures.

In our laser experiments we have studied primarily Ne-Xe-HCl and Ar-Xe-HCl mixtures in the pressure range of 1 to 4 atm. System considerations favor operation with a medium pressure near 1 atm. We studied laser performance as a function of current density, to find the limits produced by fuel

Fig. 8-42. Schematic of the e-beam-pumped cell used to study RGH lasers with variable current density.

Fig. 8-43. Faraday cup signals of the current collected in the gas cell vs time.

burn-out and electron deactivation of the excited states. Using the three-dimensional, Monte Carlo computer code SANDYL⁴³ to calculate e-beam energy deposition, we found the optimum medium efficiency for the XeCl laser to be about 4.5%, with an Ar-Xe-HCl mixture operating near 1 atm.

Experimental Setup. A diagram of the MEG I experimental setup is shown in Fig. 8-42. The nominal e-beam pumping current was $490 \pm 100 \text{ A/cm}^2$, as measured by a carbon-collector Faraday cup placed within the gas cell. The Faraday collector, positioned immediately behind the 3-mil Inconel cell foil, was apertured by a 1- by 10-cm opening in a stainless steel plate, to measure the current in a central slice of the nominally 2.5- by 10-cm beam. The error range given above represents the variation in current observed for measurements over several months time, and is much larger than the shot-to-shot variation observed during one day. We presume the long-term current-density variation was caused by slight cathode changes, resulting in different transport efficiencies of the high-current beam. The e-beam was quite uniform over 1 cm in the vertical direction, and varied about 20% in the horizontal direction. The diode voltage was approximately 1 MeV and the pulse duration was 50 ns.

We placed different perforated carbon attenuator plates within the drift region, as shown in Fig. 8-42, to attenuate the beam current without loss of voltage or uniformity. The e-beam current was varied over the range from 5 to 500 A/cm^2 using this method. The Faraday cup signals in Fig. 8-43 show the time-dependence of the

current at two current densities. Figure 8-43(a) is for the nominal full current of approximately 500 A/cm²; Fig. 8-43(b) is at an attenuated current of approximately 150 A/cm². The small negative spike, occurring at the final division in each picture, is a fiducial mark which is temporally locked to the current pulse, allowing accurate positioning of the signals relative to the e-beam pulse.

In order to achieve a homogeneously excited laser volume, we reduced our normal cell volume with a solid, stainless steel insert having a 1-cm-wide slot running its full 10-cm length. As viewed from the ends, the laser mirrors completely covered the optically accessible area of the slot (≈ 1 cm²).

Photographs demonstrated that the laser beam profile filled this square aperture and exhibited a uniform-intensity cross section. A small hole was drilled transversely through the slotted insert, to provide a viewport for observing fluorescence from within the laser volume (see Fig. 8-42). Pinholes were used in all fluorescence measurements to restrict the viewing region of the vacuum photodiodes, so that only gas regions and not metal surface were observed.

Determining the energy deposition by the electron beam was a major uncertainty in laser efficiency measurements. We designed our experiment so that the excitation volume and the e-beam current incident on the gas were well defined. Energy deposition data were provided by SANDYL, which was used with the Faraday cup measurements to give a normalization of the electron flux (see Fig. 8-44). In the SANDYL calculations, the diode foil (5-mil titanium) and the cell foil (3-mil Inconel) were placed together at the entrance of the cell to produce the appropriate scattering angle to the e-beam. In separate calculations we determined the fraction of the e-beam current incident on the foils which should be collected by the Faraday cup as experimentally apertured, and we determined the fraction of the energy of the 1-MeV beam which was deposited in various regions. As shown in Table 8-3, roughly 1% of the energy incident on the slot area was deposited in the gas. Moreover, presumably because of backscatter from the walls, the deposition was roughly twice that expected from the one-dimensional stopping power of Berger and Selzer.⁴⁴

We also used a pressure-rise (ΔP) technique to determine the energy deposition. For a thermally-isolated monatomic gas, ΔP is related to the deposited energy, ΔE , according to

$$\Delta P = \frac{2}{3} \frac{\Delta E}{V_{\text{cell}}} \quad (12)$$

Again we assumed the deposition volume was limited to the slotted region.

The ΔP technique, first devised at LLNL several years ago, has become a standard measurement of e-beam energy deposition within the laser community. This method suffers from several unknowns, however: whether energy in the gas is gained from or lost to the hot foil or cell walls; how to account for energy lost to radiation; and

Deposition region	Energy deposition (%)
Backscatter	9.7
Diode foil	13.8
Cell foil	15.9
Gas	1.07
Walls (top and bottom)	35.8
Wall (back)	22.8
End windows	0.91

^a Assuming a 1-MeV beam and a mixture of 1600-Torr Ar and 100-Torr Xe.

Table 8-3. E-beam energy deposition in MEG I experiments, based on SANDYL calculations.^a

Fig. 8-44. Comparison of calculated and measured energy deposition in the laser volume. ▼

whether energy (in the case of molecular gases) is contained in dissociation products and, perhaps, in molecular vibrations. Despite our reservations, we employed the technique using a fast-response, piezo-resistive transducer to measure ΔP 10 ms after the e-beam shot.

Results of the SANDYL calculation and the ΔP measurement are shown in Fig. 8-44, for two Ne-based and two Ar-based mixtures which were used in laser experiments. The higher set of points corresponds to the full current, 490 ± 100 A/cm²; the lower set corresponds to the attenuated current, 145 ± 20 A/cm². The solid lines connect deposition values determined with SANDYL calculations and ΔP measurements for the same conditions. Given the ranges of uncertainty in the e-beam currents, and the questions about interpreting the

pressure-rise, the agreement between these techniques is remarkably good, particularly for the Ar case.

XeCl Fluorescence Studies. We studied the time-dependence of the XeCl B-state density using a 10-nm-wide, 310-nm filter to isolate the 308-nm, B-X fluorescence. To reduce the number of variables in the experiments, we employed Ne-based mixtures with total pressures of 1700 and 3400 Torr and Ar-based mixtures with total pressures of 850 and 1700 Torr. As indicated in Fig. 8-44, the energy deposition was approximately equal in the Ne and Ar cases, because of their comparable mass densities. We determined the effect of amplification on the observed spontaneous emission by comparing the temporal dependence of fluorescence in the longitudinal and transverse directions. For current densities above 100 A/cm², we found it necessary to reduce the excited length to 2.5 cm (using e-beam-blocking plates) to eliminate the effects of amplification.

Figure 8-45 compares the peak fluorescence signals as a function of Cl-donor gas pressures for HCl, Cl₂, CCl₄, and CCl₃F, for a common mixture of 3400-Torr Ne and 66-Torr Xe, and a standard current density (145 A/cm). The XeCl amplitudes presumably declined at high donor pressure because of quenching. Based on the experimental data on XeCl quenching by HCl and Cl₂ (Ref. 41), the quenching rate is comparable to the spontaneous radiation rate for HCl and Cl₂ pressures of 2.4 and 3.2 Torr, respectively. Quenching data for CCl₄ and CCl₃F were not available. Note also that the fluorescence intensity was greatest using HCl. This supports our contention that the greater production rate of XeCl is the major benefit of using HCl rather than Cl₂ as the Cl donor.

The peak XeCl fluorescence as a function of Xe pressure is shown in Fig. 8-46, for a mixture of 3.3-Torr HCl and 3300-Torr Ne excited by an e-beam current of 145 A/cm². Observe that the fluorescence was relatively insensitive to Xe for pressures near the optimum at about 70 Torr. The XeCl quenching rate by Xe (k_q) equals the spontaneous radiation rate (k_{rad}) at a Xe pressure of approximately 200 Torr.

We also studied the XeCl fluorescence in Ar and Ne mixtures, and as a function of diluent pressure. Roughly normalizing the

Fig. 8-45. Peak XeCl fluorescence vs Cl-donor gas pressure.

Fig. 8-46. Peak XeCl fluorescence vs Xe pressure, for a mixture of 3.3-Torr HCl and 3300-Torr Ne excited by an e-beam current of 145 A/cm².

energy deposition according to the Berger and Selzer stopping power, we found no major differences in the fluorescence from Ar-based and Ne-based mixtures. Also, the XeCl fluorescence efficiency simply showed a gentle decline for diluent pressures above 1 to 2 atm.

Figure 8-47 shows time-dependences of XeCl B-X fluorescence from mixtures of 3300-Torr Ne, 66-Torr Xe, and 3.3 Torr of either HCl (left) or Cl₂ (right), excited by three different e-beam current densities. The oscilloscope sensitivity per division is shown in the upper right-hand corner of each picture; the timing can be compared directly to the current pulses in Fig. 8-43. Note first that the fluorescence always lags behind the current pulse; this lag is particularly noticeable at low currents, and persists in Cl₂-mixtures to the highest current. In the Cl₂ case there is also a long fluorescence tail. In the case of the HCl mixture at the higher currents, the fluorescence responds more rapidly to the current pulse. Also with HCl, at the highest current density the fluorescence declines significantly during the pump pulse. Notice generally there is a characteristic increase in fluorescence which occurs at high current as the e-beam terminates.

Most of the features in Fig. 8-47 can be understood qualitatively by noting that the main bottleneck to XeCl production is the electron attachment rate to form Cl⁻. According to published attachment coefficients for Cl₂ and HCl (approximately 1 to $2 \times 10^{-10} \text{ cm}^3 \text{ s}^{-1}$ for an electron temperature around 1 eV),^{45,46} the attachment time is about 50 ns at 3.3 Torr. This bottleneck is responsible for the generally slow response of XeCl fluorescence. In the case of HCl at higher current densities, vibrational excitation causes an increased attachment rate and thus a faster XeCl response. This leads to higher signal levels for the HCl mixtures compared to Cl₂, and ultimately leads to the burnout of HCl, causing the fluorescence to decline during pumping. We attribute the fluorescence peak at the end of current pulse to a reduction in electron-quenching^{47,48} of XeCl while its production persists in the afterglow.

Figure 8-48 shows the peak XeCl fluorescence as a function of e-beam current. We have plotted the peak fluorescence divided by the current, to focus on its

Fig. 8-47. Variation of ▲ XeCl-fluorescence time-dependence with Cl donor and current density.

Fig. 8-48. Maximum XeCl fluorescence normalized to the current density, as a function of ▼ e-beam current.

nonlinear behavior. The data were taken with mixtures of 3300-Torr Ne, 66-Torr Xe, and 3.3 Torr of either HCl, Cl₂, or CCl₄. The rise in normalized fluorescence with current for the HCl case is presumably caused by increased attachment resulting from vibrational excitation. The general decline in the Cl₂ case is presumably caused by electron deactivation of XeCl. This is particularly problematic with Cl₂ because its electron attachment rates are low. There is no clear explanation for the behavior with CCl₄.

XeCl Laser Experiments. We were interested in determining the XeCl laser

Fig. 8-49. XeCl laser output energy vs HCl pressure for two e-beam current densities.

Fig. 8-50. Time-dependence of XeCl laser intensity of a 1700-Torr Ar-Xe-HCl mixture.

performance as a function of current density, total gas pressure, and as a function of the diluent gas. At each condition, one needs in principle to optimize both the gas mixture and the mirror reflectivities to produce maximum output. Given the large number of variables, however, we obviously had to make some choices to reduce the parameter space. As previously explained, we chose to study Ne-Xe-HCl mixtures at total pressures of 1700 and 3400 Torr and Ar-Xe-HCl mixtures at 850 and 1700 Torr. The choice of HCl was based on our fluorescence studies and previous work. We also used two current densities, 490 and 140 A/cm², which produced reasonable gains and hence rapid laser-intensity buildup.

An example of the optimization of the laser output with HCl pressure is shown in Fig. 8-49. The gas mixture also contained 63 Torr of Xe and sufficient Ar to give 1700 Torr total pressure. The laser cavity consisted of a maximum-reflectivity mirror and a 30%-reflectivity mirror. Additional HCl appears to be necessary at higher current densities, because of HCl burnout. This was also apparent from the laser intensity time-dependence shown in Fig. 8-50 for a 1700-Torr Ar-Xe-HCl mixture including 63 Torr of Xe and either 5 Torr (a) or 9 Torr (b) of HCl, pumped by 490 A/cm². The intensity scale is constant. Note the premature laser termination at low HCl pressures.

The summary in Fig. 8-51 shows the optimum HCl pressure vs deposited energy, as calculated by SANDYL. Data were taken for 490 A/cm² and 145 A/cm² with Ar mixtures at 1700 Torr and Ne mixtures at 3400 Torr. The central curve plots the optimum HCl pressure which produced maximum output. The upper and lower curves plot the HCl pressure where we obtained an energy 90% of the optimum. The optimization with Xe was somewhat more complex because the energy deposition changed when we added Xe. To account for this effect we normalized the laser output to the relative stopping powers given by the Berger and Selzer values. A summary of Xe optimizations is given in Fig. 8-52.

The dependence of the laser output on HCl and Xe is generally similar to the fluorescence behavior. A minimum HCl pressure is required for attachment, while more HCl becomes necessary at high

currents because of depletion. At the same e-beam stopping powers, Ne and Ar mixtures show roughly the same dependences on mixture components. The optimum Xe pressures for laser operation are somewhat larger than those found in fluorescence studies. The slow increase of the optimum Xe pressure with pumping current has no obvious explanation. Given these variations, it is clear there is no standard XeCl laser mixture. The optimum mixture depends both on pumping current and pulse length.

Using optimized gas mixtures, we varied the mirror reflectivity to maximize the output energy. Figure 8-53 shows the output energy as a function of the round-trip output coupling, $1 - R_1 R_2$. Optimized HCl and Xe pressures were chosen using the data in Figs. 8-51 and 8-52 for Ar mixtures at 850 Torr and Ne mixtures at 1700 Torr. The optimum output coupling was obviously quite different, depending on the gain and loss of the medium as altered by the total pressure and the pump current.

The final results of these several optimizations are summarized in Table 8-4. We found a maximum specific energy for the XeCl laser of 11.6 J/liter, for a 1700-Torr, Ar-based mixture. The output fluence from each end of the cavity was approximately 11.6 MW/cm². Based on the energy-deposition values calculated with SANDYL, the optimum XeCl laser efficiency was 4.5%, for an Ar-based mixture operating at 850 Torr and pumped by 145 A/cm². The efficiency of Ne mixtures was clearly lower. Note that there was a small but clearly distinguishable penalty in efficiency caused by using the larger pumping current of 490 A/cm².

We also determined the degree of XeCl B-state saturation produced by various intra-cavity fluences. To observe this saturation, we compared the XeCl sidelight with and without cavity mirrors. Figure 8-54 shows that installing maximum reflectivity mirrors caused only about a 50% depletion of the sidelight. We varied the intra-cavity fluence by changing the reflectivity of one laser mirror, and measured the intra-cavity fluences using a calibrated photodiode and the known transmissions of the mirror and neutral-density filter.

Figure 8-55 shows the sidelight intensity (B-X) as a function of the intra-cavity power. The solid lines represent approximate

fits to a homogeneous saturation model. The gas mixtures were 8-Torr HCl, 200-Torr Xe, and 1500-Torr Ar excited by 490 A/cm² (upper curve), and 4-Torr HCl, 125-Torr Xe, and 1570-Torr Ar excited by 145 A/cm² (lower curve). Note that the observed

Fig. 8-51. Optimum HCl pressure for maximum XeCl laser output vs calculated energy deposition.

Fig. 8-52. Optimum Xe pressure for maximum XeCl laser output vs calculated energy deposition.

Fig. 8-53. Laser output vs round-trip output coupling.

Total pressure (Torr)	Current density (A/cm ²)	Energy deposition ^a (mJ/cm ³)	Output energy (mJ/cm ³)	Medium efficiency (%)
Ar-based mixtures				
1700	490	270	11.6	4.3
1700	145	270	3.44	4.3
850	490	149	4.84	3.2
850	145	44	1.98	4.5
Ne-based mixtures				
3400	490	323	10.4	3.2
3400	145	96	3.71	3.9
1700	490	165	3.62	2.3
1700	145	47	1.61	3.4

^aBased on SANDYL calculations.

Table 8-4. Summary of XeCl laser results.

saturation is unlike that expected for saturation of a homogeneous line. Because a similar sidelight dependence has been observed for the KrF laser³⁰ (see Fig. 8-71, in the following article), we are investigating

the possibility that incomplete vibrational coupling may generally play a role in the extraction efficiency of RGH lasers.

Absorption in E-Beam-Excited Gases. To contribute to our kinetic studies of RGH lasers, we have developed during the past year a probe laser technique to measure both the spectrally "on-line" gain and spectrally "off-line" absorption. A schematic diagram of the experimental set-up is shown in Fig. 8-56. An injection-locked unstable-resonator KrF laser produced 100 mJ at 249 nm in a well-collimated beam. A 50-cm-focal-length lens focused this beam to the center of a 75-cm cell containing either H₂ or D₂ at approximately 7 atm pressure. Stimulated Raman scattering (SRS) in the gas produces four Stokes lines and one anti-Stokes line, having energies of several percent of the input.⁴⁹ The wavelengths of the various vibrational Raman lines are given in Table 8-5.

We used the total of ten SRS lines obtained from H₂, D₂, and the KrF pump light to probe the absorption in e-beam-excited gas. The SRS output was first

Fig. 8-54. Change in XeCl sidelight caused by installing maximum reflectivity mirrors.

collimated at the exit of the Raman cell, and then passed through a 0.5-cm-diam pinhole defining a narrow beam. Part of the beam was split off for reference purposes and detected with a vacuum photodiode attached to a spectrometer to select the line of interest. The main fraction of the probe beam was attenuated, passed through the e-beam-excited gas cell, and detected with a second spectrometer-photodiode combination. To obtain proportional signals on the transmission and reference photodiodes, we found it necessary to sample the entire probe beam using collecting lenses, rather than to sample merely the small areas defined by the spectrometer slit openings.

Time-dependent absorption values were determined by appropriately delaying the discharge laser. The KrF laser and the MEG I e-beam operated with a relative jitter of approximately ± 10 ns. The pulse lengths produced by the probe source were approximately 15 ns for the pump line and 10 ns for the SRS lines.

As shown in Table 8-5, the SRS-probe wavelengths varied from 220 nm to 420 nm. A spectral range of at least this width is necessary to observe the extremely broad, bound-free absorption features resulting from the dissociative absorption continua of rare-gas molecular ions, rare-gas excimers, and triatomic RGH molecules, as well as the photoionization of excited rare-gas atoms and halogen negative ions. These species have been identified in the laser modeling as the dominant absorbers in RGH laser media (see Tables 8-7 and 8-8 in the following article). Because the predicted absorption bands are very broad, this "comb" of widely-spaced lines should be adequate to determine the overall spectrum of absorption.

The SRS probe has several advantages. First, it has a large bandwidth. Second, the SRS lines are sufficiently bright that absorption can be measured even in the presence of strong, e-beam-produced fluorescence. Although a flashlamp offers good wavelength coverage, it is generally not bright enough to probe e-beam-excited media. On the other hand, frequency-doubled dye lasers are bright enough and have the advantage of continuous tuning, but can only cover the required wavelength range with difficulty, using numerous dyes and laser optics.

To provide a basis for planned absorption studies of RGH laser mixtures, we studied the absorption properties of pure Ar and Ne. Figure 8-57 gives absorption coefficients

D ₂ scattering wavelength (nm)	H ₂ scattering wavelength (nm)
231.3	225.2
248.5 (pump)	248.5 (pump)
268.5	277.1
291.9	313.2
319.8	360.0
353.6	423.3

Table 8-5. Probe laser wavelengths obtained by stimulated Raman scattering of KrF laser light in H₂ and D₂.

Fig. 8-55. XeCl sidelight intensity vs intracavity power, as viewed through one mirror.

Fig. 8-56. Experimental setup to measure absorption in e-beam-excited gases.

obtained for Ar at 6.4 atm, excited by 145 A/cm^2 ; the coefficients were determined at the end of the 50-ns excitation pulse. Note that the peak absorption coefficients were about 0.35 cm^{-1} , allowing approximately 3% transmission through the 10-cm gas cell. Such large net absorptions reduce the uncertainty in determinations of absorption coefficients caused by uncertainty in the input probe intensity. The absorption spectrum of Ar has a peak at $295 \pm 10 \text{ nm}$ and a FWHM of 100 nm. The general shape of the absorption spectrum is

Fig. 8-57. Measured absorption coefficients in e-beam-excited Ar, and the calculated Ar_2^+ cross section, vs wavelength.

Fig. 8-58. Measured absorption coefficient in e-beam-excited Ne and the calculated Ne_2^+ cross section, vs wavelength.

in good agreement with calculations by Wadt of the $2\Sigma_u^+ \rightarrow 2\Sigma_g^+$ transition of Ar_2^+ (Ref. 50), shown as a solid line in Fig. 8-57.

The corresponding absorption spectrum of Ne is shown in Fig. 8-58, for the same conditions as in Fig. 8-57. Note that the Ne absorption coefficients are about a factor of 4 smaller than those of Ar. To obtain these data we used a three-pass arrangement (with a 30-cm path length), giving net absorptions comparable to the Ar case. The solid line is the $2\Sigma_u^+ \rightarrow 2\Sigma_g^+$ cross section of Ne_2^+ calculated by Wadt.⁵⁰

Despite the similarity of our data to the calculated molecular-ion absorption spectra, there are questions remaining about their interpretation. According to ab initio calculations done here at LLNL,⁵¹ the absorption spectrum of the Ar_2^+ excimer should closely resemble Ar_2^+ . In a simple view, this follows because Ar_2^+ can be viewed as a Rydberg electron attached to an Ar_2^+ core; the Ar_2^+ core retains the absorption properties of the free molecular ion. More detailed calculations on the absorption properties of Ar_2^+ are now being done.

To aid in the identification of these absorption features, we investigated both the time-dependence and pressure-dependence of the e-beam-excited absorption at 268.5 nm. The time-dependence of the Ar absorption for the conditions of Fig. 8-57 is shown in Fig. 8-59. This time-dependence disagrees qualitatively with that expected for Ar_2^+ . Note that the observed absorption increases during the e-beam, and reaches steady state only near the end of the current pulse. According to the following argument, however, we expect the Ar_2^+ density to reach steady state quickly and to follow the pump closely. Wadt's cross section and the measured absorption imply an Ar_2^+ density of 10^{16} cm^{-3} . At an electron temperature of 1 eV, the electron recombination rate coefficient is approximately $6 \times 10^{-8} \text{ cm}^3 \text{ s}^{-1}$ (Ref. 52). Thus for an electron density of 10^{16} cm^{-3} steady state should be achieved in less than 2 ns.

The pressure-dependence of absorption also disagrees with that expected for Ar_2^+ . Figure 8-60 gives the pressure-dependence of the absorption coefficient in Ar at 268.5 nm, at current densities of 145 and 24 A/cm². Data were taken for both 10-cm and 30-cm path-lengths. The solid lines merely connect the data points; the dotted lines are the calculated absorptions based on the steady-

state molecular-ion density. The absorption shows a roughly p^2 dependence at low pressures, and a slower dependence at high pressures.

To interpret these data we used a simple kinetic model to describe the molecular-ion density:

$$\frac{dX^+}{dt} = \beta IX - kX^2X^+ \quad (13)$$

$$\frac{dX_2^+}{dt} = kX^2X^+ - \alpha n_e X_2^+ \quad (14)$$

and

$$n_e = X^+ + X_2^+ \quad (15)$$

where X^+ and X_2^+ are the atomic-ion and molecular-ion number densities, and n_e is the electron concentration. The rate of power-deposition calculated from the Berger and Selzer stopping power for Ar implies a value for β of $22 \text{ s}^{-1} \text{ cm}^2/\text{A}$, assuming 27 eV is required per electron-ion pair. The three-body formation rate, k , was assumed to be $2.5 \times 10^{-31} \text{ cm}^6 \text{ s}^{-1}$ (Ref. 53), and the recombination rate, α , to be $5.9 \times 10^{-18} \text{ cm}^3 \text{ s}^{-1}$ (Ref. 52). This model neglects the production of ions by electron-impact ionization of excited states, a secondary process which can be important at high-excitation densities.

Solving Eqs. (13), (14), and (15) for steady state, we find

$$X_2^+ = \frac{-\beta I}{2kX} + \frac{I}{2} \left[\left(\frac{\beta X}{I} \right)^2 + \left(\frac{4\beta IX}{\alpha} \right) \right]^{1/2} \quad (16)$$

Although the solution is hardly transparent, it is easy to show that at high pressures (where $n_e = X_2^+$) we have

$$X_2^+ = \left(\frac{\beta IX}{\alpha} \right)^{1/2} \quad (17)$$

At low pressures (where $n_e = X^+$) we have

$$X_2^+ = \frac{kX^2}{\alpha} \quad (18)$$

Thus at low pressures the steady-state molecular-ion density is independent of current and varies as pressure squared while at high pressures it scales as the square root of current times pressure.

Although the experimental absorption data appear qualitatively to fit a p^2

dependence at low pressures, the $p^{1/2} I^{1/2}$ dependence is not established. For comparison with experimental data, the model predictions of the steady-state absorption are shown in Fig. 8-60 as dotted lines. The observed absorption appears to

Fig. 8-59. Time-dependence of the absorption coefficient in e-beam-excited Ar.

Fig. 8-60. Pressure-dependence of the absorption coefficient in e-beam-excited Ar.

grow more strongly with pressure and current than expected.

We now believe that the greater-than-expected absorption at high pressure is caused by the contribution of Ar_2^* , primarily in the triplet state. We have recently calculated that the time dependence of the Ar_2^* triplet closely resembles that of Fig. 8-59. Moreover, if we lump together the absorption from both Ar_2^+ and Ar_2^* and assume equal cross sections of $3 \times 10^{-17} \text{ cm}^2$, the predicted absorption curve is parallel to the experimental one and about 50% larger than measured. Given the uncertainties, we consider this fit to be quite good. In 1981 we plan to model the rare-gas data more completely and measure the absorption spectra of RGH laser mixtures.

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Rare-Gas-Halide Kinetics Theoretical Studies

In the continuing search for fusion-laser candidates, we have investigated the characteristics of xenon chloride (XeCl) and argon fluoride (ArF) lasers. These systems are similar to krypton fluoride (KrF) lasers, but they each possess characteristics that may be improvements. XeCl operates at a longer wavelength (308 nm vs 248 nm for KrF), where optical components should be more resistant to damage. Some experiments have yielded XeCl laser efficiencies comparable to KrF. Although ArF operates at a shorter wavelength than KrF (193 nm vs 248 nm) and hence may present more severe optics problems, the quantum efficiency of ArF should be higher and the nonsaturable losses in the medium should be smaller than for KrF. Hence, ArF may be a substantially more efficient laser than KrF.

In addition to studying XeCl and ArF as potential fusion lasers, we have investigated physics issues associated with the formation of the upper laser level in RGH lasers. The lasing species (KrF, ArF, XeCl) are formed primarily by recombination of

positive and negative ions. Recombination is a complex pressure-dependent and plasma-dependent process for which we have no direct experimental data; however, we have been able to simulate this process computationally and study its dependence on pressure, temperature, and ion density. Once the excimer species has been formed by recombination, it must relax through hundreds of vibrational states before it can contribute to the laser field. The amount of energy available for extraction from the energy that goes into the formation of the excimer species depends upon the vibrational relaxation kinetics of the system. We have investigated this problem for a special system and have been able to contribute to its understanding.

XeCl Laser Kinetics. Figure 8-61 shows schematically what we believe to be the XeCl kinetic chain, and Table 8-6 lists the rate coefficients of the reactions in this chain. We believe, in general, that the formation kinetics of XeCl are dominated by ionic processes. The three-body charge transfer process,

has a very fast rate, $3.9 \times 10^{-30} \text{ cm}^6/\text{s}$ (see Table 8-6), and is a key step in the formation of Xe^+ . We believe the molecular ion NeXe^+ to be so weakly bound that it has little influence on this reaction scheme.

The formation rate of XeCl appears to be dominated by the attachment characteristics of the Cl^- donor. Cl_2 has a small electron-attachment cross section,⁵⁴ so that attachment is the rate-limiting step in the formation of XeCl. Cl_2 can, however, undergo a harpooning reaction,

which helps to contribute to XeCl formation. HCl, which has a small attachment cross section for the $v = 0$ state, has very large cross sections for vibrational excitation⁵⁵ and for dissociative attachment to vibrationally excited states.⁵⁶ Dissociative-attachment cross sections for HCl ($v = 0, 1, 2$) that have been deconvolved from the measurements by Allan and Wong⁵⁶ on heated HCl are shown in Fig. 8-62.

Powell and Poli⁵⁷ have found experimentally that at low e-beam current density ($\sim 10 \text{ A/cm}^2$) the time-dependence and

intensity of XeCl fluorescence is about the same with either Cl₂ or HCl donors. The attachment rates for Cl₂ and for HCl ($v = 0$) are nearly the same. At high current densities ($\lesssim 100 \text{ A/cm}^2$), however, XeCl is a much better fluorescer with HCl than with Cl₂ (the experimental data are given in "Rare-Gas-Halide Kinetics Experiments" later in this section). Because of the slow Cl₂ attachment rate, electron densities should be higher than those obtained using HCl as a donor. We believe the decline in fluorescent intensity at high current in the mix with Cl₂ is due to electron quenching of XeCl. The peaking of the XeCl fluorescence after the termination of the e-beam suggests electron quenching during the e-beam pulse. The fluorescent intensity obtained at high currents with HCl appears to be degraded by burnout of the Cl⁻ donor. Our computer model of XeCl kinetics, which includes the reactions and rate coefficients of Fig. 8-61 and Table 8-6, agrees well with the experimental data and supports the conclusions discussed above.

With regard to production rate and fluorescence efficiency, XeCl should perform less well than KrF because of lower quantum efficiency, a smaller attachment rate (with subsequent electron quenching), and the reduction of available neutral channels. Laser efficiency is affected not only by these considerations but also by the characteristics of the absorbers. In RGH lasers at low current density, we expect the absorbers typically to be the halogen-negative ion and the rare-gas dimer ions.⁵⁸ What constitutes "low" and "high" current densities depends upon the electron density, which in turn depends upon the halogen donor and its attachment rate. At low current densities in XeCl laser mixtures, the absorbers are probably Ne₂⁺ and Cl⁻. This interpretation is consistent with measurements that have been made.⁵⁹ We list the absorption cross sections for ionic absorbers at 3080 Å in Table 8-7. Computer models indicate that at high current densities, the major absorbers will be rare-gas excited states—in the np manifold and above. Table 8-7 lists some of these cross sections. Calculations have shown that the absorption cross sections of the lowest rare-gas ns states are small.⁶⁰

Although our discussion of XeCl deals only with the kinetics of Ne–Xe–HCl or Cl₂ mixtures, Powell⁵⁷ has found that the XeCl laser works as well with an Ar or Kr buffer as it does with Ne (see the preceding article).

The kinetics with Ar or Kr buffers may be somewhat different, however, with respect to the Ar₂⁺ and Kr₂⁺ charge transfer reactions, the role played by ArXe⁺ or KrXe⁺, and the absorptions of the dimer ions Ar₂⁺ and Kr₂⁺

Fig. 8-61. Schematic of XeCl kinetic chain.

Reaction	Rate coefficient	Reference
$\text{Ne}^+ + 2\text{Ne} \rightarrow \text{Ne}_2^+ + \text{Ne}$	$4.5^{-32} \text{ cm}^6/\text{s}$	a
$\text{Ne}_2^+ + \text{Xe} + \text{N} \rightarrow \text{Xe}^+ + 3\text{Ne}$	$3.9^{-30} \text{ cm}^6/\text{s}$	b
$\text{Xe}^+ + \text{Cl}^- + \text{Ne} \rightarrow \text{XeCl} + \text{Ne}$	Function of pressure	c
$\text{Xe}^* + \text{Cl}_2 \rightarrow \text{XeCl} + \text{Cl}$	$7.2^{-10} \text{ cm}^3/\text{s}$	d
$\text{Xe}^* + \text{HCl} \rightarrow \text{prod.}$	$5.6^{-10} \text{ cm}^3/\text{s}$	d
$\text{Xe}^{**} + \text{Cl}_2 \rightarrow \text{XeCl} + \text{Cl}$	$7.2^{-10} \text{ cm}^3/\text{s}$	Assumed
$\text{Xe}^{**} + \text{HCl} \rightarrow \text{XeCl} + \text{H}$	$5.6^{-10} \text{ cm}^3/\text{s}$	Assumed
$\text{XeCl} \rightarrow \text{Xe} + \text{Cl} + h\nu$	6.2^7 s^{-1}	e
$\text{XeCl} + \text{Cl}_2 \rightarrow \text{prod.}$	$6.0^{10} \text{ cm}^3/\text{s}$	e
$\text{XeCl} + \text{HCl} \rightarrow \text{prod.}$	$8.0^{-10} \text{ cm}^3/\text{s}$	e
$\text{XeCl} + \text{Xe} + \text{Ne} \rightarrow \text{Xe}_2\text{Cl} + \text{Ne}$	$1.0^{-31} \text{ cm}^6/\text{s}$	Guessed
$\text{XeCl} + h\nu \rightarrow \text{Xe} + \text{Cl} + 2h\nu$	3.6^{-16} cm^2	f
$\text{XeCl} + e^- \rightarrow \text{prod.}$	$\approx 2^{-8} \text{ cm}^3/\text{s}$	Estimated from kinetic model and experimental data

^aO. J. Orient, *Chem. Phys. Lett.* 23, 579 (1973).

^bC. B. Collins and F. W. Lee, *J. Chem. Phys.* 72, 5381 (1980).

^cJ. N. Bardsley and J. M. Wadehra, *Chem. Phys. Lett.* 72, 477 (1980).

^dJ. H. Kolts, J. E. Velazco, and D. W. Setser, *J. Chem. Phys.* 71, 1247 (1979).

^eC. H. Fisher, R. E. Center, and J. P. McDaniel, 32nd Gaseous Electronics Conference, University of Oklahoma, Norman, Okla. (1979).

^f $\sigma_{\text{tr}} = 50 \text{ Å}^2\text{-ns}$; C. A. Brau and J. J. Ewing, *J. Chem. Phys.* 63, 4640 (1975).

Table 8-6. Rate coefficients of reactions in XeCl kinetic chain.

Fig. 8-62. HCl and Cl₂ dissociative-attachment cross sections.

Fig. 8-63. Dependence of XeCl amplifier parameters on e-beam pump power density.

(see Table 8-7 for calculated cross sections at 3080 Å for the dimer ions). Due to its greater stopping power, an Ar buffer allows us to work at lower total pressure. Although the dimer-ion absorption is greater using Ar or Kr buffers, the production efficiency may be greater and thereby compensate for the former effect.

Figure 8-63 displays model calculations of the dependence of XeCl amplifier parameters on e-beam pump power density. We used a total pressure of 1 atm and optimized Xe and HCl partial pressures. Note that the production efficiency, η_{prod} , rises to a peak as HCl vibrational excitation and dissociative attachment rates, which limit the XeCl production rate, increase; η_{prod} then decreases due to HCl burnout at high current densities.

The production efficiency for XeCl is much smaller than the values we have predicted for ArF or KrF. This is due to the smaller halogen attachment rate and the reduction in neutral channels available for XeCl production. The fluorescence efficiency, however, generally decreases due to electron quenching of XeCl. The nonsaturable loss, γ_{ns} , due primarily to rare-gas excited states, increases almost linearly with pump power. The gain-to-loss ratio is very large, due to the large XeCl stimulated-emission cross section. This large cross section, the long radiative lifetime, and the

Species	Cross section at 3080 Å (cm ²)	Reference
Cl ⁻	2.1 ⁻¹⁷	a
Ne ₂ ⁺	1.9 ⁻¹⁷	b
Xe ₂ ⁺	2.5 ⁻¹⁷	b
Xe ₂ ⁺	1.4 ⁻¹⁷	c
Ne(3p)	2.3 ⁻¹⁸	d
Ne(3d)	0.6 ⁻¹⁸	d
Xe(6p)	7.6 ⁻¹⁸	d
Xe(5d)	9.3 ⁻¹⁸	d
Ar ₂ ⁺	4.5 ⁻¹⁷	e
Kr ₂ ⁺	5.4 ⁻¹⁷	e

^aD. E. Rothe, *Phys. Rev.* 177, 93 (1969).

^bL. A. Levin, S. E. Moody, E. L. Klosterman, J. J. Ewing, and R. E. Center, 33rd Gaseous Electronics Conference, University of Oklahoma, Norman, Okla. (October 1980).

^cC. A. Brau, "Rare Gas Halide Lasers," in *Excimer Lasers*, C. K. Rhodes, Ed. (Springer-Verlag, Berlin, 1978).

^dC. Duzy and H. A. Hyman, *Phys. Rev. A* 22, 1878 (1980).

^eH. H. Michels, R. H. Hobbs, and L. A. Wright, *J. Chem. Phys.* 71, 5053 (1979).

Table 8-7. Absorbing species in XeCl lasers.

longer wavelength combine, however, to result in a much smaller saturation intensity for XeCl than for KrF. I_s does increase slightly with pump power, due to electron quenching of XeCl.

We conclude then that, in comparison to ArF or KrF, the XeCl laser is severely limited in efficiency and, due to its small saturation intensity, cannot be scaled to high intensity with large extraction efficiency. As a final note of optimism, however, we point out that Powell⁵⁷ has some experimental evidence of XeCl saturation intensity much higher than we have calculated here. We do not yet know the reason for the discrepancy between our calculated value and Powell's experimental data.

ArF Laser Kinetics. We expect the ArF kinetic scheme in an Ar-F₂ mixture to be very similar to that found in Ar-Kr-F₂ mixtures⁶¹ up through the formation of ArF via either the ionic channel or the neutral channel. The production efficiency of ArF should be some 28% higher than that of KrF, due to the higher quantum efficiency of ArF. Because of the shorter radiative lifetime of ArF (4 ns vs 9 ns for KrF), we expect ArF to be a more efficient fluorescer at higher current densities. The shorter lifetime and slightly larger stimulated-emission cross section ($\sigma\tau = 12 \text{ Å}^2\text{-ns}$ for ArF vs $17 \text{ Å}^2\text{-ns}$ for KrF) gives ArF a larger saturation intensity parameter than that of KrF. These characteristics of ArF, coupled with the absorption physics, will determine the performance characteristics of ArF as a laser.

In Table 8-8 we list the species that we expect to contribute to absorption of the 1930-Å ArF radiation (see also Fig. 8-64). In an Ar-F₂ mixture, the only dimer ion present will be Ar₂⁺; the wavelength of 1930 Å lies in the far UV wing of the Ar₂⁺ absorption profile, so the cross section of the dimer ion should be small. We expect F₂, F⁻, and the Ar(4p) states to be the major absorbers.

There is a dearth of experimental information in the literature on ArF as a laser. Recent experiments by Chang⁶² yield an ArF fluorescence efficiency of 13% in Ar-F₂ mixtures and 24% in Ne-Ar-F₂ mixtures. ArF works best at low Ar partial pressures, and its performance is only weakly dependent upon Ne pressure. This is not surprising, since Ar is likely to be the most

efficient rare-gas quencher of ArF. Chang has obtained the following quenching rate coefficients for ArF:

- $k_q(\text{Ar}) = 2.2^{-11} \text{ cm}^3/\text{s}$
- $k_q(2\text{Ar}) = 2.7^{-31} \text{ cm}^6/\text{s}$
- $k_q(\text{Ne}) = 2.0^{-12} \text{ cm}^3/\text{s}$
- $k_q(2\text{Ne}) \leq 10^{-32} \text{ cm}^6/\text{s}$

With respect to the quenching process, it is advantageous to use Ne as a buffer, provided it has adequate stopping power. With the addition of Ne, we would expect the kinetics of a Ne-Ar-F₂ mixture to be different from an Ar-F₂ mixture. NeF*, which is bound by 6.36 eV and which should be a precursor to ArF*, emits at 1080 Å and has a radiative lifetime of only 2.6 ns (Ref. 63). Because the

Table 8-8. Absorbing species in ArF lasers.

Species	Cross section at 1930 Å (cm ²)	Reference
F ⁻	6.5^{-18}	a
Ar ₂ ⁺	1.6^{-19}	b
Ar(4s)	1.0^{-20}	c
Ar(4p)	2.7^{-18}	d
F ₂	1.5^{-20}	e
Ne ₂ ⁺	3.6^{-19}	b
Ne(3s)	1.3^{-19}	d
Ne(3p)	3.5^{-19}	d
Ne(3d)	1.0^{-19}	d

^aA. Mandl, *Phys. Rev. A* **3**, 251 (1970).

^bH. H. Michels, R. H. Hobbs, and L. A. Wright, *J. Chem. Phys.* **71**, 5053 (1979).

^cK. J. McCann and M. R. Flannery, *Appl. Phys. Lett.* **31**, 599 (1977).

^dC. Duzy and H. A. Hyman, *Phys. Rev. A* **22**, 1878 (1980).

^eR. K. Stenkeberg and R. C. Vogel, *J. Am. Chem. Soc.* **78**, 901 (1956).

Fig. 8-64. Dependence of ArF amplifier-medium parameters on e-beam pump power density.

Fig. 8-65. Schematic of potential-energy curves relevant to ionic recombination in KrF* (Ref. 65).

Fig. 8-66. Recombination-rate coefficient for Kr⁺ and F⁻.

Figure 8-64 shows model calculations of the dependence of ArF amplifier-medium characteristics on e-beam pump power density. We performed these calculations for an Ar-F₂ mixture at 1 atm total pressure. If we used Ne as a buffer, we would expect the production efficiency to remain about the same and the other parameters to change only slightly. The fluorescence efficiency should increase at low pump powers, due to reduced neutral quenching, but not at higher pump powers, where electron quenching dominates. The production efficiency, gain-to-loss ratio, and saturation intensity of ArF are all very high, especially when compared to XeCl or even to KrF. Note also that the nonsaturable loss is relatively small. The decline in production efficiency, η_{prod} , is caused by F₂ burnout, but it might be improved by the addition of F₂ if quenching and absorption limitations are not yet reached. On the basis of the theoretical modeling completed to date, ArF appears to be a very efficient laser medium, scalable to high output powers.

Ion-Ion Recombination. In e-beam-pumped RGH lasers, the upper laser level is formed primarily by recombination of positive rare-gas ions with negative halogen ions, according to

where X represents a rare gas (such as Ne, Ar, Kr, or Xe), Y represents a halogen (such as F or Cl), and M represents any of the rare gases used as buffers in these lasers. Figure 8-65 shows the relevant potential curves for KrF. To describe the recombination process, we think of the positive and negative ions as approaching each other along the attractive ($Kr^+ - F^-$) curve shown in the figure. The ions lose energy by collisions with neutral atoms as they are drawn toward each other; when their relative energy is less than $\approx 3 kT$, the probability that they will separate again is so small that we say they have recombined. Much work has been done in the past to describe this process analytically. More recently, we have developed numerical Monte Carlo simulations of the recombination process. We described the analytical theories, the limitations of these theories, and the principles of the numerical simulations in last year's annual report.⁶¹

We have used Monte Carlo simulations to study the effects of plasma screening on the ionic recombination rate. The details of the computational techniques are described in Refs. 61 and 66. Figure 8-66 shows the results of these calculations for

Plasma shielding clearly results in a significant reduction in the ionic recombination rate. In an RGH laser plasma, this reduction in the ion loss rate will cause a shift in the equilibrium ionic densities to higher values. Since the rare-gas dimer ions and the halogen-negative ion are the principal absorbing species, the most serious consequence of the reduced recombination rate is to increase the nonsaturable loss in the laser medium.

Using our simulation code, we have also investigated the temperature dependence of the ionic recombination rate. Figure 8-67 displays some results from these calculations for the reaction

In this reaction, the $\text{Kr}^+ - \text{Kr}$ collisions involve charge exchange, which is a very efficient means of transferring energy. Figure 8-67 clearly shows that the recombination rate decreases markedly with increasing temperature. We see that if the performance of RGH lasers is to improve at elevated gas temperatures, those processes that are deleterious to laser performance must have temperature dependences that are even more dramatic than this.

Accessibility of KrF(B) State to Laser Photons. The use of the KrF laser for ICF applications requires high energy-deposition rates and relatively short pulse lengths. These conditions require that vibrational equilibrium be established rapidly so that the energy can be extracted from the low-lying levels. A significant fraction of the laser energy can be lost to processes such as electron quenching or heavy-particle destruction of the KrF^* , if a large population at higher vibrational levels either has sufficiently long radiative lifetimes or is not spectrally accessible to the 248-nm extraction pulse. We are currently using our model for the vibrational quenching of KrF^* ,⁶⁷ along with the calculated

A-coefficients for the individual levels, to determine the steady-state vibrational distribution during extraction.

Figure 8-68 shows the calculated A-coefficients for the lowest five vibrational levels of the B-state of KrF . In each case, the

Fig. 8-67. Temperature dependence of recombination-rate coefficient for Kr^+ and F^- in Kr .

Fig. 8-68. Frequency-dependent A-coefficients for the vibrational levels of the B-state of KrF .

frequency-dependent A-coefficient is superimposed with the A-coefficient for $v = 0$ (in red) to give some measure of the homogeneity of the level. The same scale is used for $v = 0, 1, 2, 3, 4$, and 5 . Figure 8-68 shows that $v = 0, 1, 2$, and 3 can all be extracted with about the same efficiency. Levels 4 and 5 shift to shorter wavelengths such that at $v = 5$ for $\lambda = 248$ nm, $A = 6.5 \times 10^4 \text{ s}^{-1}\text{cm}$, which is only 26% of the value for $v = 0$. Table 8-9 summarizes the values of the A-coefficients for each level at $\lambda = 248$ nm and at the wavelength of the peak value. Table 8-10 gives the vibrational energies relative to $v = 0$ and the integrated A-coefficients and lifetimes for each level.

We can draw several conclusions from the above information.

- Narrow-line extraction will not efficiently access the population in levels above $v = 4$.
- The vibrational relaxation of the higher levels is an important process in the operation of the KrF laser.

Table 8-9. Frequency-dependent A-coefficients.^a

Level	$A_v(\lambda = 248 \text{ nm})$ (s^{-1}cm)	λ_{max} (nm)	$A_v(\lambda = \lambda_{\text{max}})$ (s^{-1}cm)
0	2.5×10^5	—	—
1	2.2×10^5	—	—
2	2.4×10^5	—	—
3	2.2×10^5	—	—
4	1.6×10^5	—	—
5	6.5×10^4	247.0	2.4×10^5
6	7.5×10^3	246.6	2.2×10^5
7	5.4×10^3	246.0	1.9×10^5
8	3.9×10^4	245.5	1.9×10^5
9	7.0×10^4	244.9	1.7×10^5
10	7.5×10^4	244.4	1.5×10^5

^aFrequency-dependent A-coefficients are for the lowest 10 vibrational levels of the B-state of KrF at $\lambda = 248$ nm and for the peak values of the higher levels.

Level	Vibrational energy (cm^{-1})	Integrated A-coefficient	Lifetime (s)
1	0	1.43×10^8	7.0×10^{-9}
2	335.1	1.40×10^8	7.1×10^{-9}
3	663.9	1.38×10^8	7.2×10^{-9}
4	988.0	1.35×10^8	7.4×10^{-9}
5	1308.6	1.33×10^8	7.4×10^{-9}
6	1626.2	1.30×10^8	7.7×10^{-9}
7	1940.8	1.27×10^8	7.9×10^{-9}
8	2252.4	1.25×10^8	8.0×10^{-9}
9	2560.7	1.22×10^8	8.2×10^{-9}
10	2865.8	1.19×10^8	8.4×10^{-9}
11	3167.6	1.17×10^8	8.6×10^{-9}

Table 8-10. Energies and A-coefficients for levels of KrF B-state.

- The large peak values for the A-coefficients of the higher levels suggest that the populations of these levels are accessible to photons of wavelengths other than 248 nm. If a substantial population is bottlenecking in the higher levels, then multiple-line extraction may significantly improve the performance of the KrF laser.

In last year's annual report we presented calculations of the rate coefficients and average energy transfers for vibrational deexcitation of KrF* by charge-transfer collisions.⁶¹ These calculations, which used a hard-sphere model for charge transfer, demonstrated that vibrational relaxation can occur quickly with large energy transfers per collision. We have used this information, along with the data presented above, in a model for vibrational relaxation of KrF* in a Kr buffer. Using rate equations, we have treated 120 states or groups of states of KrF*, including recombination (feeding the highest levels of KrF*), vibrational deexcitation (with detailed balancing for excitation), spontaneous emission, and stimulated emission. This model is very similar to models used in studies of electron-ion collisional radiative recombination.

Figure 8-69 shows some results from this analysis of vibrational relaxation. We plot the relative populations KrF* levels (with pressure as a parameter) vs energy. We used 20 discrete vibrational states and 100 groups of levels up to the dissociation energy in this particular calculation. On this semilog plot, a thermal population distribution would appear as a straight line; the dashed line represents the thermal distribution for $T = 300$ K. The vibrational temperatures of the lowest states decrease with increasing pressure, indicating the more rapid vibrational relaxation. These vibrational temperatures, though, are greater than the gas temperature. For the lowest levels, it is the radiative losses that shift the equilibrium from thermal.

Figure 8-70(a) depicts this type of equilibrium schematically. The higher levels have a very high vibrational temperature, due to the large charge-exchange energy transfers. Because ΔE is large, the Boltzmann factor, $\exp(-\Delta E/kt)$, is small, and the equilibrium is determined by the deexcitation rates [Fig. 8-70(b)]. A graph of the deexcitation rates for each state would show a minimum at some vibrational level

($v = 14$ for the 1-atm case). The rate of deexcitation across this level (or "bottleneck," as it was named early in the study of collisional radiative recombination) determines the effective vibrational relaxation time of KrF^* . Figure 8-69 gives some of these relaxation times.

The saturation intensity associated with this RGH kinetic system at 1 atm is $\phi_s = 1.2 \text{ MW/cm}^2$. If we apply a stimulating fluence, ϕ , and plot the fluorescent intensity (with $\phi = 0$) divided by the intensity (with $\phi > 0$), we obtain the plot shown in Fig. 8-71. The dashed straight line represents $I_0/I = 1 + \phi/\phi_s$, which is what we expect for a homogeneously saturating medium. The solid line shows our results using the model described above. This deviation from homogeneous saturation is similar to that observed by Jacob et al.,⁶⁸ whose data for a KrF laser are shown in red in Fig. 8-71, and by Powell⁵⁷ for XeCl. This deviation is a manifestation of the finite relaxation time and the decreasing stimulated-emission cross section at 248 nm. We did not attempt to model the Jacob data, because their argon-rich mixture involves more complex relaxation physics than does simple charge transfer. Nevertheless, it is clear from Fig. 8-71 that relaxation physics plays an important role in the ultimate performance of RGH excimer systems as amplifiers.

Jacob et al.⁶⁸ have shown, using a simple model, that this deviation from homogeneous saturation can be described by the formula

$$\frac{I}{I_0} = \frac{1 + \phi/\phi_s}{1 + \phi/\phi_s} \quad (24)$$

Here, $\beta = 1 + \theta_0\tau/\tau_v$, where τ is the KrF^* radiative lifetime, τ_v is the relaxation time, and θ_0 is the Boltzmann factor. Note that, in the limit, $\tau_v \rightarrow 0$, and $\beta \rightarrow \infty$. Equation (24) gives the usual scaling of I/I_0 . Their experimental data correspond to $\beta = 4$ to 5, and the calculations presented here correspond to $\beta = 3.6$. If we consider the gain,

$$\alpha = \sum_v \sigma(v) [\text{KrF}^*(v)] \quad (25)$$

we find that it saturates normally; i.e., $\alpha/\alpha_0 = (1 + \phi/\phi_s)^{-1}$. We can show that the effect of the finite relaxation time is to reduce the small-signal gain, $\alpha_0 = \sigma R\tau$, where R is the KrF^* formation rate. Now we have

$$\alpha_0 \cong \left(1 - \frac{1}{\beta}\right) \sigma R\tau \quad (26)$$

Fig. 8-69. Vibrational populations and relaxation times for KrF^* in Kr.

Fig. 8-70. Equilibria: (a) collisional-radiative equilibrium of low-lying vibrational levels and (b) collisional (de)excitation equilibrium of high vibrational levels.

Fig. 8-71. KrF sidelight ratio vs cavity fluence.

so that α_0/γ is approximately a factor of $(1 - 1/\beta)$ smaller than what would be expected based on the gross formation rate of KrF^* . Thus, for $\beta = 4$, there will be a 25% reduction in α_0/α . This reduction will, of course, lead to a corresponding reduction in the predicted laser extraction efficiency.

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Dissociative Electron Attachment and Vibrational Excitation of F_2

A knowledge of the magnitude and energy-dependence of the cross sections for a variety of electron- F_2 scattering processes is important to us in the kinetic modeling of rare-gas fluoride excimer lasers.^{69,70} In particular, the dependence of the processes

and

on both the incident electron energy and the initial vibrational state of F_2 plays a key role in determining the electron energy distribution and F^- concentration in e-beam-excited KrF lasers.

Although the results of several experimental studies of electron attachment to ground-state F_2 are in qualitative accord, we are currently very interested in the dependence of the attachment rate [Eq. (27)] on the vibrational temperature of F_2 , since recent experiments have shown a substantial increase in excimer fluorescence efficiency with increased gas temperature.⁷¹ At present there are no experimental data available on the vibrational-excitation cross sections of F_2 [Eq. (28)]. Therefore, we undertook and completed a theoretical determination of cross sections for both vibrational excitation of, and dissociative electron attachment to, F_2 (Ref. 72).

Our work is novel for several reasons. It is the first theoretical study of dissociative electron attachment that treats both the

electronic and the nuclear motions from first principles and which uses no experimental cross sections as input. We employ rigorous resonance scattering theory, formulated within the Born-Oppenheimer (or the adiabatic nuclei) approximation, to describe the dynamics of the nuclei, and we make only two additional assumptions to simplify the theory. We extracted the required electronic-resonance parameters from large, multiconfiguration wave functions for $\text{F}_2(^1\Sigma_g^+)$ and $\text{F}_2(^2\Sigma_u^+)$, which include both target polarization and electron-correlation effects.

The wave equation governing the motion of the dissociating nuclei ($\text{F} + \text{F}^-$) at total energy, E , is

$$\left[-\frac{1}{2\mu} \frac{d^2}{dR^2} + V(R) + W_E(R) - E \right] \phi_{E v_i}(R) = -U_{v_i}(E - E_{v_i}, R) \quad (29)$$

where $\phi_{E v_i}$ is the nuclear wave function associated with the electronic resonance state, $V(R)$ is the "unshifted" potential energy curve of the negative ion, and $U_{v_i}(\epsilon, R)$ is the matrix element coupling the resonance to the nonresonant scattering solution associated with a vibrational level, v_i . $U_v(\epsilon, R)$ is defined as

$$U_v(\epsilon, R) = \left[\frac{\Gamma(\epsilon, R)}{2\pi} \right]^{1/2} X_v(R) \quad (30)$$

where $\Gamma(\epsilon, R)$ is the electronic resonance width. We assume $U_v(\epsilon, R)$ to be real. $W_E(R)$ is a complex, energy-dependent, nonlocal potential defined by

$$W_E(R) \phi_E(R) = \sum_v \mathcal{P} \int_0^\infty \frac{d\epsilon U_v(\epsilon, R) \int dR' U_v^*(\epsilon, R') \phi_E(R')}{E - \epsilon - E_v} - i\pi \sum_v^{\text{open}} U_v(E - E_v, R) \int dR' U_v^*(E - E_v, R') \phi_E(R') \quad (31)$$

where $E = 1/2 k_v^2 + E_v$ and $1/2 k_v^2$ is the (asymptotic) kinetic energy of the electron. The energy and wave function of the v^{th} vibrational level of the neutral target molecule, F_2 , are denoted by E_v and X_v , respectively. Note that the second sum on the right side of Eq. (31) is restricted to the energetically open channels.

The cross section for vibrational excitation, starting in level v_i and ending in level v_f , is given by the expression

$$\sigma_{v_i \rightarrow v_f} = \frac{(2\pi)^3}{k_{v_i}^2} \left| \langle U_{v_f} | \phi_{E v_i} \rangle \right|^2 \quad (32)$$

For dissociative attachment, the integrated cross section starting in level v_i is given by

$$\sigma_{v_i \rightarrow \text{DA}} = \frac{2\pi^2}{k_{v_i}^2} \frac{K}{\mu} \lim_{R \rightarrow \infty} \left| \phi_{E v_i}(R) \right|^2, \quad (33)$$

where $K^2/2\mu$ is the asymptotic kinetic energy of the dissociative fragments with reduced mass, μ . We constructed the required solution to Eq. (29), $\phi_{E v_i}$, subject to purely outgoing boundary conditions.

To solve Eq. (29), we needed the potential energy curves of both F_2 and F_2^- , as well as the resonance width, $\Gamma(\epsilon, R)$. We used the method of configuration interaction (CI) to calculate these electronic parameters for a number of internuclear distances between 2.4 and 100 a_0 , as described in last year's annual report.⁷³ Our calculated dissociation energies of $F_2(^1\Sigma_g^+)$ and $F_2(^2\Sigma_u^+)$ are 1.77 and 1.19 eV, respectively, which compare favorably with the experimental values of 1.68 and 1.29 eV.

Our wave functions, on the other hand, underestimate the electron affinity of atomic fluorine by 0.4 eV (2.99 eV calculated vs 3.40 eV experimental). Since a very large CI expansion of the wave functions is required to obtain an accurate value for this quantity, we chose instead to shift our calculated potential curve of F_2^- by 0.4 eV to obtain the correct asymptotic limit at infinite internuclear separation. With this adjustment, $F_2(^2\Sigma_u^+)$ becomes autodetaching for $R \leq 2.63 a_0$, whereas the calculated equilibrium internuclear distance of $F_2(^1\Sigma_g^+)$ is 2.68 a_0 . We calculated the resonance width with our CI wave functions, using Stieltjes moment-theory techniques.⁷⁴

Figure 8-72 shows the dissociative-attachment cross sections that we have calculated for the four lowest vibrational states of F_2 . The ground-state cross section rises steeply from threshold, reaching a peak value of 7.1 $\text{\AA}^2$ at 0.16 eV. With increasing initial vibrational excitation, the peak moves to higher electron energies, while the maximum value of the cross section increases from $v = 0$ to $v = 1$ but decreases thereafter.

At an electron energy of 1.0 eV, the effective cross section averaged over the vibrational population increases by 18% as T_{vib} increases from 300 to 500 K, in reasonable accord with experiment.⁷⁵

Figure 8-73 compares the present attachment cross section for the ground vibrational state with the experimental data of Chantry⁷⁶ and the semi-empirical calculations of Hall⁷⁷ and Bardsley.⁷⁸ Hall employed the so-called "boomerang," or complex-potential model,^{79,80} to describe the motion of the nuclei and to calculate both dissociative-attachment and vibrational-excitation cross sections for F_2 . In very

Fig. 8-72. Calculated dissociative-attachment cross sections for the lowest vibrational states of F_2 .

Fig. 8-73. Comparison of theoretical and experimental dissociative-attachment cross sections for F_2 .

recent work, Bardsley et al. used a more rigorous treatment of nuclear motion and explicitly took into account the energy dependence of the resonance width.⁸¹ Between 0.15 and 1.5 eV, our cross section is about 1.5 to 2 times larger than the experimental data of Chantry. In this energy region, the semi-empirical calculations manage to reproduce the experimental values closely by appropriate adjustment of

Fig. 8-74. Comparison of theoretical and experimental dissociative-attachment coefficients for F_2 .

Fig. 8-75. Calculated elastic and vibrational-excitation cross sections for F_2 .

the resonance parameters, based on either the measured attachment rate⁷⁷ or a part of the cross section.⁷⁸ At electron energies below 0.15 eV, the theoretical and experimental cross sections differ significantly.

In interpreting the region below 0.15 eV, researchers should note that all the calculations (including the present one) treated only the resonant, $^2\Sigma_u$, contribution to dissociative attachment, and that this partial cross section must behave as $E^{1/2}$ at threshold. Both the present calculations and the results of Bardsley et al.⁷⁸ show this behavior and agree well with each other near threshold. The cross section calculated by Hall⁷⁷ is quite similar to the experimental data below 0.15 eV, but it does not satisfy the $E^{1/2}$ threshold law required for the $^2\Sigma_u^+$ channel. We also have calculated the dissociative-attachment cross section, using our present ab initio resonance parameters in the boomerang model^{79,80}, we obtained results with the same energy-dependence as Hall's data.

Based on the observations in the previous paragraphs, we conclude that the apparent agreement near threshold between the boomerang model and the experiment is an artifact due to a breakdown of the assumptions underlying this model at very low electron energies.^{81,82} We also believe that no critical comparison between theory and experiment is possible in the region from 0 to 0.15 eV until the nonresonant $^2\Sigma_g^+$ contribution to dissociative attachment is calculated. The $^2\Sigma_g^+$ partial cross section diverges as $E^{-1/2}$ at threshold, but this channel is coupled to $F_2(^2\Sigma_u^+)$ by only small, non-Born-Oppenheimer terms in the Hamiltonian.

Figure 8-74 shows the attachment coefficients calculated from our cross sections with a Boltzmann code, for two different gas mixtures (N_2 and Ar); the figure also includes all of the available experimental data.^{76,83-86} Separate data shown for different buffer gases are color-coded: red data points are used for an Ar buffer, green points for N_2 , and black points for a buffer of He. The present rate coefficients are in excellent agreement with the data of Schneider and Brau,⁸⁶ but they are approximately a factor of 2 higher than the values derived from Chantry's absolute cross sections.⁷⁶

Figure 8-75 shows both the elastic cross section and vibrational-excitation cross sections for the first four excited states of F_2 . The shapes of the curves are all similar, rising to their peak values at an incident electron energy approximately half an electron volt above threshold. The peak cross-sectional values decrease monotonically with increasing vibrational excitation. There are no experimental values available for comparison with our results at present.

All of the calculated cross sections are now being included in our RGH kinetics codes, to model the electron-energy distribution and F^- concentration in rare-gas fluoride lasers.

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Advanced Concepts

Free-Electron Lasers

We have conducted a preliminary analysis which indicates that a free-electron laser (FEL) amplifier might achieve efficiencies in excess of 10% at an installed cost (including the building) of approximately \$110 million (FY-80 dollars) per optical megajoule. In contrast to other lasers, the FEL offers the possibility of significant energy recovery, with a resultant increase in system efficiency, because much of the unused energy is in the form of a high-quality electron beam. Furthermore, the FEL generates output-laser fluences high enough (500 kJ/cm^2) to allow the laser radiation to be directed onto the fuel pellet without the need for any final-focusing optics, which are a source of potentially significant losses in other laser drivers. Direct coupling of the target with the intense, narrow beams generated by an FEL also permits a reduction, by four orders of magnitude, in the solid angle subtended by the laser output aperture at the pellet. This reduction effectively eliminates many of the difficulties associated with protecting optical surfaces from x rays, neutrons, and assorted other reactor-chamber debris.

A free-electron laser converts the kinetic energy of a high-quality relativistic electron beam into coherent radiation. An FEL

amplifier consists of an electron beam, an input laser beam, and a periodic magnetic field structure, usually called the "wiggler," which is illustrated in Fig. 8-76. The x-directed magnetic field gives the electrons a velocity component parallel to the laser field so that the laser radiation field and the electrons can exchange energy. The FELs that have been successfully operated⁸⁷⁻⁹⁰ use a constant-period wiggler and are, therefore, only capable of extracting a small percentage of the electron-beam energy as laser energy. If we build a wiggler with a tapered magnetic field, however, we should be able to extract a large percentage (about 30 to 40%) of the electron energy as laser radiation.^{91,92}

The large extraction efficiency is accomplished by trapping and bunching the electrons in moving potential wells, known as buckets, and then gradually slowing down the buckets to extract energy from the entire electron bunch. Greater than 80% of the electrons can be trapped in bunches, and 50% deceleration of each bunch is possible before detrapping begins. This process of bunching and decelerating is identical to the process responsible for trapping electrons in accelerating buckets in a linear accelerator. The size of the potential well, and the resultant intensity of the input laser required to create the buckets, is directly related to the quality (energy spread and emittance) of the input electron beam. Higher-quality beams are easier to trap and decelerate than low-quality beams.

Figure 8-77 is a rendering of how one half of a 3-MJ, 250-nm FEL fusion driver might be configured. The laser amplifier consists of 15 wigglers, each 120 m long and

Fig. 8-76. Energy is transferred in the FEL by giving the electrons a velocity component parallel to the laser field.

Advanced Concepts

producing an output of 100 kJ. We arranged the FEL amplifiers in a cone-shaped distribution around a pellet injector. The electron source consists of a 1.1-GeV betatron and a 50- to 100-MeV linear induction accelerator (linac) that serves as the injector for the betatron. An alternative source is a 1.1-GeV linac. The long-output (270 ns) electron-beam pulse from the betatron is transported to the electron-beam switch ring, where fast (3-ns rise time) magnetic switches divide the pulse into 15 flat-topped electron-beam pulses, each 15 ns long. These pulses are then magnetically guided into the wigglers, and the output radiation from the FEL amplifier proceeds directly to the pellet.

The only optical components in the system are in the laser preamplifiers and in the injector telescopes, which are over 130 m from the reactor vessel. Fast-acting valves

will be required to prevent reactor contaminants from traveling back through the wiggler to the preamplifier window. The principal characteristics of each beamline, as determined from our one-dimensional numerical simulation,⁹² are listed in Table 8-11, which also gives the specifications of the required electron source.

The projected FEL system consists of the following components:

- FEL amplifier-wiggler.
- Electron accelerator and electron-beam switch ring.
- Laser oscillator and input optics.

In the following subsections, we give a brief synopsis of elements of the projected system and then provide a summary of driver costs and efficiencies.

FEL Amplifier (Wiggler). Thus far we have based our analysis of the FEL amplifier on a one-dimensional (no transverse field

Fig. 8-77. FEL reactor architecture (one half of a 3-MJ driver).

Fig. 8-78. Laser intensity growth, electron-energy loss, wiggler-field strength, and integrated optical phase in a 250-nm tapered-wiggler FEL.

variations) numerical model. This model, which follows 500 electrons through a given wiggler, allows us to estimate the basic parameters of the FEL amplifier. Currently, the model indicates that 30 to 40% of the electron-beam energy can be transferred to the laser beam.

Figures 8-78 and 8-79 illustrate the predicted operation of the 250-nm FEL described in Table 8-11. Figure 8-78 is a composite graph showing

- Laser intensity (black line) and electron-beam energy (red line) as a function of distance along the wiggler, with the latter curve showing that the electrons lose about one-half of their initial energy to the photon field. On the average, 80% of the electrons are trapped, thus yielding a 40% extraction efficiency.
- The required wiggler magnetic field strength (green line).
- The integrated optical phase change (equivalent to the B-integral used to describe most high-power fusion lasers) expected for this FEL design (blue line). The relatively large value of ϕ raises critical questions about the laser-beam quality, which must be closely studied.

Figure 8-79 illustrates the cumulative effects of the combined wiggler and laser fields on a representative bunch of 500 electrons that is one optical wavelength (λ_s) long. Note that about 80% of the electrons are trapped within a bucket (solid line) and are bunched so that they occupy only a portion of an optical period; these electrons have been decelerated to $\sim \gamma = 1700$. The remainder of the electrons are located uniformly and are not decelerated.

The FEL performance predicted in these figures is encouraging, but several major physics issues remain to be resolved, including

- Three-dimensional phenomena, including laser-beam quality and propagation.
- Generation of Raman sidebands with resultant loss of phase stability.
- Propagation of electron beams.

The effects of several three-dimensional phenomena, such as variations in a transverse magnetic field, can be estimated by defining an equivalent one-dimensional energy spread, which can then be incorporated into our present FEL model. Unfortunately, not all three-dimensional problems can be treated this simply. Most important among these more difficult

problems is the effective index of refraction resulting from electron motion. This index is proportional to the ratio of the beam-current density to the laser electric field (J_B/E_S), so that an electron-beam current-density profile with a peak in the center may lead to self-focusing. This self-focusing, of course, leads to an increased laser field in the center of the electron beam; the laser field, in turn, reduces the self-focusing. Although our analysis indicates that this feedback mechanism stabilizes the amplifier against small-scale self-focusing, whole-beam self-focusing has not yet been analyzed in any detail. We have made very crude calculations that provide an estimate of 100 μ rad for the whole-beam self-focusing angle. This estimate is to be contrasted with a 50- μ rad diffraction spreading.

The approximately equal magnitude of the self-focusing and diffraction effects leads to a propagation problem that is difficult to solve analytically. A judicious choice of the beam's electron-current density profile and input-laser profile, however, could allow balancing of the two effects. This balancing

Fig. 8-79. Electron energy vs relative position 84 m down a 250-nm tapered-wiggler FEL.

Table 8-11. Characteristics of a 250-nm FEL fusion driver.

Amplifier	
Length	120 m
Wiggler period	15 cm
Wiggler field (variable)	2.5 to 5 kG
I_{in}	10 GW/cm ²
I_{out}	41 TW/cm ²
E_{out}	100 kJ (15 ns, 0.5-cm diam)
Accelerator	
Voltage	1.1 GeV
Current	20 kA
Energy spread	0.5%
Emittance (normalized)	0.1 rad-cm
Pulse length	0.5 to 0.6 μ s

Advanced Concepts

would be extremely beneficial, as it would permit greater amplifier lengths and, consequently, higher laser-beam extraction efficiencies. In addition, the laser beam may acquire some focusing as it exits the amplifier, thereby simplifying laser-beam transport to the target. A quantitative analysis of the propagation problem is critical to a realistic evaluation of the FEL as a fusion driver; such an analysis will not only indicate whether the beam can be directed onto the pellet, it will also determine optimum fill factors and, hence, the overall efficiency. We have therefore begun two-dimensional numerical modeling of the laser-beam propagation through an FEL amplifier.

Raman instability is another factor for our consideration. As the electrons in an FEL are decelerated, they acquire a low-frequency velocity modulation. This modulation may impose sidebands on the main laser frequency. As the sidebands grow, they may disrupt the potential wells that are actually decelerating the electrons.

If stable-phase electron orbits are to be achieved in the FEL wiggler, the electron beam must be very "cold," and have a low normalized emittance (0.1 rad-cm). Without the use of focusing magnetic fields, however, even this extremely bright beam will not propagate the length of the wiggler. Although we can solve this problem in principle, bunched high-current electron beams are prone to many instabilities that we must analyze.

The wiggler magnet itself is within current technological capabilities. This magnet consists of a vacuum pipe for the electron

beam, surrounded by rare-earth-cobalt permanent magnets capable of producing a maximum field of 5 to 10 kG at the pole faces. We estimate the cost for 30 wigglers to be about \$10 million for the vacuum system and \$15 million for the magnetic materials.

Electron Accelerators. An FEL laser-fusion driver requires the simultaneous delivery of a 1.1-GeV, 20-kA, 15-ns electron-beam pulse to each of the 30 wiggler magnets that form the free-electron laser. It is more cost-effective to accelerate a longer beam pulse and divide it into shorter pulses with electron-beam magnetic switches and delay lines than to build 30 separate accelerators. In comparing different accelerators, we assumed that the parameters of these individual pulses are constant; we made adjustments in accelerator characteristics (such as pulse length) to account for differences in accelerator efficiencies. We imposed further requirements on the electron beam, including an energy spread of $\Delta E/E \approx 0.5\%$ and a normalized emittance (brightness) of 0.1 rad-cm (or 45 μ rad-cm at 1.1 GeV). Several different types of electron sources are discussed below.

Although only the electrostatic accelerator (Van de Graaff generator) can be categorically ruled out (because of high-voltage insulation requirements), the linac clearly involves the least extensions beyond present capabilities to satisfy the stated requirements. The state of the art is represented by the advanced test accelerator (ATA), currently under construction at LLNL, that will produce a 50-MeV, 10-kA, 50-ns electron-beam pulse. The required increase in electron energy and pulse length (one 540-ns pulse for a linac vs one 300-ns pulse from each of two betatrons, with estimated electron-beam switching losses) can be obtained by a relatively straightforward increase in the amount of magnetic material, pulsed power, and accelerator length. The increase in current is modest, but the required improvement in emittance at 20 kA will require further development. Our projected cost of \$235 million for an FEL linac is dominated by pulsed-power and magnetic-material costs, with only a modest potential for improvement. We give the cost breakdown of a 1.1-GeV linac in Table 8-12.

The betatron is potentially an attractive alternative to a linac because of lower costs. It is conceivable that we could obtain

Table 8-12. Nominal 3-MJ FEL fusion driver costs (millions of FY-80 dollars).

	Linac	Betatron
Accelerator		
Magnetic material	60	100
Induction modules	20	—
Pulsed power	110	20
Accelerator beam line	30	10
Injectors	—	20
Wigglers (30)	15	15
	235	165
Magnets	15	15
Pumps	10	10
Oscillators	5	5
Building	75	75
Total	\$340	\$270
Cost per MJ	\$110	\$90

betatron acceleration times as long as 500 μ s, which would allow the use of more benign and less expensive pulsed-power sources (possibly a compensated alternator). In exchange for this simplification, the betatron introduces three major problems:

- Injection of low-energy electrons into the ring. This problem can be addressed by injecting at sufficiently high energy (50 to 100 MeV) from a linac with a pulse length equal to the transit time around the betatron ring.
- Extraction of the high-energy beam from the betatron ring. The beam could be extracted with a magnetic switch, but the strength and volume of magnetic field change required would limit the rise time (~ 30 ns) of the switch and lead to a loss of beam electrons.
- Maintenance of a stable beam in the betatron ring during acceleration. Instability of the circulating beams is the most critical problem, with the resistive wall instability being the most serious aspect of this problem.

The cost of the betatron is dominated by the cost of the magnetic material used for the induction core. Our cost estimate of \$165 million for the electron-beam system includes \$90 million for a core composed of 0.01-in. sheets of Si-Fe. If longer acceleration times or thinner laminations were used to reduce eddy currents, mild iron could be used with a savings of $\sim$ \$60 million. We show the cost breakdown for this accelerator in Table 8-12.

Radio-frequency (rf) sources operate at relatively high impedance and are quasi-continuous-wave. If the source is to be used efficiently, the accelerator should match these characteristics. A circular accelerator (synchrotron) would match the quasi-cw nature, but not the high-impedance nature, of the rf source. In addition, the need to maintain phase matching between the beam bunches and the rf field is a severe constraint in the face of potential instabilities; in fact, this need introduces additional instability mechanisms. A linear rf accelerator would require the use of an electron ring to store and accumulate the low-current, long-pulse output of the accelerator. Such a ring shares the stability problems of betatrons and synchrotrons. In addition, beam stacking or accumulation may not be compatible with the required beam emittance.

Thus, for fusion applications of the FEL that require short, intense electron-beam pulses, we feel that the most viable electron source with the least technological risk is a linear-induction accelerator with a fast magnetic switch. We take this system to be the baseline electron source. The systems advantages of the betatron, however, are great enough that further research on instability control is warranted.

To produce 30 separate 15-ns electron-beam pulses from a single accelerator pulse, it is necessary to chop the beam into individual segments and delay all except the last segment so that all segments can be injected into their respective wiggler magnets simultaneously. The delays can be provided by the system that distributes the beam to the wigglers. If the inputs to the wigglers are arranged in a circle such that the electron transit time between wiggler inputs is equal to the sum of the pulse length (15 ns) and switch rise time (~ 3 ns), then the electron beam can be steered around this circle until the ring is filled. At that point, all beam switches (one per wiggler) can be thrown simultaneously, to divert individual beams into bending magnets that direct them into the wigglers. These electron-beam switches can be built with much faster rise times than the betatron extractor switch, because the electron-beam switches can have much smaller transverse dimensions, allowing smaller volumes.

Laser Front End and Input Optics. As indicated in Table 8-11, the FEL amplifiers require a collimated input beam of 10 GW/cm² or 150 J/cm² across a 0.5-cm-diam aperture. By assuming optical-damage thresholds of 10 J/cm², we imply that any input windows or mirrors must be 300 m from the amplifier input. A grazing-incidence reflecting system can be used, however, to increase the effective surface area to a point where the optical fluence is less than the damage threshold. At a damage threshold of 2 J/cm², the grazing angle must be less than 0.6°; at 10 J/cm², the angle can be about 3°. The corresponding lengths of the reflecting surface are 150 cm and 30 cm, respectively. The telescope can be slightly defocused if it becomes necessary to provide the FEL amplifier with a divergent input beam to compensate for self-focusing.

System Summary—Cost and Efficiencies. Table 8-12 summarizes the costs of a 3-MJ

induction linac-based FEL amplifier. These costs are estimated in FY-80 dollars and do not include any contingency funds or Engineering, Design, Inspection and Administration (EDIA) funds. We based the building cost estimate both on the size of the concrete structure required to house the wigglers and on the cost of comparable laboratory structures. Cost estimates derived in a similar manner for a betatron-driven FEL amplifier are also shown in Table 8-12.

Table 8-13 gives the efficiencies of each stage of the linac and betatron amplifiers.

