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MASTER Inertial Fusion Research

**Annual
Technical
Report
1979**

KMS fusion, inc.

Work Done Under
U.S. Department of Energy
Contract Number
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Annual Technical Report 1979

T.M. Henderson

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Inertial Fusion Research

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PRODUCTION SUPERVISOR: G.H. Wuttke

KMS fusion, inc.

**3621 S. State Rd.
Ann Arbor, MI 48106**

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Interferometric images of a single cryogenically cooled laser fusion target. The blue light is used for illumination in an optical system which characterizes the uniformity of the target. The green light is flare from a heating laser used to heat the target and vaporize the previously frozen DT fuel layer. The target as viewed here is held in a gaseous helium environment at $\sim 8^{\circ}\text{K}$.

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Inertial Fusion Target Technology

Fuel Shell Containers

Polymer Research

Cryogenics

Advances in Target Fabrication Technology

Physical and Chemical Analytical Support

Inertial Fusion Target Technology

Introduction

T. M. Henderson

At KMSF the development of inertial confinement fusion (IF) target technology has two goals: (1) to produce targets with the quality and characteristics required by current IF experiments and (2) to develop techniques to produce IF power reactor targets. The first goal entails not only target fabrication but the development of accurate characterization procedures and the presentation of targets for irradiation.

The vast majority of present targets are DT filled glass Microshell® containers mounted on fine fibers. Consequently, our principal efforts have been to improve the quality of glass shells, to understand the mechanisms for shell blowing, and to determine methods of presenting targets which either eliminate the fiber mount or which will allow us to estimate the extent of its effects.

Along with the activities in glass shell technology we expanded our capability to produce polymeric Microshell® targets and coatings. Previous disadvantages in using polymeric shells, such as low tensile strength, inadequate quality, and unacceptably high permeability were overcome by the development of new polymer systems and cryogenic processing technology.

The approach in the area of IF reactor target technology consisted of determining technologies potentially capable of meeting its high volume production requirements. The major thrust of this effort was to demonstrate the capabilities of semiconductor processing technologies to produce multishell configurations. Cryogenic technology was also committed primarily to solution of multishell problems, e.g., the production of cryogenic fuel layers within double shell targets. Our remaining efforts in the area of reactor target technology were to extend our glass shell fabrication to produce shells approaching the dimensions required by reactor targets.

1.1

Fuel Shell Containers

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All current and reactor target designs require the deuterium tritium fuel mixture to be contained within a thin-walled spherical container. The proposed container materials fall within a wide range of densities and average atomic number. At present the only materials which approach the stringent requirements of quality, strength, and gas-holding capability are glass and, to a lesser extent, polymers. Consequently, we concentrated our efforts on shells of these two materials with the goal of supplying the targets required for current feasibility demonstration experiments. These efforts emphasized our need to improve shell quality and to understand the mechanisms of shell formation.

1.1.1 Glass Shell Quality

The ability to achieve the high fuel densities required for thermonuclear fusion is greatly dependent on the quality of the fuel container, i.e., its sphericity, the uniformity of its wall, and the smoothness of both its interior and exterior surfaces. Of these geometric characteristics only sphericity is currently well in hand for both glass and polymeric shells. For glass shells a great deal of effort was spent in improving the latter two characteristics, which required extending our understanding of the mechanisms and chemistry of shell formation.

Wall Uniformity. Our previous experience indicates that many parameters affect wall uniformity.¹ A study was made of the effects of the following: (1) glass composition, (2) tower furnace temperature, (3) gel preparation technique, (4) gel treatment, and (5) furnace atmosphere.

To characterize the wall uniformity of a batch of shells, the most useful information is the percentage of shells which are highly uniform. The alternative of using a batch averaged uniformity can be misleading, since a few very nonuniform shells can inordinately lower the average, possibly masking significant changes. Consequently, wall uniformity was judged by the relative percentages of highly uniform shells in each batch. Uniformity of the individual shells was determined primarily by interferometry.

"Highly uniform" describes shells with wall non-uniformity (*WNU*) less than twenty percent [$WNU = (W_{max} - W_{min})/W_{min}$]; they are designated below as class A shells. To reduce analysis time, uniformity was estimated from a single interferometric examination, i.e., without rotating the shells and examining them in several planes. Using shells appearing "nearly perfect" from a single interferometric examination as the criterion for

"high uniformity", microradiography and roll sorting showed that the preponderance of these shells have the requisite wall nonuniformities (i.e., less than twenty percent).

Figure 1.1-1 shows the uniformity obtained with several gel compositions for constant furnace temperature and gel size. As can be seen, the uniformity of calcium free shells does not change greatly as the sodium or potassium content varies. This is also evident from the data in Table 1.1-1 taken for furnace runs in which gel sizes and furnace temperature were varied. However, for calcium containing gels the uniformity of the resulting shells is much more sensitive to the amount of non-silica components. Specifically, shells from the calcium containing gels highest in sodium (gels numbered 1100-135, 1100-151, 1100-75) have the lowest wall uniformity. The addition of 1 percent cobalt oxide to the potassium soda lime composition resulted in a significant improvement in wall uniformity. Cobalt was added to increase radiant heat transfer. The cobalt may have reduced the viscosity of the glass, but the mechanism by which this leads to increased uniformity is not yet clear.

The effect of furnace temperature on wall uniformity for a particular glass composition is shown in

Fig. 1.1-1. Wall uniformity for several gel compositions at constant particle size (75 μ m to 90 μ m) and temperature (1450 $^{\circ}$ C). Hash lines indicate 80% confidence intervals.

Table 1.1-1
Effect of Gel Composition on Shell Wall Uniformity

Gel No.	Shell No.	Oxide-Equivalent-Weight Percent in Gel						Gel Size (μm)	Temp. ($^{\circ}\text{C}$)	Class A Shells*
		SiO ₂	Na ₂ O	K ₂ O	B ₂ O ₃	CaO	CoO			
0112-113	T2R 152	80	20					90-106	1400	61-87
0112-26	262b	73	12	8	7			150-180	1500	69-83
0112-107	236b	73	3	17	7			125-150	1500	84-96
1100-75	TR 75	65	20		5	10		53-63	1150	0-13
1100-210	170i	70	6	9	5	10		53-63	1255	17-44
0112-123-AV	T2R 207c	70	6	9	5	10		125-150	1500	10-29
0112-13	232	72	6	10		11	1	125-150	1500	60-79

*80 percent confidence interval for mean

Figure 1.1-2. Within the temperature range of our investigation, uniformity increases as furnace temperature increases. The large increase in average uniformity between 1360° and 1435° may reflect a marked decrease in glass viscosity. Results with other gels and gel sizes show the same correlations of uniformity with temperature.

The effects of the remaining three considerations, gel preparation, gel treatment, and tower furnace atmosphere, were studied using a single potassium soda lime glass composition. The principal results are shown in Figure 1.1-3. Four gel batches with identical compositions but prepared differently were blown into shells under various combinations of gel treatments and concentrations of water vapor in the tower furnace tube. The effects of two gel preparation techniques were studied: (1) the effect of air drying the gel, as opposed to immediate vacuum drying, and (2) the effect of the inclusion of 2,4-pentanedione in the gel. The gels were hydrolyzed by exposing gel particles to air at 80 percent relative humidity at room temperature for five hours (thermogravimetric analysis indicated that no further reaction occurred after this time period). The

gel particles were vacuum dried after hydrolysis. The tower furnace atmosphere was changed by increasing or decreasing the water vapor content. Reproducible water vapor concentrations were achieved by either collecting shells dry or in boiling water.

Fig. 1.1-2. Effect of furnace temperature on shell wall uniformity for constant gel particle size (106 μm to 125 μm). (Gel composition identical to that of gel number 1100-210 given in Table 1.1-1.)

Drying Technique				Processing Conditions		
Gel	Air	Vac.	2,4-Pentanedione	Condition	Gel Treatment	Furnace Atmosphere
A		X	Yes	1	None	Dry Air
B		X	No	2	None	Air + H ₂ O(g)
C	X		Yes	3	Hydrolysis	Dry Air
D	X		No	4	Hydrolysis	Air + H ₂ O(g)

Fig. 1.1-3. Wall uniformity for several gel synthesis parameters and processing conditions. (Shells are from 125 μ m to 150 μ m gel, blown at 1500°C.)

The experiments on gel preparation were partially prompted by our concern that these two frequently employed gel preparation techniques are potential sources for segregation of gel formers: (1) slow removal of solvent by air drying the gel could promote the migration and crystallization of non-silica glass components, and (2) the presence of 1,4-pentanedione with its proven ability to form stable complexes with metal ions may also enhance formation of crystals of sodium, potassium, or calcium salts of 2,4-pentanedionates in the gel as solvent is removed.

Referring again to Figure 1.1-3 and the accompanying table, vacuum dried gels with no hydrolysis (untreated) produce more highly uniform shells than air dried gels. Hydrolysis of the gels appears to have lessened the effect of vacuum drying on samples containing 2,4-pentanedione. Similar results are obtained regardless of tower furnace shell collection conditions. The gel sample containing 2,4-pentanedione and air dried, sample C, is the worst overall with respect to production of uniform shells. Clearly, wall uniformity, which is fairly independent of processing conditions, can be enhanced by eliminating 2,4-pentanedione from the gel and by direct vacuum drying. Wall uniformity showed no consistent trends relative to gel treatment or water vapor content in the tower furnace.

technique may cause shells above $\sim 700 \mu\text{m}$ to appear more uniform than they are in fact. Shell runs 3 and 4 differed only by vacuum drying the gel. The WNU reproducibility was excellent. Other parameters which may influence shell size, such as furnace atmosphere and temperature profile, will be systematically investigated next year.

Surface Finish. Improvement in surface finish is necessary to satisfy both present and future requirements of IF target designs. Our earlier work demonstrated that surface defects on glass shells are related to the chemical durability of the glass, which led us to fabricate glass shells that are highly resistant to weathering in a normal laboratory ambience. The results obtained from varying the glass composition are shown in the scanning electron microscope (SEM) micrographs in Figure 1.1-4. Glass A has a relatively high alkali content (the chemical compositions are given in Figure 1.1-7); it is severely weathered after a four month exposure to ambient laboratory conditions. Glass B has both a lower total alkali content and a higher potassium content and exhibits less weathering than glass A. Glass C contains calcium in addition to the alkali and exhibits little weathering in the four month period.

Owing to our initial successes in improving

Table 1.1-2
Shell Wall Uniformity Change with Gel Particle Size

Run No.	1	2	3 and 4 (Vac. Dried)	5	6	7
Gel Size (Sieve Cut, μm)	212-250	350-300	300-355	355-425	425-500	500-600
Average Shell Size, μm	500	540	650	650	725	800
% Shells in Sample with WNU $\leq 20\%$	83	81	70 (68)	51	37	25

Wall uniformity obtained for several gel particle sizes is given in Table 1.1-2. The uniformity was assigned on the basis of the shell rolling technique.² Because interactions with the rolling surfaces produce irregular rolling trajectories not accounted for by the model, the WNU values of the larger shells can be in considerable error; in particular, the

weathering resistance by modifying the glass composition, we shifted our attention to developing methods for accelerating the weathering and quantifying the chemical durability of the glasses. We also investigated high temperature surface treatments for further improving chemical durability and weatherability, and the effects on surface

finish of blowing shells in an atmosphere high in water vapor content.

Accelerated Weathering and Tests of Chemical Durability. A relatively simple and effective means of accelerating the formation of surface features by weathering is to store the shells in a warm, humid environment. Figure 1.1-5 shows SEM micrographs of the exterior surfaces of glass shells of the same compositions as those in Figure 1.1-4 but maintained at 37° C and 95 percent relative humidity for ten days. Very large surface features of 10 μ m to 50 μ m were formed on the shells together with smaller features evident only at higher magnifications. The weathering of these shells is again most severe for glass A and least severe for glass C.

Because the formation of surface features by weathering is related to chemical durability, we developed an aqueous extraction test to determine the relative chemical durabilities of glass shells. This extraction test consists of agitating shells submerged in 93° C water of high purity (< 10 parts per billion [ppb] of the elements of interest) for 2.5 hours. The extraction is performed using materials and procedures that minimize contamination of the extract solution. Relatively small quantities (~ 10 mg) of glass shells are used in the tests, therefore only small amounts of material are extracted. This necessitated the development of analytical procedures for determining the small concentrations of the elements of interest in the extract solutions. By using atomic absorption and optical transmission methods and meticulous chemical preparation procedures, calcium and potassium were precisely and reproducibly determined in extract solutions down to a concentration of 10 ppb, sodium to a level of 30 ppb, and silicon to a level of 500 ppb.

Results of the extract test expressed as a percentage of the total elemental content of the shells undergoing extraction are shown in Figure 1.1-7. The absence of an elemental value in Figure 1.1-7 indicates that less than 0.5 percent of the element was extracted from the shells. The chemical durability results of Figure 1.1-7 are comparable to the weathering results of Figures 1.1-4 and 1.1-5, in that glass A has the highest percentage of sodium and potassium extracted, and glass C, the lowest. However, glass C is observed to have significant amounts of calcium and silicon extracted. This

Glass A shells after four months of ambient weathering.

Glass B shells after four months of ambient weathering.

Glass C shells after four months of ambient weathering.

Fig. 1.1-4. SEM micrographs of representative glass shells of differing compositions after four months of ambient weathering. (Chemical compositions of the shells are listed in Fig. 1.1-7.)

Glass A shell, untreated and weathered.

Glass A shell, SO₂-treated and weathered.

Glass B shell, untreated and weathered.

Glass B shell, SO₂-treated and weathered.

Glass C shell, untreated and weathered.

Glass C shell, SO₂ treated and weathered.

Fig. 1.1-5. SEM micrographs of untreated, representative glass shells after 10 days of accelerated weathering. (Chemical compositions of the shells are listed in Fig. 1.1-7.)

Fig. 1.1-6. SEM micrographs of sulphur treated, representative glass shells after 10 days of accelerated weathering. (Chemical compositions of the shells are listed in Fig. 1.1-7.)

Fig. 1.1-7. Results of chemical durability extract tests expressed as a percentage of the total elemental content extracted by the aqueous solution. Absence of a value indicates that less than 0.5 percent of the element was extracted.

would indicate that glass C is prone to dissolution, i.e., the dissolving of a complete layer of glass rather than the exchange of ions in the glass with ions in the aqueous solution.

Effects of High Temperature Surface Treatment on Weathering and Chemical Durability. To obtain additional improvement of the chemical durability (and thereby the weatherability) of glass shells, we investigated a sulfur treatment, in which a hot glass surface is exposed to a mixture of sulfur dioxide, air, and water vapor.³ This treatment dealkalizes a thin surface layer of the glass, producing a dense, silica rich layer. This "compacted" layer is chemically durable and weather resistant. This treatment was applied to shells of the glass compositions designated A, B, and C in Figure 1.1-7. The shells were held for 1.5 hours at 475° C in the SO₂ mixture described above.

The effect of sulfur treatment in reducing weathering of the above glasses is seen in comparing the SEM micrographs of Figure 1.1-6 with those of Figure 1.1-5. Both groups of shells were subjected

to the same accelerated weathering for ten days. The sulfur treated shells show smaller and fewer surface features than the untreated shells. Glass A shells showed the greatest improvement in surface finish as a result of sulfur treatment. Indeed, sulfur treated glass A shells exhibited the best surface finish of any of the glasses, treated or untreated, after accelerated weathering.

The effect of sulfur treatment on chemical durability is shown in Figure 1.1-7, where the percentage of material extracted by hot, aqueous solutions from treated and untreated shells is compared. Again glass A is observed to have benefited most from the sulfur treatment. The sodium extracted from glass A is lower than any of the other glasses, which is consistent with the weathering results. Sulfur treated glass C shells have a greater percentage of sodium extracted than untreated shells, but much smaller percentages of potassium, calcium, and silicon. These sulfur treated glass C shells should be much more resistant to dissolution than the untreated glass C shells. Sulfur treated glass B shells have more sodium and potassium ex-

(A) Dry Air

(C) Type II Feature

(B) Type I Feature

(D) Type III Feature

Fig. 1.1-8. SEM micrographs of inner shell surfaces showing representative surface features associated with dry air (A) or air with high water vapor content (B, C, D) in the furnace during blowing.

tracted than the untreated glass B shells. Nevertheless, the sulfur treated glass B shells are somewhat more resistant to weathering than the untreated glass shells, possibly a result of structural changes in the surface layer during "compaction."

Because of these promising results, high temperature surface treatments will be investigated more fully in the future.

Effects of Tower Furnace Atmosphere. A definite correlation between the occurrence of surface features on glass shells and tower furnace atmosphere was demonstrated. We studied the effects of water vapor in the furnace for several gel preparations, gel compositions, and gel treatments. Changes in tower water vapor content had by far the greatest effect on surface smoothness.

When collected dry, calcium containing shells (e.g., glass C of Figure 1.1-3) generally show very few interior or exterior surface features. Those occurring are submicron in size. Figure 1.1-8A shows a typical feature. Shells from a gel containing 2,4-pentanedione had somewhat fewer surface features than shells made from gels not containing 2,4-pentanedione; the reason for this is not clear. Calcium free shells generally show many large ($> 1\mu\text{m}$) amorphous surface features.

When shells are blown in a furnace environment high in water vapor, the shell surface features are considerably more numerous than on the shells blown and collected in dry air. Three basic types of features are seen. Type I features, shown in Figure 1.1-8B, are irregularly shaped defects ranging from $\sim 0.1\mu\text{m}$ to $\sim 3\mu\text{m}$. Type II features, shown in Figure 1.1-8C ($\sim 0.6\mu\text{m}$ to $\sim 4\mu\text{m}$), are mound-like in shape and generally very high in calcium and sodium. Hexagonal crystals ranging from $\sim 0.6\mu\text{m}$ to $\sim 3\mu\text{m}$, sometimes appearing as multiples, comprise Type III. An example is shown in Figure 1.1-8D. these crystals are also very high in calcium and sodium.

Gel preparation, gel composition, and gel treatments affected the type and frequency of features seen on all shells, including those made with high water vapor content in the furnace. Shells from calcium containing gels (containing 2,4-pentanedione) did not have hexagonal crystals. The mounds shown in Figure 1.1-8C were by far the most common features, appearing in patches on the calcium containing shells. Shells from calcium containing gels (without 2,4-pentanedione) showed all three feature types. The hexagonal crystals from hydrolyzed gels were much larger than those on shells from unhydrolyzed gel. Hydrolyzed, air dried gel (without 2,4-pentanedione) produced shells with a preponderance of such crystals.

Shells from calcium free gels displayed mainly features I and II. When the resulting shell was high in alkali, the features were large and numerous. When the resulting shell was very low in alkali, the features were small and scattered.

In summary, both low tower furnace water vapor content and use of a low alkali glass decreases the size and number of glass shell surface features.

1.1.2 Strength

Glass strength is vital to containing high pressure fills and to rapid filling. The strength of individual shells is dependent on treatment, geometry, (e.g., wall thickness, aspect ratio), flaw density, process and storage environments, and glass composition. As a first step toward investigating this complex subject, we examined failures of two types of glass shells under compressive and tensile loads. The glass compositions chosen for testing were high silica (> 99 percent) and sodium silicate (14 percent Na_2O and 86 percent SiO_2 by weight).

Shells selected for testing had good wall uniformity based on qualitative assessment using Twyman-Green interferometry.⁴ Wall thickness measurements were also made by interference microscopy, while diameters were determined by image shearing. The characterized shells were placed in a metal holder containing hemispherical cavities and sufficiently well covered to prevent their escape while allowing them to be pressurized. The shell holder was inserted into a pressure vessel, and the shells were pressurized with nitrogen in stages. After each stage (or pressure increment) the holder was removed and the shells which failed were counted. Forty-two and fifty-three shells were tested for the silica and sodium silicate shells respectively. The failure pressure was taken as the midpoint between the last (highest) pressure step which a shell survived and the next higher pressure step at which it broke. The results were treated using a buckling model for an ideal thin spherical shell,⁵

$$P_B = \frac{8E}{\sqrt{3(1-\nu^2)}} \left(\frac{w}{OD} \right)^2$$

where P_B is the pressure at which failure occurs, ν is Poisson's ratio (taken to be 0.21), E is Young's modulus, w is shell wall and OD is outer shell diameter. There was considerable scatter in the data, i.e., shells of a given aspect ratio failed over a fairly wide pressure range. Part of this scatter is undoubtedly due to the use of shells which vary in uniformity, since observations were made of a single plane only, i.e., shells were not rolled. Also, the pressure increments were too great for the high silica shells, which resulted in large uncertainties in breaking pressures. This is reflected in a rather large

Table 1.1-3
Young's Modulus of Glass in Shells

Sample (wt. %)	Calculated from Glass Shell Buckling	Relative Standard Deviation	Literature Value ⁶
~ 99% SiO ₂	72.4 GN/m ² (10.5 × 10 ⁶ psi)	22%	73.1 GN/m ² (10.6 × 10 ⁶ psi)
86% SiO ₂ 14% Na ₂ O	43.3 GN/m ² (6.28 × 10 ⁶ psi)	9%	43.4 GN/m ² (6.29 × 10 ⁶ psi) *

*Reported value is for 84% SiO₂ and 16% Na₂O.

standard deviation for the calculated Young's modulus. Despite scatter in the data, calculated averages of Young's modulus for the glasses used in the two shells were found to give remarkably good agreement with literature values as shown in Table 1.1-3. This agreement strongly supports the idea of shell breakage by buckling during compressive loading. Buckling failure is also consistent with the conclusions reached by J. Milewski of Los Alamos National Scientific Laboratory; they are based on failures of carefully characterized individual shells under isostatic compressive loads.⁷ Thus, a knowledge of Young's modulus and shell geometry can be used to calculate an allowable upper limit for the pressure difference ΔP across shell walls during filling. This limit will be lowered by shell imperfections and by the heating required during the fill.

Once filled, the shells must survive a tensile load, given by the hoop stress equation for a spherical shell, $\sigma_t = (P/4) (OD/w)$, where P is the fill gas pressure, OD/w is the aspect ratio, and σ_t is the tensile strength of the glass. Glass tensile strengths are known to vary widely, with a strong dependency on environment and handling, which determines flaw frequency and depth. The two glasses used for the compressive strength testing are sufficiently different in composition that they are expected to show different degrees of sensitivity to environmental conditions. This fact will make them particularly useful in the next phase of our study, which is designed to quantify these effects.

Tensile testing has begun on the same types of shells tested in compression. Selected, characterized

shells were pressurized in the same apparatus used for the compressive tests. Helium was the test gas since it permeates the shell walls rapidly. The shells were heated to 200°C and pressurized in stages to avoid failure by buckling. After a given pressure was reached, the gas external to the shells was vented to produce the desired ΔP across the shell wall. Those shells which failed were recorded and another pressure step was taken. Ideally this process would continue until all the shells break, but approximately 50 percent of the high silica shells (the only ones tested thus far) survived the maximum pressure achievable with our present apparatus, 313 atmospheres (31.7 MPa). Therefore, we unfortunately are not able to calculate an average breaking (hoop) stress for the shells. Based on the aspect ratios of the surviving shells, tensile stresses of $\geq 200,000$ psi (1.4 GPa) were observed for at least 20 percent of the sample. In future tests the applied pressure range will be extended and high aspect ratio shells will be used where possible.

1.1.3 Large Shell Studies

We completed the first phase of a systematic study of the production of large glass shells from gel particles processed in the vertical furnace. The furnace temperature was held at 1500°C and sequentially fed with gel particles from six different sieve cuts in the range 212 μm to 600 μm . Shell size and wall uniformity were determined for shells with aspect ratios greater than 16 (these shells float in iso-propanol).

The yield of the above shells decreased with increased particle size from ~50 percent for the smallest particles (212 μm to 250 μm) to ~3 percent for the largest (500 μm to 600 μm). The average diameter increased from ~500 μm to ~800 μm , which is about the maximum that can be produced under this particular set of furnace conditions. The decrease in the uniformity of the shells as the gel particle size increases is shown in Table 1.1-2.

1.1.4 Shell Formation Model

Analysis of gases within the shells was used as an aid in understanding the glass shell-blowing process. The total pressure of these residual gases was determined by the technique known as gas-entrapment-in-viscous-fluid (GEVF).⁸ Relative concentrations of components of the residual gases were determined by gas chromatography. A chromatographic column consisting of a concentric mole sieve and Porapak P columns was used to separate carbon dioxide, oxygen, and nitrogen. Water was not eluted from this column under the conditions necessary for analysis of the other gases. To measure water vapor it was necessary to use a Porapak N column, which separated water, carbon dioxide, and oxygen plus nitrogen, but did not separate the latter two gases from each other. The relative concentration of water for those shells in which it was detected is based on the results from the Porapak N column, but the other gas components in the shells are reported using data from the dual (CTR) column adjusted for the amount of water.

The gas analysis and pressure test results are recorded in Table 1.1-4. Since the concentration of water vapor in the furnace tube during shell blowing influences several shell properties, a notation is made in the Table to indicate whether the water in the shell collection vessel was heated by the output of the furnace "ambient," or whether additional heating was used to raise its temperature to 95°C. In one instance, sample I, shells were collected on a dry glass plate which was attached to the bottom of the furnace tube.

The gas analyses show that carbon dioxide is the major constituent in shells produced from metal-organic gels, the expected result as organic residues are present in the gel-starting materials.

Shell samples E and F, which had relatively low concentrations of CO₂, were made with no organic constituents nor any "added" blowing agent. The possible sources of CO₂ in the shells and the effects of gel treatment (e.g., hydrolysis) on gel chemistry and shell formation has been studied in detail for the sodium silicate glass system under NASA contract No. NAS8-33103. The results are reported in the 1979 quarterly reports prepared under that contract.

It was observed above that the interior and exterior surfaces of shells blown in a furnace with a high water vapor content show an increase in the size and/or frequency of surface features. The features are generally water soluble, so that with proper collection and cleaning procedures the exterior surface features can be removed. Interior surfaces are obviously not as accessible; consequently, considerable attention should be given to preventing their formation.

The magnitude of the changes seen on the interior surfaces are dependent on glass composition. Boiling the water in the collection vessel produces water vapor which evidently permeates the shell wall. Note that no water is found inside the shells when the water vapor present in the column is less than that produced by boiling (see Table 1.1-4). When the shell cools, the water is trapped and can react with the glass to form metal hydroxides, which is the reverse of the reactions that occur when hydrated oxides are heated to a high temperature. The important reactions can be represented by the equation:

Reaction of the water with the glass surfaces on cool down reduces the residual pressure in the glass shells; this is seen in Table 1.1-5 for three glass systems (A, B, and C). A possible mechanism for this reduction is discussed below. Samples A, B, and C differ in composition. Each sample was tested with different amounts of water vapor present in the furnace tube during shell blowing. Water vapor comprises about 30 percent of the residual gases in shell samples A and B, which were blown with boiling water in the shell collection vessel. This is probably only a fraction of the amount present during shell formation, i.e., it is only the amount in

Table 1.1-4
Residual Gas Analysis of Shells

Sample	Shell-Blowing Conditions *	Relative Concentration (% Volume)			Residual Pressure (atm)	Gel Conditions
		CO ₂	O ₂	N ₂		
Na ₂ O · SiO ₂						
A	Ambient H ₂ O	71	19	10	0.19	Gel made from aqueous sodium silicate. Gel made from aqueous sodium silicate.
B	Ambient H ₂ O	79	14	7	0.175	
C	95 °C H ₂ O	68	15	16	0.15	
D	95 °C H ₂ O	81	12	7	0.15	
E	Ambient H ₂ O	21	26	53		
F	Ambient H ₂ O	21	22	58		
Na ₂ O · K ₂ O · B ₂ O ₃ · SiO ₂						
G	Ambient H ₂ O	85	11	4	0.20	Repeat runs
H	Ambient H ₂ O	88	8	4	0.20	
Na ₂ O · K ₂ O · CaO · B ₂ O ₃ · SiO ₂						
I	Dry	63	16	21	0.23	Same glass composition; different gel batches and particle size.
J	Ambient H ₂ O	56	20	24	0.21	

*"Shell blowing conditions" refers to collection technique: dry means no water was in the collection vessel; ambient H₂O means the collection vessel contained water, whose only source of heat was the hot, ambient environment of the furnace tube; 95 °C H₂O means the vessel contained water and was heated by an additional source to raise the temperature of the water to 95 °C.

**All samples except C and D were analyzed for water. In no instance was water found.

excess of that needed to react with the metal oxide species available on the glass surfaces. In the case of glass shells which are formed in furnace environments having low water vapor content, water initially in the gel may participate in the blowing process, but subsequently it either diffuses through the shell walls and/or is totally consumed by the glass surface and is thus not found in residual gas analysis (e.g., the shells in Table 1.1-4). This can also explain the low incidence of surface features on shells blown in dry air.

The data in Table 1.1-5 also reveal that the CO₂ to N₂ ratio is much lower for shells made in an at-

mosphere high in water vapor content. A change in the ratio owing to a difference in permeation rates caused by the presence of water appears unlikely. A more likely explanation for these low values is that carbon dioxide reacts with metal hydroxides on the interior surfaces to form metal carbonates or bicarbonates during or after shell cooling.

Further evidence that water permeates the shells during formation was obtained by analyzing shells made in a low water vapor furnace environment which were subsequently passed through a furnace environment containing a high concentration of water vapor. Gas analyses and pressure tests for

Table 1.1-5
Effect of Water Vapor in Furnace on Shell Residual Gases

Sample	Shell-Blowing Conditions*	Relative Concentration (% Volume)				Residual Pressure (atm)	O ₂ /N ₂	CO ₂ /N ₂
		CO ₂	O ₂	N ₂	H ₂ O			
A	Ambient H ₂ O	85	11	4	0	0.20	2.7	22
A	Boiling H ₂ O	40	15	14	31	0.031	1.1	2.9
B	Dry	63	16	21	0	0.23	.80	3.1
B	Boiling H ₂ O	25	14	28	33	0.096	0.48	.90
C	Dry					0.23		
C	Boiling H ₂ O					0.075		

* All runs were in air at ~ 1500° C.

these shells are reported in Table 1.1-6. Samples 1 and 2 are shells made from the same gel composition and frit particle size and blown at the same furnace temperature. The reduction in the residual pressures is consistent with the observations for

shells produced in the presence of high water vapor concentrations. The lower pressure on repassing sample 2 in the furnace environment compared to that from repassing sample 1 may be due to differences in water vapor concentrations. The de-

Table 1.1-6
Permeation of Water into Shells

Sample	Furnace Conditions	Relative Concentration (% Volume)				Residual Pressure (atm)
		CO ₂	O ₂	N ₂	H ₂ O	
1	Ambient H ₂ O	85	11	4	0	.20
	Repass of 1 into Boiling H ₂ O	77	14	9		.115
2	Ambient H ₂ O	88.2	8	4		.20
	Repass of 2 into Boiling H ₂ O	54	11	4	31	.06
3	Boiling H ₂ O	40	15	14	31	.03
	Repass of 3 Collected Dry	76	18	6	0	.23

* All runs were made in air at ~ 1500° C. Samples 1, 2, and 3 were made from the same gel batch and particle size (212 μ m to 250 μ m sieve cut). Water analyses were not made for two of the samples owing to insufficient quantity.

creased partial pressures of oxygen and nitrogen in the reheated shells of samples 1 and 2 may be due to the expansion of the shells caused by infusion of water vapor at a faster rate than can be compensated for by the permeation thru the shell by the air components. Water is not only capable of permeating the glass at fairly high rates at these temperatures, but it also can reduce glass viscosity significantly and consequently increase the rate of expansion of the glass shell. Whether the oxygen to nitrogen ratio increases or decreases upon repassing the shells should largely be a function of the oxygen partial pressure differential between the shell interior and the ambience, since the rate of nitrogen permeation is expected to be much slower. The larger decrease in CO₂ partial pressure upon repass of the shells, as compared to oxygen or nitrogen, is probably due to reaction of CO₂ with freshly formed metal hydroxides on the inner surfaces of the shells. Scanning electron microscopy has shown an increase in the size and frequency of surface features when sample 1 was repassed into boiling H₂O. Analysis of the surface features on these and other glass shells by EDXS has shown them to be high in alkali. When calcium is present in glass shells, some of the surface features are also calcium rich.

To determine whether water can be removed from the shells by reheating in a water deficient environment, a sample of shells originally prepared in the presence of high water vapor concentration was passed through the furnace in dry air (sample 3 in Table 1.1-6). No water was found in the residual gases of the resulting shells, but a large (approximately fourteenfold) increase in CO₂ partial pressure was observed. Since the furnace environment contains very little CO₂, the increase is caused by the release of CO₂ through decomposition of carbonates or bicarbonates on the inner surface of the shell. Carbonates or bicarbonates cannot reform, since water has been removed. The CO₂ partial pressure, 0.18 atm, is comparable to that found in shells initially made in an environment very low in water vapor concentration, e.g., sample 1 (0.17 atm). The shells from sample 3 appeared to have undergone phase separation and possibly devitrification after the exposure to the dry furnace environment, making it difficult to assess which changes in surface condition were due to decomposition of carbonates or bicarbonates. The oxygen

pressure increased relative to that found for shells formed in ambient air (within the combined uncertainty of the total pressure measurements). The nitrogen partial pressure, although increasing about three-fold, was still a factor of ~20 less than that in air. Thus, nitrogen appears to be diffusing through this glass composition, but at very low rates.

The role of carbonates and bicarbonates in glass formation, in the shell blowing process, and particularly in post shell blowing process, can be fairly complex. Some of the possible reactions involving CO₂ which can occur are shown schematically in the following set of equations:

For illustrative purposes, only alkali metal carbonates are shown, but similar behavior can be expected for the alkaline earth metal carbonates. Under commonly encountered glass melting conditions, these reactions are irreversible due to loss of water and CO₂. However, during the formation of glass shells, the gases are trapped in the central void, and the reverse reactions can occur at an appropriate temperature. The fact that the CO₂ level in repassed sample 3 (Table 1.1-6) rose to quite high levels suggests that the reverse of Reaction scheme I may not be energetically favorable. In Reaction schemes II, III, and IV, it will be noted that water is a necessary component for reversal. Since shells from sample 3 were repassed in a water free environment, the water in the shells (free and chemically bound) apparently diffused out at high temperatures and the reverse reactions could not occur upon cooling. Finally, although the reverse Reactions II-IV can explain the growth of surface features on inner glass shell surfaces, excess water can also form metal hydroxides by cleavage of M-O-Si bonds, which can grow in size without CO₂ present.

Conclusions. For shells blown from metal organic gels, CO₂ appears to be the major blowing gas in the early stages of shell formation. Although the CO₂ can react with glass components during the cooling of the shell, the temperature at which this reaction starts is probably sufficiently low that the glass has become too viscous to contract appreciably from the resultant decrease in the shell's internal pressure. At high temperatures, where such reactions do not occur, the permeability of CO₂ is sufficiently low that permeating CO₂ does not contribute to changes in shell size. The role of water as a blowing gas is more difficult to describe quantitatively, since it apparently can react readily with metal oxides during shell cooling. Water is observed as a residual gas only when there is a large quantity of water vapor in the tower furnace. Oxygen appears to permeate shells more rapidly than nitrogen, resulting in higher oxygen to nitrogen ratios than found in air. Nitrogen's presence in all glass shells, sometimes appreciable (as high as 38 percent of ambient air pressure), suggests its incorporation in the shells at a very early stage in the formation process. A possible mechanism for this incorporation is the existence of a very large, thin walled, transient state of the shell during which ambient gases can permeate rapidly. This stage is analogous to that described by workers at Lawrence Livermore National Laboratories⁹ for glass shells derived from aqueous silicate solution droplets. A porous gel-glass intermediate which traps air in its pores by rapid formation of a glass skin is another possible mechanism. Oxygen, but not nitrogen, subsequently diffuses into or out of the shell depending on the oxygen partial pressure in the furnace environment. The glass shell will expand or contract in response to this gas flux only if the glass viscosity is sufficiently low to permit flow to occur in the time available during the shell's transit through the furnace. Glass viscosity will generally tend to increase during the later stages of shell formation due to the loss of boron and alkali, and in the absence of a viscosity-reducing agent such as water. The general magnitude of such losses has been described previously.¹⁰

The partial shell blowing model just outlined accounts for most of the experimental data collected to date on shell gas analysis. More study is required to understand the early steps in the transformation of a metal organic gel particle into a glass shell.