- "Pulsed power to B-field" specifies the transfer efficiency from the wall plug to the accelerating magnetic field.
- "Transfer to electron beam" indicates how much of the field energy can be coupled into the electron beam. The betatron is more efficient in these first two categories because it has the advantage of a longer accelerating pulse.
- "Stacking" reflects the energy loss that results from the finite rise time of the magnetic switches. The betatron requires one more switch than the linac to get the beam out of the accelerator; therefore, the betatron is less efficient in this category.
- "Fill factor" is an optimistic estimate that might be strongly affected by the results of the two-dimensional propagation studies.
- "FEL extraction" is an efficiency calculated from the one-dimensional amplifier simulation.

Our estimated total system efficiencies for the linac and betatron systems are shown on the line labeled "Total." On the next line down, "Total with thermal-energy recovery of output electron beam," we show the calculated slight increase in overall efficiency that might be achieved by thermally converting the spent electron beam to

electricity (assuming a 30% efficient conversion). If direct energy conversion at 0.5 to 1.0 GeV could be done at low cost, energy recovery would be higher and overall system efficiency would increase.

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Energy-Storage Solid State Lasers—V:MgF₂

Vanadium-doped magnesium fluoride (V:MgF₂) is an attractive laser-gain medium for use in high-average-power laser drivers for fusion applications.⁹³ In last year's annual report,⁹⁴ we summarized the results of a calculational assessment of a V:MgF₂ power amplifier, based on the limited spectroscopic data available in the literature.^{95,96} We concluded that efficient ($\geq 5\%$) MJ-class V:MgF₂ laser systems could be designed and built, provided that large (30 by 60 by 3 cm) high-optical-quality slabs of MgF₂ can be grown when doped to a level of ≥ 2 wt% VF₂. Accordingly, we initiated efforts in 1980 to accurately determine critical spectroscopic and laser parameters, and to grow vanadium-doped MgF₂ crystals to determine the maximum V-ion doping levels consistent with requirements of optical quality. These efforts were carried out by researchers in the Quantum Electronics and Electronic Materials Groups, Solid State Division, of Lincoln Laboratory under SAN-L Contract 806-013. We present the highlights of our research results in this article; for more detailed information, see Ref. 97.

The growth of V:MgF₂ crystals has been carried out by means of the vertical gradient-freeze technique developed at Lincoln Laboratory for growing Ni:MgF₂ crystals.⁹⁸ A charge consisting of MgF₂ and VF₂ is placed in a cylindrical graphite crucible (see Fig. 8-80) that is tapered near the bottom, ending in an open capillary about 1 mm in diameter and 5 cm long. The crucible is closed at the top with a tapered, tightly-

Table 8-13. Projected system efficiencies for a 250-nm FEL fusion driver.

	Linac	Betatron
Pulsed power to B-field	0.60	0.80 ^a
Transfer to electron beam	0.85	0.90
Stacking	0.85	0.75
Fill factor	0.80	0.80
FEL extraction	0.38	0.38
Total	0.13	0.16
Total with thermal-energy recovery of output electron beam	0.14	0.18

^aProbably requires some recovery of the energy stored in the magnetic field.

fitting graphite plug and mounted in a furnace that has a water-cooled, stainless steel jacket. The heat source (not shown) is a cylindrical tungsten-mesh resistance element surrounded by multilayer molybdenum heat shields. The heat shields are arranged so that the temperature of the crucible decreases from the top to the bottom, which is cooled by close contact with the water-cooled base of the furnace. Figure 8-80 shows a typical thermal gradient used in our growth studies. The crucible temperature is monitored at a point somewhat above the top of the capillary by means of a thermocouple inserted into a well in the crucible wall.

To begin a crystal growth run, the furnace chamber is evacuated, the crucible is heated to about 1275 K, the chamber is back-filled with argon at atmospheric pressure, and heating is continued. When the charge begins to melt, liquid MgF_2 runs down into the capillary, where it freezes to form a plug that fills the capillary. The temperature is

- Increased to about 1875 K to ensure dissolution of the VF_2 .
- Decreased rapidly to about 1650 K (100°C above the melting point of MgF_2).
- Decreased at 2.5 to 5°C/hr so that the molten charge is gradually solidified from the bottom up.

Although the plug formed in the bottom of the capillary is highly polycrystalline, the crystal generally becomes single before the solid-melt interface reaches the top of the capillary. After the charge is completely solidified, the furnace is cooled to room temperature over a period of 3 or more days. The entire cooling cycle is programmed by means of a microprocessor that sets the control point for a proportioning temperature control.

We used the procedure described above to grow 13 crystals of V:MgF_2 for use in spectroscopic and laser experiments. We also used the procedure to generate information on the growth parameters for producing high-optical-quality moderately doped V:MgF_2 crystals.

To investigate the important issue of the solubility of VF_2 in MgF_2 , we grew two crystals of 1.25 cm diameter from charges containing 5 and 10 wt% VF_2 . Figure 8-81 shows a picture of the resulting 5-wt% V:MgF_2 crystal. X-ray fluorescence analysis

Fig. 8-80. Vertical-gradient-freeze crystal-growth crucible.

Fig. 8-81. V:MgF_2 crystal grown from a 5-wt% VF_2 charge.

(described below) and lattice constant data for these crystals show that the solubility of VF_2 in MgF_2 is at least 3.3 and 9.2 wt%, respectively, for the two crystals we grew. The lattice constants of the doped crystals varied linearly between values for pure MgF_2 and VF_2 crystals, in a manner proportional to the VF_2 concentration.

This is a significant and very positive result. The V^{2+} ions appear to be entering the MgF_2 lattice on substitutional Mg^{2+} sites in a statistically random fashion, and for ion densities up to $\sim 3 \times 10^{21}$ ions/cm³ (~ 10 wt% VF_2). This doping level is considerably larger than that (>2 wt% VF_2) which we believe will be required for efficient pumping of large fusion lasers. Of course for this application we will need to be able to grow VF_2 crystals of >2 wt% with useful optical quality.

The other crystals we grew have diameters of either 2.5 or 3.75 cm, and were grown from charges containing 0.25, 0.5, or 1 wt% VF_2 . Figure 8-82 shows a picture of one of the largest crystals grown, with a 0.5-wt% V:MgF₂ melt. We used x-ray fluorescence to

measure the V^{2+} ion density in the first-to-freeze region (tapered end) of a number of these crystals. The ratio of the VF_2 content in this region to that in the charge ranges from 1.06 to 1.41. For most of the crystals, the rich violet color that is due to V^{2+} absorption decreases somewhat (in one case becoming almost undetectable) toward the last-to-freeze end, indicating a reduction in VF_2 content. This phenomenon is observable in the crystal shown in Fig. 8-82.

We obtained high-optical-quality laser rods from two of the crystals, grown from charges containing 0.5 wt% VF_2 . Overall, however, the yield of laser-quality material has been limited; most of the crystals have cracked after we removed them from the furnace, indicating that they were strained while being cooled to room temperature. Our attempts to reduce the strain by reducing the cooling rate have not yet been successful in preventing cracking, even though in several cases we extended the cooling cycle to more than 10 days. We are currently investigating two approaches intended to solve the cracking problem:

- Reducing the temperature difference along the length of the crucible.
- Using a seed oriented so that the long axis of the crystal will be perpendicular to the c-axis, an orientation that will minimize the stress on the crystal produced by the walls of the crucible.

We are continuing crystal-growth studies in 1981, with emphasis on establishing optimum growth parameters for high-optical-quality MgF_2 crystals containing ~ 2 wt% VF_2 .

The absolute cross sections of the visible absorption bands of V^{2+} ions in MgF_2 , derived from two sets of spectral data in the literature,^{95,96} differ by a factor of 2. Since we need this cross-section value to accurately predict the excitation efficiency of V:MgF₂ as a function of the V^{2+} ion doping level, we have developed an accurate method of determining V^{2+} ion concentration in MgF_2 crystals. Using concentration data together with optical absorption spectra, we have been able to determine the absolute absorption cross sections for V:MgF₂.

As a convenient technique to determine the VF_2 content of V:MgF₂ crystals, we use a standard x-ray fluorescence spectrometer to measure the intensity of vanadium x-ray

Fig. 8-82. Large V:MgF₂ crystal grown from a 0.5-wt% VF_2 charge.

fluorescence excited by x rays from a molybdenum (Mo) tube. With this technique we can measure compositions as low as 0.2 mole % VF_2 with an accuracy of $\pm 5\%$. The only specimen preparation required is cutting and polishing the crystal to produce a flat surface for exposure to the Mo x-ray beam. The method gives the average VF_2 content within a region extending several micrometers below the surface, a depth determined by the x-ray absorption properties of MgF_2 . We use a minimum sample thickness of two or three times this depth to ensure accurate analysis.

We prepared standards for the analysis by mixing weighed amounts of MgF_2 and V_2O_5 powders. (We used V_2O_5 in order to conserve our limited supply of VF_2 powder for crystal growth studies.) We established the validity of employing powder standards by conducting a preliminary experiment in which we cut two adjacent pieces from a V:MgF_2 crystal. We polished one piece, and ground the other to pass through 100-mesh screen. We found the difference in the intensities of V x-ray radiation from the solid and powdered samples to be less than the scatter of the data.

In performing the analysis, we supported each sample or standard on a mylar film

over a circular aperture in a metal disc. We then irradiated the samples with x rays from a Mo tube operated at 40 kV and 40 mA. We analyzed the fluorescence spectrum with an LiF crystal, and measured the intensity of the vanadium K_α line with a scintillation counter and pulse height analyzer. We corrected the V x-ray fluorescence signal by subtracting the signal from a blank obtained from undoped MgF_2 powder. For a series of standards ranging in composition from 0.2 to 7.0 wt% V, corresponding from 0.24 to 8.5 mole % VF_2 in V:MgF_2 , we found the corrected signals to be proportional to the V concentration, with a scatter of $\pm 5.0\%$ of the measured values.

To establish the absolute cross sections of V:MgF_2 absorption transitions, we have recorded the polarized spectral absorbance of an oriented V:MgF_2 single crystal containing $1.26 \times 10^{20} \text{ V}^{2+} \text{ ion/cm}^3$ (determined using the x-ray fluorescence standards just described). In Fig. 8-83 we show the polarized absorbances for this crystal as a function of wavelength, with the crystal held at a temperature of 77 K. We determined the peak cross section of the σ -polarized, 900-nm $^4\text{A}_2 \rightarrow ^4\text{T}_2$ absorption band to be $1.17 (\pm 0.05) \times 10^{-20} \text{ cm}^2$. Relating the measured peak absorbance of the σ -polarized 550-nm

Fig. 8-83. Polarized absorption spectra of V:MgF_2 for $T = 77 \text{ K}$.

Fig. 8-84. Axial absorption spectrum of V:MgF₂ for T = 300 K; xenon lamp emission spectrum.

Fig. 8-85. Xenon lamp excitation efficiency of V:MgF₂ as a function of optical depth.

band at 300 K to our V-ion concentration standards, we find the σ -polarized peak cross section of the 550-nm band to be $1.02 (\pm 0.05) \times 10^{-20} \text{ cm}^2$ at a temperature of 300 K.

In Fig. 8-84 we show the room-temperature axial (σ) spectrum of a V:MgF₂ crystal in relation to the spectral emission

characteristic of a large-bore xenon flash-lamp.⁹⁹ From this figure, we see that the σ -polarized peak cross section of the 900-nm band at 300 K is $1.1 (\pm 0.05) \times 10^{-20} \text{ cm}^2$. Having established this cross-section value we can now conveniently determine the line-averaged V-ion concentration in a spectroscopic sample, simply by measuring the value of the peak axial (σ) absorption coefficient of the 900-nm band at room temperature and dividing it by the product of the sample thickness and the cross-section value just given.

Figure 8-85 shows the calculated excitation efficiency of the 4T_2 metastable level of V:MgF₂ at 300 K for a xenon flashlamp, as a function of the "optical depth" of the V:MgF₂ medium. We define the optical depth as the product of the V-ion concentration and the thickness of the material through which the lamp radiation passes; we express the optical depth in Fig. 8-85 in units of $10^{20} \text{ ions/cm}^2$ and in units of wt% VF₂-cm. The efficiency plot of Fig. 8-85 replaces and corrects a similar plot given in last year's annual report⁹⁴ whose optical depth scale is now known to be 20% too compressed.

To design large V:MgF₂ lasers, we must know the dependence of the $^4T_2 \rightarrow ^4A_2$

fluorescence lifetime on temperature and V-ion concentration. The only previously published data on this lifetime⁹⁶ cites a value of 2.3 ms at 77 K. We have measured the fluorescence lifetime of V:MgF₂ as a function of both temperature and concentration. For temperature-dependent measurements, we placed a polished section of a crystal (#3 in Table 8-14) containing approximately 0.5 wt% VF₂ on the cold finger of a helium-gas-cooled variable-temperature dewar flask. We excited fluorescence using a nitrogen-laser-pumped dye laser operating at 590 nm and we detected the fluorescence radiation after passing it through a spectrometer using a PbS detector. The detected signal was analyzed by a Nicolet 1074 signal-averaging system; the decay curves were plotted on graph paper. At crystal temperatures of 240 K and above, where the response of this detection system was too slow to accurately determine the decay time, we replaced the PbS detector with a silicon photodiode possessing an enhanced long-wavelength response. We also added a Biomation 610B transient digitizer to the Nicolet averager.

The decay time, τ , vs temperature for the sample crystal is shown in Fig. 8-86. Here τ is defined as the time for a $1/e$ reduction in signal, starting at a point where the signal had decayed to $\sim 80\%$ of the peak value. The value of τ was 2.44 ms at 77 K, in good agreement with the previous data.⁹⁶ The decay constant measured to the $1/e^2$ and to the $1/e^3$ points was 2.25 ms at 77 K, indicating only a slight deviation from a true exponential decay at this temperature.

Figure 8-86 shows that the fluorescence lifetime decreases only slowly between very low temperatures and 200 K, and then decreases rapidly at higher temperatures. We presumed the physical mechanism which causes the sharp reduction in τ above 200 K to be thermal stimulation of excitation in the 4T_2 state over to very high vibrational levels

of the 4A_2 ground state, followed by rapid nonradiative decay to low vibrational levels of the 4A_2 state. The simplest mathematical description of this mechanism is given by the Mott activation-energy formula for the nonradiative rate, R_N , such that

$$R_N = A \exp \left(- \frac{\Delta E}{\kappa T} \right) \quad (34)$$

where A is a transition rate, ΔE is the activation energy, κ is Boltzmann's constant, and T is the temperature. The observed lifetime, τ , is then given by the expression

$$\frac{1}{\tau} = \frac{1}{\tau_R} + A \exp \left(- \frac{\Delta E}{\kappa T} \right) \quad (35)$$

where τ_R is the radiative lifetime.

It is possible to determine ΔE independently from the constant A , by making measurements in the temperature region where the nonradiative process dominates. If at temperatures T_1 and T_2 we observe decay times τ_1 and τ_2 , respectively, then

Fig. 8-86. Fluorescence lifetime of V:MgF₂ as a function of temperature.

Crystal	Concentration (wt% VF ₂)	Temperature (K)	Time τ (ms)
"Optovac"	0.1	20	2.43
		77	2.41
		240	0.535
#3	0.5	20	2.48
		77	2.44
		240	0.551
#9	3.3	77	2.2
#10	9.1	77	1.8

Table 8-14. V:MgF₂ fluorescence lifetimes for various temperatures and V-ion concentrations.

$$\Delta E = \frac{\kappa \ell n \left(\frac{\tau_1}{\tau_2} \right) T_1 T_2}{T_2 - T_1} \quad (36)$$

In the case of V:MgF₂, we calculated ΔE for increments of 10 K between 250 and 280 K, and found an average value of 0.394 eV or 3180 cm⁻¹, with a 4% variation. We obtained a reasonably good fit to the data by then choosing a value for A of $3.3 \times 10^{11} \text{ s}^{-1}$, as is shown in Fig. 8-86. This value of A is somewhat lower than typical values ($\sim 10^{13} \text{ s}^{-1}$) found for phonon activation rates in crystals.

Struck and Fonger¹⁰⁰ have discussed more sophisticated treatments of temperature quenching, which rely on standard, single-configuration-coordinate models. Sturge,¹⁰¹ however, has pointed out the large and difficult-to-calculate effects of anharmonicity in such theories. These effects would undoubtedly be important in the case of V:MgF₂, since a large number of phonons are involved in the nonradiative process. The apparent complexity of properly modeling the nonradiative process has discouraged further attempts at analysis using the presently available data.

We determined values of τ for several other crystals with differing vanadium concentrations. One sample contained a trace amount ($\leq 0.1 \text{ wt\% VF}_2$) of vanadium ions, sufficient to produce a measurable fluorescence signal. We also measured fluorescence lifetimes for sections of crystals (#9 and #10 in Table 8-14) containing, respectively, 3.3 and 9.1 wt% VF₂. We summarize the results of these measurements in Table 8-14. The low-temperature fluorescence lifetime of 2.48 ms that we measured for the low-concentration sample is taken as the radiative lifetime of the ⁴T₂ level of V:MgF₂. Compared to rare-earth-

doped crystals such as Nd:YAG, we note that concentration quenching of fluorescence in V:MgF₂ is much more gradual, as indicated by the data in Table 8-14. If we assume that quenching is due to ion-ion dipolar (or higher-order) interactions, this result is not unexpected because of the substantial Stokes shift between emission and absorption, and because of the smaller oscillator strength of the V²⁺ transition compared to typical rare-earth-ion forced-dipole transitions. It is apparent that concentration quenching is not a major concern in doping of crystals up to at least 9 wt% VF₂. This is also a significant and positive result since it allows us a significant degree of flexibility in designing large V:MgF₂ power amplifiers.

To design such amplifiers we need polarized fluorescence spectra and fluorescence emission cross sections. Previously published data on V:MgF₂ fluorescence consists of one spectrum, of unspecified polarization, at 77 K. It was stated that there was "little if any difference between the σ , π and axial spectra."⁹⁶ In order to determine the validity of existing data, and to obtain results over a wider temperature range, we have measured the polarized fluorescence spectra of V:MgF₂.

We employed the same crystal and dewar flask for these experiments that we used to measure lifetime data. We excited fluorescence by pumping the ⁴A₂ → ⁴T₁ transition with a 50-mW-output, cw He-Ne laser. A portion of the laser output was sampled and the detected fluorescence signal was electronically normalized to this laser output signal, for improved accuracy and reproducibility. We passed the fluorescence from the x-ray-oriented crystal ($\sim 0.5 \text{ wt\% VF}_2$) through a Polaroid HR film polarizer and a Heath EV-700 grating spectrometer, and then detected it with a PbS photoconductor. We synchronized the stepping motors driving the spectrometer scan and chart recorder, to obtain an accurate wavelength axis. We determined the polarized wavelength response of the detection system by placing a tungsten standard lamp at the location normally occupied by the crystal. Recorded data was digitized and then corrected for the system response by a HP9830 Computer. In addition, we measured the change in absorption of the He-Ne pump laser as a function of temperature, and corrected the

Fig. 8-87. Polarized ⁴T₂ → ⁴A₂ fluorescence intensities of 0.5-wt% V:MgF₂ at T = 80 K.

fluorescence intensities at different temperatures to take the change into account.

The fluorescence spectra recorded for a crystal temperature of 80 K are shown in Fig. 8-87. We note that the fluorescence is significantly polarized, the peak of the σ -spectrum being about 2.5 times that of the π -spectrum. Absolute, polarized-fluorescence cross sections are needed to calculate the specific gain coefficient of excited V:MgF₂ material. To establish the absolute peak cross sections, we first establish the cross section of the zero-phonon line (the sharp spectral feature near 1 μ m in Fig. 8-87) which appears in absorption and emission. The cross section in absorption is determined from accurate absorbance and concentration measurements. The peak-emission cross sections are then scaled to the known zero-phonon-line cross section in emission.

Analysis of the data is still in process. Preliminary results indicate that the σ -polarized peak stimulated-emission cross section is $\sim 0.4 \times 10^{-20} \text{ cm}^2$ at $T = 80 \text{ K}$. This value is roughly one-half the value quoted in the literature.⁹⁶ If our completed spectral analysis sustains the lower cross-section value cited above, the V:MgF₂ material will be characterized by a calculated homogeneous saturation fluence $\Gamma_{\text{sat}} = hc/\lambda_p \sigma_p = 42 \text{ J/cm}^2$. This rather high value of Γ_{sat} can probably be accommodated in laser system design, due to the high damage threshold properties of crystalline MgF₂. Because of the relatively low stimulated-emission cross section and its significantly σ -polarized character, it will be necessary to orient the c -axis of V:MgF₂ Brewster amplifier disks to make it perpendicular to the electric vector of the optical extraction wave.

With the exception of the preliminary measurement of a lower stimulated-emission cross section, we have found V:MgF₂ to have very favorable properties for a fusion laser:

- Long storage time for temperatures up to 200 K.
- Low-concentration quenching up to 9 wt% VF₂.
- Good spectral coupling to xenon-flashlamp radiation.

In 1981 we will complete an accurate determination of the stimulated-emission cross section, and continue growth studies of V:MgF₂, with a goal of establishing the

optimum growth and annealing parameters to obtain good optical-quality material.

We have recently obtained cw laser action in V:MgF₂ pumped with an Argon-ion laser. We will study the oscillator and amplifier behavior of this device with the goal of confirming the measured spectroscopic parameters and achieving some insight into any excited-state absorption processes which might be present in the V:MgF₂ system. Using the spectroscopic data reported here, laser amplifier and system designs will be improved.

Authors: P. F. Moulton and R. E. Fahey (Lincoln Laboratory), and W. F. Krupke

Chemically Pumped Iodine Lasers

For many years, attempts have been made to devise a laser system for fusion applications based on the iodine-atom $^2P_{1/2} \rightarrow ^2P_{3/2}$ transition at 1.3 μ m. In early efforts the excited atom was produced by photodissociation of perfluoro-alkyl iodides: CF₃I, C₃F₇I, etc. These studies show that iodine lasers are of sufficiently high quality to be scaled to meet the requirements for fusion.¹⁰² The overall efficiency of using a flashlamp to dissociate the molecule, however, is not expected to exceed 1%.^{103,104} This unacceptably low efficiency seems to rule out the photodissociation-based laser system for fusion applications. Recently the Air Force Weapons Laboratory developed a chemically pumped cw oxygen-iodine laser which is projected to be substantially more efficient.¹⁰⁵ As part of the national ICF Advanced Lasers Program, efforts are currently underway to see if this chemically pumped iodine laser can be modified to produce high power in a short, pulsed mode.

In this article we review the basic mechanism of the O₂/I laser, and discuss the three methods which are currently being investigated for producing O₂(¹ Δ):

- Conventional chemical pumping.
- Sensitization.
- Electrical discharge.

We also put into perspective the key technical and cost issues which must be successfully addressed to render the oxygen-iodine laser a viable driver.

Advanced Concepts

We believe that the oxygen-iodine laser works via the following somewhat simplified mechanism¹⁰⁶:

and

Figure 8-88 shows a schematic of an $\text{O}_2(^1\Delta)$ chemically pumped iodine amplifier.¹⁰⁷ The nearly resonant excitation transfer in Eq. (39) is approximately gas kinetic, rapidly establishing an equilibrium in about 200 ns. The equilibrium constant for that reaction,

$$K_{\text{eq}}(T) = \frac{k_3}{k_{-3}} = \frac{[\text{O}_2][\text{I}^*]}{[\text{O}_2^*][\text{I}]} \quad , \quad (41)$$

is 2.9 at room temperature. Fractional concentrations of excited $\text{O}_2(^1\Delta)$ greater than 15% will produce an iodine inversion. In fact, a more rigorous model indicates that the necessary fraction must be greater than 25% (Ref. 105). I^* is quite resistant to quenching by most of the species present, but is very rapidly quenched by I_2 . The three-

body recombination to form I_2 and the feedback reaction [Eq. (40)] controls the length of time the energy can be stored in the medium. Bush¹⁰⁷ has predicted as possible storage times on the order of 3 ms at room temperature. Using plausible temperature-dependencies for the rate constants, he predicts storage times as high as 1 s at 100 K.

Since the equilibrium constant for Eq. (39) is small, much of the energy remains stored in the $\text{O}_2(^1\Delta)$. Several investigators suggest that additional passes of the laser pulse through the medium will extract more energy.^{106,107} Such multipass schemes require a delay time of ~ 200 ns between passes, to allow the reestablishment of the chemical equilibrium. The system described in Ref. 106 assumes a total O_2 pressure of 10 Torr (80% of which is in the $^1\Delta$ state), 100 Torr of Ar, and 0.3 Torr of I_2 . The stored energy initially is on the order of 45 J/ℓ; 20 passes extract 15 J/ℓ. The projected large signal gain is >100 , and $g_{\text{max}} = 0.027 \text{ cm}^{-1}$.

The chemical pumping scheme used to generate the $\text{O}_2(^1\Delta)$ for this laser utilizes the following reaction¹⁰⁵:

Chlorine gas is bubbled through a basic solution of hydrogen peroxide. The chlorine diffuses into the solution and is rapidly converted to $\text{O}_2(^1\Delta)$. The excited oxygen molecules diffuse back into bubbles and are released into an Ar buffer gas above the solution. Since the above reaction is exothermic by approximately 2 eV, the excess heat also causes some of the water to vaporize. It is not known how H_2O molecules affect the performance of the oxygen-iodine laser. In practice, cold traps remove most of the water from the Ar- O_2 mixture.¹⁰⁵ The $\text{O}_2(^1\Delta)$ is very stable with respect to quenching by walls so that it is possible to pipe the excited oxygen into the laser volume with very little loss of the excited-state population.

The following two-step process regenerates the reactants for Eq. (42). Electrolysis of the salt solution yields Cl_2 and H_2 gases; regeneration of the hydrogen peroxide uses the H_2 gas, via the reversible auto-oxidation of an anthraquinone¹⁰⁸:

Fig. 8-88. Schematic of $\text{O}_2(^1\Delta)$ chemically pumped iodine laser amplifier.

(43)

Since these regeneration reactions are well known commercial processes, we can obtain a reasonable estimate of their efficiencies. Fisk found that the overall regeneration efficiency is approximately 60% relative to the thermodynamic minimum energy requirement.¹⁰⁹ Thus since the minimum energy needed is ~ 3 eV, the efficiency for the generation of almost 1 eV of excitation in the $O_2(^1\Delta)$ is $\sim 20\%$. This figure assumes 100% conversion of the Cl_2 molecules to $O_2(^1\Delta)$. This gives an absolute upper limit on the efficiency of the laser of 20%, if all the excitation energy can be extracted.

As stated above, the $Ar-O_2(^1\Delta)$ is piped through at least one cold trap (to remove most of the water vapor) before it is pumped into the laser volume. Fisk reports measurements of the fraction of $^1\Delta$ downstream from the generator as high as 83% (Ref. 109). The I_2 molecules are injected into the flowing $Ar-O_2$ gas just before it enters the amplifier volume. For this chemical-pumping system we need long storage times, so that complete mixing is achieved and turbulence is minimal when the energy is extracted. The homogeneity of the medium is also important for maintaining high beam quality and efficient extraction with minimal optical loss during the multiple passes of the laser pulse through the system.

To achieve these goals, we need accurate determinations of possible storage times, to establish the required flow velocity and the optimum injector scheme. Existing modeling efforts have used some rate constants that are now known to be wrong. In particular, recent measurements by Aviles et al. show that the rate constant for Eq. (38) is an order of magnitude smaller than the value commonly used ($k_2 = 2 \times 10^{-11}$ cm³/mol-s) (Ref. 110). R. F. Heidner III at Aerospace Corporation reports that the value for k_4 may be in error also.¹⁰⁶ And, as mentioned above, we do not understand the role of H_2O vapor in the laser operation. Therefore, we cannot establish optimum operating conditions for this chemical-pumping scheme since there appears to be significant misunderstanding of the laser kinetics.

Using the available information, Fisk¹⁰⁶ projects an overall efficiency for this system of $\sim 6\%$. There are, however, many critical issues associated with the development of the chemically pumped iodine fusion laser:

- Chemical pumping and I^* storage times commensurate with short-pulse (1 to 10 Hz) laser operation.
- $O_2(^1\Delta)$ generator efficiency.
- Chemical mixing and laser medium homogeneity.
- Multipass extraction efficiency, optical loss, and laser beam quality after multipass extraction.
- Chemical reconstitution cost and efficiency.

Scalable and effective solutions to these physics and technology issues could make the chemically pumped iodine laser an important device for fusion applications.

An alternative source of $O_2(^1\Delta)$ for pumping the I-atom laser involves the use of sensitizers. Chemists use organic sensitizers to study oxygenation processes involving singlet oxygen.¹¹¹ Dye sensitizers such as methyl blue, fluorescein derivatives, porphyrins, heterocyclic organics, naphthalene, and many others absorb radiation (usually $\lesssim 3$ eV), then transfer the excitation to ground-state oxygen via the following mechanism:

and

The singlet ground state of the sensitizer is generally a broad-band absorber into several excited singlet states. The excited state rapidly decays, frequently creating significant populations in the lowest metastable triplet state. The quantum yields for producing metastables vary widely for the different sensitizers. Many are found to exceed 60%. If the excitation energy of this metastable is larger than the triplet-to-singlet energy difference in O_2 , the energy can be transferred via two simultaneous spin-forbidden transitions. If Eq. (46) above is very exothermic, the reverse reaction is very unlikely, possibly producing a substantial

fraction of $^1\Delta$. Several sensitizers generate populations of $^1\Delta$ in excess of 50%. Using perfluoronaphthalene excited by mercury or xenon flashlamps, Wasserman et al. generated 70% $^1\Delta$ (Ref. 111).

Using extraction efficiencies generated by Fisk for the chemically pumped laser, we predict an overall efficiency of 2 to 3% for this sensitizer, using conventional flashlamps as a source for exciting the sensitizer. This value could be improved if the temperature of the photon source were higher. The use of exploding wires, surface discharges, or doped flashlamps may result in a significant improvement in the overall efficiency. One of the most attractive features of using a sensitizer is that no chemical regeneration step is necessary, which may eliminate many of the problems with flowing systems. And, since chemical regeneration accounts for a significant fraction of the cost of chemically pumped iodine lasers, the use of sensitizers may also substantially reduce the total system cost. It therefore seems worthwhile to investigate more extensively the use of such sensitizers for generating excited O_2 for laser applications.

Another pumping scheme for iodine lasers utilizes an electron-beam-sustained discharge to excite the oxygen-iodine mixture. Early experiments using discharges failed to produce concentrations of $^1\Delta$ in excess of 10%, a value which is too low to induce an inversion in iodine.¹¹² High-power discharges produce a substantial number of free oxygen atoms, which rapidly quench the $^1\Delta$ molecules. Recent modeling studies by Garscadden¹¹³ indicate that discharges with lower E/N will produce a much larger fraction of $^1\Delta$. In discharges with an electric-field energy per electron in a narrow range around 10 Townsend, he predicts that free-oxygen quenching will be minimized, since little energy goes into dissociation. Garscadden also finds that the fractional energy transfer into $^1\Delta$ is 48%. He predicts that populations of 40% and 15% in $^1\Delta$ and $^1\Sigma$, respectively, are attainable with an e-beam-sustained discharge.

In order to have sufficient stopping power for the e-beam, a high-pressure Ar buffer is required. If the pressure is too high, however, heavy-particle quenching of the $^1\Delta$ may be a problem. It has been suggested that running the laser at lower temperatures

(195 K) may reduce this problem sufficiently. Although the model calculations by Garscadden are not completely self-consistent, the predicted high concentration of $^1\Delta$ is very encouraging.

Authors: K. C. Kulander and R. A. Haas

Advances in Atomic and Molecular Physics

The Theoretical Atomic and Molecular Physics (TAMP) group operates within the laser program, providing the program with timely, fundamental theoretical support for understanding and modeling the complex atomic and molecular interaction processes that impact the development of advanced lasers. There is an ongoing need for small-scale investigations of new concepts and potential candidates for advanced fusion lasers other than the rare-gas halide systems that presently comprise the baseline program. From its inception in 1975, the TAMP group has supported experimental efforts in both advanced lasers and isotope separation. We have summarized much of this TAMP work in the preceding articles of this section and in Section 10 of this annual report.

The timeliness with which we can model a complex physical process depends critically on the readiness of accurate supporting theory and computational technology. In striving to be at the forefront of new developments in theoretical methodologies, we have gained expertise in a number of areas, including atomic and molecular structure, photon-excitation dynamics, chemical dynamics, and electron scattering.

In the following articles, we summarize some key developments that took place in these areas during 1980 as a by-product of our activities supporting laser-program projects. We begin by summarizing a theoretical study of the electronic states of CF_3I , a common source for excited-iodine lasers. The next two articles deal with the interaction of laser photons with atomic and molecular systems, respectively. We next describe calculations of electron-atom collision strengths used to model a candidate x-ray laser system. This discussion is followed by two articles which describe our

theoretical advances in the study of reactive chemical systems and electron-molecule collisions. The final article updates additions that have been made to the TAMP remote computational facility.

Author: T. N. Rescigno

Theoretical Studies of the Electronic States of CF₃I

Because of its potential as a laser fusion driver, the iodine laser has been the subject of considerable research over the last decade.¹¹⁴ This laser exhibits favorable optical cross sections, stores energy for extraction in nanosecond time scales, and operates in a gaseous medium. Its low efficiency, however, is a major limitation.

The CF₃I molecule is commonly used in iodine-laser systems as a source of iodine atoms. The photodissociation into the excited channel occurs for a weak continuum located at about 4.7 eV above the ground state. A source of inefficiency in the CF₃I laser is the poor spectral coupling of the radiation source to the dissociation band. Laser pumping has been suggested as a way to improve the band overlap with the pump.¹¹⁵ This would require a highly efficient UV pump laser to maintain overall efficiency, which means sacrificing an efficient UV laser photon to obtain the stored energy of an infrared laser. Other pumping methods include electrical discharge¹¹⁶ and nuclear pumping.¹¹⁷ Both of these methods use electrons to excite the medium and produce the excited I(²P_{1/2}) atoms.

To understand the kinetics and improve the efficiency of electrically or nuclear-pumped lasers, we need to know the electronic excited states, symmetries, transition moments, and electron-impact cross sections for the fuel molecules. (Nuclear-pumped lasers are in actuality excited by the secondary electrons produced

by energetic fission fragments.) Accordingly, Srivastava and Trajmar¹¹⁸ have measured the electron-scattering cross sections for the CF₃I molecule, and we have determined the potential-energy curves and transition probabilities for the excited states. This information allows us to assign the observed features in the spectra, locate forbidden, unobserved states, and determine dissociation products.

The high-energy, small-angle electron-impact spectrum of CF₃I exhibits the 4.67-eV peak normally pumped in the iodine laser, with additional sharper peaks at higher energies. At lower energies and larger scattering angles, the spectra show peaks at 4.13 and 6.22 eV.

We have determined the wave functions for the ground and excited states of CF₃I using a (9s 5p/3s 3p)/(3s 2p/3s 3p) Gaussian basis set.¹¹⁹ We determined the ground-state wave function using the generalized valence bond method, and determined the excited states using the open-shell Hartree-Fock method.¹²⁰ These wave functions give proper dissociation behavior, but do not account for electron correlation differences between the states. We included the spin-orbit interaction among the states in the usual way.¹²¹

In C_{3v} symmetry, the low-lying valence states of CF₃I are ¹A₁ (ground state), ¹E, and ³A₁. After spin-orbit interaction is included, the ¹A₁ state gives one A₁ state; the ³E state leads to A₂, A₁, and 2E states; the ¹E state produces an E state; and, finally, ³A₁ gives A₂ and E excited states. Table 8-15 gives the relationship between the relativistic and nonrelativistic states near the equilibrium geometry, along with their excitation energies; Fig. 8-89 shows the potential-energy curves for these states. Our calculations show that the second A₁ state dissociates to I(²P_{1/2}) and is responsible for the photodissociation band. The large component of ³E in the ²A₁ wave function is consistent with the small cross section observed for this transition. The lowest E

ΔE(eV)	XA ₁ 0.0	1E 4.0	2E 4.1	1A ₂ 4.5	2A ₁ 4.6	3E 5.0	2A ₂ 6.3	4E 6.3
¹ A ₁	0.9931	0	0	0	0.1170	0	0	0
³ E	-0.1170	1.0	0.9314	0.9459	0.9931	0.1141	0.3244	0.3457
¹ E	0	0	0.2370	0	0	-0.9110	0	-0.3377
³ A ₁	0	0	-0.2764	-0.3244	0	-0.3964	0.9459	0.8755

Table 8-15. CF₃I spin-orbit wave functions and excitation energies.

state at 4.1 eV is consistent with a 4.15-eV feature observed by Srivastava and Trajmar^{117,118} for large-angle scattering. We calculated the $2A_2$ and $4E$ states to be near 6.3 eV, which is in agreement with a 6.22-eV feature that is also found in the electron-impact spectra.

To investigate the features observed at higher energies, we carried out further calculations with a basis set that included Rydberg functions. We have not included spin-orbit coupling in these calculations, due to their Rydberg character and the neglect of configuration interaction between states of the same symmetry. The agreement with experiment is quite reasonable, however. The observed higher-energy features are at 7.13, 7.93, 8.25, 8.87, and 9.65 eV. The calculations give a similar pattern, with states at 6.91, 7.34, 7.54, 8.80, and 9.6 eV. More specific assignments must await the calculation of transition strengths and spin-orbit interaction. Of particular interest are the dissociation products and the possibility of the higher-lying states cascading into the state that produces the upper laser level of the iodine laser.

Our study has given the first identification and characterization of the state responsible for producing excited iodine atoms in

photolysis. The information on the higher-lying states will be valuable in understanding the kinetics of electrically and nuclear-pumped lasers. We need to do further work on the Rydberg states to correctly include the spin-orbit and configuration interaction.

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Photon-Atom Interaction Modeling

TAMP supports the laser isotope-separation program at LLNL by providing models for the fundamental physics underlying the excitation of an atomic vapor by laser radiation. In our initial studies, we formulated a coherent excitation model based upon the rotating-wave approximation of the time-dependent Schrödinger equation,¹²² and we explored the consequences of this model for the case of monochromatic excitation.¹²³⁻¹²⁸ We also examined generalizations that would admit incoherence¹²⁹⁻¹³¹ and pulse modulation¹³² of the exciting radiation. These generalizations would also include effects of magnetic sublevels¹³³ and hyperfine structure of the vapor atoms.¹³⁴ Further details appear in previous annual reports.¹³⁵

Our investigations have shown that atomic vapors excited by coherent radiation can exhibit behavior that differs appreciably from the behavior predicted by using traditional rate equations. Our studies have provided valuable insight into this newly discovered regime. Recognizing that the excitation process, in turn, affects the exciting field (e.g., incoherent excitation acts to deplete the intensity in accord with Beer's law), we have now begun to investigate propagation of laser beams through near-resonant media. In these studies, we treat the coupled Maxwell and Schrödinger equations¹³⁶ in a self-consistent manner.

The subject of atomic motion in a resonant radiation field continues to attract attention as a rapidly developing field of research,¹³⁷ finding application to atomic trapping and cooling and atomic-beam deflection spectroscopy as well as to isotope

Fig. 8-89. Potential energy curves for CF_3I .

separation. Previous theoretical approaches to this subject have dealt with specific proposed applications and special problems.¹³⁸⁻¹⁴²

We have developed from first principles a general theory of the motion of a two-level atom in a resonant or near-resonant electromagnetic wave of arbitrary amplitude and phase, including effects of radiative relaxation due to interaction with the quantized vacuum field as well as fluctuations of the radiation force and the associated diffusion of atomic momentum due to spontaneous and induced emission and absorption processes.¹⁴³

Our approach begins with the Wigner distribution function, $f(\vec{x}, \vec{p})$, describing the quantum-mechanical state of a structureless point particle, for which (by way of Heisenberg equations) we obtain a classical Liouville equation of motion. We next admit a two-level internal atomic structure to obtain generalizations of the Bloch equations. In the quasiclassical limit obtained by neglecting terms that vanish as $\hbar \rightarrow 0$, we obtain the equations

$$\left(\frac{\partial}{\partial t} + \frac{\vec{p}}{M} \cdot \vec{\nabla} \right) f = -\frac{1}{2} \hbar \left(\vec{\nabla} \Omega \cdot \vec{\nabla} \rho U + \Omega \vec{\nabla} \theta \cdot \vec{\nabla} \rho V \right) + \sum_{i,j} Q^{ij} \frac{\partial^2 (f + W)}{\partial p^i \partial p^j}, \quad (47)$$

$$\left(\frac{\partial}{\partial t} + \frac{\vec{p}}{M} \cdot \vec{\nabla} \right) U = \left(\Delta + \frac{\vec{p} \cdot \vec{\nabla} \theta}{M} \right) V - \frac{1}{2} A U - \frac{1}{2} \hbar \vec{\nabla} \Omega \cdot \vec{\nabla} \rho f, \quad (48)$$

$$\left(\frac{\partial}{\partial t} + \frac{\vec{p}}{M} \cdot \vec{\nabla} \right) V = - \left(\Delta + \frac{\vec{p} \cdot \vec{\nabla} \theta}{M} \right) U - \frac{1}{2} A V + \Omega W - \frac{1}{2} \hbar \Omega \vec{\nabla} \theta \cdot \vec{\nabla} \rho f, \quad (49)$$

and

$$\left(\frac{\partial}{\partial t} + \frac{\vec{p}}{M} \cdot \vec{\nabla} \right) W = \Omega V - \Lambda (f + W) - \sum_{i,j} Q^{ij} \frac{\partial^2 (f + W)}{\partial p^i \partial p^j}. \quad (50)$$

Here, θ is the effective phase of the applied field, Ω is the resonant Rabi frequency, and U , V , and W are the Wigner distributions of the Bloch variables u , v , and w , where

$$u = \int d^3x \int d^3p U(\vec{x}, \vec{p}), \quad (51)$$

etc. In Eqs. (48), (49), and (50), A is the Einstein spontaneous-emission coefficient. In Eqs. (47) and (50),

$$Q^{ij} = \frac{1}{2} \hbar^{-2} \int d^3k k^i k^j Z(\vec{k}), \quad (52)$$

where

$$Z(\vec{k}) = (4\pi\hbar)^{-1} W_k \delta(W_k - W_0) |\vec{\theta}|^2 (1 - \cos^2 \theta) \quad (53)$$

and θ is the angle between the propagation vector, $\vec{k}$, and the atomic dipole moment, $\vec{\mu}$.

These quasiclassical equations [Eqs. (47) through (50)] provide a simple derivation of the optical Stern-Gerlach effect,¹³⁹ i.e., the Ehrenfest-Bloch equations describing both the motion of a centroid of the atomic wave packet¹⁴⁰ and the induced momentum diffusion. For details, refer to Refs. 139 and 140.

We have also studied theoretically the effects of quantum-mechanical fluctuations of the resonance-radiation force on atomic motion. We found¹⁴⁴ that fluctuations due to both induced and spontaneous absorption-emission processes give rise to a diffusion of atomic momentum that is describable by a Fokker-Planck equation. We showed that quantum-mechanical fluctuations of the radiation force

- Place a lower bound on the temperature that can be achieved by radiation cooling.
 - Inhibit cooling in a strong standing wave.
 - Lead to finite (and often short) confinement times for atoms in radiation traps.
- These findings extend our earlier work,¹⁴¹ wherein we found that the occurrence of strong phase fluctuations can act to damp atomic velocity.

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Effects of Strong Laser Fields on Chemical Reactions

There is considerable interest in the effects of laser radiation on chemical reactions.¹⁴⁵ In particular, several research groups are investigating the effect of a strong laser field

on a chemical reaction in which photons can be resonantly absorbed only during the collision. We have found in our research that the reaction probability can be significantly enhanced if the laser radiation couples the initial electronic surface to a highly reactive excited surface. Quantal¹⁴⁶ and classical¹⁴⁷ studies of model systems show that strong fields can induce reaction probabilities >0.05 at collision energies well below the (field-off) reaction threshold.

We recently completed a detailed comparative study of the quantal and

classical treatments of this process.¹⁴⁸ We discovered the interesting (but in retrospect not surprising) fact that it is also possible to reduce the probability of reaction by coupling the initial electronic surface with a laser field to a nonreactive excited surface.¹⁴⁹ Our studies show that the specific collision energies and laser frequencies at which the enhanceive or inhibitive effects are most pronounced depend very sensitively on subtle characteristics of the potential-energy surfaces involved. We also established that the completely classical model of Orel and Miller¹⁴⁷ can be used to predict conditions for observing maximum effects.

In Fig. 8-90, we show a schematic diagram of the laser inhibition mechanism. Curves V_1 and V_2 are potential energies for the initial and excited states, plotted as functions of a reaction coordinate along the minimum-energy path of the lower surface. The line designated E_T corresponds to the initial total energy of the reactants. If, during the course of a collision, the system absorbs a photon with energy $\hbar\omega$, then the excited surface controls the motion. The reaction will not take place if the total energy (now $E_T + \hbar\omega$) is below the reaction threshold of the upper surface.

We studied¹⁴⁹ a model of the $H + H_2$ system (defined in Ref. 145) in which the "lower" surface is a London-Eyring-Polanyi-Sato (LEPS) surface but with a Sato parameter equal to 0.20. The reaction barrier for this surface is approximately 3 kcal/mole. We chose the "upper" surface to be the Porter-Karplus surface, whose saddle point is ~ 9 kcal/mole above the asymptotic minimum.

Figure 8-91 shows the total reaction probability for reactants initially on the lower surface, as a function of collision energy, both with and without the laser field. We selected the photon energy to give an asymptotic resonance energy defect, $\Delta = \hbar\omega - (V_2 - V_1)$, of 0.03 eV. The strength of the coupling between the electronic surfaces due to the laser field, $\mu \cdot E$, has the values 0.005 and 0.01 eV. In both cases, we find the reaction probability reduced significantly, with the inhibition being most pronounced just above the reaction threshold. It is possible to find the reaction products on either of the (dressed) surfaces; however, we report only the total reactive probability. As

Fig. 8-90. Schematic diagram of the mechanism for laser inhibition of a chemical reaction.

Fig. 8-91. Reaction probabilities for the model collinear $H + H_2$.

Fig. 8-92. Classical and quantum-mechanical reaction probabilities for the model collinear $H + LiF$.

expected, most of the reactive probability appears in the lower dressed state and most of the nonreactive fluence is found in the upper state.

We studied the model system of Light and Altenberger-Siczek,¹⁴⁶

The excited electronic state for this system is highly reactive and is reached most effectively in the incident channel by photons whose energies are $\gtrsim 6.57$ eV. In Fig. 8-92, we compare the classical and quantum-mechanical reactive probabilities for a laser with photon energy of 6.2 eV and coupling strength of $\mu \cdot E_0 = 0.01$ eV. The reactant diatomic is initially in its ground vibrational state, $v = 0$. The field-free reactive probabilities are also shown from both calculations. The difference between the classical and quantum-mechanical reactive thresholds is due to the zero-point energy at the barrier in the latter calculation.

Orel and Miller¹⁴⁷ correctly explained the two narrow peaks in the classical laser-coupled probability; They are due to points on the initial surface where the trajectories have turning points in both the vibrational and relative translational motion, and where the energy difference between the two surfaces exactly matches the photon energy. The system remains in the vicinity of these points long enough to have a high probability of resonantly absorbing energy from the field, and then reacting via the upper surface. The classical vibrational motion has inner and outer turning points, whereas the quantum ground state has its maximum amplitude in the center of the potential well. Therefore, we expect only a single broad peak in the quantum-mechanical results in place of the two peaks found classically.

As Fig. 8-92 shows, this expectation precisely matches the result we obtained. In Fig. 8-92, the classical and quantum-mechanical reaction probabilities are plotted as functions of translational energy for $\hbar\omega = 6.2$ eV, $\mu \cdot E = 0.01$ eV, and the ground initial vibrational state. Dashed curves are the field-free results, and the solid curves are the results in the presence of the laser. We find the two classical peaks at

$\epsilon_T \cong 0.16$ and 0.21 eV and the quantum peak at $\epsilon_T \cong 0.23$ eV.

At slightly higher translational energies ($\epsilon_T \cong 0.25$ eV), the classical reaction probability begins to rise, shows a slight hump at ~ 0.35 eV, and remains large until it merges with the field-free reaction probability at ~ 0.45 eV. The quantum-mechanical probability shows a slight, broad bump at $\epsilon_T \cong 0.37$ eV, where it rises to 0.02 but remains at least an order of magnitude smaller than the classical result up to the field-free threshold. The effect of the laser radiation on the reaction probability is negligible for collision energies at which the field-off result is large. This is true for both the classical and quantum-mechanical models. The humps appear in both calculations at ~ 0.36 eV, where it becomes possible to resonantly absorb a photon asymptotically in the entrance channel.

We have learned that the classical model of Orel and Miller¹⁴⁷ is capable of qualitatively reproducing most of the interesting effects of the laser field on the reaction. Although this model predicts quite accurately the positions of the peaks in the enhanced probability and exhibits the laser inhibition of the reaction, it greatly overestimates the size of the observed effect. Thus, we can use with confidence the computationally much less expensive classical model to study the qualitative effects on reactions of a strong laser field.

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Collisionally Pumped X-ray Laser Studies

Many of the theoretical and computational capabilities developed as part of our fusion-laser research can be applied to the scientifically challenging problem of demonstrating an x-ray laser that might be pumped with large, high-power Nd:glass lasers or other powerful excitation sources. Along these lines, the possible amplification of VUV or x-ray radiation in a plasma containing multiple charged atomic ions is now attracting considerable attention. In particular, neon-like ions with the ground-state electron configuration $(2s^2 2p^6)1S_0$ are

interesting candidates for creating population inversions in steady state plasmas.^{150,151}

To understand how population inversions can be generated in these neon-like ions, refer to the energy-level diagram shown in Fig. 8-93. The rate for electron-impact excitation of the optically forbidden $(2p^6)^1S_0 \rightarrow (2p^53p)^1S_0$ transition is apparently quite large.^{150,151} At the same time, the $(2p^53s)$ and $(2p^53d)$ excited levels decay very rapidly

Fig. 8-93. Schematic energy-level diagram for neon-like ions.

Fig. 8-94. Electron collision strengths of the $(2p^6)^1S \rightarrow (2p^53s)^1,^3P$ transitions in Kr^{+26} .

Table 8-16. Dominant electron-excitation cross sections for Kr^{+26} .

Transition	Final state	Cross section (cm^2) at $E_e = 2.0$ keV	
$2s^22p^6 \rightarrow 2s^22p^53d$	1P_1	2.1×10^{-20}	Upper laser level
	3D_1	1.6×10^{-20}	
$2s^22p^6 \rightarrow 2s^22p^53p$	1S_0	0.70×10^{-20}	
	3P_0	0.37×10^{-20}	
$2s^22p^6 \rightarrow 2s2p^63d$	1D_2	0.46×10^{-20}	Lower laser level
$2s^22p^6 \rightarrow 2s^22p^53s$	1P_1	0.51×10^{-21}	

to the ground state via strongly allowed electric dipole transitions. Thus, population inversions may be achieved on the $(2p^53p)^1S_0 \rightarrow (2p^53s)^1,^3P_1$ transitions under appropriate plasma conditions.

Electron-impact excitation rates of neon-like ions are needed to model the proposed lasers; however, except for low-Z elements and Fe^{+16} , these rates are not available. To support upcoming experiments on Kr^{+26} and to provide essential rates for modeling,¹⁵² we have calculated both the optical and the electron-impact excitation rates for the $(2p^6)^1S_0$ state of Ne-like krypton.

In our calculations we used a single-configuration description of the $(2p^6)^1S_0$ ground state. In calculating the $3s$, $3p$, and $3d$ excited levels, we accounted for both configuration mixing and spin-orbit coupling among the levels. Our calculated oscillator strengths for the electric dipole transitions $(2p^6)^1S_0 \rightarrow (2p^53s)^1,^3P_1$ and $(2p^53d)^1,^3P_1, ^3D_1$, in both Fe^{+16} and Kr^{+26} , are in good agreement with previous theoretical results.¹⁵³ The same atomic wave functions were used in a distorted-wave approximation¹⁵⁴ to calculate electron-impact excitation cross sections for Kr^{+26} , from threshold (~ 1.7 keV) to 5.0 keV.

We found that configuration mixing among the $n = 3$ excited levels has a significant effect on the calculated electron collision strengths (see Ref. 155 for the definition of electron collision strength). Figure 8-94 illustrates this finding for the $(2p^6)^1S \rightarrow (2p^53s)^1,^3P$ transitions; the black and red lines denote results obtained with and without configuration mixing, respectively. The configuration-mixing effect is most pronounced for the optically allowed transitions, and is also evident in the calculated optical oscillator strengths.

Table 8-16 summarizes the dominant electron-impact-excitation cross sections for the ground state of Kr^{+26} . Our calculations show that, for the anticipated plasma

temperatures (1.5 to 2.0 keV), the optically allowed transitions, $(2p^6)^1S_0 \rightarrow (2p^5 3d)^3D_1$ and 1P_1 , have the largest cross sections (2.1 and $1.6 \times 10^{-20} \text{cm}^2$, respectively). The cross section for exciting the upper laser level $(2p^5 3p)^1S_0$ is about three times smaller ($0.7 \times 10^{-20} \text{cm}^2$) but is still significant. The cross sections for collisional pumping of the lower laser levels $(2p^5 3s)^1, ^3P_1$ from the ground state are smaller still ($\sim 6.0 \times 10^{-22} \text{cm}^2$).

In the course of our work, we also calculated the rates for excitation transfer between levels belonging to the same configuration, e.g. $(2p^5 3p)$, as well as rates for the emission of electric quadrupole radiation from the excited states. In the latter case, we found a surprisingly large rate, $1.1 \times 10^{11} \text{s}^{-1}$, for the $(2s 2p^6 3d)^1D_2 \rightarrow (2p^6)^1S_0$ quadrupole transition. This theoretical result raises the possibility of using this line at 6.0 Å, along with the calculated radiative and collisional rates, to diagnose some of the plasma parameters in upcoming experiments.¹⁵²

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Quasiclassical and Quantum Wave-Packet Studies of Reactive Scattering

Classical trajectory methods are used far more widely than other methods for studying the dynamics of collisions between molecules. Using classical methods, we can obtain accurate results for highly averaged quantities, such as reaction-rate constants. If we need more specific quantum-state-to-quantum-state reaction probabilities, however, we must use quasiclassical theories to extract quantized information from the trajectory results. Precise kinetic models of gas-laser systems require such detailed information.

The simplest quasiclassical theory assigns quantum states by binning the trajectories; i.e., quantum numbers are identified by finding the quantum state whose eigen-energies are closest to the asymptotically partitioned energies of the trajectory. We call this binning procedure the quasiclassical trajectory histogram (QCTH) method.^{156,157}

Results obtained by binning trajectories generally disagree with quantum calculations. Improved quasiclassical schemes use energy moments of the final-state distributions to obtain more accurate results. It is important to know whether the classically generated moments are equal to the exact quantum moments.

During the past year we carried out extensive classical-trajectory and quantum-mechanical wave-packet calculations on a model collinear system chosen to represent $\text{H} + \text{H}_2$, to establish whether the classical and quantum moments of the final vibrational-energy distributions agree.¹⁵⁸ We performed calculations for three different initial vibrational states of the diatomic and for a wide range of collision energies, including energies above the dissociation threshold.

The standard close-coupling methods used in quantum reactive scattering are inapplicable to higher-energy collisions because of the large number of open channels. We use the time-dependent wave-packet method, which works particularly well at higher energies.¹⁵⁹ Using this approach, we are able to treat exactly the three-body breakup that becomes possible when the energy exceeds the dissociation threshold. The quantum methods we have used are very similar to those developed by McCullough and Wyatt,¹⁶⁰ which they used to study $\text{H} + \text{H}_2$ at collision energies much lower than those in which we are interested.

We define an initial wave packet that has a Gaussian distribution in relative translational momentum, and a particular vibrational state of the diatomic molecule. We follow the evolution of this wave packet through the collision until it returns to the asymptotic region, then evaluate transition probabilities by projecting the final wave packet onto separable translational-vibrational eigenstates. By defining a grid of points that spans the important region of configuration space, and by using standard finite-difference techniques, we evaluate the effect of the Hamiltonian operator on our wave packet. Then, we solve the time-dependent Schrödinger equation to determine the evolution of the wave packet. Thus, from a single wave-packet calculation, we obtain transition probabilities from a single initial vibrational level, n_1 , to all possible final levels, n_2 , for a range of collision energies.

Fig. 8-95. Potential-energy surface for collinear $\text{H} + \text{H}_2$ system.

 Fig. 8-96. (a) Initial and (b) final wave-packet densities for $n_1 = 4$.

Fig. 8-97. Quantal and classical distributions for final vibrational state.

Fig. 8-98. Quantal and classical rms moments over a range of collision energies.

Fig. 8-99. Quantal and classical relative kurtoses over a range of collisional energies.

We used a potential-energy surface of the extended LEPS form, which is described fully in Refs. 158 and 159 and is represented in perspective in Fig. 8-95. The reactants are initially in the valley at the left in Fig. 8-95, and are moving toward the interaction region. If reaction occurs, the products move out in the partially obscured valley to the right. At the collision energies considered (3 to 12 eV), we find that for the $\text{H} + \text{H}_2$ system the reaction probability is negligibly small, and the wave packet returns to the incident valley after the collision. In Figs. 8-96(a) and 8-96(b), we show wave-packet densities before and after collision for the case $n_1 = 4$. In the final wave packet, the five peaks of the initial vibrational wave function are evident, indicating a substantial elastic scattering component. By comparing the positions of the initial and final wave packets, we see that the system is predominantly nonreactive.

We use standard methods to carry out our classical trajectory calculations.^{157,158} We obtain normalized vibrational transition probabilities, P_{n_1, n_2}^N , as a function of collision energy, E_{rel} , from both quasi-classical and quantum calculations. In the

classical case, n_1 is an integer, but n_2 is a continuous variable. We present probabilities that include only the nonreactive and nondissociative final states, normalized such that

$$\sum_{n_2} P_{n_1 n_2}^N(E_{\text{rel}}) = 1 \quad (54)$$

We define the k th moment of the normalized distribution by

$$\langle n_2^k \rangle = \sum_{n_2=0}^{16} P_{n_1 n_2}^N n_2^k \quad (55)$$

From these elementary moments, we calculate the rms moment,

$$\sigma = \langle (n_2 - \langle n_2 \rangle)^2 \rangle^{1/2} \quad (56)$$

the relative skewedness,

$$\alpha_3 = \langle (n_2 - \langle n_2 \rangle)^3 \rangle / \sigma^3 \quad (57)$$

and the relative kurtosis,

$$\alpha_4 = \langle (n_2 - \langle n_2 \rangle)^4 \rangle / \sigma^4 \quad (58)$$

The moments of Eqs. (56) through (58) are preferred to $\{\langle n_2^k \rangle\}_{k=2}^4$ because the former have a well-established significance in statistics.¹⁶¹ The sums in Eqs. (54) and (55) are to be understood as integrals in the classical case (see Ref. 158 for a more precise definition).

We show a comparison between quantal and classical vibrational distributions in Fig. 8-97, where the classical distribution is generated using the standard binning procedure (QCTH). As stated above, the classical results differ substantially from the quantal results in detail but agree well in an average sense.

We found that the lower moments of the classical distributions converged very rapidly. We show in Figs. 8-98 and 8-99 the quantal and classical (QCT) rms moments and the relative kurtoses for the three different initial states considered. Thus, we established that the lower moments are in quite reasonable agreement over the entire energy range but that the disagreement between the higher moments is small only for the case $n_1 = 4$. Our results indicate that only

the lower moments of the classical distributions should be used to quantize the results in improved quasiclassical theories.

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New Theoretical Methods in Electron Scattering

In the past few years, considerable progress has been made in the development of theoretical methods for studying low-energy electron-molecule collision problems. Yet, there continues to be an active interest in the investigation of new techniques, since fully ab initio calculations on all but the simplest systems remain a formidable computational undertaking.¹⁶² The principal difficulties encountered in low-energy electron-molecule scattering are the noncentral nature of the interaction potential and the necessity for accurately treating nonlocal exchange forces.

Discrete-basis-set methods based on multicentered expansions have received a great deal of attention in the recent literature as a result of the success of such approaches in molecular bound-state calculations. The techniques that have been developed and applied most extensively are the methods based on separable expansions in the Lippmann-Schwinger integral equation (T-matrix methods)¹⁶³⁻¹⁶⁵ and the analytic R-matrix approach of Schneider.¹⁶⁶ These techniques were originally implemented by projecting the electron-molecule interaction potential onto a finite set of normalizable functions, $\{\psi_\alpha(\vec{r})\}$, such that

$$U \approx \sum_{\alpha, \beta} \psi_\alpha(\vec{r}) U_{\alpha\beta} \psi_\beta(\vec{r}) \quad (59)$$

The approximation of Eq. (59) leads to a particularly simple expression for the scattering amplitude, since all required integrals can be obtained in closed form when Cartesian Gaussians are chosen for the expansion basis.

It has been found, unfortunately, that errors due to truncation of the potential can be made small only by the use of very large

basis sets. This is particularly true at low energies and for polar molecules and is an obvious consequence of the difficulty in describing long-range forces with short-range functions. To circumvent this difficulty, we developed an expansion method that retains much of the simplicity of the original T-matrix method but that does not suffer from the convergence problems inherent in any scheme that uses short-range functions to describe long-range interaction forces. Our idea was to formulate the separable approximation in a distorted-wave picture that allows us to first isolate the direct interactions that fall off slowly at long range and then treat them exactly. It is only the nonlocal exchange interactions, which fall off rapidly outside the charge distribution of the target atom or molecule, that are projected onto a discrete set of basis functions.

This approach represents something of a compromise between purely numerical and purely analytic techniques. Our technique is simply illustrated by the case of single-particle potential scattering described by the radial Schrödinger equation

$$\left(\frac{d^2}{dr^2} - \frac{\ell(\ell+1)}{r^2} + k^2 \right) \phi_\ell(r) = \int U(r, r') \phi_\ell(r') dr' , \quad (60)$$

where U is taken to be a spherically symmetric potential operator. We now assume that the potential operator, U , can be split into two parts

$$U(r, r') = U_I(r) \delta(r - r') + U_{II}(r, r') , \quad (61)$$

where $U_I(r)$ contains all the local (long-range) components of the interaction and U_{II} is the nonlocal part of the potential. We define a standing-wave Green's function, $G_I^\ell(r, r')$, for the potential U_I according to

$$\begin{aligned} G_I^\ell(r, r') &= -1/k \phi_\ell^I(r) f_\ell^I(r'), \quad r' > r , \\ &= -1/k \phi_\ell^I(r') f_\ell^I(r), \quad r' < r , \end{aligned} \quad (62)$$

where $\phi_\ell^I(r)$ is a solution for scattering by the local potential, U_I (which is regular at the origin), and $f_\ell^I(r)$ is an irregular solution with the asymptotic form

$$f_\ell^I(r) \sim \cos(kr - \ell\pi/2) \quad r \rightarrow \infty \quad (63)$$

Our object is to formulate a computational scheme for evaluating the tangent of the phase shift

$$K_\ell = \tan \delta_\ell = -\frac{1}{k} \langle \phi_\ell^0 | U | \phi_\ell \rangle , \quad (64)$$

where ϕ_ℓ^0 is the regular Riccati-Bessel function and ϕ_ℓ is a regular standing-wave solution to Eq. (60). The so-called "two-potential formula"¹⁶⁷ can be used to express K_ℓ as

$$K_\ell = K_\ell^I - 1/k \langle \phi_\ell^I | U_{II} | \phi_\ell \rangle , \quad (65)$$

where K_ℓ^I is the tangent of the phase shift produced by the potential U_I alone, and ϕ_ℓ is the wave function for the complete potential.

In the computational scheme we propose, the second term on the right side of Eq. (65) is approximated by making a separable approximation for the potential U_{II} :

$$U_{II}(r, r') \approx \sum_{\alpha, \beta=1}^N \psi_\alpha(r) \langle \psi_\alpha | U_{II} | \psi_\beta \rangle \psi_\beta(r') \quad (66)$$

With this approximation, it is possible to derive the following result for the tangent of the phase shift:

$$\begin{aligned} K_\ell &= K_\ell^I - 1/k \sum_{\alpha, \beta} \langle \phi_\ell^I | \psi_\alpha \rangle \\ &\times [(\underline{1} - \underline{U}_{II} \underline{G}_I^{-1}) \underline{U}_{II}]_{\alpha\beta} \langle \psi_\beta | \phi_\ell^I \rangle \end{aligned} \quad (67)$$

The double bars denote matrices obtained by projection onto the finite set of basis functions, $\{\psi_\alpha\}$.

It is important to realize that, for Eq. (67) to be a good approximation, it is only necessary that the basis set used in Eq. (66) adequately span the region of space over which the short-range potential, U_{II} , acts. We applied this modified separable approximation to a number of electron-atom potential scattering problems and found that it did, indeed, converge quite rapidly.¹⁶⁸ We also found that adequate results could be obtained in many cases with as few as three basis functions.

We expect that the technique we have just described will be quite useful in low-energy electron-molecule scattering problems. It is the long-range forces of multipolar form encountered in electron-molecule collisions that lead to the slow rate of convergence of the usual separable expansion or T-matrix method, especially at very low energies. By

choosing the reference or distorted-wave potential appropriately, we can apply the separable approximation to expand only the important short-range exchange interactions. The latter can be quite strong and, in general, highly aspherical, but they can be quite adequately represented in terms of a discrete set of multicentered, square-integrable functions.

The most appropriate choice for the reference potential would be the static electron-molecule interaction, in which case construction of the reference Green's function would involve the numerical solution of a set of coupled ordinary differential equations. The modified separable-expansion technique would still obviate the need for single-center expansions of exchange terms, which is by far the most time-consuming aspect of purely numerical approaches to electron-molecule scattering. Currently, the modified separable-expansion technique is being implemented for solving realistic electron-molecule scattering problems.

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TAMP Remote Computational Facility

The remote computational facility (RCF) of the TAMP group was designed to enhance the group's computational power for analyzing laser-related problems. To accomplish this with a minicomputer, we attached a Floating Point System AP-120B array processor to the PDP-11/55 host. The RCF layout is shown in Fig. 8-100. As demonstrated by the work on Monte Carlo glass simulations,^{169,170} we have been able to achieve performance speeds comparable to the CDC-7600 computer; however, the PDP-11/55 data storage was limited. To increase the data storage, we added in 1980 an extended RP03 fixed disk and a set of dual-density floppy disk drives. The extended RP03 disk provides 80 megabytes of storage for the user.

To expand our time-sharing capability, we also purchased an additional partition of memory, giving 96K words of storage on the PDP-11/55 host. An additional terminal was

Fig. 8-100. Schematic of the TAMP group's remote computational facility.

added to the system, and several cables were pulled and made available for future terminals. The LSI-11 and the new terminal were consolidated into a new users' room.

Software upgrades during 1980 included the installation of the VERSATEC graphics library, the NAG extended mathematics library, and the AP mathematics library for the array processor. The host system was upgraded to RSX-11M V3.2, and the LSI-11 system was improved to RT-11 V4.

The RCF has continued to see a high level of activity despite the numerous modifications. The extensive classical trajectory

study on the $I + HBr$ reaction was completed, and the Monte Carlo glass simulation studies continue.¹⁷⁰ In 1981, we plan to begin several new dynamics studies using classical trajectories. Hardware modifications will include addition of the final partition of memory, and we will bring several new terminals on-line for the host PDP-11/55.

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TAMP Remote Computational Facility 8-81

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Section 9 Energy and Military Applications

Energy and Military Applications

Contents

Introduction	9-1
Accomplishments Associated with the HYLIFE Chamber Concept	9-4
Operating Parameters of Inertial Fusion	9-4
Parametric Studies	9-9
The Influence of Chamber Pulse Rate on Structural Stress and Lithium Flow	9-9
Blanket-Shield Design Options	9-15
The Effect of Chamber Ports on Neutron Leakage and Blanket Performance	9-21
Materials Considerations	9-25
Solubility of Hydrogen Isotopes in Lead-Lithium Solutions	9-25
Corrosion of 2-1/4 Cr-1 Mo in Lithium and Lead-Lithium Solutions	9-26
Fluid and Pumping Considerations	9-30
Transport Processes in the HYLIFE Reaction Chamber	9-30
Vortex-Flow Slanted-Jet Improvement to HYLIFE	9-42
Heavy-Ion-Beam Ballistic Propagation in Liquid-Metal-Wall Reaction Chambers	9-43
Experimental Plan to Measure Effects of Gas Blowing Through Jets	9-47
Experimental Investigation of the Stability with Liquid-Sheet Jets: A Comparison with Linear and Nonlinear Theory	9-51
Vacuum Pumping	9-56
Electromagnetic-Pump Considerations	9-58
Lithium Safety Considerations	9-61
Build-up of Tritium in a Liquid Lithium Breeding Blanket	9-63
Component Design	9-68
Stress Analysis of the Alternate Helical Bar First Structural Wall (FSW)	9-68
Nozzle Plate	9-71
Fatigue Crack Growth in HYLIFE Components	9-76
Allowable Driver and Target Factory Cost	9-79
Engineering Test Facility (ETF) Planning	9-88
A New Liquid Metal Chamber of High Repetition-Rate Operation: the JADE Chamber	9-91
Ultra-High-Pressure Shock-Wave Experiments	9-99
References	9-101

Energy and Military Applications

Introduction

The current emphasis in the inertial fusion program is directed toward such near-term military objectives as understanding the

- Fundamental physics of implosion and thermonuclear burn.
- Properties of matter at high temperatures and pressures (e.g., equation-of-state and opacity measurements).
- Eventual use of the fusion output for the simulation of nuclear weapon effects.

The current research and development effort is also applicable to the long-term goal of energy production. Starting from the current base of experimental results, we try to foresee both the future energy and military applications of inertial fusion. We estimate the potential benefits and risks of each application, evaluate alternative development strategies, create conceptual designs of test facilities and power plants, compare the competitive approaches, and then iterate this process. This procedure ensures that when we do commence the engineering phase of the development process we will be seeking the highest potential for each application while avoiding unnecessary costs and risks.

The generation of electric power is particularly attractive. The earliest such plants would burn deuterium, which would be extracted inexpensively from seawater, and tritium, which would be bred from lithium within the plant. A power plant (see Fig. 9-1) would consist of a laser driver (or, perhaps, an ion-beam driver); a reactor chamber; a factory to manufacture the fusion targets, which are often called pellets; and the balance of plant, containing heat-transfer

equipment and electrical generators. Pellets containing deuterium-tritium (D-T) fuel would be injected into a reaction chamber at the rate of 1 to 10 per second, and each pellet would be ignited with lasers at a precise point in the injection cycle. The reactor chamber must be designed to

- Absorb the products of the fusion reaction (x rays, debris from the fuel capsule, and neutrons), which are emitted in extremely short, intense bursts.
- Convert the pulsed-energy releases into a steady flow of heat at a temperature suitable for conversion to electricity.
- Breed new tritium for fuel.

The reaction chamber, the heart of the power plant, is the most challenging portion of the overall design. This chamber must withstand the damage of the incident energy fluence, minimize radioactive wastes, perform with high reliability, and have a lifetime that is long enough to minimize operating costs. We have, therefore, concentrated on chamber designs and innovations that are not only technically feasible, but promise significant benefits over alternative methods of generating electrical power. Shown in Fig. 9-1 is an artist's conception of a future inertial fusion power plant, showing the HYLIFE concept for a liquid metal wall (LMW) reaction chamber.

The purposes of our current development efforts and conceptual design studies are to demonstrate the scientific feasibility and to explore, to the greatest extent possible on paper, the engineering feasibility questions. But there are larger issues involved.

Imagine for a moment that fusion power has been demonstrated to be feasible in every way. What are the consequences? One benefit is that we would have an energy source free from the ever rising costs that accompany a depleting resource. Although fusion power plants will most likely have higher capital costs than comparable coal-fired plants or light-water reactors, the negligible fuel costs of an inertial fusion plant would ensure that sometime in the future it would indeed provide electricity at costs competitive with those of power plants that are dependent on the cost-escalating fuels of gas, oil, coal, and uranium.

As important as fuel costs are, the full range of potential advantages of inertial fusion far transcend the resource argument. The real opportunity that fusion power will provide, from the moment it is introduced, is freedom from the significant social costs inherent in the fossil- and fissile-fuel cycles. What would society pay for electric power that is free from the penalties of CO₂ production, acid rain, open-pit mines, uranium tailings, high-level fission-product wastes, or the transportation of plutonium? What value shall we attach to the number of lives saved that would otherwise be lost in the mining and transportation accidents associated with electricity generated from fossil and fissile fuels? While the current focus of our work is on an examination of the technical feasibility of inertial fusion, our vision of the future must be on these larger issues.

The potential of our conceptual designs should, therefore, be judged not only on whether they will work, but also on features that presage significantly reduced human and environmental impact. The examples follow:

- The fuel cycle is closed within the plant, with the exception of a few cubic metres of heavy water needed per year. There will be no significant effluents.
- The use of thick layers of liquid metal to protect the chamber wall allows the use of ordinary steels and results in very moderate amounts of low-level radioactive waste (and then only upon decommission of the power plant).

- Because inertial fusion is designed to routinely (or repetitively) operate with the maximum conceivable energy release, accident safety should be significantly enhanced.
- Since the laser driver can be separated from the reactor building, and since the chamber wall may never need to be replaced, the maintenance required and, therefore, the occupational radiation hazard to workers can be quite small. The reader should keep in mind these larger issues while judging the aspects of inertial fusion reactor design described in this section.

In discussing the accomplishments of the past year, we begin by describing the expected operating parameters of a 1-GW_e inertial fusion electric generating plant. The range of allowed operating points, in terms of yield, repetition rate, driver energy, and recirculated power fraction, is constrained by the achievable pellet gain and driver efficiency, the desired output power of the plant, and the yield/repetition-rate limits of different reactor chamber concepts. While the final choice of operating points will be made on an economic basis, it is valuable to synthesize our current understanding of how the technical aspects lead to certain expected combinations of fusion pellet, driver, and reaction chamber.

With the stage thus set, we then discuss progress made in the conceptual design of the HYLIFE concept for a liquid metal wall (LMW) chamber. Through parametric tradeoff studies, we demonstrate how the structural stress constraints of the chamber influence the allowed yield and repetition rate, the lithium flow rate, and the pumping power. We can control both the steady neutron-damage rate and the transient wall stresses that result from the pulsed nature of the energy release by suitable design of the liquid metal configuration inside the chamber. We have also completed studies of tritium buildup and inventory, and of neutron leakage from beam ports.

The presence of liquid metal in the HYLIFE concept dominates many aspects of the design. For example, we have examined the solubility of hydrogen isotopes in liquid lithium and lead, and we have examined the prospects for heavy-ion beam propagation in mixtures of metal vapor (from the liquid metal), helium, deuterium, and tritium. Through the use of water simulations, we have obtained experimental results on the stability of sheet-type jets in a vacuum; combined with a known correspondence between the behavior of water and liquid metals, findings give us confidence that liquid metal can be injected into the chamber in a stable manner. Experimental work on the corrosion of ferritic steels by mixtures of liquid metals is also reported. We have examined electromagnetic liquid-metal pumps to determine the projected level of efficiency, since this type of pump obviates the need for extending rotating shafts into the radioactive liquid metal, and is resistant to fluid-borne shocks. Finally, we discuss the prospects for obtaining an acceptable level of safety with the use of liquid metal.

In the area of component design, we present results of structural and stress analyses of the first structural wall, the nozzle plate, and a novel “first wall” made of a dense helical winding of bars or cables, which may be useful for reducing the liquid-wall impact forces.

Looking to the future, we use previous cost estimates for a HYLIFE power plant, in conjunction with some pellet-gain and electricity-cost scenarios, to determine the allowable sum of costs for the driver and the pellet factory. We have also investigated some options for an inertial fusion engineering test facility—particularly the question of determining the minimum thermal size of a

facility to integrate a reactor and a laser into a repetitively pulsed system. In both the cost study and the test-facility study, our intention is to explore alternative future options for planning purposes.

Although HYLIFE has many attractive features, we are interested in extending the LMW concept to increase versatility and lower costs by examining different irradiation geometries, different liquid metals, and higher repetition rates. A reactor concept which shows particular promise is described.

Most of the military applications of lasers are described in a companion classified report. However, in part of this section, we report on the use of lasers to obtain equation-of-state information through the generation of ultrahigh-pressure shock waves. Otherwise, experiments at the 1-TPa pressure level required the use of nuclear weapons.

The development of inertial fusion is a move to substitute an energy source which depends principally on an innovative new technology for energy sources which depend principally on the consumption of depletable resources. Inertial fusion will be increasingly attractive as these traditional technologies become more costly, both in economic and social terms, in the future. In the near term, we can obtain some types of weapons design and effects data in controlled and repeatable experiments which are less expensive than underground tests, and which complement the underground data. In the longer term, we have the potential to generate useful energy from nuclear fusion commercially, relatively independent of the increasingly unacceptable human and economic costs of mining and transportation, and the burning of fossil and fissile fuels.

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Accomplishments Associated with the HYLIFE Chamber Concept

Operating Parameters of Inertial Fusion

An inertial confinement fusion (ICF) power plant can be configured from a wide variety of driver and reaction chamber concepts, e.g., a laser driver with a wetted-wall chamber, or a light-ion-beam with a gas-filled chamber.

Four major considerations determine the set of allowable configurations and operating parameters for an integrated system.

- The fusion energy gain (the product of driver efficiency and pellet gain) largely determines the system efficiency of the power plant through its effect on the

fraction of gross electrical power which must be recirculated to the driver.

- The desired size of the power plant, combined with the fusion energy gain, determines the range of yield, repetition rate, and driver energy required of a chamber-driver combination.
- Technical feasibility. Thermal and dynamic stress limits, ablation temperature limits, and the time required to reestablish ambient conditions (consistent with adequate reliability and lifetime) determine the technically feasible combinations of yield and repetition rate for various chamber concepts.
- Driver and reaction chamber combinations. The overlap between the propagation constraints of the driver and the achievable chamber background density largely determine viable combinations of driver and reaction chamber.

These four considerations facilitate selection of acceptable power plant configurations and operating regimes within a set of technically feasible options. The selection of any single operating point,

however, will be ultimately determined by the capital and operating costs of the system.

Power Plant Considerations. The fusion power of the plant is the product of the yield per pulse (Y) and the repetition rate (w). The yield of each fusion pellet is the product of the driver energy (E) required to ignite the pellets and the pellet gain [$G(E)$], which is a function of driver energy:

$$P_f = Yw = E G(E)w \quad (1)$$

Recent estimates¹ of the pellet gains we might be able to achieve in the future can be plotted as total fusion yield as a function of driver energy. In Fig. 9-2, three gain regions are labeled. The optimistic gain band is based on ideal implosions of double-shell reactor pellets. Based on the current thermonuclear physics data base, however, it is prudent to estimate the gain more conservatively² by designing pellets to implode to twice as high an ignition temperature as calculated by the computer codes, and use a burn efficiency degraded to 70% of the calculated value. Bands of uncertainty are given which correspond to differences in laser wavelength or ion energy, the degree of pulse shaping, irradiation uniformity, and other factors.

As the commercialization of inertial fusion is approached, we expect to learn enough to reach a performance level significantly within the intermediate gain region. The upper bound of each band in Fig. 9-2, gives the maximum achievable yield, while the minimum acceptable bound will be set by the plant efficiency.

For fusion technology to be commercially competitive, both the total capital cost required to generate the thermal power and the overall system efficiency of the plant must be competitive. If too large a fraction of the power must be recirculated to ignite the fusion reaction, the entire cost of the large thermal energy source must be recovered through the sale of the reduced net electric power, consequently the cost of electricity will be high.

To determine the dependence of overall system efficiency on the pellet gain and driver efficiencies, refer to the power balance diagrammed in Fig. 9-3. The gross electric power is allocated to the net electrical power

Fig. 9-2. Estimates of the pellet yield as a function of driver energy.

Fig. 9-3. The power balance of an inertial fusion power plant.

at the busbar (P_e), to the driver, and to the auxiliary needs of the plant. The thermal power generated (P_t), is the product of the driver input power, the driver efficiency (η_d), the pellet gain, and the effective energy multiplication of the neutron blanket, M . The auxiliary power requirement of the plant can be viewed as an effective reduction in thermal efficiency, when it is expressed as a fraction (f_a), of the thermal power. Because of the recirculated power, the overall system efficiency is less than the gross thermal-to-electric conversion efficiency by the following fraction:

$$\eta_s/\eta_t = 1 - f_a/\eta_t - 1/\eta_t MG\eta_d \quad (2)$$

Clearly, it is desirable to keep the auxiliary power fraction low, and to operate with high pellet gains and driver efficiencies. How high? The answer to this can best be determined by plotting Eq. (2) as a function of the fusion energy gain, $Q = G\eta_d$, which is the significant figure of merit for an inertial fusion system.

Actually, the ordinate of the graph drawn in Fig. 9-4 is the inverse of Eq. (2), since the cost of electricity is inversely proportional to

Fig. 9-4. Reduction of cost due to increasing fusion energy gain, as recirculated power is reduced.

the overall system efficiency. Thus, Fig. 9 4 shows how the relative cost of inertial fusion power drops as the fusion energy gain is increased, thereby reducing the recirculated power fraction. In Fig. 9-4, we have fixed the blanket energy multiplication and the auxiliary power needs of the reactor. An effective blanket gain of $M = 1.15$ due to neutron reactions in lithium, steel, etc., is reasonable. An auxiliary power fraction of $f_a = 0.03$ is also very typical. About 3% of the thermal power, or 7 to 8% of the electrical power, is needed in a reactor for liquid metal pumping of the primary and secondary loops, for vacuum pumping, boiler feed and condensate pumps, heating ventilation and cooling, tritium removal, and other in-plant services. A light-water fission reactor has comparable needs.

Two thermal-electric efficiencies are shown, $\eta_t = 39\%$ for a superheated steam system typical of a high temperature liquid-metal-cooled system ($T_{hot} = 500^\circ\text{C}$), and $\eta_t = 34\%$ for a saturated steam system typical of a water-cooled reactor ($T_{hot} = 350^\circ\text{C}$). A performance band is shown for systems with these two efficiencies. It would take very high fusion energy gains to asymptotically approach the cost of a system having only auxiliary power needs. There seems to be a knee in the curve in the vicinity of a fusion energy gain of $Q = 20$ to 50.

A fusion energy gain of 10 is needed to bring the cost of electricity to within 40% or so of the value at zero recirculated power. We consider a fusion energy gain of $Q = 20$ to be a more desirable goal for ICF, since it corresponds to the onset of diminishing returns between reducing cost and increasing the effective plant efficiency. In this example, a fusion energy gain of 20 corresponds to a total recirculated power fraction of 19%, bringing down the electricity cost to within 23% of the value at zero recirculated power, and within 15% of the cost asymptote represented by the auxiliary power needs. If thermal efficiencies are lower, such as the case of $\eta_t = 34\%$, higher energy gains are needed (perhaps $Q = 25$). For illustrative purposes, we therefore adopt a goal of $Q = \eta_d G = 20$ as desirable for inertial fusion to approach its potential.

Multiplying both sides of this criterion by the driver energy gives a relation for an effective lower limit on yield (Y), which is

dependent on the driver efficiency and energy.

$$Y \geq 20 E / \eta_d \quad (3)$$

Combining this minimum yield criterion with the bands of expected yields from Fig. 9-2, gives the allowed regions of yield of Fig. 9-5. The lower boundaries correspond to several different driver efficiencies. The upper boundary of the intermediate gain region is indicated as a realistic upper limit. The pulse repetition rate corresponding to each yield is given on the right hand ordinate for the special case of a single chamber 1-GW_e plant with a system efficiency of 32% and a blanket multiplication of $M = 1.15$. Slightly higher repetition rates would be needed for lower thermal efficiencies than the 39% assumed, and a 300-MW_e power plant would need only 30% of the indicated repetition rate.

The major conclusion we obtain from Fig. 9-5, is that by increasing pellet gain and/or increasing driver efficiency, we have greater freedom to choose a combination of driver, pellet, and chamber to minimize cost. If pellet gains in the intermediate gain band are eventually realized, and drivers having 10 to 15% efficiency can be developed, we will have a reasonable area of parameter space in which to work. If only 5% efficient drivers can be developed, quite large gains are required, near the optimistic gain boundary, resulting in relatively large yields (>500 MJ) and low repetition rates ($\lesssim 5$ Hz). Conversely, the development of high-efficiency drivers tends to open up the region of smaller yields and higher repetition rates and may allow the use of lower driver energies. If only conservative gains can ever be achieved, 10% efficient drivers will require 6-MJ energies, and 20% efficient drivers will require approximately 2 MJ per pulse.

The performance map of Fig. 9-5 does not yet include upper bounds on repetition rate. There may be a repetition rate limit (presently unknown, but somewhere above 10 Hz) due to uneconomical or infeasible rates of pellet production or delivery.

Chamber and Driver Considerations. For commercial applications, the reaction chamber must be designed to have an economic lifetime. Therefore, damage from each pulse must be kept very low. The

chamber must be surrounded by lithium containing regions (blankets), which are designed to collect and multiply fusion energy, and to breed tritium for the fuel cycle. (It is essential that a fusion economy be self-sufficient in tritium. This means that, for each fusion reaction, at least one atom of tritium must be produced by the nuclear transmutation of lithium.) Although the momentum of the pellet debris is small, large stresses can result from high rates of energy deposition in the blankets and structural materials. Energy deposition by soft x rays and debris occurs at, or very near, free surfaces of incidence in both structural and coolant materials. However, the kinetic energy of 14-MeV neutrons is deposited throughout relatively large material volumes.

The most challenging reactor design consideration is protecting the chamber wall from the various forms of rapid energy deposition. Diverse methods for providing

Fig. 9-5. Performance map of the economic operating parameters of a 1-GW_e power plant.

this protection have led to a number of reaction chamber types: dry-wall (includes both bare-walls and sacrificial-liners), magnetically protected walls, gas filled chambers, wetted walls, and thick liquid metal walls. Furthermore, since no one chamber concept is appropriate for the entire range of yields and repetition rates, or for all drivers, it is necessary to investigate many different concepts. An overview of the feasible regimes for some of these concepts is illustrated in Fig. 9-6. The ordinate used in Fig. 9-6 is the same as that used in Fig. 9-5. The chamber types can be matched to different drivers by comparing the chamber background density with the propagation windows for the three driver types (lasers, light ions, heavy ions). The abscissa of Fig. 9-6 is the conventionally used corresponding background pressure (at standard temperature).

Lasers are limited to operation in chamber environments with pressures of less than a few Torr, by breakdown and subsequent beam refraction.³ Depending on the background gas composition, in gases having pressures ranging from 5 to 50 Torr

(Ref. 4), light-ion beams can propagate through an ionized channel created by a laser beam. It is most likely that heavy ions could be used in the range of pressures below 10^{-3} to 10^{-4} Torr (Ref. 5). But, there is also some hope that a window exists between 10^{-1} and 1 Torr, where a pinched propagation mode might be possible.⁶

The performance map of Fig. 9-6 is an estimate of the physical constraints on chamber design, which are set by wall protection parameters, and the necessity to reestablish ambient chamber conditions. We assumed a 10-m upper limit on the chamber radius. The representations of boundaries in the map are useful only for classifying the operating parameters of each chamber type. Actually, they are extremely fuzzy and, in some cases, unknown. Only limited regions of this parameter space have been explored for each concept. The correct limits will not be known until a reactor pellet is specified, and experiments can be performed. Nevertheless, the map provides a useful overview of the applicability and overlap of the different chambers and drivers that are under consideration.

Dry-wall chambers are yield-limited, due to the effects of energy deposition by x rays and pellet debris, which might result in the evaporation and/or sputtering of material surfaces. To minimize structural damage in reactor concepts having cavity walls that are exposed to surface evaporation and sputtering, there are tradeoffs possible between relative x-ray and debris energy yields, and their energy spectra. These tradeoffs lead to different optimum fusion-pellet designs for different reactor cavity concepts. In general, heating and vaporization from x-ray deposition increases as x-ray energy decreases, due to range effects.

The effects of pellet output on sputtering erosion rates are more complex. These depend on the atomic number (Z), mass, velocity, angle of incidence of the pellet debris, and on the Z -number of the target-surface material.⁷ In general, erosion rates increase with increasing mass and energy yield. But, they can either increase or decrease with an increase in the pellet material Z -number, depending on the kinetic energy of the pellet debris upon incidence. Sputtering erosion decreases as the atomic number of the target surface material decreases. Sputtering erosion is important in

Fig. 9-6. Performance map showing the operating regime of yield and background pressure allowed by chamber design and driver propagation.

the design of bare-wall and sacrificial-liner first-wall concepts. Desirable properties for sacrificial liner materials used for protecting the chamber wall include:

- Low Z-number (sputtering yields decrease, and x-ray penetration depths increase, as the atomic number decreases).
- High thermal conductivity and heat capacity.
- High-temperature resistance (to maximize heat transfer, and minimize evaporation, during energy deposition).
- Low cost.
- The ease of fabrication.

These properties appear to be best satisfied by graphite, which has therefore been chosen for studies of sacrificial-liner concepts.

The approximate yield limits, as shown in Fig. 9-6 for steel and refractory-metal bare-wall chambers, and for carbon, are on the high side. However, they are typical of results obtained with heavier fusion pellets, such as those which might be used with heavy-ion beams.

Performance might also be enhanced by either adding gas to the chamber, or redirecting the ions by using a magnetic field. In the magnetic protection scheme, the ions are deflected from the cylindrical wall by a solenoidal magnetic field, and channeled toward end plugs designed for ease of replacement and refurbishment. In the gas-protection scheme, the x rays and ions are partially absorbed in the gas, and their energy is mainly reradiated back to the wall over a longer time. In each case, the flux of short-range radiation to the walls is reduced, allowing a greater yield or smaller radius for a given wall surface. Higher gas pressures will cause the absorption to occur in a centrally located fireball, which will drive a blast wave. The resultant overpressure on the wall is limited by the acceptable transient mechanical stress. The gas-filled chamber will also have a repetition rate limit (presently unknown) resulting from either the maximum rate at which the very hot gas can be pumped out between shots, or by the injection rate of cryogenic pellets into this gas.

If the wall surface is a liquid, it can ablate and be reestablished without damage. A wetted-wall chamber, having a porous transpiration-cooled surface, can take higher yields than a dry wall. But, the underlying

structure is still eventually limited by the pulsed mechanical stresses due to the reaction forces resulting from fast ablation.

The thick LMW concepts are designed to absorb the neutrons as well. The maximum acceptable pellet yield of this concept is limited, although at a very high level, by the pulsed mechanical stress from both liquid and gas impact on the wall. For both the wetted-wall and the LWM concepts, the background pressure is limited on the low end by the vapor pressure of liquid metals in the 400 to 500°C range, and on the upper end by aerosol formation, which would interfere with beam propagation.

The performance limits, as sketched here, will undoubtedly shift as we learn more from detailed calculations and experiments. Further innovations will almost certainly lead to new concepts with broader applicability. Yet, we already have the impression that the characteristic freedom of chamber design in inertial fusion gives the diversity needed for covering a wide operational range of yields and repetition rates with various drivers. Gas-filled chambers seem particularly appropriate for both light and heavy ion beams, when the gas is necessary to form propagation channels, and when higher driver efficiencies will allow lower gains and yields. All concepts, except gas filled chambers, might be appropriate for lasers and heavy-ion beams in the low-pressure range. If sufficiently high fusion energy gains can be reached with lasers, so that lower yields and high repetition rates are accessible, we might have a wide choice among liquid metal walls, wetted walls, magnetically protected or gas-protected walls.

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Parametric Studies

The Influence of Chamber Pulse Rate on Structural Stress and Lithium Flow.^{8,9} The HYLIFE¹⁰ reaction-chamber concept for an inertial confinement fusion (ICF) power plant uses a thick liquid metal wall (LMW) to absorb the energy of the short-range fusion products (soft x rays and pellet debris) and to moderate and attenuate the energy of the fusion neutrons (Fig. 9-7). The LMW is a

distributed array of free-falling lithium jets arranged so that the material ablated by the short-range fusion products can vent through the array. Thus, the outward liquid momentum and the impact stress in the first structural wall (FSW) are reduced.

The response of the LMW to fusion pulses is dependent on its geometry. In a jet array, the short-range fusion products vaporize the line-of-sight surfaces of the jets; the vapor moves into the center of the cavity, stagnates, and then expands outward. Meanwhile, the neutron energy pressurizes each jet, causing disassembly of the jets into high-velocity liquid drops moving at right angles to the original jet surfaces. The expanding gas moves through the droplets and accelerates them outward toward the FSW. The droplets coalesce into a condensed-phase slug having an outer radius just beyond the original LMW outer radius. The slug thickness remains constant while the density decreases in flight due to the diverging geometry. The reduced slug density at impact results in a lower impact pressure, an increased impact duration, and a reduced impact stress, as compared to the impact of normal-density liquid.

If the pressure distribution is smeared to a uniform level through the liquid, we can find the wall stress analytically^{8,10} by

$$\sigma_{\max} = 2 \left(\frac{R\hat{w}}{\delta} \right) \left(\frac{\eta}{1-\eta} \right) \rho_0 \nu \quad (4)$$

where $R\hat{w}/\delta$ is the ratio of wall radius to wall thickness, ρ_0 is the normal liquid density, η is the ratio of expanded density to normal density, and ν is the impact velocity. The impact velocity ν

$$\nu = \alpha \left(\frac{\Gamma}{\rho_0 c} \right) \left[\frac{f [1 - \exp(-\mu \ell)]}{2\pi R_0 \ell (2R_0 f + \ell)} \right] E_n \quad (5)$$

where Γ is the Grüneisen constant; c is the liquid sound speed; f is the liquid-jet packing fraction; μ is the neutron-energy attenuation coefficient; ℓ is the effective LMW thickness to neutrons; R_0 is the LMW inner radius; E_n is the fusion neutron energy, including exothermic interactions in the LMW; and α is a normalization constant.

The normalization constant, $\alpha = 0.28$, is determined from a one-dimensional

Fig. 9-7. Cross section of the HYLIFE reaction chamber at the horizontal midplane.

multiple-zone calculation that considers both jet disassembly induced by neutrons and velocity imparted to the liquid by the expanding gas.^{11,12} Neither the simplified single-zone model nor the multiple-zone normalization calculation consider stress due to impact of the expanding gas on the structure. This additional load may be offset by the axial dissipation of the gas during the blowoff and expansion phases.

The material used for the FSW in this study is 2-1/4 Cr-1 Mo ferritic steel. The physical properties of this steel as a function of temperature were obtained from Ref. 13. The high-cycle fatigue-strength data¹⁴ were furnished by Booker¹⁵ and have been submitted for inclusion in ASME Code Case N-47. The material properties used do not include the effects of radiation damage or liquid-metal corrosion.

Some of the fusion energy penetrates the liquid-metal array and is deposited in the FSW. This energy raises the wall temperature, and the spatial temperature distribution results in a thermally induced wall stress. We calculated the temperatures and resultant stress distributions as functions of the radii and thicknesses of the wall and the array.¹⁶ Material physical properties were evaluated at the peak calculated temperature.

The FSW is subjected to a steady stress and to four dynamic stresses. The steady stress is a thermal stress arising from the temperature profile in the wall. The dynamic stresses include a thermal stress from the sudden deposition of energy in the wall by each fusion pulse; a change in the steady thermal stress due to startup, creep, and shutdown; an impact stress from the LMW after each fusion pulse; and an impact stress from the expanding gas. The transient stress from each fusion event is small compared to the steady thermal stress, and the transient thermal stress from startup, creep, and shutdown will occur only about 100 times during the lifetime of the reaction chamber. The transient effect of a misfire on the thermal stress is also expected to be small.

Leaving out the transient stresses that are expected to be small, we are left with three stresses of concern: the steady thermal stress, and the dynamic liquid and gas impact stresses. The steady thermal stress is a function of the fusion power, while the

dynamic liquid impact stress is a function of fusion yield. The gas impact stress is not included in this analysis.

There are many accepted models for designing with combined stresses, including the Goodman, Modified Goodman, and Soderberg models.^{17,18} We use an ASME-approved model,¹⁹⁻²¹

$$\sigma_1 \leq \frac{S_f}{2} \quad (6)$$

and

$$\sigma_1 + \sigma_t \leq S_y \quad (7)$$

where S_y is the yield strength of the structural material and S_f is the fatigue strength at 10^{10} load cycles (10^9 fusion pulses with 10 rings per pulse). Both properties are evaluated at the structural-wall centerline (peak) temperature.

We have developed a methodology for determining the design-point space of FSW thickness and jet-array inner radius for a chamber of fixed yield, pulse-repetition frequency, and array packing fraction. This methodology is shown graphically in Fig. 9-8.

The impact stress in the FSW is shown as a function of thickness in Fig. 9-8(a) for a fixed inner radius (0.75 m) of the array. Thermal stress, fatigue strength, and yield strength of the FSW are also plotted. The minimum structure thickness that satisfies the impact stress criterion of Eq. (6) is 23 mm, and the maximum thickness that satisfies the total stress criterion of Eq. (7) is 55 mm. Thus, any FSW thickness between 23 and 55 mm satisfies both design criteria.

When the calculations shown in Fig. 9-8(a) are repeated using other inner radii for the array, the results form a Venn diagram as shown in Fig. 9-8(b). The sets of allowed LMW inner radii and FSW thicknesses that satisfy each structural design criterion are shown for a wall radius of 5 m, an effective jet-array thickness of 1 m, and 1-Hz fusion pulses of 2700 MJ. The intersection of the two sets is the set of allowable design points.

The upper bound of the FSW thickness is limited both by loss of mechanical strength at high temperature and by high thermal stress. The lower bound is due to the increase in impact stress with decreasing FSW

Accomplishments Associated with the HYLIFE Chamber Concept

thickness. The lower bound on the array inner radius is set by the high impact velocity of the liquid, and the upper bound is set by the lack of liquid cavitation when the array is too close to the wall. The design point selected is the minimum inner radius of the jet array that satisfies both stress criteria, because this point minimizes the required flow rate of liquid metal.

The loci of design points as a function of the effective array thickness for a given structure radius can be developed by a series of Venn diagrams, as shown in Fig. 9-8(c). The family of design points in the space of LMW inner radius and FSW thickness is shown in Fig. 9-8(d) as a function of LMW effective thickness and FSW radius. As either the FSW radius or the effective array

thickness is decreased while the other variable is held constant, the design-point FSW thickness decreases because the increased neutron-energy density in the structure increases the structure temperature and thermal stress. Likewise, the design-point array radius increases to reduce the array energy density so that the impact-stress criterion is satisfied.

The loci of design points shown in Fig. 9-8(d) is plotted again in Fig. 9-9, at a larger scale, and with reversed axes. The minimum mass of structural material for a given chamber radius occurs at the minimum design-point FSW thickness, which is achieved by operating at the radiation damage limit of the FSW. At a radiation damage limit of 165 displacement per atoms

Fig. 9-8. Construction of ASME design-point parameter space.

(dpa), the design-point FSW thicknesses and array radii are relatively invariant with the FSW radius.

Figure 9-10 shows the normalized LMW flow rate for constant midplane velocity at the design points that minimize the array inner radius and satisfy both ASME design criteria. The combination of array equivalent thickness and FSW radii that result in a radiation damage limit of 165 dpa in the FSW are also shown. The minimum flow rates are at the radiation damage limit array thickness for a given FSW radius; at this limit, the array flow area decreases with increasing FSW radius. Operating the chamber with the minimum allowable array thickness, however, increases the fusion energy deposited in the graphite reflector; thus, the flow rate through the reflector must be increased, and the total system flow rate will not decrease as markedly as shown in Fig. 9-10 (Ref. 22).

The effect of LMW injection velocity on the flow rate, packing fraction, and inner radius of the array is as shown in Fig. 9-11 for a structure radius of 5 m and for a radiation-damage-limited effective array thickness (0.74 m). A chamber height to diameter ratio of 1.6 is used, giving an 8-m chamber height. Very low injection velocities result in a low packing fraction, a large inner radius, and a high flow rate for the array. As the injection velocity increases, the flow rate goes through a minimum. The optimum injection velocity for the array, which minimizes the array flow rate, occurs between 6 and 8 m/s for the 1-Hz system. At repetition rates above 1.5 Hz, the minimum clearing velocity is larger than the velocity of the flow-rate minimum.

For a fixed fusion-power system, a HYLIFE design with a single chamber and a low pulse-repetition rate is appropriate for high-gain pellets and a reasonable driver output energy (<5 MJ). If high-gain pellets are not available, a multichamber design using a single driver may be desirable. This multichamber system would more fully utilize the repetition-rate capability of the driver.

At a fixed output power, increasing the pulse-repetition frequency reduces the fusion-pulse energy. On the bases of the scaling laws and design criteria described above, the flow rate, pumping power, inner

Fig. 9-9. Loci of design points for FSW radii of 5, 6, 7, and 8 m.

Fig. 9-10. LMW flow rate vs array inner radius.

Fig. 9-11. LMW conditions as a function of injection velocity for one reactor geometry and pulse rate.

Fig. 9-12. LMW flow rate and pumping power as functions of chamber radii for single- and multichamber power plants.

radius, injection velocity, and midplane packing fraction of the LMW were calculated as a function of driver repetition rate. The LMW flow rate and pumping power for a 100% efficient pump are shown in Fig. 9-12 as a function of the pulse

repetition rate for several chamber radii. The effective LMW thickness was set by the radiation damage limit of 165 dpa, and the LMW was designed to minimize the flow rate.

For a single-chamber system, the minimum LMW flow rates occur between 1.5 and 2.0 Hz (1800- to 1350-MJ yield) for chamber radii between 4 and 6 m. While the minimum LMW flow rate increases with a decrease in the chamber size, the power required to pump the LMW at the minimum flow rate is virtually invariant with chamber size.

For low yields (<700 MJ) and higher driver repetition rates (>3 Hz), a four-chamber plant was used. The minimum LMW flow rate again occurs at repetition rates between 1.5 and 2.0 Hz per chamber (450- to 340-MJ yield). The minimum LMW flow-rate and pumping power for the four-chamber plant is within 10% of that for the single-chamber configuration with the same 4-m chamber radius. The LMW pumping power and minimum flow rate both increase by about 50% with a reduction in the chamber radius from 4 to 3 m in the four-chamber configuration.

For the cases investigated, the following observations can be made for the system operating at the minimum flow rate:

- The injection velocity is that required for the LMW to reestablish its integrity.
- An FSW thickness between 30 and 40 mm meets the design criteria.
- The midplane packing fraction is between 50 and 60%.
- The inner radius and thickness of the LMW are dependent on the chamber radius and the fusion power.

For a given fusion power, a multichamber plant can be compared with a single-chamber plant on the basis of the results shown in Fig. 9-12. At the driver repetition rate that gives the minimum LMW flow rate in each case, the 5-m radius single-chamber plant has roughly the same flow rate and pumping power as the 4-m-radius four-chamber plant. In the latter case, the number of pumps and the pumping power remain constant, but the four slightly smaller chambers require a considerably larger building, a more complicated driver-beam transport system, and a quadrupling of the driver recirculating power. If the chamber radius of the four-chamber design is reduced

to 3 m to reduce the building size, the number of pumps and the pumping power increases, and the additional area required for the additional pumps offsets the area reduction realized by using smaller chambers.

The results given above are tentative because the material properties do not include the debilitating effects of radiation damage and corrosion and because the scaling laws are simplified and normalized to a single calculation with a multizone model. The normalization calculation also contains many uncertainties, including lithium opacities, effects of axial dispersion, the difference between jet and annuli disassembly, and the effect of the gas impact on the FSW. Nevertheless, the results do allow us to examine trends and to make a more intelligent choice of a design point for further analysis.

The array in the HYLIFE concept should be designed so that the FSW operates at the radiation damage limit (e.g., 0.75-m effective LMW thickness for a 5-m wall radius). The wall thickness should be between 35 and 40 mm, and, for a 1.5-Hz repetition rate, the jet array should be placed about 0.5 m from the pellet. A saving in mass flow rate is possible if the wall radius is increased, but the higher injection velocity, the higher chamber and building cost, and the greater amount of radioactive waste all point toward selection of a reasonably small wall radius. The liquid-metal flow can also be reduced by redesigning the graphite blanket-shield region between the FSW and the vacuum vessel (as discussed in "Blanket-Shield Design Options," below). By increasing the driver repetition rate to 6 Hz, divided among four chambers, a lower-yield (~450 MJ), lower-gain (<200) system can be designed.

When similar scaling relationships for target and driver costs are included in our calculations, the effects of repetition rate and yield on the cost of electricity can be determined (see "Allowable Driver and Target Factory Cost," below).

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Blanket-Shield Design Options. The function of the blanket-shield region in the HYLIFE reaction chamber is to reduce the number and energy fluences of the neutrons that leak from the vacuum vessel. This reduction in

neutron-energy leakage is desirable because it will maximize the amount of energy converted to useful thermal power and it will minimize the formation of radionuclides in the region of the primary containment building surrounding the vacuum vessel.

To limit neutron leakage in the current HYLIFE reaction-chamber designs,^{23,24} we use a blanket-shield region comprising a heterogeneous composite of graphite and liquid lithium. The neutron energy is moderated in both the graphite and lithium, while the number fluence is reduced via ${}^6\text{Li}(n,\alpha)\text{T}$ reactions. In this article, we explore several design options for the blanket-shield region to

- Increase the energy density in the blanket-shield.
- Decrease the fluence of neutrons, and decrease their energy leakage from the vacuum vessel.
- Reduce the amount of circulating lithium in the primary coolant system.

These options, if implemented, may increase the thermal-energy conversion efficiency of the system and decrease the capital cost of the primary coolant pumps.

The functions of lithium in the blanket-shield region of the HYLIFE reaction chamber are to

- Cool the first structural wall.
- Cool the graphite blanket.
- Shield the vacuum vessel and exterior from neutrons.
- Cool the vacuum vessel.
- Protect the first structural wall from erosion.

The most important function of the lithium is to serve as a heat-transport medium, since cavitation of the liquid-lithium layer by x-ray energy deposition may cause erosion of the FSW, even though the lithium is intended to protect the structure from erosion. We can use other materials as a substitute for lithium to shield the vacuum vessel and exterior from neutrons. Meier²⁵ proposed substituting boron carbide for graphite in the reaction-chamber reflector as one option for reducing the tritium-breeding ratio in a HYLIFE reaction chamber.

Some relevant properties of lithium and boron carbide are compared in Table 9-1. The absorption of a neutron by ${}^6\text{Li}$ will release about 70% more energy than the absorption of a neutron by ${}^{10}\text{B}$. However, higher concentrations of the absorbing

isotope can be obtained with boron carbide than with natural lithium, although pure ^6Li gives a greater concentration of neutron absorber per unit mass than either boron carbide or natural lithium.

A measure of the probability of a neutron-absorption event is the macroscopic cross section of the absorbing material. This measure takes into consideration both the microscopic cross section and the atom concentration of the absorbing material. The macroscopic cross sections for n, α reactions for boron carbide and natural lithium were determined²⁶ as a function of neutron energy based on the cross section curves of Evaluated Neutron Data Library ENDL.²⁷

The probability of an exothermic n, α reaction in a given volume of absorber material is proportional to the product of the neutron fluence at a given neutron energy and the macroscopic cross section of the absorber for that neutron energy. Since the range of α particles is small compared to the range of energetic neutrons, the α particle energy can be assumed to be deposited at the site of the n, α reaction. Thus, the probability of the exothermic energy being deposited in a given volume of absorber is simply the exothermic energy released per reaction multiplied by the probability of a reaction taking place in the volume. In Fig. 9-13, we show the probability of the exothermic n, α energy deposition per unit fluence in natural and enriched boron carbide and in natural and enriched lithium, all as functions of neutron energy. We see that the highest absorbed-energy density from the n, α reaction generally occurs with enriched boron carbide. Enriched lithium and natural boron carbide have essentially

the same energy densities in the v^{-1} region of the cross section, while enriched lithium clearly has a higher absorbed-energy density in the resonance region of the cross section. Natural lithium has a lower absorbed-energy density than natural boron carbide, even though there is about 70% more energy released per n, α reaction in the lithium than in the boron carbide.

From Fig. 9-13, we can see that the softer the neutron spectrum, the more effective lithium or boron carbide become as neutron shields. For HYLIFE, we can soften the spectrum either by increasing the LMW thickness, with an undesirable increase in the required lithium flow rate, or by increasing the thickness of the graphite blanket. By substituting natural boron carbide for the existing lithium neutron shield behind the blanket, we can reduce both the inventory and flow rate of lithium in the HYLIFE reaction chamber. In addition, our use of boron carbide will result both in higher absorbed-energy densities in the blanket-shield region and in an increased system-energy release per fusion event. However, we can determine the magnitude of the reduced inventory and flow rate of lithium, and increased energy density and energy release to be realized by using boron carbide rather than lithium only if we know the neutron energy spectrum and if we consider other high-energy reactions in the absorber (e.g., n, p). We can better determine the effect of substituting boron carbide for lithium in the blanket-shield region by performing complete neutronic calculations.

We believe that boron carbide should be used as the neutron shield to protect the vacuum vessel in the HYLIFE reaction chamber. The substitution of boron carbide for the lithium in the shield will

- Decrease the lithium inventory.
- Decrease the lithium flow rate.
- Decrease the radioisotope production in the vacuum vessel.
- Increase the energy density in the neutron shield.

The total mass flow rate in the HYLIFE reaction chamber is dependent on the flow rate in the array and on the flow rate in the blanket-shield region.²⁸ The minimum flow rate of the jet array occurs when the FSW operates at the radiation damage limit, the magnitude of which is dependent on the

Table 9-1. A comparison of lithium and boron carbide properties.

Absorber material	Lithium	Boron carbide
Density (g/cm^3) ^a	0.5	2.5
Absorbing isotope	^6Li	^{10}B
Absorbing reaction	$^6\text{Li}(n, \alpha)\text{T}$	$^{10}\text{B}(n, \alpha)^7\text{Li}$
Energy release (MeV/absorption)	+4.8	+2.8
Isotopic concentration ($^a\text{X}/\text{X}$)%	100 ^b	7.42 ^c
19.78 ^c 10^{23} atoms/g	1.00	0.074
0.093 10^{23} atoms/cc	0.50	0.037
		1.16
		0.23

^aAssumed to be invariant with isotopic composition.

^bEnriched.

^cNatural.

injection velocity, the pulse-repetition frequency, and the yield of fusion events for a fixed fusion power.

The lithium flow rate in the blanket-shield region depends on the

- Functions assigned to the lithium in the blanket-shield region.
- Relative energies deposited in the array and in the blanket-shield region.
- Temperature restrictions placed on the system.
- Fluid-flow circuit in the reaction chamber.

We postulate here that the only function of the lithium in the blanket-shield region is to remove the energy generated in the region. The relative energies deposited in the array and in the blanket-shield region are functions of the array thickness. The maximum temperature rise in the lithium circuit is limited by corrosion considerations to a maximum exit temperature of 773 K, and the minimum inlet temperature is 450 K, the melting point of lithium.

We can identify four flow circuits through the reaction chamber. For the two models shown in Fig. 9-14, we assume that the flow through the blanket-shield region is parallel to that in the array. For the model shown in Fig. 9-14(a), we adjust the flow through both regions such that the temperature rise of the lithium through the blanket-shield region is the same as that in the array. In Fig. 9-14(b), we also use parallel flow between the blanket-shield region and the array. The flow through the former is adjusted such that the mixed-mean exit temperature of the liquid at the discharge end of the converter can be adjusted for, say, corrosion considerations. Correspondingly, since the array is in contact only briefly with the structure, the array temperature can rise to a value above the corrosion limit, while the temperature of the liquid at the bottom of the chamber is held at the corrosion limit by mixing with the cooler liquid from the blanket-shield region.

Two cooling circuits in a series arrangement are shown in Fig. 9-15. The model shown in Fig. 9-15(a) has the array operating at a lower temperature, while part of the array-circulation flow is diverted to be heated to higher temperatures in the blanket-shield region. A portion of the blanket-shield flow is then circulated through the primary heat exchanger and returned to the array circulation system. Although this scheme

has a more complex fluid circuit than the parallel arrangement, the series arrangement allows us to achieve the low chamber temperature and low vapor pressure needed for heavy-ion beam propagation, while retaining the higher thermal efficiency associated with the higher intermediate heat exchanger (IHX) inlet temperature. In addition, the series scheme allows the total flow rate of the chamber to be equal to that in the array.

A fourth model, shown in Fig. 9-15(b), uses the sodium in the intermediate loop to cool the blanket-shield region. Since the array energy is transferred to the sodium before the sodium enters the blanket-shield region, the energy transferred to the sodium

Fig. 9-13. Probability of exothermic (n, α) energy deposition in enriched and natural boron carbide and in enriched and natural lithium.

Accomplishments Associated with the HYLIFE Chamber Concept

from the blanket-shield region “tops off” the intermediate heat-transfer loop. We can operate the intermediate loop at a higher temperature than with the IHX alone, and, thus, we can increase the thermal efficiency of the system. Alternatively, we can reduce the lithium array temperature, and with it the chamber pressure needed for ion-beam propagation, and use the heat dumped into the sodium from the blanket-shield region to maintain the system thermal-conversion efficiency. While this model has the potential for decreasing the volume of lithium in the system, for increasing the system thermal-conversion efficiency (or decreasing the chamber pressure), and for decreasing the lithium flow rate, the price is a more complex intermediate loop that is activated from $^{23}\text{Na}(n,\gamma)^{24}\text{Na}$ reactions, for which the gamma decay energies are 1.38 and 2.75 MeV, respectively, and the half life is nearly 15 h.

The proposed models are compared in two parts. The first comparison, presented in Table 9-2, is of the HYLIFE,²³ modified-HYLIFE,²⁴ and parallel-coolant-flow models. In the comparisons, we set the maximum lithium temperature at 773 K, based on corrosion considerations. The

chamber pressure is based on the vapor pressure of lithium at the inlet temperature. The array packing fraction is 0.5 at the chamber midplane.

From Table 9-2, we draw the following conclusions:

- The HYLIFE²³ and modified HYLIFE²⁴ use higher flow rates than necessary to cool the system.
- Thinning the array to the radiation damage limit (1-B vs 1-A) results in a less dramatic savings in total system lithium flow rate than in the array flow rate²⁹ due to the increased energy deposited in the blanket-shield region; a higher temperature rise in the array due to the increased average energy density in the array; and a reduction in total power due to reduced $^6\text{Li}(n,T)\alpha$ reactions.
- Increasing the radius of the FSW (1-C) and operating at the radiation damage limit results in a less dramatic savings in total system lithium flow rate than in the array flow rate due to the increased energy deposited in the blanket-shield region; a higher temperature rise in the array due to the increased average energy density in the array; a lower chamber pressure due to the lower inlet temperature for the fixed exit

Fig. 9-14. Parallel-flow models for lithium in HYLIFE reaction chamber.

System balance

$$\begin{aligned}
 E_a &= \dot{m}_a c (T_2 - T_1) \\
 E_b &= \dot{m}_b c (T_2 - T_1) \\
 (E_a + E_b)_y &= X(\dot{m}_a + \dot{m}_b) (T_2 - T_c) \\
 &= (\dot{m}_a + \dot{m}_b) c (T_2 - T_1) \\
 \dot{m} &= \dot{m}_a + \dot{m}_b \\
 T_c &> 450 \text{ K} \\
 T_2 &< 773 \text{ K}
 \end{aligned}$$

System balance

$$\begin{aligned}
 E_a &= \dot{m}_a c (T_{2a} - T_1) \\
 E_b &= \dot{m}_b c (T_{2b} - T_1) \\
 E_a + E_b &= (\dot{m}_a + \dot{m}_b) c (T_2 - T_1) = X(\dot{m}_a + \dot{m}_b) c (T_2 - T_c) \\
 \dot{m} &= \dot{m}_a + \dot{m}_b \\
 T_2, T_{2b} &< 773 \text{ K} \\
 T_c &> 450 \text{ K}
 \end{aligned}$$

Accomplishments Associated with the HYLIFE Chamber Concept

	HYLIFE ^{1,2}	Modified ² HYLIFE	Parallel system				
			1-A	1-B	1-C	2-A	2-B
FSW radius (m)	5.0	5.0	5.0	5.0	8.0	5.0	5.0
Array inlet velocity (m/s)	4.4	5.5	7.1	7.1	7.1	7.1	7.1
Nozzle-plate packing fraction	—	0.95	0.80	0.80	0.80	0.80	0.80
Midplane array equivalent thickness (m)	1.0	1.0	1.0 ^a	0.74 ^{a,b}	0.43 ^{a,b}	0.74 ^{a,b}	0.74 ^{a,b}
Power (MW)							
Total	3190	3190	3190	3140	3020	3140	3140
Array	2980	2980	2980	2770	2310	2770	2770
Temperature (K)							
Inlet	761	762	760	755	743	762	750
ΔT array	12	11	13	18	30	18	18
ΔT blanket-shield	12	5	13	18	30	4	23
Exit mixed-mean	773	773	773	773	773	773	768
Vapor pressure (Torr)	2×10 ⁻³	2×10 ⁻³	2×10 ⁻³	2×10 ⁻³	1×10 ⁻³	2×10 ⁻³	2×10 ⁻³
Flow rate (m ³ /s)							
Array	93	130	111	79	38	79	79
Blanket-shield	36	18	8	11	12	79	1
Total	129	148	119	90	50	158	80

Table 9-2. Comparison of parallel flow systems with HYLIFE and modified-HYLIFE models.

Fig. 9-15. Series-flow models for coolant in HYLIFE reaction chamber.▼

(a) Lithium-cooled blanket-shield region

(b) Sodium-cooled blanket-shield region

System balance

$$\begin{aligned}
 T_2 &= T_1 + \alpha \\
 X(T_H - T_2) &= X T_H - (T_1 + \alpha) = \beta \\
 yX(T_H - T_c) &= \alpha + \beta \\
 y &\leq 1 \\
 X &\leq 1 \\
 \alpha &= \frac{E_a}{\dot{m}c} ; \beta = \frac{E_b}{\dot{m}c} \\
 T_H &< 773 \text{ K} \\
 T_c &> 450 \text{ K}
 \end{aligned}$$

System balance

$$\begin{aligned}
 &\text{Lithium} \\
 &\dot{m}_a c_{Li} (T_2 - T_1) = E_a \\
 &X \dot{m}_a c_{Li} (T_2 - T_c) = E_a \quad X \leq 1 \\
 &\text{Sodium} \\
 &\dot{m}_b c_{Na} (T_{Ni} - T_{Nc}) = E_a \\
 &\dot{m}_b c_{Na} (T_{NH} - T_{Ni}) = E_b \\
 &\text{Coupling} \\
 &T_{Ni} = T_2 - 18 \\
 &T_{Nc} = T_c - 25 \\
 &T_c > 450 \text{ K} \\
 &T_2 < 773 \text{ K}
 \end{aligned}$$

Accomplishments Associated with the HYLIFE Chamber Concept

temperature; a reduction in total power due to reduced ${}^6\text{Li}(n,T)\alpha$ reactions; and a lower fusion-power density.

- Operating with a large array temperature rise and a small blanket-shield lithium temperature rise (case 2-A) to keep the mixed-mean exit temperature to a specified value results in high system flow rates.
- Operating with a low array-temperature rise and a high blanket-shield temperature rise (case 2-B) reduces the flow rate; however, this reduced flow rate is achieved at the price of reducing the system Carnot thermal conversion efficiency due to the lower mixed-mean exit temperature set by the maximum temperature in the blanket-shield region.

The parallel-flow example (case 1-B), with identical fluid-temperature rises in the array and in the blanket-shield region, and operating with the array thickness set by the radiation damage limit, will be used as a comparison for the investigation of other blanket-shield models, as shown in Table 9-3.

When we compare the series-flow model with the parallel-flow model in case 1-B, we see that a substantial reduction in lithium flow rate and chamber vapor pressure can be achieved with the series flow. The intermediate-loop sodium system for the

parallel-flow system and the lithium series flow system through the blanket-shield region are based on the Bechtel work.³⁰

For the series flow of lithium through the array and the blanket-shield region, the lower limit on the array inlet temperature is based on the requirement that the minimum temperature of the lithium exiting from the IHX must be at or above the melting temperature of lithium; consequently, the lithium flow rate through the IHX must be less than or equal to that through the blanket-shield region. For a minimum lithium temperature of 47°C above the melting point of lithium, and for a maximum lithium temperature of 773 K due to the corrosion consideration, the lowest chamber pressure is 7×10^{-4} Torr, as given by case 3-A. Reducing the lithium temperature at the chamber inlet to achieve a lower chamber pressure will result in a corresponding drop in the peak lithium temperature (case 3-B), thereby resulting in a drop in the Carnot thermal-conversion efficiency of the system.

When we use sodium as the blanket-shield region coolant, we can achieve several orders-of-magnitude reduction in the chamber vapor pressure (case 4-A) with an increase in the system Carnot thermal-conversion efficiency because the sodium IHX exit temperature is 773 K, compared to 755 K in the lithium-cooled blanket. The

Table 9-3. Comparison of series flow system with a parallel-flow system.

	Parallel 1-B	3-A	Series 3-B	4-A
Temperature (K)				
Lithium				
Inlet	755	723	620	620
ΔT array	18	18	18	18
ΔT blanket-shield	18	32	18	—
Exit to IHX	773	773	656	638
Exit from IHX	613	500	500	500
Sodium				
Exit from IHX	755	755	638	620
Entering IHX	588	475	475	475
Peak	755	755	638	773
Lithium vapor pressure (Torr)	2×10^{-3}	7×10^{-4}	1×10^{-5}	1×10^{-5}
Lithium flow rate (m^3/s)				
Array	79	79	79	79
Blanket-shield	11	5.8	10	—
IHX	9.8	5.8	10	10
Total	90	79	79	79
Sodium flow rate (mg/s)	15	9.0	15	11

Note: For fusion power of 2700 MW; maximum lithium temperature = 773 K; array packing fraction = 0.5 at midplane; array equivalent thickness = 0.74 m at midplane; ASME design point at radiation damage limit; inlet array velocity = 7.1 m/s; FSW radius = 5.0 m.

price we pay, as mentioned above, is increased radionuclide production in the blanket and a more complex piping system. In addition, the increase in the Carnot thermal-conversion efficiency may be offset by real considerations of pressure-drop losses.

Monsler³¹ has noted that there are additional variations of the design options presented above, including using the array energy to drive a saturated steam cycle and using the blanket-shield-region energy to drive a potassium-vapor topping cycle. These variations are not analyzed here but should be investigated.

We conclude that the current HYLIFE designs are off-optimum in the blanket-shield region. We can improve the design by

- Using boron carbide as a shield material.
- Using lithium only as a coolant.
- Changing the coolant-flow circuit.

We anticipate that these changes will increase the power density in the blanket-shield region and decrease the required number of primary coolant pumps.

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The Effect of Chamber Ports on Neutron Leakage and Blanket Performance. A high negative current at the inner surface of the first structural wall (FSW) has been noted in neutronic calculations for ICF chambers.³² If chamber penetrations, such as beam ports and vacuum-pumping ports, are not included in the neutronics calculation, the results will overestimate key system-performance parameters, even if the results are modified by the fraction of the coverage that is lost. The reason for this overestimation is that the neutron leakage through a hole in the chamber will be higher than one would estimate based only on its solid angle, since neutrons reentering the fusion chamber from the blanket will have a finite probability of escaping through the hole.

As a result of the higher neutron leakage, system-performance parameters such as the tritium-breeding ratio and the fissile-fuel production rate in hybrid blankets will be degraded by more than the loss represented by the fractional solid angle. The additional neutron loss will be most apparent for weakly moderating blankets with high neutron multiplication, and the effect will be

most critical for systems with marginal performance, such as those with a tritium-breeding ratio close to 1.0.

The effects of inhomogeneities such as chamber penetrations point out the importance of doing two-dimensional calculations when major portions of the blanket perform different functions. For example, the results of a two-dimensional calculation of a cylindrical hybrid chamber with fissile breeding only in the radial zones may be significantly different than the solid-angle-averaged results of two one-dimensional problems. Preliminary calculations with TARTNP, a coupled neutron-photon Monte Carlo transport code, have been made to determine neutron leakage and blanket performance as a function of the size of chamber penetrations.

The problem geometry is shown in Fig. 9-16 and is described in Tables 9-4 and 9-5. The fusion chamber is spherical, with two diametrically opposed loss cones through the blanket. The blanket consists of a 3-cm thick ferritic steel wall at 5.0 m followed by a 70-cm-thick lithium region, a 17-cm-thick graphite reflector, and a 5-cm-thick steel vessel wall. A point-source of 14.1-MeV neutrons is located at the center of the evacuated chamber. There are no target-attenuation effects.

Fig. 9-16. Neutronic model showing loss cones and blanket geometry.

Accomplishments Associated with the HYLIFE Chamber Concept

TARTNP tallies the number of neutrons entering zones 1, 2, 3, 4, 9, and 10. Neutrons can enter a zone across any boundary of the zone. (An exception is that neutrons cannot cross a boundary to enter a zone from the leakage zones—zones 9 and 10—since neutrons cannot leave leakage zones.)

The resulting neutron tallies for both the number of 14-MeV neutrons and the total number of neutrons entering the indicated zones are tabulated in Table 9-6. Results are given for total loss-cone fractions of 0.02, 0.04, 0.10, and 0.20. For example, with two 1% loss cones, 0.019 14-MeV neutrons go straight out the holes, passing through zones 2 and 4 and ending up in zone 9. The remaining 0.981 neutrons pass through zone 3 and enter the blanket. (With perfect statistics, these would have been 0.020 and 0.980, respectively.)

Table 9-4. Zone description.

Zone	Description	Material	R ₁ (m)	R ₂ (m)	ΔR (m)
1	Vacuum chamber	None	—	4.99	4.99
2	Entrance to loss cone	None	4.99	5.00	0.01
3	Test zone	None	4.99	5.00	0.01
4	Blanket loss cone	None	5.00	6.05	1.05
5	Chamber wall	Ferritic steel	5.00	5.03	0.03
6	Breeding blanket	Lithium	5.03	5.73	0.70
7	Reflector	Graphite	5.73	6.00	0.27
8	Vessel wall	Ferritic steel	6.00	6.05	0.05
9	Loss cone leakage	None	6.05	—	—
10	Blanket leakage	None	6.05	—	—

Table 9-5. Chamber materials. ▶

Material	Isotopic composition	Density (g/cm ³)
Lithium	7.42% ⁶ Li 92.58% ⁷ Li	0.485
Ferritic steel (2-1/4 Cr-1 Mo)	96.994% Fe-natural 2.423% Cr-natural 0.583% Mo-natural	7.86
Graphite	100.0% ¹² C	1.70

Table 9-6. Neutrons entering a zone, per source neutron.

Zone	Loss-cone fractions							
	0.02		0.04		0.10		0.20	
	Two-1% loss cones 14-MeV group	Total	Two-2% loss cones 14-MeV group	Total	Two-5% loss cones 14-MeV group	Total	Two-10% loss cones 14-MeV group	Total
1	0	1.322	0	1.248	0	1.076	0	0.836
2	0.019	0.052	0.041	0.100	0.103	0.224	0.198	0.382
3	0.981	3.600	0.959	3.404	0.897	2.934	0.802	2.295
4	0.019	0.073	0.041	0.129	0.103	0.261	0.198	0.425
9	0.019	0.043	0.041	0.092	0.103	0.218	0.198	0.386
10	0+	0.057	0+	0.055	0.001	0.052	0+	0.042

Note: 14.1-MeV point source—no target attenuation.

The total neutron leakage from the holes is 0.043 in the case cited above. This leakage exceeds the fractional solid-angle loss by over a factor of 2 because many secondary neutrons (i.e., those from ⁷Li(n,n'T)α reactions) and scattered neutrons reenter the vacuum chamber. The total number of neutrons reentering the chamber is the tally for zone 1. For the case of two 1%-loss cones, the result is 1.322 neutrons per source neutron. Note that it is possible for the same neutron to enter the chamber more than once.

A neutron must pass through zone 2 or zone 3 to enter zone 1. Since zone 1 is a vacuum zone, any neutron entering it must also leave, passing through zone 2 or zone 3 once again. Therefore, tallies for zones 1, 2, and 3 are related by

$$N_1 = \frac{1}{2}(N_2 + N_3 - 1) \quad (8)$$

The minus one accounts for the source neutron that must enter zone 2 or 3 to begin with.

The neutron leakage as a function of the loss cone fraction of the total solid angle is shown in Fig. 9-17. We would expect the total leakage per source neutron to be equal to the sum of the fractional solid-angle loss (to account for direct leakage of 14.1-MeV neutrons) plus a term that is a function of both the number of neutrons reentering the chamber from the blanket and the size of the holes. We find that the data are nicely fit by the expression.

$$N_L = \frac{\Omega}{\Omega_T} \left[1 + N_{01} \left(1 - N_{01} \frac{\Omega}{\Omega_T} \right) \right] \quad (9)$$

where N_{01} = number of neutrons entering zone 1 for the case of no loss cone. For the blanket used in these calculations, $N_{01} = 1.38$. Note that this expression is restricted to small values of Ω/Ω_T since the boundary condition that $N_L = 1$ at $\Omega/\Omega_T = 1$ is clearly not met.

If Eq. (9) holds for other types of blankets, the neutron leakage as a function of the loss-cone size could be predicted from a single calculation used to determine N_{01} (i.e., a calculation with no penetrations). A preliminary comparison for a Pb-Li blanket gave $N_{01} = 6.89$ for an 8-m-radius chamber and $N_L = 0.24$ for a 5-m-radius chamber with $\Omega/\Omega_T = 0.04$. This is exactly the value predicted by Eq. (9). These results must, however, be reevaluated with a consistent geometry. We are also concerned that homogenizing the Pb-Li blanket may overpredict the neutron leakage; this will also be investigated.

The performance characteristics of the lithium blanket are tabulated in Table 9-7 for the various loss-cone sizes. The tritium breeding by isotope, the neutron-energy leakage to zone 9, and the total neutron energy deposited in the blanket are listed.

The tritium breeding results, normalized to the result with no penetrations, are shown in Fig. 9-18 as a function of the loss-cone fraction. The breeding from ${}^7\text{Li}$ falls off as $(1 - \Omega/\Omega_T)$, the fraction of the total solid angle covered by the blanket. That is,

$$T_7 = T_{07} (1 - \Omega/\Omega_T) \quad (10)$$

where T_{07} is the number of ${}^7\text{Li}(n,n'T)\alpha$ reactions with no penetrations. This indicates that, in general, the leakage neutrons, other than source neutrons, are below threshold for the ${}^7\text{Li}(n,n'T)\alpha$ reaction. The breeding from ${}^6\text{Li}$ is fit by

$$T_6 = T_{06} (1 - N_L) \quad (11)$$

where T_{06} is the number of ${}^6\text{Li}(n,T)\alpha$ reactions with no penetrations and N_L is given in Eq. (9). In other words, there is a one-to-one correspondence between the total neutron leakage and the decrease in tritium breeding from ${}^6\text{Li}$ for this particular blanket. Since ${}^6\text{Li}$ accounts for 65% of the total breeding with no penetrations, the total tritium breeding falls off as

$$T_T = T_{0T} [0.65(1 - N_L) + 0.35(1 - \Omega/\Omega_T)] \quad (12)$$

The total energy deposition in the blanket falls off more rapidly than $(1 - \Omega/\Omega_T)$ due to the changing ratio of exoergic ${}^6\text{Li}(n,T)\alpha$ reactions to endoergic ${}^7\text{Li}(n,n'T)\alpha$ reactions. These two reactions dominate the net neutron-energy multiplication for this type of blanket. As shown in Table 9-7, the

Fig. 9-17. Increase in neutron leakage to zone 9 with increasing loss-cone fraction.

Fig. 9-18. Decrease in tritium breeding from ${}^6\text{Li}$ and ${}^7\text{Li}$ with increasing loss-cone fraction.

Table 9-7. Change in tritium-breeding performance and energy deposition with loss cone fraction. ▼

Total loss-cone fraction	Tritium breeding ^a			Energy deposition	
	T_6	T_7	T_T	Zone-9 (MeV)	Blanket ^b (MeV)
0	0.947	0.514	1.461	0	16.19
0.02	0.898	0.508	1.406	0.297	15.76
0.04	0.860	0.495	1.355	0.642	15.23
0.10	0.736	0.463	1.199	1.601	13.88
0.20	0.583	0.413	0.996	3.062	11.84

^a $T_6 = {}^6\text{Li}(n,T)\alpha$
 $T_7 = {}^7\text{Li}(n,n'T)\alpha$
 $T_T = T_6 + T_7$
^bIncludes energy leakage to zone 10, which is < 0.1 MeV.

total neutron energy deposited in a blanket with no penetrations is 16.19 MeV per source neutron. Accordingly, the net energy gain from the two tritium-breeding reactions is $[0.947(4.8 \text{ MeV}) - 0.514(2.5 \text{ MeV})] = 3.26 \text{ MeV}$. As the loss-cone fraction increases, the number of 14.1-MeV neutrons lost increases and the net gain from breeding reactions changes. Considering only these reactions, the total energy deposited would be given by

$$E = 16.19 - 14.1(\Omega/\Omega_T) + (4.8T_6 - 2.5T_7 - 3.26) \quad (13)$$

where T_6 and T_7 are the number of breeding reactions from ${}^6\text{Li}$ and ${}^7\text{Li}$ at the particular value of (Ω/Ω_T) . This expression, normalized to the result with no penetrations, is plotted in Fig. 9-19, using the tabulated values of T_6 and T_7 from Table 9-7. Also plotted in Fig. 9-19 are actual values of energy deposition

Fig. 9-19. Decrease in blanket energy deposition with increasing loss-cone fraction.

Fig. 9-20. Average energy of neutrons leaking through loss cones.

for the various loss-cone fractions. Clearly, the fit is excellent.

As previously noted from Fig. 9-18, T_6 and T_7 are closely fit by Eqs. (10) and (11). Using these expressions in Eq. (13), we find that the energy deposition as a function of Ω/Ω_T can be expressed totally in terms of the results from the reference calculation with no penetrations.

The neutron-energy leakage to zone 9 is the sum of the direct leakage of 14.1-MeV neutrons plus the contributions from scattered and secondary neutrons. Figure 9-20 shows the average energy of the leaking neutrons, excluding the 14.1-MeV leakage. This average energy is obtained by taking the total energy deposited in zone 9 (Table 9-7), subtracting 14.1 MeV times the number of 14-MeV neutrons entering zone 9, and dividing by the difference between $N_9(\text{total})$ and $N_9(14\text{-MeV group})$, as given in Table 9-6. For example, with $\Omega/\Omega_T = 0.1$, we have

$$\langle E \rangle = \frac{1.601 - 14.1(0.103)}{0.218 - 0.103} = 1.29 \text{ MeV}$$

As indicated, the average energy of the scattered and secondary neutrons leaking through the loss cones increases linearly, but moderately, with increasing loss-cone fraction for the range of Ω/Ω_T investigated here.

In summary, we have investigated the dependence of neutron leakage and blanket performance on the size of penetrations in a simple ICF chamber configuration. For this particular configuration, the total neutron leakage was found to be a factor of 2 larger than the fractional solid angle of the loss cones. The total leakage was related to the size of the loss cones (Ω/Ω_T) and to N_{01} , the number of neutrons per source neutron reentering the chamber from a blanket without penetrations. The tritium-breeding performance and total-blanket energy deposition were degraded faster than the loss of blanket coverage. The accelerated decrease in tritium breeding was found to be related to the decline in ${}^6\text{Li}(n, T)\alpha$ reactions. The decrease in blanket energy deposition was related to the size of the loss cones and to the number of tritium-breeding reactions in ${}^6\text{Li}$ and ${}^7\text{Li}$. The average energy of leakage neutrons, other than 14.1-MeV neutrons, was found to be a relatively weak function of the loss cone fraction. For the simple

geometry that we considered, the neutron leakage, tritium-breeding performance, and blanket energy deposition as a function of the loss-cone fraction of the total solid angle could all be related to N_{01} and Ω/Ω_T .

These results illustrate the importance of properly treating the effects of chamber ports in determining the nuclear performance of reactor concepts. These effects will be most critical for chambers with high values of N_{01} and with marginal performance characteristics.

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Materials Considerations

Solubility of Hydrogen Isotopes in Lead-Lithium Solutions.³³

Lead-lithium (Pb-Li) solutions are of interest to designers of liquid metal wall (LMW) ICF reactors because Pb may be present to some extent in both heavy-ion-beam- and laser-driven ICF targets, and therefore, present as an impurity in a flowing lithium wall. In addition, a Pb-Li solution containing ~80 atomic percent (at.%) Pb is a strong candidate for a heavy-ion-beam-driven HYLIFE reaction chamber and is also a viable alternative to a "pure" Li wall for a laser-driven reaction chamber. The solubility of hydrogen isotopes in the LMW will determine whether tritium must be reprocessed from the liquid loop, the vacuum system, or both.

For a 2700-MW_f, 600-m³ chamber with a 30% tritium burn fraction and a 1.75 tritium-breeding ratio, 19.5 mg/s of tritium are added to the reaction chamber. If tritium is to be recovered from only the liquid, a maximum loss through the vacuum system of 1% (0.2 mg/s) is assumed to be acceptable. With a laser driver,³⁴ the vacuum system might operate at ~10⁻² Torr with about 10% of the vapor being pumped per second. Then, about 2.6 × 10⁻⁵ Torr of tritium, or 10 times the pressure associated with 0.2 mg tritium, is allowable.

With a heavy-ion-beam driver,³⁴ the vacuum system might operate around 10⁻³ Torr, with about 50% of the vapor being pumped per second. Then, about 5 × 10⁻⁶ Torr of tritium, or twice the pressure associated with 0.2 mg tritium, is allowable if only the liquid is to be reprocessed.

For vapor-only tritium recovery, 99% of the tritium production (or 19.3 mg/s) must be processed by the vacuum system, and the required pressure is ~100 times the allowable pressure for liquid-only recovery. Therefore, with a laser driver, tritium vapor pressure between 2.6 × 10⁻⁵ and 2.6 × 10⁻³ Torr requires tritium recovery from both the liquid-circulation and the vacuum systems. For the heavy-ion-beam driver, the corresponding pressure range is 5 × 10⁻⁶ to 5 × 10⁻⁴ Torr. The actual tritium vapor pressure will depend on both the vacuum system design and the solubility of tritium in the LMW.

Efficient tritium recovery cannot begin until the tritium concentration builds up in the liquid. A start-up tritium inventory of about 1.4 kg is required for each day of start up. For a 2700-m³ system of Li₉₉Pb₁ (99% lithium, 1% lead), about a day is required to reach 1 wppm (1.7 kg) tritium. For a 1500-m³ system of Pb₈₃Li₁₇, about eight days are required to reach 1 wppm (14-kg tritium).

As of December 1980, no data have been generated that directly give the hydrogen isotope pressure above Pb₈₃Li₁₇ liquid in the 400 to 500°C temperature range. However, data at other temperatures and concentrations can be extrapolated to the needed temperatures and compositions. At a given composition, Sievert's constant (K_s) is related to temperature by

$$\log K_s = m/T + b \quad (14)$$

where m and b are temperature-independent constants and T is temperature. If K_s vs composition is known at two temperatures, the constants m and b can be calculated at each composition and K_s can be calculated at all temperatures and compositions.

The data points of Ihle, Neubert, and Wu³⁵ for the solubility of deuterium in several Pb-Li compositions are shown in Fig. 9-21. The solid lines at temperatures of 677 and 767°C are interpolations between data points. The gap in the 677°C curve is due to the presence of a solid phase. The data have been extended to the deuterium-in-lithium data from Argonne³⁶ and the hydrogen-in-lead data from Opie and Grant.³⁷ The hydrogen-in-lead data were converted to deuterium-in-lead data by assuming that the ratio of Sievert's constants

for hydrogen and deuterium is constant (independent of the Pb-Li composition) and by using the Argonne data for hydrogen and deuterium in lithium. The adjustment was relatively small, with a 9% increase at 767°C and a 20% increase at 477°C.

Using the curves at 677 and 767°C and Eq. (14), we constructed curves at 477 and 577°C; however, at Pb concentrations above 50 at.%, the 677°C curve is quite uncertain. Figure 9-22 presents a second plausible

Fig. 9-21. Sievert's constant for solubility of deuterium in Pb-Li.

Fig. 9-22. A second possible extrapolation of the data on solubility of deuterium in Pb-Li.

677°C curve and the resulting lower-temperature curves. The deuterium data plotted in Figs. 9-21 and 9-22 can be used to determine whether hydrogen isotopes must be recovered from the liquid phase, the vapor phase, or both. The conclusions apply to all three hydrogen isotopes, since the isotopic correction is fairly small.

For a "pure lithium" LMW, Pb target debris can build up to ~1 at.% before cold trapping of binary Pb-Li compounds is effective. If deuterium is removed from the liquid to a concentration of 1 wppm (4.5×10^{-6} atom fraction of deuterium in alloy), the diatomic deuterium pressure at 477°C is $\sim 5 \times 10^{-10}$ Torr, and hydrogen isotopes need not be recovered in the vacuum system for either laser or heavy-ion-beam drivers.

For the $\text{Pb}_{83}\text{Li}_{17}$ LMW, no conclusion can be reached from the data. Using the extrapolation of Fig. 9-21, and liquid recovery of hydrogen isotopes down to 1 wppm each (8.7×10^{-5} atom fraction of deuterium in alloy), we find that the diatomic deuterium pressure at 477°C is $\sim 3 \times 10^{-8}$ Torr, and no vacuum-system hydrogen-isotope recovery is required. However, if the Fig. 9-22 extrapolation is used, and if deuterium is extracted from the vapor to a pressure of 2.6×10^{-3} Torr (laser driver) or 5×10^{-4} Torr (heavy-ion-beam driver), the corresponding deuterium concentrations in the liquid are 0.034 and 0.015 wppm, respectively, and no liquid recovery of hydrogen isotopes is required. To resolve this issue, researchers at Argonne National Laboratory³⁸ and elsewhere are conducting appropriate experiments.

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Corrosion of 2-1/4 Cr-1 Mo in Lithium and Lead-Lithium Solutions. A serious concern in fusion technology is the ability of standard engineering alloys to withstand corrosive attack by liquid lithium at approximately 500°C. Economics and design feasibility make the use of presently coded materials highly desirable, yet very limited data exist on the lithium containability of even the most common engineering alloys. Consequently, at the request of LLNL, the Metallurgical Engineering Department of

the Colorado School of Mines began a long-range experimental test program in 1978 to systematically develop the data base needed to select laser fusion materials.

The initial experiments^{39,40} were designed to compare the stabilities of the two most likely candidate alloys, 304 stainless steel and 2-1/4 Cr-1 Mo pressure-vessel steel, in both lithium and an alloy of 17.6 Pb-82.4 Li. (Numerals are weight percents.) Considerable industrial experience with 304 stainless steel in liquid sodium favored that alloy, while appreciable savings in cost and conservation of a strategic material (chromium) favored 2-1/4 Cr-1 Mo steel. The test matrix for the initial study is shown in Table 9-8.

Compatibility of 2-1/4 Cr-1 Mo steel in liquid lithium at temperatures up to 500°C was evaluated by using two separate concepts in specimen design. In the first concept, bent test specimens, spot welded by the gas-tungsten-arc (GTA) method, were used to evaluate the effects of tensile and compressive residual stresses in materials up to yield stress levels. In the second concept, rectangular thin-sheet test coupons were used to determine weight-loss kinetics of the test materials in liquid lithium.

The 2-1/4 Cr-1 Mo steel can be heat treated to many different microstructures. Screening tests were made on two heat treatments:

- Water quenched and tempered 1 h at 693°C.
- Isothermally transformed (2 h at 740°C) and tempered 10 h at 760°C.

The water-quench-and-temper heat treatment produced a tempered bainite structure, while the transform-and-temper heat treatment produced a ferritic-bainitic structure. Previous data have shown that cementite (Fe_3C), which is readily attacked by lithium, can be eliminated by a long temper of 2-1/4 Cr-1 Mo steel at 760°C, which replaces cementite with M_{23}C_6 and M_6C . The water-quench-and-temper heat treatment, therefore, produces the most susceptible microstructure with regard to lithium-corrosion resistance, while the transform-and-temper heat treatment produces the least susceptible microstructure.

Weldments of 2-1/4 Cr-1 Mo steel in the as-welded condition were readily attacked by liquid lithium at the heat-affected zone.

Intergranular penetration at the prior austenite grain boundaries was observed.

Weight loss of isothermally transformed and tempered 2-1/4 Cr-1 Mo steel, tested at 500°C with a nitrogen level of ~500 ppm in the lithium, was found to be parabolic with time, with an observed weight loss of 0.3 mg/cm² after 400 h. Base-metal weight losses were found to be significantly less for 2-1/4 Cr-1 Mo steel in liquid lithium than for 304 stainless steel in 17.6 Pb-82.4 Li. A marked sensitivity to the nitrogen content of the lithium was observed; when nitrogen concentration was increased from ~500 to ~4000 ppm, the liquid lithium attack on 2-1/4 Cr-1 Mo steel increased dramatically.

We decided to concentrate on the lithium corrosion susceptibility of 2-1/4 Cr-1 Mo steel weldments in low-nitrogen 17.6 Pb-82.4 Li for the second year of study.⁴¹ Base-metal stability of properly heat-treated 2-1/4 Cr-1 Mo steel in liquid lithium had been found to be good, and we felt that lithium penetration of as-welded heat-affected zones in 2-1/4 Cr-1 Mo steel could be alleviated by

Table 9-8. Colorado School of Mines (CSM) screening tests. (a) 1978 and (b) 1979.

(a)		Environmental parameters [N] ≥ 500 ppm, T = 485°C	
		Pb-free lithium	17.6 Pb - 82.4 Li
304, annealed	GTA weld coupon	X	X
	Base-metal coupon	X	X
2-1/4 Cr-1 Mo, quenched and tempered	GTA weld coupon	X	
	Base-metal coupon		
2-1/4 Cr-1 Mo, isothermally transformed and tempered	GTA weld coupon		
	Base-metal coupon	X	
(b)			
Test conditions			
● 17.6 Pb-82.4 Li.			
● Heat treatment: austenitized (927°C), cooled at 50°C/h to 704°C, isothermally transformed at 704°C (2 h), cooled at 50°C/h to 500°C, furnace cooled, tempered at 760°C for 10 h.			
● Preheat: 200°C.			
● GTA weld coupons.			
● Exposure time: 1600 h.			
Alloy	Postweld heat treatment		
	None	10 h at 710°C	10 h at 760°C
2-1/4 Cr-1 Mo, Nb-stabilized	—	—	—
2-1/4 Cr-1 Mo, unstabilized	—	—	—

postweld heat treatments (PWHT) or by the use of Nb-stabilized grades. The test matrix for the second year of study is shown in Table 9-9.

The regular-grade 2-1/4 Cr-1 Mo steel with no PWHT experienced the most rapid corrosion, especially during initial stages of attack (<100-h exposure). A weldment subjected to an 18-ksi tensile stress did not corrode more rapidly than the unstressed weldments.

A dramatic improvement in lithium-corrosion resistance is afforded the weldments of regular grade 2-1/4 Cr-1 Mo steel by a high-temperature PWHT. The 760°C PWHT was found to be significantly better than the 710°C PWHT, probably because Fe_3C and Mo_2C were more completely converted to M_{23}C_6 and M_6C at the higher temperature.

Weldments of the Nb-stabilized 2-1/4 Cr-1 Mo steel were quite resistant to corrosion by low-nitrogen (~50 ppm) 17.6 Pb-82.4 Li. Only slight penetration (near the limit of measurement) was observed after 1400 h, and the heat-affected zones of the welds were not preferentially attacked.

The Nb-stabilized alloy was extremely resistant to intergranular attack, even when no PWHT was applied. This corrosion resistance can be attributed to the formation

of stable refractory carbides upon cooling. Since these Nb carbides form in preference to cementite, a PWHT is not required for good corrosion resistance. The disadvantages of the Nb-stabilized alloy involve weldability, availability, and the undesirable nuclear properties of two of the alloying elements: nickel and niobium.

We postulated two possible causes for lithium attack on the heat-affected zone (HAZ): microstructural changes and residual stresses caused by welding. In an attempt to isolate the effects of these causes, we heat treated coupons of regular-grade 2-1/4 Cr-1 Mo steel to produce homogeneous samples of the coarse-grained bainitic microstructure found in the corrosion-susceptible HAZ of GTA weldments. The HAZ microstructure was best approximated by a 1300°C austenitizing treatment followed by an oil quench.

These simulated weldments corroded intergranularly, as did the HAZ of actual weldments. Although these results did not eliminate the possibility of stress-assisted lithium corrosion, they did establish the correlation between microstructure and lithium attack. Auger spectroscopy of selected samples revealed high concentrations of sulfur in those grain boundaries not yet attacked by lithium and high

Fig. 9-23. Corrosion of simulated weldments.

concentrations of oxygen in those grain boundaries already penetrated. The full implications of these results are not yet known.

Simulated-weldment specimens were also subjected to 17.6 Pb-82.4 Li at various temperatures to establish a temperature dependence for the intergranular penetration. These data are shown in Fig. 9-23. Note that penetration is parabolic with time at both low and high temperatures but changes from a high initial rate to a slower intermediate rate at intermediate temperatures (440 to 530°C). An Arrhenius plot of these data is shown in Fig. 9-24. The initial penetration is believed to be controlled by liquid diffusion of carbon in lithium; however, the secondary penetration is not yet clearly understood.

The need for careful PWHT of 2-1/4 Cr-1 Mo welds had been established by the studies of the previous year, and research was extended to examine two other possible parameters:

- High concentrations of lead in the lithium.
- Velocity effects.

The test matrix for these experiments is given in Table 9-10.

Preliminary results of the lead-concentration study are shown in Fig. 9-25. The effectiveness of a 760°C PWHT was

again illustrated; none of the specimens given this PWHT exhibited observable penetration up to exposures of 1000 h. However, the effect of lead concentration was puzzling. We found that, while 17.6

Table 9-9. Colorado School of Mines (CSM) screening tests, 1980.

Heat treatment	Velocity ^a			Lead concentration ^b		
	Zero	Max/2	Max	Zero	17.6 w/o	99.5 w/o
1300°C oil quench	X	X	X	X	X	X
1300°C oil quench 10 h at 760°C	X	X	X	X	X	X

^aPb concentration = 17.6 w/o. Exposure time = 400 h.

^bExposure time = 1000 h.

Fig. 9-24. Arrhenius plot of the corrosion-rate constant for simulated weldments.

Fig. 9-25. Penetration in 2-1/4 Cr-1 Mo steel.

weight percent (wt%) Pb apparently increased the penetration rate over that found for Pb-free lithium, no penetration was observed with 93.3 wt% Pb, even when no PWHT was given to the specimens. Clearly, much remains to be learned concerning the effect of lead concentration.

The velocity-effects study is still in progress. Results of this work will be discussed in the annual final report to LLNL.

In summary, we find that GTA weldments of 2-1/4 Cr-1 Mo steel are susceptible to severe intergranular attack in the HAZ by 17.6 Pb-82.4 Li because the cementite and epsilon carbide in the HAZ are thermodynamically unstable in lithium below $\sim 700^{\circ}\text{C}$. Corrosion resistance of 2-1/4 Cr-1 Mo weldments can be greatly improved by a PWHT that converts the unstable carbides to stable carbides, such as M_{23}C_6 . An insufficient postweld temper will produce needle-like Mo_2C precipitates, which are susceptible to lithium attack. Niobium-stabilized 2-1/4 Cr-1 Mo steel forms stable, refractory carbides upon cooling from the austenitic range; a PWHT is not required for corrosion resistance of this steel.

The HAZ of GTA weldments can be adequately simulated by an appropriate heat treatment. The simulated weldments, along with an actual GTA-weldment HAZ, corrode at an initially high rate. This rate has an activation energy of 5.3 kcal/mole and is probably controlled by the rate of carbon diffusion through the lithium in the grain boundary. When epsilon carbide is converted to cementite upon tempering in the lithium, the penetration rate is reduced. The secondary penetration rate, which has an activation energy of 28 kcal/mole, is most likely reaction controlled rather than diffusion controlled.

The corrosion rate is not a strong function of cooling rate from the austenitic range, but the initial corrosion rate increases with increasing austenitizing temperature. It is believed that the larger grain size produced at a higher austenitizing temperature speeds the corrosion by providing an easier diffusion path for carbon through the corroded grain boundary.

High concentrations of sulfur and oxygen have been found at corrosion-susceptible grain boundaries of 2-1/4 Cr-1 Mo steel, and lithium oxide has been found at

corroded grain boundaries. The roles of oxygen and sulfur at corrosion-susceptible grain boundaries are not currently known; however, the presence of lithium oxide as a corrosion product is consistent with the proposed mechanisms of corrosion.

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Fluid and Pumping Considerations

Transport Processes in the HYLIFE Reaction Chamber

Introduction. In the HYLIFE inertial confinement fusion reaction chamber,^{42,43} a close-packed annular array of lithium jets surrounds the D-T pellet. The jet array supersedes earlier designs in which a concentric annulus⁴⁴ or a series of concentric annuli,⁴⁵ separated by gaps, was proposed. With either of these latter arrangements, hot, high-pressure lithium plasma, deriving from low-energy x-ray and ionic target-debris deposition, was confined within the cavity formed by the inner radius of the innermost annulus. This caused the plasma to expand against, and impart momentum to, the liquid. Rayleigh-Taylor instability would eventually cause the plasma to break through the liquid barrier(s), but calculations showed that this was not likely to occur before unacceptably high radial velocity was attained by the liquid. This, in turn, could inflict catastrophic damage on the first wall in the more efficient, high-energy-density designs. The jet array circumvents the problem by allowing the plasma to flow around or in between the jets and thereby vent the cavity. In order for the discrete jets to function properly, however, the drag exerted on them by the venting plasma must not cause too great an impulse to be delivered. Moreover, the plasma must transfer energy to the liquid at a rate sufficient to preclude excessive gas pressure buildup on the cavity wall. Finally, the

plasma must eventually cool and condense on the liquid surfaces, or be removed from the cavity, before the next fuel pellet is injected (typically in a second or less).