1.2

Polymer Research

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Radiative energy losses in implosion experiments are strongly dependent upon the average atomic number ($\bar{Z}$) of the ablator material. A shift to shorter wavelengths would shift the head of the plasma conduction front to higher densities and thereby may increase radiation loss.¹ To reduce these losses, most recent complex IF target designs^{2,3} use low- $\bar{Z}$ materials, e.g., Be, Li, LiH, and organic or organometallic polymers for the fuel container and/or for the surrounding layers.

Compared to most low- $\bar{Z}$ materials, polymeric shells and coatings have a much wider range of $\bar{Z}$ and permeabilities to fuel and diagnostic gases. Adjustable quantities of high- Z elements can be included in a polymer matrix and used to reduce electron preheat. Deuterium and tritium can be substituted for hydrogen in polymers to be "burned" as additional fuel. The high fuel gas permeability of polymers greatly reduces filling time, which results in substantially smaller tritium inventories for a fusion power reactor.⁴ In addition, polymeric target components should produce fewer radioactive by-products,⁵ while biologically based polymers such as cellulose would constitute a renewable resource.

Previous drawbacks to the use of polymeric shells, such as low tensile strength, inadequate uniformity, and unacceptably high gas permeability, have been significantly reduced through the introduction of new polymer systems and cryogenic storage technology.⁶ Consequently, we have continued to develop further polymeric Microshell® fuel container fabrication, to synthesize applicable new polymers, and to develop polymeric coating technology.

1.2.1 Polymer Synthesis

Whenever possible, materials used in IF target designs are selected from those which are available commercially. Often, however, new materials are required in order to satisfy both theoretical specifications and practical considerations. We

have concentrated on two types of unavailable materials: totally deuterated polymers and low- Z polymers with chemically bound high- Z elements.

Poly (p-iodostyrene) was synthesized and formed into shells which were used in target interaction experiments in 1978. Deuterated poly (p-iodostyrene) was synthesized during 1979 by the following pathway:

A multi-step pathway is unavoidable owing to the unavailability of deuterated intermediates. Care was taken in all steps to maximize product yield.

Two iodine substituted para-xylene derivatives were synthesized by the following method:

Both reactions were accomplished with greater than 70 percent yield. These compounds were used to produce iodine-containing polymer films by glow-discharge polymerization, described below in the polymer coatings section.

The iodination of (2.2) paracyclophane by a

similar technique was attempted for Los Alamos National Scientific Laboratory.

Iodination occurred, but was accompanied by polymerization, which rendered the product insoluble.

There is considerable interest in polymers which contain both low- Z atoms, such as carbon and hydrogen, and high- Z atoms, such as tungsten or tantalum.^{7,8} Our work in this area has concentrated on organometallic polymers of hafnium ($Z=72$.)

Dicyclopentadienyl hafnium dichloride is produced by the reaction of sodium cyclopentadienide with hafnium tetrachloride.⁹

A direct copolymerization of this compound with styrene produces polymer with only traces of hafnium included.

Dicyclopentadienyl hafnium dichloride can be chemically reduced to form hafnocene by reaction with sodium naphthalide in THF.¹⁰ This reaction was repeated at KMSF.

Direct copolymerization of hafnocene with styrene produced polymers with only traces of hafnium included.

Another approach to the synthesis of a hafnium containing polymer involves the attachment of a polymerizable vinyl group to a cyclopentadienyl ring of hafnocene. This is the approach taken in the synthesis of the iron and ruthenium analogs, vinyl ferrocene, and vinyl ruthenocene.

The first step in this approach is to accomplish a Friedel-Crafts acylation of hafnocene.

The NMR spectrum of the product was very complex, with multiple peaks at 2 ppm, which indicates a decomposition and polymerization of hafnocene.

Two alternate acylation methods were also attempted, using BF_3 and HF as Lewis acids. Very similar results were obtained.

Another approach is to acylate sodium cyclopentadienide before reaction with hafnium tetrachloride.

This reaction was attempted. The product is a solid which is insoluble in all solvents attempted, making NMR analysis impossible.

Acylation of dicyclopentadienyl hafnium dichloride was also attempted. This also resulted in the destruction of the cyclopentadiene ring.

The failure of attempts to acylate cyclopentadienyl derivatives of hafnium indicates that bonding of cyclopentadiene to hafnium is weak, and that there is not sufficient π bonding to permit reactions typical of aromatic systems.

1.2.2 Polymeric Coatings

Current target designs and target interaction experiments require low $\bar{Z}$ layers of a thickness rang-

ing from hundredths of microns to several hundred microns. In general, layers must be uniform to better than 5 percent, with surface finishes approaching 100 Å. In some instances it is necessary to vary the density and $\bar{Z}$ of the layer either abruptly or smoothly.

To satisfy this broad range of requirements we are evaluating several coating techniques. This year we have investigated glow discharge polymerization (GDP), Parylene vapor deposition, and solution-coating. The major effort was expended in characterizing the coatings obtained with GDP. Preliminary experiments were conducted with solution coating of poly(vinyl alcohol), and a chamber for Parylene deposition was completed.

Glow Discharge Polymerization. Based on our previous experience with GDP deposition¹¹ we decided that a helical coil rf induction glow discharge would provide the cleanest and most versatile environment for organic film deposition. After initial studies using a vertical deposition chamber design, a horizontal configuration was finally chosen for utilization as our preliminary system (see Figure 1.2-1).

Fig. 1.2-1. System for glow discharge polymerization development.

The deposition chamber is a 30 cm x 2.54 cm diameter pyrex tube connected to inlet and outlet ports via quick connect fittings. This straightforward design allows for maximum versatility in testing alternate coil and chamber configurations. A Lepel rf induction generator operating at 13 MHz

Fig. 1.2-2. KMSF Metallomeric System III

drives the glow discharge through a helical coil. An argon sputter etch is used to preclean the substrates and initiate the rf glow discharge.

The energetic species in the glow discharge strongly affect the deposition process. Substrates placed outside the visible glow obtain negligible polymer coating. Substrates placed within the energy input (coil) region receive excess ultraviolet radiation, and their organic coatings appear strongly stressed and highly discolored. In addition, turbulent flow conditions are desirable for uniform coverage of nonlinear substrates.

Based on the results of a parametric study of the GDP system, an advanced, organometallic deposi-

tion system, the Metallomeric System III, has been designed. Parts for phase 1 of this system have been procured, and their incorporation into a bell jar type deposition chamber is shown schematically in Figure 1.2-2.

The principal features of the Metallomeric System III are: (1) "blind" rf energy input region beyond line-of-sight of the deposition area to preclude UV degradation of the coatings; (2) dc sputter head for codeposition of metallic elements to obtain variable compositions and gradients; (3) three channel, proportional flow controller for copolymerization of organic/organometallic vapors and other species such as Ar, H₂O, D₂, etc.; (4) aux-

iliary, parallel-plate, capacitive discharge fixture for deposition of low density, gas phase, nucleated polymer "snow"; (5) minishaker type levitation system for bulk coating of unmounted targets; (6) an inlet port for the parylene pyrolysis system; (7) eight port, feedthrough collar for capacitance type pressure transducers, thermocouples, deposition monitors, and cooling lines (when needed).

The p-xylene (C_8H_{10}) molecule has been the principal organic compound used as a precursor in our glow discharge polymerization studies. This compound has a low ionization potential, which permits a stable glow discharge under a wide variety of system pressures, precursor flow rates, and reaction chamber configurations.

IR absorption spectrum of GDP p-xylene

IR absorption spectrum of GDP iodo-dimethyl benzene

Fig.1.2-3. IR absorption spectra of Glow Discharge Polymerized poly (p-xylene) and poly (iodo 2,5 dimethyl benzene).

A typical IR spectrum of the resulting GDP poly(p-xylene) is shown in Figure 1.2-3. The broad absorption bands are typical of the highly cross-linked structure of GDP polymers. Also typical is the continuous oxygenation of these coatings during exposure to air.¹² Chemical analysis of our GDP poly(p-xylene) has shown 1.6 atom percent O_2 after five days and 4 atom percent after 100 days.

Three methods are being investigated for use in "seeding" small amounts of high-Z elements into a GDP low-Z coating. In the first method, we have

substituted an iodine atom for hydrogen in the p-xylene molecule to form 2,5-dimethyl iodobenzene. This molecule undergoes glow discharge polymerization in a manner somewhat similar to p-xylene, producing a film colored a very deep red. Its physical properties appear similar to its precursor, as does its IR spectrum as shown in Figure 1.2-3. An energy dispersive x-ray spectrograph (EDXS) of a 10-20 μm thick film clearly indicates the presence of iodine. Chemical analysis shows a composition of $C_{52}H_{47}I_{1.25}$ with negligible traces of oxygen even after 30 days exposure to air. A scanning electron microscope micrograph of this polymer is shown in Figure 1.2-4.

Fig. 1.2-4. SEM micrograph of GDP poly(iodo 2,5 dimethyl benzene).

A second method of high-Z seeding of a low-Z matrix is ion implantation. Using the GDP development system shown in Figure 1.2-1, argon has been implanted in a silicone grease during a 15 minute argon glow discharge. Figure 1.2-5 shows a scanning electron micrograph of the moderately cross-linked structure after a 15 day anneal at 70° C. Amazingly, the EDXS scan shows a persistent argon signal, which indicates tightly bound argon within the organic matrix.

The third method of high-Z incorporation into organic coatings relies on sputter deposition during glow discharge polymerization. Initial investigations using the GDP development system suggest that the high input energy required for efficient

Fig. 1.2-5. SEM micrograph of argon "sputter etched" silicone grease.

sputter deposition will necessitate using a magnetically confined sputtering plasma or a separate sputter head somewhat removed from the polymer deposition region. This technique will be further investigated using the appropriately designed Metallomeric system.

Adhesion of GDP coatings to most substrates is historically very strong, approaching that of chemical bonding.¹³ However, 1 μm to 6 μm GDP coatings on flat glass substrates have been found to buckle upon extended exposure to air. This is apparently due to internal stresses created in GDP films during deposition.¹⁴ Upon exposure to air, water vapor can diffuse to the interface and release the expansively stressed coating from the substrate. Figure 1.2-6 shows a scanning electron photomicrograph of a 5 μm GDP poly(p-xylene) film which has buckled from the glass cover slip. In one area the buckling film has removed a portion of the glass substrate, indicating strong interfacial adhesion. Films thinner than $\sim 1 \mu\text{m}$ do not appear to have sufficient internal stresses to rupture the interfacial adhesion since they do not display pronounced buckling. Films thicker than 6 μm also do not buckle, perhaps because they are stronger or more resistant to moisture.

Several approaches were taken to reduce the delamination of GDP coatings. We found that the interfacial adhesion of these coatings was increased by ultrasonic precleaning of the substrates, by argon sputter etching of the substrates immediately

Fig. 1.2-6. GDP poly(p-xylene) film delaminating from glass cover slip substrate.

prior to deposition, and by not interrupting the glow discharge throughout the etch deposition process. We also found that increasing the hydrophobicity of the substrate decreases the interface's sensitivity to moisture. Internal stresses can apparently be lowered by running the deposition at the lowest input energy needed for a visually uniform glow throughout the chamber. Higher power input does not appear to increase the polymer deposition rate.

Organic GDP coatings have been deposited with mean thicknesses ranging from 0.02 μm to 125 μm . To obtain appreciable deposition rates requires that the substrate be located directly in the region of glow. Substrates coated within direct sight of the rf coil (energy input region) appear highly colored due to ultraviolet radiative degradation of the crosslinked polymer structure. Thickness uniformity studies using a surface profilometer have shown that this coloration has no direct relationship to the overall thickness of the applied coating.

Figure 1.2-7 illustrates a mounted target coated with 5 μm of GDP poly(p-xylene) in a glow region outside the line-of-sight of the energy input region. Under these conditions we have found that the bending of the thin fiber stalk due to internal stresses in the coating is a good qualitative indicator of axial thickness uniformity. Thickness uniformity studies indicate that GDP deposition is a line-of-sight/line-of-flow process similar to that of classical sputter deposition. These studies suggest that magnetic field confinement of the plasma and/or use of specially configured "monomer" inlet ports

Fig. 1.2-7. Mounted target (100 μm) before and after deposition of 5 μm thick GDP poly (p-xylylene).

may be necessary to achieve highly uniform coatings on spherical substrates.

Initial surface structure, deposition contaminants, and debris are the controlling elements in the final surface finish of the coating. We find the use of a horizontal system has greatly reduced any *in situ* contamination of the growing film. Figure 1.2-8 shows a high magnification SEM photograph of a mounted target coated with 3 μm of GDP poly(p-xylylene). The only recognizable surface features appear to be weathering structures common to uncleaned glass shells.

Surface features of a flat boron target, polished to 0.1 μm RMS surface finish and partially coated

Fig. 1.2-8. SEM micrograph of GDP poly (p-xylylene) on mounted glass target.

with 3 μm of GDP poly(p-xylylene), are shown in Figure 1.2-9. The high magnification photographs of the coated and uncoated surfaces show that the

Fig. 1.2-9. SEM micrographs of flat boron target with GDP poly (p-xylylene) coating (step coverage).

coating enhances the defects present on the substrate surface, as has been reported by Letts et al.¹⁵ We are investigating means of reducing this effect by introducing argon and oxygen continuously and/or periodically into the reaction chamber in order to preferentially sputter-etch the high points on the surface.¹⁶

Solution Coating. Coating targets by immersing them in a liquid would permit the rapid application of thick or thin coatings of a wide variety of organic or inorganic compounds. As the liquid coating hardened, surface tension forces could (in favorable cases) produce a layer having a uniform thickness and possibly a smooth, defect free surface. Hardening can be accomplished by polymerization, freezing, or solvent evaporation. The basic concept has been experimentally verified.

For demonstration, polyvinyl alcohol (PVA) was chosen as a convenient and potentially useful coating material. It is available in liquid form as an aqueous solution which can be hardened by dehydration. Isoamyl alcohol is a suitable dehydrating agent, as it will dissolve an appreciable amount of water but not PVA. To support a droplet of aqueous PVA solution during dehydration, the

Fig. 1.2-10. Solid sphere of PVA after dehydration with isoamyl alcohol.

density of the amyl alcohol was increased by adding a second liquid of higher density with which it is miscible. The droplet floats stably at the level of matching densities. The floating droplet assumes a spherical shape. Subsequent dehydration produces a solid sphere of PVA with a smooth surface, as shown in Figure 1.2-10.

When a glass shell was inserted into the PVA droplet, the sphericity of the dehydrated PVA was lost due to the buoyancy of the hollow glass shell. Further work is being undertaken to circumvent this problem.

Parylene Deposition. A Parylene vapor deposition system has been designed and constructed, based on discussions with Dr. R. Leipens of LANSL. Figure 1.2-11 shows the completed system, and Figure 1.2-12, a schematic diagram. The system consists of

Fig. 1.2-11. Parylene system.

a sublimation furnace, a pyrolysis furnace, and a glass cross deposition chamber cooled by liquid nitrogen.

This chamber holds seven stalk mounted targets and allows for maximum utilization of the flowing monomer. Pinhole ports along the walls of the chamber can be used to add inert gases to induce turbulent mixing for better coating uniformity. Reactive gases for copolymerization can also be introduced through these ports. The target holder can be pulled back from the chamber during deposition for visual observation and telescopic determination of the overall diameter of the coated target. The system is currently undergoing parametric system analysis.

Fig. 1.2-12. Schema of Parylene system.

1.2.3 Polymeric Shell Fabrication

The dried droplet method for polymeric shell fabrication was further developed this year (see KMSF 1978 Annual Report, 1-27-1-29). In this method droplets of polymer solution are heated while nearly in free fall to drive off the solvent and create discrete polymeric bubbles. The principal difficulty with this procedure is making droplets of constant size which are adequately separated. A device using acoustic modulations of a fluid jet (AMFJ) to make closely spaced, uniform droplets up to a millimeter in diameter was built and tested in 1978. Initial work on a circuit to charge droplets selectively and remove them from the mainstream was also done in 1978. Figure 1.2-13 shows the droplet deflection using the experimental charging circuit. A schematic drawing of the circuit which is presently being fabricated is shown in Figure 1.2-14. A study to produce large poly (methylmethacrylate) and polystyrene shells was done using polymer droplets dispensed from a syringe as a simulation of the droplet generator. But the size

Fig. 1.2-13. Acoustically modulated, fluid jet droplet generator, showing electrostatic droplet deflection.

distribution of the droplets was too large to obtain adequate droplet-shell correlations. Consequently, these studies were terminated until the droplet generator is completed.

Fig. 1.2-14. Schematic diagram of acoustically modulated, fluid jet droplet generator.

1.3

Cryogenics

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During the past year, we directed some of our efforts toward preparing for the cryogenic target designs proposed for future implosion experiments.

The target designs in the open literature were evaluated to identify common features and to uncover any requirements which are incompatible or unfeasible. A primary intent of this review was to define the practical options available, so that we might better focus our efforts in our developmental cryogenic program. At the same time, an area of continuing responsibility is support of the ongoing experimental program at KMSF. The basic technology and associated techniques for current experiments require further study in order to better understand the detailed response of a cryogenic target. In addition, a great deal of effort is still

directed toward providing, on a routine basis, cryogenic targets for current implosion experiments.

1.3.1 Multi-Shell Cryogenic Targets

A relatively large number of designs have been proposed for high gain IF targets.^{1,2,3} One concept present in a number of them is "vacuum" layering, i.e., a spherical fuel container separated from one or more larger spherical shells by a low pressure gas, or "vacuum," layer. In some instances vacuum implies nearly zero mass density, while in others it can refer to relatively high mass density but at reduced pressures.

Vacuum layer target designs present several fabrication problems, but the targets probably can be produced in quantity by variations on techniques already introduced.⁴ However, for cryogenic designs in which the fuel is a uniform solid layer on

pressure exchange gas in the vacuum layer, however, thermal conduction and convection are now available for the formation and maintenance of the fuel layer. We have used He gas to form and maintain uniform solid fuel layers. The small fuel container is immersed in 10^2 mTorr to 10^3 mTorr of He gas under our shroud system^(6,7,8). This experimental arrangement is thermally equivalent to a double shell target where the outer shell, the shroud), is many times the diameter of the fuel container. This geometry is of no interest as a commercial IF target, but it does demonstrate that uniform fuel layers within such targets can be achieved.

Vacuum Layer Model. A model⁹ has been developed to evaluate the thermal characteristics of multi-shell target designs to determine if the required uniform solid fuel layer can be produced at high repetition rates.

The tasks of first producing and then maintaining the uniform solid fuel layer require vastly different rates of heat extraction from the fuel container. Once the solid fuel is formed, it is stable as long as the heat from β -decay and any external source is removed in a manner that maintains temperature uniformity. However, maintaining the fuel layer typically requires cooling rates that are orders of magnitude lower than that desired for the formation process.¹⁰

To obtain sufficient heat transfer to form the solid fuel layer uniformly, the vacuum may be He exchange gas.¹¹ Non-condensable down to ~ 4 K even at 1 atm, helium can provide the thermal transport necessary to form an acceptable fuel layer. The minimum required He pressure within the vacuum layer is determined by thermal transport characteristics of the He exchange gas and is the minimum value necessary to produce the desired uniformity of the fuel layer during freezing. At this pressure, the mass density within the vacuum layer may be too high to be acceptable to the target designers.

The other gases that condense in the temperature range 14K to 25 K (H_2 , D_2 , T_2 , and Ne) offer additional options as exchange gas. In particular, they make possible the circumvention of the above prob-

Fig. 1.3-1. Double shell vacuum layer target

the inner surface of the fuel container (see Figure 1.3-1), the vacuum layers impose an additional problem. In the case of a zero mass density vacuum layer, heat transfer from the fuel container to the outer shell is limited to radiation. The heat transfer rate by this mechanism is insufficient to maintain a DT filled target below the freezing point of the fuel due to β -decay heating.⁵ With the presence of a low

lem through the construction of a mass density switch for the vacuum layer.

As an illustration of this mass density switching, consider a target of the geometry shown in Figure 1.3-1, where the vacuum layer is filled at room temperature with 1.0 atm of H₂ gas. As the target cools, the pressure of the H₂ exchange gas follows closely that defined by the Ideal Gas Law. At the temperature where the isochore of the Ideal Gas Law intersects the vapor pressure curve, condensation occurs. As the temperature is lowered further, the pressure within the vacuum layer will attempt to follow the equilibrium vapor pressure curve, eventually reaching equilibrium vapor pressure when the target is at its final temperature. If the mass density within the vacuum layer as specified by the target design is less than that which can be achieved with a particular apparatus utilizing H₂, then D₂, T₂, or Ne may be used. These gases will provide lower mass densities while still offering the possibility of forming a uniform solid fuel layer within the fuel container.

Besides acting as a mass density switch, a condensable exchange gas also acts as a thermal switch. As the vapor pressure of the gas decreases with temperature, the thermal conductance of the gas also decreases. In practice, this can preclude rapid condensation of the fuel, or prevent the extraction of β -heating from certain targets. When this occurs, additional options are available. Consider the case where the minimum mass density within the vacuum layer is desired. The minimum amount of He exchange gas necessary to extract β -heating from the specific target can be calculated. To increase the thermal transport during formation of the fuel layer, a condensable gas such as Ne can be added to the He. This mixture now provides thermal transport for rapid formation of a uniform fuel layer which is subsequently maintained with the minimum mass density in the vacuum layer.

Although the preceding discussion is useful in illustrating the concept of an exchange gas as a mass density switch, it provides no direct information that can be used to evaluate the probability of forming a uniform fuel layer within a specific target geometry. The time necessary to condense and freeze the fuel layer is the significant parameter.¹² This information is obtainable from the model by predicting the temperature histories of various

target components as they are cooled.

Figure 1.3-2 shows an example of a characteristic target cooling for He exchange gas. The times to condense and freeze the fuel layer are also listed with the figure.

Fig. 1.3-2. Double shell target temperature history (He exchange gas)

Figure 1.3-3 shows results for the same target as illustrated in Figure 1.3-2 when H₂ is the exchange gas. The times to condense and freeze the fuel layer are also listed with the figure. As can be seen by comparing Figure 1.3-2 and 1.3-3, the times to condense and freeze the fuel within the fuel container are close. This implies that the uniformity of a fuel layer formed within both targets should be about the same. Since the times listed are close to those observed experimentally for single shell cryogenic targets, it can be concluded that these targets should be highly uniform. The difference between the targets of Figures 1.3-2 and 1.3-3 at 8K is that the exchange gas density for He is $\sim 1.6 \times 10^{-4}$ gm/cm³, while for H₂ it is 5.2×10^{-7} gm/cm³.

1.3.2 Experimental Program

Production of uniform, solid fuel layer targets requires several pieces of specialized equipment.¹² Figure 1.3-4 shows one in schematic form, the

Fig. 1.3-3. Double shell target temperature history (H_2 exchange gas).

cryogenic shroud retraction device. This device retracts the shroud 10 cm in 60 msec, exposing the solid fuel layer target for only 4 msec before laser irradiation. During the operation of this device it is imperative that the target not vibrate or move from its prealigned position. Since any impulse applied to the target chamber could displace or vibrate the target, the retraction device is mounted on the ceiling, not directly on the target chamber. The only mechanical coupling of the entire retraction device to the target chamber is through the two inch diameter bellows.

The physical relocation of the KMSF laboratory facilities during 1979 necessitated some engineering changes for this device. Ceiling clearance above the target chamber in the new facility required the construction of a support structure to locate this retraction device properly above the target chamber.

A temporary support structure suspended from the ceiling of the new facility was constructed to determine experimentally the stability of the roof joints under the dynamic loading experienced during shroud retraction. Results of this series of experiments show that the total lateral motion of the shroud during retraction is less than 0.5 mm. This is well within tolerable levels. Also, we determined that retraction of the shroud does not couple energy

into a vibration of the target. Target vibrations are less than 2μ m. Installation of this temporary structure over the experimental target chamber will allow cryogenic experiments to be resumed in the new facility.

Fig. 1.3-4. Shroud retraction device.

Several design changes were incorporated concurrently into the Z-motion mechanism for the shroud retraction device. These changes will improve the utility of this control during the actual target shot sequence of adjustments.

A new illumination scheme consisting of a lens system for focusing the heating laser¹² onto a cryogenic target within the triple bounce illuminating system¹³ was evaluated. This system is capable of delivering approximately 75 percent of the available laser energy into a 40μ m diameter spot at the target plane. Although not as versatile as the system currently used in off line experiments, this new system will be capable of delivering sufficient energy to process a cryogenic target within the new target chamber configuration.

Another requirement for the production of uniform cryogenic fuel layers is the attainment of a maximum target cooling rate.¹² In order to accom-

plish this, several adjustments were made to the cryogenic shroud system (see Figure 1.3-5), which had previously been used to produce single shell cryogenic targets. The orifice through which the

Fig. 1.3-5. Gaseous helium shroud.

cold He gas enters the volume surrounding the target was increased to $15\mu\text{m}$ diameter. This is approximately the optimally sized orifice for our system. A balance between the cold He gas throughput of this orifice and the conductance of our vacuum system results in a background He pressure of $\sim 3\text{mTorr}$. A larger orifice with higher pressures would result in too much heat transfer from the chamber to the shroud system, so that the minimum temperatures would be unattainable. This throughput rate of cold He gas is one of two parameters which can be used to maximize the cooling rate of the target.

The second parameter which controls the cooling rate is the conductance between the shroud volume and the vacuum chamber. The controlling factor here is the clearance or annular gap between the bottom of the shroud and the lower cryostat extension. By decreasing this conductance, the pressure increases within the shroud volume containing the target. This directly increases the cooling rate of the target.

Discussions with J. Miller of LANSL¹⁴ suggested an adaptation of a baffle sealing system¹⁵ could be used to increase significantly the impedance of the He gas flow through this annular gap without causing target vibrations. Figure 1.3-6 illustrates the geometry currently in use. Very light contact is

Fig. 1.3-6 Baffle impedance system.

maintained between the shroud end and a series of Mylar sheet baffles in order to avoid displacing the prealigned target. The shroud-baffle contact does not result in a leak tight seal. The resultant increase of He pressure around the target position has significantly decreased the time required to condense and freeze the fuel within DT targets. Condensation and freezing times of less than 15 msec have been observed experimentally.

The cryogenic groups at KMSF and LANSL agreed to perform a set of experiments jointly. The objective of the experiments was to understand the differences observed in the quality and reproducibility of the solid fuel layers formed at the two facilities. KMSF supplied LANSL with a set of specific targets for initial evaluation. At each facility the experiments are either videotaped or filmed so the direct observations can be exchanged.

The first experiment conducted at LANSL was a direct comparison of a LANSL and a KMSF target mounted side by side. In general the LANSL target exhibited a highly repeatable interference pattern over a wide variation of He exchange gas pressures (from ~ 150 mTorr to ~ 850 mTorr). The interference pattern of the KMSF target did not show the same degree of reproducibility. Even with up to 400 mTorr of He exchange gas the pattern was very erratic. As the He pressure was increased to ~ 800 mTorr the reproducibility improved but never reached the level exhibited by the LANSL target. Since both targets were not of the same size nor of the quality normally desired for fusion experiments, there is some ambiguity about these results. The attempt to ship these targets to KMSF for corroborating experiments was unsuccessful. Both targets were lost in handling. However, future direct comparison experiments are scheduled.

A second set of joint experiments involved one particular KMSF target. Contrary to others observed at LANSL, this target gave very uniform and repeatable interference patterns over a wide range of He exchange gas pressures. The target was then returned to KMSF along with a videotape of the experimental procedure. Equivalent experiments with the KMSF shroud system confirmed that this particular target does indeed give uniform and repeatable interference patterns over a wide range of operating parameters. Films of these experiments were then sent to LANSL.

At present it is not known why this particular target stood out in its cryogenic response. From the total shift in the interference pattern during condensation of the fuel and the relatively low number of counts observed with an x-ray counting system,¹⁶ it is speculated that the fuel fill of this target is below the expected level of 100 atm. A reduced fill pressure gives a thinner condensed solid layer within the target, which may make the target more

Fig. 1.3-7. Film mounted double shell target.

Fig. 1.3-8. Fiber mounted double shell target.

Fig. 1.3-9. Vapor phase of fuel in double shell target.

amenable to a uniform response. Further experiments will follow using this and other targets.

The results of this joint experimental program to date are still tentative. Not enough data have been accumulated to draw any firm conclusions about the source of the nonuniformities often observed in KMSF cryogenic targets. However, the initial results do indicate that the general differences which have been observed in the past at LANSL and KMSF may not be due to the differences in experimental apparatuses.

New manufacturing techniques¹⁷ have led to the production of very high quality transparent hemishells for the construction of multilayer targets (see Figures 1.3-7 and 1.3-8). Routine construction of these targets will now allow us to test experimentally some of the theoretical predictions calculated for vacuum layer cryogenic targets. Figures 1.3-9 and 1.3-10 illustrate some results from our first experimental efforts with these targets. Figure 1.3-9 is the interferometric image of a target like that shown in Figure 1.3-8 when the fuel is held in the vapor phase by the heating laser. Figure 1.3-10 is the interferometric image of the same target with the fuel condensed out as a uniform solid layer.

Fig. 1.3-10. Solid phase of fuel in double shell target.

Summary. Several specific modifications were made to existing experimental equipment in order to increase its versatility and operational usefulness. Continued use of this equipment in a jointly conducted effort with LANSL has helped in defining and understanding those parameters which affect the cryogenic response of our glass targets. Increasing operational control of the uniformity and reproducibility of solid fuel layers will allow cryogenic targets to become a routine option for laser interactions experiments. A general theoretical model of complex double shell targets was initiated and developed to predict the cryogenic response of proposed target designs in the future. Several complex targets composed of a fuel container inside two SiO_2 hemispheres were tested experimentally. Although the experimental parameters were not directly comparable to the analytical model, continued routine production of these complex targets will be used to test and verify the validity and accuracy of the model.

1.4

Advances in Target— Fabrication Technology

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The development of technologies for target fabrication at KMSF has a double thrust. The first is to expand and develop the techniques for the manipulation, characterization, and fabrication of targets, which will assure a continuing supply and variety of targets and diagnostic aids in support of near term fusion research experiments. The majority of these targets involve a glass Microshell® fuel container mounted on a fiber. The second and longer term research effort is to establish a base of advanced manufacturing technologies which can meet the stringent requirements of targets for future fusion reactors while being compatible with the high volume production required by future "target factories." A major goal in this effort has been to demonstrate the capabilities of semiconductor processing technologies to produce high quality IF targets and components.

Proposed high gain IF targets consist of a fuel container surrounded by multiple coatings and spherical shells. Technologies to produce these more complex targets are just beginning to be developed. Several techniques being evaluated are unique and exclusive to KMSF.¹ Two basic approaches are being examined. The first is to form the multiple shell target by assembling hemispherical components. The second is to build the multiple shells *in situ*, that is, to form a shell within a shell directly.

In the first approach, semiconductor processing technologies commonly used in the manufacture of integrated circuits are being applied to the production of IF targets and components. Methods based on integrated circuit technology (ICT) show potential not only for providing fabrication schemes with structural design flexibility for the near term, scien-

tific feasibility target study, but also for realizing the high volumes, high quality, and low cost required for fusion reactor targets.

The hemispherical assembly approach with the double shell target design is illustrated in Figure 1.4-1. The fuel pellet is supported by a membrane or filament in a position concentric within the enclosing hemispheres.

Fig. 1.4-1. Exploded view of double shell target.

Work on the second approach to producing complex, multiple shell targets—the shell within a shell concept—was recently begun. Several unique fabrication schemes are being evaluated. At this early stage of development these approaches are being tried without regard to the concentricity of the inner and outer shells. Our initial interest is to identify fabrication and characterization problems.

Highlights of Applied Fabrication Technology. The fabrication of targets for fusion experiments requires not only maintenance of skills for "conventional," Microshell® fuel containers on a stalk, but also the development of techniques for producing diagnostic and uniquely configured targets. The innovations required to produce, handle, and assemble these special targets adds to the KMSF fabrication repertoire for future needs.

Nearly 800 targets were selected, characterized, and mounted for the KMSF laser interaction experiments. The largest portion of these were DT-filled glass shells mounted on glass fibers. Some other interesting target configurations assembled

for specific experiments within KMSF and other IF laboratories are described below.

KMSF continues its effort to produce and characterize cryogenic targets. Unique double shell targets were recently assembled for the cryogenic experiments (for discussion of cryogenic considerations, see Section 1.3). One is shown in Figure

Fig. 1.4-2. Double shell target assembled from SiO₂ hemispheres and glass Microshell® fuel container.

1.4-2. The targets consist of two hemispheres of SiO₂ mated and sealed around a Microshell® fuel container. The Microshell® fuel container is held in the center of the sphere by thin membranes stretched equatorially across the hemispheres (schematically illustrated in Figure 1.4-1). Details of the hemisphere fabrication and the formation and characterization of the cryogenic fuel layer within the Microshell® fuel container are described in other Sections of this report.

Diagnostic targets and targets for special experiments require innovative fabrication, since most are unique configurations used only for one specific experiment. The target shown in Figure 1.4-3 is in this latter category. This target consists of a glass hemisphere (100μm in diameter) placed in

Fig. 1.4-3. Glass hemisphere positioned into a hole in a metal foil.

holes drilled thru a thin metal foil. The hemisphere was formed by carefully cutting a glass Microshell® target in half. The assembled target was used in one-beam transmission/absorption experiments. One of a kind targets were also produced for other DOE laboratories (e.g., PMMA slabs mounted on polyethylene and boron wafers, polished and over-coated with GDP p-xylene polymer). For KMSF, pinhole arrays and dual slits etched in gold were formed for the x-ray camera and x-ray spectrograph. The coded zone plate fabrication procedure was also improved by incorporating ion milling in the etching step. This modification in the fabrication procedure will enable the number of zones to be increased.

The capability to produce and handle large, spherical glass capillary bulbs filled with DT was achieved. Pyrex (Corning 7740) and lead glass (Corning 0010) bulbs were produced with diameters of 5 mm to 7 mm and wall thicknesses of 10 μm to 30 μm. Internal gas pressures of 1 atm to 5 atm were achieved either by filling through the capillary bulb stem or by permeation. A screening technique based on overpressurization was developed to select the bulbs which would survive processing and handling. We also developed a nondestructive technique, described in Section 1.5, to measure the tritium content of these targets.

Fig. 1.4-4. Fabrication sequence for double shell targets using integrated circuit processing technology.

1.4.1 Application of Integrated Circuit Technologies to Target Fabrication

For the multiple shell structures proposed for IF targets, solid state processing technology is proving to be a powerful tool for the formation of the shells. While complete, sealed, spherical prototypes of this structure have been fabricated, most work during this period was directed toward fabrication of the hemispherical shells. In addition, we expended considerable effort on the comprehensive characterization of hemishells and cavities in terms of circularity, wall thickness, sphericity, and surface finish. The results of these investigations will be discussed and illustrated in the remainder of this Section.

One possible fabrication sequence for the manu-

facture of the double shell target using integrated circuit technologies is illustrated in Figure 1.4-4. Conventional semiconductor processing techniques such as masking, photolithography, dopant diffusion and selective etching are used for many of the process steps. One advantage of this fabrication scheme is the large number of identical targets or hemispheres that can be produced at one time, as illustrated in Figure 1.4-5. About 20,000 hemispheres each 400 μm in diameter are produced from wafers 67mm in diameter. The wafers in Figure 1.4-5 are about 33mm in diameter.