Earlier, we devised a quasi-one-dimensional method for calculating the drag impulse exerted on the jet array by the venting plasma.⁴³ The method was based on the topological correspondence between an impulsive divergent flow constricted by the presence of jets and an equivalent flow between inclined wavy walls. The impulse delivered to the jets is due partly to skin friction and partly to profile drag, and both of these effects were accounted for, as was radiation flow in the plasma. Heat and mass transfer between the plasma and the liquid were excluded, however, and subsequent calculations were planned to explore this influence. These calculations have now been completed and are the main subject of this report. (See Table 9-10 for nomenclature.)

Model Description. An end-view of the jet array in a typical HYLIFE design is shown in Fig. 9-26. A discussion of the sequence of events following the energy deposition from the microexplosion of the pellet is given in Ref. 43. Here, we assume that the inner cavity, bounded by $r = R_0$ contains a plasma at a temperature $T(r,0)$, density $\rho(r,0)$, and particle velocity $u(r,0)$; typical initial profiles are shown in Ref. 43. For $R_0 < r \leq R_2$, $\rho(r,0) = T(r,0) = u(r,0) = 0$ (i.e., the jets are initially surrounded by a vacuum, which is the condition of the outer, jet-free region as well). The plasma will then flow in between the jets and eventually contact the cavity wall (taken to be rigid insofar as the flow is concerned) at $r = R_2$. Actually, there will probably be a thin layer of liquid lithium flowing axially down the wall, but this need not concern us here.

Taking the fusion target to be a line-source—rather than a point-source—enables the crossflow to be treated as quasi-one-dimensional, with cylindrical symmetry. Axial pressure relief is thereby ignored, and if the energy density is taken at the plane of the fusion pellet, an upper bound to the wall bounding will then be obtained. (Higher-dimensional effects will be discussed in subsequent work.)

The equations of motion differ somewhat from our earlier model. For mass conservation,

$$\frac{d}{dt}(\rho A) + \rho A \frac{\partial u}{\partial r} = f(1 - \zeta)/H_v, \quad (15)$$

where

$$A = A(r) = 2\pi r, \quad r < R_0, \quad r > R_1 \\ < 2\pi r, \quad R_0 \leq r \leq R_1, \quad (16)$$

Table 9-10. Fluid and pumping nomenclature.

a	radiation density constant
A	flow cross section (per unit axial length)
c	speed of light
C_f	skin friction coefficient
C_v	heat capacity
D	jet (or jet fragment) diameter
e	specific internal energy
e_m	material contribution to specific internal energy
e_r	radiation contribution to specific internal energy
E_K	total plasma kinetic energy in cavity
E_T	total plasma (kinetic + internal) energy in cavity
f	rate of energy delivery to jet surface (by radiation and convection from the plasma) per unit area per unit time
F_0	Fourier modulus
h	convective heat transfer coefficient
H_v	enthalpy of vapor at jet surface temperature
I_T	total impulse delivered to jets
k	thermal conductivity of (liquid) jets
K	photon-frequency-averaged opacity (K_R is Rosseland mean)
ℓ	characteristic beam length for radiation transport between plasma and jets
L	radiant energy flux in the plasma
m	mass of plasma in cavity
P	pressure (P_{gw} is pressure exerted on cavity wall)
P_m	material contribution to pressure
P_r	radiation pressure [defined by Eq. (22)]
Q	energy flux from plasma to jets interior [defined by Eq. (27)]
r	radial coordinate
R_0	inner radius of jet array
R_1	outer radius of jet array
R_2	cavity radius
t	time
T	plasma temperature (T_{gw} is plasma temperature at cavity wall)
T_s	jet (or jet fragment) surface temperature
T_v	saturation temperature of jet surface
T_c	critical temperature
u	particle velocity
V	specific volume
α	thermal diffusivity
δ_w	cavity wall thickness
ϵ	emissivity coefficient [defined by Eq. (18)]
ζ	fraction of energy delivered to jet surface that diffuses into the interior; $(1 - \zeta)$ is the fraction that results in vapor generation
λ	radiation mean free path
ρ	plasma density
σ_θ	hoop stress in cavity wall
τ	characteristic time
ϕ	cumulative jet surface (per unit axial length)

and

$$f(r,t) = h(T - T_S) + \frac{ac}{4}(\epsilon T^4 - T_S^4) \quad (17)$$

The emissivity coefficient, ϵ , is defined by

$$\epsilon = \rho K \ell \leq 1 \quad (18)$$

where

$$K = K(\rho, T) \quad (19)$$

is an appropriate photon-frequency-averaged opacity and ℓ is the characteristic beam length, in general a function of the diameter and spatial distribution of the jets. Any mass vaporized from the jet surface is assumed to instantaneously mix with the plasma, so that local thermodynamic and mechanical equilibria are maintained.

Since the jet velocity is negligible with respect to that of the plasma, momentum conservation can be expressed as in Ref. 43, i.e.,

$$\rho \frac{du}{dt} + \frac{\partial}{\partial r} (P_m + P_r) + C_t \rho u |u| \frac{1}{A} \frac{d\phi}{dr} = 0 \quad (20)$$

where

$$P_m = P(\rho, e_m) \quad (21)$$

$$P_r = aT^4/3 \quad (22)$$

The symbol ϕ identifies the cumulative jet surface, which is obtained by sweeping outward along a radius vector from the innermost jets; ϕ , therefore, varies from 0 at $r = R_0$ to πD (where D is the jet diameter) at $r = R_1$. For the jet array depicted in Fig. 9-26, the functions $A(r)$ and $\phi(r)$ are shown in Ref. 43. The impulse delivered to the jets is due partly to form (or profile) drag and partly to skin friction. The former derives from the normal stress or pressure, whereas the latter is the time integral of the shearing stress taken over the surface of the jets. The last term in Eq. (20) is intended to model the shear retardation of the plasma and was deduced for this purpose from quasisteady parallel flow theory.⁴⁶

The rate at which the specific internal energy of the plasma changes is given by

Fig. 9-26. End view of the jet array in a typical HYLIFE ICF reactor design.

$$\frac{de}{dt} = -(P_m + P_r) \frac{dV}{dt} + C_f u^2 |u| \frac{1}{A} \frac{d\phi}{dr} - \frac{1}{\rho} \frac{\partial}{\partial r} (L + Q) , \quad (23)$$

where

$$e = e_m + e_r = e_m(\rho, T) + aT^4/\rho , \quad (24)$$

$$V = 1/\rho , \quad (25)$$

$$L = - \frac{c}{3} \frac{1}{K_R \rho} a \frac{\partial}{\partial r} T^4 , \quad (26)$$

$$\frac{\partial Q}{\partial r} = \frac{1}{A} \frac{d\phi}{dr} \zeta f . \quad (27)$$

We note that both C_f in Eqs. (20) and (23) and h in Eq. (17) are, in general, functions of the local Reynolds and Mach numbers; h depends weakly on the Prandtl number, as well.

The equation for overall energy conservation is readily derived by taking the scalar product of Eq. (20) and u , the particle velocity, and then adding the result to Eq. (23). It thereby becomes apparent that whenever $\zeta \neq 1$, energy is dissipated at the rate $(P_m + P_r)f(1 - \zeta)/\rho^2 AH_v$ per unit mass—a consequence of the assumption of immediate mixing of the vapor generated at the jet surface with the free-stream plasma.

Equations (15) to (27) are not closed: neither T_S nor ζ has yet been prescribed. If $T_S = T_S^0 = T_S(r, 0) = \text{constant}$, two limiting cases are of interest: i.e., $\zeta = 1$ and $\zeta = 0$. The former implies maximum heat transfer from the plasma to the jets so that the residual energy per unit volume—and, hence, the average pressure—in the plasma will be a minimum. With $\zeta = 0$, on the other hand, we have the case of immediate vapor formation and mixing with the plasma, resulting in maximum mass transfer; the energy per unit volume in the plasma will remain invariant, but the average temperature (simply related to the energy per unit mass) may be lower. In general, however, we expect that $T_S \neq T_S^0$ for $t > 0$ and that $0 < \zeta < 1$. Physically, the surface temperature of the jets, or of their disassembled parts, will probably increase until the local vapor pressure at the surface exceeds the local ambient pressure of the

plasma. (In the time period of interest, at least the inner rows of jets will break apart as a result of isochoric neutron heating.⁴³⁻⁴⁵)

At this point, we assume that vaporization will commence and that the interface will continue to recede until the local heat-transfer rate can be accommodated by the rate at which heat can diffuse into the interior liquid.

For the heat conduction, and as long as the jet or particle radius is not too small ($F_0 = 4\alpha\tau/D^2 \ll 1$), it can be shown that T is given by the Duhamel integral⁴⁷:

$$T_S(r, t) - T_S^0 = (1/k) \sqrt{\alpha/\pi} \times \int_{t^0(r)}^t f(r, t - \tau) d\tau / \sqrt{t - \tau} = (1/k) \sqrt{\alpha/\pi} \times \int_{t^0(r)}^t f(r, \tau) d\tau / \sqrt{t - \tau} , \quad (28)$$

Equation (28) is assumed to hold as long as

$$T_S(r, t) < T_v [P(r, t)] \leq T_c , \quad (29)$$

where T_v is the saturation temperature (the boiling point at the local free-stream pressure) and T_c is the critical temperature. We can generalize Eq. (28) to include the Stefan problem as follows:

$$T_S(r, t) - T_S^0 = (1/k) \sqrt{\alpha/\pi} \times \int_{t^0(r)}^t \zeta(r, \tau) f(r, \tau) d\tau / \sqrt{t - \tau} . \quad (30)$$

Whenever the inequality (29) is satisfied, $\zeta(r, t) = 1$, and $T_S(r, t)$ results directly from the quadrature; otherwise, we set $T_S = T_v$ and invert Eq. (30) to find $\zeta(r, t)$.

Numerics. We employed operator splitting to obtain finite-difference analogs of Eqs. (15) to (30), with the Lagrangian hydrodynamics, advection, heat transfer, and radiation diffusion performed in sequence. The formulation of the difference equations (neglecting plasma-jet heat transfer) is discussed in Refs. 43, 44, and—in more detail—in Ref. 48. Also, the equation of state and opacity data for lithium are outlined in Ref. 44.

When heat transfer is included, the main problem is evaluation of the integral in Eq. (28), or, more generally, in Eq. (30). A

difficulty arises in that the kernel requires not only calculation at each spatial location, but storage in time as well. To minimize temporal storage, we subdivided the time interval from initial heating to the present, $\Delta t^n(r_j) \equiv t^n - t^0(r_j^n)$, into $k - 1$ equal segments on whose boundaries k values of ζf were to be defined. The first $k - 1$ values of the kernel were interpolated from the k stored values at t^{n-1} , with the right-hand (k th and current) value derived from Eq. (17). We note here in passing that, since the spatial increments defining the plasma were not, in general, congruent with those defining the jets, it was necessary to map the surface temperature of the jets onto the plasma in order to obtain f (in Eq. 17). Similarly, f was required to be remapped on the jets' surface to find T_S .

With $\zeta(r_j^n, t^i) f(r_j^n, t^i)$, $i = 1, \dots, k$ defining a piece-wise linear approximation to the kernel, it is easy to show that the exact solution to Eq. (30) is

$$\begin{aligned} T_S(r_j^n, t^k) - T_S^0 \\ = 2\sqrt{t^k - t^{k-1}} \sum_{i=1}^{k-1} [\sqrt{k-i} \\ - \sqrt{k-i-1}] \{ \zeta(r_j^n, t^i) f(r_j^n, t^i) \\ + [\zeta(r_j^n, t^i) f(r_j^n, t^i) - \zeta(r_j^n, t^{i+1}) f(r_j^n, t^{i+1})] \\ \times [(k-i) - \{2(k-i) - 1 \\ + \sqrt{(k-i)(k-i-1)}\}/3] \} \quad , \end{aligned} \quad (31)$$

where

$$\begin{aligned} t^k &\equiv t^n \text{ and } t^k - t^{k-1} = t^{k-1} - t^{k-2} \\ &= \dots = t^2 - t^1 = t^1 - t^0(r_j^n) \quad . \end{aligned}$$

In practice, the first $(k - 2)$ terms in Eq. (31) are summed separately, with the last term then tentatively added. If after evaluating T_S , the inequality (29) is satisfied, the results stand; if not, the value of $\zeta(r_j^n, t^n) f(r_j^n, t^n)$ needed to balance Eq. (31) with the left-hand side set equal to $T_v[P(r_j^n, t^n)] - T_S^0$ is easily found and $\zeta(r_j^n, t^n)$ is determined. We note that it is sometimes the case that the right-hand side of Eq. (31) exceeds $T_v - T_S^0$, even with $\zeta = 0$. Since the saturation temperature is a function of the local pressure, a drop in the former combined with, say, a constant or increasing heat flux from one time to the next could cause this to happen (the heat flux is only weakly coupled to the surface temperature when $T \gg T_S$).

When this situation arises, Eq. (31) is only weakly enforced, in the sense that $\zeta(r_j^n, t^n)$ is set exactly equal to zero. The numerical problem is magnified somewhat by interpolation errors derived from the approximation of the kernel in Eq. (30) by only k -discrete values. These errors are minimal, however, since the time domain $t^n - t^0(r_j^n)$ was allowed to increase only very slowly [e.g., $(t^n - t^0)/(t^{n-1} - t^0) < 1.02$ for $t^n \gg t^0(r_j)$]. The solutions appear to converge quite rapidly with increasing k ($\ll n$). For example, doubling the number of intervals from 10 to 20 in a test problem where n exceeded 20 000 resulted in the mass ratio (initial plasma mass + mass vaporized from jets/initial plasma mass) changing by only 0.3%. Similar comparisons were made with other key parameters. Subsequent problems were then run with 10 intervals ($k = 11$).

Parameter Study. A number of parameters in our model are not well-known (and, in some cases, extremely difficult to estimate). Some of these can be eliminated from concern on the basis that they have little influence on the results. For example, it has previously been shown that the viscous contribution to the crosswise momentum imparted by the plasma to the jets was negligible in comparison with the effect of profile drag.⁴³ Therefore, fairly large uncertainty in the friction coefficient, C_f , can probably be tolerated. Similarly, it was

Fig. 9-27. Variation of the plasma kinetic energy in the cavity with time.

found in the present study that, in the time period for the maximum hoop stress to develop in the first wall, convective heat transfer was insignificant when compared with radiative heat transfer, so that errors in estimating the convective coefficient are much less important than errors of the same order in the emissivity.

As defined in Eq. (18), the emissivity is essentially the ratio of two lengths, the characteristic beam length, ℓ , and the radiation mean free path, $\lambda = (\rho K)^{-1}$. Hottel and Sarofim⁴⁹ have suggested that an appropriate average beam length for radiant interchange between a gas volume of any given shape and its bounding surface is approximately 3.5 times the quotient of volume by surface, i.e., $3.5 r_H$, where r_H is the hydraulic radius. For the array geometry shown in Fig. 9-27, $\ell = 216$ mm ($D = 200$ mm). We note, however, that this is certainly a very crude estimate and that, in general, $\ell = \ell(r,t)$, implying that the spatial distribution of jets, as well as their size, should be considered in calculating ℓ .

Closely related to the spatial distribution of jets is the flow cross section, $A(r)$. As mentioned earlier, even before the plasma blowback has come in contact with the jets, isochoric neutron heating has caused considerable jet disassembly in at least the first row of jets; by 3 ms, the particulation is fairly widespread throughout all the jets.⁴⁵ This means that in reality $A = A(r,t)$. More important, the manner in which A changes in time is unknown and, in general, will depend on the disassembled particles' size and spatial distribution. The same is true for the jet surface ϕ . Once the plasma begins to interact with the particles, of course, drag will alter the particle trajectories, and this may require coupling the particle motion to that of the plasma. For the moment, and lacking even meager experimental data, we have arbitrarily chosen A to be time-independent and to assume a uniform distribution of particle sizes. Given these assumptions and the fact that the liquid (jet) mass is conserved, $\phi \sim D^{-1}$, e.g., if the 200-mm, initial-diameter jets were to uniformly subdivide into 20-mm particles, their total surface would be 10 times the initial surface. The importance of these approximations is shown below.

The Case of Maximum Heat Transfer.

Setting $\zeta(r,t) = 1$ and $T_S(r,t) = T_S^0$

(equivalent to $k \rightarrow \infty$) maximizes the heat transfer from the plasma to the jets, which in turn minimizes the average pressure in the cavity. Figures 9-27 to 9-30 contrast these results, with those obtained without heat transfer. (Table 9-11 provides an identification key for the letter code appended to each curve in Figs. 9-27 to 9-37.) In the latter case, the calculation

Note: Same cases as in Fig. 9-27.

Fig. 9-28. Energy transported from the plasma to the jets as a function of time for the cases of Fig. 9-27.

Note: same cases as in Fig. 9-27. See also Table 9-12.

Fig. 9-29. Pressure exerted on the cavity wall as a function of time for the cases of Fig. 9-27.

Note: same cases as in Fig. 9-27.

Fig. 9-30. Temperature of the plasma adjacent to the cavity wall as a function of time for the cases of Fig. 9-27.

discussed in Ref. 43 was extended in time from 360 μ s to 3 ms after pellet implosion.

Figure 9-27 shows the variation in the plasma kinetic energy in the cavity, $E_K(t)$,

Table 9-11. Description and parameters for cases studied.

Case	Description	$D(\text{mm})$	$T_s(\text{K})$	ζ
(a)	No heat or mass transfer	200	—	—
(b)	Maximum heat transfer	200	750	1
(c)	↓	2	750	1
(d)	Maximum mass transfer	2	750	0
(e)	Combined heat and mass transfer	200	Variable	Variable ($0 \leq \zeta \leq 1$)
(f)	↓	20	↓	↓
(g)	↓	2	↓	↓

Fig. 9-31. The pressure-time history on the cavity wall; a case of maximum heat transfer compared with a similar calculation assuming maximum mass transfer. See Table 9-11 for identification key and further details. ▶

Fig. 9-32. The ratio of the plasma mass in the cavity (including the component vaporized from the jet surface) to the initial plasma mass, for maximum mass transfer and three cases of combined heat and mass transfer. See Table 9-11 for identification key and further details.

Fig. 9-33. The temperature-time history for the plasma adjacent to the cavity wall for the cases of Fig. 9-31. ~

with time; the results are normalized to E_T^0 , the total plasma energy in the cavity just prior to plasma contact with the jets at $t = 60 \mu$ s. At this (initial) point, $E_K/E_T^0 = 0.1$. A rapid increase then occurs as the plasma blows through the jet array until, at about 250 μ s, the leading edge of the plasma contacts the cavity wall. Without heat transfer to the jets, more than half of the total plasma energy is in the form of kinetic energy at this point. The kinetic energy is then quickly dissipated, decreasing approximately to its initial value and then slowly to zero. (The oscillation period corresponds to the time for shocks to reflect both from the cavity wall and the jet array.) Each such interaction dissipates additional energy. In principle, frictional heating contributes to the dissipation as the plasma passes back and forth through the jets, but for the cases considered here, frictional heating is negligible.

The situation with heat transfer allowed to the jets differs in some important respects. For the same initial jet diameter the kinetic energy gained by the plasma in the first pass through the jets is about 15% less. Also, the kinetic energy is dissipated much more rapidly and is practically nonexistent after the first shock reflection from the outer row of jets. Figure 9-28 shows the energy transport to the jets. By 350 μ s the main body of plasma has passed through the jet array and about 25% of the total plasma energy has been lost. Most of this has been to the first row of jets (>70%) and most has been transported by radiation (~97%). After rebounding from the cavity wall, the plasma, now moving much more slowly, continues to radiate to the jets so that by 3 ms 50% of the total initial energy has gone into the jets.

The pressure exerted on the cavity wall, P_{gw} , is strongly influenced by the heat transfer. Again, for the same jet diameter, the initial peak pressure is seen in Fig. 9-29 to be less by a third and subsequent spikes are very much more attenuated. If the cavity wall is to remain elastic under load, the maximum hoop stress is easily obtained by integrating the equations of motion for an internally pressurized cylinder using $P_{gw}(t)$ as the forcing function.⁴⁴ The results are shown in Table 9-12. With maximum heat transfer to the jets, and for the same jet diameter, the peak hoop stress is less than half the value obtained without heat transfer.

The temperature of the gas adjacent to the wall is shown in Fig. 9-30. On first contact, a level of just under 4 eV is attained, which halves with the first rebound and remains more or less at 2 eV for the heat-transfer case. An almost-stagnant layer of hot, low-density gas then rests against the cavity wall, which will require a liquid lithium layer between the gas and wall for thermal insulation. The thickness of this layer is determined by the time required for the main liquid fall to move through the gas and splash up against the wall. In Ref. 45 we have estimated this time, τ , to be of the order of 100 ms; taking α , the thermal diffusivity of lithium, to be $0.0242 \text{ mm}^2/\text{ms}$; the layer thickness required $\sim \sqrt{\alpha\tau} = 1.6 \text{ mm}$. A layer several times this dimension was envisioned in any case to protect against erosion damage from long-term, fragmentary liquid impact.

Size Scaling, Opacity, and the Emissivity.

For the same lithium fall mass, uniformly decreasing the jets' diameter by two orders of magnitude, from 200 to 2 mm, increases the energy transport from the plasma to the jets from 50% to over 80% (Fig. 9-28). Also, peak temperature (Fig. 9-30) and especially the pressure at the cavity wall (Fig. 9-29) and the peak hoop stress (Table 9-12) are significantly decreased. If radiation is the dominant mechanism of heat transfer, however, these results are seemingly anomalous. In our model, we have taken the emissivity to be proportional to the hydraulic radius and hence $\sim D$. According to Eq. (17), this means that the heat-transfer rate per unit area is then $\sim D$. But, for fixed jet mass, the surface area $\sim D^{-1}$ so that the overall heat-transfer rate should be invariant with D . The explanation lies in the inequality sign in the emissivity definition (4). When K is sufficiently large, so that the radiation mean-free path is small compared with the characteristic beam length, the emissivity must be constrained, otherwise the emission would exceed that of a perfect radiator. The flux limiter in Ref. 45 is then merely an artifice to account for the fact that the emission approximation is no longer strictly valid. A more elaborate treatment of the gas-particle radiation transport is precluded here because of the complexity and increased computational costs.

It is apparent in Figs. 9-27 to 9-30 that the flux limiter has been called upon for much of

the calculation. This allows an alternate interpretation of the curves labeled (b) in the figures. Since the heat transfer depends on the quotient of the emissivity by particle size, the solution obtained for 200-mm D jets and $\epsilon = 1$ is approximately the same as that for 2-mm D jets and $\epsilon = 10^{-2}$ (the two solutions would actually be slightly different because viscous drag and convection influence the latter a bit more than the former).

In most of our calculations, Planck mean opacities were used for the gas-particle heat transfer. In the range of temperature and density of interest, these data are typically higher than Rosseland mean values at the same conditions by a factor of 10 (see Ref. 44). The dashed curve in Fig. 9-29 shows the effect on the pressure exerted on the cavity

Fig. 9-34. Energy transport from the plasma to the jets as a function of time for the three cases of combined heat and mass transfer.

Fig. 9-35. The pressure-time history on the cavity wall for the three cases of combined heat and mass transfer.

Fig. 9-36. The temperature-time history for the plasma adjacent to the cavity wall for the three cases of combined heat and mass transfer.

wall of substituting Rosseland for Planck data when $D = 2$ mm. With the lower opacities, the peak wall pressure is increased by 18% and the wall pressure at 3 ms from 0.1 to 0.65 MPa. The peak hoop stress (Table 9-12) shows an increase of 14%. Using Rosseland opacities for the $D = 200$ mm case gives virtually no difference from the curves labeled (b) because, in this case, the increased beam length invokes the flux limiter in the temperature-density range of interest.

The Case of Maximum Mass Transfer. Setting $\zeta(r,t) = 0$ (equivalent to $T_v = T_s^0$) maximizes the mass transfer from the jets to the plasma. All of the heating potential of the plasma goes toward vaporizing the liquid surface. The mass generated thereby is returned to and mixed with the plasma so that the total energy of the plasma does not change. The added mass will reduce the average temperature, but this will only affect the equilibrium pressure if the average specific energy falls below the cohesive energy or near the ionization potentials. The peak pressure on the cavity wall is, however, much less than if there were no heat or mass transfer. Figure 9-31 contrasts the pressure history at the cavity wall for the limiting cases of maximum heat [case (c)] and maximum mass [case (d)] transfer, with the jet diameter fixed at 2 mm. With maximum mass transfer, the peak value of P_{gw} is reduced by a third in comparison with the maximum heat-transfer case, but Table 9-12 shows that the peak hoop stress is increased by about the same amount. The reason is that the loading time is only about half the natural period of the wall (6.2 ms for the

5-m-radius steel wall) so that the impulse delivered (proportional to the integral under the curves in Fig. 9-31), rather than the peak pressure, determines the maximum hoop stress.

In Fig. 9-32 we see that with maximum mass transfer from the jets, the plasma mass in the cavity at $t = 3$ ms is more than four times the initial mass, most of the mass transfer having occurred in the first 800 μ s. This is the period in which the main body of plasma has passed through the jets twice, once on the way out and once on the rebound. Figure 9-33 shows that in the initial pass, the added jet vapor has resulted in a much lower gas temperature adjacent to the cavity wall than in the maximum heat-transfer case. The initial temperature jump at the wall results from the impact of the plasma, which converts its kinetic energy to internal energy. The temperature thereby attained is roughly proportional to the specific kinetic energy and hence to the square of the velocity. The mass added by the vaporized jet surface is assumed to leave the surface at the local jet condition, and in particular, with zero velocity. Having no momentum of its own, the vapor then mixes with the local plasma and lowers its velocity so that when contact with the wall is finally made, the temperature jump is much lower than would otherwise be the case.

The temperature drops faster however when the heat transfer is maximized so that by $t = 3$ ms the temperature at the wall is lower for case (c). As seen in Fig. 9-27, by this time the total kinetic energy in the plasma is negligible in comparison with the internal energy and mechanical equilibrium is therefore established. With case (c), Fig. 9-28 shows that the residual energy in the plasma is $0.19 E_T^0$, so that the average specific energy is $0.19 E_T^0/m^0$. On the other hand, as we have noted in Fig. 9-32, the "equilibrium" mass for case (d) is $4.03 m^0$ so that the average specific energy for this case is $E_T^0/4.03 m^0 = 0.25 E_T^0/m^0$ or 24% higher than for case (c). Again referring to Fig. 9-33, we see that T_{gw} is $\sim 21\%$ higher for case (d), in fairly good agreement with the average specific energy argument.

The Case of Combined Heat and Mass Transfer. When ζ is allowed to vary in accordance with Eq. (30), heat diffuses unimpeded into the jets interior ($\zeta = 1$) only as long as the local surface temperature

Table 9-12. Peak hoop stress in first wall for cases described in Table 9-11.

Case I/D	$(\sigma_{\theta} \delta_w)_{\max}^a$ (MPa-m)
(a)	57.9
(b)	25.8
(c)	12.1 ^b
(d)	17.0
(e)	31.9
(f)	25.8
(g)	16.0

^aSolutions to the equation of motion are of the form $\sigma_{\theta} \delta_w = f[P_{gw}(t)]$, where δ_w is the wall thickness. The wall was assumed to be made of steel.

^b13.8, when the dashed pressure-time history in Fig. 9-29 was used (based on replacing Rosseland- with Planck-mean opacities).

remains below the boiling point. Thereafter the rate of energy loss from the plasma is controlled by the jet surface temperature; increased radiation flow from the plasma beyond that required to maintain the surface at boiling merely results in increased surface vaporization ($0 \leq \zeta < 1$). When the radiation flow from the plasma recedes and is no longer able to support boiling, the surface temperature falls and vaporization ceases ($\zeta = 1$).

Figure 9-32 compares the cavity mass ratio for three cases of combined heat and mass transfer ($D = 200, 20$, and 2 mm) with our previous result obtained with $\zeta = 0$ everywhere in the cavity. We see that there is relatively little variation for the three combined cases; at 3 ms m/m^0 varies from 1.45 for $D = 2$ mm to 1.61 for $D = 200$ mm, as compared with 4.03 with $\zeta = 0$ and $D = 2$ mm.

After the first pass through the jets by the plasma, the combined results exhibit a curious inversion, more vapor being generated with the larger jet diameters. The inversion occurs after the main body of plasma has passed through the jets and rebounded. Figure 9-34 shows that a jump in interior jet heating also occurs at this point. Both are connected with the contact made between the main plasma and the outer rows of jets upon reentry.

In comparing the heat lost to the jets with variable ζ (Fig. 9-34) to the maximum heat-transfer case (Fig. 9-28), we see that only half as much heat is lost when surface vaporization is taken into account and jet disassembly is minimal ($D = 200$ mm). On the other hand, when jet disassembly is widespread, the heat loss to the jets is much less affected, the two cases (c) and (g) differing only by about 15%.

The pressure exerted on the cavity wall with combined heat and mass transfer is depicted in Fig. 9-35 and the resulting maximum hoop stress is given in Table 9-12. For the same jet particle size, the combined heat and mass transfer model results in wall stress about 20 to 25% higher than in the "best possible" case (i.e., maximum heat transfer). Moreover, comparing Fig. 9-30 with Fig. 9-36 shows that, at most, the gas temperature at the wall with the combined model exceeds that obtained with maximum heat transfer by about 10%.

The Impulse Delivered to the Jets. Figure 9-37 shows the total impulse delivered to the jets by the plasma as a function of time for all cases considered here. The data are normalized to the no-heat-or-mass-transfer case (a) at $t = 360 \mu\text{s}$ for comparison with Ref. 45. The impulse monotonically increases until the main body of plasma rebounds through the jets, with each subsequent rebound providing small perturbations. At 3 ms the combined heat and mass transfer cases (e), (f) and (g) give virtually identical results and these agree as well with the solution obtained without heat or mass transfer. The impulse is about 20% higher than reported earlier. With maximum heat transfer, the results are also more or less independent of jet size in the range considered and the impulse is about 25% lower than reported earlier. The maximum mass transfer case, however, shows a huge 250% increase as a result of the "source drag" induced by the large increase in system mass (cf. Fig. 9-32).

A Note on Condensation. We have seen that mechanical equilibrium is attained by the plasma in the cavity by ~ 3 ms from implosion time. A simplified physical picture of the cavity at this point is that of a nearly stagnant and isothermal plasma annulus, with temperature, T^0 , between 1 and 2 eV and pressure, P^0 , of perhaps 1 to 2 MPa, resting between the outer row of jets and the insulating liquid adjacent to the cavity wall. There is also a much cooler layer of lithium gas entrained in the region of the jets. Under the combined influence of plasma drag and neutron heating the jets, and the

Fig. 9-37. Total impulse delivered to the jets by the plasma, I_T , as a function of time normalized to I_T^* , the result obtained for the no-heat-or-mass transfer case at $360 \mu\text{s}$.

disassembled parts thereof, move radially outward. The initial spatial energy distribution from the fusion pellet causes the inner row of jets to move more rapidly than the outer so that multiple collisions occur and by ~ 15 ms from implosion time, the jets are more or less compacted into a near-normal density annular packet moving outward at a uniform velocity of perhaps 30 to 50 m/s.⁴⁵ The inner and outer radii of the packet at this point are approximately 1.8 and 2.3 m, respectively. Since its velocity is uniform throughout, the annular packet will maintain uniform thickness as it coasts toward the wall, which in the present design is placed at a radius of 5 m. Continuity demands that as the field diverges the density must drop, so that we expect that the liquid will once again come apart into fragments or drops. As the fragments pass through the plasma, the latter will be isochorically cooled until the boundary of the (liquid-vapor) coexistence region is reached, at which point condensation will begin. Further cooling will take place along the liquid-vapor saturation curve (significant supercooling should be inhibited by the presence of non-condensibles such as He and T, which will act as nucleation sites) so that condensation will continue at least until the leading edge of liquid fragments impacts the wall. We wish to estimate when condensation will first begin and to what extent it will have progressed by wall impact time. Cooling will occur by radiation and forced convection, but once the temperature of the plasma

drops much below 1 eV, convection will dominate. Consistent with the very much simplified physical picture described above, we will ignore spatial gradients and adopt a lumped-parameter approach. It is then easy to show that the plasma-cooling rate is

$$\frac{dT}{dt} = -f\phi/mC_v, \quad (32)$$

where m is the plasma mass, C_v its heat capacity, ϕ is the total particle surface $= n\pi D^0(D^0/D)$ and $f = h(T - T_s)$. n is the initial number of jets of diameter D^0 ; these are assumed to have uniformly subdivided into cylindrical fragments of diameter D . T_s is defined by Eq. (28), which implies $\zeta = 1$. (We expected the fragment heating rate to be sufficiently low so that $T_s(t) < T_v$, even for the largest fragments under consideration; this was verified by the calculated results.) The heat-transfer coefficient, h , was based on Hilpert's experiments with cylinders in air crossflows,⁵⁰ corrected for Prandtl number. All thermophysical data derived from our equation of state (EOS) for lithium⁴⁴; e.g., $C_v = C_v(\rho, T)$.

For all $D \leq D^0$ and $T \leq T^0$, the calculated flow was laminar. It can then be readily demonstrated that $f \sim D^{-1/2}$ so that the rate at which the plasma cools by convection $\sim D^{-3/2}$. Figure 9-38 shows the results of three calculations with $D = 200$ mm (D^0), 20 mm and 2 mm. The diameter scaling appears to work quite well.

The assumptions made in deriving Fig. 9-38 imply that all of the plasma is in contact with all of the jet fragments, a situation not attained until the fragments have arrived near the wall. At 50 ms, the characteristic time for the leading edge of the fragments (at $r = 2.3$ m) to arrive at the wall (at $r = 5$ m) is $\zeta = (5 - 2.3)/50$ or 54 ms. To interpret Fig. 9-38: this represents an approximate time scale for which the solution is meaningful.

The vacuum condition required for laser beam propagation in the cavity is on the order of 0.1 Torr (13.3 Pa)—see Ref. 51, corresponding to a saturation temperature of 893 K. For $T^0 = 2$ eV (2.321×10^4 K), $T_s^0 = 750$ K and $T = 893$ K, so that $\theta = (T - T_s^0)/(T^0 - T_s^0) = 6.36 \times 10^{-3}$. Fig. 9-38 gives $t(D^0/D)^{3/2} = 1.2$ s. If we substitute the wall arrival time τ for t , and put $D^0 = 200$ mm, we find that $D = 25.3$ mm. Thus, if the average fragment size is less than ~ 25 mm (a not unlikely condition), it seems reasonable

Fig. 9-38. Calculated average plasma temperature in the space between the jets and the cavity wall as a function of scaled time.

to expect that the plasma will have cooled and condensed to the beam propagation limit in the 54 ms it takes for the jet fragments to arrive at the cavity wall.

We have assumed that $T^0 = 2$ eV and that $P^0 = 2$ MPa, which—according to our lithium EOS—gives a plasma density of 2.87×10^{-5} Mg/m³ (note that this is a factor of 2.5 less than would be the case for an ideal gas, and neglecting ionization). If the plasma cools isochorically to the saturation curve, it can be shown that the latter is intercepted at a pressure of 393 Torr (5.24×10^4 Pa), corresponding to a saturation temperature of 1525 K. Substituting this temperature in Fig. 9-38, $\theta = 3.45 \times 10^{-2}$, so that $t(D^0/D)^{3/2} = 0.75$ s. With $D = 25$ mm, we then find that the time for condensation to first begin is ~ 33 ms.

Recognizing that the above analysis is rather crude, it nevertheless appears that the cooling and condensation of the plasma occurs an order of magnitude faster than is required for the envisioned pulse repetition rate of 1 Hz.

Summary and Conclusion. The quasi-one-dimensional method previously derived for calculating transient, compressible, viscous flows through a complex array of jets was extended to include heat and mass exchange with the jets.

The extended method was applied to the fusion reactor problem in which a lithium plasma is initially contained within the cylindrical core formed by a close-packed annular arrangement of liquid lithium jets. When the plasma is released, it flows around or in between the jets and emerges to eventually impact on the (first) wall of the reactor. Our main object was to determine the impulse imparted by the plasma to the jets and to calculate the stress induced in the wall due to plasma impact. A secondary objective was to estimate the effectiveness of the jets, and disassembled fragments thereof, in cooling and eventually condensing the plasma.

We found that more than 70% of the initial energy contained in the plasma will diffuse into the liquid jets, not including an additional 7 to 10% that will go towards jet surface vaporization. This vaporization will initially increase the mass of gas in the cavity 40 to 60%. The drag impulse delivered to the jets will exceed that previously estimated (neglecting heat and mass transfer)

by about 20%, but this is of no great consequence, since the drag impulse on the jets represents only a small fraction of the (neutron-induced) total.

Depending on conditions in the cavity, the peak hoop stress in the wall due to plasma impact was calculated to be up to four times less when heat and mass transfer were accounted for. One of the most important parameters here is the jet/fragment size and spatial distribution in the cavity. Up to 50% of the total energy transport to the jets occurs in the first row, which has completely disassembled via neutron heating even before contact has been made with the plasma. The spatial and particle size distribution can only be guessed. If the particle size distribution remains more or less uniform and the cross section blocked by the jets is invariant, we found that changing the particle size by a factor of 100 resulted in a factor-of-2 change in the peak stress. Also important are the opacity and emissivity of the plasma since most of the energy transport is by radiation. Varying the emissivity by a factor of 100 also resulted in a factor-of-2 change in the peak stress. It should be noted here that the computed magnitude of the peak stress in these cases is of less significance than the comparison of one case with another, since axial flow and pressure relief have been neglected. The energy density per unit length at the plane of the fusion pellet (used in our calculations to represent the “worst-case” performance) is approximately 10 times the overall-average-energy density in the cavity. Our results then only provide an upper bound on the true stress; a more accurate estimate will require higher dimensional calculations, which we are in the process of making.

Cooling and eventual condensation of the plasma on the lithium jet fragments is governed mainly by forced convection. We found that as long as the average jet fragment size is less than ~ 25 mm, the plasma will cool and condense to the (laser) beam propagation limit in the time it takes for the jet fragments to arrive at the cavity wall, this time being at least an order of magnitude less than required for the envisioned pulse repetition rate of 1 Hz.

Clearly, all of the above results rest on many assumptions, some of which cannot be justified without experiment. The relevant experiments (e.g., measurement of the

particle size and spatial distribution derived from isochorically heated liquid lithium as a function of specific energy) are difficult and costly and, thus far, there has been insufficient support for their undertaking. The quasi-one-dimensional method developed here has allowed us to establish upper bounds, useful for design purposes, and also to determine the sensitivity of our design to key parameters.

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Vortex-Flow Slanted-Jet Improvement to HYLIFE. A problem plaguing the HYLIFE design has been the formation of high-energy gas that causes a high pressure on the first structural wall (FSW) and generates a significant outward radial momentum of the jet array. A vortex-flow modification diminishes both of these effects.

The basic idea of the vortex-flow modification is to slant the passages through the jet array so that an angular velocity is imparted both to the gas blown off from the jets and to the jets themselves. Previous HYLIFE liquid metal wall designs incorporated concentric annuli, an array of circular cross-section jets, or an array of rectangular cross-section jets. Figure 9-39 shows a profile of the desired jet array just before a fusion pulse. The black gridwork

overlaid on the previous design shows the metal rib section of the nozzle plate used to form the outer rows of jets. The red areas show the desired jet profile at the midplane.

The shape of the inner row of jets is such that the left side of each jet is exposed to x-ray and neutron fluence at a slant angle. The right side of each jet is either radial or parallel to the left side, as shown. The x-ray-ablated material moves normal to the jet surface, reducing the radial-velocity component of the gas. About 80% of the x-ray-ablated material impacts on the adjacent jets, boils off more mass, and, thus, decreases the specific energy. Gas from each jet interacts with gas from adjacent jets to form a vortex flow; hence, stagnation pressure at the vertical centerline of the chamber is reduced, and the transit time for the gas to pass outward through the jets is increased. Cooling of the gas in contact with the jets is enhanced both by the increase in contact time and by the increase in liquid surface area with time due to neutron-induced disassembly. Gas stagnation pressure on the FSW is decreased by both the reduced initial specific energy and the increased cooling.

The inner row of jets replaces the weir flow of the previous design; consequently, some of the webs on the nozzle plate will not be adequately shielded from neutrons, and these webs may have to be replaced periodically. The shape of the inner row of jets is important because this is where essentially all of the x-ray and debris energy, and a large portion of the neutron energy is deposited. Neutron-induced disassembly of the inner jets will also be nonradial, leading to the potential for more dissipation than would arise from the concentric annuli used in current calculations. The corners of all jets will be rounded due to surface tension.

Circular or rectangular cross-section jets also improve dissipation of momentum, as compared to concentric annuli. However, with circular or rectangular cross-section jets, the angular momentum in the counterclockwise direction equals that in the clockwise direction so that, as the gases interact, the momentum not lost in dissipation is converted to radial momentum. The advantage of the proposed vortex-flow slanted-jet design over circular or rectangular cross-section jets is that the angular velocity imparted to the gas is

Fig. 9-39. Vortex-flow slanted-jet pattern overlaid on the HYLIFE nozzle plate.

mostly in the same direction, and the resulting momentum, after dissipation, is in both radial and circumferential directions. Only the radial momentum is of concern, and its magnitude is reduced if the proposed design replaces the circular or rectangular cross-section jets of the current design.

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Heavy-Ion-Beam Ballistic Propagation in Liquid-Metal-Wall Reaction Chambers.⁵²

For a heavy-ion-beam driven, liquid metal wall (LMW) ICF reaction chamber, the liquid metal composition and temperature must not preclude transport and focus of the heavy-ion beam on the target. For a near-vacuum system ($<10^{14} \text{ cm}^{-3}$), both the vapor pressure and the stripping effectiveness of the chamber gas must be known before it can be determined if ballistic propagation of the beam is possible. These conditions are discussed in this article for two LMWs: Li contaminated with 1 atomic percent (at.%) Pb, which equals 23 weight percent (wt%) Pb, and the Pb-Li eutectic containing 83 at.% (99.3 wt%) Pb. For a higher-pressure system, other beam-propagation modes are under investigation but are not discussed here.

The vapor pressure above a Pb-Li solution is

$$P = \gamma_{\text{Li}} N_{\text{Li}} P_{\text{Li}} + \gamma_{\text{Pb}} N_{\text{Pb}} P_{\text{Pb}} \quad (33)$$

where N_{Li} and N_{Pb} are the atom fractions of Li and Pb in the liquid, and P_{Li} and P_{Pb} are the vapor pressures of the pure elements at the given temperature. The temperature- and composition-dependent functions, γ_{Li} and γ_{Pb} , are activity coefficients; they describe the ideality of the solution. If the functions are 1.0, the solution is "ideal" (Raoult's Law); if they are greater than 1.0, the solution has "positive deviation from ideality" and a higher-than-expected vapor pressure. In the case of the Pb-Li system, experimental data⁵³⁻⁵⁵ indicate a profound negative deviation from ideality, probably due to conversion from metallic to ionic bonding.

The vapor pressures of pure Pb and Li are shown as a function of temperature in Fig.

9-40 and Tables 9-13 and 9-14. The vapor pressure of pure Pb is several hundred times less than that of pure Li in the 400 to 500°C temperature range. (Higher temperatures are discouraged by a decrease in mechanical properties of ferritic steel and by Li corrosion; lower temperatures are discouraged by brittleness of ferritic steel and by thermal efficiency considerations.)

For $\text{Li}_{99}\text{Pb}_1$, the Li activity coefficient, γ_{Li} , is about unity. However, data^{53,60} at 527 and 659°C indicate that the Pb activity coefficient is less than 2×10^{-6} , so that the

Fig. 9-40. Vapor pressures of pure Pb, pure Li, and $\text{Pb}_{83}\text{Li}_{17}$ mixture.

Table 9-13. Range of vapor pressures (in Torr) in the literature for pure lithium. ▼

Source	Temperature (°C)	
	400	500
Cowles and Pasternak (Ref. 56)	1.026 10^{-4}	3.56 10^{-3}
Nesmeyanov (Ref. 57)	9.11 10^{-5}	3.26 10^{-3}
Hultgren (Ref. 58)	8.60 10^{-5}	3.18 10^{-3}
JANAF (Ref. 59)	8.45 10^{-5}	2.95 10^{-3}
Fit to av of Refs. 57-59	8.57 10^{-5}	3.02 10^{-3}

Notes:

1. An equation to fit to the Douglas curve (Fig. 6 of Ref. 56) was used in this article:

$$\ln P_{\text{Li}} (\text{Torr}) = 18.221 - \frac{18,444}{T(\text{K})}$$

2. The data from Hultgren, Nesmeyanov, and JANAF were averaged at several temperatures between 327 and 627°C. An equation fit to the averages is

$$\ln P_{\text{Li}} (\text{Torr}) = 18.164 - \frac{18,526}{T(\text{K})}$$

	Temperature (°C)	
	400	500
Nesmeyanov (Ref. 57)	1.928 10^{-7}	1.504 10^{-5}
JANAF (Ref. 59)	1.995 10^{-7}	1.57 10^{-5}
Hultgren (Ref. 58)	2.38 10^{-7}	1.75 10^{-5}
Fit to av of Refs. 57-59	2.09 10^{-7}	1.60 10^{-5}

Note: An equation fit to the average of the three references at several temperatures between 400 and 550°C is

$$\ln P_{\text{Pb}} (\text{Torr}) = 18.14 - \frac{22,560}{T(\text{K})}$$

Table 9-14. Range of vapor pressures (in Torr) in the literature for pure lead.

Pb-component partial pressure is 10 orders of magnitude less than that of the Li component for $\text{Li}_{99}\text{Pb}_1$.

For high-concentration Pb solutions, such as $\text{Pb}_{83}\text{Li}_{17}$, the activity coefficients must be carefully evaluated. Demidov's γ_{Pb} data⁶⁰ at 527°C was used to determine activity-coefficient values, which ranged from 0.932 at 77 at.% Pb to 1.0 above 88.3 at.% Pb. If γ_{Li} is of order 1, the Li-vapor component would dominate the Pb vapor by a factor of 50. However, as we show below, γ_{Li} is less than 10^{-3} at temperatures below 500°C, and the Pb component dominates the vapor.

The Li-component activity coefficient, γ_{Li} , can be determined from electromotive force (emf) data taken by Saboungi⁵³ at temperatures of 497, 539, 596, and 659°C in Pb-Li mixtures with Pb atom fractions above 45%. In Fig. 9-41, γ_{Li} is shown as a function of $(1/T)$ at five liquid compositions. The lines in Fig. 9-41 are least-square fits to allow extrapolation to the temperature region of

Fig. 9-41. The ratio of actual to ideal Li partial pressure above Pb-Li liquids.

interest to ICF reactor designers, 400 to 500°C.

Since the emf data did not involve direct pressure measurements, we searched for experimental verification involving either a direct or indirect pressure reading of Li and Pb above a binary solution. Knudsen-cell work by Ihle, Neubert, and Wu⁵⁴ provide data that are also shown in Fig. 9-41. Later Knudsen-cell work by Neubert, Ihle, and Gingerich⁵⁵ is shown in Fig. 9-42, which plots the Pb and Li partial pressures above $\text{Pb}_{50}\text{Li}_{50}$ liquid as a function of $(1/T)$. Figure 9-42 also shows the earlier Knudsen-cell work for Li, the emf data, and least-square fits to the emf data. The emf data for Pb uses, at all temperatures, the γ_{Pb} value of 0.4 shown in Fig. 9-42 of Ref. 53 for a temperature of 659°C. If a value of 1.0 is used for γ_{Pb} at all temperatures, the emf extrapolation falls in the middle of the higher-temperature Knudsen-cell Pb data.

The total vapor pressure over 497°C Pb-Li liquid is shown in Fig. 9-43 as a function of Pb concentration. The pressure above $\text{Pb}_{83}\text{Li}_{17}$ (1.2×10^{-5} Torr) is a factor of ~ 50 lower than ideal and a factor of ~ 270 lower than the pressure of pure Li (3.2×10^{-3} Torr) at the same temperature. The vapor is almost pure Pb (~ 96 at.% Pb) compared to almost pure Li (~ 2 at.% Pb) for an ideal solution.

The vapor above a Pb Li melt could interfere with heavy-ion-beam transport. Total pressure within the reactor cavity is not a definitive parameter, since beam transport is sensitive to the atomic numbers of the elements in the chamber vapor. Beam-propagation calculations for the case of a $\text{Li}_{99}\text{Pb}_1$ LMW vs a $\text{Pb}_{83}\text{Li}_{17}$ LMW are required to determine allowable LMW temperatures and, hence, reactor efficiency. In each case, the vapor species are Li, Li_2 , Pb, Pb_2 , LiPb, He, hydrogen-isotope molecules, and gaseous lithium hydride (including deuteride and tritide). The significant species are Li, Pb, He, and hydrogen-isotope molecules.

For the $\text{Li}_{99}\text{Pb}_1$ LMW, the total chamber pressure ranges between 2×10^{-4} and 7×10^{-3} Torr as liquid temperature increases from 400 to 500°C (Fig. 9-44). Monatomic Li is the primary condensible constituent. Diatomic Li partial pressure is small compared to monatomic Li pressure. The Pb

partial pressure is at least a factor of 600 000 lower than Raoult's law predicts due to the nonideality of Li-Pb solutions; hence, the Pb partial pressure is negligible. If the hydrogen-isotope concentration in the liquid metal is maintained below one weight part per million (wppm) by chemical processes outside the chamber, and if the time for gaseous H_2 to dissolve into the liquid metal is much less than the interpulse time, then the hydrogen-isotope partial pressure is also negligible. The He partial pressure depends on both the He generation rate and the vacuum-system design. Helium partial pressure cannot be easily maintained much below the Li vapor pressure since, as the chamber is pumped to remove He, more Li vaporizes. The assumption used in Fig. 9-44 (He partial pressure equals the sum of the Li, Pb, and H_2 partial pressures) should be tested by vacuum-pump designers. Although lower pressures are unlikely, considerably higher He pressures could well be required by vacuum-engineering constraints. (See article on vacuum pumps in this section.)

For the $Pb_{83}Li_{17}$ LMW, the total chamber pressure ranges between 3.5×10^{-7} and 2.7×10^{-5} Torr as temperature increases from 400 to 500°C (Fig. 9-45). The Li partial pressure is negligible; the H_2 partial pressure cannot be specified, as discussed in "Solubility of Hydrogen Isotopes in Pb-Li Solutions," above, but two alternatives are plotted. The lines labeled Case 1 in Fig. 9-45 correspond to high hydrogen solubility. In this case, H_2 pressure is negligible, as it is with the other fluid. Helium pressure is then also treated as before; i.e., it is set equal to the total of the Pb, Li, and H_2 pressures. The lines labeled Case 2 in Fig. 9-45 correspond to low hydrogen solubility. In this case, H_2 is removed from the chamber through the vacuum system, and the sum of the H_2 and He pressures are set equal to the sum of the Pb and Li pressures. Again, vacuum-system design is required. In either case, H_2 solubility has almost no effect on the total chamber pressure.

To compare the effects of the two liquid metals on heavy-ion-beam transport, a method of converting the pressure of one element to an effective pressure of another element is needed. According to Gillespie,⁶¹ Pb vapor pressure (and pressures of other elements) can be converted to equivalent Li vapor pressure if it is assumed that stripping

of +1 ions in a heavy-ion beam to a higher ionization state is the main beam-degradation effect. When Gillespie's multipliers of 1/3, 4/9, 1/2, and 113 are applied to the partial pressures of H_2 , He, diatomic Li, and Pb, respectively, each pressure is converted to equivalent (monatomic) Li pressure. (Uranium was used as the heavy-ion beam in these calculations, although a nonfissionable ion

Fig. 9-42. Partial pressures of Li and Pb above $Pb_{50}Li_{50}$ liquid.

Fig. 9-43. Vapor pressures of Pb-Li mixtures.

will almost certainly be the eventual choice for a heavy-ion-beam fusion reactor.) Fig. 9-46(a) shows the results for both fluids. For

Fig. 9-44. Partial pressures of vapor constituents above $\text{Li}_{99}\text{Pb}_1$ liquid.

Fig. 9-45. Partial pressures of vapor constituents above $\text{Pb}_{83}\text{Li}_{17}$ liquid.

Fig. 9-46. (a) Equivalent Li pressures for heavy-ion-beam stripping above $\text{Li}_{99}\text{Pb}_1$ and $\text{Pb}_{83}\text{Li}_{17}$ liquids. (b) Maximum allowable LMW temperature vs stripping effectiveness of Pb. ▼

$\text{Li}_{99}\text{Pb}_1$, Li and He dominate the equivalent pressure for stripping; these components are shown with dashed lines. Two cases are shown for $\text{Pb}_{83}\text{Li}_{17}$; the lower line uses the minimum 10^{-7} to 10^{-5} He pressure set by the Pb-Li vapor pressure, and the upper line uses the 10^{-4} to 10^{-3} He pressure calculated for the other fluid. The higher He pressure seems more likely to be based on the vacuum-pumping design considerations discussed in "Vacuum Pumping," below.

Figure 9-46(a) can be used to compare three cases: $\text{Li}_{99}\text{Pb}_1$; $\text{Pb}_{83}\text{Li}_{17}$ with 10^{-4} to 10^{-3} Torr of He; and $\text{Pb}_{83}\text{Li}_{17}$ with 10^{-7} to 10^{-5} Torr of He. To make the comparison, draw a horizontal line at the maximum allowable Li pressure. (This value will eventually be provided by the beam-propagation community; current estimates range from 10^{-4} to 10^{-3} Torr.) Then compare the liquid-metal temperatures for the three cases. For example, at 5×10^{-4} Torr allowable, the maximum liquid-metal temperatures are 431, 450, and 472°C for the three cases.

Recently, it has been suggested⁶² that beam stripping might be significantly decreased if the beam ion and background gas species were identical. If so, Pb would make an excellent candidate for a heavy-ion beam. As shown in Fig. 9-43, the vapor above $\text{Pb}_{83}\text{Li}_{17}$ is ~96 at.% Pb (considering only the Pb and Li). If the stripping effectiveness of Pb is less than 100, He will be

the dominant component of the background gas with respect to beam propagation. In addition, a ^{208}Pb heavy-ion beam has high mass; Pb may also be a target material, and Pb has a very low neutron-activation probability. However, isotope separation will be required to separate ^{208}Pb from natural Pb.

The Pb stripping multiplier of 113 that was used in Fig. 9-46(a) is the maximum predicted theoretically, and a value of one-half that is not unlikely, according to the theoreticians. However, German experimental work⁶³ using U^{+8} ions interacting with nitrogen background gas indicates that the theoretical maximum cross section predictions may be too low. If these surprising experimental results are confirmed, the Pb stripping multiplier for U^{+1} could be as high as 400. The data of Fig. 9-46(a) are cross plotted in Fig. 9-46(b) as a function of the Pb stripping multiplier. Clearly, the allowable liquid-metal temperature is most sensitive to the allowable equivalent Li vapor pressure, and the plotted range of 2×10^{-4} to 10^{-3} Torr is less than one-half the range of estimates made by the beam-propagation community. The Pb stripping multiplier is also very uncertain, but it is of secondary importance in determining the maximum allowable temperatures for the two liquid metals.

We can compare $\text{Pb}_{83}\text{Li}_{17}$ with $\text{Li}_{99}\text{Pb}_1$ by using the difference in allowable temperature based on heavy-ion-beam propagation. If the $\text{Pb}_{83}\text{Li}_{17}$ has a 30°C higher liquid temperature, the net plant efficiency is increased by about 0.5%, and about 20 MWe additional power is produced.⁶⁴ In addition, the higher neutron-energy multiplication associated with Pb-Li could produce about 150 MWe additional power in a 1000 MWe-net power plant.⁶⁵ (The ratio of thermal to fusion power for an 80 volume percent (vol %) Pb_{4}Li - 20 vol % steel LMW is 1.29, by a one-dimensional calculation, as compared with 1.16 for the original HYLIFE design, by a two-dimensional calculation.) However, the higher density and lower heat capacity of $\text{Pb}_{83}\text{Li}_{17}$ require about 130 MWe additional pumping power for the LMW⁶⁶ and about 100 MWe pumping power for the primary loop.⁶⁷ Hence, from a rough power-flow viewpoint, the use of $\text{Pb}_{83}\text{Li}_{17}$ results in a loss of 60 MWe net power. A full evaluation of $\text{Pb}_{83}\text{Li}_{17}$ must also consider, on the negative side, increased flow-loop structural

requirements, increased fluid radioactivity, plating of activated Pb on plant surfaces after a spill, and larger primary heat exchangers and, on the positive side, additional recovered pump heat and elimination of liquid-metal fire hazards. At this point, both fluids are of interest to designers of LMW fusion reactors.

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Experimental Plan to Measure Effects of Gas Blowing Through Jets. The HYLIFE reaction chamber has an annular array of liquid lithium jets that serve as a heat-exchange medium for the energy of the D-T reactions, that provide a tritium-breeding mechanism, and that protect the structure. During a reactor pulse, x rays and pellet debris generate a hot, high-pressure lithium plasma at the inner radius of the jets that first implodes and then flows radially outward in and around the jets toward the FSW.

A quasi-one-dimensional analytical model and a computer code have been developed to describe plasma flow through an array of jets and the subsequent impulse on the reaction-chamber FSW.⁶⁸ Complete experimental verification of the code would be desirable, but generation of the high temperatures and energies involved, as well as the use of liquid lithium, make such an experiment complex, time consuming, and expensive. A major portion of the code, however, can be easily verified by experiments in the regime of low-energy flow over simulated jets, where little disturbance would be imparted to the jets in the form of radiative transfer, erosion, or vaporization. This section describes the design of such an experiment as it would be run in the Grumman research shock tunnel.⁶⁹

The schematic in Fig. 9-47 shows the main features of the Grumman facility. Argon is chosen as the test gas because its ratio of specific heats ($\gamma = 5/3$) is very close to that of the lithium vapor in the HYLIFE reaction chamber; furthermore, argon is ideally suited for use in shock-tube facilities because the gas has high molecular weight and is inert and nontoxic. To avoid instrumentation

Fig. 9-47. Schematic of shock tunnel and HYLIFE test nozzle.

Fig. 9-48. Reflected stagnation temperature in the shock-tube. ▶

Fig. 9-49. Reflected stagnation pressure in the shock-tube.

problems with ionized flows, stagnation temperatures in Region 5 (T_5) in Fig. 9-47 are held below 8000 K. Since the shock tunnel operates in the reflected-shock mode, the incident-shock Mach number (M_s) should not exceed approximately six. Figure 9-48 illustrates the variation of T_5 as a function of M_s , and Fig. 9-49 shows the variation of stagnation pressure (P_5) as a function of M_s . Note that T_5 is independent of the initial prerun pressure (P_1) in the tube, because argon behaves as an ideal gas in the temperature range of interest, while P_5 is clearly a function of both P_1 and M_s . The present structural limitation on the driver tube and its gas-supply system (plumbing, pumps, and reservoirs) precludes operation at driver pressures over 30 MPa, as shown in Fig. 9-49.

Figures 9-48 and 9-49 also show an operating point called the "tailored-interface condition." This term refers to the interaction of the reflected shock wave with the driver-gas/driven-gas interface. Generally, when a shock wave interacts with a gaseous interface, a wave is transmitted across the interface and another wave is reflected. If the acoustic impedance mismatch at the interface is zero, the transmitted wave continues with no change in strength from the impinging wave, and the reflected wave is infinitesimally weak, i.e. a Mach wave. In a shock tube, for any given driver-gas/driven-gas initial conditions, there is only one incident M_s for which the impedance jump across the interface is zero. This condition, called the tailored-interface condition, has a significant effect on the available testing time in a reflected-shock tube, as illustrated in Figs. 9-50 and 9-51.

Figure 9-50 is a standard shock-tube distance-time diagram drawn for the Grumman shock tube (6-m driver-tube length and 10-m driven-tube length) at the tailored-interface running condition. Note that the reflected shock wave passes through

the interface with no change in speed and with no reflected disturbance propagating back into Region 5. The testing time (about 10 ms in this case) is terminated by the arrival at the end wall of the reflected expansion wave from the driver. At any other speed of the incident shock wave, a new wave, propagating back into Region 5, would be generated at the point where the reflected shock wave intersects the interface, considerably shortening the available testing time. Figure 9-51 illustrates this point quantitatively. The curve gives the testing time available in a helium/argon reflected-shock tube with a 10-m-long driven tube. At one particular Mach number for the shock wave ($M_s = 3.75$ in this case), the interface is tailored and there is a significant jump in the available testing time.

The stagnant, shock heated, and pressurized gas in Region 5 bursts a thin secondary diaphragm at the end of the shock tube and then expands throughout a two-dimensional convergent/divergent nozzle at supersonic speed. The array of jets (steel rods for simulation) is located downstream of the nozzle in the test section and is terminated by an end wall, from which the flow (shock waves) reflects and propagates back upstream. It is the passage of the flow through the array of jets and the following reflected-flow phenomena that we propose to study experimentally. Figure 9-52 illustrates the proposed physical flow test model. If the initial test section is highly evacuated, the first pulse of gas will move through the nozzle at the limiting unsteady thermodynamic velocity (U_{LIM}) given by:

$$U_{LIM} = \frac{2}{\gamma - 1} a_5 \quad (34)$$

where γ is the ratio of specific heats of the gas and a_5 is the sonic velocity of the gas in Region 5. Following this initial pulse, there is an unsteady flow of gas in which the pressure increases and the velocity slows down to the values given by the steady-flow conservation equations.

The calculation of the wave that reflects from the test-section end wall (simulating the HYLIFE FSW) is used for experiment planning and for selecting instrumentation. A simplified analysis for this purpose was made by assuming that the wave speed was that obtained for a wave advancing into a

uniform region where mass flow equals throat mass flow with steady-flow properties behind the wave (i.e., no axial gradients between the wave and the end wall). The flow in the test section was modeled as a cylindrical supersonic source flow.

Point designs are shown in Table 9-15. Although other test sections were

Fig. 9-50. Distance-time diagram for tailored-interface condition.

Parameter	Values		
Shock number	3.75	5.0	3.6
Test-section expansion half-angle, θ (deg)	15	15	15
Mach number at jets	4.5	4.5	3.0
Shock-tunnel available test time ^a (ms)	10.0	2.3	3.4
Test-section test time (ms)	1.2	0.95	—
Time of initial pulse at test-section end wall (ms)	0.2	0.15	—
Time of reflected shock over jets (ms)	1.4	1.1	—
Shock tunnel stagnation temp, T_5 (K)	3230	5665	2800
Shock-tunnel stagnation pressure, P_5 (MPa)			
$P_1 = 10$ kPa	0.85	1.6	—
$P_1 = 100$ kPa	8.5	16	—
$P_1 = 200$ kPa	17	—	28
^a At region 5.			
^b $P_1 = 400$ kPa			
Test-section length	0.61 m		
Test-jet location (center of array)	0.305 m		
Driver gas	Helium		
Driven gas	Argon		
Initial driver-gas temp	290 K		
Initial driven-gas temp	290 K		
Driver-tube length	6.1 m		
Driven-tube length	10.1 m		

Table 9-15. Shock-tunnel parameters for verification experiment on HYLIFE reaction chamber.

Accomplishments Associated with the HYLIFE Chamber Concept

considered, the final selection had a half-angle of 15° . The test section was 60 cm long, and shock Mach numbers of 3.75 and 5.0 were considered. The test-section test times vary between 0.8 and 1.2 ms, and the Region 5 stagnation pressures and temperatures remained constant throughout the test-section test time. Three Region 5

stagnation pressures for each stagnation temperature correspond to three different driven-tube initial pressures (P_1).

An extensive program was conducted to examine laser doppler velocimeter (LDV) and infrared measurements of velocity and temperature downstream of the throat. To accommodate these measurements, a third shock-tube running condition was considered in which the incident M_s was 3.6, the driven-tube pressure was 400 kPa, and the resulting Region 5 pressure was 28 MPa.

The primary measurements would be of pressure at several locations on the end wall of the test section as a function of time. These measurements would be completed without jets present in the test section and then with two different diameters of jets present. Comparison of the results would quantify the effectiveness of the jets in reducing pressure on the end wall (simulated HYLIFE FSW).

Two options are available for the test program. In Option 1, the stagnation pressure and temperature in Region 5 of the driven tube are used as the initial conditions for the code verification. The stagnation pressure and temperature are derived from measurements of incident-shock speed and stagnation pressure. In Option 2, the initial

Fig. 9-51. Reflected shock-tube testing time.

Table 9-16. Option 1 . test matrix.

Mach No.	3.75	3.75	3.75	5.0	5.0	5.0
Pressure 10^5 Pa	0.1	0.5	1.0	0.1	0.5	1.0
	Number of runs					
Without array	1	1	2	1	2	1
1.27-cm jets	1	1	2	1	1	2
1.905-cm jets	1	—	1	—	1	1

Fig. 9-52. HYLIFE supersonic-flow test section.

conditions for the code verification are obtained from measurements of velocity, pressure, and temperature at a test station located between the test-section throat and the inner radius of the jet array.

The Option 1 test matrix of 20 tests is shown in Table 9-16. Tests will be conducted at Mach numbers of 3.75 and 5.0 and at initial driven-tube pressures of 10, 50, and 100 kPa. At all Mach numbers and pressures, tests will be conducted as calibrated runs without the array. Most tests will be run with a 13-mm-diam jet array, but some tests will have a 19-mm-diam jet array to determine the effects of jet diameter. Four tests will be duplicated to check the repeatability of the instrumentation: two without the array, and two with the 13-mm-diam array. Timespan for completion of Option 1 is 27 weeks, including construction of components, setup, and testing.

Under Option 2, two proof-of-principle experiments would be run to determine the feasibility of using LDV and infrared emission measurements for velocity and gas temperature downstream from the throat of the test section. Five test runs with the LDV would be required to establish a velocity profile, and one separate test run would be needed to establish a temperature profile. The test section would be designed and built in the same manner as for Option 1. The schedule for Option 2 runs for 33 weeks and again is composed of component construction, set-up, and testing. A 13-week test period includes provisions for one setup and a shakedown procedure and 50 runs.

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Experimental Investigation of the Stability with Liquid-Sheet Jets: A Comparison with Linear and Nonlinear Theory. A potential problem arises in the application of liquid jets in the HYLIFE reaction chamber. The vacuum vessel and jet nozzles could sustain large amplitude vibrations from "ringing" of the entire structure caused by the pulsed nature of laser fusion microexplosions. Researchers at the University of California,

Davis, have theoretically and experimentally analyzed the stability of jets under forced vibrations as a part of the HYLIFE design study. Jet breakup would allow neutrons, ions, and x rays to stream directly onto the structure and to thereby drastically shorten the useful life of the chamber.

Experiments were performed on planar sheet jets of water flowing vertically downward. The nozzles were vibrated transversely to excite the sinuous mode of instability. The sheet jets were photographed from the edge view to obtain wave growth parameters. The data are compared both to a linear⁷⁰ and a second order nonlinear theory.⁷¹

The stability of the jets was studied and compared to the linear temporal stability theory of Squire,⁷² as well as Hagerty and Shea.⁷³ It was shown that the linear theory could explain only the initial portion of the wave envelope, as expected.⁷⁰ A second-order nonlinear theory was then developed by Asare,⁷¹ based on perturbation methods such as described by Nayfeh.⁷⁴ Nayfeh and Saric⁷⁵ have solved the somewhat similar problem of the Kelvin-Helmholtz instability on the free surface of a semi-infinite liquid sheet in the presence of a gas flow over the surface, but in their case the gravity was perpendicular to the free surface. Asare has solved the case of the inviscid, incompressible planar sheet jet with two free surfaces in the absence of gravity forces. Specifically he solved for the sinuous mode of instability which results from the relatively large-amplitude, transverse forced vibration imposed on the nozzle.

The method of derivative expansion using multiple time and length scales was used by Asare⁷¹ to solve for the second-order terms in the sheet jet wave growth. As in the case of the linear stability theory,^{70,72,73} the liquid jet and surrounding air are assumed to be incompressible, inviscid, and initially irrotational. The continuity and transverse momentum equations are solved for each of the three regions (Fig. 9-53) subject to the usual kinematic and dynamic boundary conditions at the two interfaces. The linear theory reveals that the sinuous mode induced by the transverse forced nozzle vibration will be unstable for forcing frequencies below a critical frequency which can be determined by the solution of the

characteristic equations obtained from the linearized equation and boundary conditions.

Fig. 9-53. Edge view of a liquid sheet jet showing the sinuous mode of instability.

Fig. 9-54. Spatial growth rates predicted by the linear theory for constant jet velocity for flows from two nozzles used in the experiments.

The sinuous mode (sketched in Fig. 9-53) is shown in Ref. 73 to have a higher growth rate than the dilational mode where the two surfaces are 180° out of phase. In addition, it was shown in Ref. 70 that the growth rates predicted by the linear temporal theory are essentially the same as those predicted by the linear spatial theory using the transformation that the spatial growth rate is approximately equal to the temporal growth rate divided by the initial jet velocity. The same transformation has been shown by Asare to be valid for the nonlinear theory.⁷¹ The spatial growth rates predicted by the linear theory are shown in Fig. 9-54 for the two nozzle flows discussed later in this work.

The solution for the dimensionless wave surface amplitude was assumed to be of a form containing first and second order terms using an expansion in terms of a small dimensionless parameter, ϵ . The velocity potential is assumed to have an analogous form. After substitution into the equations of motion for each of the three regions of Fig. 9-53 and into the boundary conditions, and removing secular terms from the second and third order solutions, ϵ and the coefficients were found. The coefficients can be expressed in terms of a parameter assumed to be a weak function of x and t .

This parameter was then evaluated by solving an approximate second-order differential equation obtained as a consequence of suppressing the secular terms.⁷¹

Results for this nonlinear theory AD2 of Ref. 71 (referred to as NT hereafter) are compared to the linear theory (referred to as LT hereafter) for one typical set of run conditions for Nozzle A in Fig. 9-55. The constant-velocity LT (shown by the solid curve) predicts that the sheet jet is unstable at these conditions, and that the mean wave envelope amplitude, A_e , will grow essentially exponentially with the growth rate shown in Fig. 9-54 (Ref. 70). It can be seen from Fig. 9-55 that the NT (shown by the dash-dot curve) begins to deviate significantly (>5%) from the LT at values of A_e/A_0 of around 10 for this relatively small initial forcing amplitude, A_0 , of only 0.2 mm.

The second-order nonlinear term can be seen to be stabilizing; i.e., it causes the wave envelope amplitude to reach a maximum and then decrease toward zero very rapidly. The stabilizing effect of the nonlinear terms

implies that the stabilizing surface tension force increases faster than the destabilizing aerodynamic force as the wave form begins to depart from a near-sinusoidal shape. It should be noted that the second-order NT is no longer valid as the wave envelope amplitude approaches zero, because this implies that the magnitude of the second-order term is approaching that of the first order term.

The NT predicts the fall distance x_m where the maximum wave amplitude will be reached and also the fall distance x_{co} at which the wave amplitude is zero (the cross-over point). These predictions will be compared later to the actual behavior of the sheet jets.

Since the sheet jets were directed vertically downward in our experiments, the jets accelerated and the jet velocity, U , and thickness, b , did not remain constant as assumed above. An attempt was made to include these effects in both the LT and the NT using the following approximation method (the results are termed "variable-velocity" and are designated LTX and NTX, respectively)⁷¹ The equations were integrated for a small Δx starting at the nozzle exit. Then the conditions at the end of the integration step were used as new initial conditions for evaluating all the parameters (such as growth rate, etc.) for the next integration step. In this way a few approximate variable velocity predictions were obtained at the end of our research program. The LTX is shown in Fig. 9-55. The difference between the LTX and constant-velocity LT for this high initial jet velocity of 5.0 m/s can be seen to be small. Difficulties were experienced in calculating the NTX for most of the unstable cases as well as for the cases where the jets were initially stable but accelerated into the unstable regime. We did not resolve this problem.

Extensive results using the experimental apparatus described elsewhere^{70,76} with three different nozzles were obtained in the course of this research.^{70,71} Due to space limitations, only a small but representative sample of these experimental results will now be compared to the nonlinear theory.

Nozzle A produced the thinnest sheet jets with initial thicknesses, $b_o = 2h_o = 1.59$ mm and an initial sheet width, $w_o = 78$ mm. The linear theory predicts that all these jets will

be unstable at 70 Hz as can be seen from Fig. 9-54 where the critical frequency for Nozzle A is about 332 Hz. The sheet jets all tend to grow rapidly initially, but then appear to reach a "saturation" amplitude. The jets actually started to break up when they first reached the "saturation" amplitudes and then the assembly of fluid ligaments and drops just continued to fall, giving the appearance of a nearly constant or saturated amplitude at large fall distances.

The nondimensional data for the mean wave envelope amplitudes, A_e/A_o , are plotted in Fig. 9-56 along with the LT. The constant-velocity LT can be seen to agree with the data for short fall distances only, as expected. The approximate linear theory with variable velocity (LTX) only improves the agreement slightly for these jets with high initial velocities of 5 m/s.

It is somewhat surprising at first to note that the data show that the higher the forcing amplitude, A_o , the smaller the maximum wave amplitude ratio, and that this maximum or saturated amplitude occurs at

Fig. 9-55. Comparison of two versions of the linear theory with the constant-velocity nonlinear theory for one typical set of experimental conditions for nozzle A.

smaller fall distances for larger A_o . The constant-velocity NT curves can be seen to agree with these trends in data. The NT tends to predict the fall distance at which the jets first reach the maximum amplitude quite well. However, the experimental data all show a higher amplitude than predicted by the NT.

Beyond the maximum point the constant-velocity nonlinear theory predicts that the wave envelope should decrease rapidly toward zero. It is hypothesized that the sheet jets cannot actually follow these rapid changes imposed by the nonlinear effects and begin to break up as a result.

The fall distances for breakup read from many photographs for Nozzle A⁷¹ and also

for earlier runs using Nozzle C⁷⁰ (with $b_o = 2.5$ mm and $w_o = 76$ mm) are compared to x_m and x_{co} in Fig. 9-57 in terms of the appropriate dimensionless parameters from the nonlinear theory. The agreement can be seen to be reasonably good with jet breakup tending to occur between x_m and x_{co} .

Data at three different ambient pressures are compared to the theories in Fig. 9-58. For all three cases, the LT predicts that the sheet jets will be unstable, although at $P_a = 0.17$ atm the jet is only very slightly unstable (i.e., the growth rate predicted by the linear theory is quite small⁷⁰). The NT can be seen to predict the observed fall distance to the maximum points reasonably well, but it does not predict the observed maximum amplitudes at breakup very well. As in Fig. 9-56, the experimental data show a higher amplitude than predicted by the nonlinear theory.

Nozzle B produced thicker jets than Nozzle A. The initial thickness, b_o , of Nozzle B jets was about 7.9 mm and the initial width, w_o was 75 mm. For this large a flow area, the maximum velocity at which we could obtain reliable data was about 2 m/s. This had two potentially important consequences. First, the jets accelerated sufficiently to almost double the jet velocity in the 60 cm fall distance of our apparatus. Thus variable-velocity effects may have been much more important for Nozzle B than for Nozzle A (where U increased only about 20% in 60 cm). Secondly, for forcing frequencies above the initial critical frequency of about 23 Hz, the linear theory predicts that the sheet jets will be stable at an inlet velocity of 2 m/s. However, as the sheet jet accelerates, the critical frequency increases and the jet soon enters the unstable regime where the spatial growth rate is positive.

The LT predicts that the wave envelope will be sinusoidal for the cases where the jets are initially stable. This is due to the fact that the wave amplitude grows initially due to the *initial transverse momentum* imparted to the jet by the transverse forced vibration, not due to any instability. After a certain time the surface tension forces (which dominate the aerodynamic forces for the stable case) succeed in slowing and reversing this transverse motion of the fluid particles. Since the theories are inviscid, the wave envelope is

Fig. 9-56. Comparison of the linear and nonlinear theories with the experimental data for nozzle A.

Fig. 9-57. Comparison of the predictions of the nonlinear theory (NT) for the fall distance to the maximum wave envelope amplitude, x_m , and to the zero cross-over point, x_{co} , with the data points where the sheet jets begin to break up.

predicted to continue to grow and decay sinusoidally indefinitely according to linear theory.

The NT introduces second-order effects which tend to increase the effectiveness of the surface tension forces and flatten the wave envelope. The agreement with the data in Fig. 9-59 can be seen to be reasonably good for short fall distances. However, the constant-velocity nonlinear theory is for a sheet jet which remains in the stable regime. Since these jets enter the unstable regime at a fall distance, x_i , of about 14.6 cm, the variable velocity effects should be important for large fall distances.

The approximate effect of variable velocity on the linear theory is shown by the dashed curve (LTX) on Fig. 9-59. It can be seen that when the jet goes unstable at $x_i \approx 14.6$ cm, this theory shows the wave envelope growing rapidly, contrary to the trends exhibited by the data.

Attempts were made to estimate the variable-velocity effects on the nonlinear theory using the approximation method described above, but the computer code always experienced difficulties near the transition point, x_i . As mentioned previously, there was not sufficient time or funds to solve this problem.

It should be noted that we did not see breakup of the Nozzle B jets in the observable fall distance of 60 cm. The rather flat maximum of the wave envelope and gradual variation of the wave envelope amplitude with fall distance may result in effectively delaying the breakup, but this is only a speculation.

The nonlinear theory developed by Asare^{71,77} agrees only in part with the experimental results. In particular, for the thinnest sheet jets from Nozzle A:

- The NT correctly predicts that the second-order term is stabilizing.
- The NT also predicts the fall distance at which the wave envelope amplitude reaches its maximum value.
- The breakup of the jet tends to begin at the fall distance where the NT predicts the wave amplitude will be at maximum. (The breakup may be induced partly by the very rapid decrease in the wave envelope amplitude predicted by the theory after the maximum point.)
- The actual wave amplitudes at breakup

tend to be significantly larger than predicted by the NT.

For the thickest sheet jets from Nozzle B (which experience large velocity increases with fall distance):

- In the unstable regime, at low forcing frequencies below the critical frequency, the sheet jet behavior is best described by

Fig. 9-58. Comparison of the linear and nonlinear theories with the experimental data for nozzle A.

Fig. 9-59. Comparison of the linear and nonlinear theories with experimental data for nozzle B for two nozzle forcing amplitudes. ▼

the approximate version of the nonlinear theory which accounts for the variable jet velocity with fall distance (i.e., NTX).

- In the higher forcing frequency regime where the jets are initially stable, the agreement between the NT for constant jet velocity and the data is reasonably good for short fall distances. The jets quickly accelerated into the unstable regime, and hence the NT would not be expected to give good agreement. Also, the effective initial forcing amplitudes were very large, causing the waveform to become highly nonsinusoidal in shape. The resulting zig-zag shape would be expected to cause the aerodynamic and surface tension forces and their distribution on the sheet jet surfaces to be strongly affected, so that even a second-order NT might be inadequate. Additional theoretical studies are required to fully explain the nonlinear behavior of these rapidly accelerating liquid jets.

Scaling of these results to the condition expected in the HYLIFE reaction chamber between microexplosions indicates that no sheet jet breakup would be expected in the fall distance of the lithium jets.

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Vacuum Pumping. Vacuum pumping is an integral part of all fusion reaction chambers and, as such, its design and that of the reaction chamber should be carried out in parallel so that the requirements unique to each particular fusion confinement concept are identified early. Each concept imposes specific constraints on the vacuum system

options. Conversely, vacuum system considerations, especially pumping capacity, can impact the selection of the fusion device operating point. This article describes a preliminary analysis⁷⁸ of the vacuum pumping requirements for the HYLIFE inertial confinement fusion reactor concept. The pumping system removes constituents, principally hydrogen isotopes and helium, from the reaction chamber cavity during a portion of the time between fusion reactions, so that initial conditions are restored.

HYLIFE design criteria selected for the vacuum pumping system of this study are based on a continuous series of 1-Hz fusion reactions. At time, $t = 0$ when a fusion reaction occurs, the cavity equilibrium temperature is 500°C. A time span of 0.25 s was arbitrarily selected during which the lithium was heated, partly vaporized and recondensed. At $t = 0.25$ s, the lithium vapor temperature and that of the gases evolved as a result of the fusion reaction is assumed to be 1000°C. This assumed temperature value is not expected to be exceeded when pumping is started at $t = 0.25$ s. More likely, the temperature at this time will have recovered closer to its initial value, and the assumption is conservative. The time period from $t = 0.25$ to 1.0 s is available for cavity pumping.

The gases generated after each fusion reaction are hydrogen isotopes and helium (proportions shown in Table 9-17). Two levels of hydrogen were evaluated to allow for uncertainty in the amount which might diffuse from the secondary intermediate loop. Initial cavity pressures were chosen at values ranging from 10^{-4} to 10^{-1} Torr (excluding lithium vapor pressure), so that pressure levels dictated by any particular beam injection method would be considered. Vacuum system data were elaborated at 10^{-3} Torr to provide insight into possible system configurations. The cavity volume to be pumped, which is the difference between the cavity total volume and the volume occupied by the liquid lithium, was estimated at 700 m³.

The total effective pumping speed requirement (i.e., at the vacuum vessel wall) was derived using the relationship for "pumpdown" of a cavity at a fixed pumping rate.⁷⁹ The required partial pressure change from $t = 0.25$ to 1.0 s was the summation of partial pressures calculated for the gas

Table 9-17. Design criteria for constituent quantities and pressures.

Gas	Generation g/pulse	Partial pressure increase, Torr
Deuterium	0.0074	2.09×10^{-4}
Tritium	0.0196	3.69×10^{-4}
Helium	0.0176	4.99×10^{-4}
Hydrogen	0.0054 or 0.152	3.05×10^{-4} or 8.6×10^{-3}
Total	0.05 0.195	—
Initial pressure conditions		
Lithium	3.6×10^{-3} Torr	
Other gas	10^{-4} , 10^{-3} , 10^{-2} , 10^{-1} Torr (four cases)	

constituents from Table 9-17. The assumed vacuum system configuration (Fig. 9-60) consisted of numerous circular pumping ducts extending radially outward from the vacuum vessel wall, with each duct having one or more pumps. The required pump speed at 25°C was determined from the basic relationship for pumping at constant temperature, modified for the elevated cavity temperature in the cavity where the effective pumping speed is specified.

The basic relationship for a constant temperature system is

$$\frac{1}{S_{\text{eff}_p}} = \frac{1}{C} + \frac{1}{S_p} \quad (35)$$

where

S_{eff_p} = effective pumping speed (cfm or m^3/s), at the vessel wall for a cavity temperature equal to the gas temperature at the pump

C = duct conductance

S_p = pump speed

The cavity effective pumping speed (S_{eff}), for a cavity temperature which is different from pump temperature, is derived from the root-temperature pumping relationship which is valid through molecular and transition regimes of interest in this study:

$$\frac{S_{\text{eff}}}{S_{\text{eff}_p}} = \left(\frac{T_c}{T_p} \right)^{1/2} \equiv \text{TR} \quad (36)$$

where T_c is the temperature in the cavity and T_p is the temperature of the constituents passing through the pump (25°C).

Combining the Eqs. (35) and (36) and defining K as the ratio of duct conductance (C) to pump speed (S_p) gives

$$S_p = \frac{K+1}{K} \cdot \frac{S_{\text{eff}}}{\text{TR}} \quad (37)$$

The results of the vacuum pumping systems analysis are shown in Table 9-18 for the selected design point (initial partial pressure of the constituents equals 10^{-3} Torr) and in Fig. 9-61 for the parametric evaluation (initial partial pressure range from 10^{-1} to 10^{-4} Torr). The assumed configuration for the design point analysis consisted of 40 ducts of 900-mm-diam, 3-m long. Duct conductances for this

configuration were derived for the cases of high and low hydrogen content using average molecular weights. To satisfy the requirement to return to initial pressure at $t = 1.0$ s, the pumping requirement is one 25 000-cfm ($12\text{-m}^3/\text{s}$) pump per duct for the

Fig. 9-60. HYLIFE vacuum system schematic showing one typical location.

Fig. 9-61. Required initial pressure vs pumping capacity.

Accomplishments Associated with the HYLIFE Chamber Concept

low hydrogen content case, and two 30 000-cfm (14-m³/s) pumps per duct for the high hydrogen content case. Roots blower-type pumps are commercially available in these sizes.

The parametric study results (Fig. 9-61) clearly show the effect on the pumping system of varying the required initial pressure, and allude to the difficulty of restoring low operating pressures around the 10⁻⁴ Torr range in the allotted 0.75 s because of the high pumping requirement. For this set of data, the value of the duct conductance to pump speed ratio (K) was chosen at 0.7, which implies smaller duct dimensions at the expense of larger pumps. The design point with forty 900-mm-diam ducts points out the need for a reduced duct size; however, reducing the duct size below 600-mm diameter would more than double the number of pumps.

During the first quarter-second after each fusion reaction, the vacuum system would be isolated from the cavity by a conductance limiter such as a valve qualified for high temperature service. This would prevent appreciable lithium vapor from passing through the vacuum system pumps. The bulk of the lithium vapor intersecting the first structural wall would condense on its relatively cool surfaces. Lithium entering the pumping ducts would be intercepted by condensers (baffle heat exchangers) in the ducts. The condensers would serve the dual purpose of removing the lithium and of precooling the other gas constituents before they enter the vacuum pumps.

The major vacuum system components are: a reasonably fast shutoff valve for high temperature service with lithium and tritium

compatibility; a duct heat exchanger to condense lithium for separation and to cool gases from 1000°C for pumping; a medium range roots blower-type vacuum pump; and hydrogen isotope/helium separation system for tritium recovery and accountability. All items except the vacuum pumps will require development. Note that this system is not the primary tritium recovery system and a system must be devised for tritium recovery from the lithium in the primary loop.

The design approach for this vacuum pumping system study was to apply pumping directly to the vessel cavity, and to maintain the pressure required by the particular driver in the cavity. The study has shown the problem of this approach, illustrated by the high pumping requirements and concomitant large system.

It may be more appropriate to pump the beam and pellet injection tubes directly and to satisfy an integral line density limit to be specified for beam propagation to the pellet. If pumping only the beam and pellet injection tubes turns out to be a feasible concept, the size of the vacuum pumping system could be reduced below that determined by this study.

In conclusion, we have found that the vacuum pumping system configuration and its ability to satisfy the design criteria are clearly affected by the selected initial cavity pressure. Large vacuum systems are required (e.g. forty 25 000-cfm roots blower-type pumps or more at an initial pressure of 10⁻³ Torr) if the required initial pressure is low. Reasonable size vacuum systems are anticipated if the required initial pressure is in the order of 10⁻¹ Torr or if the beam and pellet injection tubes can be selectively pumped.

Table 9-18. Vacuum pumping system definition for selected design points.

Parameters	Low hydrogen	High hydrogen
Evolved mass per pulse, g	0.05	0.20
Average molecular wt	4.1	2.3
Pressure rise, Torr	1.4×10^{-3}	9.7×10^{-3}
Initial pressure (excluding lithium), Torr	10 ⁻³	10 ⁻³
Effective pumping speed (1000°C), CFM (m ³ /s)	1.6×10^6 (755)	4.4×10^6 (2000)
No. of pumping ducts (assumed)	40	40
Duct configuration	3-ft diam × 10 ft (0.9 m diam × 3 m)	3-ft diam × 10 ft (0.9 m diam × 3 m)
Required pump speed/duct (25°C), CFM (m ³ /s)	25 000 (12)	62 000 (29)
Assumed pumps/duct	1 × 25 000 CFM	2 × 30 000 CFM
Total no. of pumps	40 at 25 000 CFM	80 at 30 000 CFM
Duct conductance/pump speed	3.4	6.3

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Electromagnetic-Pump Considerations. The 1979 Laser Program Annual Report⁸⁰ included a comparison of mechanical and electromagnetic (EM) pumps. Here, we describe additional details of the physics of EM pumps and present a preliminary design for a full-scale HYLIFE pump; i.e., a flow

rate of $8 \text{ m}^3/\text{s}$ (128 000 gpm) and a developed head of 82 kPa (12 psi).

An isometric sketch of a helical-rotor pump is shown in Fig. 9-62. As Figs. 9-63(a) and (b) show, direct current passing through the series-wound helical windings on the surface of the rotor (Fig. 9-64) produces a

◀ Fig. 9-62. Helical-rotor electromagnetic pump.

Fig. 9-64. Rotor for a $0.13\text{-m}^3/\text{s}$ (2000-gpm) sodium pump.

Fig. 9-63. Operating concept of the electromagnetic pump.

radial magnetic field in the liquid metal. The field lines pass outward through the liquid-metal annulus from the north magnetic poles of the rotor, along the outer flux-return path, and return inward through the liquid-metal annulus to the south magnetic poles of the rotor. In Fig. 9-63(b), the radial (horizontal) axis has been enlarged for clarity. The dashed lines are the magnetic poles, i.e., the centers of the windings. Blue indicates that

magnetic fields are passing outward from the rotor through the liquid metal. The magnetic field varies with location, due to the positions of the windings; therefore, as the rotor spins, the field at any location varies in time, as shown in Fig. 9-63(c). (Blue indicates that the field is changing from inward to outward.)

Electrical currents are induced in the liquid metal as shown in Fig. 9-63(d) according to Faraday's law of induction,

$$\oint \vec{E} \cdot d\vec{\ell} = - \frac{d}{dt} \int \vec{B} \cdot d\vec{A} \quad (38)$$

where $\vec{B}$ and $\vec{E}$ are the magnetic and induced electric field vectors and $d\vec{\ell}$ defines the contour surrounding the area $d\vec{A}$. These induced currents parallel the helical windings on the rotor, except that they are offset by one-half the pole spacing.

The magnetic field interacts with the induced currents to produce a driving force on the liquid metal as shown in Fig. 9-63(e) according to the Lorentz relation,

$$\vec{F} = i\vec{\ell} \times \vec{B} \quad (39)$$

where i is the induced current and $\vec{F}$ is the driving force. Because the force has both axial and radial components, the liquid metal will move around and through the annulus without flow guides or passages.

Detailed design procedures and the effects of construction materials and techniques for helical-rotor EM pumps are discussed in Ref. 81.

Helical-rotor EM pumps are an established technology for pumping liquid metals, although they have not been scaled up to the size required by the HYLIFE concept. In 1962, Atomics International built and operated a $0.13\text{m}^3/\text{s}$ (2000-gpm), 240 kPa (35-psi) molten-sodium pump.⁸¹ Figure 9-64 shows the 0.53-m-diam rotor of that pump. It was interesting to note that the measured efficiency of the pump (26%) was higher than the calculated efficiency (21%).

Until recently, we have assumed that ferromagnetic steels, such as 2-1/4 Cr-1 Mo, could not be used in the EM pump.

Fig. 9-65. Electrical resistivity of lead-lithium solutions at 800°C.

Fig. 9-66. Efficiency of a helical-rotor superconducting electromagnetic pump as a function of the load on the cryogenic heat removal system.

However, the thinness of the annular-passage wall and the fact that, at 500°C, this steel is not highly ferromagnetic indicate that the steel may be acceptable as a pump-construction material. If so, elimination of dissimilar-metal transition joints in the primary loop will be a welcome improvement in the HYLIFE design.

The efficiency of an EM pump is inversely related to the liquid-metal resistivity. The alloy $Pb_{83}Li_{17}$ has a fourfold higher resistivity than pure lithium (Fig. 9-65). Since the mass-flow rate for $Pb_{83}Li_{17}$ is also much higher (tenfold) than for pure lithium, the pump efficiency will have high leverage. Hence, EM pumps will probably not be appropriate if $Pb_{83}Li_{17}$ is used as the HYLIFE working fluid.

Some very preliminary designs of conventional and superconducting pumps were produced for the HYLIFE operating conditions (Table 9-19), and a study of the tradeoffs between pump size, cost, and efficiency is in progress. Several areas of the design procedure were arbitrary; e.g., 11-mm-thick pipe walls were used irrespective of pump size, and, for the superconducting design, the efficiency of the cryogenic refrigeration system was not considered. If only 0.1% of the pump losses must be removed by the cryogenic system, an overall pump efficiency of 60% falls to ~40% (Fig. 9-66).

Both mechanical and EM lithium pumps are being considered for the HYLIFE system. Each type of pump has distinct advantages and weaknesses. Neither system has been demonstrated at even one-half the scale of HYLIFE, but large mechanical pumps are being developed for the sodium-cooled liquid-metal fast-breeder reactor (LMFBR). To permit future design flexibility, we are continuing to investigate the feasibility of EM pumps, since they are not being actively pursued by either the LMFBR or Magnetic Fusion Energy (MFE) programs. Our first-cut design point for a normal-conducting EM pump appears attractive from the viewpoints of simplicity, maintainability, and capability for handling the pulsating, low-suction-head HYLIFE flows; however, the efficiency and size of this design are both marginal. The initial design point of a superconducting pump is as efficient as a mechanical pump, and is very

Parameter	Normal windings	Superconducting windings
Efficiency (%)	41	53
Weight (Mg)	94	34
Lithium inventory (m ³)	34.4	9.2
Annulus width (m)	0.30	0.40
Rotor length (m)	6.10	1.65
Rotor diameter (m)	1.56	1.03
Pump diameter (m)	2.59	2.49
Rotor speed (rpm)	49	77
Peak flux density (T)	0.34	0.59
Drive motor power (MW)	1.5	1.3

compact, but the cost and efficiency effects of the cryogenics remain to be calculated.

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Lithium Safety Considerations. The presence of a large amount of molten lithium in the HYLIFE power plant concept concerns scientists interested in fusion reactor safety. Reports and media releases on fusion safety issues often make the questionable assumption that huge amounts of molten lithium immediately and completely react with air or water. The heat of reaction is further assumed to be well coupled to activated structural materials, vaporizing them and making them available for dispersal to the environment if the containment building fails.

In the HYLIFE concept, a number of passive engineering design features should prevent these events. Low activation materials have been used in the reaction chamber. All lithium circuits are in a containment room filled with inert gas that will not react with molten lithium. The lithium loop is separated from the steam loop by an intermediate sodium loop. Concrete surfaces are lined with steel to prevent lithium-concrete reactions. Surfaces below liquid metal components are sloped to deep gallery cells such that large spills of lithium drain away from the activated structural components. The galleries are

Table 9-19. Preliminary HYLIFE pump-design points.

tall and narrow to minimize the exposed surface area of any spilled lithium. A layer of hollow graphite microspheres will be prepositioned in each gallery so that they will float on the top of spilled lithium to prevent or smother potential fires.

Hot molten lithium can react violently with either air or water. Cool molten lithium may not react significantly with dry air, although it can react violently with either wet air or water. Cold solid lithium does not react violently with either air or water, (e.g., it fizzes like a Seltzer tablet when dropped in water).

After a lithium fire, a finely divided solid lithium (dust) that is coated with a reaction product may remain. Although the reaction product prevents further reaction with wet

or dry air, water can dissolve the coating, and a violent reaction with the lithium dust could result. However, mixtures of reaction product and lithium dust have a characteristic coral color which can alert a decontamination crew to avoid the use of water.

Radioactive materials in a HYLIFE power plant fall into three categories: activated structure; activated impurities in the molten lithium; and tritium in the molten lithium, the vacuum system, and the target fabrication factory. The structural materials have been chosen to minimize activation, particularly long-lived activation products. Activated impurities in the lithium will be present due to target debris, corrosion products, and impurities in the original lithium. Care is being taken to minimize these impurities and their activation hazards.

Tritium will be held to below one weight part per million (wppm) in the molten lithium by a molten salt-electrolysis process. Because of the high solubility of tritium in lithium, only about 40 ng of tritium is present in the reaction chamber as a vapor. This amounts to about 400 μCi which can be diluted to below the maximum permissible concentration for 168 hours of exposure per week by either 200 m^3 of air (one-third the chamber volume) or 12 litres of water. In the event of a spill of more than a few litres of lithium, the 200-fold larger volume of the room will be contaminated by tritium that outgasses until its vapor pressure is reached. The tritium in the building volume will then amount to about 8 μg , but the concentration will still be only one-third the maximum permissible concentration in air. The solubility of hydrogen isotopes in lithium increases with decreasing temperature; hence less tritium will be released by a small, rapidly cooling spill.

As a further precaution, an inert cover gas, such as argon, will fill the containment room. Even a 10- m^2 hole in the room cannot result in a large air concentration without a considerable delay time (Table 9-20).

Makeup cover gas can be used as an active safety measure to delay the buildup of air near the lithium circuit. A temporary plastic bubble is an additional active safety measure that can be used.

In addition to preventing direct lithium-water contact in the event of a spill, the intermediate sodium loop minimizes the

Table 9-20. Time for air concentration to reach 10% if the 110 000- m^3 containment room is breached.

Model 1: Air diffuses through a hole in the 1-m-thick wall.

$$t = \frac{LV}{DA} \ln \left(\frac{1}{1 - C} \right)$$

where t is the time, L and A are the hole thickness (1 m) and area, V is the room volume, D is the diffusion coefficient (1920 m²/s), and C is the final air concentration (0.1).

Hole size (m ²)	Time (days)
1	1000
10	100

Model 2: This model requires breach of both the exterior building wall and the interior containment room wall. Wind blows into the room through a 1 to 10 m² hole until the interior static pressure equals the exterior stagnation pressure. When the wind slows or changes direction, a mixture of air and inert gas flows out the hole until the interior and exterior static pressures are equal. The cycle repeats until the interior air concentration is 10%.

$$\Delta n = \frac{(\rho v^2)}{2} \frac{V}{RT}$$

where Δn is the additional mass of air added to the room, ρ and v are the air density and velocity, V is room volume, R is the gas constant, and T is the temperature.

Perpendicular wind velocity		Wind cycle time	Time (hours)
m/s	(mph)		
0 to 4	(0 to 10)	2 min	24
4 to 8	(10 to 20)	2 min	6
0 to 9	(0 to 25)	2 min	4
0 to 9	(0 to 25)	30 s	

Model 3: Both the exterior building wall and the interior containment room wall are breached in two separate locations. If perfect mixing takes place, a 4-m/s wind through a 5-m² hole could increase the air concentration to 10% in about 10 min. However, the mixing time will probably be longer. In the regions of the gallery drainage tanks, a stagnant gas region should delay mixing for many days or weeks.

release of radioactive elements in the event of a leak inside the steam generator. The intermediate loop also takes advantage of the extensive sodium pump and steam generator development work in the Liquid Metal Fast Breeder Reactor (LMFBR) program.

The steel lined concrete cell concept confines spilled lithium to a small area. The steel lining prevents extensive reaction of lithium with concrete. Lithium-concrete reactions could be further reduced by the inclusion of a sacrificial material between the concrete and the steel. In the event of a spill and cracks in the steel lining, the sacrificial material would alloy with the lithium and protect the concrete. The alloying reaction should be endothermic and rapid although a rapid, mildly exothermic reaction may be acceptable.

A large scale reaction of lithium and concrete is precluded in the region of the activated steel because the steel lined concrete cells quickly drain into deep narrow galleries. Minor spills and the residue left after major spill drainage will solidify rapidly in the form of thin sheets, which can be easily rolled up and removed.

The floor areas below the lithium loop components slope into a number of deep narrow galleries. These are double walled tanks that can accept large lithium spills while maintaining a minimal lithium exposed surface area. Because of the large heat capacity of lithium, the gallery contents could remain molten for several weeks. An active measure consisting of heat pipes or heat exchanger tubes containing a nonreactive coolant could be used to cool and solidify the spilled lithium in a much shorter time.

The small exposed surface area of spilled molten lithium in the galleries minimizes the probability and size of fires. The presence of inert gas and floating hollow graphite microspheres at the exposed surface further reduces the probability of fires. The localization of the exposed lithium surface also facilitates the injection of additional inert gas and microspheres to smother a fire. Finally, because of the short flame of burning lithium, and because of the distance between the galleries and any activated structure, a fire in a gallery will not release large quantities of radionuclides into the containment building atmosphere.

In summation, passive measures designed into the HYLIFE power plant do not allow large scale reaction of molten lithium with either water or air. Most spilled lithium quickly drains to containers designed to minimize both the possibility and the consequences of a fire. A fire that is ignited by any means will probably be small scale. The prepositioned graphite microspheres and inert gas should smother fires with no operator assistance. Simple active measures can further reduce the fire probability and extinguish any fire that starts, but these measures are the secondary rather than the primary line of defense against harm to the environment by the HYLIFE plant.

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Build-up of Tritium in a Liquid Lithium

Breeding Blanket. The design of the fusion chamber must incorporate lithium for the purpose of breeding new tritium to replace the tritium burned in the targets. Initially, the fusion chamber and all other reactor components will be free of tritium. Once operation begins, tritium will be generated within the chamber from two sources:

- Unburned fuel from the target.
- Bred tritium.

Tritium must be recovered from the reactor to provide a continuous supply to the target factory. Depending on the chamber concept, the recovery of unburned tritium and bred tritium could require either two separate systems or only a single system. With dry-wall chambers, unburned tritium would most likely be recovered from the exhaust gas while bred tritium would be recovered from the lithium blanket. With HYLIFE⁸² or a wetted-wall chamber,⁸³ essentially all tritium would be recovered from liquid lithium.

In this article, expressions relating the build-up of tritium in the breeding material to the rates of tritium generation and recovery are developed; tritium breakeven time and breakeven inventory are defined and related to the various tritium flow rates; and an expression for the steady-state inventory is given for a liquid-lithium blanket with a molten salt extraction system. Only the tritium inventory associated with

the breeding material is dealt with here. It is assumed that mechanical pumps can be used to exhaust the fusion chamber; if this is the case, the tritium inventory associated with exhausting unburned tritium will be negligible compared to the inventory of the breeding material.⁸⁴ In addition, the tritium inventory associated with the molten salt used in the extraction process is small compared to the inventory in liquid lithium. Inventories associated with target production, target storage, and tritium storage are discussed by Sherohman.⁸⁵

Since liquid lithium is commonly chosen as the breeding material in ICF reactor designs, we consider the buildup of tritium in a liquid lithium blanket assuming the recovery rate is proportional to the tritium concentration. The rate of change of the tritium inventory is equal to the difference between the generation rate and recovery rate, and can be expressed as

$$\frac{dI}{dt} = \dot{m}_g - \dot{m}_r = \dot{m}_g - \gamma I/I_L, \quad (40)$$

where

$$\begin{aligned} I &= I(t) = \text{tritium inventory in lithium (g),} \\ t &= \text{time (s),} \\ \dot{m}_g &= \text{constant tritium generation rate which may include contributions from unburned fuel as well as bred tritium (g/s),} \\ \dot{m}_r = \dot{m}_r(t) &= \text{tritium recovery rate from lithium (g/s),} \\ \gamma &= \text{proportionality constant relating the recovery rate to the tritium concentration (g/s), and} \\ I_L &= \text{total lithium inventory (g).} \end{aligned}$$

Under steady-state conditions, the recovery rate equals the generation rate and is given by

$$\gamma I_{ss}/I_L = \dot{m}_g, \quad (41)$$

where I_{ss} = steady-state tritium inventory (g). Solving for γ and substituting into (40) gives

$$\frac{dI}{dt} = \dot{m}_g - \dot{m}_g I/I_{ss} \quad (42)$$

Solving Eq. (42) gives the tritium inventory as a function of time,

$$I = I_{ss} [1 - \exp(-\dot{m}_g t/I_{ss})], \quad (43)$$

from which the recovery rate is

$$\dot{m}_r = \dot{m}_g [1 - \exp(-\dot{m}_g t/I_{ss})] \quad (44)$$

The tritium inventory and the tritium recovery rate follow the same type of $(1 - e)$ growth curve as they approach their respective steady state values. A similar analysis was carried out by Hickman for magnetic confinement fusion reactors.⁸⁶ The analysis for ICF reactors is different in that for ICF unburned tritium may contribute to the generation rate in the breeding material.

The tritium generation rate in lithium can be expressed as

$$\dot{m}_g = \beta \dot{m}_b + \dot{m}_i - \dot{m}_b - \dot{m}_e, \quad (45)$$

where

$$\begin{aligned} \beta &= \text{tritium-breeding ratio,} \\ \dot{m}_b &= \text{tritium burn rate (g/s),} \\ \dot{m}_i &= \text{tritium injection rate (g/s), and} \\ \dot{m}_e &= \text{tritium exhaust rate (g/s).} \end{aligned}$$

The tritium burn rate, in g/s, is related to the fusion power of the chamber by

$$\dot{m}_b = 1.78 \times 10^{-6} P_f, \quad (46)$$

where P_f = fusion power (MW_f). The injection rate is equal to the burn rate divided by the burn fraction, f_b

$$\dot{m}_i = \dot{m}_b / f_b \quad (47)$$

Tritium exhaust is unburned tritium that is pumped directly from the chamber and not absorbed in lithium. As previously stated, the inventory of the exhaust system is assumed to be negligible; therefore, exhausted tritium is available to the target factory at a constant rate, $\dot{m}_e$, from the start of operation.

Equation (45) can be rewritten as

$$\dot{m}_g = \dot{m}_i [\beta f_b + \alpha(1 - f_b)] \quad (48)$$

where $\alpha = (\dot{m}_i - \dot{m}_b - \dot{m}_e)/(\dot{m}_i - \dot{m}_b)$ = ratio of unburned tritium absorbed in lithium to total unburned tritium.

In a manner similar to Hickman,⁸⁶ we define the time at which the tritium flow returning to the target factory equals the flow from the factory into the fusion chamber as the breakeven time, t_{be} . That is,

$$\dot{m}_r + \dot{m}_e = \dot{m}_i \text{ at } t = t_{be} \quad (49)$$

Substituting $\dot{m}_r$ from Eq. (44) and solving for t gives

$$t_{be} = -(I_{ss}/\dot{m}_g) \ln [1 - (\dot{m}_i - \dot{m}_e)/\dot{m}_g] \quad (50)$$

Solving Eq. (45) for $(\dot{m}_i - \dot{m}_e)$ and substituting gives

$$t_{be} = -(I_{ss}/\dot{m}_g) \ln [\dot{m}_b(\beta - 1)/\dot{m}_g] \quad (51)$$

In terms of the parameters described in Eq. (48), this becomes

$$t_{be} = \frac{(I_{ss}/\dot{m}_i)}{[\beta f_b + \alpha(1 - f_b)]} \ln \left[\frac{\beta f_b + \alpha(1 - f_b)}{f_b(\beta - 1)} \right] \quad (52)$$

The dimensionless breakeven time, t_{be}/τ , where $\tau = I_{ss}/\dot{m}_i$, is shown in Fig. 9-67 as a function of the tritium-breeding ratio and the fraction, α , of unburned tritium absorbed in the lithium. The target burn fraction is taken as 30%. The characteristic time, τ , is the operating period required for the injection of an amount of tritium equal to the steady-state inventory, I_{ss} . The breakeven time decreases with increasing breeding ratio and increasing fraction of unburned tritium absorbed in lithium.

As a specific example, consider a reaction chamber with a steady fusion power of 3000 MW_f. With a 30% target burn fraction, the required tritium injection rate calculated from Eqs. (46) and (47) is 1.54 kg/d. The characteristic time is thus 0.65 days per kg of tritium inventory. If the tritium-breeding ratio is 1.4 and all the unburned tritium is absorbed ($\alpha = 1$) we see from Fig. 9-67 that $t_{be}/\tau = 2$. Therefore, $t_{be} = 1.30$ d for a 1-kg steady-state tritium inventory. For $\alpha = 0$ at $\beta = 1.4$, $t_{be}/\tau = 3$ and $t_{be} = 1.95$ d.

Of more interest is the breakeven tritium inventory which is defined as the net tritium absorbed by the system before the total removal rate equals the injection rate.⁸⁶ The breakeven tritium inventory is equal to the integral of $(\dot{m}_i - \dot{m}_e - \dot{m}_r)$ from $t = 0$ to $t = t_{be}$

$$I_{be} = \int_0^{t_{be}} (\dot{m}_i - \dot{m}_e - \dot{m}_r) dt \quad (53)$$

Integrating with substitutions from Eqs. (44), (45), (48), and (51) gives

$$I_{be}/I_{ss} = \frac{f_b(\beta - 1)}{[\beta f_b + \alpha(1 - f_b)]} \ln \left[\frac{f_b(\beta - 1)}{\beta f_b + \alpha(1 - f_b)} \right] + \left[\frac{f_b + \alpha(1 - f_b)}{\beta f_b + \alpha(1 - f_b)} \right] \quad (54)$$

The ratio of the breakeven inventory to the steady-state inventory is shown in Fig. 9-68. The breakeven inventory decreases both with increasing breeding ratio and with

Fig. 9-67. Breakeven time as a function of tritium-breeding ratio.

Fig. 9-68. Ratio of breakeven tritium inventory to steady-state tritium inventory.

decreasing fraction of unburned tritium absorbed in lithium.

Figures 9-69 and 9-70 illustrate the relationships between the various tritium flow rates, the breakeven time, and the breakeven inventory. Figure 9-69 is for the case where all unburned tritium is absorbed in lithium

Fig. 9-69. Relationship between various tritium flow rates, breakeven time and inventory for $\alpha = 1$.

Fig. 9-70. Relationship between various tritium flow rates, breakeven time and inventory for $\alpha = 0$.

Fig. 9-71. Molten salt extraction scheme. ▼

($\alpha = 1.0$ and $\dot{m}_e = 0$), and Fig. 9-70 is for the case where all unburned tritium is exhausted [$\alpha = 0$ and $\dot{m}_e = (1 - f_b) \dot{m}_i$]. In both cases, the burn fraction and breeding ratio are 0.3 and 1.4, respectively. The shaded area in both figures is the breakeven tritium inventory. While the breakeven time is longer for $\alpha = 0$, the breakeven inventory is less as illustrated in Figs. 9-67 and 9-68.

The steady-state tritium inventory in the lithium will depend on both the lithium inventory and capability of the recovery system. We consider here the molten-salt extraction process as an example. Studies at Argonne National Laboratory^{87,88} indicate that a molten-salt extraction process based on LiF-LiCl-LiBr (22-31-47 mol%) is an attractive technique for recovering tritium from liquid lithium fusion reactor blankets. It is estimated that with such a processing scheme tritium levels in the blanket can be maintained at concentrations less than one weight-part-per-million (wppm).

The molten-salt extraction process is illustrated in Fig. 9-71. A stream of lithium from the fusion chamber is contacted with the salt which has a higher affinity for tritium than the lithium does; consequently, tritium is extracted from the lithium by the salt during the mixing process. The two liquids are then separated, the lithium returns to the fusion chamber, and the salt is circulated to an electrolytic recovery tank. The T_2 evolved in the electrolysis process can be returned to the target factory.

In a system where equal volumes of salt and lithium are contacted, the tritium recovery rate is proportional to the mass flow rate of tritium entering the extraction system. The relationship⁸⁷ is

$$\dot{m}_r = [\epsilon D_v \eta / (\epsilon D_v \eta + 1)] \dot{m}_{ri} \quad (55)$$

where

- ϵ = efficiency of tritium recovery from the salt,
- D_v = volumetric distribution coefficient of tritium between the lithium and the salt,
- η = efficiency factor that accounts for nonequilibrium tritium distribution during contacting,
- $\dot{m}_{ri}$ = tritium mass flow rate into the recovery system (g/s).

Equation (55) can be rewritten in terms of the lithium flow rate through the extraction system as

$$\dot{m}_r = [\epsilon D_v \eta / (\epsilon D_v \eta + 1)] \dot{m}_L I / I_L \quad (56)$$

where

$\dot{m}_L$ = lithium mass flow rate (g/s),
 I = tritium inventory in lithium (g),
 and
 I_L = total lithium inventory (g).

For fixed process parameters and lithium flow rate through the system, the recovery rate is proportional to the tritium concentration I/I_L ; therefore, Eqs. (40) through (54) are indeed applicable. Comparing Eq. (56) to Eq. (40), we note that, for the molten salt extraction process,

$$\gamma = [\epsilon D_v \eta / (\epsilon D_v \eta + 1)] \dot{m}_L \quad (57)$$

At steady state, $\dot{m}_r = \dot{m}_g$, and the steady-state tritium inventory is given by

$$I_{ss} = \dot{m}_g I_L / \gamma \quad (58)$$

If all other parameters are held constant, the tritium inventory is inversely proportional to the lithium flow rate through the extraction system.

As Calaway points out,

An important feature of the presently conceived molten-salt extraction process is that it appears to be applicable throughout the entire range of tritium concentrations that would be desired. No limiting behavior in the electrochemical processing technique was found down to the detectable limit of the experimental apparatus (corresponding to 0.06 wppm of tritium in lithium). The processing requirements are quite small at 5 to 0.5 wppm of tritium in lithium. They increase linearly as the tritium concentration decreases so that the processing requirements at 0.05 wppm are a factor of 10 higher than those for 0.5 wppm. Furthermore, the molten-salt extraction process lends itself to a modular design in which relatively small units are connected in parallel to the liquid-lithium blanket. This design

allows additional flexibility for accommodating changes in the allowable tritium inventory of the blanket. If future safety regulations, economic constraints, or design difficulties require a reduction in the tritium inventory, additional processing units could be easily attached to meet the new requirements.⁸⁸ As an example, consider the lithium flow through the tritium extraction system required to maintain the tritium concentration at 1 wppm in a typical reactor. Reasonable values for ϵ and D_v are 0.9 and 2 respectively.⁸⁸ The efficiency factor that accounts for nonequilibrium tritium distribution during contacting is not yet known but researchers at Argonne National Laboratory believe 10 to 50% could easily be achieved. Consider the following parameters:

Fusion power	$P_F = 3000$ MW
Tritium-breeding ratio	$\beta = 1.40$
Total lithium inventory	$I_L = 10^6$ kg
Unburned fraction T absorbed	$\alpha = 1.0$
Desired tritium inventory	$I_{ss} = 1.0$ kg
Extraction efficiency	$\epsilon = 0.9$
Volumetric distribution coefficient	$D_v = 2.0$
Distribution efficiency	$\eta = 0.3$

For a target burn fraction $f_b = 0.3$, the tritium-generation rate in lithium is $\dot{m}_g = 19.9$ mg/s. Substituting these values into Eqs. (18) and (19) and solving for $\dot{m}_L$ gives 56.7 kg/s or ~ 0.12 m³/s. This is about 1% of the lithium flow rate required for heat removal if lithium is used as the primary coolant.

The buildup of tritium in a liquid-lithium breeding blanket for an ICF chamber has been examined. The breakeven time is found to decrease both with increasing tritium-breeding ratio and increasing values of α , the fraction of unburned tritium absorbed in lithium. The breakeven inventory also decreases with increasing breeding ratio but increases with increasing values α . For a molten-salt extraction process, the steady-state tritium inventory in lithium is inversely proportional to lithium flow rate through the extraction system.

Author: W. R. Meier

Component Design

Stress Analysis of the Alternate Helical Bar First Structural Wall (FSW). The HYLIFE reaction chamber contains a First Structural Wall (FSW) that absorbs pulsating loads generated by the fusion reactions so that only a smaller steady secondary load is imposed on the vacuum vessel. The radial momentum imparted to the liquid metal jet array after each fusion reaction is absorbed in the FSW. Ideally, the FSW will last the entire 30-year lifetime of the reaction chamber, but if not, it can be replaced for a fraction of the cost to replace the vacuum vessel.

A simple FSW design involves welding 50-mm-thick ferritic steel plates to form a cylindrical shell. Dispersed liquid from the jet array is stopped abruptly as it hits this slab-type FSW. The rapid transfer of energy

Fig. 9-72. Tapered HYLIFE reaction chamber with alternate helical bar first structural wall.

results in high stresses, which, depending on final determination of input parameters, could restrict the life time of a slab FSW below desirable limits.

Because we desire a FSW with a 30-year lifetime an alternative design was conceived that could possibly improve lifetime characteristics.⁸⁹ This alternative FSW design consists of multiple layers of bars wound in a helical fashion. Spaces between the bars result in the radial momentum of the jet array being transferred more gradually than with a slab-type FSW; hence, stresses are conceptually lower. The multiple layers of the helical bar FSW dissipate fluid kinetic energy by promoting splashing interactions of two-phase flow, as the liquid penetrates the array of bars. The number of layers of helical bars needed to dissipate most of the energy is determined so that the fluid exits with an insignificant radial velocity.

This article outlines the structural analyses which were performed on the alternative helical bar-type FSW.⁹⁰ The analyses confirmed the conceptual feasibility of this helical bar-type FSW. Also, the results highlighted some questions regarding adequacy of the natural damping mechanisms for the helical energy absorbers.

The helical bar FSW is designed as a continuous, 19-mm-diam, 2-1/4 Cr-1 Mo ferritic steel bar, as shown schematically in Figure 9-72. The inner helical layer is formed by winding the bar material from top to bottom, and the adjacent helical layer is wound from bottom to top to produce a crisscrossing pattern. The winding continues until the desired number of layers is obtained. Center-to-center spacing between the bars in a single layer is 38 mm, which can be adjusted as desired to accommodate fluid dynamic considerations. The helical layers are supported rigidly at the top. A radial sway restrictor is positioned at the bottom to prevent any contact of the bars with either the vacuum vessel or the graphite reflector. The helical layers are held separate from each other by spacers located at discrete positions around the circumference of the FSW.

The FSW deflects vertically under its own weight, and is arranged inside the vacuum vessel so that the lower end reaches the desired position when statically deflected. It deflects elastically a slight amount as liquid lithium impacts the bar layers after each

fusion reaction, which helps to absorb the energy of each impact.

Only the innermost helical layer of bars was analyzed because this layer has the highest loading, and because definition of the forcing function grows in complexity for the other helical layers. Figure 9-73 indicates the structural model used to idealize a portion of a turn on the innermost helical layer. Because the pitch is small compared to the radius and the forcing function is axisymmetric, only radial motions were considered. In reality, tangential motions are also possible because the forcing function varies with height, consequently tangential motions can appear at the top and bottom of the helix. However, these were neglected by the study. Thus, the equations used to model the helix are identical to those for a thin elastic hoop. The dynamic motion of an elastic element is given by the following equation (see Fig. 9-73):

$$(RP\dot{d} - T - CR\dot{U}) = m\ddot{U}R \quad (59)$$

where

- R = helix radius
- C = damping constant for cable of diameter d
- P = applied pressure
- d = diameter of helical cable
- $\dot{U}$ = radial displacement
- m = cable mass per unit length
- T = helix tension

The associated cable stress-strain deflection law yields a relationship between cable tension and displacement as follows:

$$T = A\sigma = EA\epsilon = EA \frac{U}{R} \quad (60)$$

where

- A = cable cross sectional
- E = modulus of elasticity
- ϵ = strain, $\dot{U}/R$ area
- σ = stress, $E \cdot \epsilon$

Therefore, the resulting dynamic equation becomes:

$$m\ddot{U} + C\dot{U} + \frac{EA}{R^2}U = Pd \quad (61)$$

The forcing function used to model the lithium impact upon the helical coil was that

of a plane liquid wall of finite width, traveling at a constant velocity (V_f), and striking a solid circular cylinder of diameter d, with axis normal to the fluid velocity. The width chosen for the liquid wall is one-half the sum of the jet widths. This assumption was based partially on the work of Glenn and Young⁹¹ which indicates that the neutron deposition associated with the D-T reaction resulted in a spallation of half the lithium thickness for a solid fall. For multiple jets, which are included in the present design, each jet will split and many of the internal jet splittings will result in net momentum cancellations. It was assumed a liquid width equal to half the sum of the individual jet widths (0.5 m total) would impact the helical coil. However, the peak loading occurs shortly after the initial impact, and would not change appreciably if the total sum of the individual jet widths had been used.

Kossom⁹² presents empirical data for the pressure associated with a finite width wall of fluid impacting a fixed circular cylinder with axis normal to the velocity of the fluid. The initial impact pressure upon the cylinder (P_{max}) has been empirically determined to be given by the relationship

$$P_{max} = \frac{7}{2} \rho V^2 \quad (62)$$

where ρ is the fluid mass density, and V is the relative velocity between fluid and cylinder.

- P = applied pressure
- $\dot{U}$ = radial displacement
- T = tension in helix
- d = diameter of helix
- R = radius (helical pitch assumed very small compared with R)
- t = time
- θ = as shown

Fig. 9-73. Structural dynamics for helical bar FSW.

Accomplishments Associated with the HYLIFE Chamber Concept

The associated impact (I_p) imparted to the cylinder is given by

$$I_p = K_H \frac{\rho \pi V d}{4} \quad (63)$$

where K_H has been empirically determined to lie between 0.2 and 0.27. The time (τ) for the fluid to travel around the cylinder defines the impact time. This together with the drag pressure (P_D) after impact are given by

$$\tau = \frac{2I_p}{P_{\max}} = \frac{K_H}{7} \frac{\pi d}{V} \quad (64a)$$

and

$$P_D = C_D \frac{1}{2} \rho V^2 (1 + K_i) \quad (64b)$$

where K_i is a spacing factor for flow past more than one cylinder (lined up normal to the direction of flow), and is given by

$$K_i = \frac{0.2d}{S - d}, 0.05 < \frac{S - d}{d} < 2, \quad (65)$$

Fig. 9-74. Typical response of a helical bar FSW.

Table 9-21. Cable dynamic stresses for modified forcing function (p') and structural damping (ζ). ▼

V_F , m/s	ζ^*	$\sigma_{\max}$, 10^7 Pa (PSI)	Time for T_{peak} to drop to 1% of $\sigma_{\max}$, s
10	0	1.213 (1 759)	∞
	0.001	1.211 (1 757)	2.956
	0.003	1.207 (1 751)	0.985
	0.010	1.193 (1 731)	0.296
20	0.001	3.909 (5 669)	2.956
	0.003	3.896 (5 651)	0.985
	0.010	3.853 (5 589)	0.296
40	0.001	14.08 (20 425)	2.956
	0.003	14.04 (20 364)	0.985
	0.010	13.90 (20 153)	0.296

$\zeta^* = C/C_c; P' \equiv P \left[\frac{V_F - \dot{U}}{V_F} \right]^2; R = 3.3 \text{ m}; d = 2 \text{ cm}; \text{ and}$
 $E = 2 \times 10^{11} \text{ Pa } (30 \times 10^6 \text{ PSI})$

where S is the center-to-center spacing of the cylinders.

It should be noted that the forcing function becomes quasi-static when the relative velocity ($V = V_F - \dot{U}$) is used (where V_F is the velocity of the wall of fluid), and that this assumption for V was used in conjunction with the data presented in Ref. 92.

It is anticipated that fusion reactions will occur with a frequency of one per second. Thus it is important to ascertain that there are sufficient mechanisms to dampen the dynamic response prior to initiation of each successive excitation. To investigate the helical damping phenomenon in a preliminary manner, two modes of energy dissipation were conceived and a numerical study was performed. The first dissipation mechanism considered was associated with nonlinear fluid viscous damping given by:

$$C \dot{U} \leq C_D \frac{1}{2} d \rho | \dot{U} - V_F | (\dot{U} - V_F) = F_D \quad (66)$$

and the second was structural damping with

$$C = 2\zeta \sqrt{\frac{mEA}{R^2}}, \quad (67)$$

where ζ is the (assumed) critical damping ratio, F_D is the drag force, and C_D is the drag coefficient for which various values between 0 and 2 were assumed.

A computer program for numerical solution of the governing equations [Eq. (60) and (61)] for purely radial FSW motions was written in FORTRAN and solved on the IBM 370-168 machine. A forward integration scheme, coupled with zero radial displacement and velocity initial-conditions, was used in each case. Typical results are plotted in Fig. 9-74 where the upper portion shows the pressure applied to the surface of the helical bar as a function of time. The bottom portion shows the resulting stress in the helical bar. The value T_{fp} on this figure is defined as the time for one-half meter of fluid to pass over the first row of helix turns. As may be seen, most of the impact pressure occurs over a time interval which is very short compared to the natural period of the helix motion, and that this natural period is short compared to T_{fp} (one quarter of a period occurs in a time equal to about 0.03 T_{fp}). Thus, damping is not an important

factor in limiting peak FSW helix stresses. Various values for C_D between 0 and 2 were tried, but the results for $\sigma_{\max}$ remained essentially fixed. Furthermore, it was determined that fluid damping alone was hardly capable of attenuating the helix motion between successive pellet firings.

The effect of structural damping upon helical motion decay is displayed in Table 9-21. As may be observed, a minimum of 0.3% is required to permit pellet firings of one per second without significant FSW motion buildup over successive fusion pulses. Reference 93 indicates that locating steel materials with such natural damping ability is difficult, but possible.

It should be noted that increasing the helix radius (R) will decrease the peak stresses. The reasoning behind this statement is as follows: the dynamic stresses, which are proportional to dynamic radial motions (U_{dynamic}), are produced by impact loading conditions and, therefore, are directly proportional to $\omega^2 U_{\text{static}}/R$; where ω is the natural frequency of the motions U_{dynamic} , and U_{static} are the static deflections associated with the peak FSW pressures if they were statically applied. For the governing FSW helix structural equations, $\omega^2 \sim 1/R^2$, and $U_{\text{static}} \propto R^2 P_{\max}$. Thus, $\sigma_{\text{dynamic}} = (EU_{\text{dynamic}}/R) P_{\max}/R$. Glenn⁹⁴ has shown that $P_{\max}$ decreases rapidly with increasing R . Hence, a minimum suitable radius can be chosen so that stresses in the helix may be maintained within acceptable limits.

In conclusion, we find that our structural dynamics studies on the helical-coil-type First Structural Wall (FSW) for the HYLIFE reaction chamber show that the maximum impact stress varies roughly with the inverse of the wall radius squared, and is virtually independent of the amount of damping present. Because of the rapid decrease of peak stress with increased radius, the wall radius can be selected to maintain the stress within acceptable limits for a given impact velocity. For example, the allowable fatigue limit of 60 MPa for 2-1/4 Cr-1 Mo ferritic steel is reached with an impact velocity of 25 m/s and the wall placed at a radius of 3.3 m.

Selection of a fusion pulse frequency requires consideration of FSW damping. The study shows the effect of structural damping on the time to recover to a low

stress value. Structural damping of about 0.3% will provide recovery within the 1-Hz pulse rate. Steel materials with this damping characteristic can be found, although not readily. If higher pulse frequencies are mandated, the FSW design probably will be more elaborate than that evaluated here. Additional investigation is desirable to assure that adequate damping can be obtained under HYLIFE operating conditions.

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Nozzle Plate. The basic concept of the High Yield Lithium Injection Fusion Energy (HYLIFE) reaction chamber⁹⁵ involves a falling liquid metal (lithium) jet array that absorbs 90% of the energy released from inertial confinement fusion reactions. The jet array disassembles once per second after the energy from each fusion reaction is absorbed. The key element of the chamber that produces the jet array is the nozzle plate. In this article we describe the concepts and the analyses of a nozzle plate that permits the jet array to be reestablished between fusion reactions and withstand the associated static and dynamic loads.

The nozzle plate (shown in Fig. 9-75) is a 0.5-m-thick, 5-m-radius circular disk that is simply supported at its outside edge by a series of gussets attached to the first structural wall. It is constructed of 33-mm-thick, extruded angle plates welded together. The part of the plate through which the jet array passes contains constant cross section vertical channels. The structure of the outer portion is similar, except for the bottom and top which are closed by flat plates that prevent flow-through of lithium and give additional strength. Some lithium circulation does occur in this outer region with the result that the maximum metal temperature remains below 800 K. This circulation flow can be discharged to cool the first structural wall, graphite reflector, and vacuum vessel as well. A central hole allows rapid equalization of pressure from each fusion reaction on the top and bottom of the nozzle plate. An annular lithium weir flow exists over the edge of the central hole

so that the inner areas of the nozzle plate are protected from x rays and neutrons. Intentional geometric perturbations in the weir area are shown schematically in Fig. 9-75. The optimum shape of these perturbations would be established empirically to offer maximum shielding and yet force the weir to break up into jets before reaching the reaction chamber midplane.

The shape of the vertical channels is generally rectangular. Radial webs between the channels are located so that a line of sight from the location of the fusion reactions to the first structural wall passes through four or more jets. The flow separates at the entrance of each channel, and then falls under the influence of gravity through the reaction chamber. Hence, fluid flow is decoupled from the structural considerations that dictate the nozzle plate thickness. There are additional jets (not shown in Fig. 9-75)

that shield the laser-beam and pellet-injection regions.

A major advantage of this nozzle plate design is its flexibility. The pattern of the jet array will undoubtedly change as more analytical results and experimental data become available. Any reasonable jet pattern can be accommodated with minor revisions to the nozzle plate inner jet flow region. The plate thickness can thus be adjusted to assure adequate structural performance. In fact, changes in the jet pattern can be incorporated with comparatively little difficulty even after the reaction chamber has been operated, by merely lifting the nozzle plate out of the chamber and replacing it with an alternate version.

We established the basic design of the nozzle plate to yield a jet array (including the weir flow) that has a 550-mm inner radius and a 2.5-m radial thickness at the reaction chamber vertical midplane. The packing fraction at the midplane is $\sim 35\%$, and results in ~ 0.8 m effective radial thickness for shielding purposes. We could have chosen a larger inner radius, but this would require a larger flow rate and pumping power for the same shielding protection. A smaller inner radius is not permissible because this would result in excessive stresses in the first structural wall.⁹⁶

We selected a 5.5-m/s downward jet velocity at the bottom of the nozzle plate so that the jets are reestablished throughout the chamber before the next fusion reaction occurs 1 s later. In actuality, the jets only need to be reestablished in the upper 7 or 8 m of the chamber prior to the next fusion reaction because the lower portion contains a lithium pool. This conservatism allows for additional disruptions of the jets that we have not yet fully analyzed (c.g., effects of gas pressure below the nozzle plate on the initial jet velocity in the first ten milliseconds or so following a fusion reaction). The initial velocity dictates the head of lithium that must be present above the nozzle plate. Under ideal conditions this head would be equivalent to the dynamic head. A dynamic head of ~ 1 m corresponds to a velocity of 4.5 m/s at the top of the nozzle plate, which, because of gravity, gives the desired 5.5 m/s at the bottom of the nozzle plate. The actual magnitude of the head of lithium will need to be established empirically because of the complex flow pattern that

Fig. 9-75. Elevation and plan views of the HYLIFE nozzle plate.

exists in the upper pool. We do, however, round the entrances to the jets to minimize both losses and any contractions of the jets. Because jet-area contraction coefficients of unity are typical for such a rounded entrance, the jet would have the same dimensions as the channel at the top of the nozzle plate. The thickness of each jet decreases as it falls. Using conservation of energy and the continuity relationship we can show that a radial jet thickness of 0.3 m at the nozzle plate becomes 0.15 to 0.2 m at the midplane, depending on the aspect ratio of the jet. Radial webs in the nozzle plate that separate the flow channels are positioned so that a direct line from the fusion reaction to the first structural wall passes through at least four of the six rows of jets. Additional curved jets are located around the beam port areas forming an arched pattern to shield the metal portions.

Experiments in progress at the University of California at Davis show that rectangular shaped jets oscillate in time with the major axis becoming the minor axis (and vice versa).⁹⁷ These experiments use water and relatively small jets. Preliminary theoretical and experimental results suggest that with the large lithium jets designed for HYLIFE, no appreciable variation in the jet axis will occur before the lithium reaches the bottom of the reaction chamber.

Breakup of the jets into droplets has also been found not to exist under HYLIFE conditions. Hoffman^{97,98} and Kusui⁹⁹ have shown that the breakup length is related by

$$L/D = 55 + 1.085 (We)^{0.5} \quad (68)$$

where L is the breakup length, D is the characteristic diameter of the jet, and We is the Weber number. For 0.3-m-thick lithium jets having an initial velocity of 4.5 m/s, the breakup length is 45 m. This is well beyond the 8 to 10 m working length of the reaction chamber, and it indicates the jets will remain intact.

Dynamic loads are imposed on the nozzle plate following each fusion reaction.¹⁰⁰ A downward pressure induced by the sudden deposition of neutron energy in the upper lithium pool during the first 300 μ s varies with radius as shown in Fig. 9-76. In addition, short-ranged x rays and pellet debris vaporize and superheat lithium in the

central area of the reaction chamber, which results in an upward pressure on the nozzle plate. This upward pressure also varies with space and time, as shown in Fig. 9-77. Additional static loads due to the dead weight of the nozzle plate itself and the liquid lithium are shown in Fig. 9-78. In the static load values we include a stress concentration factor of 3, which we consider to be conservative for this type of welded construction. Effects of the steady lithium-fluid acceleration from the upper lithium pool into the flow channels was considered when determining these static loads.

Fig. 9-76. Neutron energy-induced pressure on top of the nozzle plate.

Fig. 9-77. Superheated gas upward pressure loads on the nozzle plate.

Fig. 9-78. Static loading on the nozzle plate model.

We used the HEX^{101,102} three-dimensional element of the PLANS¹⁰³ finite element computer program to model a six-degree wedge of the nozzle plate. Orthotropic properties, which vary with radius, were used for 50 eight-noded, six-sided, isoparametric elements used in the study. These properties were based upon local material thickness ratios in the various element directions.

We circumvented the need to conduct a full, standard dynamic solution (by either normal-mode or transient-response methods) by using an approximate model approach. The method starts with the governing finite-element dynamic plate equations

$$[M] \{\ddot{\delta}\} + [K] \{\delta\} = \{F(t)\} \quad (69)$$

where $[M]$ and $[K]$ are the plate mass and stiffness matrices, respectively, $\{\delta\}$ are the time-dependent nodal displacements, and

$\{F\}$ are the transient nodal forces. We generated the mass and stiffness matrices by a standard finite element computer program once we entered the appropriate geometry and material properties.

We adopted the following solution procedure: the forcing function $\{F\}$ is decomposed into a series of products of spatial, $\{P_i\}$, and temporal, $f_i(t)$, functions

$$\{F\} \sim \sum_i \{P_i\} f_i(t) \quad (70)$$

We sought approximate solutions of the form

$$\{\delta\} \sim \sum_i \{\delta_i\} X_i(t) \quad (71)$$

where $\{\delta_i\}$ are assumed modes that satisfy the statical equations, i.e.,

$$[K] \{\delta_i\} = \{P_i\}, i=1, 2, \dots \quad (72)$$

Our use of static solutions assumes that the dynamic stresses and deflections may be represented as a linear combination of these distributions, where the dynamic stresses, $\{\sigma(t)\}$, are given by the stress matrix, $[S]$:

$$\{\sigma(t)\} = \sum \{\sigma_i\} = [S] \sum \{\delta_i\} X_i(t) \quad (73)$$

We substitute Eqs. (70), (71), and (72) into Eq. (69) and equating in terms of equal i to obtain

$$[M] \{\ddot{\delta}_i\} + [K] \{\delta_i\} \sim \{P_i\} f_i = [K] \{\delta_i\} f_i, \quad i=1, 2, \dots \quad (74)$$

Multiplying Eq. (74) by $\{\delta_i\}^T$ and defining ω_{Ri}^2 as the Rayleigh quotient gives the transpose of the nodal displacement vector:

$$\omega_{Ri}^2 = \frac{\{\delta_i\}^T [K] \{\delta_i\}}{\{\delta_i\}^T [M] \{\delta_i\}} = \frac{\{P_i\}^T \{\delta_i\}}{\{\delta_i\}^T [M] \{\delta_i\}}, \quad i=1, 2, \dots \quad (75)$$

and dividing Eq. (74) by $\{\delta_i\}^T [M] \{\delta_i\}$ yields

$$\ddot{X}_i + \omega_{Ri}^2 X_i = \omega_{Ri}^2 f_i(t), i=1, 2, \dots \quad (76)$$

Starting with the forcing functions shown in Figs. 9-74 and 9-77, we obtained statical-mode shapes $\{\delta_i\}$; Rayleigh quotients ω_{Ri}^2

Fig. 9-79. Static stresses near the bottom surface.

Fig. 9-80. Dynamic stress vs time for the inner radius of the nozzle plate: Top and bottom surfaces.

were obtained through solution of Eqs. (72) and (75), respectively. Then we solved Eq. (76) by a computerized numerical-integration scheme subject to the quiescent initial conditions

$$X_i(0) = 0 \quad , \quad (77a)$$

$$\dot{X}_i(0) = 0 \quad , \quad (77b)$$

and we combined the results as indicated by Eq. (71) to obtain an approximation for $\{\delta(t)\}$ and the stresses.

Because basic parameters of the present nozzle-plate configuration will change as the conceptual design develops, it would be convenient to identify the influence that such changes would have upon primary stresses. As may be seen from Eqs. (71), (72), and (76), the dynamic plate stresses (σ_d) are proportional to the Rayleigh quotient ω_R^2 , and the modal stresses, (σ_a) associated with the "assumed" trial mode shape, i.e.

$$\sigma_d \sim \omega_R^2 \sigma_a \quad , \quad (78)$$

where

$$\sigma_a \sim q_{\text{dynamic}} \left(\frac{a}{h} \right)^2 \quad . \quad (79)$$

In Eq. (79), q_{dynamic} is a measurement of the dynamic-loading intensity, h is the plate thickness, and a is the plate outer radius.

For a given plate design

$$\omega_R^2 \sim \frac{Eh^2}{\rho a^4} \quad , \quad (80)$$

where E is the effective stiffness modulus of the plate material, and ρ is the mass density.

The static stresses (σ_s), however, depend primarily on the loading parameter (q_s) and the geometric parameters a and h as follows:

$$\sigma_s \sim q_s \left(\frac{a}{h} \right)^2 \quad . \quad (81)$$

Therefore, increasing the outer plate radius from a to a' will have the effect of increasing the static stresses in the ratio of $(a'/a)^2$, while decreasing the Rayleigh quotient in the ratio $(a/a')^4$. This reduction in ω_R^2 would result in a decrease in σ_d by $(a/a')^2$, since σ_a [see Eq. (79)] will increase in the ratio $(a/a')^2$.

Similarly, an increase in h to h' will cause σ_a and σ_s to decrease in the ratio $(h'/h)^2$. However, it will leave σ_d unchanged, since the decrease in σ is offset by a corresponding increase in ω_R^2 . We point out that if the increase in h to h' increases q_{static} , this should be accounted for as well when computing σ_s .

We generated results for a nozzle plate of slightly different dimensions than those described here; we adjusted these by use of the stress-scaling laws. We obtained a peak static stress of 60 MPa at the inner radius near the bottom of the nozzle plate, as shown in Fig. 9-79. Circumferential (θ direction) static stresses near the top are essentially the negative of those shown. The absolute values of radial and axial stresses (r and z directions) on the top surface are equal to, or less than, those shown. The peak dynamic stress of 10 MPa is reached (as shown in Fig. 9-80) at the top and bottom surfaces of the inner radius. This dynamic stress includes the effect of the nonhomogeneous reduced area in the plate.

Static and dynamic peaks both occurred at the same location on the nozzle plate (i.e., at the inner radius of the plate, near the top and bottom surface), and were both circumferential components of stress. The spatial stress distributions shown in Fig. 9-79 are for a uniform pressure distribution, but all the loadings used gave similarly distributed stresses. In fact, all the trial modes yielded Rayleigh quotient frequencies of around 60 Hz.

Choice of nozzle plate thickness other than 0.5 m would have resulted in a different magnitude of static stress. (Dynamic stress magnitudes are independent of nozzle plate thickness.) A thickness of 0.5 m was selected because of limitations in fatigue crack growth (see "Fatigue Crack Growth in HYLIFE Components"). A static stress of 60-MPa coupled with a fluctuating dynamic stress of ± 10 MPa results in an allowable initial crack length of about 20 mm. This value can be easily detected by ultrasonic techniques. A nozzle plate thickness less than 0.5 m would have resulted in increased static stresses and marginal ability to detect flaws in the multitude of welded areas present.

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Fatigue Crack Growth in HYLIFE

Components. Major components of the HYLIFE reaction chamber will undoubtedly be of welded construction. Flaws in the weld areas (or in the base metal itself) are sites of local stress concentration which form cracks. These cracks can grow under the fatigue loading conditions present in the HYLIFE reaction chamber. On one hand, if the initial crack size is large enough, a critical crack length can be reached before the end of the desired lifetime, resulting in structural failure. On the other hand, a threshold value

of crack length (or stress-intensity-factor) exists, below which a crack will not propagate at all.

Application of fatigue-crack-growth analysis, with documented properties for the 2-1/4 Cr-1 Mo ferritic steel material used in HYLIFE, shows that the initial crack size necessary to cause structural failure within the desired lifetime can be detected with available inspection techniques. Hence, the major HYLIFE components can be manufactured as welded structures without fatigue crack growth limitations. However, manu-

Fig. 9-81. Effect of stress ratio on fatigue crack propagation of 2-1/4 Cr-1 Mo steel at 510°C in air.
[Reproduced from *Nucl. Tech.* 28, 501 (March 1976).]

fracture of the first structural wall will need to be completed with special care to ensure that defects are below allowable limits.

The basic fatigue-crack-growth equation¹⁰⁴ has the form

$$\frac{da}{dN} = A(\Delta K_{eff})^n, \quad (82)$$

where

- a = crack length, m;
- N = number of cycles;
- K_{eff} = fluctuation of the effective stress-intensity-factor;
- A and n = constants.

The effective stress-intensity-factor for elastic materials in general and in particular 2-1/4 Cr-1 Mo ferritic steel is related by¹⁰⁵

$$\Delta K_{eff} = \Delta\sigma \sqrt{\frac{\pi a}{1-R}}, \quad (83)$$

where

- $\Delta\sigma$ = peak stress fluctuation in a cycle, Pa;
- R = ratio of minimum to maximum stress in a cycle.

Under equilibrium conditions $\Delta\sigma$ and R are constants so that

$$\frac{da}{dN} = B a^m, \quad (84)$$

where

$$B = A \Delta\sigma^n \left(\frac{\pi}{1-R} \right)^m, \quad (C)$$

$$m = n/2.$$

Although Eq. (84) is nonlinear, it may be solved¹⁰⁶ by setting

$$a = z^\beta \quad (85)$$

and

$$\frac{da}{dN} = \beta z^{\beta-1} \frac{dz}{dN}. \quad (86)$$

Substituting into Eq. (84)

$$\frac{dz}{dN} = \frac{B}{\beta} z^{\beta(m-1)+1}. \quad (87)$$

If we set the exponent equal to zero, Eq. (87) is linear with the solution

$$z = B(1-m)N + C \quad \text{with} \quad \beta = \frac{1}{1-m}, \quad (88)$$

where C is a constant of integration.

Combining Eqs. (85) and (88) gives the final solution

$$a = [B(1-m)N + C]^{\frac{1}{1-m}}. \quad (89)$$

This solution can be confirmed by substitution into Eq. (84). The constant of integration is evaluated by the initial condition

$$a_{N=0} = a_0, \quad (90)$$

which gives

$$C = a_0^{1-m}. \quad (91)$$

The values of A and m are found from experimental data,¹⁰⁵ with B determined for specific values of $\Delta\sigma$ and R from its definition.

A difficulty with this formulation concerns the applicability of Eq. (83) under HYLIFE conditions. A plot of experimental data at 510° C for 2-1/4 Cr-1 Mo steel is reproduced from Brinkman et al.¹⁰⁵ as Fig. 9-81. The data on the left side is plotted using a stress-intensity-factor

$$\Delta K = \Delta\sigma \sqrt{\pi a}. \quad (92)$$

Table 9-22. Tabulation of results of the two-step process to calculate the initial crack length (see p. 9-78).

$\Delta\sigma$ (Pa)	R —	Step 1		Step 2		
		$\Delta K_{eff} > 18 \times 10^6$ Pa $\sqrt{m}$		$\Delta K_{eff} < 18 \times 10^6$ Pa $\sqrt{m}$		
		a_{knee} (m)	N_1 (cycles)	N (cycles)	a_0 (m)	$a_{threshold}$ (m)
120×10^6	-1	14.3×10^{-3}	1.6×10^5	10^8	3.6×10^{-3}	2.8×10^{-3}
120×10^6	0	7.16×10^{-3}	7.8×10^4	10^8	1.7×10^{-3}	1.4×10^{-3}
120×10^6	0.5	3.58×10^{-3}	3.9×10^4	10^8	8.2×10^{-4}	7.0×10^{-4}
120×10^6	0.9	7.16×10^{-4}	7.8×10^3	10^8	1.5×10^{-4}	1.4×10^{-4}
120×10^6	-1	14.3×10^{-3}	1.6×10^5	10^9	3.1×10^{-3}	2.8×10^{-3}
120×10^6	0	7.16×10^{-3}	7.8×10^4	10^9	1.5×10^{-3}	1.4×10^{-3}
120×10^6	0.5	3.58×10^{-3}	3.9×10^4	10^9	6.9×10^{-4}	7.0×10^{-4}
120×10^6	0.9	7.16×10^{-4}	7.8×10^3	10^9	1.2×10^{-4}	1.4×10^{-4}
20×10^6	-1	8.59×10^{-2}	5.6×10^6	10^8	1.6×10^{-1}	1.0×10^{-1}
20×10^6	0	4.30×10^{-2}	2.8×10^6	10^8	8.0×10^{-2}	5.1×10^{-2}
20×10^6	0.5	2.15×10^{-2}	1.4×10^6	10^8	3.7×10^{-2}	2.5×10^{-2}
20×10^6	0.9	4.30×10^{-3}	2.8×10^5	10^8	6.7×10^{-3}	5.1×10^{-3}
20×10^6	-1	8.59×10^{-2}	5.6×10^6	10^9	1.4×10^{-1}	1.0×10^{-1}
20×10^6	0	4.30×10^{-2}	2.8×10^6	10^9	6.8×10^{-2}	5.1×10^{-2}
20×10^6	0.5	2.15×10^{-2}	1.4×10^6	10^9	3.2×10^{-2}	2.5×10^{-2}
20×10^6	0.9	4.30×10^{-3}	2.8×10^5	10^9	5.8×10^{-3}	5.1×10^{-3}

The same data is plotted on the right-hand side using the effective stress-intensity-factor, k_{eff} . Brinkman's k_{eff} is identical to the ΔK_{eff} defined in Eq. (83). Inclusion of the term $(1 - R)$ in ΔK_{eff} permits a good correlation of the data when ΔK_{eff} is greater than about $2 \times 10^7 \text{ Pa}\sqrt{\text{m}}$. There is a knee in the curve at about this point for data taken at both 510°C and 590°C. (Brinkman only plots data at 590°C using ΔK , not ΔK_{eff} .) Below the knee, the correlation is not as accurate. In fact, a threshold value of ΔK_{eff} is reached below which no crack growth occurs.¹⁰⁷ For 2-1/4 Cr-1 Mo steel the threshold is at a value of ΔK_{eff} slightly below the location where the knee appears. We used the threshold value of

$$\Delta K_{\text{eff}} = 8 \times 10^6 \text{ Pa}\sqrt{\text{m}} \quad (93)$$

Fig. 9-82. Variation of tolerable initial crack length, a_0 , with the ratio of minimum to maximum stress, R , for various peak stress fluctuations, $\Delta\sigma$, and lifetimes, N .

for three reasons: first, this value is below the knee at values of R up to 0.5 and temperatures of both 510 and 590°C; second, it is consistent with data at room temperature;¹⁰⁷ and third, it appears to be conservative for an environment that contains no oxidizing elements.

Quantitative values of both A and m were determined from the data of Ref. 105 at 590°C and 40 Hz to be

$\Delta K_{\text{eff}} < 18 \times 10^6$	$\Delta K_{\text{eff}} > 18 \times 10^6$
$A = 4.34 \times 10^{-215}$	2.15×10^{-27}
$m = 14.6$	1.876

(units of Pa and m)

The maximum initial crack size that can be tolerated without failure is determined by a two-step process. First, for given values of $\Delta\sigma$ and R , the value of crack length, a_{knee} , is calculated at the knee where $\Delta K_{\text{eff}} = 18 \times 10^6 \text{ Pa}\sqrt{\text{m}}$. Using a_{knee} as the initial crack length, the value of N is calculated that makes the bracketed term $[B(1 - m)N + C]$ in Eq. (89) with constants for $\Delta K_{\text{eff}} > 18 \times 10^6$ equal to zero. This value of N , labeled N_1 , corresponds to the number of cycles necessary for a crack of initial length, a_{knee} , to reach an infinite length. The value of $(1 - m)$ is negative and the solution is undefined for $C < B(1 - m)N$. A failure criterion where $C > B(1 - m)N$ could have been utilized, resulting in a finite crack length at the end of the lifetime for the component. However, the difference in a_{knee} for a reasonable end of lifetime crack length as compared to that calculated herein is insignificant and well within the uncertainty of the data.

Second, N_1 is subtracted from the HYLIFE total fatigue lifetime (i.e. 10^9 cycles). We then calculate C from Eq. (89) by using this result and setting $a = a_{\text{knee}}$. The initial crack length, a_0 , is then found using Eq. (91). The result is compared to the threshold value of a_0 , and the larger selected as the final result.

Selected results are tabulated in Table 9-22 (p. 9-77). The left-hand portion is associated with crack growth where ΔK_{eff} is larger than $18 \times 10^6 \text{ Pa}\sqrt{\text{m}}$ (above the knee). This portion of the calculation results in the number of cycles, N_1 , after the crack had reached a length, a_{knee} , until failure occurs. The center portion is associated with crack growth below the knee. For a total desired

lifetime, N , taken to be either 10^8 or 10^9 cycles, the initial crack length a_0 is determined. An initial crack length less than a_0 would not propagate to failure within the lifetime indicated. Values of threshold crack length below which no crack growth occurs are indicated on the right-hand side. Note that $a_{\text{threshold}}$ is less than, or very close to, a_0 in every case, indicating that the value of a_0 may be used as the maximum tolerable initial crack length.

This same information is shown in graphical form on Fig. 9-82, where maximum tolerable initial crack length is plotted as a function of the ratio of minimum to maximum stress. The effect of a decreased lifetime to 10^8 cycles can be seen by comparison to the red lines. An asymptote of $a_0 \rightarrow 0$ is reached as $R \rightarrow 1$ and the validity of the results in this neighborhood is questionable. The magnitude, however, is conservative. The fact that this assumption is conservative can be recognized by considering the case of a nearly steady load where the value of R would be near unity, yielding a maximum tolerable initial crack length approaching zero. In actuality, a finite crack length can be tolerated as long as the magnitude of the steady load is below the creep strength (or other failure criterion) of the material.

Placed in Fig. 9-82 are two circles that represent the design points for the HYLIFE first structural wall and nozzle plate. A detectable crack length of about 1 mm is feasible for 2-1/4 Cr-1 Mo steel using ultrasonic techniques. If the tolerable initial crack length is larger than that which can be detected, any flaw can be repaired before the component is placed in service. This is the case as can be seen in Fig. 9-82.

The difficulty in manufacturing an acceptable part is another matter. Parts with defects less than 10 mm long can be manufactured with little difficulty, but to manufacture a part with defects less than 1 mm requires substantial care and extra expense. It can be concluded that the nozzle plate may be manufactured with normal attention to welding processes, but that the first structural wall will require meticulous care.

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Allowable Driver and Target Factory Cost

Accurately predicting the future cost of fusion generated electricity is a difficult task from the current stage of fusion development. It is particularly sobering to recall that history has demonstrated that pathbreaking technologies, such as fusion generally turn out to be more expensive than originally conceived. Mindful of this tenet, we do not attempt here to predict future electricity costs for ICF; rather, the purpose of this study is to determine the allowable capital costs corresponding to the electricity costs that seem reasonable. We assume the cost of electricity is dominated by capital requirements for the plant, with operation, maintenance, and fuel cycle costs (other than the target factory) being a small factor. The cost of electricity as used here is the cost at the busbar and does not include transmission and distribution.

If we knew nothing about the cost of an ICF power plant this study would consist of a single equation relating the total allowable cost to the cost of electricity. Fortunately, we do have a basis on which to establish a worthwhile parametric study. In 1978, Bechtel National completed a cost analysis based on the conceptual design of HYLIFE.¹⁰⁸ Their study included only the reactor plant, that is the HYLIFE chamber and balance of plant; it did not include the driver or target factory.

The HYLIFE reactor plant is, in many respects, quite similar to a liquid-metal-fast-breeder reactor (LMFBR). France has already built a LMFBR demonstration plant and is currently constructing SUPERPHENIX, a commercial size system.¹⁰⁹ Bechtel National was involved in cost studies for a similar reactor, the Prototype-Large-Breeder Reactor. Bechtel estimates that the HYLIFE chamber and associated lithium circulation system will account for ~20% of the total reactor plant cost. The remaining balance-of-plant structures and components are comparable to those in an LMFBR. For these reasons, we feel the HYLIFE reactor plant cost basis is quite good.

If we are willing to evaluate ICF economics on the basis of a power plant that incorporates a HYLIFE reaction chamber,

we can then determine the allowable capital cost for the driver and target factory as a function of how much we are willing to pay for electricity. Similar economic analyses can be performed for any ICF reactor concept for which the cost scaling relationships are known. We consider not only the allowable cost but also try to gain some insight into the best operating point for the plant in terms of the driver energy, pulse repetition rate and target yield. Clearly, this is not simply the point of maximum allowable cost since both the driver and target factory costs scale with the varying parameters.

Methodology. Bechtel's cost analysis for the HYLIFE reactor plant was based on a point design, that is a specific set of operating parameters. Some of these are listed in Table 9-23. This past year, Hovingh and Blink¹¹⁰ developed scaling laws for the HYLIFE chamber that allow us to determine the lithium flow rate as a function of the chamber pulse repetition rate and target yield. Assuming the cost of the lithium circulation system is proportional to the flow rate then, we can scale the reactor plant cost in relation to repetition rate and yield using Bechtel's point-design cost as the basis.

The chamber fusion power is held constant at 2700 MW_f. The resulting thermal power and gross electric power are 3130 MW_t and 1220 MW_e. By holding the power level constant we fix the cost of the balance of plant (excluding the lithium circulation system) and thus avoid the possibility of scaling this major cost factor incorrectly. Since the fusion power is constant, the chamber repetition rate is inversely proportional to the target yield.

$$R = 2700/Y \quad (94)$$

The size of the HYLIFE chamber is also fixed with a 10-m diameter and 8-m height from the nozzle plate to the lithium pool at the bottom of the fusion chamber. The effective thickness of the lithium jet array is 0.74 m at the chamber midplane. The HYLIFE scaling relationships determine the

position of the jet array and the repetition rate determines the flow velocity to give us the dependence of flow rate and pumping power on repetition rate and yield.

The next step is to determine the driver beam energy required to produce the desired target yield and the electric power required by the driver. This requires a relationship between the driver energy and target gain G(E). We have selected a gain verses driver energy curve that is midway between the conservative gain and optimistic gain estimates. The yield is simply

$$Y = E \cdot G(E) \quad (95)$$

and the electric power consumed by the driver is

$$PD = R \cdot E/\eta \quad (96)$$

where η is the driver efficiency.