The critical aspects of the fabrication sequence studied this year involve the cavity formation by isotropic etching, the shell wall formation by boron doping, and the formation of free standing hemispheres by anisotropic etching.

The formation of hemispherical cavities in silicon

Fig. 1.4-5. Arrays of circular mask openings and hemispherical cavities on silicon wafers.

using batch photolithography and isotropic silicon etching is a key step in the formation of shells by use of solid state process technology. Other cavity

forming techniques will be investigated in the coming year.

The surface quality and geometric hemisphericity of the cavity depend on the etch composition, the etch agitation rate, and the manner in which the etch is exposed to and removed from the etching surface. In this cavity formation process, illustrated in Figure 1.4-6, a circular opening is engraved in a metal/oxide mask to expose the silicon surface. The wafer is then subjected to an isotropic silicon etch. The etch used is 10 percent hydrofluoric acid and 90 percent nitric acid. The dissolution of silicon in this HF-HNO₃ mixture is known to be rate limited by the transport of HF to the silicon surface, where the HF then removes the SiO₂ produced by nitric acid oxidation.^{2,3} Since the etch rate is limited by mass transport, it is proportional to both the bulk HF concentration and to the velocity of the etchant normal to the etching surface. Thus, both factors

Fig. 1.4-6. Fabrication sequence to produce hemispherical cavities.

must be accurately controlled to obtain uniform, reproducible cavities. With the appropriate conditions we have obtained surfaces completely free of features at $10,000\times$ magnification.

Circularity of Isotropically Etched Cavities in Silicon. The circularity of the cavities was measured with a Quantimet image analyzer. The Quantimet consists of a video camera which interfaces with an inspection microscope, a PDP-11 minicomputer, and a teletype output device. The image is also displayed on a CRT. The screen is divided into 525 lines of 800 pixels each. The cavities were examined individually at a magnification sufficient to fill most of the display screen (a $375\text{ }\mu\text{m}$ diameter cavity corresponds to about 400 pixels). The resolution of the Quantimet was verified to be reproducible to ± 1 pixel over the field of view.

More than 800 cavities were examined on eight wafers. For wafers processed identically, the reproducibility in average cavity diameter was better than ± 1 percent from wafer to wafer. On a given wafer, 90 percent of the cavity diameters were within ± 2 percent of the average. A slight anisotropy was noted in the etch, with the $\langle 100 \rangle$ directions (45° and 135°) being about 1 percent larger than the $\langle 110 \rangle$ directions. These results are summarized in Figure 1.4-7, where the directional distributions and cavity departures from circularity are shown. It should be noted that the circularity of these cavities approaches the resolution limits of the measuring technique. These studies demonstrate the attainability of high yields of cavities whose diameters are uniform (with direction) and reproducible to 1 percent or better.

Sphericity of Isotropically Etched Silicon Cavities. The shape or sphericity of the etched cavities was determined by cross sectioning the cavity and by a radiographic technique developed for these studies.

The radiographic approach shows promise for characterizing large numbers of cavities, if computer assisted evaluation were available. The technique is illustrated in Figure 1.4-8. The method uses high resolution film to capture the image of the cavity produced by the enhanced penetration of the x-rays in the thinner (i.e., the cavity) regions of the wafer. Densitometer traces of the image can then be compared to those expected from an ideal hemispherical cavity.

The cavity shown in Figure 1.4-9 is thus too shallow. The point by point variations in the data are due to the noise in the densitometer trace and do not reflect surface roughness.

The sphericity of an etched cavity is a function of the mask opening size relative to the final cavity diameter and of the velocity etch rate of the etchant. Intuitively one would suspect that cavities could never be hemispherical owing to the finite mask opening size. However, surface diffusion layers associated with the acid transport and agitation produced velocity flow patterns more than compensate for finite mask openings. These conditions result in cavity cross sections which can vary from bathtub to bullet shapes. Resulting cross sections are illustrated in Figure 1.4-10.

The effects of mask opening size and agitation rate are summarized in Figure 1.4-11 for a fixed etching time of 40 minutes. As expected, cavity size increases with the mask opening and agitation, since both increase the supply of HF to the etching surface. Etch rate is a linear function of agitation and is independent of etch time over the ranges employed here. A sphericity factor is defined as the cavity radius to depth ratio. The data show that at low to moderate agitation rates, an increase in mask diameter leads to a decrease in sphericity, producing more bathtub shaped cavities, as would be expected in an isotropic etch. However, at high etchant velocities, the sphericity not only increases with velocity for a given mask size, but also increases with increasing mask size. This may be due to the decreasing influence of lateral flow and fluid shear near the cavity center. This dependence also suggests the existence of a velocity at which sphericity is independent of mask opening size.

Formation of Hemispherical Shell Walls and Free Standing Hemispheres. Once a high quality, hemispherical cavity is formed, it can be used as a mold, or the inner cavity surface can be modified to form the shell wall. When used as a mold, a variety of materials can be deposited using a host of techniques (e.g., chemical vapor deposition, sputtering, plating, ion beam deposition, etc.) to form the shell wall. Most of our work, however, has concentrated on the second approach, in which single crystal silicon shells are formed by modifying the cavity wall so that it is retained after dissolution of the

Fig. 1.4-7. Distribution of cavity diameters on one wafer measured on the Quantimet Image Analysis System.

Fig. 1.4-8. Sphericity measurement by x-ray radiography.

supporting wafer.

Several etchants are known to show anisotropic behavior depending on one or all of the following conditions: the chemical state, the boron concentration, and/or the crystallographic orientation of the silicon. Of particular utility has been an etchant composed of ethylene diamine, pyrocatechol, and water which is sensitive to all three conditions. The etch rates drop to nearly zero for boron levels which exceed $5 \times 10^{19} \text{ cm}^{-3}$ (about 0.1% boron in the silicon lattice).⁴ The etch rate ratios exceed 100 for silicon versus silicon dioxide or different crystallographic orientations, such as the (100) versus the

Fig. 1.4-9. Cross section of hemispherical cavity.

Fig. 1.4-10. Cavity profiles produced by isotropic etching.

Fig. 1.4-11. Cavity diameter and sphericity as a function of mask diameter.

(111) directions. We have made high quality, free standing shells from boron doped, single crystal silicon and from silicon dioxide using these anisotropic features.

Boron is introduced into the silicon lattice in amounts approaching 1 percent of the total atoms for silicon hemisphere formation. Although there are several techniques available for doping silicon (conventional surface diffusion, ion implantation, or epitaxy), our studies have been limited to conventional diffusion. A layer of boron is deposited on the wafer surface at high temperature and allowed to diffuse substitutionally into the lattice. The boron can be introduced into an ambient gas steam via solid, liquid, or gaseous sources, or can be applied to the wafer directly via liquid spin on solutions. The spin on sources were found unsuitable, and all subsequent work was done with solid boron oxide source wafers. These wafers are positioned adjacent to the silicon wafers, and boron is volatilized at high temperature. The wall thickness was precisely controlled by controlling the time and

temperature of the diffusion. Measured wall thicknesses were shown to agree well with the predictions of classical diffusion theory. On a given hemisphere, variations in wall thickness are expected to be extremely small, since the diffusion process is highly isotropic. The variation from run to run can be estimated, assuming a furnace temperature variation of $\pm 1^\circ\text{C}$ and a diffusion time variation of ± 10 seconds. These variations lead to an expected worst case, run to run wall thickness variation of ± 1.3 percent on $1.3 \mu\text{m}$ thick shells ($\pm 168 \text{ \AA}$). In actual practice, furnace temperatures are often held to $\pm 0.5^\circ\text{C}$, so that a corresponding reduction in wall reproducibility should be expected. The time variation assumed here accounts for only 11 percent of the total wall variation at $\pm 1^\circ\text{C}$.

Surface Quality. The quality of the surfaces produced on the hemispherical shells is of great importance in target implosion dynamics, and it depends

on the etching and diffusion procedures employed in cavity and shell formation. The interior surface quality of the free standing, hemispherical shells is dependent on the isotropic etching procedures employed in cavity formation, on the doping process, and on the ability of the boron doped, silicon surface to withstand the anisotropic wafer dissolution step without degrading. Similarly, the exterior shell surface quality is determined by the quality of the initial interior surface, the doping process, and the interaction of the anisotropic etchant.

Examples of the high quality surfaces that can be achieved by optimizing the conditions which influence the surface are shown in Figures 1.4-12 and 1.4-13. Figure 1.4-12 shows an SEM view of a free

Fig. 1.4-12. Free standing silicon hemisphere.

standing, hemispherical shell. Except for the flanges, both the inner and outer surfaces of the hemisphere are essentially featureless within the resolution of the SEM at 10,000× magnification. This result is typical of nearly all interior and exterior locations and is reproducible from shell to shell and run to run. This indicates that smooth, featureless surfaces can be etched, doped, and then preserved throughout the dissolution of the wafer without degradation.

Silicon dioxide shells are produced by thermal oxidation of the silicon cavities and by direct oxidation of free standing silicon hemispheres. Figure

Fig. 1.4-13. Array of SiO₂ hemispheres.

1.4-13 shows an array of hemispherical shells produced by the former process. The surface and clarity of the silicon dioxide shells were of a quality sufficient for use in cryogenic experiments on double-shell targets (see Section 1.3). Figure 1.4-2 shows a double shell target which was assembled using these hemispheres.

Variations of the semiconductor processing schemes which take advantage of the particular properties of silicon and the anisotropic etchants were also explored. Nonhemispherical structures such as

Fig. 1.4-14. Free standing pyramidal shell.

that shown in Figure 1.4-14 can be generated in a manner similar to hemispherical cavities.

During the past year the idea of using solid state process technology for IF target production was developed and prototype structures were characterized. A process for producing hemispherical shells having excellent reproducibility, circularity, and interior and exterior surface finish was defined. The process produced hemispheres 150 μm to 400 μm in diameter and having walls from 0.5 μm to 4 μm thick. The influence of more than a dozen independent process parameters was quantified and partially optimized. Nevertheless, the fabrication sequence used thus far relies on the more commonly existing processes. A variety of more advanced but still widely available processes and approaches remains untapped. In nearly all areas of the present shell fabrication process there exists a potential to increase the design flexibility for more complex target structures without losing the high volume potential of the basic approach. At least four areas for work during the coming year have been defined: (1) refinement of the current fabrication process, (2) evaluation of advanced shell fabrication techniques, (3) development of a process to form the support membrane, and (4) development of a batch sealing process.

1.4.2 Target Injection/Suspension

A first step toward a target injection system is development of the ability to illuminate with laser pulses targets without mechanical support. This was accomplished at KMSF laboratories last year using the retracting pedestal, inertial suspension system.⁵ This apparatus permits study of the effects of target asymmetry on implosion dynamics. The absence of a stalk which normally supports the target removes an expected source of plasma interference.

Rapid withdrawal of the support pedestal is accomplished by an hammer-anvil assembly. When activated, a solenoid plunger (hammer) accelerates to its maximum velocity just before striking the pedestal assembly (anvil), which then recoils downward from under the target. Because the pedestal assembly is approximately 1/10 the mass of the plunger, a velocity multiplication occurs at impact (see Figure 1.4-15). The pedestal moves with a constant velocity at a value which, for a nearly perfect elastic interaction, approaches twice the velocity of the plunger.

The major part of this year's work was directed toward understanding the interactions between the target and the retracting pedestal. The target is attracted to the pedestal by forces such as static elec-

Fig. 1.4-15. Solenoid and pedestal velocity of retracting pedestal system.

Fig. 1.4 -16. Typical target pedestal separation obtained by retracting pedestal system, recorded with strobe light assembly.

Fig. 1.4 -17. Target pedestal separation obtained by retracting pedestal system, recorded with high speed motion picture camera.

tricity, van der Waals forces, and chemical bonding. Consequently, when the platform is withdrawn, an initial velocity is imparted to the shell. For laser interaction experiments, the pedestal must be withdrawn beyond the diagnostic field-of-view prior to the shot. If too large an initial velocity is imparted to the shell, it can move beyond accep-

table limits of target position during the pedestal's withdrawal.

Two instrument packages were used to photograph the target platform interaction. The first is an MP-4 Polaroid camera with a strobe light assembly. Multiple images are produced with intervals between strobe flashes as short as $10 \mu\text{sec}$ by using several strobe lights tied into a common triggering unit. A typical photograph taken with this system set at $5 \mu\text{sec}$ flash intervals is shown in Figure 1.4-16. Note that only one image of the graphite pedestal is present. The pedestal is completely withdrawn from the picture by the time successive images are taken. Also, observe that the images of the falling microsphere are spaced farther and farther apart due to gravitational acceleration.

The second device we have for observing target-pedestal separation is a high speed motion picture camera, an 8 mm Fastax WFI model. This camera is capable of filming 8,000 frames/sec ($125 \mu\text{sec}$ intervals). A sequence of frames taken with this Fastax camera is shown in Figure 1.4-17.

These instruments allowed us to determine that an initial velocity is consistently being imparted to the Microshell® targets by the pedestal withdrawal. This velocity varied from 2 cm/sec to ~ 50 cm/sec, depending on the shell size, shell cleanliness, and the amount of surface moisture present.

The importance of moisture in causing shells to adhere to surfaces was illustrated dramatically this year. A recurring problem was that targets would not fall freely from the loading tube onto the support pedestal (see Figure 1.4-18). To test the hypothesis that surface moisture might be the problem, a heating element was attached to the tube exit port. Heating the tube to 100°C for a few minutes alleviated this difficulty. Targets could be loaded onto the pedestal routinely in five minutes.

This year, targets as small as $125 \mu\text{m}$ diameter were loaded on the pedestal successfully. A glass capillary tube was drawn to a diameter opening of $250 \mu\text{m}$ and glued in place over the end of the loading tube to serve as a funnel for loading small shells onto the pedestal (see Figure 1.4-19). Continued progress in solving the problems of target-shell interactions will allow smaller targets to be used with this system. Incorporation of a heater in the pedestal assembly is expected to contribute to this progress.

Fig. 1.4-18. Shell loading apparatus.

Fig. 1.4-19. Small shell modification.

1.5

Physical and Chemical Analytical Support

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Introduction. Physical and chemical analytical support consists of non destructive and destructive fuel analysis, characterization (i.e., wall thickness, diameters, wall uniformity, surface features, etc.), elemental analysis of materials, statistical analyses, strength tests, general analytical chemistry support, and so on. Many of the chemical and physical analyses to be performed often require new techniques and technology.

An adequate and readily usable non destructive analysis of the fuel does not exist. In 1979 we demonstrated the use of Laser Raman Spectroscopy (LRS) to determine non destructively the partial pressures of D₂, T₂, DT, DH, and HT, with detectability limits less than 100 pg per species. We are continuing development of the technique.

Destructive fuel assay procedures are important for consistency tests and calibrations required for non destructive analysis development. This last year we have improved our gas chromatographic technique by column development. We have found that a nitrogen de-activated alumina column, instead of the carbon dioxide de-activated alumina, produces excellent gas species resolution.

New methods and techniques were designed for silicon, boron, aluminum, and hafnium quantitative analyses.

Quantitative surface analysis of submicron structures via Scanning Electron Microscopy (SEM) and Energy Dispersive X-ray Spectroscopy (EDXS) was improved. The Coates and Welter Quickscan was upgraded by incorporation of an electrostatic lens, which provides a resolution approaching 90 Å. The Kevex EDXS system was fitted with a PDP 11/10 minicomputer and software.

Two rapid inspection techniques for wall uniformity and aspect ratio, the Rollatron and the Ameba, were evaluated and further developed. The

Rollatron utilizes the rolling of a spherical shell to determine its uniformity. We have utilized this device routinely to qualify batches of shells fabricated in the vertical tower furnace. Batch-uniformity measured by this technique correlates well with interference measurements.

The AMEBA system, which determines wall thickness and diameter, is based on a wall effect phenomenon. The device is still in the developmental stage, but tolerances of ± 5 percent on the diameter will be attainable with measurements at about one per second. Wall thickness can be determined by the vertical velocity of the shell through a fluid once its diameter is known. Preliminary tests of the experimental apparatus indicate a sensitivity of $\sim 0.05 \text{ cm/s}/\mu\text{m}$ for the shell diameters currently being used.

1.5.1 Non Destructive Fuel Analysis

A non destructive method for fuel analysis of laser fusion targets via Laser Raman Spectroscopy has been demonstrated by KMSF and the National Bureau of Standards. The results have been published¹ under joint KMSF-NBS authorship.

Typical Microshell[®] targets (100 μm to 200 μm in diameter, with 1 μm to 2 μm wall thickness and 10 atm to 70 atm of fuel) were irradiated with an argon ion laser. Rotational Raman scattering lines were used to analyze for mole fractions and absolute partial pressures for all hydrogen isotope species (H_2 , HD, HT, D_2 , DT, T_2). Accuracies of 0.01 (mole fraction) were obtained from spectra requiring 20 minutes of scan time. Individual species detectability limits were less than 100 pg. Absolute partial pressure measurements were accurate to ± 20 percent.

While analysis times of 20 minutes per shell were observed in this work, application of current technology can reduce this to a few seconds or less per target. Thus, this method may be useful for fuel analysis of targets in a fusion power plant.

TAXE. This year a Tritium Assay by X-ray Emission (TAXE) system was implemented as a nondestructive method for determining the fuel contained

within fusion targets.² The energy of the beta particles emitted from the decay of tritium in the fuel gases is partially absorbed by the glass shell. The

Fig. 1.5-1. X-ray spectrum of tritium filled Microshell[®] target from 1.7 keV to 18 keV.

energy is reemitted as bremsstrahlung radiation over the range 0.7 to 18 keV (Figure 1.5-1). Since the β -particle to X-ray conversion efficiency is constant for a given composition and thickness of the glass shell, a measurement of the emitted x-rays is a measure of the tritium fuel fill.

The design of this system and its required instrumentation is that developed by LANSL.³ The target sits within 4.8 mm from two 1 mm thick,

Fig. 1.5-2. Pressure x diameter vs x-ray cps per ng tritium for glass Microshell[®] targets.

Fig. 1.5-3. TAXE system modified to assay large capillary bulbs, NaI detector positioned above plastic cylinder which encloses capillary bulb (not pictured).

cleaved NaI crystals, which are fitted with .05 mm Be windows. These are optically coupled to photomultiplier tubes which feed a multichannel analyzer.

Calibration of the system (Figure 1.5-2) is accomplished by obtaining the x-ray spectrum from a shell with the multichannel analyzer and then measuring the shell's tritium content by rupturing it in an ionization chamber. With the calibration, a measured x-ray count rate, and the fuel composition (D:T:H ratio), the total pressure can be calculated. This total pressure agreed well with the total-pressure measurement by the GEVF pressure test⁴ for several batches of shells. Future automation of this system will allow assay services with little direct labor input.

Shells from 55 μm to 300 μm in diameter with walls 0.6 μm to 2.5 μm thick and containing 2 atm to 130 atm of fuel have been routinely assayed with this system. An adaptation (see Figure 1.5-3) has been used to measure the tritium (0.1 Ci to 0.5 Ci) contained in 6 mm diameter capillary glass bulbs.

1.5.2 Cryogenic Gas Chromatography

Two methods for improving the cryogenic gas chromatograph⁵ were examined: reduction of analysis time using the packed alumina column, and use of a glass capillary column with an activated inner surface as the chromatographic substrate.

We experienced only limited success with the latter approach. The glass capillary columns treated were typically 50 m to 70 m long, 0.6 mm to 1.0 mm O.D. and 0.25 mm to 0.40 mm I.D. Both pyrex and soda-lime glass were investigated. The inside diameter of the column was activated by etching with either ammonium hydroxide or sodium hydroxide, or by causing crystalline growth with a chlorofluoro ether.⁶ We find pyrex glass capillary columns cannot be activated to separate hydrogen isotopic species at liquid nitrogen temperatures. Some soft glass columns were made active, others were not. The chromatographic activity appears to be dependent on glass composition. Since the glass capillary columns which separated the isotopic species also produced an ortho-para separation for both the hydrogen and deuterium species (undesirable for quantitative purposes), we stopped work on glass columns and returned to packed alumina columns.

In order to provide an acceptable separation of all hydrogen isotopic species in a fuel mixture, the alumina in the gas chromatography column must first be strongly heated to activate the surface of the packing, and then selectively deactivated to provide the degree of resolution useful for analytical purposes. This selective deactivation was a time-consuming step involving careful pressurization and equilibration of CO_2 over the alumina packing. Gaseous N_2 can also be used for deactivation. Using N_2 in the deactivation step, we can bring the column into working condition in approximately two hours. Figure 1.5-4 illustrates a separation of fuel-gas molecular species obtained using N_2 as the deactivating gas.

1.5.3 Elemental Analysis

The analytical group developed methods of analysis for aluminum and hafnium and extended the elemental analysis of low concentrations of silicon two orders of magnitude (from 100 ppm to 1 ppm).

Fig. 1.5-4. Gas chromatographic separation of molecular species in KMSF fuel mixture.

Aluminum in glass can now be determined by atomic absorption spectroscopy both in the flame mode (for normal concentrations) and the flameless, carbon rod mode (for trace levels). Hafnium is determined gravimetrically in polymers by decomposing the polymer and precipitating hafnium oxide with ammonium hydroxide. As a major constituent, silicon is normally measured by flame atomic absorption spectroscopy. Since the lower concentration level for practical use here is 70 ppm, we have implemented a method for silicon analysis based on that published by the International Commission of Glass. The method is colorimetric, depending on development of a silicomolybdenum blue color, which provides sensitivity to silicon in the 0.5 ppm to 3.0 ppm range.

1.5.4 Shell Characterization/Analysis

The commercialization of inertial fusion will require high volume, low cost manufacture of power reactor targets. An important aspect of the manufacturing process is the quality assurance of target components such as the fuel containers. Where current requirements allow, we are attempting to meet this need by developing techniques for rapid and

cost effective characterization and sorting of Microshell® targets.

Wall Uniformity of Shells. Current theory predicts that shell wall nonuniformity (WNU) better than 1 percent is needed to assure ignition. WNU is presently measured by microradiographic⁷ and interferometric techniques.⁸ These techniques are sufficient for the low volume production of experimental targets. However, both are inadequate for the high speed, inexpensive target inspection rates that will be required by future IF power reactors. To meet this requirement, we developed a rolling technique, the Rollatron System, which characterizes and sorts shells for wall uniformity in a simple, relatively precise, and rapid manner.

The principle underlying the operation of the Rollatron can be stated succinctly: when a sphere rolls down an inclined plane, it will veer away from the centerline path if its center of mass does not coincide with its geometric center. The trajectory depends on the location of the sphere's center of mass with respect to its geometric center, and thus, on wall non uniformity, the determinant factor in mass asymmetry.

A set of glass shells selected for various degrees

Fig. 1.5-5. Schema of shell rolling apparatus (Rollatron).

Fig. 1.5-6. Comparison of shell nonuniformity measurements from shell rolling technique to those from interferometry.

of WNU as measured by interferometry was studied. The measured roll trajectories were similar to those predicted by the first order model⁹ of a sphere having an offset center of mass. The apparatus and results of this experiment are displayed in Figures 1.5-5 and 1.5-6 respectively. Measurements made by interferometry and the rolling technique provide strong support for the validity of the simple mathematical model.

We improved the performance of the system by adding a heating element to the shell feeder trough. Heating the shells just prior to rolling drives off

moisture and decreases the tendency for shells to agglomerate or stick to the surface. As in the case of the retracting pedestal, the results indicate the strong role moisture plays in agglomeration.

The WNU of a batch of shells can be determined readily from the distribution pattern produced by rolling representative samples (see Figure 1.5-7). Experimental batches of glass shells are now being characterized by the Rollatron. In the future this technique should be useful for quality assurance in large scale processing.

Fig. 1.5-7. Distribution of rolled shells (KMSF batch T2R-11) resulting from differences in wall nonuniformity.

Diameter and Wall Thickness. In the AMEBA (Automatic Microshell® Evaluation by Buoyancy Analysis) device, the velocity of a spherical shell passing through a fluid is attenuated by the close proximity of a surface. In this device, the buoyancy of a shell forces it upward through a cylindrical tube containing a fluid (e.g., distilled water). The shell's velocity is attenuated by the tube's walls. If the shell's velocity is measured in two tubes with different inside diameters, a set of two equations in two unknowns (shell diameter and wall thickness) can be used to describe the experiment. The vertical velocity (V_{free}) of the shell is equal to the terminal velocity in an infinite medium when the I.D. of the tube is much greater than the shell's diameter. Frictional drag resulting from wall effect significantly reduces the shell's vertical velocity ($V_{restricted}$) when the second tube has an I.D. similar to the diameter of the shell.

In either tube the upward force on the Microshell® is the force of buoyancy, F_B . After the shell reaches constant constant velocity, F_B is equal to the forces of gravity, F_G , and drag, F_D :

$$F_B = F_G + F_D \quad (1)$$

EXPERIMENTAL SET-UP TESTING
AMEBA* CONCEPT

*AUTOMATIC MICROSHELL® EVALUATION
BY BUOYANCY ANALYSIS

Fig. 1.5-8. Schema of AMEBA device (Automatic Microshell® Evaluation by Buoyancy Analysis).

Since F_B and F_G are independent of the tube diameter, the drag forces in each tube are equal.

The drag force, F_{D1} , acting on the shell in a tube whose diameter is large compared to the shell, is:

$$F_{D1} = 3\pi\eta \cdot V_{free} \cdot D \quad (2)$$

where D is the diameter of the spherical pellet, and η is the viscosity of the fluid. The drag force, F_{D2} , in the narrow tube can be written,

$$F_{D2} = 3\pi\eta \cdot K_w \cdot V_{restricted} \cdot D \quad (3)$$

where K_w is the wall effect factor (the degree of velocity attenuation caused by the walls). Equating (2) and (3) gives:

$$K_w = \frac{V_{free}}{V_{restricted}} \quad (4)$$

VELOCITY ATTENUATION OF KMSF MICROSHELL® PELLETS

Fig. 1.5-9. Velocity attenuation of glass shells in AMEBA.

A measurement of a shell's velocity in each tube is sufficient to determine the diameter of the shell, since, to first order, K_w is a monotonic function of the shell diameter only.

The prototype apparatus is displayed schematically in Figure 1.5-8. The results, shown in Figure 1.5-9, indicate shell diameter can indeed be measured using this technique, even utilizing manually timed events.

The value of K_w along with V_{free} can be used to calculate the effective wall thickness once the diameter of the shell is determined. Equations for standard particle dynamics relate the terminal velocity of a sphere to diameter and sphere density, which is a measure of mean wall thickness.

The correlation of wall thickness with shell velocity was demonstrated in an experiment where seven shells having identical diameters but different wall thicknesses were AMEBA analyzed. The vertical velocities for shell wall thicknesses as measured by interferometry are plotted in Figure 1.5-10.

Precise measurement of shell diameter and wall thickness necessitates precise velocity measurements. Skanomatic sensors have been incorporated into a second version of the system to provide accurate and automatic timing. Each sensor produces a light beam "gate" within the glass tubes. When a shell breaks the beam, a timer (frequency counter) is activated, which runs until the shell breaks the next light gate. The distance between

Fig. 1.5-11. The AMEBA target handling technique.

gates being known, a precise, contact free measurement of average velocity is obtained.

An inherent attribute of this technique is the gentle handling received by the shell. The drag forces associated with the wall effect tend to center the shells as they rise in the tube. In principle the shell need never touch a surface in order to be characterized.

New handling techniques were required to transfer shells in and out of the AMEBA tube assembly. A simple yet effective technique is illustrated in Figure 1.5-11. A small, fluid filled bottle contains the shell which has floated to the apex of the liquid meniscus. The shell transfer to the tube assembly is accomplished by joining the meniscus of the bottle to that of the tube. Fluid in the tube is prevented from flowing out by two teflon stopcocks in series. Opening one valve at a time permits the shell to rise to the apex of the meniscus on the top of the tube. Transfer of the shell back to the bottle is accomplished by inverting the bottle and again joining the menisci.

With continued development, the AMEBA should find immediate use as an analytical tool for statistical analyses of batches of shells. At a later date, automation may allow unattended sorting of shells for specific diameters and wall thicknesses. Ultimately, arrays of AMEBA units could function at quality control stations within IF target manufacturing facilities.

Fig. 1.5-10. Wall thickness vs velocity of shells having nearly equal diameters (O.D. of $514 \pm 3 \mu\text{m}$).

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Laser Fusion Experiments

Long Pulse Experiments

Short Pulse Backscatter from a Gas Target

Target Experiments at $0.53\ \mu\text{m}$

Implosion Physics

Diagnostic Development

Laser Fusion Experiments

Introduction
F.J. Mayer

The fusion experiments program has continued to emphasize spherically symmetric implosion experiments. Our motivation is to keep multidimensional effects to a minimum while keeping high fuel densities as a goal. We have begun to examine the more ablatively driven laser implosions by going to lower irradiances and shorter wavelengths to reduce hot electron production and, thereby, the "explosive" nature of the shell implosions. In our long pulse experiments, low irradiances at longer pulses (~ 700 psec) have shown some results which are quite different from our previous short pulse experiments. In particular, x-ray pinhole photographs are of a different character, and the ion flow has not shown the expected steady state nature. The analysis of these results and their implications regarding flux limitation and ablative transport is still underway.

In the past year we have continued our experimental efforts to investigate the effects of shorter wavelength irradiations. Our $0.53 \mu\text{m}$ "green" experiments are already showing some expected (and beneficial) coronal behavior. The hot electron spectrum appears to be cooler, and fast ion energy losses are smaller. These results definitely suggest that proceeding to $0.35 \mu\text{m}$ "blue" implosion experiments is an important next step.

Because reactor size pellets will be large, therefore necessitating long density gradients, we have begun experiments to examine Brillouin backscatter from a long density gradient gas target. These target experiments may prove to be a test bed for various absorption and transport experiments in the future.

In our continuing efforts to understand quantitatively the spherical shell implosion process, we are developing both simple models and more complete hydro code simulations of the implosion process. Section 2.4 gives the results from some of our efforts in these areas during the past year.

Our ability to diagnose the target plasmas has in-

creased considerably in the past year with the development of holographic interferometers, time resolved x-ray measurements, plasma measurements, and a new method for determining fuel qR .

Finally, we believe that the fusion experiments program has produced important new results for the laser fusion community. Our "green light" experiments and the comparison with $1.05 \mu\text{m}$ experiments definitely seem to be giving the "green light" to "blue light" experiments. We will proceed in this important new area of shorter wavelength experiments in the coming year.

2.1

Long Pulse Experiments

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A series of 62 target shots using ~ 1 ns laser pulses was performed prior to the upgrade and restaging of the laser. This series was designed to study laser performance for long pulses and to obtain absorption and implosion data in a new parameter regime which will be useful in the design of future experiments.

The long pulse was generated by stacking a series of up to 22 individual pulses. The individual pulses are produced from the single 40 psec FWHM pulse output of the passively mode locked oscillator.¹ When the individual pulses are in proper phase, the resulting pulse rises and falls sharply and is relatively flat topped. Figure 2.1-1 shows the streak camera record of a representative pulse. Occasionally interference effects produce pathological pulse shapes. The maximum laser energy on target for these shots was 428 joules. We saw no evidence for

Fig. 2.1-1. Laser power vs time on target shot 3986 from analysis of a streak camera photo.

pinhole closure in the laser spatial filters in any of these experiments.

The targets were spherical glass shells ranging in diameter from 109 μm to 290 μm and in wall thickness from 0.8 μm to 2.5 μm . The irradiance, obtained by dividing the laser power (averaged over the FWHM pulse length) by the target surface area, ranged from $2 \times 10^{13} \text{ W/cm}^2$ to $1 \times 10^{15} \text{ W/cm}^2$. The targets were filled with $\sim 2 \text{ mg/cm}^3$ of DT.

The laser system was operated in two distinct modes during this series, either with or without the plasma spatial filter (PSF) device.² The PSF changes the character of the laser beam in three ways. First, it smoothes the time averaged, spatial intensity pattern, effectively removing the hot spots typical of an unfiltered beam. Second, it broadens the bandwidth of the light incident on the target from $< 5 \text{ \AA}$ to $> 30 \text{ \AA}$. (This broadening varies from shot to shot.) Finally, the beam is slightly less focusable with the PSF installed, so that the volume enclosing 80 percent of the laser energy is about 30 μm on average.

Absorbed energy measurements were presented in the 1978 KMSF Annual Report. They showed increased absorption at lower irradiance when the plasma filter was installed. We will highlight here several distinctive features of this experimental series. Fortunately all of these features can be found in the results of target shot 3986, whose experimental parameters are given in Table 2.1-1,

this shot will be used to illustrate each of the subsections below.

2.1.1 Ion Flow Measurements

Measurements were made of the ion flow using ordinary charge collectors to examine the effects of pulse length on the ion flow and, in particular, the fast ion group, which we have studied and modeled extensively in short pulse experiments.^{3,4} A typical charge collector trace for the long pulse experiments is shown as the dashed line, curve (c), in Figure 2.1-2. It appears to be quite similar to the

Table 2.1-1

Shot 3986

115 μm diameter, 1.6 μm wall	408 Joules, 1.05 μm
1.7 mg/cm^3 DT fill	760 psec flat pulse

IRRADIANCE:	$1.3 \times 10^{15} \text{ W/cm}^2$
ABSORPTION:	122 Joules (30 %)
HOT TEMP:	8 KeV
NEUTRONS:	3.7×10^7
FAST ION ENERGY:	$\approx 20 \%$

Fig. 2.1-2. Ion charge collector voltage vs time for shot 3986 from: (a) TRHYD simulation with classical thermal conduction, (b) TRHYD simulation with flux limited transport, and (c) experiment.

traces from short pulse experiments, except that the total amount of recovered ion mass is higher. The total energy in the fast ion group is somewhat lower, however. We have examined the dependence of the fraction of absorbed energy appearing as

Fig. 2.1-3. The fraction of the absorbed energy in the fast ion group vs incident laser intensity.

fast ion energy as a function of target irradiance. The results are shown in Figure 2.1-3. We find that the scaling of $\epsilon_f = E_f/E_{abs}$ is not inconsistent with $\epsilon_f \sim I_0^{0.4 \text{ to } 0.66}$ (the straight lines shown in the Figure) over the dynamic range of the experiments. Notice that there does not seem to be a systematic difference between shots with and without the plasma spatial filter.

In these experiments we expect to set up a more nearly steady ablation process, and since the intensities are lower, we might expect to observe the "ablation signature" of Matzen and Morse.⁵ This signature, which is a strongly peaked $\delta N/\delta V$ versus V curve, or equivalently, a substantial amount of ablated mass moving at about the same velocity, was *not* observed. We have examined this "ablation signature" effect in code calculations (without using any superthermal electrons) of both classical and inhibited (flux limiter = 0.05) condi-

tions. Curves (a) and (b) of Figure 2.1-2 are plots of our preliminary attempts to find agreement between the experimental charge collector traces and our code calculation. Notice that in both cases a very large ion flow spike is predicted, quite different from the observation.

In experimental shot number 3986 (see parameters in Table 2.1-1), curve (a) shows a classical transport, code calculated, charge collector trace, and curve (b) shows a flux limited ($f = 0.05$), code calculated, charge collector trace. Curve (c) is the experimental charge collector. It is obvious that both classical and flux limited runs show the expected "ablation signature," but the experiment does not. Presently we believe that by incorporating some combination of superthermal electron acceleration of fast ions and transport inhibition on the thermals, we may be able to model the long pulse ion flow experimental results.

2.1.2 Holographic Measurements of Plasma Density

A double pulse holographic camera with the probe pulse light at $0.53 \mu\text{m}$ wavelength was used to record interference fringes due to plasma expansion during irradiation. The optical layout of this diagnostic technique was shown in Section 2.4 of the 1978 KMSF Annual Report. Interferograms were made for several shots in the long pulse series, varying the timing between the main pulse and the probe pulse.