The cost of electricity is proportional to the total capital cost of the power plant divided by the net electric power. The total cost is the sum of the reactor, driver, and target factory costs. The net electric power is equal to the gross electric power minus the recirculating power requirements. The recirculating requirements consist of the auxiliary power needed to run the plant, the pumping power for circulating lithium, and the power consumed by the driver.

The busbar cost of electricity can be expressed as

$$COE = \frac{2.45 (TRC + TDC + TTFC)}{(PG - PA - PP - PD)} \text{ cents/kWh} \quad (97)$$

where

- TRC = Total reactor cost, \$M
- TDC = Total driver cost, \$M
- TTFC = Total target factory cost, \$M
- PG = Gross electric power, MW_e
- PA = Auxiliary power, MW_e
- PP = Pumping power, MW_e
- PD = Driver power, MW_e

The constant 2.45 is a conversion factor based on an annual fixed charge rate of 15% and a capacity factor of 70%.

We are interested in relating the allowable sum of the total capital cost of the driver and target factory (TDC + TTFC) to the cost of electricity. This relationship will be a function of more than one parameter. The reactor cost, TRC, and lithium pumping power, PP, can be related to a single parameter; either the driver energy or chamber repetition rate which are coupled

Table 9-23. Point design parameters for Bechtel cost study.

Target yield	2700 MJ
Pulse repetition rate	1.0 Hz
Fusion power	2700 MJ
Jet array equivalent thickness	1.0 m
Jet array inner radius	0.5 m
Total lithium flow rate	129 m ³ /s

Accomplishments Associated with the HYLIFE Chamber Concept

by Eqs. (94) and (95). This follows from the fact that the reactor cost and lithium pumping power are determined by the lithium flow rate which we previously noted was a function of yield and repetition rate. The driver power, PD, is a function of the driver energy (or repetition rate) and driver efficiency as seen from Eqs. (94) through (96).

The gross electric power (PG) is constant, and auxiliary power, which is proportional to PG, is therefore also fixed. Consequently, we investigate the allowable cost of the driver and pellet factory in terms of three parameters; the driver energy (or chamber repetition rate), the driver efficiency, and the desired cost of electricity. We also examine the economic advantages of sharing a single driver and target factory between two full size HYLIFE reactor plants, and the effects of using a more conservative target-gain/driver-energy relationship.

HYLIFE Cost Basis. A breakout of the *direct* capital cost of the HYLIFE reactor plant is given in Table 9-24. The costs from the Bechtel study¹¹¹ were escalated by 13.8% to convert them from third-quarter 1978 dollars to first-quarter 1980 dollars. This assumes an annual cost escalation rate of 9%.

Included in the HYLIFE primary heat transfer system costs are the pumps and piping required to circulate lithium through the chamber and to the intermediate heat exchangers for heat removal. The flow through the chamber is primarily for establishing the jet array and is much larger than the flow required for heat removal. Therefore, we subdivide the 200.7 \$M cost element into two factors: 159.2 \$M is associated with circulating lithium, and 41.5 \$M is assigned to the primary heat removal circuit. (The 41.5-\$M figure was chosen by noting that the pumps, motors, and piping for the secondary heat transfer system cost 41.5 \$M.)

The total capital cost, which must be used in determining the cost of electricity, is the sum of the direct capital costs, indirect capital costs, and time related costs. When all these factors are included we find that the total capital cost for the reactor plant in 1980 dollars is 2.05 times the direct capital cost.¹¹¹

Care must be taken in comparing capital costs. The allowable driver plus target factory cost discussed in this study is the total capital cost. To determine the allowable direct capital cost, the total capital cost must

be divided by a factor that depends on the assumptions made for indirect and time related costs. For example, we estimate that the total cost for these systems will be 1.8 times the direct cost. This is different than the factor of 2 used for the reactor plant since we assume a shorter construction period (5 years compared to 10 years for the reactor plant) and a lower escalation rate (7% compared to 9%) apply to these systems.

The variation in the lithium flow rate with chamber repetition rate and target yield is shown in Fig. 9-83. A minimum flow rate of 82 m³/s is obtained at 1.5 Hz and 1800 MJ.

Table 9-24. Direct capital costs for HYLIFE reactor plant systems.

	\$M
STRUCTURES AND SITE FACILITIES	297.4
Site improvement	12.1
Reactor building	137.7
Reactor auxiliary building	59.3
Turbine building	31.9
Miscellaneous buildings	56.4
REACTOR PLANT EQUIPMENT	552.2
Reactor vessel and internals	70.7
Primary heat transfer system	
Pumps, motors and piping	200.7
IHX	59.1
Secondary heat transfer system	
Pumps, motors and piping	41.5
Steam generators	129.1
Other reactor plant equipment	51.1
TURBINE PLANT EQUIPMENT	174.6
Turbine-generator	79.9
Heat rejection system	28.3
Condensing system	14.5
Other turbine plant equipment	51.9
ELECTRIC PLANT EQUIPMENT	77.0
MISCELLANEOUS PLANT EQUIPMENT	17.0
TOTAL DIRECT COST	1118.2

Fig. 9-83. The minimum lithium flow rate is 82 m³/s at 1.5 Hz and 1800 MJ.

Accomplishments Associated with the HYLIFE Chamber Concept

The point design on which Bechtel based their cost analysis had a lithium flow rate of $129 \text{ m}^3/\text{s}$ for a chamber repetition rate of 1.0 Hz and an equivalent lithium array thickness of 1.0 m at the midplane. The $82\text{-m}^3/\text{s}$ flow rate shown here is considerably lower since the effective array thickness has been reduced to 0.74 m and the lithium flow in the blanket region has been reduced from $\sim 36 \text{ m}^3/\text{s}$ to $\sim 13 \text{ m}^3/\text{s}$.

The inner radius of the annular-shaped lithium jet array must be increased with increasing target yield.¹¹⁰ This results in an increased flow area to maintain an equivalent lithium thickness of 0.74 m at the midplane. When the target yield is increased

above 1800 MJ , the increase in flow area outweighs the decrease in injection velocity, which is possible with the lower repetition rate, and the flow rate therefore increases.

When the target yield is decreased below 1800 MJ , the injection velocity required to reestablish the array in the shorter interpulse time increases faster than the flow area decreases. Again the result is an increase in flow rate.

Figure 9-84 shows how the total capital cost of the lithium circulation system varies with the chamber repetition rate and target yield. As previously noted, it is taken to be directly proportional to the flow rate. The remaining reactor plant cost ($1960 \text{ \$M}$) is constant. The lithium circulation pumping power is a function of both the flow rate and injection velocity. The high injection velocities required to reestablish the jet array at high repetition rates cause the pumping power to increase sharply beyond a few hertz.

We see that both the cost of the HYLIFE reactor plant and the recirculating power required for lithium pumping minimize at $\sim 1.5 \text{ Hz}$. If the cost of electricity was based solely on these factors, it too would be minimized at 1.5 Hz . At this point, the total reactor plant cost is $2170 \text{ \$M}$, and the electric power available is 1130 MW_e (the auxiliary power is 76 MW_e , and the pumping power is 14 MW_e). This represents a busbar electricity cost of 4.7 cents/kWh (1980 dollars). Inclusion of the costs and power requirements of the driver and target factory will result in incrementally higher electricity costs.

Figure 9-85 shows the base from which we must build. It shows the resulting cost of electricity as a function of repetition rate when the cost and power requirements of the driver and target factory are set to zero. This base is already higher than projected electricity costs from coal and nuclear power plants in the year 2000. The Electric Power Research Institute (EPRI) reports coal-electric power at $\sim 4 \text{ cents/kWh}$ and nuclear-electric power slightly less than 3 cents/kWh (in 1979 dollars).¹¹² New energy technologies such as fusion, solar, and fission-breeders are all expected to be more expensive than current systems. We argue, however, that society will be willing to, and eventually may be required to, pay more for these inexhaustible energy sources. Taking 1.5 to

Fig. 9-84. The lithium circulation system cost is proportional to the lithium flow rate. The pumping power is a function of flow rate and injection velocity.

Fig. 9-85. Cost of electricity due only to the HYLIFE chamber and balance of plant.

2 times the electricity cost from a coal-burning power plant (4 cents/kWh) implies that costs up to the 6 to 8 cents/kWh range are of interest. At the same time, it is clear that we must continue to innovate the HYLIFE reactor plant design with a view toward reducing costs. This work is currently underway in conjunction with Bechtel.

Net Electric Power. To determine the recirculating power required for the driver we need to establish a relationship between target gain and driver beam energy on target. Figure 9-86 is a plot of target gain vs driver energy.¹¹³ The national ICF program expects to achieve the conservative levels initially and then follow a learning curve to more optimistic levels. These initial levels are based on conservative assumptions about implosion instabilities, thermonuclear ignition and burn, assuming efficient beam-target coupling (with ions or short-wavelength lasers). It should be pointed out that target gain depends on factors other than just the driver energy. Implicit in these curves are additional requirements on driver parameters such as pulse shape, pulse width, and target design parameters.

For this study we have selected a gain curve, referred to here as the reference gain, that lies between the conservative and optimistic estimates. The reference gain is equal to 4 times the conservative gain. In the range of 1 to 4 MJ, it is very close to the average of the optimistic and conservative gain curves shown here. For example at 2.2 MJ, the conservative gain is 100, the optimistic gain is 700, and the reference gain is 400.

The net electric power for a HYLIFE power plant using the reference gain curve is shown in Fig. 9-87 as a function of the driver energy and repetition rate for different driver efficiencies. Indicated in this figure is the gross electric power of 1220 MW_e and the auxiliary power requirement of 76 MW_e. This leaves 1144 MW_e from which the driver power and lithium pumping power must also be subtracted.

At low driver energies the net electric power rises sharply with an increase in energy as the recirculating power requirements for both the driver and the lithium pumping decrease rapidly. The lithium pumping power reaches a minimum at a driver energy of ~3.4 MJ which corresponds to 1800 MJ at 1.5 Hz. The

driver power continues to decrease but less rapidly as the gain curve tends to flatten out beyond 4 MJ. As a result, the net electric power curves are very flat beyond 4 to 6 MJ. For the 25% and 10% efficient drivers, the net electric power actually decreases beyond 4 and 6 MJ, respectively, as the increased pumping power outweighs the decrease in driver power. At 4 MJ (1.15 Hz) the net electric power produced using a 25% efficient driver is 9% more than that produced using a 4% efficient driver. It is clear that the 2- to 4-MJ range for driver energies will be the region of interest for different driver efficiencies shown here (and for the assumed

Fig. 9-86. Reference gain curve used in this study is midway between the optimistic gain and the conservative gain.

Fig. 9-87. Net electric power from the HYLIFE power plant as a function of driver energy and chamber repetition rate for various driver efficiencies. ▼

reference gain curve). This corresponds to chamber repetition rates of ~ 4 to 1 Hz.

Driver Plus Target Factory Allowable Cost.

If the reactor plant cost, which has a component that varies as indicated in Fig. 9-85, is coupled with the net electric power curves from Fig. 9-87, we can determine the allowable total capital cost of the driver plus target factory for various electricity costs as presented in Eq. (97). The results are

Fig. 9-88. Allowable total capital cost for the driver plus target factory to give the indicated cost of electricity.

shown in Fig. 9-88. The allowable total capital cost in billions of dollars as a function of driver energy and repetition rate is given for various electricity costs and a range of driver efficiencies. The curves all peak in the range of 3 to 4 MJ (1.9 to 1.2 Hz). This is expected since the reactor plant cost is a minimum at 3.4 MJ, and the net electric power is over the knee of the curve at 4 MJ. At a busbar cost of 6 cents/kWh, the peak allowable costs are in the 0.3 to 0.5 \$B range, and at 8 cents/kWh they are in the 1.1 to 1.5 \$B range. Recall that these are total capital costs and must be divided by a factor of ~ 1.8 to convert to direct capital cost in current dollars.

If we take a vertical cut through these curves, we can plot the allowable total capital cost as a function of driver efficiency as is shown in Fig. 9-89. The three bands are for busbar electricity costs of 6, 7, and 8 cents/kWh. The upper bound of each band corresponds to a driver energy of 3.5 MJ and a chamber repetition rate of 1.4 Hz. The lower bound is for 2.5 MJ and 2.6 Hz.

At low driver efficiencies the curves rise steeply with increasing efficiency. The knee of the curve is in the 5 to 10% range. In general, the additional allowable cost in going from 5 to 7.5% is double the additional allowable cost in going from 7.5 to 10%. For example, using the upper bound of the 8 cents/kWh curve, the allowable cost at 5% is ~ 1.2 \$B; at 7.5% it is 0.1 \$B more or ~ 1.3 \$B. Increasing the efficiency further from 7.5 to 10% increases the allowable cost half as much, or 0.05 \$B for a total of 1.35 \$B.

The target gains at 2.5 MJ and 3.5 MJ are 415 and 550, respectively. Therefore, the corresponding fusion energy gain (the product of target gain and driver efficiency) in the 5 to 10% driver efficiency range is ~ 21 to 42 on the lower bound of the bands and ~ 28 to 55 on the upper bound.

The economic advantages of sharing a single driver and target factory with more than one full-size reactor plant are clear. First, there are savings in the construction cost of the reactors. For dual unit (two 1-GW_e) nuclear stations the second unit is costed at 88% of the first.¹¹⁴ Thus, the average cost of the reactor plant is reduced by 6%.

Also, it is generally held that the cost penalty for increasing the driver repetition rate at constant beam energy is small, if not

Fig. 9-89. Allowable total capital cost for the driver plus target factory for the indicated cost of electricity.

negligible, at least in the 1- to 10-Hz range. Therefore, the driver cost per unit of electric power is reduced by ~50% for a dual unit plant. There will, however, be some incremental cost added for the systems required to alternately switch the driver beams from one chamber to the other.

Finally, the cost of the target factory is expected to scale less than linearly with production rate.¹¹⁵ A target factory producing the same target at twice the rate may cost only 50% more. If so, this would result in a 25% reduction in the target factory cost per unit of electric power.

The primary disadvantage of the dual-unit plant is that the utility must make a larger initial investment. Capital formation for such plants may prove difficult especially in the early stages of deployment when the utilities are introducing a new technology.

The allowable total capital cost is shown in Fig. 9-90 for a driver and target factory that are shared in a dual unit power plant. The allowable cost is plotted as a function of the driver efficiency using the same upper and lower bounds as in Fig. 9-89. The repetition rate in each chamber is the same as before, 1.4 Hz at 3.5 MJ, and 2.6 Hz at 2.5 MJ. The driver repetition rate and target production rates are, of course, double these numbers. Note that the 6 cents/kWh band has shifted up dramatically. For example, the increase for the upper bound of the band at 7.5% is \$0.7 billion.

Operating Point. While Figs. 9-88 through 9-90 indicate the allowable cost for the driver plus target factory, they give little guidance as to the optimum design point. For example, the allowable cost at 3.5 MJ and 1.4 Hz is greater than at 2.5 MJ and 2.6 Hz, but a 3.5-MJ driver costs more than a 2.5-MJ driver, while the target factory is expected to cost less at 1.4 Hz than at 2.6 Hz. Which operating point is better? We cannot answer this question explicitly without specifying the individual costs of the driver and target factory and how these costs scale with parameters such as driver energy, repetition rate, and target yield.

We can, however, get some indication of the appropriate operating point by making some simplifying assumptions. First, we assume the cost of the target factory is constant. Secondly, we assume the cost of the driver scales with driver energy as

$$C = C_1 E^n \quad (98)$$

where

$$C_1 = \text{cost at } E = 1 \text{ MJ}$$

$$n = \text{scaling exponent} \leq 1.0$$

The cost of the target factory is subtracted from Fig. 9-82 to give the allowable cost of the driver as a function of energy. This is set equal to C in Eq. (98), and we solve for C_1 for various values of the scaling parameter, n . The plot of C_1 vs E shows a maximum which corresponds to the optimum operating

Fig. 9-90. Sharing a single driver and target factory between two full size reactor plants increases the allowable cost for a given cost of electricity.

Fig. 9-91. Maximum allowable total capital cost per MJ peaks at operating-point driver energies in the range of 2.6 to 3.3 MJ for a 4% efficient driver.

point. It is the point that allows the maximum cost per megajoule.

As examples, we have taken the 8 cents/kWh curves from Fig. 9-88 for the 4% and 25% efficient drivers. The target factory cost was set to zero to give an upper bound on the allowable driver costs.

Figure 9-91 shows the allowable total capital cost per megajoule for the 4% efficient driver with scaling exponents

Fig. 9-92. Maximum allowable total capital cost per MJ peaks at operating point driver energies in the range of 2.0 to 2.6 MJ for a 25% efficient driver (see p. 9-86).

Driver efficiency %	Scaling exponent	Operating point driver energy ^a	Maximum ^b allowable cost per MJ
4	1.0	3.2 MJ	0.25 \$B
	0.8	3.3	0.32
	0.6	3.4	0.41
	0.4	3.6	0.52
10	1.0	2.6	0.40
	0.8	2.8	0.49
	0.6	2.9	0.60
	0.4	3.1	0.75
25	1.0	2.4	0.48
	0.8	2.5	0.57
	0.6	2.7	0.69
	0.4	3.0	0.85

^aTo the nearest 0.1 MJ.

^bTarget factory cost = 0. Cost given is the total capital cost. Direct capital costs are typically a factor of 1.8 smaller.

Table 9-25. Twin unit plant operating point and maximum allowable driver cost—6 cents/kWh—reference gain.

ranging from 0.4 to 1.0. Clearly, the more favorable the scaling, the higher the allowable cost per MJ. The arrows indicate the optimum operating points. These range from 2.6 MJ, for a driver with a cost that scales linearly with energy, to 3.3 MJ for $n = 0.4$. The tick marks on the curves indicate the points at which the allowable cost is at 95% of the maximum value. They give an indication of the range of operating point driver energies over which the allowable cost would be quite close to the optimum. The size of this range increases slightly with decreasing values of n . Using a nonzero cost for the target factory or including a cost that increased with increasing production rate would tend to shift the optimum operating points to higher driver energies and lower repetition rates. For example, with a constant target factory cost of 200 \$M, the operating point for $n = 0.6$ shifts slightly from 3.0 to 3.1 MJ. The allowable cost per megajoule decreases 0.1 \$B from 0.54 to 0.44 \$B. These represent direct capital costs of 0.30 \$B and 0.24 \$B, respectively.

Figure 9-92 shows the results for a 25% efficient driver. The allowable costs are, of course, higher for the same 8 cents/kWh electricity cost. The operating points have also shifted to lower driver energies compared with the 4% efficient driver. Again, including the target factory would tend to shift the operating points to higher driver energies.

The results for a dual unit plant at a lower cost of electricity, 6 cents/kWh, are tabulated in Table 9-25 for comparison. The table presents the maximum allowable total capital cost per MJ and the corresponding operating driver energy for various driver efficiencies and scaling exponents. The target factory cost is again taken to be zero. Operating point driver energies range from 2.4 to 3.6 MJ. For a given driver efficiency, the operating point energy increases as the cost scaling becomes more favorable (lower values of n). For a given scaling exponent, the operating point energy increases with decreasing driver efficiency.

Table 9-26 shows the effect of lower target gains on the optimum operating point and maximum allowable driver cost. Table 9-26 gives the results for a twin unit plant, at 8 cents/kWh, using the conservative gain

curve rather than the reference gain. We note that the operating points have shifted to higher driver energies, from 4.5 to 10 MJ, while the corresponding chamber repetition rates are ~ 3.7 Hz and ~ 1.3 Hz, respectively. Comparing these results to Figs. 9-91 and 9-92, which are for the same cost of electricity but with single unit plant, and with Table 9-25, which is for a dual unit plant but lower electricity cost, gives an indication of the impact of the lower gain curve on the allowable driver costs. The effects are, of course, most critical with the less efficient driver.

Summary. Using the HYLIFE reactor concept as a basis, we have examined the allowable total capital cost for the driver and target factory of an inertial confinement fusion power plant. The allowable total capital cost was evaluated in terms of the cost of electricity one is willing to pay, the driver efficiency, driver energy, and chamber pulse repetition rate. We found that the system parameters corresponding to the location of the maximum allowable cost fall within the range of $E = 3$ MJ, $R = 1.86$ Hz, and $Y = 1450$ MJ to $E = 4$ MJ, $R = 1.15$ Hz, and $Y = 2350$ MJ for the electricity costs and driver efficiencies that we examined. This location is strongly influenced by the fact that the HYLIFE chamber cost is a minimum at $R = 1.5$ Hz and $Y = 1800$ MJ. It also depends on the target gain relationship; higher gains move the location of the peak toward lower driver energies and lower gains move it toward higher driver energies.

The magnitude of the allowable cost increases with increasing driver efficiency and increasing cost of electricity. For the reference gain used here, the knee of the allowable cost vs driver efficiency curve is in the range of 5 to 10%. At $E = 3$ MJ, this corresponds to fusion energy gain, η_G , of 24 to 48.

An indication of the trends in location of the optimum operating point was given as a function of how the driver cost scales with energy. The optimum operating point was defined as the driver energy corresponding to the maximum allowable cost per MJ assuming a constant target factory cost. For zero target factory cost and 8 cents/kWh, the optimum operating point for a 4% efficient driver was found to be in the range

of 2.6 to 3.3 MJ (2.4 to 1.6 Hz) for cost scaling exponents $n = 1.0$ to 0.4, respectively. The optimum operating point energy decreases with increasing driver efficiency. The operating points for a 25% efficient driver ranges from 2.0 to 2.6 MJ for $n = 1.0$ to 0.4. Including the target factory cost will shift the optimum driver energy to higher driver energies and lower pulse repetition rates.

The economic implications of twin unit plants and lower target gains were also examined. Sharing a single driver and target factory between two full-size reactor plants results in significantly higher allowable costs for a given cost of electricity or lower electricity costs for a given driver plus target factory cost. On the other hand, it requires a larger initial capital commitment on the part of the utility. With more conservative gains, higher driver energies are required and the allowable cost is less. The decrease in allowable cost with decreasing target gain is most significant for low driver efficiencies.

In this study, we have taken the approach of determining the "allowable" cost for the driver and target factory assuming the reactor plant cost is given. This should not be construed as implying that reductions in the reactor plant cost are not being sought. Such cost reductions would either allow for lower electricity costs or permit larger allowable costs for the driver and target factory. Clearly, we need continued innovation in all the plant subsystems to

Driver efficiency %	Scaling exponent	Operating point driver energy ^a	Maximum ^b allowable cost per MJ
4	1.0	8.0 MJ	0.13 \$B
	0.8	8.5	0.20
	0.6	9.0	0.31
	0.4	10.0	0.48
10	1.0	4.5	0.39
	0.8	5.0	0.54
	0.6	5.5	0.71
	0.4	6.5	1.07
25	1.0	4.5	0.54
	0.8	4.5	0.73
	0.6	5.0	0.99
	0.4	5.5	1.36

^aTo the nearest 0.5 MJ.

^bTarget factory cost = 0. Cost given is the total capital cost. Direct capital costs are typically a factor of 1.8 smaller.

Table 9-26. Twin unit plant operating point and maximum allowable driver cost—8 cents/kWh—conservative gain.

minimize the ultimate cost of electricity and to contain cost growths which traditionally occur as such systems are developed.

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Engineering Test Facility (ETF) Planning

The engineering development required to harness energy output from inertial confinement fusion (ICF) is currently being addressed on a very small scale. However, experience with the development of other new technologies indicates that substantial engineering-development time will be required for ICF energy-conversion systems.

To begin planning such a development program, we invited representatives from much of the ICF reactor community to LLNL in July 1980 to define an ICF engineering test facility, and to discuss the reactor technology development that will be required before the engineering test facility (ETF) can become a reality. This meeting was notable for the spirit of cooperation and for the agreement on technical issues.

The following design goals and issues were established:

- There must be systems integration of all the major ETF systems: repetition rate driver, beam propagation, automated target fabrication, target injection, coolant circulation, and chamber exhaust. This goal is listed first because it is a prerequisite for accomplishment of the other goals.
- The key design issue is thermal and mechanical response to the fusion pulses of the first wall protection material (gas, liquid, or solid), and the subsequent loading of the structure. Although theoretical analyses will continue to be performed, full-scale experimental verification is essential because of the complex synergistic phenomena involved. We expect that some current chamber concepts will be eliminated through analysis and nonnuclear experiments before the ETF is completed. The four different wall concepts described below have different thermal/mechanical response objectives.
 - Dry wall. The amount of wall ablation or debris buildup and the transient structural loads must be determined, and the first-wall integrity must be verified. If the wall is also magnetically protected, the interaction of the target debris with the magnetic field must be determined.
 - Gas-protected wall. For a variety of target designs, gas compositions, and gas pressures, we must determine the gas deposition profiles for x rays and debris, the blast-wave loading of the structure, and the transient heat transfer through the gas.
 - Wetted wall. The magnitude and duration of the pressure pulse transmitted to the structure from the fusion-heated liquid must be verified, and the subsequent motion of the liquid wall and ablated gas must be tracked.
 - Liquid-metal wall. Jet disassembly characteristics, liquid-liquid interactions, ablated-gas behavior and gas/liquid loads on the structure must all be measured.
- The chamber atmosphere must be restored to acceptable levels of pressure and temperature before the next target is injected and the next driver pulse is propagated. Early in the interpulse time, the chamber atmosphere includes extremely hot material, e.g., debris ions, heated chamber gas, and vaporized liquid- or solid-wall material. The vacuum systems must cope with these materials or be capable of interrupted operation. The return to nominal first-wall conditions must be demonstrated.
- Heat removal must be verified. The structure is heated both directly and by thermal contact with its protective material (gas, liquid, or solid). The ETF must verify the scheme used to cool the structure.
- The pressure pulses produced in blanket liquids by pulsed neutrons could be measured by the ETF, as will the streaming of neutrons down beam lines and vacuum ports. Although facility lifetime precludes full-scale testing, the ETF will be capable of producing observable radiation damage in small

samples mounted near the source.

Between 10^{18} and 10^{21} n/cm² are required depending on the material and the effect. At a 10-cm distance from a 25-MW source, 10^{18} neutrons/cm² will be accumulated in ~ 140 s. About 10^5 s (10% of the facility life) would be required to produce 10^{21} n/cm² at a distance of 10 cm from a 25-MW source. These experiments, which will contrast pulsed neutron damage with steady neutron damage from facilities such as the Rotating Target Neutron Source (RTNS) or the Fusion Materials Irradiation Test facility (FMIT), will be essential for the design of commercial ICF power plants.

- The chemistry of contaminants (target debris, ablated wall materials, and corrosion products) in the chamber gas and coolant could be studied with the ETF to verify the capability of remotely processing these fluids. Because of the facility's low duty cycle, this area of interest should be explored between experiments.
- The safety and remote operation and maintenance of a facility containing radioactive materials and, possibly, large quantities of liquid metal will be a key demonstration goal of the ETF.

Sizing the ETF. The size of a magnetic-fusion ETF is determined by the required plasma parameters; e.g., the Tokamak ETF will be 600 to 750 MW_t. In ICF, however, the plasma is decoupled from the reactor, and the ETF power is based on duplicating reactor parameters such as first-wall fluence and blanket power density. Because some design issues depend on power, while others depend on fusion-pulse energy, it is appropriate to design the ETF to operate at the full reactor repetition rate. There is general agreement that a 25-MW_f facility, operating with a 1.0- to 1.5-MJ driver and gains in the range of a few to a few tens at 1 to 10 Hz, is all that is required. These operating conditions are compatible with conservative gain-curve estimates (Fig. 9-93). The shaded area is an extrapolation of the conservative gain curves, evaluated for driver energy above 1 MJ).¹¹⁶

First wall fluence is the key design parameter for the dry-wall, magnetically-protected wall, and wetted-wall reactor concepts; the wall radius required to achieve

full-scale first-wall fluences in a 25-MW_t ETF is between 0.5 and 1.0 m for these three concepts. For the gas-protected wall, the key parameter could be either first-wall fluence or chamber power density. For the 20-Hz SOLASE gas-protected-wall reactor concept,¹¹⁷ these reactor-scale parameters are duplicated at radii of 0.6 and 1.2 m, respectively, for a 25-MW_f ETF.

For the HYLIFE liquid-metal-wall reactor concept, the key design parameters are the neutron energy density in the inner portion of the jet array, and the soft-x-ray and debris fluence on the inner surface of the jet array. For a 100-MW_f ETF, full-scale reactor midplane conditions are duplicated when the inner radius of the jet array is 0.1 to 0.15 m. Thus, a 25-MW_f ETF would be about one-quarter scale for HYLIFE.

The engineering impact of a low-power ETF is considerable, and we do not have to utilize or wait for the development of high-gain targets (which can proceed in parallel). Low power minimizes the heat transfer and vacuum pumping requirements, as well as the costs. In addition, low power also means that the facility will be easier to operate and maintain, and that tritium fueling costs will be reasonable ($\sim \$20\,000/\text{yr}$ at $\$8/\text{mg}$).

Because all of the current reactor candidates require about the same size chamber, and because they all could use the same heat-removal and other supporting

Fig. 9-93. Fusion yield as a function of driver energy for a 25-MW_f facility.

Accomplishments Associated with the HYLIFE Chamber Concept

systems, sequential testing of multiple reactor concepts in the ETF may be feasible and cost-effective.

Cost of the ETF. The conservative target-gain curves indicate that a 1.0 to 1.5-MJ on-target short-wavelength laser or heavy-ion beam could drive a 25-MW_f ETF at reactor repetition rates. Because of the low thermal power and driver energy, unit costs for individual subsystems such as the driver (\$/MJ), liquid metal pumps (\$/MW), or target factory (\$/target) can be much higher than allowable unit costs in a commercial plant and still result in an ETF cost that is reasonable. A cost scoping study for an ICF-ETF has been initiated.

Characteristics of a HYLIFE ETF. The first-cut design parameters for an ICF-ETF using the HYLIFE reactor concept are listed in Table 9-27. The target gain is ~17 at 1.5 MJ; a combination which should be available in the early 1990s. Although the facility power is only 0.8% of a 1-GW_e plant, the ETF chamber is about 50% scale, and the lithium pumps are full-scale. The intermediate sodium loop should be very small, with no steam or electricity

generation. The buildup of tritium and target debris in the lithium would take hundreds of operating hours, and the buildup of helium in the chamber vessel would take several operating hours; hence, vacuum pumping and liquid-metal cleanup could be done off-line, between experiments. If 20 minutes are required to reach thermal equilibrium, about 1200 targets would be required per experiment.

Summary. We have enlisted the advice of a broad spectrum of fusion-engineering experts to begin the planning for an engineering test facility for inertial confinement fusion. We concentrated on setting the objectives for such a facility, listing the development requirements for each reactor concept, and defining the appropriate size and time scale.

A small (25-MW_f) device, operating at 1 to 10 Hz and with modest target gains in the range of 3 to 30, appears to be sufficient to attain engineering objectives comparable to those in the Tokamak program, which will require a 600- to 750-MW_f device. Therefore, a High Gain Test Facility (HGTF), designed to develop gains ≥50 need not precede the ETF. In addition, many ETF-type experiments could be performed with a single pulse driver such as a portion of the HGTF.

An ETF that does not require high gains or high power levels can provide potentially significant savings in both development time and cost and is, therefore, an important opportunity for ICF.

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Table 9-27. Characteristics of a HYLIFE ETF.

Yield = 25 MJ	Driver energy = 1.5 MJ
Gain = 17	Repetition rate = 1 Hz for up to 20 min
Chamber Dimensions	
2-m first-structural-wall inner radius; 3-m vessel outer radius	
Bottom of nozzle plate to exit pool = 6.4 m ($\frac{H}{R} = 1.6$)	
Total vessel height ~13 m	
Liquid-Metal Parameters	
500°C Li array inner radius ~10 cm	
0.5-m effective thickness and 50% packing fraction at midplane	
80% packing fraction at nozzle	
6.3 m/s inlet velocity and 10.1 m/s midplane velocity	
2.1 m injection head and 13.5 m total head	
Flow rate = 19 m ³ /sec in array; 30 m ³ /s max	
Clear time for array reestablishment = 0.67 sec (max repetition rate = 1.5 Hz)	
Total Li inventory ~61 Mg (10 × chamber inventory)	
Tritium buildup time to 1 wppm: 323 operating hours (30% burnup, 1.6 tritium breeding ratio)	
Helium buildup time to 0.1 Torr: 6 operating hours	
Pb buildup time to 1 a/o (23 w/o): 500 operating hours at 10 g/target (10 g is assumed to obtain an order-of-magnitude buildup time; no actual target design has been used)	

A New Liquid Metal Chamber of High Repetition-Rate Operation: the JADE Chamber

We have developed a new design concept¹¹⁸ for absorbing the pulsed energy release of inertial fusion directly in a liquid metal. This concept should work at repetition rates in the range of 4 to 10 Hz, using either lithium or lead-lithium over a wide range of temperatures, with either laser or heavy ion beam drivers, and with a wide variety of illumination geometries. The development of such a chamber would be extremely important in providing the flexibility needed to adapt to future developments in drivers and capsules. The success of the new concept rests on the feasibility of simultaneously:

- Increasing the residence time of the liquid metal in the chamber to reach a high repetition rate without prohibitive pumping power.
- Minimizing the time required to recondense ablated vapor.
- Reducing the forces resulting from induced fluid motion to the low levels required for long fatigue life of ordinary materials.

These issues have been addressed conceptually; but further analysis is required before their feasibility can be assessed.

The Design Strategy. The most general way to think of the function of an ICF reaction chamber is not only as an energy-conversion process,¹¹⁹ but also as an entropy-generation process. Chamber design is in some sense exactly opposite to target design. In igniting a fusion capsule, we seek to compress matter as nearly isentropically as possible. In designing the chamber, we seek to convert the fusion energy to heat with the largest possible entropy-generation rate. We must take a highly directed outflow of high-grade energy (energetic particles and photons) and degrade this energy to heat at 500°C in the brief time between pulses. If this energy is absorbed in a solid material, the material suffers rapid damage in the form of vaporization, displaced atoms, and

thermal fatigue.¹²⁰ If, instead, the radiated energy were absorbed in a fluid, avoiding permanent damage, momentum is created as the fluid expands to relieve the high energy densities and pressures. The resultant gas pressure and impact of liquid on the first wall could cause excessive transient stresses in the wall material, but through clever design, we now have control in the rate at which momentum is delivered to a wall. By changing the fluid and wall configurations we can reduce pressure amplitudes and extend interaction times beyond the structural response times, thereby transforming potential impulsive loadings into effectively steady loadings. The first key to successful LMW chamber design then is to use the fluid to greatly extend wall interaction times and thereby lower stress amplitudes.

The second key to successful LMW chamber design is to dissipate the fluid kinetic energy by damping out the fluid motion, generating heat as quickly as is consistent with first-wall stress limitations imposed by safety and lifetime parameters. This entropy generation is critical to the energy conversion process. Our new chamber design must both result in low wall stress and damp out fluid motion quickly, to safely transform the fusion pulses into a steady flow of heat.

The new design employs a large spherical chamber lined with a cellular or honeycomb structure made of layers of porous, wick-like metal fibers. Imagine a spherical shell a few metres thick and from 7 to 10 m in radius that is effectively an enormous metal sponge. The entire sponge-like structure would be heavily laden with liquid metal, which would migrate from top to bottom of the reactor by gravity and capillarity, picking up a temperature differential of 160°C. A structure like this might hold eight to ten times its volume of liquid metal in liquid metal. The total flow rate of liquid metal in this structure would be about 10 m³/s, approximately one-tenth that of HYLIFE.¹²¹ The surface of the structure that faces the target would be the open cells of a deep honeycomb, somewhat resembling the waveguide lens of Fig. 9-94, to provide a very large surface area that would significantly reduce reaction forces caused by ablation of the fluid. Cells would be several

centimetres across; the inside radius of the structure would be 6 to 8 m.

The inner portion would be constructed in a way that would provide an efficient absorber of x rays and debris,¹²² and the structure would also be an effective neutron absorber.¹²³ The entire blanket structure would be made of layers of a porous material (e.g., fibers of steel as shown in Fig. 9-95), formed into a honeycomb structure a metre or two deep. When heavily laden with liquid metal, the structure should be an excellent absorber of neutrons. The structure would use very little solid metal, would provide internal voids required for fluid expansion, and would also provide for rapid dissipation of fluid kinetic energy both by splashing and by drag forces.

We have the flexibility of using either lithium or an 83% lead-17% lithium eutectic

alloy as a protective working fluid.^{118,124} Lithium has better heat capacity and tritium-breeding capability, but the eutectic alloy requires only 40 to 50% of the thickness for the same neutron attenuation. The eutectic alloy also has better energy multiplication and does not burn in air, which may be a great safety advantage, but requires greater pumping power to perform the same function because its density is 20 times that of lithium. It will also have a significantly greater induced activation than lithium. Until either of these liquid metals is shown to be clearly superior in further studies, we consider both metals to be viable options. Finally, the honeycomb concept lends itself to a reactor-chamber system that is conceptually very simple, easy to assembly, and maintainable with minimum downtime.

Soft x rays and debris have a very short range in matter and deposit their energy very near the surface,^{125,126} producing energy densities high enough to ablate a portion of the surface, or even to spall off portions of the surface as the pressure is relieved, on each shot. A natural fix for this problem is to, first, reduce the flux of short-range radiation by increasing the surface area, and, second, coat the wall with a liquid that can be replenished on each shot to replace the vaporized material, thereby leaving the wall undamaged. Several wetted-wall designs have been suggested^{127,128} for ICF reactors, but there are four technical issues of concern in these designs.

- Can the porous structure replenish the wetted surface without dry spots between pulses?
- Will energy deposition by hard x rays or neutrons in the porous metal structure itself cause a kickoff or shedding of liquid?
- Will the impulse delivered by the fast ablation to the porous structure be tolerable?
- Can the vapor (ablated liquid) be self-pumped by condensation quickly enough between pulses to ensure that the vapor pressure in the chamber will be sufficiently low for beam propagation?

While we do not have definitive answers for any of these questions (which will probably require experimental resolution) we suggest some appealing design approaches. We also point out that our fiber metal structure differs from the usual transpiration-cooled surface (or wetted-wall)

Fig. 9-94. Cellular structure similar to absorber concept.

Fig. 9-95. A 0.5-in.-thick fibermetal layer of 12% solidity, courtesy of the Technetics Division of Brunswick Corp. ▼

by being far more porous. A fiber metal material of only 10% solidity will not feel the same effect as an 80% solid wall that oozes liquid.

Consider the deposition process. The x rays arrive at the surface first, deposit within a few micrometres of the surface, and cause a very fast ablation. The debris, also easily stopped in matter, arrives sufficiently later to be absorbed in the blowoff vapor. Because the soft x rays and debris are absorbed in such a thin layer, this absorption is unlikely to be interfered with by the fluid motion induced by the deposition of hard x rays or neutrons. The real concern relating to the deposition dynamics is whether the liquid metal can be held between the fibers of the structure after neutron deposition: expansion of the liquid due to isochoric heating by neutron deposition may cause the liquid to leave the pores at a rate greater than that by which it can be reestablished by capillarity in the time between fusion pulses.

The impulse delivered to the first wall by rapid ablation is reduced in two ways: by increasing the wall area and by the indirect coupling resulting from drag forces.

The surface area can be increased by more than a factor of 10 by using a large-radius chamber with a properly structured, involutioned surface. An example is the early suppressed-ablation reaction chamber¹²⁷ developed by LLNL.

Our most recent conceptual design for obtaining a large area absorbing surface is simply to form a cellular or honeycomb structure of porous fiber metal, as shown in Fig. 9-96. The surfaces of the honeycomb structure would be slanted at a small angle off the radial direction, and the entire wick-like structure would be heavily laden with liquid metal. The x rays and debris would penetrate deeply into the cells of the structure and would be absorbed over a large area of wetted wall. Because the entire structure would be perfused with liquid, each ablating surface would face a cold, liquid, condensing surface. Perhaps the vapor, in making contact with the wet walls in the many interactions required to leave the deep cells, would condense before exiting the cellular structure.

Experiments have shown that liquid metals can cling very effectively to metal screens; lithium does not drip even when the screens are horizontal.¹²⁹ Temperatures

would be kept uniform in such a wick-like structure by the combined effect of condensing vapor and capillary flow of liquid. In fact, the entire structure would act very much like a heat pipe.

Lithium heat pipes^{130,131} are very efficient, operate up to 1600°C, and have been used to carry axial fluences of 10 to 20 kW/cm². In comparison, the average heat flow that the fusion-reactor structure must carry is less than 1 kW/cm². We know of no other application of heat pipes in which the pipes are effectively open to a pulsed energy source and where the vapor is produced by ablation. Nevertheless, the condensation dynamics and the liquid flow by capillarity in the fusion-reactor arrangement should be very similar to those in a liquid metal heat pipe.

Porous structures made of sintered steel fibers have been in commercial use for many years. Formed into a rectangular honeycomb configuration, these structures should be very strong. The fibers used in the fiber metal structure should have both good corrosion resistance and low activation. Pure iron or titanium are excellent candidates.

This type of structure is extremely effective at taking the sharp impulse due to ablation and transforming it into very low forces

Note: Cells are canted slightly to the direction of X rays and debris.

Fig. 9-96. Honeycomb structure with canted cells.

applied over very long times. This is done both through increasing the absorbing area and through the use of the drag force. We can estimate these forces by first imagining the reference case of a simple sphere, coated on the inside with a thin layer of lithium.

The maximum pressure^{132,133} that can be exerted on this surface, if all the reaction to the ablation is transmitted through the layer to the wall, is

$$p_{\max} = \frac{fY}{4\pi R^2 \tau \sqrt{2L_v}} \quad (99)$$

where p = pressure, Y = yield, f = fraction of yield in x rays and debris, R = chamber radius, L_v = heat of vaporization per unit mass, and τ = duration of x-ray pulse. ($L_v = 1.96 \times 10^7$ J/kg for lithium, $L_v = 8.57 \times 10^5$ J/kg for lead.) Using a 400-MJ yield with 30% in x rays and debris, incident in a 1- μ s pulse on an 8-m chamber coated with a thin layer of lithium, we obtain a maximum pressure of 240 atm for the microsecond. However, if we have the honeycomb or open-celled configuration, the pressure is far less. Imagine cells 3 cm², extending for 1 m between $R = 8$ and 9 m. The cells are canted with respect to the radius so that the entire inner side of one cell is illuminated; this works out to an angle of about 2° (0.03 rad), and an area enhancement of a factor of 33. The pressure is reduced during the pulse to about 7 atm.

This is an estimate of the force on a wetted solid surface; however, the force on a 10% solid fiber metal structure is considerably lower. Please refer to Fig. 9-97. The impulse from the ablation will drive a wave of

Fig. 9-97. Ablation force drives wave through liquid metal, exerting drag on fiber-metal structure.

velocity v and length τ into the fluid, exerting drag on the mass of fibers. The force on the structure is much lower than the original ablation force, and the duration time far longer. We can estimate the coupling coefficient as follows:

$$\begin{aligned} \text{C.C.} &= \frac{1}{2} \frac{C_D \rho v^2 A_{\text{fibers}}}{p_{\max} A_{\text{surface}}} \\ &\times \left(\text{fraction of fibers perpendicular to flow} \right) \\ &\times \left(\text{fraction of fibers in wave} \right) \end{aligned} \quad (100)$$

In a layer of fiber metal of thickness Δ and solidity s , made up of fibers of diameter d , the ratio of side area of fibers to surface area of the layer is

$$\frac{A_{\text{fibers}}}{A_{\text{surface}}} = \frac{4s}{\pi} \left(\frac{\Delta}{d} \right) \quad (101)$$

As the wave travels through the layer, only the fraction $(v\tau/\Delta)$ of the volume experiences drag, and only 2/3 of the fiber side area is perpendicular to the flow. The flow velocity of the wave is related to the ablation pressure by $\rho v^2 = p$. The coupling coefficient is then

$$\text{C.C.} = \left(\frac{4}{3\pi} \right) s C_D \left(\frac{v\tau}{d} \right) \quad (102)$$

Assuming a solidity of $s = 0.1$, a drag coefficient of $C_D = 1.0$ appropriate for the Reynolds number range of 10^3 , and a velocity of 37 m/s corresponding to the pressure of 7 atm on lithium which was calculated above, we obtain a coupling coefficient of $\text{C.C.} = 0.01$. Only 1% of the applied reaction force is transmitted to the structure by drag, and the force is applied over a much longer time scale, as the wave travels back and forth in the fiber metal layer filled with 90% lithium. This is exactly the effect we sought. All of this fluid kinetic energy is eventually lost by viscous dissipation due to drag, in a time scale long compared to the original pulse duration, but short compared to the time between pulses.

There is some concern that it would be difficult to keep the surfaces facing the pellet

explosion sufficiently wetted. One option is to inject a few jets of liquid metal into the chamber, which would then be blown into cells on each shot. If capillarity forces are sufficient, this should not be needed.

Absorption of Neutrons. The neutrons are absorbed in the liquid metal in a time short compared to the hydrodynamic response time, and the increased energy density manifests itself as a substantial rise in pressure, but with only a very small temperature rise. Since there are free surfaces, as in any LMW concept, expansion waves will move inward from the surface, converting the pressure to kinetic energy of the expanding fluid. In the HYLIFE design, we arranged the liquid metal jets in a pattern that had a packing fraction of 50%, thereby allowing room for pressure relief. Glenn¹²³ has found that only a small relief volume is required, just enough to provide sufficient room so that two adjacent layers will not collide before an expansion wave has had time to traverse the entire layer and relieve the pressure. We also learned in the same study that splitting the neutron blanket into many layers promoted a great deal of collision between two-phase expanding flows. There is significant dissipation of kinetic energy involved in the splashing of two-phase flows, just as there is in the compression of porous solid media¹³⁴ (e.g., such energy-absorbing processes as the crushing of styrofoam or honeycomb in impact). Therefore, our strategy in constructing a neutron blanket of liquid-impregnated porous material includes provisions for having layers of neutron absorber alternating with open volumes to accommodate both pressure relief and splashing of the accelerated liquid. A honeycomb construction not only has the open cells that are required but also has the strength needed for supporting the porous metal (see Fig. 9-98).

For a specific example, imagine a structure in which the layers of porous metal are 1 cm thick, have a solidity of 20%, and are formed into cells 3 cm across. When perfused with liquid metal, this structure would be approximately 44% void, 44% liquid metal, and 12% solid metal (steel). We have calculated the breeding ratio for such a structure, using lithium and 83% lead-17% lithium. Preliminary calculations using a one-metre thickness of such a structure

indicate that a tritium breeding ratio of 1.26 will result with natural lithium and a porous steel fiber metal of 20% solidity. Using lead-lithium, the tritium breeding ratio varies between 1.08 and 0.91 as the solidity of fiber metal is increased from 10 to 20%.

As a bonus, this fiber-metal structure introduces an additional energy-dissipation mechanism: flow through a porous medium. The drag encountered by the neutron-heated fluid as it expands through the fibrous medium, and again as it splashes back into the structure, is certain to contribute to the rate of reduction of kinetic energy between pulses. Note also that no liquid would be lost during this process; it would simply splash around in the honeycombed or layered structure of porous plates and end up back in place for the next shot. Arterial channels could be provided to supply the porous structure with sufficient fluid or to serve as a drain.

An extremely important feature of a neutron blanket of this type is that the impact forces would be self-cancelling, so that there should not be significant net impact forces transmitted to the surrounding support structure. To see this, consider the dynamics of the response to neutron deposition of a honeycomb-core structure, as shown in Fig. 9-98, and focus on the

Fig. 9-98. Illustration of absorber wall performance.

behavior of the liquid from two adjacent layers. When the neutrons are deposited, the liquid is heated volumetrically to a high pressure in each layer. The liquid expands perpendicularly to the surface, heading toward a collision with fluid similarly expanding from the facing surface. A velocity of about 20 cm/s typically results when the neutron energy from a 400-MJ pulse deposits in liquid lead at a distance of 8 m from the pulse source. When the two expanding liquid flows meet at the center, liquid to liquid splashing occurs, dissipating some of the kinetic energy. The liquid flows then return to the porous layers from whence they came. Because this process happens between each layer, every porous-metal layer is impacted symmetrically and simultaneously from both sides. Therefore, there will be little net force acting on the structure as a whole. Of course, there will be drag forces on the metal fibers making up the fiber-metal layers as the fluid moves through. We have yet to determine either the fatigue life of the fibers or the dissipation rate of the fluid energy resulting from this drag.

There is concern that the fibers would be damaged by neutrons. Conventional fusion neutron damage limits for pressure vessels have no meaning for this application however; they are far too conservative. The fibers can swell and distort without losing their function. Embrittlement may be a limitation though.

Chamber Construction. The blanket of the spherical chamber would be constructed of individual building blocks in the shape of longitudinal sections of a sphere; these sections would fit together like slices of an orange or a melon. Each section would have

enough strength to resist the reaction forces from each shot, but a strong outside supporting structure must be provided to support the chamber against gravity, since the weight of the structure would be considerable when loaded with liquid metal. The longitudinal sections would be attached to the outer structural strongback by interleaved ribs.

In Fig. 9-99, we see more clearly how the longitudinal sections would be stacked together to form a sphere. The primary structural members of the strongback framework would be rings that attach to the ribs, forming the latitudinal support. Holes large enough to admit the beams of photons or heavy ions would be cut in a few sections of the sphere. Because the entire supporting structure would be external (an exoskeleton), it would be effectively protected against neutron damage by 50 to 100 cm of liquid metal and would last the 30-year life of the power plant without replacement. The entire spherical structure could be assembled quite easily by building it in a jig. Because there would be no vacuum connections or close tolerances, either a bolted or welded construction could be used (whatever is most appropriate to the remote disassembly and maintenance techniques of the day).

Reactor Configuration. The reaction chamber must be encased by shielding and must be furnished with a supply of liquid metal. The simplest way to do both is to use a containment pit similar to a pool-type fission reactor, as shown in Fig. 9-100. Imagine a circular concrete pool with a hemispherical bottom and a large concrete cover, all lined with steel. A dense spray of liquid metal would infiltrate the porous structure, seep through the structure to a bottom sump, and then be pumped first through the intermediate heat exchanger and then to the top of the chamber. The chamber could be disassembled piece by piece and removed from this pool by an overhead crane. Because the curvature of the pool bottom would match that of the chamber, no additional support should be necessary.

The liquid metal flow loop is a hybrid loop/pool design. Flow from the cold side of the intermediate heat exchanger would be introduced at the top as a very heavy spray that would be designed to heavily saturate the wick-like structure and to run down through arterial channels so that the liquid

Fig. 9-99. The fiber-metal absorbing structure.

would be uniformly distributed by gravity throughout the spherical shell. Therefore, the large-scale flow would be by gravity fed through channels, and the small-scale flow would be by capillarity.

Note that the entire pool volume above the chamber, in which there would be a spray of the cooled liquid, would be available to condense vapor that originates in the chamber, because the entire chamber structure would be open to vapor flow. In fact, because the average temperature of the droplets would be 100°C lower than that of the liquid in the blanket, the spray would also act as an effective condensing vacuum pump.

The cold-leg temperature is constrained by the melting temperature of the lead-lithium eutectic alloy (340°C), and the hot-leg temperature is constrained by the strength and the temperature-sensitive corrosion resistance of ferritic steels (500 to 550°C). The liquid would migrate through the structure and pour out of the bottom into a sump, from which it would be pumped by either mechanical or electromagnetic pumps to the intermediate heat exchanger.

In the HYLIFE design, the total liquid metal flow rate is 10 times greater than that

required for the transfer of heat, since a great excess of liquid metal is required simply to set up the thick cascade. In the fiber-metal design, on the other hand, only that flow required for heat transfer is required. The pumping power should not exceed 20 MW_e required for HYLIFE even if the lead-lithium alloy is used.

A schematic plan view of the reactor is shown in Fig. 9-101, where we depict two-sided laser-beam irradiation. Of course, there is nothing in this concept that precludes the use of heavy-ion beams or of illumination geometries that are closer to symmetrical illumination, such as the double high-cone-angle scheme of the Nova laser. In fact, although we do not show it here, we also envision the use of this reactor concept with neutralized light-ion beams. The most natural geometry in this latter case would be to array the diodes around the equator and to have porous liquid-metal-filled structures form upper and lower hemispheres. The central pool that would hold the reaction chamber would be surrounded by the liquid metal pumps, the intermediate heat exchangers, and the vacuum pumps needed to pump such noncondensable gases as helium and unburned D-T.

Fig. 9-100. The reaction chamber in a pool-type containment pit.

The JADE Chamber

Comparison of the schematic reactor shown in Fig. 9-101 to a Liquid Metal Fast Breeder Reactor (LMFBR) such as SUPERPHENIX, and to the HYLIFE plant costed by Bechtel¹³⁵ gives indications that such a reactor, exclusive of the driver, could be constructed at a cost equal to or less than that of an LMFBR of comparable thermal size. Of course, any real cost-estimating of this concept must await a credible engineering design.

Conclusions. This conceptual idea for an ICF reaction chamber is suggested as a way to realize the advantages of an LMW system (e.g., protection of the first-wall structure, energy absorption and heat transfer in a single fluid, minimized wall replacement and induced activity) in a high-repetition rate system without incurring the penalties of high pumping power. Additionally, this concept should offer the flexibility of operating with either a laser or heavy-ion-beam driver, with either lead-lithium or lithium fluid, and with a variety of irradiation geometries.

Notwithstanding these desirable features, there are many serious issues that must be investigated to assess the feasibility of this concept, including:

- What is the lifetime of these fiber-metal structures, immersed in liquid metal and subjected to transient forces, in the fusion environment?
- What is the largest feasible size of honeycomb-type cellular structures that can be fabricated of fiber-metal material? Which low-activity materials can be used?
- What is the appropriate neutron-damage limitation that should be applied to such a fibrous structure, as compared with a pressure vessel?
- What is the corrosion life of a structure having this much surface area?
- What is the tritium breeding ratio for a structure with this configuration?
- What are the forces applied, first to the fibers and then to the structure, when the reaction force from the fast ablation caused by x rays and debris is applied to the surface of a liquid metal filled structure of 20% solidity?
- Can engineering designs be developed to ensure that such large wick-like structures will be uniformly wetted on all the layers and surfaces?
- What would happen if a dry spot should occur? Would it lead to only a small portion of the structure harmlessly

Fig. 9-101. Schematic plan of the reactor shown for two-sided laser-beam irradiation.

burning away, or would it rapidly lead to the destruction of the structure?

- Can an engineering design be found that can assure both surface rewetting in the time between pulses and a flow rate through the structure that will provide an average temperature rise that can be matched to the steam cycle?

If these questions can be resolved by analysis, experiment, and ingenuity, we will have a fusion reactor of great functional simplicity. This reaction chamber would be nothing but a large, wet, metal sponge supported on the outside by a simple framework that can be sufficiently protected from neutron damage by the liquid metal. There would be no critical tolerances, no welded coolant tubes that must be leakproof, no high voltage, no superconducting coils—in fact, no cryogenic liquids, electricity, or water within the primary shielding of the containment structure. From this standpoint, the present chamber concept shows great promise for the development of a practical fusion power plant. Therefore, we must begin to address the critical issues associated with this idea and to seek practical engineering designs.

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Ultra-High-Pressure Shock-Wave Experiments

Accurate equation-of-state (EOS) information is of major importance to several laboratory programs, and is of general scientific interest as well. Conventional shock-wave experiments using impactors provide usable data at pressures only up to about 0.5 TPa (5 Mbar). Above roughly 10 to 100 TPa, theoretical models are typically used to determine material behavior. In the range 0.5 to 10 TPa, current theoretical models cannot be used to accurately determine material properties. Experimental data in this pressure range is sparse, since the experiments have required nuclear explosions to generate high pressures, up to about 6 TPa. The goal of our program using laser-generated shockwaves is to measure, in a laboratory environment, the equations-of-state of

various materials in this ultra-high-pressure range. The use of lasers offers the advantage of experimental flexibility, and many experiments can be done in a short time, and for a relatively modest cost.

We have achieved pressures up to 2 TPa in aluminum using the Janus laser. During the past year, we have used Shiva to generate a pressure of 3 TPa in a gold sample. That experiment is described in detail in “Ultra-High-Pressure Laser-Driven Shocks in Gold” in Section 7 of this annual report. In those experiments, only shock velocity was measured, which alone is insufficient to determine the EOS. We are currently pursuing two approaches to measuring the EOS of materials on the Hugoniot curve (the locus of final thermodynamic states which are reached when shock waves of varying strength pass through a material). The first approach will employ a technique for measuring the mass velocity of material behind the shock front. This is still in the design stage; we will report it later.

The second method for finding the EOS is to use the impedance-matching technique, which has proved very useful in impact-generated shockwave experiments. This method involves the passage of a shock wave between two materials, one of which is usually a standard material with known EOS, the other with an unknown EOS. When the shock velocities in both materials are measured, a point on the Hugoniot curve for the unknown can be determined. In the pressure regime of our experiments (roughly 0.3 to 3 TPa), there are no materials with a sufficiently well-known EOS to serve as a standard material. However, a large number of accurate impedance-matching experiments using two materials such as aluminum and copper may sufficiently constrain the calculated EOS for these materials to qualify one or the other as a “standard” material. Aluminum and copper have been chosen for our experiments because they are both particularly amenable to calculation, and they can be easily fabricated into targets.

We have performed a preliminary series of impedance-matching experiments using the Janus laser, with the goal of assessing target designs and to identify those problems which limit our accuracy. Figure 9-102 shows a typical target design used in these experiments. The targets were fabricated at Los

Alamos National Laboratory using vapor deposition. Mechanical knife-edge masks were used to form the steps. We use an ultra-fast streaking camera to record the arrival of the luminous shock front at the two surface levels.¹³⁶ The average shock velocity in the steps is found by dividing the step height (measured with a sensitive stylus gauge) by the time difference of shock arrival at the base and top of the step.

Measured shock velocities in aluminum and copper steps are plotted in Fig. 9-103 for a number of experiments. The data agree (within about 10%) with calculations using extrapolations of the Hugoniot data of Mitchell and Nellis,¹³⁷ approximating the reflected Hugoniot of aluminum by the mirror reflection of the aluminum principal Hugoniot. These experiments achieved

inferred shock pressures of 0.2 to 0.6 TPa in aluminum, and 0.4 to 1 TPa in copper, or about three times higher than have been obtained for these materials in impactor experiments using a two-stage gas gun. The pressure in the aluminum steps follows the approximate scaling law $P = 0.5 I^{0.6}$, where P is pressure in TPa, and the incident laser intensity, I , is expressed in units of 10^{14} W/cm^2 . In the work shown here, the laser intensity varied between $5 \times 10^{13} \text{ W/cm}^2$ and $5 \times 10^{14} \text{ W/cm}^2$.

While the data are clearly not accurate enough to serve as a starting point for theoretical calculations, the apparently random scatter of the points is somewhat reassuring in that it indicates that random errors rather than systematic ones were dominant. We have identified target deformation from flatness, nonuniformity of the incident laser beam, shock decay, and streak camera problems (nonlinearity and poor time and space resolution) as mainly responsible for the errors. All of these problems appear solvable. Using flatter targets and more uniform laser irradiation focused in a narrower convergence angle (these experiments used $f/1$ illumination) will improve the results. Recent improvements in streak camera technology will make shock velocity measurements more accurate. Shock decay is a problem when the shock transit time through the target is significantly longer than the laser pulse (300 ps FWHM in this work). Using a longer laser pulse or pulse shaping may help, and multiple steps of each material on the rear surface will allow shock decay to be characterized for each shot. Furthermore, LASNEX calculations^{138,139} indicate that inclusion of a 1- to 2- μm -thick gold layer imbedded in the target within about 2 μm of the front (laser-irradiated) surface will reduce the effects of both shock decay and hot-electron preheat of the target. We are now testing this design in our experiments at Janus.

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Fig. 9-102. A typical aluminum-copper impedance-matching target.

Fig. 9-103. Aluminum-copper impedance match data.

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Operating Parameters of Inertial Fusion
9-4

The Influence of Chamber Pulse Rate on Structural Stress and Lithium Flow
9-9

Blanket-Shield Design Options
9-15

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The Effect of Chamber Ports on Neutron Leakage and Blanket Performance 9-21

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Solubility of Hydrogen Isotopes in Lead-Lithium Solutions 9-25

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Corrosion of 2-1/4 Cr-1 Mo in Lithium and Lead-Lithium Solutions 9-26

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Heavy-Ion-Beam Ballistic Propagation in Liquid-Metal-Wall Reaction Chambers 9-43

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Nozzle Plate 9-71

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Section 10 Advanced Isotope Separation

Advanced Isotope Separation

Contents

Introduction	10-1
AIS Facilities	10-5
Current Facilities	10-5
Planned Future Facilities	10-9
Technical and Systems Highlights	10-13
Laser Research and Development	10-17
Introduction	10-17
Copper-Vapor-Laser Introduction	10-18
Copper-Vapor-Laser Performance	10-18
Copper-Vapor-Laser Modeling	10-21
Cross Sections for Electron-Impact Excitation of Copper	10-26
Fire Suppression in Laser Dye Media	10-27
XeCl Excimer Laser Injection-Locking and Amplifier-Energy Extraction	10-28
Atomic-Vapor Physics	10-30
Introduction	10-30
Associative Ionization in Alkaline Earth Elements	10-30
Collisions Between Uranium Atoms and High-Energy Electrons	10-31
Electronic Excitation and Ionization in Intermediate-Energy U-U Collisions	10-33
Laser-Induced Fluorescence Studies of U II Produced by Photoionization of Uranium	10-36
Spectroscopic Data Acquisition	10-42
Beam-Propagation Theory	10-43
Near-Resonant Propagation in Sodium	10-44
Computational Facilities	10-47
Introduction	10-47
Laboratory Computer Systems	10-47
Alternative Applications of Technology	10-51
Introduction	10-51
Laser Cleanup of Pt-Group Metals	10-52
Vulcan Very-High-Temperature Gas (VHTG) Heater	10-54
Planning Methodology	10-58
Introduction	10-58
AVLIS Design Criteria	10-58
Optimal Resource Allocation and Decision Analysis	10-60
References	10-68

Advanced Isotope Separation

Introduction

The purpose of the DOE Laser Isotope Separation (LIS) Program is to demonstrate a uranium-isotope enrichment process that is more efficient and less expensive than present technologies for producing light-water-reactor (LWR) fuel. In the LIS Program at LLNL we are working toward this goal by developing an Atomic-Vapor Laser Isotope Separation (AVLIS) process, taking advantage of shifts in spectral-line frequencies between different isotopes (typically 5 to 10 GHz for atomic uranium). Because these isotope shifts exceed the bandwidths attainable with tunable laser sources, we can preferentially excite and then ionize a desired isotope (^{235}U) using powerful lasers. We then use electromagnetic forces to separate the ionized atoms.

Figure 10-1 shows the essential steps of our AVLIS process:

- An electron-beam (e-beam) heats a crucible of molten uranium metal to produce a sheet-like stream of atomic uranium vapor.

Fig. 10-1

- The atomic vapor flows into a separation zone where it encounters pulsed laser beams. These lasers, tuned precisely to visible transitions of ^{235}U , selectively photoionize the ^{235}U .
- Electromagnetic fields draw the ions onto charged product-collector plates, leaving the bulk of the vapor to accumulate on a tails collector.

The ability of the AVLIS process to obtain a large separation factor (and hence enrichment rate) in a single stage gives AVLIS cost-effectiveness advantages over the conventional gaseous-diffusion process or gas-centrifuge processes. Figure 10-2 presents our estimated comparison of capital costs,

Fig. 10-2

	Gaseous diffusion	Gas centrifuge	AVLIS
Capital	High (~\$600/SWU/year)	High (~\$800/SWU/year)	Low (~\$60/SWU/year)
Energy	High (~2400 kWh/SWU)	Low (~100 kWh/SWU)	Low (~100 kWh/SWU)
Operating	Low	Moderate	Low

operating energy costs, and other operating costs for these three processes. We express production in conventional units of separative work units (SWUs), a value function related to the free energy of “unmixing” required to separate a feed stream of specified assay into product (enriched) and tails (depleted)

Introduction

streams, also of specified assays. In the case of LWR fuel, as here, these assays are 0.7% ^{235}U feed, 3.2% product, and 0.25% tails. A 1-GW LWR consumes about 10^5 SWU/year.

The introduction of AVLIS technology into the LWR fuel cycle offers the following benefits:

- **High Single-Stage Enrichment.** Enrichment plants based on our AVLIS process can produce uranium fuel for LWRs using as feed either natural uranium (0.7% ^{235}U assay) or the depleted tails from conventional gaseous-diffusion plants (0.25% ^{235}U). Alternatively, the AVLIS process can enrich the tails from conventional facilities so they become feed for a conventional or AVLIS enrichment plant. Thus AVLIS offers a means of significantly increasing the effective uranium supply or, equivalently, reducing the effort and investment necessary to mine and mill uranium.
- **Low Operating Costs.** The low energy requirements for operating the AVLIS process suggest using it to replace existing gaseous-diffusion capacity as a means of reducing domestic energy demand.
- **Low Capital Costs.** The low capital cost of AVLIS enrichment makes it a logical candidate for meeting future needs for new enrichment capacity, whether it is the capacity to meet present worldwide commitments to nuclear power plants or for future growth. The substantial reduction in capital investment extends from enrichment plants to electrical-generating capacity and ultimately to mining and milling.

Our economic analyses, based on major technology advances in the last five years, show that AVLIS offers the unique advantage of both low capital and low power costs. On this basis any new uranium-enrichment capacity should use AVLIS after demonstration in a pilot facility. More significantly, there is a clear economic benefit to replacing existing gaseous-diffusion capacity with AVLIS facilities, even in the absence of a growth in demand.

The current level of AVLIS technology points to commercial deployment in the early 1990s. The interim development program for meeting this date must be an aggressive one involving parallel systems developed in the 1980s, as delays that defer commercialization until the mid-1990s or later will erode the benefits of introducing AVLIS. With this in mind, we have structured the AVLIS program at LLNL into four progressive stages, as indicated in Fig. 10-3:

Fig. 10-3

- **Technology Evaluation.** In this phase we used research equipment to examine the basic physical processes involved in AVLIS and to obtain the basic data needed for adequately defining the AVLIS process. Using these data we selected our mainline process. We investigated process-parameter scaling to determine the best approaches for developing system components.
- **Technology Demonstration.** During this stage we address issues of engineering scale, collect system-performance data, optimize process performance, and begin to resolve the problems associated with system lifetime.
- **Engineering Demonstration.** This phase calls for us to increase the component and subsystem lifetimes, reduce hardware costs, improve component reliability and maintainability, and optimize the system performance under plant operating conditions.
- **Economic Demonstration.** This final phase marks the actual use of AVLIS technology in producing cost-effective separative work in a production plant. We are now nearing the end of the technology-evaluation phase.

Spectroscopy experiments at LLNL early in 1974 showed the scientific feasibility of selectively isolating desired isotopes of uranium through photoionization. Since 1974 we have concentrated on developing laser, vaporizer, and extractor systems that can be scaled to large size. We have measured the relevant scientific parameters for a completely integrated process and have demonstrated uranium enrichment in a number of process experiments.

Our program is now moving from the technology-evaluation to the technology-demonstration phase at engineering scale. We have begun increasing component sizes and addressing the full-scale system-integration issues on which a full-technology demonstration depends.

To the extent that program schedules and funding permit, we will continue to explore parallel approaches in the areas of lasers, vaporizers, extractors, and collectors. While these options give the program flexibility and underwrite the performance and cost goals of future enrichment systems, it should be emphasized that the system components for the mainline AVLIS process already exist.

As Fig. 10-1 indicates, our AVLIS process integrates two major systems: a laser system, with associated optical subsystems; and a separator, incorporating vaporizer and extractor subsystems enclosed in a common vacuum tank.

Figure 10-4 shows the present integrated process-development facilities at LLNL. The SPP II system, first operated in 1976, provides primary laser power and waveform generation. In 1978 we added the Regulis separator. During 1979-80 we combined these systems for enrichment runs. In 1980 we integrated the Venus laser power-amplification system into our process-simulation experiments, thereby achieving a substantial increase in the illumination volume of vapor. The Mars higher-throughput vaporizer and process chamber, started up in 1980, will provide a factor-of-10 increase in uranium-vapor production.

In Fig. 10-4 we schematically indicate the relationship of the major components that form our integrated experimental-process test bed, representing two generations of both laser and process-chamber hardware.

In the remainder of this section we summarize in "Current Facilities Inventory" the various components of our present program: the laboratories, simulators, and test units that facilitate our R&D effort. This inventory provides a detailed overview of our program by defining the main constituents. "Planned Future Facilities" describes the additional buildings planned to handle future program growth.

Subsequently, we discuss some of the important "Technical and Systems Highlights" of 1980: the construction of the Mars separator facility; advances in the Venus laser system; and development of a large-bore CVL, as we moved along our planned development path from technology evaluation into technology demonstration.

We next present selected details of our AVLIS program. The articles emphasize supporting research, rather than component-development studies.

We have divided the technical articles into five categories:

"Laser Research and Development" articles deal with our R&D in CVLs, the dye lasers, and the RGH excimer lasers.

Fig. 10-4

Next, under "Atomic-Vapor Physics" we discuss some of the many processes that occur in the atomic vapor. We discuss measurements of the dimerization energies in "Associative Ionization in Alkaline Earth Elements," calculations of collision cross sections in "Collisions Between Uranium Atoms and High-Energy Electrons" and in "Electronic Excitation and Ionization in Intermediate-Energy U-U Collisions," the use of laser-induced fluorescence diagnostics in "Laser-Induced Fluorescence Studies of U II Produced by Photoionization of Uranium," the processing of spectroscopic data in "Spectroscopic Data Acquisition," theory and experiments on photon-atom interactions in "Beam-Propagation Theory" and laser-light propagation in "Near-Resonant Propagation in Sodium."

Under "Laboratory Computer Systems" we describe our computation system. Then we discuss two examples of technology developed in our AVLIS process that have alternative applications: laser isotope separation of platinum metals in "Laser Cleanup of Pt-Group Metals" and a high-temperature gas heater in "Vulcan Very-High-Temperature Gas (VHTG) Heater."

The final articles are grouped under "Planning Methodology," where we discuss the development of process-design criteria ("AVLIS Design Criteria") and computer-aided decision analysis ("Optimal Resource Allocation and Decision Analysis").

More detailed descriptions of all aspects of AVLIS at LLNL appear in classified publications.

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AIS Facilities

Current Facilities

Introduction. Over the last few years, as we progressed from technology evaluation to process selection, we have continuously and significantly altered the LIS facilities at LLNL. We identified problem areas and established laboratories and facilities to address critical issues. As we completed these investigations, we redirected our efforts and facilities to newer, more pressing issues. Thus the history of our physical facilities reflects significant changes from year to year in the use and the titles of these facilities. The following pages summarize our current facilities.

Copper-Vapor-Laser (CVL) Facilities. We devote six laboratories to studies related to CVL development.

In the **CVL Research and Development Laboratory** we perform experiments to understand the physics and bore-diameter scaling issues of high-power CVLs. We equipped the laboratory with the diagnostics necessary for monitoring important CVL operating parameters, including temporal pulse shapes, current and voltage waveforms, and relative power in the two CVL-output wavelengths. We operate two laser stations in this laboratory to allow experiments over a range of bore diameters and to permit oscillator-amplifier extraction studies. Early in 1981 we will install a sophisticated copper-level population diagnostic based on the Rozhdestvenski "hook" method. This system will provide for both temporal and spatial measurement of the densities of a number of excited states in a CVL discharge. We will use the data obtained, together with the results of a careful parameterization of the large-bore copper laser, to improve the computer model of the CVL now under development.

In the **CVL Engineering Development Laboratory** we are conducting advanced engineering development of large-bore cylindrical-discharge CVLs. Following the demonstration of the feasibility of the large-bore laser, we began experimental work to develop a reliable low-cost version. The laboratory is equipped with a 15-W CVL for measuring the amplifier performance of the high-power lasers located in an adjacent station. Here during 1980 we conducted the first successful life test of a high-power CVL. To assist in further development work, we have nearly completed a minicomputer-based data-acquisition system.

In the **CVL Performance and Reliability Laboratory** we develop the engineering advances required for reliable high-performance CVL operation. Our experimental work has involved the development of advanced small-bore CVL assemblies as well as the characterization of some large-bore CVL construction concepts. Our results provided modifications to CVL construction that significantly improved both the useful lifetime and beam quality of the small-bore CVLs operating in SPP-II and Venus.

The **CVL Kinetics Laboratory** contains a high-temperature copper oven, a single-mode cw dye laser, and a high-energy pulsed dye laser. In this laboratory we determine the collisional lifetimes of the $^2P_{3/2,1/2}$ and $^2D_{5/2,3/2}$ states of atomic copper in various gaseous mixtures of interest to the CVL system. We use a frequency-doubled Nd:YAG-pumped dye laser to selectively excite copper atoms. We employ a cw dye laser to monitor the excitation and relaxation of the excited states.

We established the **Advanced CVL Electronics Laboratory** in 1980 to test and evaluate new designs of pulse-power circuits required by high-power CVLs. We also use the laboratory as a life-test facility for the

thyatrons used in these circuits. In its role as a test bed for new designs, considerable changes in apparatus occur from month to month. At present we are testing a laser pulser constructed entirely of solid-state components using magnetic-switch technology. The pulser when completed may eliminate entirely the need for thyatrons for CVL switching. We expect to continue experimental work on this concept throughout 1981.

Late in Fiscal Year 1978 we established a **CVL Tube-Assembly Facility**, comprising a building and group of people dedicated to CVL tube manufacturing. We set aside a 140-m² room to house the facility. We divided this room into two areas: one area for machine-tool operations and a cleaner area for material cleaning and assembly. We designed and built several permanent tooling fixtures to facilitate tube-assembly manufacturing. These include devices for: inserting wicks; press-fitting electrodes; measuring alumina-tube diameters, eccentricity, and bow; pulling bellows without damage; and tack-welding heat shields and wicks. The technicians working in the facility are trained in all phases of tube manufacturing.

Dye Laser Facilities. We have three laboratories devoted to various aspects of dye lasers.

The dye laser is the principal means of frequency conversion from a laser pump (CVL or XeCl) to the wavelengths and line widths required for atomic isotope separation. We can tune and match the laser to the atomic absorption line. The laser is also inexpensive to operate and it is scalable. The active medium is a solution of an organic fluorescent dye. Using appropriate chemical synthesis we can engineer the dye's molecular structure to specific requirements. Thus we can build into the dye molecule such properties as absorption and lasing wavelengths, Stokes shift, tuning range, photochemical stability, solvent compatibility, and conversion efficiency. Chemical synthesis provides the most certain approach for progressing from our present effective dye systems to the needs of an AVLIS plant. We use two laboratories in the program's dye synthesis and evaluation effort. In the **Dye Evaluation and Test Laboratory** we characterize the optical properties of the dye.

In the **Dye Synthesis Laboratory** we synthesize the organic molecules used as laser dyes. The laboratory contains apparatus for chromatography, hydrogenation, crystallization, rotary evaporation, and for most routine organic chemical operations. Diagnostic equipment includes a Varian EM-360 NMR spectrometer, a grating infrared spectrophotometer, and other facilities of LLNL's Chemistry Department. A number of organic chemical intermediates come within the "toxic irritant" or "cancer-suspect agent" categories. Consequently we have built a ready-access negative-pressure cabinet in which to store them.

We established the **Dye-Laser Amplifier Development Laboratory** in 1977 to test and evaluate the performance of dye amplifiers of high average power. In this laboratory we determine and demonstrate conditions for efficiently converting the pump laser energy to a usable and reliable tunable radiation. We also investigate amplifier staging configurations, optimize the pump and signal optical systems, and run experiments to confirm and improve the developed amplifier models.

Electro-Optics Facilities. We have two laboratories devoted to electro-optics development.

The **Electro-Optics Laboratory (EOL)** is a newly renovated facility for developing dye-laser-system hardware and laser-system diagnostics. It enables us to divorce dye-laser research and development from the Laser Operations Facility (SPP-II-Venus) whose primary function relates to AVLIS process experiments. In the EOL we develop, test, and package new technology for advanced AVLIS systems for ultimate use in the Laser Operations Facility.

In the **Electro-Optics Control Systems Laboratory** we develop both supervisory and servo-type control systems for application to the next generation of laser and electro-optical systems. The laboratory emphasizes the electronics, computer, and transducer development for specific control-systems applications. In 1980 we developed a laser pointing and centering system that controlled the position and angle of a laser to high precision ($\pm 3 \mu\text{m}$, $< \pm 10 \mu\text{rad}$). This system uses an analog control loop with a 100-Hz closed-loop bandwidth. We recently demonstrated equivalent performance with a

digital system based on using an LSI-11. We will use either or both of these basic approaches to satisfy control requirements of next-generation systems. In the coming months we will be developing the electronics for a waveform test facility to be constructed in 1981.

Rare-Gas Halide (RGH) Laser Facilities. We have three laboratories devoted to RGH laser development.