A reconstruction of the interferograms for shot 3986 is shown in Figure 2.1-4. The interference fringes follow lines of uniform plasma density. They are approximately circular, as would be expected for a uniformly illuminated target. The actual plasma density profile can be determined by an Abel inversion of the fringe pattern. The inversions have been done by Dr. Donald Sweeney of Purdue University and are shown in Figure 2.1-5. The solid lines (curves 1, 2, and 3) indicate density profiles along different radial spokes. Figure 2.1-5 also shows two profiles computed by the hydro code TRHYD in a simulation which matches many of the experimental parameters. For densities greater than 10^{19} cm^{-3} (1 percent of critical density), they are in close agreement with TRHYD calculations. The

Fig. 2.1-4. Reconstructed holographic interferogram for shot 3986.

Fig. 2.1-5. Electron density profile obtained from the interferogram along three radial directions compared to computed profiles at two different times.

probe pulse which produced this interferogram was delayed by 60 psec relative to the start of the main laser pulse. The TRHYD simulation cannot reproduce the observed profile at 60 psec. The plasma simply does not move fast enough to match the data. There are two possible explanations for the discrepancy. First, the relative timing calibration may be in error. The optical paths are quite long and include a large number of optical components, so that subtle changes could gradually accumulate to produce timing errors between the main beam and the probe pulse. The calculated profile at 200 psec matches the experimental profile quite well. The second possible explanation would be the presence of some prepulse energy from the laser, producing plasma blowoff prior to arrival of the main pulse. Other code simulations suggest that 10-40 millijoules of pre-energy absorbed over a few nanoseconds could produce the observed profile. Such a prepulse would be near the threshold for detection in our prepulse monitor.

2.1.3 X-ray Pinhole Photographs and Code Simulations

Pinhole camera pictures of x-radiation emitted by a laser imploded spherical shell normally show

an outer ring which extends somewhat beyond the original diameter of the target. The optical density and width of the outer ring is a function of the shell thickness, the laser power, and the pulse length. If the temperature attained in the central core upon final implosion is sufficiently high, a second, inner ring is normally seen marking the interface between the shell and the fuel. Regions of enhanced x-ray emission are expected to result from relatively slow moving regions of high density plasma at elevated temperature. The traditional inner and outer rings correspond to the positions at which the shell material's velocity is low and its density is high.

Some of the pinhole images from the long-pulse experiments showed a unique feature: a ring of intense x-ray emission from a region at approximately half the original target dimension. The intermediate ring had never been observed in short-pulse experiments. Two pinhole camera photographs from shot 3986 are shown in Figure 2.1-6. These images were recorded from orthogonal viewing directions, and with different thicknesses of beryllium filters in front of the film. The left hand image shows the intermediate ring along with x-ray emission out to the original target diameter. The right hand image, taken behind thicker foil, shows the inner ring structure of the central core region more clearly.

We believe the intermediate ring, which may be seen when the laser pulse length is longer or only slightly shorter than the implosion time, indicates the position at which the conduction front stagnates during the implosion. This conclusion is based on studies using our fluid hydrodynamics code TRHYD, to simulate this phenomenon in the x-ray images. The ring arises because the spherically diverging shock from the implosion stops the inner motion of the conduction front and increases its density. The location of this ring provides quantitative information on the location of the conduction front at implosion time. The separation between this "conduction front ring" and the implosion ring is a measure of the residual tamper thickness at implosion time. Figure 2.1-7 shows plots of the position of the shell-fuel interface and conduction front as a function of time from the simulation of shot 3986. The fuel glass interface shows the implosion time to be ~ 400 psec, while the laser pulse lasted 760 psec. The temperature

Fig. 2.1-6. X-ray images from target shot 3986.

Fig. 2.1-7. Computed r - t plots and ρ - t plot for shot 3986.

profile in the glass shows a sharp rise from a few tenths of a keV over a very short distance. This transition region is called the conduction front, and its trajectory is also plotted. In this simulation, it reaches a minimum radius just after the time of peak compression, and the minimum radius is about $20\text{ }\mu\text{m}$. Also plotted is the density just behind the conduction front as a function of time. Just after peak compression time, the rebound shock ex-

Fig. 2.1-8. Comparison of experimental and computed microdensitometer traces along a diameter of pinhole camera image of shot 3986.

panding away from the center raises the density at the conduction front. This set of conditions produces a region of enhanced x-ray yield at the radial position of the stagnated conduction front. Figure 2.1-8 shows a microdensitometer trace across a diameter of one pinhole image, along with a trace computed from the simulation code. The qualitative agreement is good. For this simulation run, flux-limited thermal transport was imposed, with the

flux limit parameter set at 0.06 around the critical density surface and 0.1 inside. The film density in the central region of the pinhole image could only be matched by allowing ~ 10 percent of the absorbed energy to preheat the entire target.

The significance of this result is that, despite the thin target wall ($1.6 \mu\text{m}$) and the moderately high laser intensity ($1.3 \times 10^{15} \text{ W/cm}^2$), this target clearly does not fit the exploding pusher model of a uniformly heated shell. It shows a behavior better characterized as a mildly preheated, ablatively driven implosion.

2.2

Short Pulse Backscatter from a Gas Target

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Laboratory)*

The Brillouin backscatter mechanism^{1,2,3} has been recognized as a possibly severe problem for some laser fusion schemes, since it could strongly reduce laser energy deposition in fusion targets. Recent laser target interaction experiments^{4,5} have produced strong evidence that stimulated Brillouin backscatter is playing a role in longer gradient length plasma experiments. Laser plasma interaction experiments using CO₂ lasers and gas targets^{6,7} have produced clear evidence of Brillouin backscatter, since a direct measure of the ion acoustic-frequency shift was obtained. We were motivated by these experiments to examine the backscatter from a gas target using the Nd:glass laser which, operating with a short pulse, could perhaps minimize plasma motion effects, e.g., Doppler shift in short pulse, solid target experiments and ionization wave effects in long pulse, CO₂ laser, gas target experiments. We were particularly interested in: (1) obtaining a clear signature of Brillouin backscatter in spectral measurements, (2) observing saturation of the backscattered energy with laser power, and (3) observing any systematic spectral changes with increasing laser power.

2.2.1 Experimental Set-Up

Our experimental set-up, shown in Figure 2.2-1, is quite simple. We used a puffed nitrogen gas jet as the laser target. High pressure nitrogen (~ 10 atm) was allowed to leak through a $100 \mu\text{m}$ diameter pinhole, producing a supersonic nozzle flow into a vacuum chamber. The gas density falls off approximately as $\rho \sim \rho_0(r_0/r)^2$ in this free expansion.⁸ Here r_0 is the pinhole radius, ρ_0 the gas density at the pinhole, which was high enough to produce a critical density surface when the nitrogen was partially ionized. The laser beam was directed along the jet's axis of symmetry, up the density gradient. The beam could be focused at various positions along the axis. The electron density at the lens focus was not measured, and is crudely estimated at about 10^{21} cm^{-3} . The experiments were performed at the focal position, which maximized backscatter.

The output of the laser was approximately a 100 psec square pulse with a maximum energy of 700 mJ and a beam diameter of 40 mm. The energy was focused onto the gas jet with an f/3.5 lens. We measured prepulse energy (less than 10^{-6} of the main pulse) and both the incident and reflection energy with a fast, calibrated photodiode. Spectral measurements of the backscattered light were made with a Czerny-Turner grating spectrograph having $1.2 \text{ \AA}$ resolution. Far field focal spot measurements showed that 90 percent of the energy was contained within a $20 \mu\text{m}$ diameter, which gave a peak focused intensity of $2 \times 10^{15} \text{ W/cm}^2$.

2.2.2 Results and Discussion

Figures 2.2-2a through 2.2-2d show the backscattered spectra obtained as the incident pulse energy is increased from 50 mJ to 700 mJ, corresponding to peak focused intensities from $1.6 \times 10^{14} \text{ W/cm}^2$ to $2.2 \times 10^{14} \text{ W/cm}^2$. The backscattered energy in Figure 2.2-2a is seen to be red shifted by about $4.5 \text{ \AA}$ with no energy at the incident frequency ω_0 . This shift is typical of all the backscattered spectra even at higher incident powers. The red shift clearly indicates the Brillouin backscatter mechanism, and is in qualitative agreement with the measurements of Offenberger.^{6,7} For an in-

Fig. 2.2-1. Experimental configuration of gas target reflection and spectral measurements.

homogeneous, infinite plasma, the intensity threshold condition to be satisfied⁹ for Brillouin backscatter is given by

$$\left(\frac{V_{os}}{V_{th}}\right)^2 > \left(\frac{n_c}{n}\right) \left(\frac{8}{k_o L_t}\right) \quad (1)$$

where L_t is the temperature gradient length, V_{os} is the electron oscillatory velocity in the electric field of the wave, V_{th} is the electron thermal velocity, n is the electron density, n_c the critical density, and k_o is the vacuum wave number. For typical values of $n_c/n \approx 5$ and $L_t \approx 50 \mu\text{m}$, this threshold occurs at $(V_{os}/V_{th})^2 \approx 0.1$. The electron temperatures are estimated to be $\sim 100\text{eV}$, based on the spectral shift of the backscattered line. Since $(V_{os}/V_{th})^2$ is exceeded for all our intensities at these electron temperatures, Brillouin backscatter should occur in these experiments.

In Figure 2.2-3, we present the backscattered energy fraction versus incident pulse energy (or power, since the pulse length is constant). As can be seen, the backscatter fraction increases with incident laser power and appears to reach a maximum

at about 40 percent. These results agree well with those of Ng, et al.⁷ However, a few anomalous shots gave backscattered fractions up to 60 percent, as shown in the Figure. Because Brillouin backscatter (also sidescatter) could fall in angles outside our collection aperture (up to $f/2$), misalignment of the gas jet and optical axis may be responsible for the saturation levels being below the anomalous points. Further investigation of this point is planned. The saturation value of ~ 40 percent is also in agreement with the measurements of Ripin, et al.⁴

From the measured shift of the backscattered line, we can roughly estimate T_e , T_i from the relationship between the frequency shift and the ion acoustic speed given by

$$\frac{\omega_s}{\omega_o} = 2[(1-n/n_c)(\bar{Z}T_e + 3T_i)/m_i c^2 (1+k^2 \lambda_D^2)]^{1/2} \quad (2)$$

where ω_s is the frequency shift, ω_o is the incident frequency, Z is the ionization state, T_e , T_i are the electron and ion temperatures, m_i is the ion mass, c is the speed of light, and $k\lambda_D \ll 1$. Following Of-

Fig. 2.2-2 (a), Microdensitometer traces of back reflected spectra. Wavelength increases to the right. (b), (c), & (d). Note the location and the spectral width of the incident $1.05 \mu\text{m}$ laser energy (the lower traces). The incident power increases from $2 \times 10^{14} \text{ W/cm}^2$ in (a) to $2 \times 10^{15} \text{ W/cm}^2$ in (d).

fenberger, and assuming that the line broadening is due entirely to ion Landau damping, we can determine the ratio of $\bar{Z}T_e/T_i$ from the measured ratio of the width to the line shift. Simulations confirm that the $\text{FWHM} \approx \omega_s$. We have also assumed no other broadening mechanisms are operative, such as time-dependent temperatures, collisions, and so on. Figure 2.2-4 shows the measured line shifts and widths as a function of laser energy (or power). There is

considerable data scatter. However, at the lower powers we have approximately

$$\frac{\Delta\omega_{\text{FWHM}}}{\omega_s} = 1/3$$

which corresponds to $\bar{Z}T_e = 1.3T_i$. Using this result in Equation (2), along with the assumption $n/n_c = 0.2$, we find,

$$\bar{Z}T_e = 1.3T_i = 250 \text{ eV}.$$

Simple ionization rate estimates show that the plasma cannot be fully ionized ($Z = 7$). For our 100 psec pulse length and at these low temperatures, we estimate $Z \approx 2$. Consequently, $T_e \approx 125 \text{ eV}$ and $T_i \approx 190 \text{ eV}$. As the laser power is increased, the shift of the backscattered line does not increase as might be expected. Apparently as the power is increased, backscatter, sidescatter, and critical surface reflection increase so as to keep the energy absorbed nearly constant.

Figures 2.2-2b and 2.2-2c are typical spectra obtained as the laser power is increased. Some energy in these cases is being reflected at ω_o . Presumably the reflection is from the critical density surface, which is not moving appreciably during the pulse, since the line would thereby be Doppler shifted. Also, the Brillouin red shift does not increase at higher power, apparently indicating that the ion acoustic sound speed is not substantially increased, as might be expected. However, a further broaden-

Fig. 2.2-3. Total back reflected energy fraction. (Brillouin-shifted plus unshifted) as a function of incident laser energy (in joules).

Fig. 2.2-4. Measured Brillouin shift and line width (FWHM) as a function of incident laser energy.

ing of the backscattered feature occurs at the higher powers (see Figure 2.2-4), indicating a decrease in $\bar{Z}T_e T_i$. Qualitatively, the backscattered line is always skewed to the long wavelength side. In some shots (Figure 2.2-2d) we observed a nearly periodic, modulated structure on the long wavelength side. A similar result has been observed in particle simulation studies of the stimulated Brillouin process¹⁰ but has not been previously observed experimentally. In the simulations (1½ dimensional) the plasma was uniform at $n/n_c = 0.5$, 40 wavelengths deep, with a vacuum at the left boundary and a perfect conductor or perfect absorber at the right boundary. Other simulation parameters were: $m_i/m_e = 100$, $V_{os}/V_{th} = 0.9$, and $\bar{Z}T_e/T_i = 30$. Figure 2.2-5 shows the Fourier spectrum of the incident and reflected electric field for the case of a perfectly conducting boundary (solid lines) and the case of a perfectly absorbing boundary (dotted lines). The Brillouin reflection in both cases was about 60 percent. The multiple lines on the red side of ω_o are separated by the ion acoustic frequency. The light can also scatter off the critical surface (solid lines in Figure 2.2-5) with no shift in this simple model. The significance of multiple scattering is: (1) for each scattering the light has another chance to be absorbed by inverse bremsstrahlung; (2) each scattering off the critical surface gives the light another chance to be absorbed there; (3) each scattering

Fig. 2.2-5. Amplitude of incident and reflected spectrum (in frequency space) from the particle simulation code.

heats the ions, and (4) increases bandwidth, which acts to reduce Brillouin scattering.¹¹ We have made no attempt to achieve a quantitative fit between the experimental spectra and those from simulations, but the observed periodic pattern on the red side of ω_o is in qualitative agreement with the simulation results.

It is particularly noteworthy that energy is being reflected from the critical surface, since inverse bremsstrahlung in our low temperature plasma should be nearly 100 percent efficient in absorbing the energy which is not Brillouin backscattered. We believe this reflected energy at ω_o is not stray light, because in shots at similar laser power but at gas pressures of insufficient density to produce a critical surface we found no energy at ω_o . Also, experiments with no gas showed no reflection. The appearance of the reflected light at ω_o at the higher powers may be due to profile steepening at the critical surface. The steepened profile with a short gradient length would explain the reduced inverse bremsstrahlung. From our measured backscatter shifts, we estimate $(V_{os}/V_{th})^2 \approx 2$ at the higher powers which produce profile steepening. Further experiments are required to understand this result.

In conclusion, we have observed the clear signature of the Brillouin backscatter in a nearly stationary, long gradient length plasma. We have found the backscatter to saturate at about 40 percent. At intensities of 10^{15} W/cm^2 , we have observed reflection from the critical surface as well as the Brillouin shifted backscatter; both may be the

result of profile steepening at the critical surface. Finally, we observe increased ion heating and multiple Brillouin scattering at our highest intensities. We believe these long gradient length, short pulse measurements may isolate the absorption and back-scatter physics effectively.

2.3

Target Experiments at 0.53 μm

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Target experiments using 0.53 μm light have been performed at KMSF periodically during the last several years.¹ These experiments demonstrated a small increase in absorbed energy fraction and a reduction in fast ion energy fraction compared to similar experiments with 1.05 μm light. Recent theoretical work has indicated several potential benefits of utilizing submicron wavelength laser light in a number of processes which affect target response. A recent summary² has concluded that hot electron production, stimulated Brillouin scattering, transport inhibition, filamentation, and hydrodynamic efficiency are all sensitive to wavelength, and that in each instance wavelengths less than 1 μm are preferred.

We are presently doing a systematic experimental comparison of target response to 0.53 μm and 1.05 μm light at three intensities ($1 \times 10^{15} \text{ W/cm}^2$, $2.5 \times 10^{15} \text{ W/cm}^2$, and $6 \times 10^{15} \text{ W/cm}^2$) and three pulse lengths (100 psec, 400 psec, and 1,000 psec). Although these experiments are not yet complete, the strong interest in short wavelength data has prompted us to present some early results here.

A schematic representation of the optical diagnostics is shown in Figure 2.3-1. The frequency doubling crystals are 14 cm circular aperture KDP, cut for type II angle phase matching. They are mounted in index matching fluid between AR-coated windows. The unconverted 1.05 μm light is allowed to pass into the target chamber, but is not focused on the target by the highly chromatic 80°

focusing lenses. A silver overcoating on the ellipsoidal mirrors provides high reflectivity at 0.53 μm . The incident 0.53 μm beam profile is recorded in near and far field cameras. An Imacon 600 streak camera records the temporal profile of the incident and reflected green light. A spectrograph is used to record the spectrum of the incident and reflected 0.53 μm light, as well as the incoming 1.05 μm beam. Total energy, prepulse, and ASE measurements are made for each beam. A far field camera is also available to record light reflected from the target. The target diagnostic data presented here include absorbed energy from a plasma calorimeter, hot electron temperature from hard x-ray spectra, and ion energy from charge collectors. Data from several other instruments have not yet been fully analyzed.

The targets for these experiments were spherical glass shells filled with approximately 2 mg/cm³ of DT gas. For all the shots presented here, the targets were $80 \pm 5 \mu\text{m}$ in diameter with wall thickness of 0.9 μm to 1.2 μm . The high-intensity shot series ($6 \times 10^{15} \text{ W/cm}^2$) will use 55 μm diameter targets. The green light pulseshape was Gaussian in time and approximately 90 psec FWHM. On target peak power of 0.6 TW at 0.53 μm was typical. The target illumination optical system consisted of two lens-ellipsoidal mirror pairs, which provided reasonably uniform illumination at near normal incidence over almost all the spherical target surface.

Absorbed energy data are shown in Figure 2.3-2. The absorbed energy is determined from a single differential calorimeter viewing the target from above. The total absorption calculation assumes uniform spherical expansion of the plasma. Because of the small solid angle of the calorimeter, the absolute magnitude of the absorbed energy is not well determined, but the data scatter is small enough to allow detection of small changes in the absorbed energy fraction. Previous experience with more than one plasma energy detector has shown no large deviations from spherical symmetry in the plasma expansion, except around the target mounting post at the bottom of the target. Data from a number of shots have been condensed in Figure 2.3-2 by averaging results obtained from five shots at $(0.6 \text{ to } 0.8) \times 10^{15} \text{ W/cm}^2$, eleven shots at $(1.0 \text{ to } 1.3) \times 10^{15} \text{ W/cm}^2$, and thirteen shots at $(2.0 \text{ to } 3.0) \times 10^{15} \text{ W/cm}^2$. In addition, one point represents an

Fig. 2.3-1. Layout of optical components for 0.53 μm experiments.

Fig. 2.3-2. Absorbed energy fraction at 0.53 μm and 1.05 μm vs irradiance for 80 μm diameter, spherical glass shell targets and 90 psec pulselength.

Fig. 2.3-3. Hot electron temperature deduced from hard x-ray measurement vs $I\lambda^2$, where I is the laser intensity at the target and λ is the laser wavelength. The line is extrapolated from a fit to data taken between 10^{15} W/cm 2 and 10^{16} W/cm 2 .

eight shot average of absorption data from 1.05 μm shots around 2.8×10^{15} W/cm 2 . The error bars indicate one standard deviation from the average. The irradiance is defined as the peak laser power divided by the target surface area. The 0.53 μm data show a clear trend toward decreasing absorption with increasing irradiance. The 1.05 μm point is not significantly different from the 0.53 μm data.

The hot electron temperature determined from the hard x-ray spectrum is plotted as a function of $I\lambda^2$ in Figure 2.3-3. The x-ray spectra are measured with a ten channel, k-edge filter, PIN diode array. The diode data are fitted to a two temperature bremsstrahlung spectrum. The results shown in the Figure represent weighted averages over the specified number of shots, where the weighting factor is the inverse of the variance of the least squares fit. The error bars represent one standard deviation from the mean. The small uncertainty reflects the consistency of the temperature measurement from shot to shot. Also shown is a line extrapolated from a fit to glass shell target data taken at Livermore³ for $I\lambda^2$ values between 10^{15} and 10^{16} W/cm 2 . Our limited data are reasonably consistent with this line and support the choice of $I\lambda^2$ as the appropriate variable for scaling hot electron temperature.⁴

We have observed a marked change in the fraction of the absorbed energy which is carried away by fast ions in comparing 0.53 μm and 1.05 μm experiments. This difference is shown most clearly in the charge collector measurements. Charge collectors measure the plasma ion current as a function of time at a large distance from the target. Without an accurate knowledge of the ionic charge to mass ratio distribution they provide only a qualitative measure of the energy distribution. However, they do provide a convenient means to observe changes in the ion energy distribution as the laser wavelength is changed. In high irradiance experiments such as these, charge collector traces typically show two components: the early-arriving, fast ions with velocities greater than 10^8 cm/sec, and the more numerous but less energetic, "thermal" group, which peaks at about 7×10^7 cm/sec.⁵ An empirical relation between the average charge-to-mass ratio and the ion velocity is used to convert the current trace into an approximate energy-arrival rate.⁶ The integrated energy is divided between fast and slow ion groups, with the division made at the minimum of the current-versus-time trace. The fast ion energy fraction is plotted in Figure 2.3-4 as a function of irradiance. The fast

Fig. 2.3-4. Estimated fast ion energy fraction of the total absorbed energy as a function of irradiance.

ion energy fraction is defined here as the charge collector, fast ion energy (averaged between two collectors) divided by the total absorbed energy measured by the plasma calorimeter. Each point represents a single shot. Despite substantial scatter in the $1.05 \mu\text{m}$ data, it is clear that considerably less energy goes into the fast ion group in the $0.53 \mu\text{m}$ experiments. In the $2\text{--}3 \times 10^{15} \text{ W/cm}^2$ range,

Fig. 2.3-5. Relative slow ion energy fraction of the total absorbed energy vs irradiance. The absolute slow ion energy scale is not well determined.

Fig. 2.3-6. X-ray energy fraction of the total absorbed energy vs irradiance. X-ray energy is measured by unfiltered TLD chips behind a magnetic field which deflects all charged particles.

roughly three times as much energy is carried away by fast ions for targets illuminated with $1.05 \mu\text{m}$ light.

A similar plot has been made here for the slow ion component of the charge collector data using an analogous definition for the slow ion energy fraction. For these low velocity ions, the assumed charge-to-mass ratios are highly uncertain, and only a relative energy variation can be determined. Data shown in Figure 2.3-5 are from the same shots used in Figure 2.3-4. A larger share of the absorbed energy appears in the slow ion group at $0.53 \mu\text{m}$ than at $1.05 \mu\text{m}$. These two plots indicate that at the shorter wavelength illumination, the energy distribution changes in favor of the slow ion group.

The total x-ray energy yield is sampled using calibrated TLD chips behind a magnet.⁷ The available data are plotted in Figure 2.3-6. The x-ray energy fraction is the ratio of the x-ray energy determined by the TLD's to the total absorbed energy measured by the plasma calorimeter. The data at both laser wavelengths fall in the 5-10 percent range, with no apparent wavelength dependence.

The results to date appear to verify the expectations of (1) $I\lambda^2$ scaling of the hot electron temperature, (2) a reduction in the fast ion energy fraction

for the shorter wavelength, and (3) minimal change in absorbed energy fraction for short pulse lengths. We anticipate learning much more about stimulated scattering and transport inhibition as we continue to longer pulse lengths, add more 1.05 μm comparison experiments, and study results from other diagnostic instruments.

2.4

Implosion Physics

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This section describes a simple model of spherical implosions and our high density experiments using cryogenic targets.

2.4.1 Spherical Implosion Model

Most calculations of the implosion of spherical shell, inertial confinement targets are performed using extensive hydrodynamic computer codes such as LASNEX. With most complex models the calculations have difficulty producing physical insight and, in some cases, scaling information. In this Section we describe a simple model of the ablative implosion of a high aspect ratio (i.e., shell radius to shell thickness ratio) spherical shell. The model is similar in approach to that of Rosenbluth's¹ snowplow model. We determined the scaling of the implosion time in terms of the ablation pressure and the shell parameters, such as diameter, wall thickness, and density, and compared these to complete hydrodynamic code calculations. We also found the energy transfer from ablation pressure to shell implosion kinetic energy to be highly efficient. It may be possible to "attach" a simple heat transport calculation to our implosion model to describe the laser driven ablation implosion process.² The model may also be useful for determining implosion scaling for other drivers (e.g., ion beam).

The Model Equations. Consider the implosion of a thin spherical shell caused by the reactive force of ablating material (the "rocket effect") and the force

due to the external pressure resulting from the deposition of energy. We assume: (1) a high aspect-ratio shell, (2) a constant ablation density (ρ) and ablation velocity (V_a) in a coordinate system moving with the shell, and (3) the shell remains thin, i.e., it can be represented by an infinitely thin, finite mass shell at $R(t)$. In addition, the fuel gas is adiabatically compressed and "turns around" the imploding shell ($p_g V_g^\gamma = \text{constant}$).

Newton's Law is written as,

$$\frac{d}{dt} (M_s \dot{R}) = (\dot{R} + V_s) \dot{M}_s + 4\pi R^2 (p_g - p_a) \quad (1)$$

where M_s is the shell mass at time t , R is the location of the shell, γ is the adiabatic constant, p_g the internal gas pressure and p_a is the ablation pressure. The shell loses mass according to $\dot{M}_s = -4\pi R^2 \rho_a V_a$. Making these equations non-dimensional, we find the characteristic implosion time τ is given by,

$$\tau = \left(\frac{M_o}{4\pi R_o p_a} \right)^{1/2} \quad (2)$$

and the differential equations become,

$$\eta \ddot{y} = -(M^2 + 1 - \beta y^{-3\gamma}) y^2 \quad (3)$$

$$\text{where } \eta = 1 - \int_0^{t'} M \alpha y^2 dt', \quad (4)$$

$$\eta(t) = \frac{M_s(t)}{M_s(o)}, \quad y(t) = \frac{R(t)}{R(o)}, \quad t' = t/\tau,$$

$$M_s(o) = 4\pi R_o^2 \delta \rho_{solid}, \quad \alpha = \left(\frac{R(o)}{\delta} \frac{\rho_a}{\rho_{solid}} \right)^{1/2},$$

$$\beta = \frac{P_g(o)}{P_a}.$$

M is the Mach number and C_a is the atmospheric sound speed ($C_a^2 = p_a / \rho_a$), δ is the shell thickness, and ρ_{solid} is the shell density. The initial conditions are $\eta(0) = 1$, $y(0) = 1$, $\dot{y}(0) = 0$.

We have numerically integrated this system of equations for various choices of input parameters M and α . Figure 2.4-1 shows an example for which $M = 0.2$, $\alpha = 3.48$, $\beta = .002$. Notice also that the implosion time, t'_{imp} , depends on the parameters in τ and on the parameters M and α . This Figure in-

Fig. 2.4-1. Normalized radius vs time plot, y vs t' , $\dot{y}$ vs t' , and η vs t' for $M=0.2$, $\alpha=3.48$, $\beta=0.002$.

indicates that ~ 50 percent of the shell mass was ablated in the interval $t' = t'_{imp}$. Figures 2.4-2 and 2.4-3 show the dependence of the unablated shell mass and t'_{imp} on M and α . Notice that the parameter a contains the shell aspect ratio, $R(o)/\delta$.

The Energy Integrals and Pulse Shape. It is useful to compute a few energies associated with this simple implosion system. We can multiply Equation (1) by $\dot{R}$ and integrate to find the following energies (we have set $p_g = 0$): the shell kinetic energy,

$$E_{k.e.}^{shell} = \frac{3}{2} E_o \left(\frac{dy}{dt'} \right)^2 \eta = \frac{1}{2} M_s \dot{R}^2$$

the "blow off" kinetic energy,

$$\begin{aligned} E_{k.e.}^{blow-off} &= \frac{1}{2} \int_0^t \dot{M}_s (\dot{R} + V_a)^2 dt \\ &= \frac{3}{2} E_o \int_0^{t'} \dot{\eta} \left(\dot{y} + \frac{M}{a} \right)^2 dt', \end{aligned}$$

and the energies which can be considered as *input* energies, which are negative. The mechanical energy produced by the ablation pressure,

$$E_{mech} = - \int_0^t 4\pi R^2 p_a \dot{R} dt = -3E_o \int_0^{t'} y^2 \dot{y} dt',$$

Fig. 2.4-2. Plot of t'_{imp} and $\eta(t'_{imp})$ vs M for $\alpha = 1$.

Fig. 2.4-3. Plot of t'_{imp} and $\eta(t'_{imp})$ vs α for $M = 0.2$.

and the "rocket exhaust" energy,

$$E_{exh} = \frac{1}{2} \int_0^t \dot{M}_s V_a^2 dt = \frac{3}{2} E_o M^2 / a^2 \int_0^{t'} \dot{\eta} dt'.$$

In the above, the scale energy E_o given by,

$$E_o = \frac{4\pi}{3} R_o^3 p_a$$

Fig. 2.4-4. Plot of ϵ , E_{exh} , $E_{k.e.}^{pusher}$, $E_{k.e.}^{blow off}$, E_{mech} and the pulse shape P vs t' for $M = 0.2$, $\alpha = 3.48$.

is the energy needed to fill the original shell volume at the ablation pressure p_a . The energy conservation equation reads,

$$E_{k.e.}^{shell} + E_{k.e.}^{blow off} = E_{mech} + E_{exh}.$$

It is interesting, and perhaps unexpected, to find that the transfer of energy from the "mechanical" input (i.e., from ablation pressure) to the kinetic energy of the shell is rather high in almost all cases. Physically this means the fraction of energy turned into atmospheric pressure at the ablation surface is effectively utilized in the shell implosion process. If we define the energy transfer efficiency relative to the ablation surface as

$$\epsilon = \frac{E_{k.e.}^{shell}}{E_{mech} + E_{exh}}$$

we can compute this efficiency for various choices of parameters M and α . Figure 2.4-4 shows plots of the various dimensionless energies and the efficiency as a function of time. In these Figures we also present the power input pulse shape required to produce our assumed constant ablation pressure p_a . This pulse shape is simply computed from the expression $P = E_{mech} + E_{exh}$. By allowing $p_a = p_a(t)$, it would be possible to determine the ablation pressure pulse shape which would produce a perfectly isentropic compression³ of an imploding shell.

Fig. 2.4-5. Plot of the locus of points M , α for which half the shell mass is ablated, $\eta = 0.5$ at $t = t_{imp}$.

Comparisons to Hydro Code Calculations. To compare this simple model with complete heat transport, hydrodynamic code calculations, we need to specify the values of M and α . It is interesting to note that the Mach number of the ablating material is a measure of the relative contributions of the "rocket reaction" force and the ablation pressure forces driving the shell. This can be seen by examining Equation (3). For smaller values of M (which are expected on the basis of steady rate, spherical hydro considerations),⁴ the ablation pressure force dominates; at $M=1$, both forces contribute about equally. In many cases of interest it is found from hydro code runs that about 50 percent of the mass is ablated from an imploding shell at the peak compression time, if the energy deposition pulse is longer than the implosion time. We are therefore most interested in the M , α points for $\eta=0.5$ at $t'=t'_{imp}$. Figure 2.4-5 is the plot of these points found from the model. We might expect different values of M for different transport models of the ablation process (e.g., strong or weak heat flux inhibition).⁵ Also shown in Figure 2.4-5 is the implosion time in units of τ .

We have chosen to fit our simple model to a hydro code calculation of the strongly flux limited ($f=.05$) implosion of an 80 μm diameter spherical glass shell with a 2 μm wall, which absorbed 50 joules of 1.05 μm laser light in 250 psec. X-ray preheat was "turned off" for this calculation, as it

Fig. 2.4-6. A comparison of the simple model R vs t and the complete hydro code calculation.

tends to further expand the shell. Figure 2.4-6 shows the agreement between the simple model (dots) and the complete hydro code run (straight lines). The lines from the hydro code indicate the location of the inward moving shell material. the simple model parameters were $M=0.2$, $\alpha=3.48$, and $\beta=.002$. Perhaps more important is the ablation pressure comparison. The o's give the peak values of the ablation pressure versus time from the hydro code, whereas the simple model had constant ablation pressure shown by the straight line. Notice that the peak pressure from the code increases ($\sim 1/R$) at the later stages of the implosion. This pressure increase due to spherical convergence is not contained in our simple calculation.

The ablation density in the model is also constant and is compared in Figure 2.4-7 to the ablation density from the code calculation. The "fitting point" in the hydro code was found to be a best match with the simple model if we chose the peak pressure point in radius as the effective ablation point. Figure 2.4-7 shows the hydro code pressure, density, sound speed ($C_s = \sqrt{P/\rho}$), and flow Mach number (in the frame moving with the shell). It appears that the R versus t trajectory from the simple model can be made to agree rather well with the full hydro code calculation for our parameter values.

Fig. 2.4-7. Hydro-code density, pressure, sound speed, and flow velocity profiles at $t = 60$ psec showing the model fitting point.

Implosion Time Scaling. The scaling of the implosion time depends on τ as well as on α and M . For small M (which seems close to the full hydro simulations), it is interesting to extract the scaling of t'_{imp} with α at fixed M . Figure 2.4-3 shows the plot of t'_{imp} as a function of α for $M=0.2$. Also plotted is the shell mass at peak compression, $\eta(t'_{imp})$. Resulting scaling gives $t'_{imp} = \tau(1.7 - 0.8\alpha)$. The implosion time scaling with aspect ratio ρ_a/ρ_{solid} , C_a , and δ is

$$t'_{imp} = \frac{\delta}{C_a} \left(\frac{\rho_{solid}}{\rho_a} \right) a [1.7 - 0.8a].$$

This Equation shows that the implosion time is a multiple of the acoustic transit time (at the ablation surface sound speed) through the shell thickness.

Compression Scaling. The simple model can be used to examine the scaling of peak compression upon the amount of preheat in the core gas. In Equation (1), we set

$$p_g = p_o \left(\frac{R_o}{R(t)} \right)^{3\gamma} = p_o y^{-3\gamma}$$

and $\beta = p_o/p_a$. Figure 2.4-8 shows the peak compression, i.e., $C = (I/y_{min})^3$ as a function of the initial gas pressure. If the gas is preheated at close to $t=0$, or before the implosion has proceeded very far, this curve shows how the compression scales with the

Fig. 2.4-8. Peak compression, C vs $\beta = p_o/p_a$, $\gamma = 5/3$, and $\gamma = 1$.

preheat pressure. The curve shows that $C \sim \beta^{-1.4}$ for an assumed adiabatic index of $\gamma = 5/3$. If we take an isothermal index, $\gamma = 1$, we find the compression scaling given in Figure 2.4-8, which is not a simple power law but indicates smaller compressions for the same preheat pressure ratio.

Conclusions. We have found that this simple, spherical, ablative compression model can reproduce some aspects of the implosion of thin-walled, spherical shell targets. In particular, the implosion time scaling with shell parameters, the R - t trajectory, and shell velocity agree well with the more complete hydro code calculations. Also, it is clear from the model that ablation pressure most strongly influences target implosion, and the energy delivered to the shell from the mechanical ablation pressure is the efficient energy transfer mechanism.

Future work on a model incorporating heat-transport, from which the ablation pressure is determined from the corona conditions, may lead to further understanding of the implosion of thin shell targets.

2.4.2 High Density Experiments and Code Simulations

In our experiments to produce high compression, we have used cryogenic targets and included a plasma spatial filter in the laser configuration. The two principal diagnostics used to characterize target

compression were alpha particle detection and x-ray photography. Alpha particle spectra are downshifted and broadened in high density compressions as a result of the fuel and tamper $\langle \rho R \rangle$. x-ray streak photography is employed to verify that the alpha particle emission occurs after the laser pulse and therefore is not influenced by the laser produced electrostatic fields in the corona. In addition to these two diagnostics, radiation chemistry can be used to determine the $\langle \rho R \rangle$ of the tamper. For an accurate measurement of tamper $\langle \rho R \rangle$ in the range of interest, approximately 10^{-3} g/cm², a neutron yield greater than 10^8 is required. With the alpha particle, x-ray, and radiation chemistry measurements, we can determine the fuel density, the fuel $\langle \rho R \rangle$, the tamper $\langle \rho R \rangle$, and the ion and electron temperatures, thereby completely describing conditions in the imploded core.

A schematic drawing of the fuel and tamper regions is shown in Figure 2.4-9, where δ refers to the thickness at peak compression of the glass tamper, which confines the fuel. Alpha particles will traverse different amounts of material, depending on the angle θ . Integrating over θ to obtain an alpha particle energy spectrum, it is found that a significant broadening of the spectrum results from transit of the alpha particle through the fuel portion of the core. The energy loss has been calculated as a function of the $\langle \rho R \rangle$ of the fuel and tamper for different electron temperatures. The alpha particle energy spectrum for a given electron temperature and for the geometry of Figure 2.4-9 is shown in Figure 2.4-10. The downshift in the high energy end of the spectrum is caused by the $\langle \rho R \rangle$ in the fuel. When the density of the fuel compared to that of the tamper is increased, the profiles are skewed toward the high energy region. Since these relative values of density cannot be easily measured, this effect is potentially very useful, provided the spectra can be measured accurately. Various laboratories have investigated the radiochemistry involved in determining the $\langle \rho R \rangle$ of the tamper. It should be emphasized that it is possible to determine exactly the $\langle \rho R \rangle$ in the tamper by measuring the neutron yield and the number of activated particles.

For the cryogenic experiments, we used a target filled at room temperature with 100 atm of DT gas. The glass shells were typically 80 μ m to 90 μ m in diameter with a wall ~ 1 μ m thick. The condensed

Fig. 2.4-9. Schematic drawing for calculating the energy loss of an alpha particle passing through a distance f in the fuel and a distance g in the tamper.

fuel is approximately $1.5 \mu\text{m}$ thick when uniformly distributed on the inside of the shell. Helium gas flows past the target in the shroud surrounding the target, thereby maintaining a cold environment. Sapphire windows in the shroud allow a YAG laser to heat the shell, vaporizing the fuel. When the laser is turned off, the fuel freezes quickly and uniformly on the inside of the shell. The shroud surrounding the target is retracted just before the laser is fired. As the entire shroud must be out of the focal volume before the laser beam arrives at the target, the retraction begins approximately three milliseconds before the target shot. The retraction device is attached to the ceiling and must still be in motion when the laser hits the target to avoid target vibration. Figure 2.4-11 shows an interferogram of the fuel layer in a cryogenic target. There is a slight nonuniformity of the cryogenic layer in the target shown, but the variation is less than 10 percent in fuel thickness.

One feature very important in understanding high compression experiments is why cryogenic targets achieve higher densities than gas filled

Fig. 2.4-10. Alpha particle energy spectra.

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Fig. 2.4-16. Alpha particle spectra with the form of the fuel as (a) gaseous, (b) cryogenic liquid layer, and (c) cryogenic solid layer.

Fig. 2.4-17. Calculated alpha particle energy spectrum compared to the measured spectrum (points) of shot 3723.

density is shown in Figure 2.4-14. The cryogenic simulation shows the compressed fuel was contained within a region of $7.6 \mu\text{m}$ radius, and for the noncryogenic case the compressed fuel radius was $14 \mu\text{m}$. The ion temperature is lower in the cryogenic case, since the internal energy is approximately the same. Although higher fuel densities and lower fuel temperatures were calculated for the cryogenic case, the neutron yields turn out to be very similar in both target configurations.

The implosion in cryogenic target experiments is determined from a time integrated x-ray pinhole photograph. A densitometer trace of the x-ray image is shown in Figure 2.4-15. This measurement shows the compressed fuel region was not completely spherical, probably a result of the cryogenic layer not being completely uniform. The image of the implosion shows that the compressed fuel region was $13.6 \mu\text{m}$ long in the horizontal direction and $9.5 \mu\text{m}$ long in the vertical direction. This gives a volume reduction of approximately 350 for an average fuel density of 7 g/cm^3 .

In addition to the x-ray images, alpha particle spectra were taken for these experiments. As shown in Figure 2.4-16, the spectrum for shot 2833 was produced in a cryogenic target experiment in which the fuel was in gaseous form. In this case we

observe the alpha peak to be at 3.5 MeV. The alpha spectrum shown for target shot 3046 is from a cryogenic target with a liquid layer having a factor of 2 variation in thickness. This resulted in a nonuniform implosion as measured by the x-ray camera. The downshift in the alpha spectrum is approximately 5 percent of the initial energy. The alpha spectrum for target shot 3723 was taken from the cryogenic solid layer target experiment of Figure 2.4-15. It produced a downshift of 9 percent of the initial energy, the other feature of interest being the width of the spectrum (375 keV). This is much too large to ascribe to thermal broadening alone, as it would indicate an ion temperature of 4.5 keV. The corresponding neutron yield would then be 2 orders of magnitude higher than measured. We believe this broadening effect is due to the large $\langle \rho R \rangle$ in the fuel, which we estimate from the x-ray image to be approximately $5 \times 10^{-3} \text{ g/cm}^2$. Figure 2.4-17 shows a calculation of the broadening of the alpha spectrum due to a $\langle \rho R \rangle$ of $5 \times 10^{-3} \text{ g/cm}^2$ in the fuel; it is based on the path length variation in the alpha particle trajectory, as described previously in Figure 2.4-10.

2.5

Diagnostic Development

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Efforts to extend our diagnostic capabilities centered on x-ray streak camera measurements, $\langle \rho R \rangle$ measurements by detection of elastically scattered deuterons, and holographic measurement of pre-energy. The measurement instruments and procedures, together with their utility and accuracy are discussed below.

2.5.1 X-ray Streak Camera Measurements

The x-ray streak camera, which was obtained from G.E.A.R., Inc., was used successfully to record time-resolved, x-ray pinhole photographs and x-ray spectral distributions. The camera's specifications and alignment procedures have been reported previously in the 1978 Annual Report.

SHOT 4177

KMSF MICROSHELL[®] TARGET

78 μm dia., 1.0 μm wall

10 atm DT fill

44 nanograms

LASER ENERGY ON TARGET

46.1 joules in an 80 psec pulse

Fig. 2.5-1. Time resolved x-ray spectral distribution (λ vs t) for shot 4177: (a) Reproduction of film recording with time scans through 25 μm , 75 μm , and 275 μm Be (the blackening of the lower part of the film is an artifact of the camera); (b) Densitometer measurements of the scans indicating electron temperature, T_e , at peak compression of 0.8 keV to 1.1 keV.

An x-ray image typical of the results obtained for spectral distributions is shown in Figure 2.5-1a. In Figure 2.5-1b we have plotted the optical density obtained in scanning the image along the lines designated as scans A, B, and C in Figure 2.5-1a. These densitometer scans A, B, and C cross the regions in which the target's x-ray emission was filtered by 25 μm Be, 75 μm Be, and 275 μm Be respectively.

Two maxima in the density plots of scans A and B and a single maximum in scan C can be seen in

Table 2.5-1
KMSF x-ray Streak Camera Sensitivity

Non-Imaging Mode (Without Streaking)

X-ray source emission (with peak energy at ~ 2 keV)	= 0.1 joule in 4π
X-ray target-to-photocathode distance	= 39.8 cm
Measurable average optical density, above fog, on Kodak 2475 film developed in DK-50 for 15 min with a 2-stop-push using factor 8 [®]	= 0.1
Resulting (static) sensitivity	~ 13 ergs/cm ²

Imaging Mode (With Streaking)

Target-to-pinhole distance	= 0.8 cm
Pinhole diameter	= 12 μm
Magnification	= 49
X-ray source emission (with peak energy at ~ 2 keV)	= 0.2 joule in 4π
Compressed x-ray emitting diameter (obtained from an associated pinhole camera image)	= 15 μm
Sweep rate	= 25 psec/mm
Average optical film density at peak compression, for Kodak 2475 film developed in DK-50 for 15 min with a 2-stop-push using factor 8 [®]	= 0.3
Corresponding camera sensitivity	~ 66 ergs/cm ² /8psec
or, for an optical film density of 0.1	~ 30 ergs/cm ² /8psec

95 μm dia., 1.3 μm wall glass Microshell[®] target
6.9 atm DT fuel
92.3 joules on target in 99 psec
24 joules absorbed energy
1.97 joules x-ray energy in 4π

25 μm dia. pinhole imaging
Magnification = 65
No added filters
Inverse sweep velocity = 25 psec/mm
No. of usable resolution elements = 4

Fig. 2.5-2. X vs t recording for shot 4210 showing a measurable implosion time of ~ 144 psec.

Fig. 2.5-3. Comparisons of measured implosion times with predictions from the implosion model of Section 2.4.1.

Figure 2.5-1b. The maxima occurring late in the implosion (i.e., at ~ 220 psec to 230 psec) are produced by the target's x-ray emission during the interval in which the motion of the glass tamper is being reversed. The earlier maxima, not resolved in scan C, are produced by emission during the initial laser irradiation of the thin stationary target. The electron temperatures shown in Figure 2.5-1b were determined from the ratios of the filtered x-ray fluxes at or near tamper turn around. We obtained a value of ~ 1.1 keV with a variation of ~ 10 percent.

For the time resolved spatial distribution of the target's x-ray emission, we used a pinhole in conjunction with the target alignment system to obtain target images at the camera slit. Optimal pinhole imaging is obtained by matching pinhole resolution to camera spatial resolution: pinhole diameter (mm) $\times$ magnification = [camera spatial resolution (lines/mm)] $^{-1}$. Limitations of camera sensitivity and geometry (e.g., placement of the camera slit 40 cm from the target) did not permit optimal operation. The camera, whose detector was a 210 Å gold photocathode on a 25 μm Be substrate, had a sensitivity of 13 ergs/cm 2 in the non-imaging mode and 30 ergs/cm 2 /8 psec with imaging for an optical film density of 0.1 (see Table 2.5-1). Comparisons of implosion times in both imaging and non-imaging modes with those predicted from the simple implosion model of Section 2.4.1 are shown in Figure 2.5-2.

Figure 2.5-2 depicts shot 4210 where 1.97 joules of total x-ray emission (in 4π) was sufficient to record the compression "wings" of the imploding target. To do this, a pinhole diameter of 25 μm was required, limiting the number of resolution elements obtainable to ~ 4 for the 95 μm diameter target. Figure 2.5-3 shows comparisons of measured implosion times with predictions from the model in Section 2.4.1.

2.5.2 A Fuel $\langle \rho R \rangle$ Measurement Using Elastically Scattered Deuterons

Achievement of thermonuclear burn conditions is the principal objective of inertial confinement fusion research. The reaction rate is proportional to the density ρ of the compressed DT fuel, and the confinement time is proportional to the compressed fuel radius R . The energy deposition by fusion generated alpha particles is proportional to the density radius product $\langle \rho R \rangle$. Large values of $\langle \rho R \rangle$ are clearly necessary to produce thermonuclear burn, and techniques to measure $\langle \rho R \rangle$ are required. A variety of diagnostic techniques, including x-ray imaging,¹ x-ray backlighting,² alpha particle imaging,³ x-ray line broadening,⁴ and charged particle energy loss,⁵ have been developed to measure or infer parameters of the compressed core. A radiochemistry technique, which measures $\langle \rho R \rangle$ for the compressed tamper material by means of neutron activation,⁶ has also been used to infer fuel $\langle \rho R \rangle$ values.⁷ All of these techniques were reviewed in 1977 by Pan and Larsen and none was found to be both accurate and generally applicable over a wide range of experimental configurations.⁸

This Section analyzes a technique not discussed by Pan and Larsen for directly measuring $\langle \rho R \rangle$ in the target core by counting the number of deuterons which have been elastically scattered by fusion neutrons. In a sense this technique is closely related to the radiochemistry technique. In each case, as Figure 2.5-4 illustrates, the measured quantity is a product of a nuclear reaction induced by a 14 MeV fusion neutron. The radiochemistry technique measures the product of the neutron yield and the tamper $\langle \rho R \rangle$, while the elastic recoil technique measures the product of the yield and the fuel $\langle \rho R \rangle$.

Among the many neutron induced reactions in

Fig. 2.5-4. Schematic of compressed target core showing two possible neutron reactions.

the DT fuel, deuteron elastic scattering is particularly attractive as a $\langle \rho R \rangle$ diagnostic because of its large, well measured cross section, a relatively narrow peak in the deuteron energy distribution, and the possibility of a simple detection device.

The principle of using elastically scattered deuterons to measure the fuel $\langle \rho R \rangle$ is shown in Figure 2.5-4. The number of scattered deuterons is given by the product of the neutron yield, the elastic scattering cross section, and the average integral of the deuterium density along all the possible neutron paths.

The measured total and differential cross sections for elastic scattering of 14 MeV neutrons by deuterium have been published by Horsley.⁹ The differential cross section is shown in Figure 2.5-5, converted from the standard $d\sigma/d\Omega$ versus center-of-mass scattering angle to $d\sigma/dE$ versus lab energy of the recoil deuteron. This energy spectrum has a rather sharp peak between 10 MeV and 12.5 MeV, and it rises again at low energy values. Our detection scheme uses a tantalum foil to block charged-particle debris from the target. The foil thickness can be chosen to stop all deuterons of energy less than 9 MeV, so that the detected deuteron number

Fig. 2.5-5. Elastic scattering, neutrons on deuterium, as a function of recoil deuteron energy.

will not be very sensitive to the exact cutoff of the foil.

An estimate of the number of recoil deuterons is obtained by assuming the fuel to be of uniform density ρ in a sphere of radius R . Then

$$\int \rho n_D dr = \frac{3}{4} n_D R = \frac{3}{4} (\rho R) (n_D / \rho).$$

In a 60:40 D:T fuel mix, $(n_D / \rho) = 1.5 \times 10^{23} \text{ g}^{-1}$. Integrating the cross section gives

$$\int_{10 \text{ MeV}}^{12.5 \text{ MeV}} (d\sigma/dE) dE = 0.09 \text{ barns},$$

so that finally the number of recoil deuterons is $1 \times 10^{-2} N n (\rho R)_F$ (for $E_o > 10 \text{ MeV}$), when (ρR) is given in g/cm^2 . For comparison, the number of $^{28}\text{Si}(n,p)^{28}\text{Al}^*$ reactions estimated in the same way is $2 \times 10^{-3} N_n (\rho R)_T$.

As deuterons traverse the compressed fuel and tamper, they will lose energy due to collisions modifying the energy spectrum at the detector. In the limit that the deuterons lose so much energy that they cannot be distinguished from other charged

particles, they are no longer available as a diagnostic tool. In order to study the range of $\langle \rho R \rangle$ over which this technique can be used, a Monte Carlo particle tracking code was written to compute the deuteron energy spectrum from a model target.

The compressed target is taken to be a uniform density DT sphere at fixed temperature and surrounded by a shell of uniform density, uniform temperature "glass" plasma made up of electrons and $A = 20$, $Z = 10$ ions. Neutrons are generated uniformly throughout the fuel with energies distributed according to the fuel temperature.¹⁰ Each neutron scatters a deuteron at a randomly determined point, with the scattering angle (and hence recoil energy) distributed according to the experimental cross section. The deuteron energy loss is then calculated in $0.1 \mu\text{m}$ steps along a straight-line path out of the target.¹¹ The energy loss calculation is that described in Reference 5, with two additional terms: one accounts for nuclear (non-Coulomb) collisions in the fuel,¹² and the other for energy loss due to collective excitation of the plasma.¹³

Several calculated deuteron energy spectra are shown in Figures 2.5-6 and 2.5-7. For these comparisons, the geometry and temperatures were fixed, and the densities were varied. The spectrum in Figure 2.5-6a has $(\rho R)_F = (\rho R)_T = 10^{-3} \text{g/cm}^2$. The resulting spectrum is essentially unchanged from that of Figure 2.5-5 at energies above a few MeV. When $(\rho R)_F = (\rho R)_T = 10^{-2} \text{g/cm}^2$, as in Figure 2.5-6b, the high energy peak has begun to shift down, but it is still a well defined group. Figure 2.5-7a shows the spectrum for a situation in which $(\rho R)_T$ is large (10^{-1}g/cm^2) and $(\rho R)_F$ is small (10^{-3}g/cm^2). The high energy peak remains narrow even though substantially shifted. When $(\rho R)_F$ is large (10^{-1}g/cm^2) and $(\rho R)_T$ small (10^{-3}g/cm^2), as in Figure 2.5-7b, the peak is substantially broadened due to the much greater variation in path lengths through the fuel compared to the tamper. In general, the high energy deuteron peak remains useful as a fuel $\langle \rho R \rangle$ diagnostic for combined $\langle \rho R \rangle$ values of a few times 10^{-2}g/cm^2 .

The experimental problem is to detect a small number of deuterons with energies between 9 MeV and 12.5 MeV in the presence of substantial numbers of lower energy charged particles. The tech-

Fig. 2.5-6. Deuteron energy spectra computed by the Monte Carlo code from a total of 4000 deuterons in each case: (a) $\rho_F = \rho_T = 1 \text{g/cm}^3$; (b) $\rho_F = \rho_T = 10 \text{g/cm}^3$. The number of deuterons to the right of the arrow is 14% of the total in each of the spectra in Figures 2.5-6 and 2.5-7. Electron and ion temperatures of 1 keV were used in both fuel and tamper.

nique suggested here is to use a nuclear track recording polymer placed close to the target for large solid angle and behind a suitable filter for discrimination against background particles. The arrangement is shown schematically in Figure 2.5-8. A polymer known as CR-39 was shown by Cartwright, Shirk, and Price¹⁴ to be sensitive enough to record 1 MeV protons and 6 MeV alpha particles.

Fig. 2.5-7. Continuation of Figure 2.5-6: (a) $\rho_F = 1 \text{ g/cm}^3$, $\rho_T = 100 \text{ g/cm}^3$; (b) $\rho_F = 100 \text{ g/cm}^3$, $\rho_T = 1 \text{ g/cm}^3$.

More recently, excellent tracks have been observed in CR-39 for protons at 4.3 MeV¹⁵ and even up to 6 MeV,¹⁶ well beyond the limited range of the more familiar cellulose nitrate films. According to Cartwright et al, "Its uniformity of response, high sensitivity, and superb optical quality make it ideal for identification of mass and charge of nuclear particles and for other applications." The tracks are revealed by etching the CR-39 in NaOH. Track diameters depend on etch conditions, but 10 μm to 20 μm is a typical dimension.

A simple detection system could use a 1 cm radius disk of CR-39 behind a 150- μm tantalum foil at 2.5 cm from the target. With this solid angle, the num-

Fig. 2.5-8. Schematic representation of the detection technique, showing two possible sources of background tracks.

ber of detected deuterons is $4 \times 10^{-4} N_n(\rho R)_F$. Yield $\langle \rho R \rangle$ products of $\sim 10^5$ would be required to make a satisfactory measurement.

The tantalum foil thickness is sufficient to completely stop fast ions from the target surface, fusion alpha particles, and recoil tritons from the fuel. Recoil deuterons entering the foil with energies from 11 MeV to 12.5 MeV will emerge in the range from 1 MeV to 5 MeV, which is readily recorded by CR-39. In situations of higher $\langle \rho R \rangle$, where the energy loss of the deuterons becomes significant, a thinner foil would be required.

There are two sources of background tracks indicated in Figure 2.5-8 which might or might not prove troublesome. These are elastically scattered protons from the tamper or from CR-39 itself. The total number of proton recoils in the etched layer of the CR-39 will be about 10^{-6} times the neutron yield. (The number of deuteron events is proportional to the yield — $\langle \rho R \rangle$ product.) Of these, most will produce tracks which are readily distinguished from deuteron tracks on the basis of diameter or shape. Recoil protons from a hydrogenic tamper material could be a more serious matter, depending on the detailed appearance of proton and deuteron tracks for a given set of etch conditions.

In summary, fuel $\langle \rho R \rangle$ can be measured by counting the number of high energy deuterons elastically scattered by fusion neutrons. The minimum requirement is $N_n(\rho R)_F \geq 10^5 \text{ g/cm}^2$ and the maximum $\langle \rho R \rangle$ is limited to a few times 10^{-2} g/cm^2 .

The only background problem which is anticipated is the presence of recoil protons from neutron scattering in the tamper or in the detector itself.

From a more general point of view, the recoil deuteron measurement is only one of many possible charged particle measurements that might be developed as fusion target diagnostics; another might be recoil tritons. Other possibilities are suggested by Figure 2.5-7, which shows how energy spectra for core produced particles change in shape if the ratio $(\rho R)_F/(\rho R)_T$ is far from unity.

Finally, the properties of CR-39 as a detector for low-Z ions in the few MeV energy range might be utilized in an inertial fusion context in a variety of measurements. A program to measure track characteristics of hydrogen and helium isotopes over a full range of energies would help in defining the potential uses of this interesting material.

2.5.3 Pre-Energy Measurement Using Holography

The pulses amplified to 2TW in Chroma I are now generated in the actively mode locked oscillator. The pulse width at half maximum of single pulses is nominally 80 psec. Prior to the Chroma I upgrade, the pulses generated in the passively mode locked oscillator were 40 psec to 50 psec. Observed differences in experiments involving the two lasers and their respective oscillators raised questions about the effects of the pre-energy of the pulses. Interest centered on the 500 psec preceding the peak of the first main pulse in a series of stacked pulses.

The KMSF holographic probe (described in the KMSF 1978 Annual Report, Section 2.4.6) was used at $0.264 \mu\text{m}$ in a series of high power (0.5TW to 1.0TW) shots to observe the onset of perturbations for $80 \mu\text{m}$ diameter glass Microshell® targets. This observation required simultaneous streaking of the probe and driver pulses in the target chamber. By carefully auditing the timing of the probe and the main pulse beams, we could measure the earliest time at which the target was visibly affected.

Figure 2.5-9 illustrates the approximate energy-time relationship in the target volume. In addition, four photographs of reconstructed holograms are presented (Figure 2.5-10). Photograph (a) shows the

Fig. 2.5-9. Relationship between the relative timing of the UV holographic probe pulses to the target irradiating stacked pulses (5×100 psec). (See Fig. 2.5-10.)

target before any effects of irradiation are observable. Photograph (b) is a view of a target 190 ± 40 psec prior to the peak of the first main pulse. The third photograph (c) is 100 psec later. Note that the target visibility is strongly affected by the motion of the early plasma. The presence of the stem indicates that the power on target is still very low. The time of the fourth photograph (d) is approximately 50 psec after the peak of the first 80 psec main pulse. This photograph is an interferogram resulting from two probe pulses, the first occurring ~ 2 nsec prior to the second. The first pulse records the static target before the arrival of the main laser pulse, while the second together with the first maps the extent of the plasma interferometrically at a time during the main pulse. The irradiated stem shows clearly that the peak power of the main pulse has arrived.

This time resolved series of holographic images illustrates the application of this diagnostic to the observation of pre-energy effects on glass targets.

(a)

(b)

(c)

(d)

Fig. 2.5-10. Reconstructed holograms of target shots: (a) before irradiation, (b) 190 ± 40 psec prior to main pulse, (c) at ~ 100 psec later, and (d) interferogram 50 psec after peak of main pulse.

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Laser Fusion Theory

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Transport in Steepened Profiles

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Rarefaction Shocks in Spherical Geometry

Simulation Code Development

$$Q(x_k) \sum_{i=1}^m P_k(y_i) a_{ij} = \alpha$$

$$1 + \left(\frac{2(1 + k^2 \lambda_D^2) m_i}{Z m_e} \right)^{1/2} \times$$

$$\sum_{i=1}^m \sum_{j=1}^n K(x_k, y_i) \sum_{j=1}^n (z_j) =$$

$$2 \sum_k \gamma_k^e F_o e^{-\delta} \sum_{j=1}^n \int_0^{\theta_m}$$

$$Q = \min \left\{ \sum_{j=1}^n \alpha_{kj} P_l(Z_j) \right\}$$

Laser Fusion Theory

Introduction

R.L. Berger

P.M. Campbell

Progress in laser fusion in the next decade depends on understanding plasma phenomena, as well as hydrodynamic stability problems, so that target designs will achieve predicted yields. It is clear that stimulated scattering, filamentation, hot electrons, and anomalous electron transport will be important and that *ad hoc* treatments of these processes will not be adequate to design a successful laser fusion experiment. During the past year we have addressed the issues of stimulated scattering, hot electron production, and electron transport.

Experimental data have shown clear evidence of stimulated Brillouin scattering in laser produced plasmas.¹ The linear theory of Brillouin scattering in inhomogeneous and finite plasmas is well understood.² The effects of plasma spatial turbulence³ and of moderate laser bandwidth⁴ on the thresholds of parametric instabilities have been investigated theoretically and, to some extent, experimentally using microwaves.⁵ Increased laser bandwidth has been shown to have a positive effect on absorption in KMSF experiments on glass shells.⁶ This increase in absorption has been interpreted as a reduction in the amount of stimulated backscatter. A model of stimulated backscatter, which includes laser bandwidth, ion heating, and pump depletion has been incorporated into the KMSF one-dimensional, fluid simulation code, TRHYD, and has met some success in interpreting the KMSF absorption data on long pulse experiments ($\tau_L \approx 1$ nsec).

In some experiments of interest^{5,6} the laser bandwidth was larger than allowed in the existing theories of stimulated Brillouin scattering. We have extended these treatments to the situation where the bandwidth ν is larger than the ion wave frequency ω_s , as well as the damping rate of the ion wave ν_i . We have found the growth rates for the time

averaged scattered light wave amplitude and for the corresponding ion wave amplitude. Because the growth rate for the scattered light wave, but not for the ion wave, is greatly reduced when $\nu > \omega_s$, work is continuing to calculate growth rates for the intensities.

The data from experiments on gas targets at KMSF, and data from more recent LLE experiments⁷ have shown the appearance of multiple scattering, i.e., the backscattered light is itself scattered. We have solved exactly a set of nonlinear equations that describes such multiple scattering in a finite homogeneous system for three light waves—the incident, reflected, and re-reflected waves. This solution shows that multiple scattering is a weak process unless the re-reflected wave starts from an enhanced initial level. Deliberate introduction of such a level could be a method of suppressing Brillouin scattering.

Another method for reducing stimulated Brillouin scattering and other undesirable effects resulting from high intensity illumination is to divide the energy into separate frequency components. In general, an optical system cannot simultaneously focus light waves of widely disparate frequencies on a target, however KMSF has designed optics to focus infrared (1.05 μm) and green (0.53 μm) light. In this particular case, the light waves can interact with an electron plasma wave at the infrared critical surface and scatter the green light by a Raman process.⁸ We investigated this nonlinear interaction of two large amplitude waves for normally incident light and found that the steepening of the density profile by the ponderomotive forces of the lower frequency waves limits the interaction length enough to keep the scattering small.

Another advantage of multiple frequency il-

lumination is a reduction in the hot electron temperature. A wavebreaking model predicts the hot electron temperature by calculating the maximum energy of an electron fluid element at the time wavebreaking occurs.⁹ If the individual frequencies are separated by a sufficient amount, then the maximum energy scales with total laser energy I , plasma scale length L , and number of lines N as $(I L^2/N)^{1/2}$.

There is considerable indirect evidence that transport of low velocity electrons, commonly designated as thermal electrons, is inhibited at least in the vicinity of the critical surface. This evidence comes from the necessity to use flux limits to do a successful analysis of experimental data for high intensity illumination.¹⁰ The use of a flux-limit implies that one does not understand the basic transport process. It is well known that the density scale length is small at the critical surface. We have looked at the 'classical' transport of thermal electrons whose mean free paths are long compared to the density scale length at critical. This analysis leads naturally to an intensity dependent flux-limit.

In a previous publication¹¹, we reported the results of a study of spherical rarefaction shocks. There was no evidence in that study of a density shelf on the low density side of the shock, whereas, other studies and certainly planar geometry results, show clear evidence of such a shelf. We have shown that this shelf exists but disappears as the radius at which the shock occurs becomes smaller.

In the process of simulating the experiments for long laser pulses, we found that the scattering was occurring predominantly near the critical surface. Moreover the simulations showed large inverse bremsstrahlung absorption of whatever light was not scattered. Since the scale length near the critical surface plays an important role in both of these processes, we have made use of a code (FRHYD) which can use a very fine mesh (with only a modest increase in computing time) and which includes ponderomotive forces in the hydrodynamics. This latter inclusion is as important to the analysis as the inclusion of the model of Brillouin scattering.

Additions were made to the description of the transport of thermonuclear reaction products in the simulation code TRHYD. The energy loss mechanisms now include bound electrons and electron plasma oscillations.

3.1

Stimulated Scattering

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It is widely believed that the success of laser fusion depends on the use of shorter wavelength light or longer laser pulse lengths at lower intensity. Particularly in the latter case, stimulated scattering processes may limit the degree to which one can pursue this avenue. In the following Sections we report investigations of two limits on the amount of Brillouin scattering, i.e., laser bandwidth and rescattering of the scattered light. We also report on a model of Brillouin scattering for use in fluid simulations of laser plasma interactions. It is thought that, although green light may provide better absorption than infrared light, Rayleigh-Taylor instabilities may more easily develop due to the closer juxtaposition of the critical surface and the ablation surface. Thus, there is some interest in simultaneous illumination of a target with both long and short wavelength light to smooth asymmetries. We report on the results of simulations of the nonlinear interaction of two imposed laser waves on the plasma.

3.1.1 Stimulated Brillouin Rescattering

Stimulated Brillouin backscattering can severely limit the deposition of light energy in a laser produced plasma.^{1,2} In this process, light is scattered back toward its source with the momentum difference carried by an ion acoustic plasmon. The process is particularly important in long pulse experiments in which the volume of the underdense plasma is large (many vacuum wavelengths) and the thresholds are correspondingly low. Under such conditions, however, it is possible for the scattered light to be scattered again so that it propagates back into the plasma. The net result of multiple scattering is a reduction in light flux back toward the source compared with the single scattering case. In this Section we investigate the special case in which the incident light is allowed to scatter twice. Specifically, we calculate the reflection coefficient

for laser light undergoing Brillouin backscattering from a plasma slab in addition to the slab thickness required to produce 10 percent transmission. The results are compared with those corresponding to the single scattering case, thereby establishing bounds for higher order scattering effects.

Consider light normally incident upon a plasma slab occupying the region $0 \leq x \leq L$. The steady state equations governing the spatial dependence of the incident, backscattered, and rescattered light intensities, denoted by the subscripts 0, 1, and 2 respectively are², in the strong ion damping limit,

$$\frac{dI_0}{dx} = -I_0 I_1 \quad (1)$$

$$\frac{dI_1}{dx} = -(I_0 - I_2)I_1 \quad (2)$$

$$\frac{dI_2}{dx} = I_1 I_2 \quad (3)$$

in which the intensities are normalized to the incident intensity $\hat{I}_0(0)$ and lengths are normalized to the growth length $L_0 = [\beta \hat{I}_0(0)]^{-1}$. The nonlinear scattering coefficient³, β , which couples the light waves via the ion acoustic waves is taken, for simplicity, to be the same for both scatters. Higher order scattering may be required when $I_2 \sim I_1$. The limit of weak ion damping has previously been considered by Karttunen and Salomaa.⁴

Equations (1) - (3) are to be integrated subject to the boundary conditions $I_0(0) = 1$, $I_2(0) = b_2$, and $I_1(L) = b_1$ where b_1 and b_2 represent background noise levels of I_1 and I_2 . The integration is facilitated by observing that Equations (1), (2) and (3) can be combined to give

$$I_0 + I_2 - I_1 = J \quad (4)$$

a constant, while Equations (1) and (3) yield

$$I_0 I_2 = b_2 \quad (5)$$

Equations (4) and (5) are employed in Equation (1) which is integrated to produce the expression for $I_0(L)$:

$$\frac{I_0(L) - I_+}{I_0(L) - I_-} = \frac{1 - I_+}{1 - I_-} \exp [-(I_+ - I_-)L] \quad (6)$$

where

$$I_{\pm} = (J/2) \pm [(J/2)^2 - b_2]^{1/2} \quad (7)$$

Equation (6) can be used to calculate the reflection coefficient $R = I_1(0) - b_1$ (note the b_1 term) or the transmission coefficient $T = 1 - R$ by employing the relation

$$2 I_0(L) = T + b_2 + [(T + b_2)^2 - 4 b_2]^{1/2} \quad (8)$$

This equation is obtained by combining Equations (4) and (5) and noting that $J = T + b_2 - b_1$. The positive square root is chosen in Equation (8) in order to obtain the correct result in the limit $b_2 \rightarrow 0$.

Figure 3.1-1 shows the reflection coefficient as a function of the slab thickness obtained from a numerical solution of Equation (6). The reflection coefficient is plotted for several values of the noise levels b_1 and b_2 with the $b_2 = 0$ curves corresponding to the single scattering case. The reflection coefficient for $b_2 = b_1 = 10^{-3}$ does not differ significantly from the corresponding $b_2 = 0$ result over the entire range of L values given and differs only slightly as the reflection coefficient approaches its maximum value. At higher noise levels, $b_2 = b_1 = 10^{-2}$, the maximum value of the reflection coefficient is reduced significantly compared with the single scattering case while the plasma length required to produce a given reflection level is increased. Rescattering is not an effective process for small noise levels because the secondary wave I_2 must grow from a small amplitude as it propagates in the direction of decreasing I_1 . There is no region where both I_2 and I_1 are simultaneously near their maximum values which, however, is the case with I_0 and I_1 .

An alternate form of the solution is shown in Figure 3.1-2 where the plasma length required to produce a given reflection coefficient is plotted as a function of the noise level b_1 . In the example given, the reflection coefficient is set to 90 percent ($T = 10$ percent) and b_2 is chosen to be a multiple of b_1 for each of the curves shown. Again, $b_2 = 0$ represents the single scattering case. The minima of the $b_2 \neq 0$ curves indicate approximately the noise level at which multiple scattering begins to be important. The curves turn up because a longer plasma is required to compensate for the reduced reflection due

to rescattering.

The largest value of $b_1 = fb_2$ for which a solution to the double scatter problem exists, with fixed T , is

$$b_1 = f(1-f)^{-2} \{1 - [1 - T(1-f)]^{1/2}\}^2 \quad (9)$$

which occurs when $I_* = I$ [see Equation (6)] and higher order scattering is required above this value. The limiting noise level is indicated by the dashed lines in Figure 3.1-2. Alternatively, Equation (9) can be employed in Equation (6) to determine the critical plasma length beyond which higher order scattering must be considered. The result is

$$L_c = 4 T_m^{-3/2} (1 + T_m^{1/2} / 2)^{-1} - [1 + (T_m / 2)]^{-1} \quad (10)$$

in the important case $f=1$. The quantity $T_m = 2b^{1/2}$ is the minimum transmission coefficient allowed in the double scattering problem. Note that the corresponding minimum transmission coefficient in the single scattering case scales only as b and we expect the $L \rightarrow \infty$ value of the transmission to lie between these values when scattering to all orders is included.

Finally we mention situations in which $b_2 > b_1$. This could be produced by introducing a small amount of radiation downshifted by twice the acoustic frequency into the incident beam. Figure 3.1-1 compares reflection curves for $b_2 = b_1 = 10^{-2}$ and $b_2 = 4b_1 = 4 \times 10^{-2}$ with the result that the limiting reflection coefficient can be reduced from about 80 percent to as low as about 62 percent. On

Fig. 3.1-1. Reflection coefficient vs plasma length for various noise levels, b_1 , b_2 .

the other hand Figure 3.1-2 shows that the plasma length for 90 percent reflection can vary considerably for different ratios of b_2/b_1 . Notably, higher order scattering must be taken into account in Figure 3.1-2 for $b_2 > b_1$ even before b_1 reaches 10^{-2} .