We use the **RGH Closed-Cycle Test Bed (CCTB)** for studying average-power scaling and operating-lifetime limits in high pulse-repetition-frequency (prf) RGH lasers and for developing practical components and subsystems of high average-power devices. Because this facility is a test bed there are frequent changes in the apparatus. The major components are a flow loop and the associated gas-handling and pumping systems, the laser head, and the excitation, control, and diagnostic electronics.

In 1980 we began fabricating hardware for the developmental **High-Power Prototype (HPP)** closed-cycle RGH laser device. When operational it will be an enlargement of our CCTB. It will be used to investigate XeCl lasers under near-plant conditions.

We use the **Pulsed Electron-Beam System** facility to investigate advanced techniques for the preionization and excitation of discharge RGH lasers with the aid of electron beams (e-beams). The major element of the system is a cold-cathode e-beam gun that produces a rectangular sheet beam.

For laser excitation we mate the e-beam gun with a discharge laser cavity. The e-beam enters the discharge cavity through a semitransparent electrode and propagates into the discharge volume along the direction of the laser electric field. By adding suitable discharge pulsers to the laser cavity we investigate the various modes of RGH discharge excitation and preionization assisted by e-beams.

Separator Facilities. We devote three laboratories to vaporizer and extractor development.

We commissioned the **Vacuum Test Unit** facility in 1978 to characterize e-beam guns in support of the Test e-beam, Regulis, and Mars projects. We assembled the Vacuum Test Unit, a $1 \times 1 \times 0.6$ -m-deep stainless-steel box, from spare components in mid-1978. We pump the system with a

trapped 15-cm oil-diffusion pump plus appropriate mechanical backing. The pump has a maximum measured pumping speed of 1200 litre/s with a base pressure of 7×10^{-7} Torr. Numerous ports on the system add greatly to its versatility as a test chamber. The system is equipped with a small e-beam power supply.

We use this facility as an offline vacuum chamber to test and debug Mars prototype tetrode carts, filament power supplies, and gun controllers, as well as to hot-test Mars e-beam guns prior to installation in the Mars facility.

We did not design the facility to meet the ventilation or containment requirements needed to handle or vaporize uranium.

We established the **Extractor Test Facility** for extractor development and began operation in the spring of 1979. We are presently using the facility to develop the electromagnetostatic (EMS) extractor. This facility consists of the experimental chamber and controls, diagnostics, and computers. In 1980 we upgraded our Test e-beam facility to improve both the e-beam gun and the diagnostics capability. We use the new facility (**Minerva**) to study plasma physics and extraction development.

Integrated Systems Facilities. We have four major facilities usable for integrated separation demonstrations.

We use the **SPP-II** facility for generating the three dye-laser wavelengths required for the three-step uranium photoionization process. Eight 15-W discharge-heated CVLs serve as pumps for the various dye-laser cells in the SPP-II system. Each of the three dye-laser chains consists of a closed-loop stabilized master oscillator followed by three stages of optical-signal amplification. The stabilization system comprises a reference interferometer that monitors the frequency stability of the three master oscillators. Frequency drifts generate error signals that drive piezoelectric crystals for adjusting laser-cavity length, or intracavity etalon tilt, to maintain stable single-frequency operation.

The **Venus** system consists of three dye-laser power amplifiers energized by an array of 32 CVLs. We transport low-power dye-laser beams from the SPP-II laboratory and inject them into the Venus dye-laser amplifiers to increase their average power. We then combine the beams onto a common

axis before transporting them to an isotope-separation chamber.

The **Regulis** facility provides the program with a complete enrichment capability and a stable well-diagnosed experimental environment for ion-extractor and process-physics experiments. The **Regulis** vacuum chamber incorporates all the major components required for laser isotope separation. We produce atomic vapor with a 15-cm-strip e-beam gun heating a crucible containing uranium. We have at our disposal a large variety of computer-controlled diagnostic systems for the large number of experiments we perform in **Regulis**. With **Regulis** we can investigate a variety of extractor configurations and can accept laser input from either SPP-II or Venus.

The **Mars** laboratory is divided into three bays. The west bay contains the power-supply tetrodes, the east bay contains control and cleanup rooms, and the center bay houses the experiment module. The **Mars** facility represents a significant scale-up in engineering design and operational capability and is intended as a demonstration of large-scale source performance and high-output enrichment capability. It will also serve as a test bed for engineering concepts and hardware to be used in subsequent more-advanced enrichment facilities. **Mars** is connected with the other AVLIS facilities by a tunnel capable of transporting laser beams. It has independent facilities for source power, cooling, and material handling within **Mars**. We designed the mechanical, electrical, and other support facilities to permit upgrades in system capability for the requirements of future experiments.

Supporting Research Facilities. Five of our laboratories are devoted to basic research supporting our AVLIS process.

In the **Atomic Lifetime and Relaxation Laboratory** we determine the lifetimes and absolute absorption cross sections for atomic and ionic transitions of interest to the LIS program. We use nitrogen-laser-pumped dye lasers (oscillator and amplifier) to selectively excite uranium atoms. We employ a photomultiplier to monitor the fluorescence signal originating from the excited state. We then digitize the complete waveform of the fluorescence signal and average the digitized fluorescence signal with a signal averager.

The **Ultraviolet Laser and Raman-Scattering Laboratory** contains the infrared-VUV stimulated Raman scattering (SRS) pulsed light source used for one- and two-step uranium spectroscopy measurements and for Raman frequency-shifting studies. Furthermore, we pipe light from the Nd:YAG-pumped dye laser/SRS generator into the Atomic Lifetime and Relaxation Laboratory to perform kinetics and lifetime measurements on atomic and ionic uranium and to investigate advanced applications of new laser technologies.

We are converting the crossed molecular-beams apparatus into a facility for the measurement of electron-impact excitation and ionization cross sections, the **Electron Physics Laboratory**. The major apparatus change is the addition of an electron gun that projects its beam through the scattering volume. We will retain mass spectroscopic and fluorescence-detection capabilities. Initial experiments will measure electron-uranium collision parameters of significance to the extractor-development program.

The **Surface-Physics Laboratory** contains a set of ultrahigh-vacuum surface-analysis systems and tunable pulsed dye lasers. The surface-analysis system includes an electron gun, an Auger electron-energy analyzer, and a sputter ion gun. We use this facility to determine the secondary-electron emission yield for uranium surfaces bombarded by uranium ions. We will use laser-induced fluorescence to study the chemistry of excited atoms in the presence of solid surfaces.

The **Process Laboratory** contains four distinct facilities for performing physics experiments on various aspects of the AVLIS process. The equipment (details below) includes three nitrogen-pumped pulsed dye lasers, five argon-ion laser-pumped dye lasers, two high-power YAG lasers, a sophisticated laser-wavelength acquisition setup, uranium- and sodium-vapor sources, a very fast transient digitizer, and a data-acquisition system.

● **Laser-Diagnostics Facility.** In the Process Laboratory we constructed the laser facilities for performing uranium e-beam source diagnostics and for producing large-volume laser ionization in any of the uranium-vaporizer or extractor-development facilities. The system for monitoring vapor conditions includes

several cw dye lasers and fiber-optics transport. By simultaneously propagating two different cw lasers down the fiber and monitoring absorption in the vapor we can determine relative populations of the metastable levels. We can also measure absolute vapor densities, flow velocities, and Mach numbers. The large-volume ionization system uses a high-power (~50 mJ/pulse) YAG-pumped pulsed dye laser to produce laser ionization over extractor-sized volumes.

- **Beam-Propagation Facility.** In the Process Laboratory we have set up an experiment for studying beam-propagation physics. We propagate a single wavelength of the high-power pulse-amplified cw dye-laser system through high-density vapor and observe the temporal, frequency, and spatial behavior of the transmitted beam. For these experiments we have available both a high-density sodium oven and a transmission tower to the high-density e-beam-generated uranium vapor. The diagnostics include scanning Fabry-Perot interferometers, fast (<1 ns) transit digitizers, and video spatial monitoring.

- **Atomic Spectroscopy Facility.** In this facility, we measure the atomic parameters needed to evaluate two- and three-step AVLIS schemes. This includes lifetimes, branching ratios, hyperfine structure, isotope shifts, and absorption cross sections. We completed much of this work in past years, but keep this facility in readiness to address issues that surface from time to time.

- **LEA Facility.** We are presently assembling a system capable of producing three single-frequency fully modulated wavelengths with plant-scale fluences at a 10-Hz repetition rate. This system is almost entirely assembled with commercially available equipment and is hence named LEA for Lasers Easily Available. The setup includes three cw single-frequency ring dye lasers and a frequency-doubled YAG laser for pulsed amplification. We obtain an output power of 15 mJ per wavelength for a total of 45 to 50 mJ per pulse. This is sufficient to pump large-scale separation volumes. We are currently using a single wavelength of this system for pulse-propagation studies.

Other Facilities. Additional facilities include Vulcan, a high-temperature gas

heater, and a laboratory devoted to plutonium spectroscopy.

The **Vulcan** high-temperature gas facilities form a complete system, initially configured for multicompression heater development, that provide long-range capability for a wide variety of LIS and high-temperature chemistry experiments. The completed heater facility will heat and handle helium, neon, or a mixture of the two gases in the 4500- to 9000-K range. At present the following equipment is operational:

- Heater drive train.
- Isolation base for process line.
- Heater ancillaries.
- Gas supply, mixing, and associated controls.
- Exhaust-gas pumping and handling systems.
- Data acquisition and experiment/process control system.
- Various offline facilities.

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Planned Future Facilities

Introduction. As our Laser Isotope Separation (LIS) program moves from the technology-development phase to the engineering-development phase, we need to expand our facilities. To meet these needs we will be constructing three major new buildings during the next three years:

- Building 178, a \$630K new general-purpose laboratory facility.
- Building 272, a \$1M Laser Electro-Optics (LEO) facility housing a large tunable laser system.
- Building 482, a \$25M Advanced Isotope Separation Facility (AISF) combining laboratories for advanced development of LIS components and office space for 175 people.

AIS Facilities

Figure 10-5 presents a site plan showing the relationship of these new facilities to existing LIS facilities.

Building 178. Scheduled for construction bidding in early 1981, Building 178 will provide 740 m² of general-purpose laboratory space, to be allocated as program needs dictate. The building will be located adjacent to Building 175 and will provide space in convenient proximity to our existing facilities.

Building 272—Laser Electro-Optical Facility. Scheduled for completion in FY 81, LEO facility (Building 272) will house a large CVL-pumped dye-laser system to provide laser light for isotope-separation experiments in other facilities. This building, shown in Fig. 10-6, will contain the laser systems, their supporting equipment, the laser-control system, a laser-maintenance area, and a small meeting room. LEO will be located near the intersection of Avenue B

Table 10-1. Characteristics of advanced isotope separation facility.

General		Laboratory	
Laboratory space	7300 m ²	Structure	Modular on 14-m spans, 85% 1-story slab on grade, 15% high bay with basement
Office space	5000 m ²	Laboratories	45
Construction period	FY 82 through FY 83	Conference room	1
Total cost	\$25M	Offices	2
Office		Special features	Clean rooms Uranium-qualified HVAC Laser inertial bases Adaptable utility-distribution system
Structure	2-story closed circulation		
Offices	140		
Personnel housed	175		
Main briefing room capacity	120		
Conference rooms	2 at 42 m ² 2 at 21 m ²		
Libraries	1 principal 2 local		

Fig. 10-5. Site plan showing relationship of future buildings to existing buildings.▼

and the Inner Loop Road. Underground pipes will transport the laser light to separator facilities in Buildings 175 and 332.

Building 482—Advanced Isotope-Separation Facility. The planned AISF shown in Fig. 10-7 is comprised of two sections: a 7300-m² laboratory wing and a two-story 5000-m² figure-eight-plan office wing housing 175 people. Construction will begin in early FY 82 and will be completed near the end of FY 83. Table 10-1 lists the characteristics of the facility.

- **Siting.** As Figure 10-5 shows, the AISF will be located east of the Nova complex now under construction, with the office building sharing an open green space with the Nova office building. The orientation and location of the AISF office will result in a functional and appealing setting. Staff will enter the building from the northwest directly from the large parking area. Pedestrians crossing the commons from other laser facilities will enter from the southwest, while visitors and those using laboratory transportation

will enter from a smaller parking area to the southeast.

The laboratory building will be located northeast of the offices with entries from the offices and directly from the northeast. The corporation yard will be north of the laboratory and will contain the utility pads providing near access to systems within the building. We will run utilities to the buildings from the west and south from existing utility bundles that will be extended as part of this project.

- **The Laboratory Building.** The laboratory will house a variety of technical activities whose emphasis will evolve as the program itself evolves. To accomodate these changes we will build the laboratory as a largely open structure with flexible space differentiation and a flexible distribution system for utilities and system interconnections. Specialized spaces will include large laser and separator high bays; clean dye-laser laboratories; maintenance and storage for laser and separator systems; laser, separator, and

Fig. 10-6. Future LEO facility, Building 272.

Fig. 10-7. Future AIS Facility, Building 482.

plasma-physics laboratories; electrical and mechanical shops; computer-control system and development laboratories; shipping, receiving, and general storage; two access-control offices; and a conference room.

The basic structure of the laboratory, including the high bays, will be modular on 14-m spans. The interior walls are not load bearing and can be moved. The regularity of this structure should minimize the cost.

- **The Office Building.** The design criteria of security and energy conservation led us to concentrate on closed-circulation (looped) floor plans with individual (rather than open) office areas (Fig. 10-8). The building shape is such that a rectilinear layout still results in a diagonal appearance that complements the Nova office building. Our two-story design will facilitate communication among the staff and provide organizational flexibility in an appealing environment.

Fig. 10-8. Floor plan for Building 482 (level 1 only).

few people (several offices) to 50 (one floor loop) or 100 (two loops) can be housed close together, with changes in group size readily accommodated. Movement of people from the building core to their offices is direct, thus minimizing trespassing. The clustering of offices breaks the hall into short lengths and provides alcoves for clerical and group functions. Wider sections will allow congregation without infringing on individual work spaces.

The closed floor plan uses courts to provide exterior space for interior offices. In this design one court is open and landscaped, while the other is covered and contains the main briefing room and the library. Stairway locations provide ready local vertical access. Common office functions are provided in the building core, while conference areas are dispersed.

This office should provide a productive, adaptable, and attractive environment as our program evolves.

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Technical and Systems Highlights

In this article we summarize some of the notable specific developments in the LIS program during 1980, as our program continued its steady advance along our planned development path.

The separation unit forms the core of our AVLIS process. Figure 10-9 shows the Mars three-chamber vacuum vessel. The central chamber, here seen in the open position, houses the separation subsystems. The end chambers house optical subsystems for directing laser beams, which enter through the flanged ports. With the introduction of the Mars facility we have advanced significantly toward demonstrating full-scale technology and have achieved scaling well beyond the level of an experimental benchtop prototype.

Fig. 10-9. Mars vacuum chamber.

AIS Facilities

Together with the separator, the laser system forms the second basic element of our LIS program. Given the goals of our AVLIS program, it is necessary not only to develop high power in single laser heads, but even more important, we must show that it is possible to build, operate, and maintain integrated systems containing many laser heads. Our most recent step in this direction, the Venus laser facility, contains 32 copper-vapor-laser (CVL) heads, each rated at 15 W and based on 1977 technology. Typically these oscillators are part of six-head master-oscillator-power-amplifier (MOPA) units, as diagrammed in Fig. 10-10. Figure 10-11 shows a typical six-head MOPA unit in the compact space frame of the Venus system. Each of the CVL heads slides out of the space frame for servicing. The modular design and quick-disconnect features make it possible to repair or replace one CVL without disrupting the operation of the rest of the system. We have packaged similarly

the electronic components for each laser in replaceable modules.

The CVLs form only one part of the AVLIS electro-optical system. An equally important part, the dye laser system, converts the fixed-frequency power into tunable radiation. Figure 10-12 diagrams schematically the dye laser system. A master oscillator feeds the desired frequency into a chain of highly saturated transversely pumped dye amplifiers, which we combine for transmission into the separator chamber. Figure 10-13 shows one of our dye amplifiers. The dye fluid, flowing nearly vertically through the unit, passes through an active volume pumped by the green light from a CVL. The dye laser here appears as a red beam.

As the AVLIS program moves towards a production plant, we must develop hardware of steadily increasing size. One of the most interesting and challenging of these development tasks is the scaling of the

Fig. 10-10. Schematic diagram of Venus 6-head MOPA unit.

Fig. 10-11. Typical Venus MOPA unit.

Fig. 10-12. Schematic diagram of dye laser system.

Fig. 10-13. Dye amplifier unit. ▼

process lasers. By combining estimates of the throughput, laser parameters, and vapor capacity, we estimate that a full-scale AVLIS plant (10^7 SWU/year) will require 10^5 W of laser power. To this level of approximation, we can regard 10^5 W as the power required of either the pump laser or the tunable laser.

At present the most promising lasers for large-scale AVLIS applications appear to be CVLs. Figure 10-14 shows development schedules for these lasers. In this figure the horizontal color band is the range of single-aperture power that will be required for a commercial plant assuming 10^2 to 10^3 units per 10^7 -SWU/year plant. The data points (open circles) represent performances already obtained at LLNL through 1980. We have now demonstrated CVLs at plant scale. Prototype development beyond this point will emphasize improving component lifetimes and reducing laser costs.

Our mainline process relies on CVLs as pumps. However, a constraint on the maximum bore diameter at a given pulse-repetition frequency (prf) limits the average output power of these lasers. The constraint on the bore diameter is usually attributed to the requirement for wall deactivation of the metastable lower laser level during the interpulse period. Although one can circumvent this constraint by using lasers that have an annular or slab cross section, such geometries lead to complex engineering problems and more difficult optical requirements. We therefore continued research

on volume-deactivated gas mixtures for CVLs.

Fig. 10-14. Schedule for pump CVL development.

During the last year we achieved a major advance in laser technology with the first successful operation of a large-bore (7.5 cm) longitudinal-discharge CVL. Our

AIS Facilities

experiments proved the practicality of volume scaling by increasing the laser aperture. Figure 10-15 shows the output beam of our large-bore copper laser, when configured as an oscillator and directed into a power meter. We have achieved over 85 W of average power and 135 W of maximum extractable power at 5 kHz prf. The uniformly filled circular output beam evident in Fig. 10-15 greatly simplifies the optics for MOPA design.

Figure 10-16 shows the stainless-steel water-cooled vacuum housing surrounding the alumina plasma tube that comprises the large-bore CVL.

Although we have identified a mainline AVLIS process, the large number of transitions in atomic uranium makes it possible to build successful processes around a number of tunable laser systems. Figure 10-17 shows the laser options under study at LLNL. They are: three-step photoionization using CVL or frequency-doubled Nd:YAG-laser-pumped dyes; two-step photoionization using XeCl-laser-pumped dyes; and one-step AVLIS with a Raman-shifted ArF laser.

Although technology for our mainline process progresses satisfactorily, we will continue technology evaluation of alternate

Fig. 10-15. Output beam from a large-bore CVL.

Fig. 10-16. Large-bore CVL unit.

process options. As an example, excimer lasers provide a possible alternative to CVLs. Figure 10-18 shows our most advanced rare-gas halide (RGH) laser, the closed-cycle test bed (CCTB). This system comprises a flow loop that continuously circulates a rare-gas halogen mixture through the discharge excitation region (the rectangular cross section of the loop), thereby permitting repetitively pulsed laser operation. Adjacent to the flow loop are power-conditioning systems for the gas-excitation discharge and for the preionizing electron gun. Operating the CCTB as a XeCl laser, we have achieved 70 W maximum average output power at a prf of 850 Hz. We are currently building a scaled-up RGH laser system, the high-power prototype (HPP), potentially capable of operating at levels of average power comparable to the Venus CVL system.

Authors: J. I. Davis and B. W. Shore

Fig. 10-17. Alternative AVLIS excitation schemes.

Fig. 10-18. RGH closed-cycle test bed (CCTB) laser.

Laser Research and Development

Introduction

Our aim in laser R&D is to develop laser systems that can provide efficient and reliable laser power at low cost.

Photoionization of uranium requires either one, two, or three excitation steps with light in the ultraviolet-visible-infrared range

to provide the 6-eV ionization energy. Dye lasers provide tunable light of wavelengths ranging from about 330 to 1000 nm and thus overlap the uranium absorption spectra of either two- or three-step excitation. Dye lasers can be excited, (or "pumped") either incoherently (with a flashlamp) or coherently (with a pump laser). Our systems currently employ the latter concept. We are evaluating two possible pump laser systems: copper-vapor lasers (CVLs) for the red light of the three-step process and rare-gas halides

(RGH) lasers for the blue light of two-step processes.

Our mainline processes, based on three-step excitation with CVL-pumped dye lasers, relies on proven technology. We have a fairly complete base of three-step spectroscopic data and the CVL engineering is proceeding well. Because a substantial fraction of AVLIS plant costs originate with the laser system, improvements in laser systems offer potential for overall cost reduction. By increasing the size of CVL pump heads we can reduce the number of components needed in a plant and reduce the cost per watt of laser power. Thus our recent CVL development effort has concentrated on increasing the size and reliability of our CVLs. The four articles beginning with "Copper-Vapor-Laser Introduction" describe some of our work toward these goals.

The dye lasers operate by stimulated emission from excited organic dye molecules dissolved in a suitable solvent. Although water is a possible dye solvent, solutions in ethanol provide broader spectral range, longer lifetime, and better conversion efficiencies. "Fire Suppression in Laser Dye Media" describes our efforts to eliminate the potential fire hazard that could exist when handling large quantities of ethanol solvent.

Our system studies have projected the possibility of a modest reduction in net cost for an AVLIS process based on two-step excitation. At present the engineering development of RGH lasers lags behind that of CVLs. We are therefore continuing to develop and evaluate RGH lasers and their system constraints to improve our understanding of their ability to accomplish a reduction in cost. Finally, "XeCl Excimer Laser Injection-Locking and Amplifier-Energy Extraction" describes a portion of our work on RGH lasers.

Author: M. L. Spaeth

Copper-Vapor-Laser Introduction

We use copper-vapor lasers (CVLs) to optically pump the pulsed tunable dye lasers used in the AVLIS process. Thus the average power, efficiency, and cost of the overall

laser system depend heavily on the properties of the CVLs. For several years we have pursued the development of the CVL to increase its average power and efficiency, to reduce its cost, and to improve substantially the laser lifetime. During 1980 we made great progress in all these areas.

In late 1979 we demonstrated that volumetric deactivation of the metastable lower laser level in copper could be exploited to build large-bore cylindrical CVLs. While we were able to produce over 55 W of output power with this approach, we were still hampered by failures of the alumina tube. By redesigning the laser head we eliminated the need for high-temperature seals. Thus we now have a reliable laser head that is capable of repeated cycling and long-life operation.

In addition to improvements in the laser mechanical design, we increased our capability to provide the electrical input power to drive the CVLs. We are successfully operating alternately fired thyratrons from a single command charge circuit, which has allowed us to double the pulsed power available (to ~ 18 kW) for CVL experiments. At the same time we are pursuing advanced switching concepts, including totally solid-state magnetic circuits. If successful, these approaches would eliminate entirely the requirement for thyratrons.

We obtained a wealth of data on the dependence of the laser performance on various operational parameters (buffer-gas pressure, copper density, repetition rate, etc.) by operating a large-bore laser on a routine basis. These data in turn provided us with a good experimental base with which to compare the results of extensive computer modeling of the CVL. In the two articles that follow we describe both the laser performance and the computer-modeling results.

Author: R. E. Grove

Copper-Vapor-Laser Performance

During 1980 we developed a reliable long-lived high-power copper-vapor laser (CVL) emitting up to 85 W as an oscillator. We life tested this laser for more than 300 h at 70 ± 5 W. We measured the maximum extractable power (MEP) from the laser

using it as an amplifier in a master-oscillator-power-amplifier (MOPA) configuration. We measured 118 W MEP under conditions in which the laser was operating at 75 W as an oscillator. Its efficiency η , which we define as the ratio of MEP to input power from the dc supply, was greater than 1.2%.

We began the development of this laser in 1979 (see Ref. 1), when it became apparent that bore-diameter scaling was possible. Earlier, Bokhan² and Smilanski³ reported laser operation in smaller-diameter tubes using neon as a buffer gas at several tens of Torr. At these pressures and tube diameters, wall collisions in the interpulse period could not depopulate the metastable ²D state in copper and therefore could not account for the observed performance. On the other hand a volumetric depopulation process (depopulation by one of the plasma species) would enable the laser to operate as described at the reported neon pressures. Based on this evidence of a volumetric depopulation process we decided to build a larger-diameter CVL and to operate it with neon as a buffer gas at several tens of Torr.

We built and tested a 73-mm-diam-tube in 1979. This laser, described in Ref. 1, produced 55 W of average power using 25 to 35 Torr of neon as a buffer gas. We

encountered many mechanical problems with the design of this laser, but the most serious one was the fracturing of the ceramic plasma tube or the ceramic sealing collars. When this happened, usually during the second or third thermal cycle, the tube could not maintain the neon gas pressure inside the tube and we had to rebuild the laser. We realized that we needed a new design.

Internally Insulated CVL Design. Figure 10-19 shows a schematic drawing of the new CVL. It has three concentric cylinders: an inner AD-998 alumina tube, a Pyrex tube, and an outer stainless-steel water jacket; it also has a ceramic break. A fibrous ceramic insulation fills the annular region between the alumina tube and the Pyrex tube. The O-rings located at each end of the Pyrex tube form a vacuum seal. The region between the stainless-steel shell and the Pyrex tube is evacuated to 10 to 30 μ , while the volume inside the glass is filled with 30 to 50 Torr neon.

A high-repetition-rate electrical discharge heats the alumina tube to $1450 \pm 50^\circ\text{C}$ in the neon buffer gas. At this temperature copper inside the tube vaporizes to the optimum gas density. In equilibrium the electrical power into the discharge equals the power conducted through the insulation and radiated from the Pyrex tube to the stainless-steel water jacket. Generally this power is

Fig. 10-19. Large-bore CVL schematic.

between 2500 and 3500 W, depending on the design. There is an additional transfer of power out the ends of the alumina tube because of thermal radiation. Consequently there is a portion of tube on each end, usually 3 to 4 tube-diam long, that is considerably cooler than the center of the tube. The hot zone, which we define as the length of the tube that is approximately 1400°C or higher, is limited in length by the cooler ends.

There are several advantages to this design over the previous one. The first is the elimination of the problem caused by failure of the ceramic tube. The O-ring seals on the Pyrex tube maintain the pressure differential between the inside of the laser and the evacuated region. Since the Pyrex is at a relatively low temperature the seal is under minimal stress and does not fail. Furthermore the stress on the alumina tube is also reduced, since it is uniformly supported by the insulation. This prevents bowing and allows the tube to expand and contract freely. With this design, complete failure of the plasma tube has no effect on laser performance.

The second major advantage of this design is the extended hot zone. As mentioned above, the ends of the alumina tube are cooler than the center because of radiation cooling. In this design the electrodes are 3 to 4 diam in length and are surrounded by the ceramic insulation. This raises the temperature of the ends of the alumina tube and effectively creates a hot zone from electrode to electrode. This in turn increases the laser power and efficiency.

Oscillator Performance. We first evaluated the performance of a 6.0-cm-bore-diam oscillator with an electrode separation of 122 cm. We employed flat-flat optics and we

used a 99% reflector at 510.5 and 578.2 nm and a 4% reflecting output coupler. We found that this laser produced 60 to 65 W average power using 25 to 35 Torr of neon and 50 to 55 W using 5 to 7 Torr of helium as a buffer gas. As shown in Fig. 10-20 these power levels (in neon) were achievable at 4 to 5 kHz, but dropped steadily as we increased the repetition rate from 5 to 7 kHz.

The second laser we evaluated had a 7.3-cm bore diam and an electrode separation of 122 cm. We used the same flat-flat resonator. This laser produced 85 W average power at 5.0 kHz pulse-repetition rate using 21 Torr of neon as a buffer gas. The input power from the dc supply was 11 kW. As shown in Fig. 10-20, this laser was also relatively insensitive to repetition rates between 4 and 5 kHz, but performance fell off significantly at higher rates.

The third laser we evaluated had a 7.0-cm bore, but we increased the interelectrode separation to 152 cm. We used the same flat-flat resonator as on the previous two lasers. We obtained 82 W average power at a pulse-repetition rate of 5 kHz using 35-Torr neon as a buffer gas. Again the output-power variations with repetition rate were similar to the previous two lasers.

Amplifier Performance. In a high-power laser system we would use the large-bore CVL as an amplifier. Consequently we measured the MEP of the 7.0-cm-bore tube using a MOPA configuration. In this optical arrangement we used a 2.5-cm-bore, 15-W CVL as an oscillator with an unstable resonator whose beam was expanded to fill the inner diameter of the 7.0-cm amplifier. We measured the output power from this oscillator-amplifier configuration as a function of the input (optical) power to the amplifier. Figure 10-21 shows the results.

In Ref. 1 we outlined how the data in Fig. 10-21 can be used to measure the MEP or stored power. Using that analysis we found the maximum extractable power from the 7.0-cm-bore tube to be 118 W under conditions in which it produced 75 W as an oscillator. This implies that there is 57% more power available from this CVL when used as an amplifier rather than as an oscillator. The amplifier efficiency η is greater than 1.2%.

CVL Life Test. We life-tested the 7.0-cm-bore CVL to gather data on the reliability of many of the components. We anticipated

Fig. 10-20. Output power vs repetition rate for CVL.

that the plasma-tube insulation-Pyrex-tube combination would operate many hours without failure. In fact, as mentioned earlier, the reliability of this combination was the main reason for using it rather than the "scarf-joint" design discussed in Ref. 1.

We believed that an important design parameter was the copper-loss rate from the laser. Rough calculations for loss caused by diffusion in 50 Torr of neon indicated a loss rate of less than 0.01 g/h per tube end. However, we suspected that copper might be transported from one end to the other by the flow of the buffer gas, in addition to diffusive transport. By operating the laser for long periods of time until all the copper had condensed on the cool surface, we were able to determine the loss rate. This became the second reason for life-testing the new design.

Although we expected the laser head to operate for long periods of time, we were uncertain about other laser subsystems. The electronic components were the major concern. We had little data on the lifetime of the thyatron or capacitors operating in this high-voltage, high-repetition-rate mode. We knew that the thyatrons had three different failure modes:

- Failure to hold off anode voltages.
- Failure to trigger with a grid pulse.
- Failure to deionize sufficiently in the interpulse period.

A life test would provide not only data on the thyatron failure rate and mode, but also give us the opportunity to modify the circuits and evaluate the effects.

The life test of the 7.0-cm-bore CVL lasted for over 300 h and was still in progress at the end of 1980. The laser operated at 70 W during 90% of this time, but we also allowed it to drift as low as 65 W and as high as 82 W. To date there appears to be no degradation of the plasma tube, insulation, or Pyrex tube, confirming that the laser head is long-lived and reliable.

All the original copper condensed at the anode end of the tube after 230 h of testing. This indicated an effective loss rate of 0.2 g/h, an order of magnitude higher than the diffusion loss rate. However, this rate is within a factor of 3 of the calculated rate because of the flow of the buffer gas in the laser. Therefore, we believe that the copper-loss rate is controlled by flow rather than diffusion. In experiments subsequent to this life test, we operated the laser sealed

Fig. 10-21. Amplifier performance of large-bore CVL (7.0 cm).

off and found no performance degradation. We presume that under these circumstances we can reduce the copper-loss rate to that of the diffusion-loss rate.

The first thyatron used in the life test failed after 70 h, when it physically broke at the grid-ceramic interface. We installed a second thyatron, and it failed after only 35 h because of inadequate hydrogen pressure. At this point we modified the electronic circuitry by installing a back-swing clipper diode. This prevented the thyatron's anode voltage from going negative and causing a reversed breakdown in the switch. We then installed the third thyatron and operated it for 130 h before it failed because of inadequate hydrogen. We installed a fourth thyatron, and it is still operational.

The life test is still in progress, but we will terminate it early in 1981. The test provided us with much-needed data and confirmed the reliability of the new design.

Author: R. S. Anderson

Copper-Vapor-Laser Modeling

Last year we put a considerable effort into developing a kinetics model of the copper-vapor lasers (CVL). We used the computer code CUKIN to help us to understand the controlling kinetic parameters of the CVL. Through this understanding we can optimize our mechanical and electrical designs, guide prototype designs, and estimate scaling performances and risks. This model is particularly valuable since the performance of the CVL is difficult to parameterize experimentally because of the

complicated relationship of the macroscopic experimental parameters to detailed kinetic processes that determine laser performance. For instance, the electrical circuit (i. e., applied voltage, capacitance, and repetition rate) and the mechanical considerations (i. e., diameter, length, and heat shield) each influence detailed kinetics such as electron density, electron temperature, gas temperature, copper density, and excited-state populations. These in turn influence the macroscopic performance parameters of the laser (i. e., average green power, yellow power, pulse length and timing, and efficiency). Using a computer model we can untangle the external parameters of the laser from the kinetic details and uncover the reasons for the observed responses of various laser devices.

Copper-Vapor-Laser Kinetics

Code—CUKIN. CUKIN is a simple rate-equation description of five electronic states of atomic copper and four of the buffer gas (helium or neon). The five copper states we

Fig. 10-22. Energy levels in copper and neon modeled in CUKIN.

Fig. 10-23. Equivalent circuit modeled in CUKIN.

considered, shown in Fig. 10-22, are the ground state, the 2D -metastable lower laser state, the 2P upper laser state, a pseudo state between the upper laser level and the ionization limit (in which recombination occurs), and the ionization level of copper. The states we considered for the buffer gas are the ground state, the metastable state, a pseudo state, and the ion level. These nine states are the minimum number required to adequately describe the important low-energy excitations in copper and in the buffer gas.

Laser radiation occurs between the 2P and 2D levels in copper. We find the time dependence of the excited-state populations in these levels by integrating the rate equations over the time of interest.

Transitions between levels occur because of electron collisions, heavy-particle collisions (such as Penning ionization), and radiation. We calculate the electron-collision rates by assuming a Maxwellian electron-energy distribution (electron-impact cross sections for copper are discussed in the next article). For fast-pulse discharges the electron-energy distribution is likely to deviate from Maxwellian form. It will be depleted in the high-energy tail, particularly above 15 eV where buffer-gas excitation is energetically allowed. However, we expect the low-energy distribution, which is important for pumping the copper levels, not to be significantly perturbed by this effect. Good agreement between code predictions and experimental results leads us to believe that the Maxwellian assumption provides a reasonable approximation for modeling the CVL.

In the code we also include a full-circuit model to drive the discharge kinetics through calculation of the applied electric field across the tube. Figure 10-23 shows the equivalent circuit. It includes a thyatron switch, which we model by an exponentially falling resistance, and a time-varying discharge resistance, which we calculate consistently using kinetic equations. The laser-head inductance we calculate from geometrical factors.

CUKIN differs significantly from previous CVL codes because of the solution technique we employ to obtain self-consistent results. Previous codes employed a guess-and-calculate technique starting from guesses for the initial conditions of

electron density, electron temperature, and copper metastable density. By running the code through the discharge pulse we obtained output powers and macroscopic parameters. However, this technique has not provided valuable parameterization of the CVL. We first employed⁴ the present scheme in a computer model for the copper halide laser, for which the time evolution between pulses of discharge species such as chlorine is an important factor in determining the output of the laser. In the pure CVL, interpulse effects also play an important role in determining the detailed behavior of the laser during the discharge pulse. Therefore our present technique starts from a first guess of the initial conditions, such as electron density and copper metastable density, and then proceeds to solve the rate equations through several discharge-pulse afterglow cycles until we reach a steady state in the initial prepulse conditions. This technique proved to be extremely powerful in providing self-consistent results for the repetitively pulsed CVL.

Figure 10-24(a) shows the time evolution of the electron density calculated by the code through the first five iterations of the discharge-afterglow cycle. We guessed the initial electron density to be $10^{12}/\text{cm}^3$, but after several pulses it builds to a prepulse value of $2.4 \times 10^{13} \text{ cm}^{-3}$. In Fig. 10-24(b) we plot the output power of the laser calculated from iteration to iteration. The code predicts that during the first iteration the laser will produce high power (85 W) because the copper metastable density is initially set to zero. However, after several cycles the metastable density builds up to its steady-state prepulse value of $6.4 \times 10^{12} \text{ cm}^{-3}$ and the laser power drops to an asymptotic value (57 W). We calculate all of the initial conditions for the kinetic equation in this manner and thus this code represents a self-consistent approach to modeling the laser.

Code-Predicted Results and Comparison with CVL Experiments. This year we collected a significant amount of data parameterizing the large-bore (6-cm i.d.) CVL. By using CUKIN we can simulate this laser and compare the results with the experimental observations. Our discussion will fall into three areas: the circuit response of the laser, the pump-pulse kinetics, and the afterglow kinetics.

The pulse-forming network for the CVL is simple (Fig. 10-23). However, the time-varying impedance of the discharge makes it very difficult to characterize the circuit response unless a discharge-circuit model and plasma model are coupled, as we do in CUKIN. Table 10-2 lists typical values of circuit parameters for the CVL.

Figure 10-25(a)–(c) shows, respectively, experimental measurements and code predictions of the voltage applied across the tube, the current through the laser tube, and the timing of the laser-light pulse. Typically the laser breaks down at an applied voltage of approximately 10 kV, after which approximately 1500 A flow through the tube. The light pulse occurs at nearly the peak current (which coincides with peak power transmission to the laser). The agreement between experiment and code is quite good, particularly in the shape of the voltage and current traces and the relative timing of the light pulse. We found that by varying the circuit parameters in the computer code the laser output power is not particularly sensitive to thyatron loop inductance, but it is a maximum when the inductances of the laser head and thyatron loop are nearly equal. We observed such behavior in the laboratory, where we found

Length	122 cm
C _{storage}	12 nF
C _{peak}	4 nF
L _{thyloop}	275 nH
L _{laser}	350 nH
P (neon)	30 Torr
T (copper)	1480° C

Table 10-2. Typical operating parameters of 6-cm-diam CVL.

Fig. 10-24. CUKIN code predictions.

that different thyatron capacitor arrangements did not significantly change the laser output power.

The CUKIN code predicts improvement in laser performance if we decrease the laser-head inductance. As we decrease both the thyatron loop and the head inductance to lower values, the amount of laser output increases for a given stored energy in the storage capacitor. Experimentally, the inductance of the head is fixed by the geometry of the discharge tube, heat shield, and vacuum envelope of the laser. If the inductance is extremely low, energy will be deposited in the thyatron because circuit response competes with thyatron switching time. This is important in the laboratory because we want reliable long-lived thyatron operation.

The CUKIN code helped us to understand a great deal about the detailed kinetics of the

discharge pulse. Electron density and temperature are extremely important in determining the pumping efficiency of the copper levels. Figure 10-26 shows code-predicted intrapulse kinetics. Figure 10-26(a) shows the electric field applied to the tube from the pulse-forming network along with the resultant electron temperature. At the peak of the applied field the electron temperature rises to the 3- to 4-eV range. This is significant since at low electron temperatures the lower 2D state is pumped more efficiently than the upper state, preventing formation of a population inversion. For lasing to occur the electron temperature must rise above approximately 2.5 eV to pump the upper level more than the lower level.

Figure 10-26(b) shows the predicted time variation of populations in the 2P and 2D copper levels. It takes approximately 90 ns for the 2P population to approach that of the 2D . In fact, the 2P population never increases over that of the 2D . Inversion occurs only because of the statistical weights of the two levels. The onset of laser action occurs when the 2P and 2D levels break from their parallel path into a ratio of 3:2 as given by the ratio of statistical weights. The laser pulse terminates when the 2D metastable level has filled up, caused by transitions from the upper laser level and by collisions with electrons of ground-state atoms.

Changes in the peak electron temperature help explain the variation of laser output power with copper density. In Fig. 10-27 we plot predicted and measured output power as a function of the copper-vapor temperature, which controls the copper-vapor density. We observe good agreement between code results and measurement. At low vapor temperatures there is insufficient copper density to achieve a large inversion. We achieve optimum output power with temperatures in the range of 1480 to 1500°C. The decrease in output power at higher copper-vapor temperature is caused by a decrease in peak electron temperature during the discharge pulse because of lower E/N . At 1525°C the CUKIN code predicts a rise in the peak electron temperature to only 2.8 eV, approaching the value where the lower metastable level is pumped as efficiently as the upper laser level. The experimental parameters that influence the peak electron

Fig. 10-25. CUKIN code predictions and measured data.

temperature are the external circuit (applied field), the buffer-gas type and pressure, and the copper density in the laser.

The afterglow period in the CVL provides the initial conditions for the next discharge pulse. During this important period the copper 2D metastable density decays so that inversion can be created on the next discharge pulse. Electron collisions in the afterglow are the kinetic processes that control the deactivation of the metastable state in large-diameter neon CVL. The parameters that control the rate of this deactivation are the electron density and temperature.

We calculate the rate coefficients for the metastable deactivation to be on the order of 10^{-9} cm³/s. At an electron density of 2×10^{13} cm⁻³ such cross sections lead to a decay rate of the copper metastable of 2×10^4 /s, which is fast enough to deactivate the metastable to the ground state during the interpulse periods, typically of 200 μ s duration. However, if the electron temperature is too high during the afterglow period, then not only are the metastables deactivated to the ground state, but ground-state atoms are in turn collisionally excited to the metastable. The ratio of the ground state to the metastable density is controlled by detailed balance. Therefore it is important for the electron temperature to decay to a value on the order of 2000 to 3000 K between discharge pulses so the metastable density can decay to the low 10^{12} -cm⁻³ range.

The parameters that influence the electron temperature and the electron density in the afterglow are important because they control the metastable deactivation. The parameters that influence the afterglow characteristics are the average power delivered to the discharge, the repetition rate at which the laser is operated, and the temperature of the buffer gas during the afterglow. The CUKIN code includes all of these kinetic processes and it should accurately model the decay of the metastable.

Figure 10-28 shows the predicted time variation of the metastable 2D density. In Fig. 10-28 we also plot the detailed balance limit of the copper 2D density. The metastable density approaches the detailed balance limit during the afterglow. For shorter interpulse periods the electron

temperature does not decay to as low a value and thus there is a relatively higher metastable density. This increase in metastable density causes the laser power to drop. Figure 10-29 shows this effect. In Fig. 10-29 we plot for the experiment and code predictions the output energy per pulse vs repetition rate. At low repetition rates the prepulse metastable density is low (the 2D density is 7×10^{11} cm⁻³ at 1 kHz), since there is plenty of time to deactivate the metastable. Consequently the single-pulse laser energy is high (~ 23.0 mJ). As the repetition rate increases the interpulse period decreases, causing the prepulse metastable density to

Fig. 10-26. CUKIN code predictions.

Fig. 10-27. Predicted and measured output power for 6-cm-diam CVL operating at 5 kHz with 35-Torr neon. ▼

become larger ($1.6 \times 10^{13} \text{ cm}^{-3}$ at 8 kHz), thereby decreasing the laser energy to 1.8 mJ.

Summary. The development and operation of a self-consistent CVL kinetics code has given us insight into the controlling kinetic parameters of the CVL. From the code we learned the importance of the pulse-forming network in coupling energy into the discharge. We now understand the importance of the peak electron temperature during the discharge pulse for determining the rate of excitation of the laser levels and understand the macroscopic parameters that influence this pumping rate. The CUKIN code also revealed the controlling kinetic parameters during the afterglow, that is, deactivation of the ^2D metastable because of electron collisions where the electron density and temperature play a crucial role in determining that quenching rate. We are using this code to optimize our designs, guide prototype designs, and estimate scaling performances and risks.

Fig. 10-28. CUKIN prediction of copper ^2D metastable density. Dashed line shows detailed balance limit.

Fig. 10-29. Measurement and CUKIN prediction of laser output-pulse energy as a function of repetition rate.

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Cross Sections for Electron-Impact Excitation of Copper

Collisions of low-energy electrons with copper atoms play an important role in the copper-vapor laser (CVL). We need to know the cross sections for the excitation and deexcitation of the low-lying energy levels to understand and quantitatively model laser performance under a variety of operating conditions. In 1974 Trajmar et al. measured^{5,6} relative excitation cross sections for the $(3d^{10} 4p) ^2P$ and $(3d^9 4s^2) ^2D$ levels between 6- and 100-eV impact energies. They placed these cross sections on an absolute scale by comparing them to the elastic cross section, which they in turn normalized to various theoretical cross sections.⁷⁻⁹ The uncertainty in the correct normalization of the cross sections was a source of confusion and concern for the kinetic modeling efforts.¹⁰⁻¹³ The cross sections normalized to the static-exchange elastic-scattering calculations⁸ were presumably the most accurate. However, recent models indicate that these cross sections may have been overestimated by a factor of 2 to 4 (see Ref. 13). We were motivated by this uncertainty to perform new calculations of the inelastic cross sections of copper atoms to improve the accuracy of the normalization of the experimental data.

We used the semiclassical impact-parameter method^{14,15} to compute the cross sections for the electron-impact excitation of both the upper laser levels $(3d^{10} 4p) ^2P_{1/2,3/2}$ and the metastable lower laser levels $(3d^9 4s^2) ^2D_{3/2,5/2}$. In this approximation the integrated cross section is given by Eq. (10) of the following article entitled "Collisions Between Uranium Atoms and High-Energy Electrons." We determined the quantity R_0 by the requirement that the Born approximation and the impact-parameter method give the same cross sections for the $4s \rightarrow 4p$ transition at a collision energy of 200 eV. The resulting value of 3.1 bohr is completely consistent with the expected radius of the

copper atom. Having determined R_0 , we calculated physically realistic excitation cross sections for the $2S_{1/2} \rightarrow 2P_{1/2}$ and $2P_{3/2}$ transitions using the experimental transition energies and oscillator strengths in the impact-parameter formula [i. e., Eq. (10), cited above]. We then normalized the data of Trajmar et al.⁶ for the $4s \rightarrow 4p$ transition to the computed cross section at 60-eV impact energy. Since the relative normalization of the measured $2S \rightarrow 2P$ and $2S \rightarrow 2D$ cross sections is correct, we renormalized the excitation cross sections of the metastable $2D_{3/2}$ and $2D_{5/2}$ levels by exactly the same factor.

Figure 10-30 compares the renormalized experimental cross sections to our theoretical values. Our present study yields inelastic cross sections for the copper atom, which are a factor of 3 smaller than those obtained in the earlier work.⁶ Figure 10-30 also shows that the calculated and the experimental cross sections exhibit different energy dependences even between 20 and 60 eV, where the agreement should be best. Very recent accurate close-coupling calculations¹⁶ found a similar discrepancy. We do not understand the source of this difficulty, and further theoretical work will be necessary to establish the correct energy dependence of the excitation cross sections, especially below 10 eV, which is the most important energy region for modeling the performance of the CVL.

Authors: N. W. Winter and A. U. Hazi

Fire Suppression in Laser Dye Media

Optimum performance and highest photochemical stability for dyes in the three-step atomic-vapor laser isotope separation (AVLIS) process occur in an alcohol such as methanol or ethanol. The alcohols are flammable, having flash points of 54° F (12.2° C) and 55° F (12.8° C), respectively, and their use in large dye amplifiers, where flow rates of 20 to 40 gpm are anticipated, could under unusual circumstances present a serious fire hazard.

Among laser-dye basic molecular structures, the ionic rhodamines offer the best chance for operation in water or water-

mixed systems. On the other hand such dyes are large hydrophobic molecules; that is, they resist wetting or solvation by water. For water-alcohol-mixed systems the vapor is richer than the liquid in alcohol, hence considerable dilution is needed for a useful flash-point elevation. For example, just to reach the flash point [105° F (41° C)] of trifluoroethanol—an earlier candidate solvent—requires an 85%:15% water-alcohol mixture.

A concentration of $10^{-4}M$, as used in our larger dye amplifiers, can be prepared from rhodamine 6G in water. As a result of molecular aggregation it shows spectral changes that are sufficient to interfere with laser action.¹⁷ Our studies on aqueous systems such as ethanol-water (1:1 volume:volume), dimethyl formamide-water (1:1 volume:volume), and ethanol-water with either dipropyl acetamide or Triton X100 additives showed that dye lasers behaved more poorly in aqueous solution than in pure alcohol solvent. Water and water-mixed solvents diminish dye photochemical stability.¹⁸

The use of quenching additives dissolved in the alcohol is an immediate answer to fire suppression. Earlier work¹⁹ described complete flame inhibition by a number of additives, for example, just 11% (volume:volume) bromochloromethane. A class of Halon fire-extinguishing agents includes materials even more effective than bromochloromethane and have high

Fig. 10-30. Theoretical and experimental electron-impact-excitation cross sections for ground state of copper atom.

photochemical stability and low toxicity.²⁰ Table 10-3 lists some of the agents. They quench by scavenging a chemical intermediate in the combustion (chain breaking) and typically require 6% of the material (by volume) in the vapor phase. Thus the more volatile the agent the less would be needed dissolved in the liquid to quench combustion. On the other hand too volatile an agent would evaporate rapidly, leaving behind the flammable alcohol. Favored additives are Halon 1202 (dibromodifluoromethane) and Halon 2402 (1,2-dibromo-1,1,2,2-tetrafluoroethane).

We measured flash-point elevations for methanol and ethanol with these two additives. We modified slightly a standard "open cup" ASTM flash-point test device to reduce radial heat losses. We ignited the mixtures by either spark or flame at the top plane of the open container to allow mixing of the vapors with air. We needed only modest concentrations of the agents to raise the flash points significantly. They were for methanol 5% by volume of Halon 1202 and 7% by volume of Halon 2402; for ethanol they were 2.5% of 1202 and 5% of 2402. Although the test apparatus may not scale directly to a laboratory environment, these Halon concentrations increased the flash points measured in the pure solvents by at least 20° C, well beyond a normal operating temperature of 73° F (23° C).

Such concentrations do not absorb the pump or fluorescence wavelengths. Thus absorption cutoffs at 210 nm for the pure alcohols are shifted to 290 nm for the Halon 1202 solutions and to 270 nm for the 2402 solutions. Also, they do not interfere significantly (<5% change) with laser performance as was demonstrated in a frequency doubled neodymium-pumped dye oscillator where the dye rhodamine 6G was the active medium.

In future measurements we will examine the effect of the Halons on prolonged laser operation; that is, on dye-medium photodegradation. We plan further study to define the effect of Halons on fire prevention. For example, planned vapor-phase measurements will address the problem of liquid sprays or suspended liquid mists.

Halon additives provide a reduction or complete removal of fire risk in large dye-laser systems such as an LIS plant. With the reduction in fire risk there is also a reduction in cost of plant construction. For example, a flammable liquid in closed containers comes under Class I, Division 2 of the category of hazardous locations of the State of California Electrical Safety Fire Code. The code specifies expensive design, construction, and equipment modifications. This technique of fire suppression also widens the option of organic solvents that may be used as dye laser media, for example in the two-step process.

Author: P. R. Hammond

XeCl Excimer Laser Injection-Locking and Amplifier-Energy Extraction

The ultimate use of rare-gas halide (RGH) lasers in high average-power applications depends on oscillator-amplifier staging. Despite considerable research on RGH lasers, many amplifier staging parameters are poorly known. For example, the level of nonsaturable loss relative to small-signal gain is one of the more significant unknowns, since even modest losses can severely limit amplifier output fluence and optical-extraction efficiency. To determine

Table 10-3. Physical properties of some halogenated fire-extinguishing agents.

Agent	Chemical formula	Molecular weight	Halon number	Boiling point (°C)	Density	Heat of vaporization (cal/g)	Agent effectiveness ^a		
							Molar	Weight	Liquid volume
Bromochlorodifluoromethane	BrClCF ₂	165.4	1211	-4	1.83	32	51	46	53
Bromochloromethane	BrClCH ₂	129.4	1011	67	1.93	—	39	45	56
Bromotrifluoromethane	BrCF ₃	148.9	1301	-58	1.57	28	100	100	100
Carbon tetrachloride	CCl ₄	153.8	104	77	1.59	46	35	34	35
Dibromodifluoromethane	Br ₂ CF ₂	209.8	1202	24	2.28	29	94	67	97
Dibromotetrafluoroethane	BrCF ₂ CF ₂ Br	260	2402	47	2.17	25	100	57	79

^aFreon FE 1301, Fire Extinguishing Agent, Dupont pamphlet B-29, E. I. Dupont de Nemours, Wilmington, DE. Values are quoted with respect to Halon 1301 on a weight-effectiveness basis and were recalculated in terms of molar- and liquid-volume effectiveness.

the important amplifier staging parameters we performed oscillator-amplifier experiments on ArF and XeCl media. Our investigations yielded information on small-signal gain, saturation fluence, energy density, and the level of nonsaturable loss.

Our results on ArF (Ref. 21) indicate a level of nonsaturable loss at least 50 times less than the small-signal gain. In addition, the measured saturation flux (1.6 MW/cm^2) was in excellent agreement with the value of 1.7 MW/cm^2 that we calculated from known values of cross section, spontaneous lifetime, and ArF* quenching rates.

We completed similar measurements on XeCl. Figure 10-31 shows the results. We varied the oscillator input-energy fluence E_1 to the amplifier with neutral density filters and measured the resultant output-energy fluence E_2 . The oscillator beam area ($1.9 \times 0.5 \text{ cm}$) closely matches the amplifier gain volume. Because we measure energy the data represent temporal as well as transverse spatial averages. We analytically model the effect of nonsaturable losses by assuming steady-state homogeneous extraction.²¹ Our assumptions are reasonable since the XeCl* lifetime ($\sim 7 \text{ ns}$, including quenching) is shorter than the optical pulse length (30 ns) and the spectral input is broadband. This model predicts a straight-line saturation characteristic on a graph such as Fig. 10-31 for low nonsaturable losses. The solid line through the data in Fig. 10-31 is a least-squares fit for low nonsaturable loss. The dashed curves show the expected effects if losses were more significant. These data, together with the model, place a lower bound on the gain-to-loss ratio of ~ 40 .

The parameters of the least-squares fit (assuming low nonsaturable loss) yield the small-signal gain ($12.3\% \text{ cm}^{-1}$), saturation fluence (6.66 mJ/cm^2), and energy density (0.820 J/litre) for the XeCl medium. The breakdown electric field and peak current densities for these measurements were 12.5 kV/cm and 790 A/cm^2 , respectively. The relatively low electric field available in these experiments accounts for the somewhat low optical-energy density. One consequence of low nonsaturable loss is that extremely high optical-energy extraction is possible. In these experiments we extracted 88% of the available energy ($72 \text{ mJ}/82 \text{ mJ}$).

As in the case of ArF, we can independently estimate the saturation flux

for XeCl from our experimental values for cross section, spontaneous lifetime ($\sigma T_s = 50 \text{ Å}^2 \text{ ns}$; $T_s = 11 \text{ ns}$) (Ref. 22), and XeCl* quenching rates for He ($1.9 \times 10^{-13} \text{ cm}^3/\text{s}$), Xe ($2.1 \times 10^{-12} \text{ cm}^3/\text{s}$) and HCl ($7.7 \times 10^{-10} \text{ cm}^3/\text{s}$) (Ref. 23). For the gas mixtures and pressures in these experiments (Fig. 10-31) the calculated saturation flux is 0.20 MW/cm^2 and for our 30-ns optical pulse width the saturation fluence is 6.0 mJ/cm^2 ; this latter value is in good agreement with our experimental value (6.7 mJ/cm^2). However, if electron-quenching rates are significant, the 6.0 mJ/cm^2 estimate is a lower bound. It is difficult to quantify this quenching contribution to saturation fluence, but at these current densities ($\sim 790 \text{ A/cm}^2$) it may be important.

Our extraction model (dashed curves in Fig. 10-31) assumes nonsaturable losses. The introduction of loss saturation in the model results in gain saturation curves that are much closer to linear than those calculated for our figure. Hence the experiments do not

Fig. 10-31. Characteristics of XeCl amplifier input E_1 and output E_2 .

exclude the possibility of significant saturable losses in XeCl media.

In any case the XeCl amplifier saturates linearly to a very high level of gain saturation, as shown by the experimental data. If the simplest interpretations are correct (low nonsaturable or easily saturable losses), XeCl amplifiers can achieve efficient optical-energy extraction ($\gtrsim 90\%$) at high saturation levels. However, our model indicates there are significant optical extraction efficiency penalties whenever the absorption length αl exceeds unity, irrespective of gain to nonsaturable loss ratio. Consequently a simple linear-amplifier scaling strategy can be used to achieve high pulse energy or average power only to the point where the total medium length is less than $1/\alpha$.

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Atomic-Vapor Physics

Introduction

Evaporated uranium, collimated into a vapor sheet as it emerges from the vapor-source region, can undergo a variety of collisions as it flows through the separation region, interacts with laser radiation, and eventually strikes either an ion or tails collector. For several years we have maintained an active theoretical and experimental research effort to help us understand the physical processes that occur during the course of vapor flow so we can operate our process most efficiently. The following articles highlight our activities in this area during the last year.

Author: B. W. Shore

Associative Ionization in Alkaline Earth Elements

In previous annual reports we described our work on associative ionization and other processes in the alkaline earths Sr (see Refs.

24 and 25), Ca, and Ba (see Ref. 26). This year we studied Mg and have extended our investigation of Ca. We completed measurements of the binding energies (dissociation energies) for the dimer ions of the alkaline earth series Mg, Ca, Sr, and Ba. In previous publications we reported satisfactory values for Sr and Ba.^{24,26} We wanted to determine if the dimer-ion binding energy in the alkaline earths showed systematics similar to those observed in the alkali metals.²⁷ For these the binding energy decreases as one ascends the series. We found the same behavior in the alkaline earths.

Our experimental procedures remain as previously described^{24,25}: One or two pulsed dye lasers, tuned to excite the atoms to the desired excited states through resonant transitions, irradiate a stream of atoms in a vacuum chamber. Depending on the process or reaction we are studying, we can vary the atom density in the interaction region from about 10^{12} to 10^{14} atoms/cm³. We detected ions with a quadrupole mass filter serving to select either the atomic ion or dimer ion of the element. Atomic ions are usually formed by photoionization but under appropriate conditions collisional and electron ionization can occur.²⁴⁻²⁶ Dimer ions are formed by associative ionization reactions.

Laser-excitation processes are

(1a)

where M is Mg, Ca, Sr, or Ba, m is 3, 4, 5, or 6, respectively, and $n \geq m$. Table 10-4 gives the λ_1 wavelengths for the four alkaline earths.

Associative ionization reactions observed in the alkaline earths are

where l is an s, p, or d electron. We had observed Reaction (4) for Ca, Sr, and Ba but we did not expect it nor did we see it for Mg. The reason Reaction (4) cannot occur in Mg

is that the energy of the $3s3p^1P_1$ level is more than half the ionization energy of Mg. In Mg Reaction (4) becomes

For associative-ionization Reactions (3) and (4) to occur, the excitation energy of the colliding pair must be lower than the ionization potential of the atom. The binding energy of the dimer ion formed must be such that the total energy exceeds or equals the ionization potential. Thus the defined conditions for associative ionization are

$$E < \text{IP} \quad (6)$$

and

$$E + \text{BE} \geq \text{IP} \quad (7)$$

where E equals total excitation of the colliding atoms, BE equals the dimer binding energy or dissociation energy, and IP equals the ionization potential of the atom.

The binding energy of the dimer can be obtained by observing the lowest Rydberg level that produces dimer ions by Reaction (3). This level and the next lowest Rydberg level that does not produce dimer ions can be used in Reaction (7) to calculate limits for the binding energy of the metals being studied. When this is done we obtain an energy value that the binding energy must equal or exceed and a value that the binding energy must be below. Table 10-5 gives these values for the four alkaline earths studied. In each case except Ca we bracketed the dimer-ion binding energy. The next lower accessible Rydberg level in Ca is $4s5p^1P_1$ at 4.554 eV, or 1.559 eV below the IP. We expect to investigate this level soon. The results given in Table 10-5 show that the dimer-ion binding energy is definitely greater for the lighter alkaline earths. This is similar to the trend found for the alkali metals where the binding energies are: Li_2^+ , 1.55 eV; Na_2^+ , 0.99 eV; K_2^+ , 0.79 eV; Rb_2^+ , 0.73 eV; and Cs_2^+ , 0.70 eV.²⁷

We plan experiments to determine if associative-ionization Reaction (3) occurs for the $4s5p^1P_1$ level in Ca. We also plan to investigate the potential excited-state-excited-state associative ionization reaction

Table 10-4. Excitation wavelengths λ_1 for Mg, Ca, Sr, and Ba.

Element	λ_1 for $ms^2\ ^1S_1 \rightarrow ms\ mp\ ^3P_1$ (Å)	λ_1 for $ms^2\ ^1S_0 \rightarrow ms\ mp\ ^1P_1$ (Å)
Mg	4571.1	2852.1
Ca	6572.8	4226.7
Sr	6892.6	4607.3
Ba	7811.3 ^a	5535.5

^aNot used in this investigation.

Table 10-5. Binding energies of alkaline earth dimer ions.^a ▼

Element	Ionization ^b potential	Lowest Rydberg Level yielding ^b		Binding energy	
		Dimers	No dimers	Greater than	Less than
Mg	7.646	5.946	5.753	1.700	1.893
Ca	6.113	4.624	— ^c	1.489	— ^c
Sr	5.695	4.920	4.766	0.775	0.925
Ba	5.212	4.434	4.261	0.778	0.951

^aAll values are in electron volts (eV).

^bIonization potentials and energy levels from the following sources: Mg, Ref. 28; Ca, Ref. 29; Sr, Ref. 30; and Ba, Ref. 31.

^cNot determined.

We tried to investigate this reaction in Ca but the production of Ca_2^+ by Reaction (4) prevented us from obtaining positive results. The Ca_2^+ signal decreased when the lasers producing $4s4p^3P_1$ levels were turned on [Reaction (1)], indicating Reaction (8) has a lower cross section than that of Reaction (4). In Mg, Reaction (4) cannot occur; thus an observation of Mg_2^+ when irradiating Mg vapor at high density with two lasers to give Reaction (1) and (1a) will show that Reaction (8) or the equivalent occurs.

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Collisions Between Uranium Atoms and High-Energy Electrons

Collisions between free uranium atoms and high-energy electrons play fundamental roles in several aspects of the atomic vapor laser isotope separation (AVLIS) process. Understanding and quantitative modeling of the microscopic-collision physics of AVLIS require cross sections for the excitation and the ionization of U I by electron impact. Since the available experimental data on electron-uranium collisions are very limited, we undertook and completed a theoretical study of two aspects of electron-uranium scattering. We calculated inelastic and momentum-transfer cross sections

associated with the excitation of U I to optically allowed (even parity) levels by electrons at impact energies between 5 and 10 keV. We also computed the total electron-impact ionization cross section of U I between 100 eV and 10 keV to extend the experimental values³² (measured up to 500 eV) to higher electron energies.

To determine the excitation cross sections we used the semi-classical impact-parameter method³³ and experimental optical-oscillator strengths³⁴ for 84 even-parity levels of the uranium atom. As shown by recent work at LLNL, the impact-parameter method is an effective and practical technique for treating the electron-impact excitation or deexcitation of optically allowed transitions in both atoms and molecules. In Ref. 35 we discussed recent developments of the method and its application to diatomic

Fig. 10-32. Calculated cross sections for electron-impact excitation (σ_{exc}) and ionization (σ_i) of ground state of uranium atom.

10 keV. The computed σ_{exc} reaches a maximum value of approximately $6.6 \times 10^{-15} \text{ cm}^2$ at 9 eV (not shown) and it decreases as $(\ln E)/E$ at high electron energies, compared to the E^{-1} falloff of σ_{mom} . Since we included only the levels lying below $31\,000 \text{ cm}^{-1}$ in our study, the computed σ_{exc} is a lower limit to the total cross section for exciting all the discrete optically allowed levels below the ionization threshold.

We also calculated the variation of the excitation cross section as a function of the energy transferred from the incident electron to the uranium atom. Figure 10-33 shows the differential excitation cross section $d\sigma/d\Delta E$ (in cm^2/eV) as a function of energy transfer ΔE for an electron-impact energy of 3.0 keV. The shape of this cross section is independent of the electron energy and bears a great resemblance to the differential optical-oscillator strength distribution of U I. This behavior is expected for electron-impact excitation of optically allowed levels at high incident energies. Unfortunately, no useful cross sections could be obtained for energy transfers between 4 and 6.2 eV, the first ionization limit. However, based on the magnitude of the threshold photoionization cross section and the expected behavior of the optical-oscillator strength distribution, one can assume that the differential cross section $d\sigma/dE$ will peak and then roll over for an energy transfer somewhere between 4 and 6 eV.

Finally, we studied the relative importance of excitation and ionization of uranium by electron impact. Halle and co-workers³² have measured the total ionization cross section σ_i , which is defined as

$$\sigma_i = \sigma_+ + \sigma_{2+} + \sigma_{3+} + \dots \quad (12)$$

Here σ_+ refers to single ionization, σ_{2+} to double ionization, etc. Since their data do not extend above 500 eV, we calculated σ_i by estimating the ratio $\sigma_i/\sigma_{\text{exc}}$ and using our calculated σ_{exc} as shown in Fig. 10-33. We can estimate the ratio $\sigma_i/\sigma_{\text{exc}}$, which is approximately independent of E for large impact energies, using optical sum rules. We computed the total transition probability from

$$M_{\text{total}}^2 = \sum_{ij}^{92} (\psi_0 \mathbf{r}_i \cdot \mathbf{r}_j \psi_0) \quad ,$$

where ψ_0 is the wave function of the ground state. We extracted the transition moments connecting the ground state to discrete excited levels from available optical data.³⁴ Figure 10-32 shows the theoretical ionization cross section along with the experimental points.³² Our values of σ_i between 200 and 500 eV agree very well with the measurements of Halle and co-workers.³² It is interesting to note that in uranium, excitation of the discrete levels dominates over ionization at high impact energies and this behavior is contrary to what is observed for most other atoms.

Author: A. U. Hazi

Electronic Excitation and Ionization in Intermediate-Energy U-U Collisions

We have generated estimates of the size and energy dependence of cross sections for electronic excitation and ionization in collisions involving two ground-state U atoms for collision energies above a few electron volts. Because of the enormous number of electronic states involved in this diatomic system, we are unable to perform detailed ab initio calculations for these processes. Fortunately we expect semiclassical methods to be valid for the energy range of interest and the high density of states actually becomes an advantage. We wish to know only the total cross sections for excitation rather than specific state-to-state transitions. Therefore we can use a complex interaction between the colliding atoms

$$V(R) = U(R) - iW(R) \quad (13)$$

The real part of this potential is known to be repulsive and can be obtained, for example, from Thomas-Fermi-Dirac or Gordon-Kim calculations.³⁹ The imaginary part, which determines the probability of an electronic transition occurring during a collision, we generate from considerations involving the molecular density of states as discussed below. The negative imaginary part of the potential reflects the fact that some of the flux in the initial channel (two ground-state U atoms) disappears in the course of a collision, i. e., irreversibly changes its electronic state and thus is "absorbed."

Using this interaction we calculate an elastic scattering or momentum-transfer cross section and an absorption cross section that represents the sum of all translational-to-electronic (T-E) processes. The momentum-transfer cross sections have been studied both experimentally^{39,40} and theoretically.⁴¹ We use these results to compare with our absorption cross sections and to calibrate the real part of our interaction potential.

We calculate the complex phase shift for the potential in Eq. (13) using a semiclassical or high-energy approximation. We write the phase shift in terms of its real and imaginary parts $\eta_\ell = \lambda_\ell + i\mu_\ell$. Mott and Massey⁴² show that the absorption cross section is given by

$$\sigma_A = \frac{\pi}{k^2} \sum_{\ell} (2\ell + 1) (1 - e^{-4\mu_\ell}) \quad (14)$$

We define the semiclassical phase shift by

$$\eta_\ell = \int_{F_1}^{\infty} F_1^{1/2} dR - \int_{F_0}^{\infty} F_0^{1/2} dR \quad (15)$$

where

$$F_0 = k^2 - (\ell + 1/2)^2/R^2 \quad (16)$$

and

$$F_1 = k^2 - 2\mu(U - iW) - (\ell + 1/2)^2/R^2 \quad (17)$$

and the lower limits of integration are the respective classical turning points. At high energy $k^2 \gg |2\mu(U - iW)|$,

$$\eta_\ell \approx -\frac{\mu}{k} \int_0^{\infty} \{U[(p^2 + z^2)^{1/2}] - iW[(p^2 + z^2)^{1/2}]\} dz \quad (18)$$

where

$$p = (\ell + 1/2)/k \text{ and } z^2 = R^2 - p^2$$

Using Eq. (18) in Eq. (14) and changing the sum over ℓ to an integral, we obtain

$$\sigma_A = 2\pi \int_0^{\infty} p \left\{ 1 - \exp \left[-\frac{4p}{k} \int_0^{\infty} W[(p^2 + z^2)^{1/2}] dz \right] \right\} dp \quad (19)$$

Note that at high energy the real part of the interaction is unimportant because contributions from high partial waves dominate the cross section. At energies of interest several thousand partial waves are important.

The imaginary part of the potential is caused by the finite lifetime of the state because of the possibility of transitions to other electronic states during the collision. In general we expect that the strength of the coupling between the molecular states will be on the order of the average spacing between these states. Thus, since W equals half the width of the state we can calculate it using the Golden Rule,

$$W = \Gamma/2 = \pi \rho_f |V_{if}|^2 \quad (20)$$

where ρ_f is the density of final states.

Assuming then that on the average $|V_{if}| \approx 1/\rho_f$, we have

$$W \approx \pi/\rho_f \quad (21)$$

where actually ρ_f should include only final states coupled to the initial state. Finally, at large separations the coupling must vanish so that two obvious functional forms for W suggest themselves, namely

$$W_1(R) = \begin{cases} \pi/\rho_f & R < R_0 \\ 0 & R > R_0 \end{cases} \quad (22a)$$

and

$$W_2(R) = \pi/\rho_f e^{-\alpha R^2}, \quad \alpha = \frac{\ln 2}{R_0^2} \quad (22b)$$

Figure 10-34 shows these potentials. The Gaussian falls to half its maximum value at $R = R_0$. Putting these into Eq. (19) we obtain the absorption cross sections

$$\sigma_A^1(E) = \frac{8\pi\mu R_0^3 W_0}{3k} \quad (23a)$$

and

$$\begin{aligned} \sigma_A^2(E) &= \left(\frac{\pi}{\alpha} \right)^{3/2} \frac{2\mu W_0}{k} \\ &= \left(\frac{\pi}{\ln 2} \right)^{3/2} \frac{2\mu R_0^3 W_0}{k} \end{aligned} \quad (23b)$$

In both cases we retained only the leading term in an expansion in powers of $\mu W_0 R_0/k$, where W_0 equals π/ρ_f . In both cases the cross section is inversely proportional to the collision velocity and to the third power of the radius of the interaction. The cross section also depends linearly on the strength of the coupling. In fact if the amount of energy transferred from kinetic to electronic excitation is of interest, we use a somewhat different derivation of the cross-section formula. Following Miller et al.⁴³ we obtain the following formula for the absorption cross section, again in the high-energy limit:

$$\sigma_A(E) = \frac{8\pi\mu}{k} \int_{\infty}^{\infty} dR R^2 W(R) [1 - U(R)/E]^{1/2} \quad (24)$$

Here the lower limit of integration is the classical turning point for the real part of the potential, i. e., the zero of the expression in the brackets. We also derive the differential cross section for producing electronic excitation in the range ϵ and $\epsilon + d\epsilon$. We assume that a transition occurring at R_t transfers $\epsilon = U(R_t)$ to electronic excitation (see Fig. 10-34), then

$$\sigma(\epsilon, E) = \frac{8\pi\mu}{k} |U'(R_t)|^{-1} W(R_t) R_t^2 [1 - U(R_t)/E]^{1/2} \quad (25)$$

gives the energy distribution of electronic excitation.

We can obtain a qualitative understanding for the cross sections for electronic excitation from the above formulas. To produce quantitative results we must assume some values for the parameters in the potentials. For collision energies above 1 eV the repulsive part of the real potential is all that is important. Again we chose two different forms:

$$U_1(R) = A/R^{12}, A = 7.7382 \times 10^9 \text{ a.u.} - a_0^{12}; \quad (26a)$$

and

$$U_2(R) = ae^{-bR}, a = 2.92 \times 10^3 \text{ a. u.}$$

$$\text{and } b = 1.865 a_0^{-1} \quad (26b)$$

Preston⁴¹ determined the first of these by equating the hard sphere and momentum-transfer cross sections at a U-U scattering energy of 1 eV. The second real potential comes from a Thomas-Fermi-Dirac

calculation.⁴⁰ For the range of the interaction we choose $R_0 = 4 \text{ \AA}$, which is somewhat less than twice the value for the 7s orbital radius, and $W_0 = 1/2000 \text{ eV}$ for the strength of the coupling.

Figure 10-34 shows the results for both real parts of the potential and only for the imaginary part $W_2(R)$. We evaluated numerically the integral in Eq. (24) for the absorption cross section. Figure 10-35 shows the resulting cross sections for a range of collision energies. We also display the momentum-transfer cross sections from the Gordon-Kim and Thomas-Fermi-Dirac potentials of Ref. 39. The absorption cross sections decrease as $E^{-1/2}$ as expected. The momentum-transfer cross section for a potential of the form U_1 decreases much more slowly, $E^{-1/6}$, as a function of collision energy.

We are also interested in the amount of electronic excitation caused by a collision. Equation (25) can be rewritten as

$$\sigma(\epsilon, E) = \frac{4\pi W_0 (2\mu)^{1/2}}{bE\epsilon} R_t^2 e^{-\alpha R_t^2} [e - \epsilon]^{1/2}, \quad (27)$$

where $R_t = (1/b) \ln(\epsilon/a)$. As expected, the distribution peaks strongly at low-energy transfer and we find that it decreases rapidly with increasing ϵ . We obtain an estimate of the cross section for ionization by integrating this differential cross section over all values of the energy transfer greater than 6.2 eV, the ionization potential of uranium:

$$\sigma_I(E) = \int_{6.2}^E d\epsilon \sigma(\epsilon, E) \quad (28)$$

We find that for both choices for the real part of the potential [$U_1(R) = A/R^{12}$ and $U_2(R) = ae^{-bR}$] the ionization cross sections are on the order of a few $\text{\AA}^2$ and vary slowly with collision energy.

Fig. 10-34. Schematic representation of real and imaginary parts of interaction potential.

Fig. 10-35. Cross sections in a_0^2 for E-T and momentum transfer as function of collision energy.

In summary, these calculations have shown that for collision energies of a few eV:

- The E-T cross sections are of the same order of magnitude as the available momentum-transfer cross sections but decrease more rapidly as a function of collision energy than the available momentum-transfer cross sections.
- The E-T collision rate is approximately independent of energy and therefore the number of collisions suffered by an atom over a fixed distance is inversely proportional to its velocity.
- The average energy transferred from the translational to electronic excitation is only a few percent of the total collision energy.
- The collision rate is approximately gas kinetic.

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Laser-Induced Fluorescence Studies of U II Produced by Photoionization of Uranium

Introduction. Over the past several years we have made substantial progress in understanding the process of photoionization of atoms by lasers. In general,

when an atom is photoionized the resulting ion may be singly or multiply charged. For elements having close-lying neighboring ion states, photoionization may also produce ions in various electronic levels.

Traditionally we can determine the initial electronic-state distribution (or branching ratio) of photon-produced ions by photoelectron spectroscopy.⁴⁴⁻⁴⁶ By studying the kinetic energy of the ejected photoelectrons we can obtain detailed knowledge regarding the binding energies of electrons in various states and the probability of ejecting electrons from these specific states. The angular distribution of the ejected photoelectrons depends sensitively on the angular momentum of the electron and on the degree of polarization of the incident radiation. Normally the photoelectron angular distribution is not isotropic. Therefore it is difficult to deduce accurate information on the initial-state distribution of photoions based on the photoelectron spectrum. In the experiments reported here we directly measured the initial electronic-state distribution of photon-produced uranium ions using laser-induced fluorescence (LIF).⁴⁷ From the relative intensities of LIF signals of U and U^+ , we determined the fraction of U atoms excited and ionized. We then applied optical U^+

detection to monitor the subsequent motion of photoions in the laser-produced uranium plasma.

The spectra of atomic U I (U) and U II (U^+) are extremely complex. The neutral uranium atom has a radon core with six optical electrons. The ground state belongs to the $5f^3 6d 7s^2$ configuration and in LS coupling is approximately 5L_6 . The ground state of the uranium ion U^+ belongs to the configuration $5f^3 7s^2$ and is approximately described as $^4I_{9/2}$. The lowest excited state of U^+ has a configuration of $5f^3 6d 7s$, $^6L_{11/2}$ and is about 289 cm^{-1} above the first ionization limit of uranium. The parities for these ground and metastable excited states are all odd so that optical transitions between these two states are therefore parity forbidden.

We measured the electronic-state distribution (or branching ratio) of U^+ ions ($^4I_{9/2}$, $^6L_{11/2}$) generated by multiphoton ionization of uranium (5L_6). In this article we give a brief account of the experimental methods and findings on radiative lifetimes, photoionization branching ratio, and photoionization cross section of uranium. In principle we can apply the experimental techniques described here to determine the

photoionization branching ratio for other elements of interest where measurements using traditional photoelectron spectroscopic methods would be exceedingly difficult.