Fig. 3.1-2. Length of plasma required to produce 90% reflection ($T = 10\%$) as a function of noise level for various ratios, b_2/b_1 .

3.1.2 Bandwidth Control of Stimulated Brillouin Backscatter

Stimulated Brillouin Backscatter (SBS) is thought to be a serious problem in high intensity, long pulse laser fusion experiments.⁵ SBS occurs when laser light interacts with ion acoustic fluctuations in a plasma. The result can be a significant fraction of the incident radiation reflected back toward its source. The reflection of laser light by SBS requires phase matching (frequency and wavelength) between the incident and reflected light waves and the ion wave. For this reason, the suggestion has been made to introduce bandwidth into the incident wave so that the resonance of the three waves is spoiled. The fraction of light reflected by SBS is thereby reduced. This effect has apparently been demonstrated in microwave simulations of the laser plasma interaction.⁶ Also, recent long pulse (approximately 1 ns) experiments at KMSF appear to show increased absorption of laser light by a Microshell® target when a plasma spatial filter is employed.⁷

Previous analyses^{8,9} of bandwidth control of SBS have considered the case in which the laser (pump) bandwidth ν , is small compared with the acoustic frequency ω_a . It was shown that the growth rate associated with SBS is essentially reduced by the factor γ_0/ν when ν exceeds γ_0 , the growth rate corresponding to the monochromatic case. (The ratio γ_0/ν compares the resonance width to the bandwidth.) Likewise, there is an increase in the threshold intensity of the pump required to produce SBS when wave damping is taken into consideration. In many situations of practical concern it is possible for γ_0 to be comparable to ω_a so that the inequalities $\nu < \omega_a$ and $\nu > \gamma_0$ cannot be satisfied simultaneously. Furthermore, it may be possible to produce $\nu > \omega_a$ so that it becomes necessary to extend the bandwidth calculations to include the case $\nu > \omega_a$.

In order to generalize the theory of the SBS instability driven by a finite bandwidth pump which includes the case $\nu > \omega_a$, equations are derived for the evolution of the ion density fluctuation and backscattered light amplitudes which retain the full time dependence of the acoustic wave. Consider a one-dimensional, homogeneous plasma containing cold ions ($T_i \ll T_e$). The coupled wave equations for light waves (vector potential A) and ion density fluctuations (δN) are given by²

$$\left(\frac{\partial^2}{\partial t^2} - c^2 \frac{\partial^2}{\partial x^2} + \omega_p^2 \right) A = -\omega_p^2 \frac{\delta N}{N_0} A \quad (1)$$

$$\left(\frac{\partial^2}{\partial t^2} - c_s^2 \frac{\partial^2}{\partial x^2} \right) \delta N = \frac{e^2 n_0}{2mMc^2} \frac{\partial^2}{\partial x^2} A^2 \quad (2)$$

in which $\omega_p = (4\pi e^2 n_0/m)^{1/2}$ is the electron plasma frequency, n_0 and N_0 are the (unperturbed) electron and ion densities respectively, $c_s = (ZT_e/M)^{1/2}$ is the sound speed and T_e is the electron temperature; also, c is the vacuum speed of light, e the electron charge, m the electron mass and M is the ion mass. The vector potential is now expressed as $A = A_0 + A_1$ with

$$A_0 = \hat{A}_0 \exp [ik_0 x - i\omega_0 t + i\theta(t)] + c. c. \quad (3)$$

and

$$A_1 = \hat{A}_1(t) \exp(-ik_1 x - i\omega_1 t) + c. c. \quad (4)$$

where A_0 and A_1 correspond to the incident light

wave (pump) and reflected light wave respectively. Pump depletion may be taken into account by letting $A_0 = A_0(t)$. The quantity ω_0 is the (mean) frequency of the pump wave with k_0 the corresponding wave number while ω_1 and k_1 are the frequency and wave number of the backscattered light. Pump bandwidth² may be modeled by letting $\theta(t)$ be a random function. The function $\theta(t)$ will be taken to be a Kubo-Anderson process which is piecewise constant, i.e.,

$$\begin{aligned} \theta(t) &= \theta_1 & 0 \leq t < t_1 \\ \theta(t) &= \theta_2 & t_1 \leq t < t_2, \text{ etc.} \end{aligned} \quad (5)$$

wherein the phases, θ_j , are uncorrelated and randomly distributed over a 2π interval. The jump times, t_j , form a Poisson process with mean rate ν . Note that the ensemble average $\langle \exp(i\theta) \rangle$ vanishes and the correlation function

$\Phi = \langle \exp[i\theta(t+t)] \exp[i\theta(t)] \rangle$ has the value $\exp(-\nu|t|)$. The ion density fluctuations are expressed as

$$\delta N = \hat{N}(t) \exp(ikx) + c. c. \quad (6)$$

with the complete time dependence contained in N and k is the wave number of the ion acoustic wave.

Now suppose A_1 varies slowly compared with $\exp(i\omega_1 t)$ and that the phase matching conditions

$$k_1 = k - k_0 \quad (7)$$

$$kc_s \equiv \omega_s \equiv \omega_0 - \omega_1 \quad (8)$$

are satisfied with $\omega_1 \equiv \sqrt{\omega_p^2 + k_1^2 c^2}$. The coupled equations for the backscattered light and ion fluctuations can now be expressed as

$$\omega_p^2 \frac{\hat{N}^*(t)}{N_0} \hat{A}_0 \exp[-i\omega_s t + i\theta(t)] = 2i\omega_1 \frac{\partial \hat{A}_1}{\partial t} \quad (9)$$

and

$$\begin{aligned} \left(\frac{\partial^2}{\partial t^2} + k^2 c_s^2 \right) \hat{N}^*(t) = \\ - \frac{k^2 e^2 n_0}{mMc^2} \hat{A}_0^* \hat{A}_1 \exp[i\omega_s t - i\theta(t)] \end{aligned} \quad (10)$$

A single integro-differential equation for the light wave can be obtained by combining Equations (9) and (10) with the result

$$\frac{\partial \hat{A}_1}{\partial t} = \gamma_0^2 \int_0^t d\tau \left\{ 1 - \exp[-2i\omega_s(t-\tau)] \right\} \times \exp[i\theta(t) - i\theta(\tau)] \hat{A}_1(\tau) \quad (11)$$

in which γ_0 is the zero bandwidth growth rate. Similarly, a single equation for $\hat{N}^*$ is given by

$$\hat{N}^* = \gamma_0^2 \int_0^t d\tau \quad (12)$$

$$\left\{ \exp[i\omega_s(t-\tau)] - \exp[-i\omega_s(t-\tau)] \right\}$$

$$\int_0^\tau d\tau' \exp[i\omega_s(t-\tau') + i\theta(\tau') - i\theta(\tau)] \hat{N}^*(\tau')$$

Equations (11) and (12) differ from the corresponding narrowband, $\nu < \omega_s$, equations by the appearance of the second term within each pair of braces.

The stochastic equations [Equations (11) and (12)] can be solved for the ensemble averaged quantities $\langle A_1 \rangle$ and $\langle \hat{N}^* \rangle$ by the following considerations. Equation (11), e.g., has the form of

$$I = \int_0^t d\tau K(t-\tau) \exp\{i[\theta(t) - \theta(\tau)]\} A(\tau)$$

and in the time interval $(0, t)$, θ can undergo any number of jumps. The contribution from n such jumps is

$$I_n = \exp[i(\theta_n - \theta_0)] \int_0^{t_1} d\tau KA$$

$$+ \exp[i(\theta_n - \theta_1)] \int_{t_1}^{t_2} d\tau KA$$

$$+ \dots + \exp[i(\theta_n - \theta_{n-1})] \int_{t_{n-1}}^{t_n} d\tau KA + \int_{t_n}^t d\tau KA$$

so that upon phase averaging only the last integral makes a nonzero contribution. Thus, the ensemble average of I becomes, upon summation over n ,

$$\langle I \rangle = (\text{Probability of no jumps}) \times \int_0^t d\tau KA + \sum_{t'} (\text{Probability that last jump occurred at } t') \times \int_{t'}^\tau d\tau KA.$$

The probability of m jumps in a time interval Δt is given by

$$P_m(\Delta t) = \frac{(\nu \Delta t)^m}{m!} \exp(-\nu \Delta t)$$

for a Poisson process so that

$$\langle I \rangle = \exp(-\nu t) \int_0^t d\tau K \langle A \rangle + \int_0^t \nu dt' \exp[-\nu(t-t')] \int_{t'}^t d\tau K \langle A \rangle$$

An integration by parts finally gives

$$\langle I \rangle = \int_0^t d\tau \exp[-\nu(t-\tau)] K(t-\tau) \langle A(\tau) \rangle$$

which corresponds to the so-called Bourret closure approximation.

Upon ensemble averaging, Equation (11) can therefore be approximated by

$$\frac{\partial \langle \hat{A}_1 \rangle}{\partial t} = \gamma_0^2 \int_0^t d\tau \left\{ 1 - \exp[-2i\omega_s(t-\tau)] \right\} \exp[-\nu(t-\tau)] \langle \hat{A}_1(\tau) \rangle \quad (13)$$

while Equation (12) may be ensemble averaged to produce

$$\langle \hat{N}^* \rangle = \gamma_0^2 \int_0^t d\tau$$

$$\left\{ \exp[i\omega_s(t-\tau)] - \exp[-i\omega_s(t-\tau)] \right\}$$

$$\int_0^\tau d\tau' \exp[i\omega_s(t-\tau) - \nu(t-\tau)] \langle \hat{N}^*(\tau') \rangle \quad (14)$$

Equations (13) and (14) may be solved by standard techniques. However, it is perhaps more useful to consider limiting cases. First, when $\gamma_0 < \omega_s$ the "narrowband" results can be recovered by taking $\nu < \omega_s$ and observing that, e.g., in Equation (13), the oscillating exponential contributes little to the integral. Then, an integration by parts results in the leading order expressions for the growth rate γ

$$\gamma = \gamma_0 \quad \nu < \gamma_0$$

$$\gamma = \gamma_0 \left(\frac{\gamma_0}{\nu} \right) \quad \nu > \gamma_0 \quad (15)$$

as expected. Identical results are obtained by employing the same approximations in Equation (14).

The broadband limit ($\nu > \omega_s$), however, produces an unexpected result. When ν is large, only values of τ within a correlation time, ν^{-1} , of t contribute to the integral in Equation (13) and a similar statement holds for the inner integral of Equation (14). Thus, near the upper limit of integration in Equation (13), the $\exp[2i\omega_s(t-\tau)]$ term contributes and partially

cancels the first term in braces. In the broadband limit, the right side of Equation (13) can be integrated by parts to obtain the growth rate

$$\gamma = \gamma_s \left(\frac{\gamma_0}{\nu} \right) \left(\frac{2\omega_s}{\nu} \right)^2 \quad (16)$$

for the backscattered light wave. Equation (14), in contrast, reproduces the growth rate given in Equation (15-b) for the ion density fluctuations in the broadband limit.

The quantities of greater interest are the ensemble averaged quantities $\langle |\hat{A}|^2 \rangle$ and $\langle |\hat{N}|^2 \rangle$. The determination of the growth rates corresponding to the intensities is considerably more difficult than for the amplitudes. By employing and extending the methods described in the preceding paragraphs, the growth rates given therein can be shown to also hold for the corresponding intensities. However, it is necessary to verify those results by more rigorous means. The Green's function method of Brissaud and Frisch¹⁰ which has been employed in the narrowband problem by Thomson⁸ and Laval⁹ et al. and by us to verify the preceding results when used with Equations (8) and (9), appears to be the most promising approach although the extension to the broadband problem complicates the calculation substantially.

3.1.3 Fluid Simulation of Brillouin Scatter

A model of Brillouin scattering has been developed and incorporated into a fluid simulation code in order to help analyze long pulse experiments at KMSF. In these experiments, the absorption fraction for 1.05 μm light was obtained as a function of the incident intensity for narrow bandwidth (1-4 $\text{\AA}$) light. A significant difference in the absorption percentage was noted as a function of the bandwidth. The difference can be explained qualitatively by the effects of bandwidth on the stimulated scattering process.⁸

The nonlinear interaction of the incident laser light with the backscattered wave can be represented by the coupled nonlinear equations,²

$$\left(\frac{d}{dx} + \gamma_a(x) \right) I_0 = \alpha(x) I_0 I_1 \quad (1)$$

$$\left(\frac{d}{dx} - \gamma_a(x) \right) I_1 = \alpha(x) I_0 I_1 \quad (2)$$

where the laser wave is incident from the left, γ_a is the total of all linear damping processes (such as inverse bremsstrahlung and ion acoustic turbulence scattering), $\alpha(x)$ is the coupling coefficient valid in the induced scattering limit, and $I_0(I_1)$ is the intensity of the incident (reflected) wave normalized to the incident intensity. In deriving Equations (1) and (2), it has been assumed that

$\nu + \nu_i \gg \text{Max}(\gamma_0, C_s I_1^{-1} |\partial/\partial x I_1|)$ where ν is the laser bandwidth, ν_i is the damping on the driven ion density perturbation, γ_0 is the growth rate for stimulated scattering when there is no damping, and $c_s \equiv (ZT_e/m_i)^{1/2}$ is the ion acoustic speed. The WKB assumption has also been used to obtain Equations (1) and (2), i.e., density gradients have been assumed small compared to the laser wavelength. Conditions for which this assumption is marginally valid will be discussed.

For small bandwidth ($\nu \ll \nu_i$), the coupling coefficient $\alpha(x)$ is given by

$$\alpha(x) = \left(\frac{x_0}{x} \right)^{s-1} \left(\frac{eE_0}{m\omega_0 v_e} \right)^2 \frac{n(r)}{\sqrt{1-n(r)}} \frac{\omega_0}{c} \beta(x) \quad (3)$$

where

$$\beta(x) = \text{Im} \left(\frac{ZT_e/T_i W(\xi_i) W(\xi_e)}{W(\xi_e) + ZT_e/T_i W(\xi_i)} \right), \quad (4)$$

and where $n(r)$ is the density in units of critical density, $v_{e,i} = \sqrt{T_{e,i}/m_{e,i}}$ is the electron (ion) thermal velocity, ω_0 is the incident light frequency, E_0 is the electric field amplitude at x_0 (a point where the interaction is negligible), and $s = 3$ for spherical geometry and $s = 1$ for planar geometry. The function $W(\xi) = 1 + \xi Z(\xi)$ where $Z(\xi)$ is the plasma dispersion function. The arguments of the dispersion function are

$$\xi_{e,i} = \frac{\omega_0 - \omega_1 + 2k_0 u(x)}{2\sqrt{2} k_0(x) v_{e,i}} \quad (5)$$

where $u(x)$ is the plasma flow velocity and ω_1 is the frequency of the backscattered light.

Provided that $\alpha(x)$ is slowly varying and that the damping terms are negligible, Equations (1) and (2) can be integrated in the WKB spirit to obtain the reflection coefficient, R , from the equation,

$$B = \frac{(1-R)R}{\exp[(1-R)z] - R} \quad (6)$$

in terms of the initial noise level in the scattered radiation relative to the incident wave, $B = I_1(x)/I_0(x_0)$, and the amplification factor

$$z = \int_{x_0}^{x_1} dx \alpha(x) \quad (7)$$

For $ZT_e/T_i \gg 1$, $\beta(x)$ is the ratio of the ion acoustic frequency, ω_s , to the ion acoustic damping rate ν_i . The quantity $\beta(x)$ is a function of the variables ξ_i and ZT_e/T_i . For fixed ZT_e/T_i , $\beta(x)$ is nearly a gaussian function of ξ_i . The maximum value of β increases with ZT_e/T_i but the full width at half maximum Δ , decreases. The reduction of the scattering rate by bandwidth has been shown to be equivalent in many cases of interest to increasing the damping rate to $\nu + \nu_i$. Thus, the value of β decreases in our model according to $\beta \rightarrow \beta/(1 + \nu/\Delta\omega_s)$. The change of plasma parameters also changes the growth rate especially through the dependence of ξ on the flow velocity (see Equation 5). This is most marked when $ZT_e/T_i \gg 1$ and $\nu/\Delta \rightarrow 0$. This dependence is taken into account by assuming a gaussian form for $\beta(x)$, i.e.,

$$\beta(x) = \frac{\beta_m}{1 + \frac{\nu}{|\omega_0 - \omega_1| \Delta} + \frac{(\xi_i - \xi_{im})^2}{\Delta^2}} \quad (8)$$

where β_m and Δ are functions of ZT_e/T_i , ξ_{im} is the value of ξ_i for which β achieves its maximum value, and ξ_i is given by Equation (5).

To calculate the reflection from Equation (6), one finds the value of ω_1 , that maximizes $\alpha(x)$ at $x = x_1$, computes the integral in Equation (7), and finds R . Knowing R , one knows the solution for $I_1(x)$ and $I_0(x)$ everywhere. The calculation is self-consistent if the assumption of negligible damping is correct for $x < x_1$. If this assumption is violated a new x_1 is chosen until self consistency is found. Because the plasma is inhomogeneous, many different waves can contribute to the total backscatter, i.e., different frequencies are scattered preferentially in different regions. In most cases of interest, one wave is dominant. This fact is evident from a calculation of z for typical plasma profiles. When this is not the case, each scattered wave has been treated as interacting independently with the laser wave according to the following prescription. Start at $x_1' < x_0$, reflect the energy according to Equa-

tion (6), and then calculate the scattering from $x_1'' < x_1'$ but with an incident intensity reduced by the previous reflection.

After each scattering calculation, the amount of energy given to the acoustic wave is calculated and deposited as ion internal energy. This heating lowers the ratio ZT_e/T_i and thus lowers the scattering rate. However, simulations have shown that cooling of the ions by hydrodynamic expansion is sufficient to keep $ZT_e/T_i > 10$ and thus prevents any substantial reduction in the growth rate. Although the ions are heated, the treatment is incorrect in that the ion-ion collision rate is probably not rapid enough to maintain a Maxwellian velocity distribution. Since the high velocity ions contribute to the damping, they may be underestimated in this model. One could define an effective "hot ion" temperature e.g., from a wavebreaking model of the ion density perturbations, for the purposes of the Brillouin scattering calculation, but without a transport treatment of these ions the utility of such an approach is suspect.

The energy conservation of the ions is not complete in that the supersonic flow of the corona requires a blue shift of the backscattered light. However, that blue shift requires a loss of momentum by the ions. This momentum change is in addition to, but similar to the zero frequency ponderomotive force exerted on the plasma by the light waves. Since the plasma density can be only a fraction of critical density, substantial changes in the flow velocity may occur locally due to these forces, and these may in turn help explain the red shifts observed experimentally. Moreover, since the dominant inhomogeneity in the growth rate $\alpha(x)$, arises from the variation of the flow velocity $u(x)$, local modification of the flow rate will also reduce the backscatter. Inclusion of these effects requires very fine resolution which is outside the scope of most fluid simulations.¹¹

3.1.4 Raman Scattering by Dual Frequency Illumination

One of the disadvantages of driving a target implosion with a laser is that all the energy is concentrated in a single wavelength, therefore the absorption must occur in a narrow region in space. As a

result collective absorption processes dominate at high intensities, and these processes have as a characteristic that few electrons carry off most of the energy. The unfavorable influences on heat transport and implosion efficiency of this type of energy deposition are well known. Other phenomena also scale unfavorably with high laser intensity such as Brillouin scattering, Raman scattering, and filamentation.

One obvious means of putting the same power on target but reducing the intensity available to drive these processes is to spread the same energy over many frequencies. Most optical systems cannot focus all of these frequencies at once but at KMSF a technique has been developed using all reflective optics to focus 0.53 μm and 1.05 μm light simultaneously. The only obvious difficulty is that the infrared (1.05 μm) light may scatter the green light (0.53 μm) at or near the 1.05 μm critical surface by Raman scattering. In this process, the ponderomotive force due to the interaction of the two laser waves drives an electrostatic wave to large amplitude. A portion of the green light energy is transferred to the infrared light and a roughly equal portion to the electrostatic wave.¹² This later portion is lost by the wave to high energy electrons either by Landau damping or wavebreaking.

It is important to note that, because both transverse waves are large in amplitude, the equations cannot be linearized. Moreover, because the interaction occurs at the critical surface the electrostatic wave and the infrared wave are at their turning points so that a WKB analysis is invalid. The Equations under consideration are

$$\left(\frac{d^2}{d\xi^2} + 4 - n\right)\hat{E}_1 = -\frac{2v_e}{c}\hat{E}_2\frac{d}{d\xi}\hat{E}_3 \quad (1)$$

$$\left(\frac{d^2}{d\xi^2} + 1 - n\right)\hat{E}_2 = -\frac{v_e}{2c}\hat{E}_1\frac{d}{d\xi}\hat{E}_3^* \quad (2)$$

$$\begin{aligned} \left(\frac{3v_e^2}{c^2}\frac{d^2}{d\xi^2} + 1 - n + iv_L\right)\hat{E}_3 \\ = \frac{v_e}{4c}n\frac{d}{d\xi}(\hat{E}_1\hat{E}_2^*) \end{aligned} \quad (3)$$

where $\xi = \omega x/c$ and the electric field amplitudes have been normalized to the plasma pressure, i.e., $\hat{E} = eE/M\omega v_e$. Normal incidence has been assumed but the equations can be easily generalized to obli-

que incidence. For obliquely incident light, the electrostatic wave is not only driven nonlinearly by the ponderomotive force but is also driven linearly, and thus the interaction may be stronger than calculated here.

Equations (1) - (3) have been solved numerically for a fixed density profile with outgoing wave boundary conditions on the low density side of the infrared critical surface. On the high density side of critical, the low frequency wave is assumed to be evanescent and the high frequency wave to have no reflected component. The equations are differenced so as to change Equations (1) - (3) into a nonlinearly coupled system of algebraic equations. In the absence of the nonlinear terms, this system can be solved efficiently since the matrix of coupling coefficients is tridiagonal. The nonlinear set is solved by iteration using the solution to the uncoupled set as the first guess. The convergence criteria is that the percentage change in the absolute value of the field at each position i between iterations j and $j + 1$ summed over all positions i be less than K , i.e.,

$$\sum_{i=1}^N \frac{|E_i^j - E_i^{j+1}|}{|E_i^j|} < K.$$

The transmission percentage of green light and the total absorption (no inverse bremsstrahlung was allowed) is shown in Figure 3.1-3 as a function of η , the percentage of energy in the infrared for a fixed total incident energy. Optimum absorption occurs when there are roughly equal amplitudes in infrared and green because this maximizes the driving term in Equation (3) for the electrostatic wave. The transmission fraction of green continues to drop after this maximum is attained because backscatter of green occurs due to the interaction of the infrared laser and the electrostatic wave and this driving term reaches a maximum when almost all the energy is initially in the infrared. When η approaches one, the transmission rises again and approaches one as $\eta \rightarrow 1$.

The absorption fraction is a linear function of the length of the plasma as shown in Figure 3.1-4. Clearly then, the ponderomotive forces of the electrostatic wave and the infrared laser will act to reduce the interaction by steepening the profile. Since the scale length varies inversely with the incident infrared intensity, (see Figure 3.1-5) the

Fig. 3.1-3. The total absorption fraction of light by the Raman Scattering Process as a function of the fraction of energy in the low frequency (infrared) light. The scale length and total incident energy are fixed to the values given by $\alpha = 10$ and $|\hat{E}_{tot}|^2 = 1$. The transmission of the high frequency light is indicated by the dashed curve.

Fig. 3.1-4. The absorption percentage of all incident energy and the transmission of high frequency light as a function of scale length.

assumption of a fixed scale length gives misleading results. In Figure 3.1-5, the absorption is shown for a scale length that corresponds to a profile at the infrared critical surface that is consistent with a normally incident light wave with energy corresponding to the infrared intensity. Clearly, there is a dramatic effect in that the absorption peaks now when almost all the energy is in the green and more significantly the absorption is negligibly small.

The results above were all done for fixed profiles. We have run a fluid simulation with isothermal Lagrangian hydrodynamics to watch the time evolution for a case where the initial profile was consistent with the infrared light intensity and the initial absorption was 11 percent. At four picoseconds the electrostatic wave has modified the profile noticeably as shown in Figure 3.1-6 and the absorption has dropped to 7 percent. At no point was the absorption greater than it was initially.

The plasma temperature plays a role in coupling in Equations (1)-(3) because of the normalization of the fields. The scale length is determined by the hydrodynamics which does not involve the tem-

Fig. 3.1-5. The absorption percentage as a function of the fraction of energy in the low frequency wave for fixed total energy and for a density profile consistent with the energy in the low frequency wave. The density profile is found for an isothermal plasma in which the flow on the high density side of critical is subsonic. The dependence of the scale length on the low frequency incident-intensity fraction is also shown.

Fig. 3.1-6. The initial density profile and electrostatic wave amplitude resulting from illumination of the plasma with a total incident energy $|\hat{E}_{tot}|^2 = 5$, 3% of which is in the low frequency wave (solid curves); dashed curves indicate the density profile and electrostatic wave amplitude after 4 psec.

perature explicitly (for the same normalization of the variables). Thus, increasing the temperature will increase the absorption and backscatter of green light if the normalized intensities are fixed. A complete answer to the role of this process in determining absorption when dual frequency illumination is used will require a self-consistent determination of the temperature, scale length and angle of incidence.

3.2

Transport in Steepened Profiles

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In the vicinity of the critical surface, the density profile scale length can be as small as ten or twenty Debye lengths due to ponderomotive force or to the momentum reaction of resonantly heated electrons.¹ The plasma adjacent to the steepened profile often has gentle density and temperature gradients. If one considers an electron traversing this region, it is clear that the scattering mean-free-path may be

short compared to the scale length for low energy electrons everywhere but in the very steep gradient region. In the steep gradient region however, the electron moves in a trajectory determined by the self-consistent fields and rarely undergoes a collision with an ion.

The procedure to derive a transport equation averaged over the fast time scale of a laser oscillation is outlined as follows. The analysis is simplified by the assumptions of normal incidence and slab geometry. The direction of propagation and the density inhomogeneity lie along the x-axis. It is useful to transform to the frame of the oscillating electron and then expand the distribution function in the perturbation series with parameter $v_e/\omega L$ where L is the minimum scale length in the plasma. Provided one can neglect electron-electron collisions ($Z \gg 1$) and one can assume that heat source terms are negligible in the steepened profile region, the Boltzmann equation describing the electron transport can be written as

$$\pm v_x \frac{\partial}{\partial x} f = v_x \nu_{ei} \frac{1}{v_{\perp}} \frac{\partial}{\partial v_{\perp}} (v_x v_{\perp} \frac{\partial}{\partial v_{\perp}} f) \quad (1)$$

where $\nu_{ei} = 4\pi n Z e^4 / m_e^2 (v_e^2 + v^2)^{3/2}$ is the electron-ion collision frequency including the average jitter velocity of an electron in the applied laser field $v_e = e|E_0|/m_e \omega$ and $v_x = \sqrt{2} (W - \Phi(x) - v_{\perp}^2/2)^{1/2}$. Equation (1) describes the evolution of the time averaged distribution function in the oscillating frame. The distribution is a function of the variables

$$w = \frac{v^2}{2} + \Phi(x), \quad v_{\perp} = (v_y^2 + v_z^2)^{1/2}$$

and x where the potential $\Phi(x) = -e\phi(x)/m_e + |eE_0(x)/m\omega|^2$ is a sum of the ponderomotive and ambipolar potentials. The shape of $\phi(x)$ can be determined if the density profile is known. For the present, the density will be taken to be a monotonically decreasing function of x . From the ion momentum and continuity equations, one can easily conclude that $\Phi(x)$ must have the shape shown in Figure 3.2-1 where the steeply rising portion corresponds to the assumed steeply dropping density.

In the middle of the steeply rising potential, $\Phi(x)$, the approximation $L \ll \lambda_{ei}$ is a good one except for the very slowest electrons (although, on the low density side of this region, the contribution to the

Fig. 3.2-1. The potential $\hat{\Phi}(x) = \Phi(x)/\Phi_0$ is plotted for the model given in Equation 32. The form of the density profile that is consistent with this model potential (Eq. 32) is shown for $0 \leq s \leq 20$, where $s = (x - x_0)/L_2$ and $\xi(x) = n(x)/n_1(x_0)$. The other parameters are given in the text. The dashed horizontal line indicates the density which would be attained asymptotically in x if $\Phi(x)$ were equal to Φ_0 for $x \geq x_1$.

collision frequency of quiver velocity improves this assumption for these electrons). Asymptotically far from this region, the opposite approximation is valid, i.e., $L \gg \lambda_{ei}$. In either of these two limits, one can easily obtain solutions to Equation (1). The problem remains to join these solutions across the region where both approximations fail, i.e., where $L \sim \lambda_{ei}$. Since the mean free path is a function of energy w , as well as space, the equality $L = \lambda_{ei}$ will occur at different spatial positions for different energies. The emphasis here is on plasma where, for a given energy, there is at least one spatial position for which $L = \lambda_{ei}$. It should be pointed out however that, in general, there are some low energy electrons for which $L > \lambda_{ei}$ always applies and some high energy electrons for which $L < \lambda_{ei}$ always applies. In what follows, we assume these electrons make no substantial contribution to the results.

Let the superscript 1(3) denote the high (low) density regions I(III) where the approximation $L \geq \lambda_{ei}$ holds. Then, to first order in λ_{ei}/L , the solution to Equation (1) is

$$f_D^{\pm, \sigma}(w, v_{\perp}, x) = \left(1 \mp \frac{v_x}{2v_{ei}} \frac{\partial}{\partial x}\right) f_{oD}^{\sigma}(w, x) \quad (2)$$

where $\sigma = 1, 3$ and $\pm$ denotes the sign of the velocity along the x direction. Let the subscript F denote the region II where $L < \lambda_{ei}$ holds. Then, to first order in L/λ_{ei} , the solution to Equation (1) is

$$f_F^{\pm}(w, v_{\perp}, x) = f_{oF}^{\pm}(w, v_{\perp}) \pm \int_{x_{\pm}^*}^x dx' L_1(x') f_{oF}^{\pm}(w, v_{\perp}), \quad (3)$$

where

$$L_1 f = \frac{v_{ei}}{v_{\perp}} \frac{\partial}{\partial v_{\perp}} (v_x v_{\perp} \frac{\partial}{\partial v_{\perp}} f) \quad (4)$$

The point $x_{\pm}^*$ will be specified in the following discussion.

In order to determine the distribution everywhere in space, one needs to join solutions (2) and (3) across a boundary layer where $L \sim \lambda_{ei}$. To simplify the analysis even further, let the potential have the special form shown in Figure 3.2-1. This form allows us to specify x^* easily. If

$$w > \Phi_0 \text{ and } v_{\perp}^2 < 2(w - \Phi_0),$$

an electron incident on region II from the left will reflect from the potential wall and return to region I. Thus the position $x_{\pm}^*$ is defined as $x_+^* = x_0$ and $x_-^* = x_1$, where x_1 is defined by

$$w - \Phi(x_1) - v_{\perp}^2/2 = 0.$$

If $w > \Phi_0$ and $v_{\perp}^2 < 2(w - \Phi_0)$, then an electron is not reflected while traveling across region II. Thus position $x_{\pm}^*$ is defined as $x_+^* = x_0$ and $x_-^* = x_1$.

It is clear from the preceding discussion that the distributions in region II are determined to first order by the boundary conditions that $f^*(w, v_{\perp}, x)$ be continuous at x_0 and that $f^*(w, v_{\perp}, x)$ be continuous at x_1 , i.e.,

$$\begin{aligned} f_F^+(x) &= f_D^{1,+}(w, v_{\perp}, x_0) + \\ &\int_{x_0}^x L_1(x') f_{oD}^{1,+}(w, x_0) dx' \\ &= f_D^{1,+}(w, v_{\perp}, x_0) \end{aligned} \quad (5)$$

and

$$\begin{aligned} f_F^-(x) &= f_D^{3,-}(w, v_\perp, x_1) ; \\ w &> \Phi_0, \quad v_\perp^2 < 2(w - \Phi_0) \\ &= f_D^{1,+}(w, v_\perp, x_0) ; \quad w < \Phi_0, \\ \text{and } w &> \Phi_0, \quad v_\perp^2 > 2(w - \Phi_0) \end{aligned} \quad (6)$$

In obtaining Equations (5) and (6), use is made of the fact that

$$\partial f_{oD}^\pm / \partial v_\perp = 0.$$

Since the distribution is completely determined in region II by using only two boundary conditions, the distribution of outward going electrons will be discontinuous at $x = x_1$ since

$$f_F^+(x_1) \neq f_D^{3,+}(w, v_\perp, x_1)$$

and $f_F^-(x_0) \neq f_D^{1,-}(w, v_\perp, x_0).$

This result is expected because the potential reflects electrons according to their angular momentum as well as their energy so that even an initially isotropic distribution will become anisotropic after passing through this region. However, the anisotropy will disappear in a distance on the order of one mean free path.

One is thus led to consider the problem where an anisotropic distribution is specified at $x_0(x_1)$ and to find its value at $x < x_0(x > x_1)$. Since the anisotropy disappears when

$$L \gg |x - x_{0,1}| \gtrsim \lambda_{ei},$$

the approximations in Equation (1) can be made that

$$\Phi(x) \approx \Phi(x_{0,1}) \text{ and } \nu_{ei} \approx \nu_{ei}(x_{0,1}).$$

With the definitions

$$u = v_\perp / \sqrt{2} (w - \Phi(x^*))^{1/2} \quad (7)$$

$$\xi = \frac{\nu_{ei}(x^*)}{\sqrt{2} (w - \Phi(x^*))^{1/2}} |x - x^*| \quad (8)$$

where $x^* = x_{0,1}$ as appropriate, Equation (1) becomes

$$\frac{\partial}{\partial \xi} f_{BL} = \frac{1}{u} \frac{\partial}{\partial u} \left(u \sqrt{1 - u^2} \frac{\partial}{\partial u} f_{BL} \right) \quad (9)$$

where the subscript denotes that this Equation applies in the boundary layer. The approximation,

$$\nu_{ei}(x) \approx \nu_{ei}(x_{0,1}),$$

is not necessarily valid because the collision frequency depends on the electric field amplitude of the laser through $v_0(x)$ and because $v_0(x)$ varies on the scale of the wavelength. The treatment here is still qualitatively correct; and, moreover, the change of variables to

$$\xi = \int_{x^*}^x dx' \nu_{ei}(x') (2w - 2\Phi(x'))^{-1/2}$$

makes the analysis valid in any case. Note the change of variable $\mu = \sqrt{1 - u^2}$ changes Equation (9) into the Equation considered by Stein and Bernstein² and earlier by Bethe, et al.⁸

The distribution function $f_{BL}^*(w, v_\perp, x)$ is given by

$$f_{BL}^+(w, v_\perp, x) = \int_0^1 d\bar{u}' h(\bar{u} | \bar{u}', \bar{\xi})$$

$$[f_F^+(x_1, \bar{u}') - f_D^{3,+}(x_1, \bar{u})] + f_D^{3,+}(x, \bar{u}) \quad (10)$$

where $\bar{u}$ and $\bar{\xi}$ are given by Equations (7) and (8) respectively when $x^* = x_1$. The distribution function is given by $f_{BL}^-(w, v_\perp, x)$

$$\begin{aligned} f_{BL}^-(w, v_\perp, x) &= \int_0^1 du' h(u | u', \xi) [f_F^-(x_0, u') - \\ &f_D^{1,-}(x_0, u')] + f_D^{1,-}(x, u) \end{aligned} \quad (11)$$

where u and ξ are given by Equations (7) and (8) respectively when $x^* = x_0$. The Green's function $h(u | u', \xi)$ is given by

$$h(u | u', \xi) = \sum_n u' U_n(u') U_n(u) e^{-\lambda_n^2 \xi} \quad (12)$$

where

$$\int_0^1 du u U_n(u) U_n'(u) = \delta_{nn'} \quad (13)$$

and $U_n(u)$ satisfies

$$\frac{1}{u} \frac{d}{du} \left(\sqrt{1 - u^2} u \frac{d}{du} U_n \right) + \lambda_n^2 U_n = 0 \quad (14)$$

It is easy to show that Equations (10) and (11) satisfy the required boundary conditions (continuity of f) at x_1 and x_0 respectively. These solutions also show that the boundary layer thickness is proportional to λ_{ei}/λ_1^2 where λ_1 is the smallest eigenvalue of Equation (14). It is interesting to note that $\lambda_1 \approx 4$ so that isotropy is obtained in a fraction of a mean free path.

Now that the distribution function is determined at least formally everywhere, one can take moments to determine the density, current, J , and heat flow, q . The distribution has been obtained in the oscillating frame but this fact can be ignored provided one sets the current $J=0$ and consistently drops terms of order $v/\omega L$.

To obtain explicit results, it is useful to specify the distribution function $f_{0\sigma}$ in regions I and III to be Maxwellians, i.e.,

$$f_{0\sigma}^\sigma(w, x) = \quad (15)$$

$$n_\sigma(x) \left(\frac{m}{2\pi T_\sigma(x)} \right)^{3/2} \exp \left[-\frac{m}{T_\sigma(x)} (w - \Phi(x)) \right]$$

where $\sigma = 1, 3$. The density, current, and heat flow are obtained by evaluating the moments,

$$n(x) = \int d^3v f, \quad (16)$$

$$J = -e \int d^3v v_x f, \quad (17)$$

$$q = \frac{m}{2} \int d^3v v^2 v_x f. \quad (18)$$

In regions I and III asymptotically far from x_0 and x_1 , the density $n(x) = n_\sigma(x)$,

$$J_\sigma = j_{D\sigma}(x) L_{n\sigma} [L_{A\sigma}^{-1} W_{3/2}(v_{0\sigma}^2) + L_{T\sigma}^{-1} W_{5/2}(v_{0\sigma}^2)], \quad (19)$$

where

$$j_{D\sigma}(x) = \frac{4ne}{3\bar{v}_{ei\sigma}} \sqrt{2/\pi} \frac{T_\sigma}{m} L_{n\sigma}^{-1}$$

and

$$W_\nu(a) = \int_0^\infty ds s^\nu (s+a)^{3/2} e^{-s} \quad (20)$$

$$\text{and } v_{0\sigma}^2 = \frac{m v_0^2}{2 T_\sigma}, \quad L_{n\sigma}^{-1} = d \ln n_\sigma / dx, \quad L_{T\sigma}^{-1} =$$

$$d \ln T_\sigma / dx \quad L_{A\sigma}^{-1} = L_{n\sigma}^{-1} - 3 L_{T\sigma}^{-1} / 2 + m d\Phi/dx T_\sigma^{-1},$$

and $\bar{v}_{ei\sigma} = 4\pi nZ e^4 / (T_\sigma)^{3/2} \sqrt{m_e}$. In the limit that

$$v_0 \rightarrow 0, \quad W_\nu(v_{0\sigma}^2) \rightarrow \Gamma(\nu + 5/2)$$

where $\Gamma(\nu)$ is the gamma function. Similarly the heat flow is given by,

$$\begin{aligned} q_\sigma(x) &= -q_{D\sigma} [L_{t\sigma} L_{A\sigma}^{-1} W_{5/2}(v_{0\sigma}^2) + W_{7/2}(v_{0\sigma}^2)] \\ &= -q_{D\sigma} [W_{7/2}(v_{0\sigma}^2) - (W_{5/2}(v_{0\sigma}^2))^2 / W_{3/2}(v_{0\sigma}^2)] \end{aligned} \quad (21)$$

where

$$q_{D\sigma} = \frac{4n}{3\bar{v}_{ei}} \sqrt{2/\pi} \frac{T_\sigma}{m} \frac{d}{dx} T_\sigma$$

and where the second form in Equation (21) follows from setting $J=0$ and using Equation (19). In the limit $v_0 \rightarrow 0$ the heat flow becomes,

$$q = -\Gamma(5) q_{D\sigma} \quad (22)$$

which is the well known classical result.

In region II, the corresponding quantities are given by

$$n(x) = n_1(x_0) \exp \left[-\frac{m}{T_1} \Phi(x) \right] \quad (23)$$

$$\left[1 - \frac{1}{2} \operatorname{erfc}(\sqrt{y_1}) \right] + \frac{1}{2} n_3(x_1) \exp(y_3) \operatorname{erfc}(\sqrt{y_3})$$

for $x_0 \leq x \leq x_1$ and $y_\sigma = m(\Phi_0 - \Phi(x))/T^\sigma$ and $\operatorname{erfc}(\sqrt{y})$ is the complement of the error function,

$$\operatorname{erfc}(\sqrt{y}) = \frac{1}{\sqrt{\pi}} \int_y^\infty ds s^{-1/2} e^{-s} \quad (24)$$

The current is given by

$$\begin{aligned} J(x) &= -\frac{e}{\sqrt{2\pi}} \left\{ n_1(x_0) \sqrt{T_1/m} \exp(-y_{01}) \right. \\ &\quad \left. - n_3(x_1) \sqrt{T_3/m} \right\} = 0 \\ &\equiv -\frac{1}{2} [j_{f1} - j_{f3}] \end{aligned} \quad (25)$$

where $y_{01} = m\Phi_0/T_1$. Terms of the order $(\lambda_{ei}/L)_{x=x^*} \ll 1$ have been neglected and will be needed to match the solutions continuously across the boundary layer. Note that the current is a constant because only transmitting electrons contribute to the current and heat flow and because the number of these electrons is inversely dependent on their radial velocity between x_0 and x_1 . The current must

be zero since there are no sources or sinks for electrons by assumption. Thus Equation (25) is a constraint equation. The heat flow $\dot{q}$ between x_0 and x_1 is given by

$$q(x) = q_{f1} - q_{f3}$$

$$q(x) = \sqrt{2/\pi} (T_1 - T_3) n_3(x_1) \left(\frac{T_3}{m}\right)^{1/2} \quad (26)$$

where the condition $J=0$ has been imposed to obtain the second expression. In Equation (26), we have defined

$$q_{f1} = \sqrt{2/m\pi} T_1^{3/2} n_1(x_0) \exp(-y_{01}),$$

and

$$q_{f3} = \sqrt{2/m\pi} T_3^{3/2} n_3(x_1).$$

The connection of the Equations (23)-(26) for the density, current, and heat flow to the Equations (19)-(22) is done by computing these quantities using the distributions given in Equations (10) and (11). These boundary layer regions are fractions of a mean free path thick, but these distances are still large compared to the distance between x_0 and x_1 . The density scale length in strongly steepened profiles is 10-20 Debye lengths whereas the mean free path is about $n\lambda_D^3$ Debye lengths. An estimate of $L(\lambda_{ei}/10)$ at $n = 10^{21} \text{ cm}^{-3}$ and $T_e \sim 1 \text{ keV}$ for a scale length of 10 Debye lengths is

$$10L/\lambda_{ei} = 10/n\lambda_D^3 \approx 0.02.$$

The heat flow will differ from the classical value (Equation 22) therefore over a region substantially larger than the critically dense region itself. To evaluate the density in the boundary layers, Equations (10) and (11) are substituted in Equations (16)-(18) with the result, for $x \geq x_1$, that

$$n(x) = n_3(x) + \left\{ n_1(x_0) \exp(-y_{01}) K_0^1(b_1(x_1), v_{01}^2(x_1), 0) - n_3(x_1) K_0^1(b_3(x_1), v_{03}^2(x_1), 0) \right\} \quad (27)$$

In Equation (27), the function $K_n^\nu(a, b, c)$ is defined by

$$K_m^\nu(a, b, c) = \pi^{-1/2} \sum_{n=1}^{\infty} a_n^0 \int_0^{\infty} dz a_n^\nu(c) (z+c)^{1-\nu/2+m} \exp\left[-z - \frac{\lambda_n^2 a}{(z+c)^{1/2} (z+c+b)^{3/2}}\right] \quad (28)$$

and

$$v_{0a}^2 = m v_0^2 / 2T_a, \quad b_a(x_*) = |x - x_*| / 4\lambda_{ei}^a(x_*),$$

and

$$\lambda_{ei}^a(x) = (T_a/m_e)^{1/2} / \bar{v}_{eia}(x).$$

The coefficients are given by

$$a_n^\nu(c) = \int_0^r du u U_n(u) (1-u^2)^{-\nu/2} \quad (29)$$

where

$$r^2 = z/(z+c) \quad a_n^0 \equiv a_n^0(0)$$

The identity

$$\sum_n a_n^0 a_n^1(c) = 1 - (1-r^2)^{1/2} \quad (30)$$

will prove useful in showing how Equation (27) relates to Equation (23).

The density clearly reduces to the correct limit in region III if the terms in parentheses in Equations (27) decay to zero as, in fact, they were constructed to do. To demonstrate the explicit decay, the asymptotic expansion of the function $K_n^\nu(a, b, c)$ is given, for $a \rightarrow \infty$, by

$$K_m^\nu(a, b, c) \xrightarrow{a \rightarrow \infty} \sum_n a_n^0 a_n^\nu(0) (2\lambda_n^2 a)^{[2-(\nu+1)/2+m]/3} \exp(c) \exp\left(-\frac{3}{2}(2\lambda_n^2 a)^{1/3}\right) \quad (31)$$

provided $(2\lambda_n^2 a)^{1/3} > c, 1$. Since λ_n increases with n , the $n=1$ contribution dominates the series in the asymptotic region.

The current and heat flow in the boundary layer are obtained by substituting Equations (10) and (11) in Equation (17) and (18). The continuity of the heat flow and current can be explicitly demonstrated in the same way as for the density but terms first order in λ_{ei}/L_e must be kept to obtain exact matching.

Before the results of this analysis can be used in simulations, the heat sources (e.g., inverse bremsstrahlung) in the critical region, the scattering terms that are proportional to the ponderomotive potential, and the fast electrons that are not diffusive anywhere but do contribute heavily to the heat flow, as well as more general self-consistent profiles must be included.

The present analysis shows that the density near critical has several scale lengths. There is the scale length determined by the ponderomotive force and, on either side of this region, there are the regions where the scale length is characteristic of a mean free path over which electron distributions are isotropized. The density that is consistent with the model potential,

$$\Phi(x) = \Phi_0 (x - x_0)/L_2; \quad x_0 \leq x \leq L_2;$$

$$\Phi(x) = \Phi_0 [1 + (x - x_1)/L_3]; \quad x_1 \leq x \quad (32)$$

is shown in Figure 3.2-1 for $0 \leq (x - x_0)/L_2 \leq 20$. The parameters

$$\gamma_{01} = 1, T_1(x_0)/T_3(x_1) = 0.25,$$

$L_2/L_3 = 0.01$, $v_{01}^2(x_1) = 1$, $L_2/\lambda_{ei}^{(3)}(x_1) = 0.5$ were chosen. In addition, the function $K_0^1(a_1, a_2, 0)$ was approximated by setting

$$a^0 = a_1^1(0) = 1; \quad a_n^0 = a_n^1(0), \quad n \neq 1; \quad \text{and } \lambda_1 = 4.$$

This approximation preserves the scaling of the exact solution, uses a value of λ_1 that is close to the exact value, and is close to the exact solution everywhere but very near x_1 since $\lambda_{n+1} > \lambda_n$ for all n . More importantly, the 'correct expression' for the heat flow is calculated in the steepened profile. In simulations, it is common to use the free-streaming expression for the heat flow, i.e., $q_s = fn_e v_e T_e$ where f is the so called flux limit parameter, whose value is typically taken to be less than one and

greater than 0.05. According to Equation (26), if $T_3 \gg T_1$, the heat flux is $q_f = \sqrt{2/\pi} n_3(x_1) T_3(x_1) (T_3(x_1)/m_e)^{1/2}$. If one interprets f as the ratio of q_f to $n_e T_e v_e$, then $f \cong n_3(x_1)/n(x)$ between x_0 and x_1 and changes continuously outside of this region to the value given using Equation (22). In a strongly steepened profile $n_3(x_1)/n(x) \ll 1$ when $n(x) > n_c$ where n_c is the critical density and, thus, $f \ll 1$ so that the heat flow would be interpreted as inhibited.

3.3

Hot Electron Production with Multiple Frequency Illumination

T. Speziale

A central concern in the laser fusion effort is the presence of so called hot electrons which possess energies in excess of a few kilovolts. The mean free path of the suprathermal electrons can exceed the target diameter and allows the electrons to deposit energy in the core region of the target. The degree to which the fuel can be compressed is thereby diminished so that the thermonuclear yield is adversely affected.

Energetic electrons are produced when electron plasma oscillations¹ are resonantly driven^{2,3} to high amplitude by an electromagnetic wave³ having a component of electric field along the plasma density gradient.^{4,5} The interaction is largely one-dimensional and a capacitor model² has been employed to estimate the energy carried by the suprathermal electrons. In such a model, the electrons (electron fluid elements) behave as independent, harmonic oscillators. When the secularly growing electron amplitude^{2,6} becomes as large as the spatial extent of the resonance, wavebreaking^{1,2} occurs and the energetic electrons leave the resonance region. If E_d is the amplitude of the (monochromatic) driver electric field and L is the length scale for density variation at the resonance point, then the characteristic energy of the suprathermal electrons is proportional to $E_d L$.

In the laser plasma interaction, therefore, the hot electron energy increases with E_d^2 (L also depends on E_d^2 but we will take it to be fixed for simplicity) while the energy absorbed^{3,5} by the target is propor-

tional to E_d^2 . Thus, the hot electron energy can be decreased by lowering E_d^2 but at the expense of reduced energy absorption (provided collisionless conditions prevail). A potential remedy to this dilemma lies in the use of multiple frequency illumination with a sufficiently large (to be defined) frequency separation between spectral components. In that way, the total laser energy available can be distributed among many spectral lines while the effective driver amplitude at each of the many critical surfaces (resonance points) is reduced. It is shown herein that if the laser energy is distributed among N lines then the characteristic energy of the suprathermals can be reduced by the factor $N^{-1/2}$ when the frequency separation is sufficient. The range of an electron passing through the background plasma scales as the square of its energy. Thus, in the present example, the hot electron range would be reduced by the factor N^{-1} .

The effect of multiple frequency illumination on the hot electron energy has been investigated using a capacitor model. This model has been employed previously and, in general, leads to the oscillator equation:^{2,3}

$$\ddot{\delta} + \omega_p^2(x_0)\delta = \mathcal{E}(t) \quad (1)$$

for the displacement, δ , of an electron (fluid element) from its initial position. The quantity $\mathcal{E}$ is the driver field which is related to the capacitor electric field via $\mathcal{E} = -(e/m) E_d$ with e/m the electron charge to mass ratio. The characteristic frequency associated with each oscillator is the plasma frequency $\omega_p(x_0) = [4\pi e^2 n(x_0)/m]^{1/2}$ in which $n(x_0)$ is the electron number density at $\tau = 0$. The quantities x_0 and τ represent Lagrangian¹ coordinates and are defined via the relations

$$x = x_0 + \int_0^\tau u(x_0, \tau') d\tau'$$

and $t = \tau$.

The coordinates x and t are the usual Eulerian variables. The mapping of one coordinate system into the other is single valued so long as $\partial\delta/\partial x_0 > -1$. Notice that the oscillator equation possesses a particularly simple form in the Lagrangian frame.

The driver field can be expressed as

$$\mathcal{E}(t) = \sum_{j=1}^N \mathcal{E}_j \cos(\omega_j \tau + \varphi_j) \quad (2)$$

which corresponds to, for example, a laser which produces N discrete frequencies, ω_j , with corresponding amplitudes $\mathcal{E}_j$. In general, the phase angles, φ_j , are uncorrelated and randomly distributed over a 2π interval. The amplitudes $\mathcal{E}_j$ are to be prescribed and, for now, the normalization

$$\mathcal{E}_0^2 = \sum_{j=1}^N \mathcal{E}_j^2 \quad (3)$$

of the total, time averaged energy density contained in the driver field is the only restriction placed on the amplitudes. This allows comparison with the equivalent, well-known problem of a single frequency ($N=1$) driver with amplitude $\mathcal{E}_0$. The solution of the oscillator equation with the more general driver, Equation (2), is straightforward. A convenient form of the solution is given by

$$\delta \cong \sum_{j=1}^N \frac{\mathcal{E}_j}{\omega_j} \frac{\sin(\Delta_j \tau / 2)}{\Delta_j} \sin\left(\omega_j \tau + \frac{\Delta_j \tau}{2} + \varphi_j\right) \quad (4)$$

in the "resonance approximation" wherein the detuning functions $\Delta_j = \omega_p(x_0) - \omega_j$ are regarded as small compared with ω_j . At this point, the object of the calculation is to first determine the time at which wavebreaking occurs which is expressed formally by the relation² $\partial\delta/\partial x_0 = -1$. Having established the wavebreaking time, τ_b , one then seeks to evaluate the electron velocity, $v_b = |\partial\delta/\partial\tau|_{\tau=\tau_b}$ in order to obtain the hot electron energy.

It is instructive to briefly consider the single driver ($N=1$) solution

$$\delta \cong \frac{\mathcal{E}_0}{\omega_0} \frac{\sin(\Delta_0 \tau / 2)}{\Delta_0} \sin\left(\omega_0 \tau + \frac{\Delta_0 \tau}{2}\right) \quad (5)$$

which has the following properties. At the critical surface $\Delta_0 \equiv 0$ and $|\delta| \sim \mathcal{E}_0 \tau / 2\omega_0$. Also, the width of the region over which δ is large is given approximately by $\Delta_0 \tau \sim \pi$. If the density profile is linear with length scale L , then the spatial extent of the resonance is found to be $\Delta x \sim 2\pi L / \omega \tau$ so that $|\delta| \sim \Delta x$ when $\tau \sim \sqrt{L/\mathcal{E}_0}$. Application of the formal wavebreaking condition gives

$$\tau_b = (8L/\mathcal{E}_0)^{1/2} \quad (6)$$

with the corresponding electron velocity at wave-breaking

$$v_b = \frac{\mathcal{E}_0 \tau_b}{2} = (2 \mathcal{E}_0 L)^{1/2}. \quad (7)$$

This is larger than the oscillatory velocity,

$$v_0 = \mathcal{E}_0 / \omega_0,$$

of a free electron in the driver field by a factor of $\omega_0 \tau_b / 2$.

Now consider a multiple frequency driver with arbitrary N . Each frequency component of the driver has a critical surface ($\Delta_j = 0$) associated with it. Also, the plasma response associated with each critical surface is localized, as in the single driver case, to a narrow region about the critical surface and is proportional to $\sin(\Delta_j \tau/2)/\Delta_j$. The simplest case to consider is that in which the separation between successive critical surfaces exceeds the spatial extent of the plasma resonances. In this case, hot electron production occurs at each critical surface without the interaction of the neighboring surface so that the problem reduces to N single driver problems. The condition that a given pair of frequencies, ω_1 and ω_2 produce hot electrons independently can be expressed as

$$\frac{\omega_1 - \omega_2}{\omega_0} > \sqrt{\frac{v_1/c}{kL}} \quad (8)$$

for a linear density gradient, $\omega_p = \omega_0 (1 + x/2L)$. In Equation (8), v_1 represents the oscillatory velocity of an electron under the influence of $\mathcal{E}_1 \approx \mathcal{E}_2$ (disparate amplitudes are not as interesting) and k is the vacuum wave number associated with a light wave of frequency $\omega_1 \approx \omega_2$ (for example). If all N driver amplitudes are about equal then $v_1 \sim v_0 N^{1/2}$. In that case, for $v_0/c \approx 0.01$ and $kL \sim 2\pi$, the right hand side of the inequality, Equation (8), becomes $0.04 N^{1/2}$. Note that v_0/c is essentially the same as that corresponding to the incident laser field [for $kL = 0(1)$] when *all* the laser light is incident at the optimum coupling angle. The driver electric field is related to the vacuum field via $E_d(\pi kL)^{1/2} = E_0(A \cos\theta)^{1/2}$ in which θ is the angle of incidence and A is the (resonance) absorption

coefficient.⁴

When the separation between successive frequencies satisfies Equation (8), the electron velocity at wavebreaking is given by Equation (7) with $\mathcal{E}_0$ replaced by $\mathcal{E}_j$ for each driver component. Thus, for example, if all the $\mathcal{E}_j$ are equal, then Equation (7) yields

$$v_b^2 = 2 \mathcal{E}_0 L / \sqrt{N} \quad (9)$$

through the use of the normalization, Equation (3). If R_0 is the hot electron range in the single driver case, then the hot electron range will therefore scale as

$$R \approx R_0 / N \quad (10)$$

in the N frequency case.

The results indicated by Equations (9) and (10) are valid provided the frequency separation between spectral lines satisfies Equation (8). When the inequality, Equation (8), is reversed, the frequency separation is small and the plasma response at one critical surface significantly affects the response at each of the others. The limit of small frequency separation has also been examined for the cases $N = 2$ and $N \gg 1$. Neither case produces a result differing significantly from Equation (7) in the limit of vanishingly small frequency separations. This result is to be expected because the plasma should respond as though driven by a single frequency driver in that limit. The small frequency separation calculations are somewhat more involved and the reader may refer to the appropriate KMSF report for details.

In conclusion, it has been demonstrated that the use of multiple frequency illumination in laser fusion experiments can reduce the hot electron energy. The coupling of hot electron energy to the fuel containing core of a fusion target depends first on the ability of an electron to reach the core. A measure of this ability is the hot electron range which scales as the square of the electron energy and should be compared with its shell thickness. Irradiation of the fusion target by light containing N discrete frequencies has been shown to be capable of reducing the hot electron range by the factor N^{-1} . This result is valid when the frequency separation is large and the individual intensities are of the same order.

3.4

Rarefaction Shocks in Spherical Geometry

D. Mitrovich

When the electrons in a plasma have a bi-Maxwellian distribution, the plasma can sustain a rarefaction shock if the temperature ratio between the Maxwellians exceeds a critical value of about 10. The theory of these shocks in planar geometry has been worked out by Bezzerides, Forslund, and Lindman.¹ Their analysis derives the appropriate Hugoniot relation for the shock, and shows that, in general, the shock attaches to the downstream plasma expansion wave via an expanding shelf of uniform density and velocity. The shock conditions were determined by requiring equilibrium flow at the upstream shock boundary, plasma neutrality on either side of the shock, and zero net charge in the shock. The shelf occurs because the plasma flow speed at the downstream shock boundary exceeds the local sound speed (determined mainly by the hot electron temperature) at which the after shock expansion wave propagates upstream.

Spherical geometry can modify the expanding plasma flow so that no constant density shelf forms. Figure 3.4-1a shows the isothermal expansion of an initially uniform density plasma ball ($Z=10$), in which the hot to cold electron temperature ratio $\alpha \equiv \theta_H/\theta_C = 100$, at a time $\omega_{pi}t = 65$. The densities and plasma ion frequency are referred to the density $\tilde{n}$, and the flow speed is referred to the sound speed at the cold electron temperature. Radii are in units of the cold electron reference Debye length λ . Both upstream and downstream expansion waves are clearly evident, and the latter is seen to attach directly to the rarefaction shock. The shock thickness by this time has increased to about 8 percent of the radius. A similar type of simulation performed by M. True² was perplexing because it did exhibit a density shelf. The shelf can be ascribed to the effects of planar geometry as can be seen in Figure 3.4-1b. Here simulation parameters are identical to those in Figure 3.4-1a except that the initial radius of the plasma ball has been increased by adding 24,000 λ to the radius, effectively making the geometry planar. Now the constant den-

Fig. 3.4-1a

Fig. 3.4-1b.

Fig. 3.4-1. (a) Density and velocity distributions at $\tilde{\omega}_{pi}t = 65$ of an isothermal, two-electron temperature spherical plasma ball in which the temperature ratio is $\alpha = 100$. The reference Debye length $\tilde{\lambda}$ and sound speed refer to the reference density $\tilde{n}$ and cold electron temperature θ_C . The ball's initial density is $N_0 = n_e/\tilde{n} = 30$. Charge separation can be seen in the shock region. (b) Same conditions except that the ball's initial radius has been increased by a factor of approximately 25.

sity, constant velocity shelf is clearly visible. The shock relations and shelf parameters are precisely those derived in Reference 1.

The simulations were performed with a finely zoned, spherically symmetric Lagrangian hydrodynamic code. Both hot and cold electron temperatures, Θ_H and Θ_C , were held fixed, and the ion temperature was assumed equal to Θ_C . Charge separation effects were included by solving the Poisson equation in the shock region, and the plasma was assumed to be quasi-neutral elsewhere. The electron populations were in equilibrium with the potential field $\Phi \equiv e\varphi/\Theta_C N_C \equiv n_C/\tilde{n} = e^*$ and $N_H \equiv n_H/\tilde{n} = e^*/\alpha$. The initial electron density of the plasma shell was taken as $n_e/\tilde{n} = 30$.

The detailed results of these simulations are related to earlier results of two electron temperature TRHYD simulations in which charge separation was ignored. Even at late times these simulations show a rarefaction shock whose behavior does not resemble that discussed by Bezzerides, et al. Figure 3.4-2 shows the velocity and density distributions at time $\omega_{pi}t = 3.97 \times 10^3$ of a deuterium-tritium plasma in which the reference electron density $\tilde{n} = 10^{21} \text{ cm}^{-3}$, the cold electron temperature $\Theta_H = 1 \text{ keV}$, and the temperature ratio $\alpha = 25$. If the rarefaction shock were to follow the conditions for a planar shock (dashed lines in Figure 3.4-2), the density at the upstream (u) and downstream (d) boundaries would be $N_{e_u} = (n_{e_u} + n_{H_u})/\tilde{n} = 2.3$ and $N_{e_d} = 0.454$, and the change in flow velocity would be $1.29 \times 10^8 \text{ cm/sec}$. The variation from the planar shock conditions appears related to how the shock attaches to the downstream plasma. The attachment is smooth, without a gradient discontinuity. It resembles somewhat the shock-to-expansion transition of Figure 3.4-1a (in which, however, $\alpha = 100$), except that the shock "tail" is considerably longer in Figure 3.4-2.

The late time density and velocity distributions of Figures 3.4-2 can be understood on the basis of one-temperature, isothermal, spherically symmetric similarity flow solutions, with the plasma assumed quasi-neutral. The continuity and momentum equations, with the definitions $c_s \equiv \sqrt{\Theta/M}$, $V \equiv v/c_s$, $s \equiv c_s t$, are

$$\frac{\partial n}{\partial s} + \frac{1}{r^2} \frac{\partial}{\partial r} (r^2 n V) = 0 \quad (1)$$

$$n \left(\frac{\partial}{\partial s} + V \frac{\partial}{\partial r} \right) V = - \frac{\partial}{\partial r} n \quad (2)$$

Defining $X \equiv \ell n n$ and the similarity variable $\xi \equiv r/s$, these equations become

$$(V - \xi) X' + V' + \frac{2}{\xi} V = 0 \quad (3)$$

$$(V - \xi) V' + X' = 0 \quad (4)$$

where the prime denotes $d/d\xi$. Combining these to eliminate X' gives

$$\frac{dV}{d\xi} = \frac{2V/\xi}{(V - \xi)^2 - 1} \quad (5)$$

This is singular when $V = \xi \pm 1$ at which points the assumption of quasineutrality breaks down. Figure 3.4-3 shows the function $V(\xi)$ (solid curves) plotted starting at different values of ξ_* at which $V = \xi_* + 1$, along with the corresponding density distribu-

Fig. 3.4-2. Density and velocity distributions at late time ($\omega_{pi}t = 3.97 \times 10^3$) from a hydro code simulation, without charge separation, of the expansion of a two-electron temperature deuterium tritium plasma ball with a temperature ratio $\alpha = 25$ and cold electron temperature $\theta_c = 1 \text{ keV}$. The reference density was chosen as $\tilde{n} = 10^{21} \text{ cm}^{-3}$. The shock jump conditions for planar geometry are shown by the dashed lines.

tions $n = \exp(X)$.³ The asymptotic behavior of the solutions follows

$$V \approx \xi + \sqrt{3}$$

$$n \approx \text{const} \times \exp(-\sqrt{3}\xi). \quad (6)$$

Letting the reference plasma temperature be the hot electron temperature θ_h , the velocity distribution of Figure 3.4-2 can be plotted in terms of the similarity variable ξ . This is depicted as the dashed curve in Figure 3.4-3, showing the similarity between the flow downstream of the shock and the similarity solutions. The simulation velocity distribution beyond the shock is displaced downward from the similarity solution, but in shape matches it almost exactly (the source of the displacement is not understood). The singularity in the similarity flow is replaced by the rarefaction shock, but the flow modifies the shock so that the jump conditions differ markedly from those of planar flow. In particular the shock can merge smoothly with the downstream expansion flow, with part of the density and velocity jumps, which in planar geometry are

Fig. 3.4-3. Similarity solutions in spherical geometry for velocity (V) and density [$X \equiv \ln(n)$] for an isothermal, single-temperature plasma, for different values of the singularity position ξ_* . The normalization for n was chosen arbitrarily to be $X_* = 7$. With the hot-electron temperature as referent, the center curves with $\xi_* = 0.89$ correspond to the distributions in Fig. 3.4-2, whose velocity is plotted as the dashed line.

concentrated in the shock, now distributed over the self-similar expansion flow. The detailed way in which this occurs has not yet been worked out.

3.5

Simulation Code Development

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As diagnostics become more sophisticated and our knowledge of laser plasma interactions expands, improvements must be made to simulation codes. Here, we report on two such developments. Section 3.5.1, reports on the use of a new numerical technique in a Lagrangian hydrodynamics code that allows locally small scale spatial resolution without sacrificing description of the entire target and without excessively increasing computational time. In Section 3.5.2, the improvements in the code to the energy deposition in plasma and transport of nuclear reaction products through the target are described.

3.5.1 FRHYD: A One-Dimension Fluid Simulation Code with Automatic Fine Zoning

Numerical simulation of laser fusion has been done almost exclusively with Lagrangian hydrodynamics. The number of zones used in the mesh affects the computer time needed for a calculation in two ways. First, since the same calculation must be performed for each zone (or class of zones), the number of times a given calculation is done is proportional to the number of zones. Second, as the zone size is reduced, stability of the difference scheme (e.g., the Courant condition) requires that the simulation time step be made smaller. In other words, the time required to complete a computer calculation goes roughly as the square of the number of zones. Thus, there is always a compromise between the detail desired in a given calculation and the size and speed of the computer available.

In some situations, such as at material boun-

daries and in the laser deposition region, the zones can be made smaller in one place to provide better resolution, leaving the zones larger elsewhere. However, this only solves the number of zones problem, the time step is still reduced. In other cases, such as shock waves, the region where better resolution is needed, while localized, travels through the mesh, and no such local fine zoning is possible.

Automatic fine zoning attempts to tackle both of the problems: first, small zones (fine zones or FZ) are added in regions where greater resolution is needed and then move as this region moves through the mesh of larger zones (macro zones or MZ). Secondly, the calculation is done separately—with separate time steps—in the regions containing fine zones.

The fine zones take care of conditions requiring small zones, so that the background mesh can consist of larger zones. Since the fine zone regions are advanced on a separate time step, the background mesh time steps can be large and the calculation using a small time step is limited to a few fine zones. This report describes a code, FRHYD, a one-dimensional laser fusion simulation code that incorporates these ideas.

General Description of the Code: The FRHYD code consists of a main line program which handles most of the logic of running the separate FZ and MZ regions. The main code calls subroutines which perform the calculations involved in solving the equations of the model. These subroutines are time step control, laser deposition, hydrodynamics and nuclear reactions, forward energy calculation, backward energy calculation, and output.

The fine zones are managed in a routine which performs several functions: (1) decides whether to add a new type of fine zones, or completely drop an existing type, (2) decides when to add a single zone to an existing type or drop a zone from an existing type, (3) manages the subscript arrays involved in keeping FZ and MZ sets separate, and (4) watches for the various pathologies that can occur when types of FZ's merge, separate, or cross MZ boundaries.

A type of fine zones is added as a block (of 4) at a region where finer zoning is desired, such as when a shock begins to form or in the laser deposition re-

gion. The code presently allows up to five independent types of fine zones to exist simultaneously. We refer to all fine zones of a particular type as a *FZ group* to distinguish them from a *FZ set* defined below.

We give here a brief description of how a shock front FZ group works: A FZ group follows a shock front by adding zones in front of the shock and dropping zones from the region behind it. In this way the region of the shock discontinuity is covered by small zones, allowing the shock to be spread out over several zones yet remain quite sharp.

The mesh of macro zones overlaid with FZ groups is collected into sets to solve the differential equation. Each FZ set consists of all the contiguous (or overlapping) FZ groups separated by less than three macro zones. The contiguous regions of macro zones not included in FZ sets form the remaining MZ sets.

The differential equations are solved for each of these sets of zones. All MZ sets are run with the same time step and all FZ sets are run with the same time step. A complete MZ time step begins with both the MZ and FZ sets at the same (problem) time. The new MZ time step is determined to give the next MZ time and the FZ sets are run until this time is reached, adjusting the last FZ time step as necessary. The MZ sets are then also run to this time and the cycle begins again. Fine zones may be added and dropped during any FZ time step. FZ groups may be added or dropped only at the end of a MZ time step when the macro zones and fine zones are at the same time.

The FZ sets are advanced in time first. A few of the reasons for this are given here. The bounding macro zones are held at constant temperature because MZ temperatures are slowly varying. If the MZ sets were advanced first, extrapolation of the bounding FZ temperatures would be required. A completely implicit solution (involving 10-20 FZ time steps and one MZ time step) is neither necessary nor easy. As the FZ sets are run, the dropping and adding of zones changes the size of the bounding macro zones. This would be difficult to handle if the macro zones were run first.

The overall flow of the main line code is now discussed. The first step is to determine the size of the next MZ time step. Both the old and new times are kept. After adding or dropping FZ groups, the

FZ running sets, if present, are advanced by the FZ loop. The FZ loop advances using its time step until the FZ time reaches the MZ time. Then the macro zones are advanced to this common time and the main loop starts again.

The hydrodynamic and nuclear equations are solved explicitly. The advancing of the fine zones for several time steps before the macro zones are advanced makes zone boundaries between FZ and MZ special. The velocities and positions of these boundary zones are advanced in the FZ loop. The most compelling reason to do it then is that the boundary zones are continually being added and dropped. Only the positions and velocities of these zone boundaries are changed; the pressure, temperature, etc. of the neighboring macro zone are not because, by design, regions of rapid change are covered by fine zones. The effective mass of the boundary zone may change as zones are added or dropped. The other boundary of each neighboring zone of the MZ set (except when the FZ set is an inner or outer boundary set) has its original position saved and is advanced using its old velocity. When the MZ sets are advanced, the saved positions of these zone boundaries are restored and then advanced. In the same way the boundaries of the MZ set (which were advanced during the FZ time step) are temporarily extrapolated backward for the MZ time step calculation.

The hydrodynamics part of FRHYD conserves energy by calculating the net work done by the internal pressure of a zone and subtracting this amount from the internal energy of that zone. Since the calculation is done at boundaries, each zone is adjusted twice. For a macro zone bordering a FZ set this means that the adjustment at one boundary gets spread over many FZ time steps while at the other boundary the adjustment is done in one MZ time step. A problem exists when a fine zone is added or dropped, for example, when the dropped FZ energy is added to the MZ energy a new temperature cannot be calculated as the MZ energy is (temporarily) out of balance. The solution is to run the bordering MZ as part of the FZ set. The alternate approach of accumulating the increment, but adding it in only at the beginning of the MZ time step was rejected because it would have required separate hydrodynamics routines for the two calculations (FZ and MZ) at the boundaries. This is one of the

problems that requires dropping a fine zone on a MZ boundary only after a MZ time step.