Experimental Conditions. Figure 10-36 is a simplified schematic diagram of the experimental setup used for U^+ generation and for the optical- and charged-particle detection of U^+ . Uranium atoms from a thermal source diffuse into a LIF chamber. A broadband pulsed dye laser (tuned to $5915 \text{ \AA}$) excites and volumetrically ionizes a significant fraction of the U atoms. Shortly after photoionization of the uranium, a second pulsed dye laser fires and selectively excites the photoion into an even-parity state near the 2 to 4 eV region. We send the corresponding LIF signals, detected by an RCA 8575 photomultiplier, to a boxcar integrator for signal averaging. A strip-chart recorder traces the output of the boxcar integrator. For radiative-lifetime measurements we digitize the complete waveform of the fluorescence-signal with a Biomation 6500 waveform recorder having 2 ns resolution. We average the digitized fluorescence signal with a Nicolet model 1170 signal averager.

Fig. 10-36. Experimental setup for simultaneous optical and charged-particle detection of U^+ .

The uranium vapor issues from a high-temperature thermal oven⁴⁸ where a closely woven tungsten-mesh filament heats metallic uranium or URe₂ alloy in a thick-walled tungsten crucible. Numerous layers of tungsten heat shields surround the crucible. We have designed the oven to operate at 2600 K or higher. We measured the oven temperature by sighting with an automatic optical pyrometer located directly below the source. We have calibrated the pyrometer to the IPTs 1968 temperature scale by comparison with an NBS reference lamp. We determined window absorption and emissivity correction factors before the experiments.

Before beginning a series of LIF experiments we evacuate the oven and the optical- and charged-particle detection chambers to about 10^{-8} Torr. We then adjust and maintain the thermal oven at some specific temperature, typically 2400 K. The vapor pressure of uranium at this temperature is about 10^{-3} Torr. At the oven temperature of 2400 K approximately 48% of the uranium atoms in the vapor are in the ground state, 28% are in the 620-cm^{-1} level, 5.7% are in the 3801-cm^{-1} level, and roughly 18.3% are in higher-lying levels. We first tune the photoionization laser (Quanta Ray YAG-pumped dye laser) to about $5915\text{ \AA}$ for excitation and ionization of ground-state uranium atoms. The typical laser performance at $5915\text{ \AA}$ was about 50 mJ/pulse, 5 ns duration, and 5 GHz in line

width. We then tune a N₂-laser-pumped dye laser (oscillator–amplifier configuration) to a predetermined wavelength to selectively excite uranium atoms or ions for subsequent fluorescence measurements.^{48,49} The beam diameters for the ionizing and probe lasers are 4 and 1 mm, respectively. The fluence of the photoionization laser at the crossing of uranium and laser beams is about 50 mJ/cm^2 , which is about 3 times the saturation fluence for ionization. We constantly monitor the flux of uranium ions with a charged-particle detector (channeltron) adjacent to the crossing. We send the corresponding ion signals to a separate boxcar integrator for signal averaging. A strip-chart recorder also traces the output of this boxcar integrator.

Radiative-Lifetime Measurements for U⁺

Figure 10-37 shows the time history of a typical ion fluorescence signal at $3859.6\text{ \AA}$. By examining a semilog plot of the fluorescence intensity vs time, we found that the decay signal is best fit with a single exponential. Thus our interpretation is that the fluorescence signal originated from a single excited state. To verify that the fluorescence does not originate in decay from an excited neutral atom, we installed a pair of ion-deflection plates in the fluorescence chamber. The signal disappeared as soon as an electrical field was applied to the system.

We found the radiative lifetime for the $26\,191\text{-cm}^{-1}$ state (a level with configuration of f^3dp , $^6M_{13/2}$ excited by the $3859.6\text{ \AA}$ laser) to be $20 \pm 5\text{ ns}$. The configuration for the lower level is f^3ds , $^6L_{11/2}$, which is 289 cm^{-1} above the ground state. We have also measured the radiative lifetime for the $24\,684\text{-cm}^{-1}$ state (a level with configuration of f^3sp , $^6I_{9/2}$ excited by the $4050.0\text{-\AA}$ laser from ground state f^3s^2 , $^4I_{9/2}$). The radiative lifetime for this excited state is $44 \pm 5\text{ ns}$. The branching ratios for radiative decay from the $24\,684\text{-}$ and $26\,191\text{-cm}^{-1}$ levels measured in Ref. 50 are 0.383 for the $4050.0\text{-\AA}$ transition and 0.403 for the $2859.6\text{-\AA}$ transition, respectively. Based on these data we calculated the gA values by the relationship $gA = gBR/\tau$. We give the results in Table 10-6. For comparison we also list the gA values reported in Ref. 51.

Initial-State Distribution of Photoions. At present we have very little information about the initial-state population of photoions

Fig. 10-37. Fluorescence decay signal of $26\,191\text{-cm}^{-1}$ state of U⁺.

Table 10-6. Uranium⁺ oscillator strengths.

Transition ($\text{\AA}$)	Energy levels (cm^{-1})	Radiative lifetime $\tau(\text{ns})$	BR ^a	$gA(10^8\text{ s}^{-1})$	
				This work	Corliss ^b
4050.0	0 ($J = 9/2$)	44 ± 5	0.383	0.86	0.42
	24 684 ($J = 9/2$)				
3859.6	289 ($J = 11/2$)	20 ± 5	0.403	2.8	1.1
	26 191 ($J = 13/2$)				

^aRef. 50.

^bRef. 51.

generated by this process. Figure 10-38 shows a simplified schematic diagram of energy levels of uranium. The excited even-parity states of uranium at $50\,701\text{ cm}^{-1}$ are about 743 cm^{-1} above the first ionization limit.⁵² Thus ionization of uranium in this region could produce both ground and metastable excited (289 cm^{-1}) ions. Since the configurations for U I states in this region are not well known, we cannot make any meaningful predictions about the initial branching ratio of this photoionization process.

In this work we attempted to determine the relative population of U^+ by monitoring the intensities of LIF signals using a narrow-bandpass filter and boxcar integrator. We set the boxcar gate to switch on a few nanoseconds after the probe-laser pulse sweeps through the radiation zone. We kept the gate width at about 30 ns. In theory we can relate the ground state ion density and the 289-cm^{-1} U II-level density by the following equation

$$\frac{U_g^+}{U_{289}^+ \text{ cm}^{-1}} = \frac{I_g(t)}{I_{289}(t)} \frac{\tau_g}{\tau_{289}} \frac{T_{289} \text{ BR}_{289} (h\nu)_{289}^{\text{FL}}}{T_g \text{ BR}_g (h\nu)_g^{\text{FL}}} \times \frac{\sigma_{289}}{\sigma_g} \frac{J_{289}}{J_g} \exp\left(\frac{t}{\tau_g} - \frac{t}{\tau_{289}}\right), \quad (29)$$

where

$I_g(t), I_{289}(t)$ = fluorescence intensities of transitions initiated from ground or 289-cm^{-1} level and monitored by the photomultiplier and boxcar integrator at time t after the probe-laser pulse swept through the uranium vapor.

τ_g, τ_{289} = radiative lifetimes of the upper ion states excited by the probe laser from the ground or 289-cm^{-1} level.

σ_g, σ_{289} = absorption cross sections for the probe transitions; we used $4050.0\text{ \AA}$ for probing the ground-state ion (σ_g) and $3859\text{ \AA}$ for probing the 289-cm^{-1} ion (σ_{289}).

J_g, J_{289} = probe-laser fluence in $\mu\text{J}/\text{cm}^2$ for excitation of ground and 289-cm^{-1} level.

t = boxcar-gate switch-on time after the probe-laser pulse swept through the uranium vapor.

T_g, T_{289} = transmission of the narrow-bandpass filter and collection optics used for monitoring LIF signals initiated from ground or 289-cm^{-1} level.

$\text{BR}_g, \text{BR}_{289}$ = fluorescence branching ratios

from the excited states initiated from ground or 289-cm^{-1} level.
 $(h\nu)_g^{\text{FL}}, (h\nu)_{289}^{\text{FL}}$ = photon energy of fluorescence light initiated from ground or 289-cm^{-1} level.

Figure 10-39 shows the fluorescence signals observed in the $4050\text{-to-}4150\text{-\AA}$ bandpass region as the probe laser is tuned to 4050 and $3859\text{ \AA}$, respectively. We kept the fluence of the ionizing laser ($5915\text{ \AA}$) at $50\text{ mJ}/\text{cm}^2$. Radiative decay from the $26\,191\text{-cm}^{-1}$ ion state branches⁵⁰ approximately 38.5% into the 1749 cm^{-1} level (emission at $4090\text{ \AA}$). Radiative decay from the $24\,684\text{-cm}^{-1}$ ion state branches⁵⁰ approximately 38.3% into the ground state (emission at $4050\text{ \AA}$) and 14.5% into the 289-cm^{-1} level (emission at $4098\text{ \AA}$).

Based on this information and the transmission of our fluorescence-detection

Fig. 10-38. Excitation and probe schemes used to obtain initial-state distribution of U^+ .

Fig. 10-39. Fluorescence signals of U^+ ions observed in the $4050\text{-to-}4150\text{-\AA}$ bandpass region as the probe laser is tuned to 4050 and $3859\text{ \AA}$, respectively.

system, we calculated the density ratio ($U^+/U_{289}^+ \text{ cm}^{-1}$) for our experiment. The ratio is approximately 1:1. Thus the initial-state distribution of U^+ produced by the three-step process using 5915 Å is approximately 50% in the ground state and 50% in the first metastable excited state at 289 cm^{-1} :

$$U_g (^5L_6, f^3ds^2) \times \frac{3h\nu}{5915\text{\AA}} \begin{cases} e + U_g^+ (^4I_{9/2}, f^3s^2) & \sim 50\% \\ e + U_{289}^+ \text{ cm}^{-1} (^6L_{11/2}, f^3ds) & \sim 50\% \end{cases} \quad (30)$$

We have not determined whether the initial-state distribution of U^+ changes with the fluence or line width of the ionizing laser.

At present we cannot draw any valid conclusions about the configurations of these excited states at 33 801 and 50 701 cm^{-1} and hence about the ionization mechanism associated with the 5915-Å process. We believe that additional information on the configurations of these states might be obtained by studying the Rydberg progressions^{52,53} in uranium using multistep photoionization spectroscopy.

Degree of Ionization and Fraction of U Atoms Removed by the Stepwise Photoexcitation Process. We have also performed two sets of LIF experiments to

Table 10-7. Lifetimes and cross sections for the 5915-Å process.

Steps (cm^{-1} to cm^{-1})	Radiative lifetimes $\tau(\text{ns})$	Cross sections $\sigma(\text{cm}^2)$
0 to 16 900	205 ± 8^a	5.4×10^{-13}
16 900 to 33 801	534 ± 19	$>> 10^{-15}$ (not measured)
33 801 to 50 701	$<< 1$ (not measured)	4.7×10^{-17}

^aRef. 55.

Fig. 10-40. Fluorescence signals of U^+ and U observed in the 4050-to-4150-Å bandpass region as the probe laser is tuned to 4061.3 and 4050.0 Å, respectively.

determine the fraction of ground state U atoms excited and ionized by the three-step photoexcitation process using a 5915-Å laser. In the first set of experiments we determined the density ratio ($U^+/U_{620} \text{ cm}^{-1}$) by monitoring the relative intensities of U^+ and U_{620} LIF signals. Figure 10-40 shows the fluorescence signals observed in the 4050- to 4150-Å region as we tune the probe laser to 4061.3 and 4050.0 Å, respectively.

Fluorescence at 4061.3 Å resulted from excitation and subsequent emission from the 25 235- cm^{-1} level (excited from the 620- cm^{-1} level). Fluorescence from the 25 235- cm^{-1} state branches⁵⁰ approximately 36.8% into the 620- cm^{-1} level. The radiative lifetime measured with LIF is 113 ± 10 ns. Using these data we calculated the density ratio (U^+/U_{620}) at approximately 0.59. We kept the U source temperature at 2400 K throughout the LIF experiment. At this temperature the populations of neutral atomic uranium at the ground- and 620- cm^{-1} levels are approximately 48% and 28%, respectively. Using this information we were able to make further calculations for the density ratio (U_{total}^+/U_g). At this temperature and ionizing laser fluence of 50 mJ/cm^2 the density ratio is approximately 0.65. Therefore during the photoionization process the 5915-Å laser ionized about 65% of the ground-state U atoms.

In a second set of experiments we measured the depletion of ground state U atoms during the photoexcitation process by monitoring the 4362.1-Å LIF intensity before and after the photoionization-laser pulse has swept through the uranium vapor. The gA value for the 4362.1-Å transition⁵⁴ (ground to 22 919- cm^{-1} state) is $1.6 \times 10^8 \text{ s}^{-1}$. We have purposely kept the probe-laser fluence at about 3 times the saturation level. We found that when the source temperature was 2400 K and photoionization-laser fluence was 50 mJ/cm^2 , the photoionization pulse removes approximately 69% of the ground state U atoms. Therefore we conclude that during our photoionization experiment using 5915-Å light, over 90% of the U atoms pumped from the ground state were converted to U^+ .

Photoionization Cross Section and Excited-State Lifetime Measurements. To further our understanding of the mechanism of the three-step photoionization process using 5915-Å photons, we measured both the

excited-state lifetimes and excitation cross sections. Earlier Miron et al.⁵⁵ measured and reported the lifetime and cross section for the first 5915-Å transition (0 to 16 900 cm⁻¹). We measured the radiative lifetime for the 33 801-cm⁻¹ state (second excited state pumped by the 5915-Å laser) by monitoring the fluorescence of excited U in the 4500-to-5800-Å region, using the setup shown in Fig. 10-36. We found the radiative lifetime for the 33 801 cm⁻¹ state is 534 ± 19 ns, in good agreement with the measurement of 560 ns.⁵⁶

We determined the photoionization cross section for the last excitation step (33 801 to 50 701 cm⁻¹) by measuring the relative U⁺ yield as a function of 5915-Å laser fluence. Figure 10-36 shows the experimental setup for the optical and charged-particle detection of U⁺.

Figure 10-41 shows a plot of relative U⁺-ion signal (monitored by a channeltron detector) vs 5915-Å laser fluence. The curve shows a clear saturation as the laser fluence increases beyond 15 mJ/cm². The effective fluence required for saturation is roughly 6.8 mJ/cm², in close agreement with the results obtained in a separate experiment using a multistep-photoionization technique.⁵⁷ We calculated the effective photoionization cross section with the formula

$$\sigma_{\text{ionization}} = \frac{h\nu}{J_{\text{saturation}}} = 4.7 \times 10^{-17} \text{ cm}^2 \quad (31)$$

Table 10-7 summarizes the cross sections and lifetimes for the excitation processes using a 5915-Å laser. We do not know the cross section for the second step (16 900 to 33 801 cm⁻¹). However, it is reasonable to assume that it must be at least 2 orders of magnitude larger than that for the photoionization process.

Ambipolar Expansion of U⁺ in a Laser-Produced Plasma. As a logical extension of this study we applied the optical U⁺-detection technique to monitor the motion of U⁺ in a laser-induced plasma. A pulsed dye laser tuned to 5915 Å volumetrically ionizes a significant fraction of uranium atoms. Shortly after the ionization laser swept through the vapor, we triggered a probe laser tuned to the 4050-Å ion-resonance line. We monitored and recorded the corresponding LIF signal. As we

described earlier, the beam diameters for the ionizing and probe lasers are 4 and 1 mm, respectively. Figure 10-42 shows the time history of the LIF signal as a function of t , which is the time interval between the ionizing and probe lasers. The flow velocity of U can be calculated from its source temperature. At 2400 K it is roughly 3.5×10^4 cm/s. We found that the density of U⁺ decreases more rapidly than that calculated by the flow velocity. We believe this more rapid decrease of U⁺ is caused by ambipolar expansion of U⁺ and electrons in a field-free plasma.

The ambipolar expansion velocity can be calculated by

$$v_{\text{a.p.}} = \left(\frac{\gamma k T_e}{M} \right)^{1/2}, \quad (32)$$

where T_e is the electron temperature of the plasma, γ is the heat-capacity ratio, and M is the mass of U⁺. For our experimental conditions the ambipolar expansion velocity is roughly 2.1×10^4 cm/s. The estimated plasma response time is roughly 10^{-9} s or shorter. For comparison we plotted the

Fig. 10-41. Relative U⁺ yield vs ionizing laser fluence.

Fig. 10-42. Time history of U⁺ in plasma.

calculated time history for U^+ assuming ambipolar ion expansion (Fig. 10-42).

Conclusions. We have demonstrated that we can directly measure the branching ratio for the atomic photoionization process using the LIF technique. From the relative intensities of LIF signals of ions and neutrals, we can also determine the fraction of U atoms excited and ionized by the stepwise photoexcitation process.

Our experimental data clearly demonstrate that photoionization of uranium through the $50\,701\text{ cm}^{-1}$ even-parity level using a broadband pulsed dye laser of 5915 Å produces roughly 50% ground-state ions and 50% metastable excited (289 cm^{-1}) ions. The effective photoionization cross section is about $4.7 \times 10^{-17}\text{ cm}^2$. The subsequent dispersion of U II after photoionization appears to be caused by ambipolar expansion of ions and electrons in the plasma. The gA values for the 4050-Å and 3859-Å transitions of U II are 0.86×10^8 and $2.8 \times 10^8\text{ s}^{-1}$, respectively.

In summary the LIF technique described here is a potential practical approach for determining transient and spatial ion densities in electrical discharges, in sparks, in coronas, and in laser-produced plasma. We anticipate that LIF of ions, carried out with continuously tunable laser sources, will be a powerful method for studying

photoionization and electron correlation in atoms.

Author: H. L. Chen

Major Contributor: C. V. Borzileri

Spectroscopic Data Acquisition

During the past year the automation of data acquisition and analysis has had a major impact on our experimental investigation of photoionization physics. These experiments usually require us to make three classes of measurements: uranium-vapor density, a set of laser parameters, and photoproduced ion density. We now make these measurements under computer control, allowing us to gather much more data than before, to systematically document the relations among the data, and to store the data in a form suitable for automated analysis.

For example, we determine neutral-uranium vapor density by measuring the absorption in the vapor of a stable cw dye laser tuned to an absorption line of known absorption strength. Because of the velocity distribution of the atoms this absorption occurs over a range of frequencies. Our equipment scans the cw laser frequency

Fig. 10-43. Measurement of photoproduced ions at three laser powers.

across the absorption profile, measures the absorption at each of 1024 points, stores these data for future analysis, and also performs a quick analysis for immediate output. Thus we measure the neutral-uranium number density to an accuracy of 10% once every few seconds.

As another example of our automated data acquisition we show in Fig. 10-43 the measurement of photoproduced ions as we vary the laser powers. Each dot represents a simultaneous measurement of ions and power level for each of the three lasers involved. We make these measurements 10 times per second and store them on a magnetic disk. The solid lines are computer-produced averages of the data.

Rapid acquisition and analysis of data in the laboratory gives us increased flexibility in varying experimental parameters in real time. The laser and vapor-generating system are large and costly to operate. Our new computer aids have greatly increased the productivity of these systems.

Author: M. A. Johnson

Beam-Propagation Theory

As a result of our theoretical and experimental studies over the last 5 years, we understand the basic detailed dynamics of coherent excitation and have a realistic model of the atom-laser interaction. We are confident that, given a prescribed laser pulse, we can predict the atomic excitation and ionization behavior. However, the atomic response to the field in turn modifies the field. The laser electric field drives atomic excitation, which creates an induced dipole moment oscillating at the population-inversion (Rabi) frequency. This dipole then acts as the source of a secondary radiation field. As the laser pulse propagates through the medium the secondary fields may significantly alter the initial pulse by delay, amplitude reshaping, or carrier-frequency shifting. We must consider these changes if we are to maintain the efficiency and selectivity of the atomic-vapor laser isotope separation (AVLIS) process. Our studies of laser propagation in near-resonant media aim to provide an understanding of

propagation effects and will ultimately influence our process design.

The fundamental equation of laser propagation is the inhomogeneous wave equation derived from Maxwell's equations in the presence of polarizable material:

$$\left[\frac{\partial^2}{\partial z^2} + \nabla_{\perp}^2 - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \right] \mathbf{E} = \frac{4\pi}{c^2} \frac{\partial^2}{\partial t^2} \mathbf{P} \quad (33)$$

Here $\mathbf{E}$ is the total time- and space-dependent electric field for a pulse propagating along the Z axis and $\nabla_{\perp}^2$ is the transverse Laplacian.

The polarization $\mathbf{P}$ is the macroscopic average (over atomic orientations, thermal Doppler shifts, etc.) of induced atomic-dipole moments. Theory yields expressions for $\mathbf{P}$ under various simplifying conditions.

The simplest situation occurs when atomic-relaxation processes bring weak excitation into temporary steady-state equilibrium. In this classical case $\mathbf{P}$ is proportional to $\mathbf{E}$ and we recover analytically the index of refraction and absorption coefficient of classical optics.

A more refined atomic model, the adiabatic following model (AFM), applies when the pulse amplitudes change sufficiently slowly. We can use it to treat population inversions and, like the classical complex refractive index, it provides us with an analytic connection between $\mathbf{P}$ and $\mathbf{E}$.

Generally we must solve dynamical equations, either the time-dependent Schrödinger equation or multilevel Bloch equations, to find the time-dependent behavior of $\mathbf{P}$ for a given $\mathbf{E}$.

Each of these (and other) approximations for the relationship between $\mathbf{P}$ and $\mathbf{E}$ allows an appropriate level of detail in treating pulse propagation.

Propagation effects fall into two broad interconnected classes. Spatial effects, caused by diffraction (present even in the absence of vapor) and refraction (enhanced dramatically by resonant excitation), alter the volume of illuminated vapor and the intensity of illumination. Temporal effects modify the pulse shape and frequency characteristics and hence alter the atomic-excitation dynamics.

In propagation through uranium vapor the dominant component is ^{238}U , for which the pulse is not resonant. Little excitation

occurs and we can therefore apply classical approximations (absorption and refraction). The desired component is ^{235}U , present in small concentration yet resonant with the pulse-carrier frequency. We anticipate appreciable excitation of ^{235}U and hence we must deal with coupled Maxwell-Schrödinger or Maxwell-Bloch equations. We have developed codes to treat each of these excitations in various approximations. For details, see Ref. 58.

As a first step toward understanding the off-resonant effects of propagation through ^{238}U we have studied off-resonant propagation in the simpler atomic system of sodium. In the following section we describe our experiment and discuss theoretical and computational modeling of such propagation.

Author: B. W. Shore

Near-Resonant Propagation in Sodium

During the last year we initiated a series of theoretical and experimental studies on near-resonant pulse propagation in atomic

sodium to help us model propagation phenomena for atomic uranium in our atomic-vapor laser isotope separation (AVLIS) process.

Propagation of low-intensity resonant or near-resonant light, too weak to initiate appreciable atomic excitation, follows a behavior that we can predict from a linearized approximation to the inhomogeneous wave equation. In the linear approximation we treat each Fourier component of the pulse independent of the others. Thus, although pulse breakup may occur within the approximation, linear theory admits no alteration of the pulse spectrum. Applied to a frequency-modulated (FM) pulse, linear theory predicts a conversion from frequency to amplitude modulation (AM), leading to a train of high-peak-intensity pulses.

At high intensities, as excitation becomes appreciable, linear theory fails. However, if the Rabi frequency is much less than the off-resonance detuning, we can treat the pulse-propagation problem by introducing an intensity-dependent susceptibility derived from the adiabatic following model (AFM). In the low-intensity limit this approximation

Fig. 10-44. Schematic of experimental apparatus used to study nonresonant pulse propagation in atomic sodium.

works well for sufficiently small propagation distances. It also works well as long as pulse bandwidths remain smaller than the limits set by pulse duration. Although the AFM fails for very large propagation depths of dispersive pulse reshaping, it nonetheless can provide us with valuable scaling laws.

We have investigated both the linear and nonlinear propagation regimes with experiments on atomic sodium. Figure 10-44

shows the elements of our experimental apparatus. We use an argon-ion laser to pump a ring dye laser, whose tunable output is passed through an optical train into an oven that produces sodium vapor. We use an assortment of frequency diagnostics (interferometers, sodium lamps, etc.) and temporal diagnostics to monitor the characteristics of the input and output radiation.

Fig. 10-45. Low-input-intensity pulse reshaping for $\delta = 4.2$, $\nu_m = 0.2$ GHz, $\text{nfL} = 5.44 \times 10^{14} \text{ cm}^2$, and $\Delta\nu = \nu_1 - \nu_0 = -6.5 \text{ GHz}$.

Pulse shape

Pulse spectra

(a) Input (offset to show noise level).

(b) Output

$\delta = 4.2$, $\nu_m = 200 \text{ MHz}$, $\Delta\nu = \nu_1 - \nu_0 = -6.5 \text{ GHz}$, $\text{nfL} = 5.44 \times 10^{14} \text{ cm}^2$ and $I = 350 \text{ W/cm}^2$. Sodium D_1 line is on right-hand side of frequency spectrum.

Fig. 10-46. Spectral and temporal profiles of near-resonant (D_1 line) input and transmitted light in Na vapor at high incident intensity $I = 350 \text{ W/cm}^2$.

To investigate the linear theory we employed an attenuated laser beam, detuned from the sodium D₁-line resonance by $\Delta\nu = -6.5$ GHz, modulated at a frequency $\nu_m = 200$ MHz, with modulation index $\delta = 4.2$. We obtained a value of the product of nfL (where n = vapor density, f = transition oscillator strength, and L = vapor path length) by measuring pulse delay in the vapor. Both linear theory and the AFM predict delay times proportional to $nfL/(\Delta\nu)^2$.

After passage through the vapor the FM cw pulse becomes AM, as seen in Fig. 10-45(a). These results agree quantitatively with predictions of linear theory, smoothed with a 400-ns time constant, as shown in Fig. 10-45(b). By removing the beam

attenuator and focusing the laser beam we achieved intensities (~ 350 W/cm²) that took us well beyond the limits of linear theory. In this regime the pulse spectrum undergoes marked changes. Sidebands appear and the overall spectral bandwidth increases, as seen from the experimental results shown in Fig. 10-46. To model this regime we employed a plane-wave treatment of coupled Maxwell-Bloch equations, in the slowly varying envelope-rotating-wave approximations. Figure 10-47 shows the predictions of this code.

A self-phase modulation causes the increased bandwidth of the propagating light field. Since in our experiment the AM is periodic, the nonlinear self-phase modulation occurs as sidebands spaced by

Fig. 10-47. Two-level Maxwell-Bloch code predictions for near-resonant propagation in Na vapor of FM input pulse using vapor condition in Fig. 10-46.

the modulation frequency with a spectral envelope determined by the frequency chirp occurring for any single pulse in the pulse train. Thus the physical significance lies not in the sidebands themselves but in the envelope of the sideband. We predicted the spectral envelope using a Bloch-equation code in Fig. 10-47(b) and it compares very well with our experimental data in Fig. 10-46(b).

The agreement between our experimental and computational results for propagation in sodium of FM off-resonant light provides a significant step toward modeling the more-complicated propagation phenomena in uranium.

Authors: J. A. Paisner and L. A. Hackel

Computational Facilities

Introduction

Our process facilities and experiments demand coordinated control of numerous electrical, optical, and mechanical subsystems as well as coordinated monitoring of detector responses. The growing complexity of these facilities and experiments provides a strong impetus for us to develop laboratory computational facilities. This section describes our efforts to design and implement a system of hardware and software units that will form a master facility for our LIS program.

Author: B. W. Shore

Laboratory Computer Systems

Introduction. Since 1978 the Laser Isotope Separation (LIS) program has employed microcomputers for acquisition of laser and plasma-physics data. This year we acquired our first minicomputer, created a flexible data-acquisition program, and formed a computer software group.

We now use 17 Digital Equipment Corporation LSI-11 microcomputers and two Hewlett-Packard HP1000 minicomputers for data acquisition, data

analysis, and program development. Of the LSI-11 machines, two are LSI-11/23 processors (which have many of the features of minicomputers), while the remainder are LSI-11/2 processors. Our HP1000 computers are linked by fiber optics and perform as a team while our other machines act alone. The computers are assigned to the programs listed in Table 10-8. We also have begun to incorporate Zilog Z8000 microprocessors into dedicated acquisition peripherals.

Our data-acquisition program has evolved from a specialized routine servicing a fixed set of instruments to a generalized set of routines that combine to operate a growing number of instruments. We can also construct a tailored experimental procedure from existing software modules. Such tailored routines are changed for each experiment.

- The major tasks of our software group are
- Creation and maintenance of programming languages, libraries, and utilities.
 - Maintenance and modification of data-acquisition software.
 - Creation of general and specific data display and analysis programs.

Our major computational-support accomplishments are

- Functional modularization of our data-acquisition program.
- Customization of experimental procedures in our standard program.
- Activation of the Mars computer system.
- Creation of dedicated Z8000 acquisition peripherals.
- Commencement of the Mars software design.
- Creation of computer-automated measurement and control (CAMAC) module test programs.

Control-Systems Development	Plasma Simulation
Dye Amplifier Development	Process Laboratory
Electron-Physics Laboratory	Regulis
Extractor Test Facility	SPP-II
Mars	Venus
Minerva	Vulcan

Table 10-8. Programs to which data-acquisition computers are assigned.

- Creation of a fast plotting program.
- Activation of a single-user program-development computer.

Many of these items appear in more detail elsewhere in this article. Our major future activities are

- Create new experiment definition methods.
- Create new data-storage algorithms for raw and processed data.
- Improve the program-development environment.

Table 10-10. CAMAC modules.

Manufacturer and model	Characteristics
LeCroy 2256AS	Digitizer; 1024 points; 20-MHz maximum digitizing rate
Jorway 237	Digitizer; 1024 points; 20-kHz maximum digitizing rate
Kinetic Systems 3361	Stepper motor controller
Standard Engineering E260	Digital-analog converter; 8 channels; 12-bit resolution
LeCroy 8100	Programmable differential amplifier; 2 channels; 0.1- to 50-V input; 0.2 to 100 gain
Joerger FT	Signal fan out; 2-channel dual input; 4 outputs per channel
Kinetic Systems 3061	General-purpose 16-bit input-output register.

Fig. 10-48. Standard LIS laboratory LSI-11 configuration. ▼

connected to the computer for lists of programs and data. We use a disk system with 10 Mbytes capacity (on fixed and removable media) for programming, operating the system, and for data storage. Some of our computer systems also have devices for data transfer and archival storage. These devices (connected by dotted lines in Fig. 10-48) may include flexible diskettes and cartridge-tape units. An independent battery-operated clock provides dates for our data files and time for the computer line clock.

Currently we use two classes of devices for data acquisition: analog-to-digital converters and CAMAC crate modules. Table 10-9 shows the characteristics of our converters. Many modules are available for CAMAC crates and Table 10-10 lists the ones we currently use. Our use of CAMAC systems provides measurement flexibility for our standard computer system.

For many of our LSI-11 computers we established the standard software environment listed in Table 10-11. This allows users and data to move from machine to machine without learning how each machine operates. This mobility assures maximum use of our machines. We hope to refine this environment and extend it to all our machines.

Data-Acquisition Program for LSI-11. Our primary data-acquisition program exists in four forms corresponding to the four projects where it is used (Minerva, Regulis, Extractor Test Facility, and SPP-II). The versions differ in minor ways and will be discussed as one program. The major components of the program are

- A description of the experiment.
- Data collection on demand or at a fixed time interval.
- Archiving commands, comments, and data.

Written in FORTRAN, this large program is constructed of functional modules contained in overlays. (An overlay is a portion of a program that resides in an area of memory used by other portions of the same program at different times.) Figure 10-49 shows the overlay structure and the major activity of the program. The main program invokes two types of processes. The first is a process that must be done periodically, such as our slow-signal acquisition, which simulates a chart recording. The other process is usually

control of CAMAC modules for data collection and control. User commands invoke these processes while periodic processes proceed automatically.

By reorganizing our originally unstructured program into the functional modules shown we were able to provide multiprocess operations. For example, a physicist may desire to obtain data at a sequence of probe positions and then return the probe to its original position. Each element of this task is present in the standard acquisition program but the task itself is not. To invoke such a task the program couples the existing routines with a control structure and places them in a separate program overlay. These may then be executed as if they were a basic routine. We customize such overlays for each project and each experiment within the project. We find these routines are a useful way of providing a changing acquisition

Table 10-9. Analog-digital instrumentation.

Manufacturer and model	Number of channels	Input range	Special features
ADAC 1030	32	± 15 V dc	Differential input programmable gain
ADAC 1112	8 or 16	Thermocouple	Cold-junction compensation

Operating system	DEC RT-11	<p>◀ Table 10-11. Computing environment of LIS LSI-11.</p> <p>Fig. 10-49. Data-acquisition-program overlay structure. ▼</p>
Languages	Pascal FORTRAN FORTH BASIC	
Utilities	CAMAC test routines General file display General data plotting

environment without causing major changes to the standard program.

Future improvement for our data-acquisition programs includes:

- Elimination of multiple-output files.
- Establishment of a tree-structured command syntax.
- Separation of the description of the experiment from the data-acquisition program.

Computer and Software for HP1000. The Mars computer-based data-acquisition system is a distributed system consisting of two processors, various peripherals, and acquisition hardware (see Fig. 10-51). The general-purpose interface bus (GPIB) interfaces the instrumentation in this system rather than CAMAC (used on our LSI-11 processors). We configured the data-acquisition system so it can expand to meet the Mars requirements as the experiment grows.

We are developing the data-acquisition-system software along two paths. Initially we constructed a minimal data-

acquisition system to record data as the experimental hardware became available. We completed this effort in December 1980 and applied it to Mars. This software, developed for experimental startup, will remain in use until the next phase of the data-acquisition software is completed.

This second phase is a design that is structured from the top down, using modern programming techniques. First we will write a requirements document that specifies as closely as possible what the user community expects of the system. We will design the system to meet these requirements and then begin the actual code development. The main advantage of a top-down structured approach is that although the design will consume more time, coding will require less time. We expect the final product to be adaptable to the rapidly changing experimental environment of the Mars facility.

Data-Acquisition Peripheral of Z8000.

The Z8000 data-acquisition and control peripheral (DACP) is a microprocessor peripheral device capable of performing data acquisition, real-time analysis, and control functions independently or under the control of a host computer. Figure 10-50 shows a block diagram of its components and its connection to a host computer. The microprocessor peripheral frees the host computer from primitive and time-consuming tasks, reduces program size and complexity, and can vastly decrease the amount of data that must be stored.

The DACP is primarily a Z8002 single-board computer packaged in a rack-mountable chassis with a 40-character display and interfaces to the LSI-11 through a CAMAC module. The inclusion of the AM9511 floating-point chip and service routines for arithmetic, keyboard, display, and CAMAC interface makes the system a powerful computing tool and provides ease in configuration to a particular application.

With the addition of commercially available digital-analog acquisition and control modules and the development of minimal-user software, the DACP is capable of performing a wide spectrum of acquisition, real-time analysis, and control

Fig. 10-50. Components and host connection of data-acquisition and control peripherals.

Alternative Applications of Technology

Fig. 10-51. Mars data-acquisition computer system.

functions. Two of these systems are in use with LSI-11 processors acting as the host. We can operate the DACP with a terminal, with a keyboard, and/or with a host computer through CAMAC. We use the terminal primarily for program development.

The DACP must perform functions defined by programs developed by users. These functions will become accessible through the keyboard and/or the host computer. Our system allows local entry of data and display of results. All function switches are software controlled and are unlimited in number and function. The system firmware provides display, keyboard entry, arithmetic, and CAMAC-interface control.

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Alternative Applications of Technology

Introduction

The technology that we develop during our atomic-vapor laser isotope separation (AVLIS) program often has applications in areas remote from uranium-isotope enrichment. The following two articles describe such applications. First we discuss the feasibility of applying an AVLIS process to the separation of platinum metals. Then we describe the ongoing development of Vulcan, a high-temperature gas heater originally intended as an alternative to our electron-beam source.

Author: B. W. Shore

Laser Cleanup of Pt-Group Metals

Because of increasing interest in chemical and fuel synthesis from syngas, we have reevaluated the feasibility of purifying Pt-group metals palladium (Pd), ruthenium (Ru), and rhodium (Rh) using laser isotope separation (LIS) technologies. Nuclear reactors produce Pd, Rh, and Ru in significant quantities as fission products.⁵⁹ Industry commonly uses these metals as catalysts for reforming petroleum products and for producing nitric oxide and glass

fibers. Currently we import about 95% of the Pt-group metals needed in the US from the USSR and South Africa. Thus it is strategically important for us to produce these materials domestically. Unfortunately, reactor-produced Pt-group metals contain small amounts of harmful radioactive isotopes. These isotopes must be removed before the metal can be used as a catalyst in industry. If we can use LIS technology to accomplish this task, it might make a significant impact on the synfuel program in the foreseeable future.

At present there is no economic incentive for extracting the Pt-group metals from

Fig. 10-52. Selective photoionization of Pd by circularly polarized light.

nuclear wastes using conventional methods. We will take Pd as an example. The mid-1979 market value of Pd was about \$3.50/g. The estimated cost⁶⁰ of nonisotopic Pd recovery from nuclear-reactor wastes is \$12.30/g, or 4 times the current value of imported Pd. However, it appears that we could remove the harmful radioactive isotopes from the reactor-produced Pt-group metals by using an LIS technique. We briefly describe the proposed technique for ¹⁰⁷Pd removal (the harmful isotope in Pd).

Earlier we made a preliminary LIS feasibility study for the Pt-group metals.⁵⁹ We found that a simple atomic-vapor photoionization process, as developed for the uranium-isotope separation program, did not appear to be directly applicable to the Pt-group metals. The isotope shifts of these metals are much smaller than the Doppler-broadened visible absorption widths. Thus it is difficult (if not impossible) to achieve selectivity without resorting to a Doppler-free technique.

The technique described here is similar to that proposed by Balling and Wright⁶¹ for separating isotopes having different nuclear spins and hyperfine structures. The selectivity of this technique relies on the angular-momentum selection rules for atomic absorption of circularly polarized light. In principle it is possible to selectively excite isotopes with nonzero nuclear spin, using two circularly polarized laser pulses, while leaving the isotopes with zero nuclear spin unexcited. We can then ionize the excited atom using a third photon of appropriate energy. Figure 10-52 is a schematic drawing of the excitation and ionization processes for selective removal of ¹⁰⁷Pd from Pd.

no further excitation.

The selection rules for atomic absorption of left-circularly polarized laser light are

- $\Delta M_J = +1$ for the even isotopes
 $I = 0$ (^{104,106,108,110}Pd),

- $\Delta M_F = +1$ for the odd isotopes

$$I = 5/2 (^{107,105}\text{Pd}),$$

where $F = I + J$. Note that the photoionization selectivity relies on the absorption cross section for the second transition (³P₁ → ³P₀ at 5210 Å), which is zero for the even isotopes.

The apparent limitations on the selectivity of this purification process are

- The degree of circular polarization that can be achieved in the two laser beams (λ_1 and λ_2). The presence of an incorrect polarization such as right-circular polarization will result in an unwanted excitation of even isotopes with $I = 0$.
- The collinearity of the two laser beams (λ_1 and λ_2). If the two beams are not perfectly collinear, an even isotope with $I = 0$ excited to the $M_J = 1$ level of the ³P₁ state would have a nonzero probability of being in the $M_J = -1$ level relative to the axis of the quantization defined by the second beam, enabling it to make a $\Delta M_J = +1$ transition.
- The oscillator strength of the second transition (³P₁ → ³P₀ at 5210 Å), which determines the minimum laser intensity required for effective removal of the odd isotopes from the ³P₁ level. If the required laser power is too high, some fraction of the even isotopes might also be excited through multiphoton processes under the influence of the intense laser field.

There are two issues pertaining to the economics of this purification process. One is the optical selectivity of the first transition, ¹S₀ → ³P₁ at 2763 Å, or the degree of unwanted laser absorption by the even isotopes ($I = 0$) and ¹⁰⁵Pd. The other is the oscillator strength of the second transition ³P₁ → ³P₀ at 5210 Å for the odd isotopes (^{107,105}Pd) with $I = 5/2$. It is important that we can effectively remove a large fraction of the ¹⁰⁷Pd isotopes in a relatively small number of stages.

Figure 10-53 is a composite atomic-absorption spectrum (λ_1 , ¹S₀ → ³P₁ at 2763 Å) for Pd recovered from the reactor waste. The initial isotopic composition is 18% ¹⁰⁷Pd, 19% ¹⁰⁴Pd, 23% ¹⁰⁵Pd, 26% ¹⁰⁶Pd, 12% ¹⁰⁸Pd, and 3% ¹¹⁰Pd.

The isotope shift for the 2763-Å transition is roughly 8 mK (10^{-3} cm^{-1}) per mass change. The nuclear spins of ^{107,105}Pd are identical ($I = 5/2$). We assume that the nuclear structure

and the nuclear quadrupole moments for $^{107,105}\text{Pd}$ are also similar (within $\pm 20\%$), which allows us to make estimates for the hyperfine coupling constants (A and B) and hyperfine splittings for ^{107}Pd .

It might be possible to attain some selectivity on the first step using a single-mode pulsed dye laser tuned to a region between the hyperfine absorption features of ^{107}Pd ($F = 5/2 \rightarrow 5/2, 5/2 \rightarrow 7/2$). Since ^{105}Pd also has a nuclear spin of $5/2$, some fraction of the ^{105}Pd will also be excited and ionized along with the ^{107}Pd . We estimated that roughly 83% of the excited ^{107}Pd (3P_1) can be further excited to the 3P_0 level and hence ionized by λ_3 . In principle we can ionize 65% of the ground ^{107}Pd atoms per stage. For achieving a final ^{107}Pd concentration of 2.2% (an acceptable level for industry), we require a minimum of two stages.

Fig. 10-53. Composite absorption spectra of Pd recovered from nuclear waste.

The oscillator strengths of the second (λ_2) and third (λ_3) transitions are not known. Clearly the first priority on ^{107}Pd cleanup is to map out transitions that have acceptable strengths for the second and third steps when using circularly polarized laser light. Assuming that we find adequate transitions, we can then make a more realistic cost analysis for the purification process.

We have made a crude estimate for the cost of ^{107}Pd cleanup using the LIS technique described. The cost increase for Pd recovery (from nuclear wastes) because of the added laser-purification process is roughly 1 to 2 \$/g, which amounts to roughly 10% of the final reactor recovery cost of 12.3 \$/g.

Author: H. L. Chen

Vulcan Very-High-Temperature Gas (VHTG) Heater

Introduction. We are developing a technology that uses mechanical compression to generate streams of contamination-free gas with temperatures up to 9000 K. Our heater has a number of intrinsic advantages over competing electric-discharge heaters. It is relatively efficient, it bulk heats the gas uniformly, and it is insensitive to the electrical properties of the gas. We patterned the equipment after well-developed industrial compressors and it should be very reliable and easy to maintain.

Once developed it will be a unique and valuable tool as a source of bulk-heated gas of known composition for wind-tunnel testing, high-temperature chemical processing, and many other research purposes.

Concept. Presently the most common method of producing very-high-temperature gas streams for large-scale industrial applications and scientific research use is electric-discharge heating. Although the conventional discharge technologies have been under development for many years, certain undesirable performance characteristics remain: plasma instability; restrictions on the volume of generated plasma; gas-temperature inhomogeneities; and strong sensitivity to the exact nature of the gas and the operating pressure. We are developing a new technology whose intrinsic

features offer significant advantages over conventional electric-discharge heating. Our method uses mechanical energy to heat gas directly to temperatures as high as 9000 K by compression as originally outlined by Weatherston.⁶²

We need a high ratio of output to input pressures to achieve the desired high ratio of temperatures. As a first approximation we consider a reversible adiabatic-compression device for which the temperature ratio is given by

$$\frac{T_2}{T_1} = \left(\frac{p_2}{p_1} \right)^{(k-1)/k} \quad (35)$$

Consider for the moment heating a gas from room temperature (300 K) to a discharge temperature of 4500 K:

$$\frac{T_2}{T_1} = \frac{4500}{300} = 15 \quad (36)$$

This would require a pressure ratio:

$$\frac{p_2}{p_1} = 15^{2.5} = 870 \quad (37)$$

for helium, neon, and other monatomic gases with $k = 5/3$.

There are two classes of compressors to be considered for this service: centrifugal or axial steady flow; and positive displacement such as a piston-type. Weatherston has pointed out the need in a VHTG heater to maximize the power per unit area to minimize the temperature loss to wall cooling.

A centrifugal steady-flow compressor has a maximum stage pressure ratio r of approximately four so that the number of stages N for

$$R = \frac{p_2}{p_1} = 870 \quad (38)$$

would be

$$r^N = R \quad (39)$$

which gives

$$N = 6.6 \quad (40)$$

or seven in whole numbers with r suitably trimmed back. The number of stages for an axial steady-flow compressor would be much higher because r is typically ~ 1.1 .

Following the development of Taylor⁶³ we can show that the power per unit of transverse area in a centrifugal impeller stage is approximately

$$\frac{P_c}{A_x} = p_d \frac{A_r}{A_x} \frac{V_r}{V_t} V_t k M_t^2 \quad (41)$$

where

P_c = stage power input;

A_x = transverse area of the stage;

p_d = stage discharge pressure;

A_r = tip gas-flow area;

V_r = tip gas-flow velocity relative to rotor;

V_t = tip rotor velocity;

k = specific-heat ratio; and

$M_t = V_t/c_d$,

where

c_d = speed of sound in tip gas.

In a piston-type positive-displacement compressor the unit power during constant-pressure delivery is

$$\frac{P_p}{A} = \frac{FV}{A} = p_d V \quad (42)$$

where

P_p = positive-displacement compressor power input;

A = piston area;

F = piston force; and

V = piston velocity.

Equating the two unit powers we obtain

$$\frac{A}{A_x} = \frac{A_r}{A_x} \frac{V_r}{V_t} k M_t^2 \frac{V_t}{V} \quad (43)$$

The first term A_r/A_x generally has a value of about 0.2, V_r/V_t is about 0.3 to 0.4, and the specific heat ratio is 5/3. For the final stage M_t is approximately 0.2. These values give

$$\frac{A}{A_x} \approx 0.005 \frac{V_t}{V} \quad (44)$$

Centrifugal stress in the impeller limits V_t to a value of about 300 m/s. Piston velocity varies between 30 m/s for piston compressors to about 150 m/s for rotary equivalents. Substituting values we obtain

$$0.01 < \frac{A}{A_x} < 0.05 \quad (45)$$

from which we conclude that there is much more working-volume surface area in a

Fig. 10-54. Heater cross section.

Fig. 10-55. Overall view of heater. ▼

centrifugal compressor than in a positive-displacement piston-type compressor for the same input power and therefore much more loss of temperature because of wall cooling. Corrections to the foregoing analysis for centrifugal diffuser areas, staging effects, and positive-displacement cylinder-wall areas will probably increase the advantage of a positive-displacement compressor.

As a result of this analysis we are developing our VHTG heater in a rotary positive-displacement form much like the familiar Roots blower. We show its cross section in Fig. 10-54 and its overall view in Fig. 10-55. The rotors counterrotate and are synchronized at exactly a 45° phase angle. In Fig. 10-54 gas is positively displaced from right to left in the wells formed by the case and rotor lobes. When a leading lobe tip reaches the discharge volume there is a back flow into the well that equalizes the pressure. Well A in the figure has opened and well B is about to open.

In normal service the Roots blower operates at an r value of about 2. In our service we come close to shutting off the flow so that r can be very large. The rotor-to-rotor and rotor to case clearances are typically a few thousandths of a cm to limit back leakage without oil or solid seals. The absence of sealing oil minimizes working-gas contamination and the absence of solid seals reduces friction losses. Oil-sprayed antifriction bearings support the rotor spindles, and labyrinth seals and pumpouts isolate them from the working volumes so the leakage and friction are very low.

Since back leakage is constant for a given pressure ratio across the heater, we sustain the desired level by running at high speed. The very high temperature of the gas requires elaborate rotor cooling, which leads to high thermal stresses in the rotor. These combined with the rotational stresses require a high-strength rotor material. Elaborate cooling is also required for the case walls and discharge nozzle.

Table 10-12 gives our estimate of the design-point operating conditions for helium working gas.

Note that the pressure-ratio requirement is considerably below the value of 870 calculated for reversible adiabatic compression. Although the cooling itself increases the pressure-ratio requirement, that increase is overwhelmed by our estimate

of the inefficiency, which in a compressor decreases the pressure ratio for a given temperature ratio.

Figure 10-56 shows one of the impellers from our first heater. Gas-feed grooves are shown radially disposed on the lobe cover plate. The bolt pattern secures the cover plate over the cooling-water holes. At the opposite end of the lobes is a similar cover plate. Water enters at one shaft end and returns at the other shaft end on each impeller.

A slight overpressure of the gas supply to the feed grooves as shown in Fig. 10-57 helps to reduce hot-gas leakage at the impeller ends. Figure 10-57 also shows the net flow of 2.2 g/s to the P_1 vacuum pump, which establishes a low inlet pressure to the impeller wells and thereby sets the pressure ratio of 450 for a discharge flow of 14.0 g/s. Vacuum pumping of labyrinth-groove 1 ensures outleakage of the working gas to prevent contamination of the main stream.

Since both gas leakage and thermal conductivity decrease for neon compared to helium, we estimate that design point operation with neon would give approximately a 9000-K discharge temperature with intermediate temperatures for various He-Ne mixtures.

Testing. We built and tested the first heater in 1978 as a cooperative venture with Weatherston and his associates representing Calspan Corporation. We conducted startup operations with a bypass connecting the discharge to the input gas supply. The bypass allowed minimum loading. Unexpected vibrations caused the impeller lobes to bump and shut the heater down with significant damage.

Test data taken during this startup, given in Table 10-13, show good agreement between the net measured power W_{net} and the calculated power W_{Roots} . These are the powers supplied to the gas so they represent the actual rise of the gas temperature. Table 10-13 shows that the ideal rise ($W_{\text{rev. ad.}}$) is only 58% of actual. This low efficiency of compression, undesirable for heat-engine cycles, is helpful for our gas heating by reducing the necessary pressure ratio as previously discussed.

We are redesigning the impellers with a new profile and a helical form that we expect will eliminate the bumping hazard. We are also revising power flow paths and stiffnesses in the drive train for better

positioning of the torsional and bending critical speeds with respect to the design-point speed.

We upgraded the original data-acquisition and control system so it relies completely on microprocessors. This system provides safety interlocks, command sequencing, status monitoring, and data storage. It samples and records 32 fast channels at a rate of 1500/s and 400 slow channels at a rate of 1/s. Diagnostics consist of sensors to measure pressure, temperature and flow of

Input power	650 hp
Cooling-power loss	170 hp
Discharge gas temperature	4500 K
Discharge gas flow	14 g/s
Pressure ratio	450
Discharge pressure	2 atm
Rotor speed	11 000 rpm

Table 10-12. Design-point operating conditions for helium.

Fig. 10-56. Vulcan impeller. ▼

Fig. 10-57. Gas-flow ▲ diagram.

Heater speed (rpm)	Power (W)				
	Test data		Calculated		
	W_{in}	W_{loss}	W_{net}	W_{Roots}	$W_{\text{rev. ad.}}$
5887	24.6	6.5	18.1	18.1	10.5
6695	31	8.4	22.6	22.8	13.2
7222	41	9.8	31.2	28.1	16.3

Table 10-13. Heater performance.

gas, structure velocity and acceleration, and shaft power. We are developing a telemetry system to monitor impeller temperatures and strains and a noncontact system to monitor rotor torsional and lateral vibrations. The console includes a real-time vibration-spectrum analyzer. We expect to begin testing of the redesigned heater in 1982.

Authors: B. Myers and J. G. Harri

Planning Methodology

Introduction

The ultimate goal of our laser isotope separation (LIS) program is to develop the technology that will permit economically attractive construction and operation of atomic-vapor laser isotope separation (AVLIS) plants. In working toward this goal we must not only develop the needed technology but also determine the appropriate operating conditions for a plant to optimize the economic benefit. The following article illustrates the kind of issues we have addressed in analyzing plant design.

In an attempt to build a quantitative foundation for decisions governing allocation of our manpower and monetary resources, we have applied techniques of decision analysis to our program. The final article describes our work on this topic.

Author: B. W. Shore

AVLIS Design Criteria

Our economic analyses represent the results of extensive modeling of the relationships between fundamental physical parameters and the relative enrichment parameters that describe a plant. Figure 10-58 schematically diagrams the components of any isotope-separation plant. The feed, whose molar fraction of ^{235}U is X_F , undergoes separation into product, whose molar fraction of ^{235}U is X_P , and tails, whose molar fraction of ^{235}U is X_T . Two parameters will, along with the mass flow, combine to determine the separative work capacity of a process. These are the relative concentration of the product compared with the feed $\beta_1 = X_P/X_F$ and the relative concentration of the tails compared to the feed $\beta_2 = X_F/X_T$. The target values of these parameters for an efficient resource-conserving process are $\beta_1 \lesssim 4$ and $\beta_2 \lesssim 3$. If a process yields smaller values of β_1 or β_2 we must restage, with product and/or tails of one stage becoming feed for a subsequent stage. We can gain an enormous advantage if we can design a process that yields target values of β_1 or β_2 in a single stage without recycling. We believe this is possible with our atomic-vapor laser isotope separation (AVLIS) process.

The conventional measure of separation effectiveness is the separative work unit per kilogram, defined as:

$$\frac{\text{SWU}}{M_F} = \left(\frac{\alpha - \beta_2}{\alpha - 1} \right) \ln \beta_2 - \left(\frac{\beta_2 - 1}{\alpha - 1} \right) \ln \beta_1 \quad (46)$$

where $\alpha = \beta_1 \beta_2$ and M_F is the mass of the feed. This defining expression simplifies if $\beta_1 \gg 1$, when it becomes

$$\frac{\text{SWU}}{M_F} \approx \ln \beta_2 \quad (47)$$

Of greater interest to an AVLIS process, when $\beta_2 \gg 1$ the expression becomes

$$\frac{\text{SWU}}{M_F} = \left(\frac{\beta_1 - 1}{\beta_1} \right) \ln \beta_2 \quad (48)$$

In a laser process very high selectivity is possible, leading in principle to very high enrichments $\beta_1 \gg 1$. However, for efficient generation of separative work we must also obtain a significant depletion of the tails

Fig. 10-58. Elements of isotope-separation processes.

stream ($\beta_2 \gg 1$), which is more of a challenge. That is, a process cannot just enrich a small part of the feed stream spectacularly well; it must extract the desired uranium isotope from a large part of the feed.

We control the enrichment parameters β_1 and β_2 through our process design by choosing basic physical parameters. The principal atomic parameters are the optical-absorption cross sections σ and the spectroscopic selectivity S , defined as the ratio of ^{235}U to ^{238}U cross sections at the frequency of peak ^{235}U absorption. Given the large number of transitions in uranium atoms we can pick transitions that match the lasers we want to use and that have excellent spectroscopic selectivity. Then we must choose values for the vapor density n and the vapor column length L , constrained by the fact that the vapor opacity cannot exceed unity. This choice ensures that all the atoms are illuminated, while at the same time it minimizes the energy and intensity that the laser must deliver. To express laser fluence we often use the saturation parameter χ , which is the laser fluence scaled to the saturation fluence of the particular atomic transition being pumped. As a rough guide, when $\chi = 1$, half of the atoms are excited in a two-level system. In practice saturation parameters considerably above 10 may be required for each transition to obtain efficient ionization of a multilevel atom.

In Fig. 10-59 we show one typical result from our modeling of the relation between the fundamental physics parameters and the relative enrichment parameters that describe a plant. Here we simultaneously consider two cascades of enrichment stages: one cascade of N_1 stages, which uses the enriched stream from the previous stage as feed for the next, and one "all tails" cascade of N_2 stages, which uses the depleted stream as feed for the next stage. We assume overall plant β_1 and β_2 values that give the standard light-water-reactor (LWR) product and tails assays, set N_2 at one of a range of values, and ask how many enrichment stages are required as a function of the spectroscopic selectivity. This is done both for an idealized process and for one where some of the ^{238}U is unavoidably collected with the photoionized ^{235}U . There is clearly a great advantage to working with a selectivity $S \gg 10$, a feature that sets laser processes

apart from aerodynamically based processes or from conventional gaseous diffusion.

Having seen in general terms how the SWU production rate relates to basic parameters, we can look in more detail at the factors that influence the choice of one of these parameters, the vapor number density n . The driving force in this analysis is the desire to minimize the cost per SWU as shown in Eq. (49),

$$\frac{\$}{\text{SWU}} = \frac{(\$/\text{kg})_M + (J/\text{kg})_L (\$/J)_L}{\text{SWU}/\text{kg}} \quad (49)$$

Here $(\$/\text{kg})_M$ is the cost per kilogram of the materials-handling subsystem, $(\$/J)_L$ is the cost per joule of the laser electro-optical subsystem, and $(J/\text{kg})_L$ is the laser energy required per kilogram. Assuming we have chosen β_1 and β_2 , which are implicit functions of n , such that the separative work per kilogram of feed is high, we then need to examine the density dependence of the remaining cost factors: the laser and materials-handling subsystems.

The materials-handling costs originate with several process elements: feed, vaporization, product and tails collection. Our analyses show that the density dependence primarily occurs in the cost of collection. This cost in turn is roughly

Fig. 10-59. Minimum number of product enrichment stages N_1 required vs selectivity S for various numbers of tails stripping stages N_2 .

proportional to the collector surface area. Therefore we want to get maximum throughput per unit collector area, corresponding to a materials-handling cost that decreases as the density increases.

The other cost component is the laser system, which irradiates an area W by H over process length L . As a first-cut analysis of the density dependence, the laser energy required per kilogram is the laser pulse energy times the laser pulse-repetition frequency (prf) divided by the mass throughput, leading to a $1/nL$ dependence for the specific laser energy. However, in this analysis L cannot be treated as an independent variable. In practice $L \propto (1/n)^\delta$ with δ in the range $1 \leq \delta \leq 2$. The requirement that the vapor opacity be less than 1, as mentioned above, gives $\delta \sim 1$. More detailed examination of laser-propagation effects, such as refraction and diffraction, indicates values for δ above 1. The result of all these considerations is the laser-cost scaling given by the expression

$$(\$/\text{kg})_L = (\$/J)_L (J/\text{kg})_L \sim (n/n_0)^\gamma, \quad (50)$$

$$(0 \leq \gamma \leq 1)$$

so that $(\$/\text{kg})_L$ increases slowly with vapor density n .

Combining the materials-handling and laser-cost models described above, we get the cost vs density curve shown in Fig. 10-60. The optimum operating density is, not surprisingly, a compromise between the limiting effects listed in Fig. 10-60.

Similar arguments can be applied to the other process variables that we are free to choose. Finding the best operating point is a multidimensional nonlinear optimization problem whose solutions are continuously being updated by refined values for the process parameters.

Authors: J. I. Davis and R. C. Stern

Major Contributor: B. W. Shore

Optimal Resource Allocation and Decision Analysis

Introduction. Any extensive R&D program requires numerous decisions as the program evolves. In our atomic-vapor laser isotope separation (AVLIS) program, decisions span issues associated with process selection, systems design, operating and funding schedules, and the setting of performance goals and milestones. Usually we must choose among several processes and systems, but we would like to maintain alternatives until we are assured success in each required category. For example, we are considering a number of alternative pump laser systems for the AVLIS process. If resources of money, manpower, and time were unlimited we would develop all available alternatives. Construction of the full-scale laser isotope separation (LIS) plant would be based on the best alternatives available at the time of design freeze. In practice we must allocate limited resources (typically money) among a few competing alternatives. Thus as the program progresses we must eliminate alternatives as soon as it becomes clear that their contribution to the expected final payoff no longer justifies the resources they consume.

We have developed two computer codes for use as decision tools that systematize the timing of decision points, the setting of funding levels, and the dropping of alternative technologies. These codes aid the decision maker in formulating strategies that are consistent both with his preferences regarding outcomes and with his subjective judgements concerning the probabilities of unknown factors or future events.

The simpler model, the static nonlinear programming (NLP) code, makes the assumption that there is only one decision point, at which time we optimally allocate a fixed total budget among the various alternative technologies in each LIS program. We maximize the overall probability of success, defined as the successful completion of at least one of the LIS projects, using the numerical-

Fig. 10-60. Dependence of cost per separative work unit on vapor density n .

optimization techniques of nonlinear programming. Our second code, the dynamic-resource allocation code SAGE, is more complex and makes use of the methods of statistical decision theory. In SAGE we allow for up to five decision points at which we reallocate budgets and where we may drop alternate technologies or LIS programs to maximize the expected (average) payoff.

Such tools of decision analysis do not actually make decisions. Only someone cognizant of all quantifiable and qualitative factors in the project, and who is also maintaining the perspective of the project context, can make wise decisions. Rather, such tools provide a language and framework for clarifying issues, focusing discussion, and developing intuition regarding sensitivities to important parameters. Using these tools we can examine the consequences of various assumptions, scenarios, and possible alternate strategies.

One-Time Allocation Problem. Our NLP code makes the simplification that there is only one decision point, at which time we optimally distribute a fixed total budget among alternative competing laser drivers (D) and uranium-handling (U) technologies to maximize the probability of eventually achieving scientific, engineering, and economic success for an LIS process. Figure 10-61 illustrates an example of such a situation. Our code can also optimally allocate a total DOE/LIS budget among a number of competing LIS programs. As presently implemented, each of three LIS programs has four drivers and four uranium-handling technologies. Thus the allocation of a fixed overall DOE budget for LIS constitutes a 24-dimensional optimization problem constrained by the prior expenditures $D\phi_i$ for each technology in each program.

Both of our decision codes use a two-parameter functional relationship for the dependence of the probability of eventual success of the i^{th} technology [e. g., copper-vapor laser, rare-gas halide, and frequency-doubled neodymium YAG (FDNY) laser technologies] on the total dollars spent (D_i) on R&D for that technology,

$$P_i = P\phi_i \left[1 - \exp \left(- \frac{D_i}{\text{SAT}_i} \right) \right] \quad (51)$$

In this expression SAT_i is the level of funding that would begin to saturate the development of the i^{th} technology. The other parameter $P\phi_i$ is the probability of achieving success if we had all the money we reasonably needed for development of the particular technology.

The various laser and uranium-handling technologies we investigated comprise a series/parallel circuit where the set of all laser technologies is in series with all uranium-handling technologies (both are necessary). However, CVL, RGH, and FDNY laser technologies are in parallel with each other (success in any one is sufficient). The logic of this type of circuit, shown in Fig. 10-61, also applies for each of the competing LIS programs. Additionally we assume that the three LIS programs are in a parallel development mode. In other words success in any one of the programs would constitute success (although with varying levels of benefit) of the overall LIS program.

Using the functional form of Eq. (51) for the dependence on funding of the probability of successful development of each technology, we can write the expression for the dependence of the overall probability of success P on the funding of each alternate technology D_i . Describing the funding allocation as a vector, we express our problem as

$$\begin{aligned} &\text{maximize } P(D) \\ &\text{subject to the constraints} \\ &\quad (a) D_i \geq D\phi_i \text{ and } (b) \sum D_i = \text{budget.} \end{aligned}$$

The NLP code performs this optimization using a modified Fletcher-Reeves conjugate-gradient technique using penalty functions.⁶⁴ The code output gives the budgetary allocation to each of the alternative

Fig. 10-61. LIS program structure.

technologies, along with the probability of success for each program. We can thus exercise the model using the latest estimates of the input parameters and the expenditures to date and find the relative proportion of the next year's funding that should be spent on R&D in each technology area. This also identifies those technologies (given the finite remaining budget) that should not be funded any further.

We have exercised the NLP code using preliminary data representing rough estimates of $P\phi_i$ and SAT_i for sets of technologies in each of the three competing LIS programs separately. The results provide an initial estimate of the optimal resource allocation in each LIS program, subject to a given program budget. We also found the corresponding optimal probability of success for each project and for the LIS program as a whole. Using the same input data, but allowing the NLP code to allocate the total budget among the three programs, we found a significant improvement. Sensitivity studies varying the total budget provide the optimal number and identity of technologies in a program (and the optimal number of programs) at each level of funding. As we reduce the funding the code begins to allocate zero resources to alternates that previously had been funded.

Dynamic Resource-Allocation

Methodology. The methodology we have developed for the sequential allocation of R&D dollars is an application and extension of the techniques of statistical decision theory. Certain concepts, problem structure, and calculations occur in any statistical decision-theory application.⁶⁵ Certain other problem elements are peculiar to the LIS R&D application where a number of alternate programs or technologies are developed in parallel. To aid in the exposition of possibly unfamiliar decision-theory principles we will begin with a brief description of a simplified problem, which we analyze within the framework of statistical decision theory.

Statistical-Decision-Theory Example. Our example consists of the R&D of two pump lasers, laser 1 (L1) and laser 2 (L2), e. g., any two of FDNY, CVL, or RGH. Although we require only one pump laser system for our development of an AVLIS process, we are not 100% certain at the outset that either laser system can be adequately developed for our needs. Therefore we hedge

our bets by developing in parallel two laser systems, with the expectation that we have significantly increased our chances of success with at least one of the alternates. After a period of parallel development we must decide (presumably based on the outcome of the development effort) which laser to drop from consideration, and which to develop to (hopefully) eventual success. For simplicity we assume that the outcome of the initial development stage (in addition to an enhanced state of development of each technology) is either a positive or negative indication for each of the laser types (e. g., L1 looks good, L2 looks bad, etc.).

We can depict the sequence of development results and decisions in a decision tree as shown in Fig. 10-62. Two types of nodes occur on the decision tree: probability (P) nodes and decision (D) nodes. These alternate along any path. Probability nodes represent the outcome of some chance event (in this example the result of development). We have no control over which branch to the right is taken, but we do have certain probabilistic knowledge regarding the likelihood of each alternative. At decision nodes we do have complete control over which path to take. In fact, it is assumed that we have some decision algorithm with which we make this choice. Our algorithm, of course, is a function of the particular path we have traversed to reach each D-node. The final P-node in our decision tree leads to either success or failure of the particular technology, with its associated benefit or penalty. In this example we have chosen all benefits equal to one and all penalties equal to zero. Since we will be maximizing the expected payoff, this choice corresponds to maximizing the probability of success. Note that even for this simple problem, with only one development phase and one decision to be made, the tree already has 29 branches representing all possible alternative sequences of decisions and outcomes.

The procedure for analyzing a problem using statistical-decision theory consists of the following four steps.⁶⁵

1. Set up the decision tree from left to right, defining the probability and decision nodes, hence specifying the logic of the program activities and the decision sequence.
2. Assign the payoffs at the terminal nodes, thus determining the figure of merit to be optimized.

3. Calculate the probabilities at the chance forks (P-nodes) from left to right. This consists of a Bayesian analysis,⁶⁵ for each possible branch, of the probabilities of success of each laser system that updates these probabilities to reflect each of the possible results of the preceding development stage. We also calculate the (predictive) probability that each of these branches will in fact be realized. These predictive probabilities appear in parentheses in Fig 10-62.

4. Average out and fold back. Working from right to left we calculate the average payoff at each probability node as it is encountered. At each decision node we select the branch that leads to the P-node with the greatest expected payoff. We associate this value of expected payoff with the D-node. This branch then represents the decision made, e. g., drop laser L2. This constitutes the optimal use of our limited control at each decision point. Continuing from right to left we construct the complete set of optimal decisions.

This analysis yields the expected payoff written at the left-most node. Note that in our example this is the same as the probability of success. In Fig. 10-62 the numbers written in parentheses on each branch emanating from a P-node are the probabilities that the respective path will be traversed. The other number written at each node is the expected payoff corre-

sponding to that position on the tree. We have crossed out the unfavored options from each D-node.

Behind our analysis of this decision tree lie calculations that require specification of a "likelihood function" representing the reliability of R&D information. An example of this function is the conditional probability P_r (predict L1 ok|L1 will be successful), which we use in the Bayesian probability analysis. For our example we have arbitrarily chosen this probability as 0.6 for both technologies. Later we will need to specify a model that determines this probability as a function of the level of R&D funding for each technology. Clearly, more dollars spent should yield more reliable R&D information. Other inputs required for this simple example include the a priori probabilities of success for each laser technology, here chosen to be 0.9 for L1 and 0.6 for L2. The right-most parentheses in Fig. 10-62 enclose the corresponding posterior (relative to the development stage) probability for each laser system.

Our analysis in Fig. 10-62 shows that whatever the outcome of the remaining development prior to a final decision, we will eventually terminate development of laser L2. This conclusion is a consequence of the relative weakness of our reliability-of-research factor 0.6 (0.5 corresponds to no usable information). We conclude that L2 should be dropped immediately, since its

Fig. 10-62. Decision tree for L1 and L2 laser development.

continuation adds nothing to the probability of success (0.9). The other choice is to provide enough additional funds during the current development stage so that the reliability-of-research factor becomes great enough to influence the final decision. This would definitely increase the probability of success.

Specific Model Assumptions. The statistical-decision-theory model just described treats only single decisions, i. e., which laser program to drop. The size of the tree increases almost exponentially with the number of decisions required. We also wish to decide how much money should be allocated to each technology at each research stage. If we were to represent the continuum of possible allocations, even divided into fairly large increments by alternate branches from the decision nodes, the problem would quickly become intractable. To deal with this problem we have derived a number of alternative budget-allocation algorithms. We use these algorithms to distribute at each D-node a fixed budget for each R&D stage among the surviving alternate programs or technologies. Depending on the option selected the budget at each stage is allocated proportional to

- The initial allocations input for the first R&D stage.

$$\frac{P\phi_1}{SAT_1} \exp \left[- \frac{D_1(\text{prior})}{SAT_1} \right] \quad (52)$$

- The gradient of the probability of overall success with respect to the dollars allocated.
- The gradient of the expected payoff with respect to the dollars allocated.

The third and fourth options locally maximize the increase in probability of eventual success and expected payoff, respectively, as they are influenced by the dollars allocated. We can also include additional allocation algorithms.

The money spent on R&D gives rise to two separate types of gain. One type has more of the "research" character. Analogous to a market survey or statistical sampling, it garners more precise information to aid in better forecasts and decision making. The other type of gain has more of the

"development" character, in which the program or technology is brought to a more mature and complete level of development. This would include construction of required facilities, attainment of milestones, defining a preferred process, and selection of operational options and procedures, etc.

Our methodology reflects the research gain by using the likelihood function. We apply this function in the Bayesian analysis for the calculation of both *posterior* probabilities of success and the probabilities of various research outcomes. The likelihood function, which describes the reliability R of the information derived from R&D during each stage, is a function of the dollar allocation to each technology during that stage. For each technology if one expected a positive or negative indication with a probability of 0.5 (irrespective of whether success is in fact eventually achieved), then clearly that would constitute no information, i. e.,

$$P_r \left[\begin{array}{c|c} \text{Positive} & \text{Eventual} \\ \text{indication} & \text{success} \end{array} \right] = 0.5 \quad (53)$$

Hence this should be the value of the likelihood function when no money is allocated to a technology, i. e., the result is no new information. The equation we have used to represent the dependence on dollar allocation is

$$R(D_1) = 1 - \frac{\exp(-\beta D_1 / SAT_1)}{2} \quad (54)$$

We choose the parameter β (generally equal to one) to rescale the saturation dollars. Large values of β correspond to very efficient research, i. e., research that generates reliable predictions.

We now briefly discuss the manner in which we treat the development aspect of R&D. Note that the a priori probabilities of success (and their subsequent Bayesian updates depending on results of prior R&D) are based on anticipated total expenditures. These have the same functional form as in the nonlinear programming model. Programmatic decisions at the D-nodes (dropping programs and budget reallocation) may result in greater (lesser) total expenditures and a consequent greater

(lesser) chance of eventual success. Hence we update the probabilities of success for each program or technology both at each P-node (because of the current indications of the most recent research stage) and at each D-node (following the latest reallocation of funds).

Finally we specify the logical sequence of activities and decisions that we wish to analyze. In starting our analysis after an arbitrary amount of funds have already been expended, we begin with the next research stage and an initial allocation for three programs or technologies being developed in parallel. (Below we extend this analysis to the case in which there are up to four programs or technologies with an immediate decision at hand to drop one.)

Figure 10-63 shows the sequence of probability and decision nodes representing the results of research and reallocations of budgets, respectively. Note that only on every other D-node is one program or technology dropped. This is because no matter how negative the latest research results, a program or technology is generally phased down gradually, rather than abruptly terminated. This allows for the possibility that it may later be ramped up as a result of significant innovations during the "backburner" stage, along with possibly diminished expectations for the alternate programs. We indicate below each node the multiplicity of branching. This decision tree, if fully drawn, would have over 23 000 branches.

Description of SAGE Code. We implemented the methodology described above in a computer code SAGE (*Sequential Allocation GEnerator*) for the CRAY computer. The code handles as many as three (four) programs or alternate technologies if the first branch point is a probability node (decision node at which one program or technology is dropped).

If the case being run requires the first node to be a decision node, at which one of the

four programs is dropped, the code provides the option of some operator interaction. This is because the selection of the program or technology to be dropped first may correspond to an actual program decision at hand. The code provides the decision maker with the implications of each of the four options corresponding to dropping each of the four programs or technologies. In addition, although the SAGE code makes decisions based on maximizing expected payoff, other considerations and figures of merit may be of interest. The code first presents the decision maker with the expected payoff, standard deviation of payoff, and probability of success for each of the four options. For a number of reasons he may choose to drop technology P4 even though dropping technology P3 leads to a slightly higher expected payoff. The decision to drop technology P4 may lead to a substantially lower standard deviation of payoff or it may result in a higher probability of success (as is true in our example). Either of these decisions could be preferable to an insignificant increase in expected benefit. Political considerations could also easily outweigh a small advantage in expected payoff. By knowing the consequences of each of the alternatives, he is free to incorporate any qualitative considerations he feels are important.

After the decision maker chooses a decision option (in our example, dropping P3) the code prints out corresponding full numerical output data at the terminal and creates the graphical output file. The terminal output includes the expected value (7.86 arbitrary units) and standard deviation of the payoff as well as the probability of success that had been printed previously. In addition the output shows two useful measures of our R&D sequence, i. e., the expected value of perfect information (EVPI) and the expected value of sample information (EVSI). Lastly the output indicates that if a final decision were forced

Fig. 10-63. Decision tree for three programs or technologies.

now, technology P1 would be selected, yielding an expected payoff of only 7.59, even though all R&D money is assumed allocated to P1.

In the case referred to here EVSI is positive. However, it is possible for it to become less than zero, and in fact this has occurred in some example cases we have run. A negative EVSI indicates that the R&D effort directed towards three parallel programs or technologies has less value than immediately reducing the program to one candidate. Such a result indicates that a selection decision was overdue. The value of EVPI printed out is a measure of the benefit one might realize from more precise forecasts of success and failure.

The detailed probability data printed at the terminal relate both to the probabilities of success and failure and to the statistics regarding the various decisions made at the decision nodes. The probability of survival and success for each program or alternate technology is equal to the product of the conditional probability of success, given

survival to the final stage, times the probability of survival to the final stage (in our example, stage 4). From changes in the probability of survival to each stage we calculate the probability that each program or technology is dropped at each stage. These probabilities all appear in the output. The code computes these probabilities from the decision-tree data by summing the path probabilities corresponding to the appropriate event, such as survival of a specific technology to stage n.

The graphical output consists of the following for the selected option:

- Bar charts of the probability of being dropped for each program or technology at each stage.
- Bar chart of the probability of success for each program or technology given that it survives to the final stage, with the a priori probability of success presented for comparison.
- Bar chart of the probability of surviving to the final stage.
- The cumulative distribution of payoff.

Fig. 10-64. SAGE graphical output.

- A bar chart displaying the contribution of each program or technology to the expected payoff.
- A histogram showing the distribution of expenditures for each program.

When the first node is a decision node the output also contains a bar chart showing the expected payoff and the probability of success for each option (corresponding to dropping each of the four programs at the first decision node). Figure 10-64 shows examples of these graphics

We observe from Fig. 10-64(a) and (c) that in this example alternate technology P2 contributes nothing to the optimal expected payoff of 7.86 (arbitrary) units. Figure 10-64(d) confirms this by showing that only when one drops P4 or P1 is the expected payoff significantly reduced. We should therefore consider dropping P2 and P3 immediately. Putting all the R&D funding into a development program comprising P1 and P4 only will result in an expected payoff greater than the 7.86 we obtained here since no funds are wasted on an option that had virtually no chance of final selection. By trying different combinations of program alternatives and parameter values we have gained numerous other insights that both confirm and extend our intuition regarding various possible tradeoffs. In particular, sensitivity studies with respect to magnitudes of the budget for each stage, keeping a fixed mission budget as in Fig. 10-65, can indicate near-optimal points in time at which to drop redundant alternatives.

We are now generating the data base for further validation and sensitivity studies. We

Fig. 10-65. Sensitivity studies, timing of decisions.

will use the codes in the direct manner described above to derive the consequences of our assumptions regarding programmatic parameters and to determine optimal decision sequences and times. We will also use the codes to quantify the implications for programmatic variables of any postulated decision sequence and assumed probability of success. By using the code in this inverted mode we can help evaluate any proposed program-decision logic.

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Work performed under the auspices of the U.S. Department of Energy by Lawrence Livermore National Laboratory under Contract W-7405-Eng-48.