FRHYD solves the energy difference equations using an implicit scheme. Each running set is solved with its own boundary conditions. The standard scheme for solving the resulting tri-diagonal equations is used: a forward pass through the mesh to evaluate the coefficients which are then used in a subsequent backward pass to evaluate the new energies. The boundary conditions are quite different for the MZ sets and the FZ sets. The FZ set boundary condition assumes that the neighboring MZ temperature remains constant during the FZ time steps. However, since the MZ time step can be many (even 10 or more) FZ time steps, this assumption might cause problems for certain types of fine zones. Present fine zone conditions tend to keep the temperature gradient low at a MZ-FZ boundary. If this proves not to be the case, the temperature of macro zone boundary FZ sets will have to be interpolated in the same manner as in the position of the zone boundary. As each FZ set is advanced, the energy flux across the boundary with the neighboring MZ sets is saved. The accumulated flux is added as a boundary condition when the MZ set is advanced. A complication arises because there are both implicit and explicit components to the flux, and they are calculated at different points in the code.

Application: Fine Zoning of the Critical Density Surface. The form of the density profile near the critical density surface is of considerable importance in laser fusion research. The density gradient affects absorption, heat flow, and Brillouin scattering rates. However, standard Lagrangian hydrodynamic codes cannot resolve the density gradient sufficiently under the effect of profile steepening due to ponderomotive force or flux inhibition. Reasons for this difficulty are that (1) Lagrangian zones expand in the vicinity of the critical density surface, (2) it is difficult to predict the position of the critical surface within the mesh, and (3) the critical surface moves through the mesh.

The introduction of a fine zone group around the critical density surface allows resolution of the steepened profile in this region and causes significant changes in the simulation results. The fine zone group used in this illustration extends a distance of five laser wavelengths on either side of the critical

density surface. The width of the fine zones at the critical surface is $\leq 0.3 \lambda_{\text{laser}}$. Figure 3.5-1 compares the density and temperature profiles obtained from simulations with and without fine zones for a $100 \mu\text{m} \times 1 \mu\text{m}$ target with laser intensity 1.6×10^{15} watts/cm². The plasma scale length at the critical

Fig. 3.5-1. Comparisons of electron density profiles (solid line) and electron temperature profiles (dotted lines) obtained from hydrodynamic simulations for a $100 \mu\text{m} \times 1 \mu\text{m}$ target with incident intensity 1.6×10^{15} watts/cm² and with (a) standard Lagrangian zoning scheme and (b) zoning including a fine-zone group extending 5λ on either side of the critical density surface.

surface is much smaller in the simulation including fine zoning. This results in a significant reduction in collisional absorption and corona temperature over what would be predicted using a normal mesh. Table 3.5-1 summarizes the effects of profile steepening due to ponderomotive force on inverse bremsstrahlung absorption, corona temperature, and Brillouin reflection in simulations including and not including fine zones. In the two cases which include Brillouin scattering the total absorption is fixed, with the energy absorbed distributed uniformly with respect to mass from $5 \mu\text{m}$ inside the critical density surface to the outside of the target. The difference in electron temperature in these two cases can be understood by noting that the mass outside the critical surface will be less when ponderomotive force is included. Ponderomotive force changes the results significantly when fine zones are included to better resolve the density gradient at the critical surface, whereas in runs not including fine zones the effect of ponderomotive force is small.

Table 3.5-1

Summary of the Effects of Profile-Steepening Due to Ponderomotive Force on Inverse Bremsstrahlung Absorption, Corona Temperature, and Brillouin Reflection in Simulations Including and not Including Fine Zones

Intensity: 1.6×10^{15} watts/cm ²				
Dimensions: 100 microns $\times$ 1 micron				
Time: 250 picoseconds				
Ponderomotive Force	Fine Zones	Absorption (Joules)	Brillouin Reflection	TEcrit (keV)
No	No	23		2.3
No	Yes	20		2.3
Yes	No	21		2.2
Yes	Yes	7.4		1.3
No	Yes	9.9	0.8	0.5
Yes	Yes	9.4	0.3	1.2

3.5.2 Thermonuclear Energy Transport in a Dense Plasma

Over the last several years we have designed, as part of our numerical simulation, an elaborate model for nuclear processes occurring within a laser fusion target. This model describes the rate of nuclear reactions, the transport of the nuclear reaction products through the target, and the transfer of energy from the nuclear particles to the plasma. We have made improvements to the nuclear energy transfer section of the model. Below we describe the energy loss mechanisms of a fast nuclear particle to the ambient plasma.

Energy Loss of a Fast Charged Nuclear Particle to a Plasma. A charged nuclear particle passing through a dense plasma gives up energy through collisions with the plasma electrons and ions. The collisions with the electrons are by Coulomb scattering. If we model the energy losses as continuum processes in a material, we can write an energy loss equation in formal terms as

$$\frac{dE}{ds} = F_{c,e} + F_{c,i} + F_{n,i} + F_{be} + F_{plas} \quad (1)$$

where dE/ds is the particle energy loss per unit length travelled, $F_{c,e}$ is the energy loss by Coulomb scattering with the free electrons in the plasma, $F_{c,i}$ is the energy loss by Coulomb scattering with the ions, $F_{n,i}$ is the energy loss by nuclear scattering with the ions, F_{be} is the energy loss to bound electrons attached to the ions, and F_{plas} is the energy loss to electron plasma oscillations. The improvements to the energy loss model are the addition of the last two terms, namely, energy losses to bound electrons and to electron plasma oscillations. Our previous approach had been to ignore the electron plasma oscillations and to make no distinction between free and bound electrons. But at cold material temperatures these effects become important.

General Formula for Coulomb Energy Loss. The general formula for Coulomb energy loss from a fast charged particle may be written¹

$$F_{coul} = -2\pi N_a e^4 Z_f^2 Z_a^2 \ln \Lambda (1/\theta_a) G(x_a) \quad (2)$$

where indices f and a denote fast and ambient particles, respectively, N_a is the ion number density, Z_a and Z_f are the particle charges, $\ln \Lambda$ is the Coulomb logarithm, θ_a is the ambient particle temperature, $G(x)$ is a dimensionless function, $x_a = [m_a E_f / m_f \theta_a]$ is the ratio of fast to ambient particle velocities. The function $G(x)$ is²

$$G(x) = \frac{1}{2x^2} \left[\operatorname{erf}(x) - \frac{2}{\pi} \left(1 + \frac{m_a}{m_f} \right) \times \exp(-x^2) \right] \quad (3)$$

The Coulomb logarithm, $\ln \Lambda$, is given by

$$\ln \Lambda = \frac{1}{2} \ln \left[1 + \frac{b_{max}^2}{b_{min}^2} \right] \quad (4)$$

where b_{max} and b_{min} are the maximum and minimum impact parameters. In some of the energy loss terms we define a maximum impact parameter b_{max} by the following relations,

$$b_{max} = R_{screen} \quad (5)$$

$$R_{screen} = \operatorname{Max}(r_i, \lambda_e) \quad (6)$$

$$r_i = \left[\frac{3}{4\pi N_i} \right]^{1/3} \quad (7)$$

= atomic cell size for ion number density N_i .

$$\lambda_e = \left[\frac{\theta_e}{4\pi N_e e^2} \right]^{1/2} \quad (8)$$

= electron Debye length for electron number density N_e and temperature θ_e . And, similarly, we define a minimum impact parameter b_{min} by the following relations,

$$b_{min}^2 = b_c^2 + b_q^2 \quad (9)$$

$$b_c = Z_f Z_a e^2 / 2E \quad (10)$$

= classical turning point for a particle in a Coulomb potential,

$$b_q = \hbar / (2mE)^{1/2} \quad (11)$$

= de Broglie wavelength of the fast and ambient particles in their c.m. system, where E = relative energy of particles f and a , and m = reduced mass of particles f and a . In atomic units we then have

$$\left[\frac{b_{min}}{a_0} \right]^2 = \frac{E_{ryd}}{E} \left[\left(\frac{E_{ryd}}{E} \right) Z_f^2 Z_a^2 + \frac{m_e}{m} \right] \quad (12)$$

where E_{ryd} = Rydberg energy for hydrogen, and a_0 = Bohr radius for hydrogen. All the Coulomb energy loss terms, including losses to bound electrons and to plasma oscillations, can be written in the general form of Equation (1), but with different meanings attached to the G function and to the Coulomb logarithm. Some of these results have already been reported.³

Energy Loss to Free Electrons. We apply the general formula of Equations (2)-(4) to the case of electrons as the ambient particles. If we factor out the electron Debye length λ_e , we can write the energy loss per unit length as

$$F_{c,fe} = -\frac{e^2}{\lambda_e^2} Z_f^2 \ln \Lambda_e G(x_e). \quad (13)$$

The deposition function $G(x_e)$ must be evaluated in its general form, Equation (3), because the argument x_e takes on both very small and very large values. The mass ratio m_e/m_f in the second term of G can be ignored, and then the function reduces to the Spitzer form.² For numerical calculations we approximate G by rational fractions adapted to the asymptotic forms at small and large arguments. The Coulomb logarithm is given by Equations (5)-(12), with the reduced mass, $m = m_e$, and relative energy

$$m = m_e$$

and

$$E = \frac{m_e}{m_f} E_f + \frac{3}{2} \theta_e. \quad (14)$$

Energy Loss to Ions. We apply the general formula of Equations (2)-(4) to the case of ions as the ambient particles. If we factor out an "ion Debye length,"

$$\lambda_i = \left[\frac{\theta_i}{4 N_i \bar{Z}^2 e^2} \right]^{1/2} \quad (15)$$

where θ_i = ambient ion temperature, N_i = ion number density, and $\bar{Z}$ = average ionic charge, we can write the energy loss per unit length as

$$F_{c,i} = -\frac{e^2}{\lambda_i^2} Z_f^2 \ln \Lambda_i G(x_i). \quad (16)$$

Because x_i , the velocity ratio of fast particle to ambient ion is large, we can approximate the deposition function

$$G(x_i) \approx \frac{1}{2 x_i^2}. \quad (17)$$

The Coulomb logarithm is given by Equations (5)-(12) with $\bar{Z}$ as the average charge of the ambient ions, and with the reduced mass and relative energy

$$m = \frac{m_f m_i}{m_f + m_i}$$

$$E = m \left[\frac{E_f}{m_f} + \frac{E_i}{m_i} \right]$$

$$E_i = \frac{3}{2} \theta_i. \quad (18)$$

Nuclear Scattering Energy Loss to Ions. The energy loss of a fast particle from nuclear scattering on the ambient ions is important in materials with low to medium atomic number Z . When the atomic number Z is high, the Coulomb energy losses, described in the previous sections, are so large that we may neglect the nuclear loss term. And so, in our simulation calculations, we consider nuclear scattering losses only in those regions that contain nuclear fuel (D and T). We write the energy loss per unit length in the form

$$F_{n,i} = -N_i \sigma^T(fi) E_f \quad (19)$$

where N_i is the ion number density, $\sigma^T(fi)$ is the transport cross section of fast particle f on ambient ion i , E_f is the energy of fast particle f . The transport cross sections σ^T are gotten from angular averages over elastic scattering differential cross sections.⁴ These cross sections are for pure nuclear scattering; the Coulomb scattering cross sections have been subtracted out. We use the following rough estimates of transport cross sections in DT fuel: $\sigma^T = 0.39$ b for α particles from the DT reaction,⁴ $\sigma^T = 0.88$ b for protons from the DD reaction,⁵ $\sigma^T = 0.63$ b for tritons or ^3He from the DD reaction.⁶

Energy Loss to Bound Electrons. For this energy loss mechanism we use a quantum mechanical calculation of H. Bethe³, corrected for screening by free electrons, and a semiclassical prescription for mean binding energies.⁷ When we factor out an electron Debye length λ_e , we have for the energy

loss per unit length,

$$F_b = -\frac{e^2}{\lambda_e^2} Z_f^2 \left[\frac{Z_a - \bar{Z}}{\bar{Z}} \right] \ln \Lambda_b \frac{1}{2x_e^2} \quad (20)$$

where Z_f is the charge of the fast particle, Z_a is the atomic number of the ambient ion, $\bar{Z}$ is the average ionization. The average ionization $\bar{Z}$ comes from the equation of state model that we use.⁸ When the material is cold and the ionization $\bar{Z}$ is small, this loss term is important. In the numerical simulations we have a lower cutoff on $\bar{Z}$; to prevent it from being exactly zero.

The term $(1/2x_e^2)$ plays the role of the deposition function G in the general formula. The Coulomb logarithm Λ_b , however, differs from the ones used in previous sections. It is expressed in impact parameter form. The minimum impact parameter b_{min} is the de Broglie wavelength of the fast particle and an ambient electron in their center of mass system. The maximum impact parameter b_{max} is a complicated formula which takes screening into account. We state the results,

$$\left[\frac{b_{min}}{a_o} \right]^2 = \frac{E_{ryd}}{E}$$

$$E = \frac{m_e}{m_f} E_f + \frac{3}{2} \theta_e \quad (21)$$

$$\left[\frac{b_{max}}{a_o} \right]^2 = \max \left\{ \left[\frac{b_{min}}{a_o} \right]^2, \min \left[\left(\frac{R_{screen}}{a_o} \right)^2, \frac{4 E_{ryd} E_f}{I_{av}^2} \cdot \frac{m_e}{m_f} \right] \right\} \quad (22)$$

where I_{av} is an average ionization potential, and all other variables have been previously defined. The average ionization I_{av} is taken from a simple model of Zel'dovich and Raizer,⁹ and the ionization energy levels are taken from Hartree-Fock calculations.¹⁰

Energy Loss to Electron Plasma Oscillations. For this energy loss mechanism we use a result that is related to work of Bohm and Pines.³ When we factor out an electron Debye length λ_e , we have for the energy loss per unit length,

$$F_{plas} = -\frac{e^2}{\lambda_e^2} Z_f^2 \left(\frac{1}{4x_e^2} \right) \ln(1 + 4x_e^2) \quad (23)$$

where all the variables have been previously defined.

Summary of Results. The model of energy loss by a fast particle moving through a plasma has been described through five terms, of which four are electromagnetic processes and one is elastic nuclear scattering. All the electromagnetic energy loss terms have the same general form. They depend on the electron and ion temperatures and densities, most strongly through the electron and ion Debye lengths. The Coulomb logarithms are written in terms of ratios of impact parameters that have obvious physical meaning, namely, de Broglie wavelengths, Debye lengths, and atomic interparticle distances (cell sizes). But we must also rely on a model for an equation of state to define an average ionization charge and an average ionization potential.

In our numerical simulations, the principal equation is Equation (1). The five specific loss terms are given by Equations: (13), (16), (19), (20), and (23). The other equations define the subsidiary variables used in the model.

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Laser & Optics

The Chroma I Nd: Glass Laser
Advanced Laser Projects

4 Laser and Optics

4.1 The Chroma I Nd:Glass Laser

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Assembly of our 2 TW Nd:glass laser, Chroma I, in the new KMSF facility began in January, 1979, and was completed in early summer. The basic 2 TW system was designed in 1977 using a lumped element computer code which assumes small signal gain in the amplifiers, and accounts for nonlinear index induced self focusing by decreasing the transmission of spatial filters according to the ac-

cumulated B-integral.¹ John Trenholme and Alex Glass of the Lawrence Livermore National Laboratory assisted in this design.

The optical schematic drawing for the Chroma I laser system is shown in Figure 4.1-1. The system includes an actively mode locked oscillator,² rod amplifiers with a maximum diameter of 64 mm, nine 10 cm disc amplifiers and four 14 cm amplifiers. To protect the system from reflected energy, three pulsed Faraday rotators are used, one 70 mm and two 105 mm in diameter.

Spatial filters are used as shown to improve the intensity distribution of the beam and to retard the growth of small scale self focusing. Pinhole

Fig. 4.1-1 Chroma I 2 TW glass laser system.

diameters are approximately 15 times greater than the diffraction limit. The Hunt Relay Technique³ is employed to relay the image of a hard aperture from spatial filter to spatial filter, with the last image plane near the target chamber. This method permits a flat intensity profile and consequently a high fill factor to be propagated through the laser chain without the growth of high intensity Fresnel ripples.

Emphasis on spectral diversification continued, as reflected in the new name for our laser, Chroma I. Frequency doubling target experiments continued with 14 cm Type II KDP crystals placed after the last spatial filter and with the old lens ellipsoidal mirror target illumination system. The all reflective triple bounce illumination system (TBIS) was fabricated in 1979. A dual-frequency (0.53 μm and 1.05 μm), high power dielectric coating was developed and successfully applied to the TBIS.

The initial device used to broaden the laser bandwidth for experiments during 1978, the plasma spatial filter,⁴ did not meet our expectations operating with the Chroma I laser. Although the bandwidth could be increased from about 1 $\text{\AA}$ to 25 $\text{\AA}$, the focusable power on target decreased from 2 TW to 0.5 TW. Most of this decrease occurred as the broadened light traveled through the disc amplifier system which now includes five spatial filters per chain. Future research will be directed at alternate bandwidth broadening devices placed at the end of the laser.

The laser has been used in this configuration for over 200 target shots delivering up to 2 TW on target at 1.05 μm . Because the transmission of the current lens/ellipsoid target chamber optics is only 65 percent, 2 TW on target represents 3 TW out of the laser. At this high output power level, several elements were damaged, therefore maximum output has been limited to the safe operating level of 1.6 TW to 1.8 TW on target. Up to these power levels, the laser performs very linearly with good transmission through all the spatial filters. We expect routine operation at the 2 TW level on target after the all reflective TBIS is installed since it will have a one-way transmission of about 86 percent.

Various components of the Chroma I laser are described in more detail in the Sections which follow.

4.1.1 Active Mode Locked Oscillator (AMLO)

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AMLO Description. An actively mode locked oscillator using Nd:YLF was designed and built to provide a stable controlled input to the Chroma I phosphate laser system. The active mode locking technique² demonstrated much better stability characteristics than the passively mode locked oscillator and the Nd:YLF rod generated a wavelength of 1.05 μm , which is close to the Q-88 phosphate gain peak. The theory and computer simulations for this oscillator were given in the 1977 Annual Report.⁵

The resonant cavity, shown in Figure 4.1-2, was made with 4 parallel invar bars for thermal stability. The acousto optic mode locker is located as close to the 10 meter end mirror as possible. Both the mode locker and Q switch use lithium niobate transducers on fused quartz substrates with Brewster angle ends for the laser beam. The output mirror is mounted on a longitudinally adjustable slide and has a reflectivity of 78 percent. This high reflectivity is required

Fig. 4.1-2 Nd:YLF active mode locked oscillator and Pockels' cell pulse selector.

since the gain of Nd:YLF is less than one half that of Nd:YAG. The pump cavity is a single ellipse with a CW krypton arc lamp. The laser rod is 6.3 mm in diameter by 89 mm long with Brewster ends. The shutter S is used to block oscillation in the cavity while permitting the constant repetition rate of the lamp and acousto optic devices.

The single pulse selector consists of a pulse generator, two single crystal KDP Pockels' cells and three Glan-air polarizers. The pulse generator (Lasermetrics model 8601B) uses a krytron driven Blumlein which is triggered by a fiber optic signal that electrically isolates it from the other electronics and the optical table.

The electronics circuitry provides the repetition rate and control of the lamp driver, shutter, mode locker and "on" period of the Q switch. It also provides the timing for the laser amplifier flashlamps and the target diagnostics. All of these external timing pulses are transmitted through LED driven fiber optics to prevent electrical noise pickup.

The fast timing is done with ECL digital counters and other logic. This circuitry synchronizes the Q switch "off" period with the 64 MHz RF and provides the 32 MHz RF for the Q switch. Fast timing is also provided, through separate adjustable ECL counters to the Lasermetrics Pockels' cell drivers. The 64 MHz RF generator must be adjustable in frequency over at least a 500 kHz range, have very small phase jitter, and be stable within 10 ppm over

long periods. The HP8640B was the only signal generator which satisfied these requirements. The 20W RF amplifiers were designed to absorb reflected RF energy which would damage ordinary transistor amplifiers.

The lamp driver consists of a commercial SCR direct current power supply, capacitor bank, two current regulators and a 2 kV, 100 mA simmer supply.

Oscillator Operation. The lamp driver was designed to operate at repetition rates up to 10 pps, but good stability was obtained at 2 pps, which is adequate for alignment. Lamp life at 60 A, 2 pps, 10 hours per day, and 5 days per week is at least 3 months. The optical homogeneity of Nd:YLF laser rods is not as good as that of Nd:YAG rods, but a satisfactory TEM₀₀ mode can be obtained with careful alignment and orientation of the laser rods.

A photograph of the relaxation oscillation envelope is shown in Figure 4.1-3. The amplitude, pulse width, and spectral width stabilize during this period when the relaxation oscillations damp out. Figure 4.1-4 shows the mode locked train with the single pulse selected from its peak. This pulse is about 0.1mJ. The limited gain of Nd:YLF prevented a narrower pulse train. Cavity dumping would provide at least a factor of 10 more output energy, but we believe its complexity and susceptibility to damage would limit reliability.

Fig. 4.1-3 Relaxation envelope, time scale 0.5 ms/div.

Fig. 4.1-4 Mode locked train with single pulse removed.

4.1.2 Disc Amplifiers

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F.J. Masters

The Chroma I laser utilizes nine 10 cm and four 14 cm disc amplifiers to amplify the output of the 64 mm rod driver. These amplifiers are modifications of the old KMSF amplifiers,⁶ with design goals including higher gain coefficients and a high degree of cleanliness. The 10 cm amplifiers each contain three discs of 2.5 cm thick, 3.0 percent Nd doped Kigre Q-88 phosphate glass. These are pumped by ten water cooled flashlamps with reflectors extending around the entire back side of each lamp.

The overall gain at 1.05 μm of each of these 10 cm amplifiers under operating conditions in the system is 2.2 (9 percent/cm),⁷ which is a considerable improvement over the old amplifier gain of 1.37 with silicate glass.

The 14 cm amplifiers each contain two discs of 2.5 cm thick 3 percent Nd doped Q-88 phosphate glass pumped by twelve water cooled flashlamps. The system gain of the fourteen cm amplifiers at 1.05 μm is 1.7 (9 percent/cm),⁷ up from the unmodified gain of 1.26 with the silicate glass.

When the flashlamps are fired, a very large light flux (up to 30 J/cm²) is incident on the glass discs. Any dirt or contaminant present can be heated to

very high temperatures, damaging the disc surface. To reduce the problem of contamination and consequent damage to the faces of the discs, the glass discs and holders are carefully cleaned in a Class 100 clean room and then enclosed and sealed in a clean pyrex cylinder.⁸ This loaded cylinder is then inserted in the flashlamp cavity in the laser bay and sealed to the tube or element on each end, so that the inside environment of the clean cylinders is isolated from the outside laser bay.

The laser discs have a newly developed thick monolithic edge cladding⁹ which has a much higher damage resistance to the intense flashlamp radiation than the old frit type thin claddings. This eliminates one of the major sources of particulate contamination from inside the sealed pyrex tubes.

The flashlamps are water cooled, primarily to extract the residual heat from the flashlamp after firing. This reduces the temperature rise of the laser glass and consequently reduces the time required for cooling between firings to 30 minutes. The loss of pump light due to the flow tube and water jacket is negligible.

The energy transmission of the Chroma I laser is very linear as illustrated by the shot data in Figure 4.1-5, taken with various numbers of amplifiers fired. There is no evidence of nonlinear small scale self focusing.

The Chroma I power transfer function data are plotted in Figure 4.1-6 for all thirteen amplifiers on.

Fig. 4.1-5 Energy transfer function under different amplifier operating conditions.

Fig. 4.1-6 Chroma I power on target vs energy out of 45mm rod with all 13 amplifiers on (80 ps pulses).

All data include the 65 percent transmission loss due to the lens/ellipsoid target illumination system. The solid curves in this Figure represents the lumped element computer code predictions, where " ϵ " is the initial noise fraction in Trenholme's approximation to spatial filter transmission functions.¹⁰ Although the error bars are large, the Chroma I system appears to operate quite linearly at this level, with an apparent noise fraction averaging between 0.01 percent and 0.5 percent.

4.1.3 Faraday Rotators

R.G. Sanderson

KMSF has operated large (up to 18 cm clear aperture) cryogenic Faraday rotators in its laser systems for 5 years. These rotators utilized superconducting magnets which were cooled by liquid helium. Although these Faraday rotators were very reliable, the liquifier system was expensive to operate and required constant attention.

A detailed analysis of the Chroma I staging indicated only two 10 cm Faraday Rotators would be required to adequately protect the laser from back reflected energy. Because pulsed Faraday rotators of this size do not require large capacitor banks, and many of the components to build them were available as spare parts from our old laser driver system, the cryogenic Faraday rotators were replaced by the more maintenance free pulsed rotators.

The Faraday rotators were designed and built at KMSF. Two 1070 μH coils and high voltage rectifier stacks were purchased, and two spare high voltage chargers and capacitor banks were modified to drive the coils. The control electronics consist of a high voltage power supply, dual high voltage triggering circuit, and a dual system safety inhibit circuit. The rotators have performed as designed, generating a 30 kG magnetic field in a 1.8 ms pulse.

4.1.4 Plasma Spatial Filter

C.D. Kinzer

W.W. Lawrence

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N.K. Moncur

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The plasma spatial filter^{4,11,12} (PSF) has been used at KMSF as a device to produce a smooth time

Fig. 4.1-7 f/0.8 reflective plasma spatial filter.

averaged laser beam intensity distribution and to increase the laser bandwidth for our studies of target absorption. In previous KMSF laser systems the plasma filter was inserted between the 45 mm rod and 64 mm rod. In this position, power into the PSF was limited to about 10 GW and the reflected light collected with the f/3.8 illumination system was less than 5 percent of the input energy. In the Chroma I laser a much faster f/0.8 plasma spatial filter was designed and placed after the 64 mm rod (see Figure 4.1-7). This allowed us to operate at higher powers (up to 40 GW) into the filter and to collect up to 40 percent of the incident light.

Fig. 4.1-8 Chroma I disc amplifier gain with and without the PSF.

The plasma spatial filter gave a smooth time averaged intensity distribution and a broad bandwidth as it had previously, but the beam did not propagate well at high intensity through the laser chain. In an attempt to improve the performance of the laser with the PSF, a slower (f/3.2) refracting illumination system was installed in the PSF, but the propagation did not improve. Figure 4.1-8 shows a comparison of the system performance with and without the PSF.

The performance of the f/3.2 PSF was measured in detail with a series of 25 full power shots. The primary diagnostics included energy measurements along the laser chain, spectral measurements of the PSF output and of the Chroma I laser output, and far field measurements of the laser output.

PSF Energy Propagation Measurements. The transmission of energy through the Chroma I laser was measured in stages, using single 80 ps pulses. The energy propagation through the first three 16 mm and the 45 mm rod amplifiers was very linear. The energy at "D" was varied between 0.02 and 0.2 J, which corresponds to about 0.24 to 2.4 J at the

PSF. The 193 mm focal length PSF lens focuses this energy on a glass target.

The effective efficiency of the PSF was inferred from the ratio of the "E" to the "D" photodiodes. The "E" photodiode measures the light energy reflected by the PSF and amplified by the 64 mm amplifier as the light passes back through this amplifier. The polarizer excludes any reflected light which is depolarized.

The laser energy next passes through a conventional vacuum spatial filter, SF₄. This spatial filter has a 1 mm pinhole and a 1520 mm focal length input lens, resulting in a 658 μ rad full acceptance angle. The transmission of this spatial filter averaged about 80 percent, as shown by the circle data points in Figure 4.1-9. This transmission did not change significantly as a function of energy. Doubling the pinhole size to 2 mm had little effect on transmission, indicating that there is negligible energy in the 658 to 1316 μ rad range.

A saturable dye cell was placed in front of SF₄ for many of these tests, since this is used for most target experiments to attenuate early light. This dye cell (containing a mixture of Eastman Kodak 9740 Q

Fig. 4.1-9 Transmission of pinhole spatial filter after the PSF (with and without the saturable dye cell).

switch dye dissolved in 1,2-dichloroethane) transmits 0.5 percent at low intensity. The "x's" of Figure 4.1-9 show that the high intensity transmission varies from about 30 percent to 50 percent as the input energy is varied from 0.4 to 1.8 J.

The small signal gain of the next stage, which includes three 10 cm disc amplifiers, is about 5.1. The measured gain with the PSF was actually higher than the gain without the PSF. This anomaly is probably due to the polarization characteristics of the anti-reflection coated beam splitter which feeds the "F" photodiode; without the PSF, the laser beam has some depolarized light due to the birefringence of the 64 mm laser rod, whereas, in the PSF configuration, the light reaching photodiode "F" is well polarized, since it was reflected from a polarizer before entering SF₄.

The small signal gain of the second disc amplifier stage is about 10.4. The measured gain with the PSF was only 6 to 7, falling 15 or 20 percent to the 5.5 level with increased drive energy.

The gain of the third disc amplifier stage is plotted in Figure 4.1-10. The pinhole was removed from the last three spatial filters (SF₄, SF₈, & SF₁₀) for most of the PSF shots, so that the far field camera could measure the laser beam divergence beyond the pinhole cut-off. (A 328 μ rad pinhole was installed for all non-PSF shots.)

Fig. 4.1-10 Gain of third disc amplifier stage.

PSF Spectral Measurements. The spectrum of the light reflected from the PSF is superimposed on the nearly monochromatic output at 1.0530 μ m from the YLF oscillator in Figure 4.1-11. A CW YAG exposure at 1.0642 μ m is used to calibrate the wavelength scale. For this particular shot, the plasma filter output bandwidth was 21.3 $\text{\AA}$ (FWHM). A linearized, smoothed, and normalized PSF spectrum is compared with the Kigre Q-88 phosphate laser glass gain curve in Figure 4.1-12. The PSF output peaks to the blue side of the Q-88 gain by about 13 $\text{\AA}$. Multiplication of these two curves would predict an output bandwidth somewhat narrower (31.8 $\text{\AA}$ vs. 35.2 $\text{\AA}$), and a peak shifted to the red to about 1.0532 μ m. (The actual measured Chroma I output for this shot was 25.7 $\text{\AA}$ with a peak at 1.0535 μ m.)

One might expect a variation of bandwidth or a frequency shift as the input energy to the PSF is varied. If the blue shifted light is due to a doppler shift from the expanding plasma, then higher intensity might produce a more rapidly expanding plasma and more blue shift. Three characteristics of the PSF spectra are plotted in Figure 4.1-13 as a function of input energy: the center wavelength (midway between the -3dB points), the peak wavelength (always toward the blue of center), and the -3 dB blue edge. While the data scatter is large,

Fig. 4.1-11 Plasma spatial filter spectrum.

Fig. 4.1-12 Comparison of Q-88 gain curve with PSF output spectrum.

Fig. 4.1-13 Variation of PSF spectra with input energy.

there is a small trend of increased blue shift of the blue edge at higher energy.

The spectral characteristics of the high power light exiting the Chroma I laser are shown in Figure 4.1-14. The most striking feature of this entire experiment is the effect of the saturable dye cell on the output bandwidth and center wavelength. The saturable dye cell broadens the average Chroma I output from $12.1 \text{ Å} \pm 2.6 \text{ Å}$ to $29.2 \text{ Å} \pm 6 \text{ Å}$. Apparently nonlinear processes in the saturable dye convert some energy to longer wavelengths, making it a laser pumped dye laser.

For eight out of the nine shots with the saturable dye cell in place, the Chroma I output bandwidth was actually larger than the input bandwidth to the disc amplifiers, despite the bandwidth narrowing effect of the Q-88 laser glass. The ratio of the output bandwidth to the input bandwidth (which is the PSF output) is plotted in Figure 4.1-15. In all cases

Fig. 4.1-14 Laser bandwidth vs center wavelength with the plasma spatial filter.

Fig. 4.1-15 Bandwidth enhancement with a saturable dye cell.

the bandwidth was narrowed when the dye cell was removed. This mechanism will be investigated further as a possible bandwidth broadening device.

PSF Far Field Measurements. The PSF effectively degrades both the temporal and the spatial coherence of the laser beam. The scattered light from the turbulent plasma has a wider far field pattern than the non-PSF light, as shown on a log scale in Figure 4.1-16 for two 100 J shots.

The size of each far field pattern was measured by integrating the linearized irradiance patterns recorded on photographic film. From each integration we extract those far field diameters which would enclose 50 percent, 80 percent and 90 percent of the total energy. Two orthogonal scans are analyzed for each shot. The results are plotted in Figure 4.1-17 as a function of the energy per beam out of the laser.

The far field spot diameters grow with output energy. In general, the addition of the saturable dye cell causes the far field patterns to reach a given diameter at much lower energy levels than without the dye cell.

Fig. 4.1-16 Far field patterns with and without the plasma spatial filter.

Fig. 4.1-17 Far field enclosed energy diameters for Chroma I with the plasma spatial filter.

Doubling Efficiency with the PSF. The output light from the Chroma I laser was frequency doubled by a 14 cm KDP crystal. The doubled green energy is plotted as a function of the incident IR energy in Figure 4.1-18. The average doubling efficiency with the PSF was 31.3 percent. For comparison, the doubling efficiency without the PSF exceeds 60 per-

cent for input irradiance above 4 GW/cm² (See section 4.1.6). The PSF decreases doubling efficiency a factor of two. This decrease may be due in part to the increased bandwidth of the PSF light. However, Figure 4.1-19 indicates very little dependence of doubling efficiency on bandwidth variations from 10 Å to 40 Å.

Fig. 4.1-18 KDP doubling efficiency with the PSF.

Fig. 4.1-19 Doubling efficiency vs laser bandwidth (with the plasma spatial filter).

4.1.5 The Triple Bounce Illumination System (TBIS)

G.C. Catella
C.J. Edmonson

R.D. Sigler

Introduction. One of the crucial problems in laser fusion research is the efficient focusing of enormous amounts of power (typically several terawatts) onto targets $\sim 100\text{ }\mu\text{m}$ diameter. With two beam laser systems such as Chroma I, the optical requirements are especially severe, because uniform illumination is considered essential to produce stable, high density implosions. An illumination system consisting of a pair of ellipsoidal mirrors and aspheric, single element lenses has been successfully used for several years at KMSF.^{13,14} There are however, some serious drawbacks to this system. The lens surfaces are so steep that fabrication of anti-reflective coatings for both polarizations is very difficult, and the uncoated lens transmission is only 80 percent, averaged over the beam. Lenses must be very thick to accommodate short radii of curvature. This results in the largest contribution to the beam breakup integral (B-integral) of any passive component in the entire laser system.^{15,16} Also the lenses are stigmatic for only one wavelength. Experimental evidence demonstrates that target absorption is increased when irradiated with harmonically generated short wavelength laser light,^{17,18} suggesting that a system which is achromatic for two or more wavelengths would be desirable for comparing target response at different wavelengths with the same illumination system. Recent experience has indicated that target absorption may be increased by illumination with broadband laser light; the best coupling may require bandwidths which exceed the chromatic focusing capabilities of conventional refractive illumination systems.

The following guidelines were considered in the development of the KMSF all reflective illumination system. Since our laser is configured for two-beam experiments, the illumination system must have bilateral symmetry. The target full cone angle must approach 180° on each side. The total angle achievable must be limited only by the need to provide an access ring for diagnostics and target insertion. While area obscuration of the central region

of the entrance pupil is permissible, it should be held to a minimum and in no case be larger than 4 percent. The target must be isolated from laser energy crossing the mirror cavity prior to reaching the final focus or from any other ray paths which would allow the energy to disturb the target prematurely. The target surface irradiance must be as normal and uniform as possible. No element of the illumination system should be exposed to power densities significantly greater than that found within the collimated input beam. The RMS spot diameter must be as small as possible (e.g., near diffraction limited) over a spectrum encompassing the primary, doubled, and tripled frequencies of neodymium glass lasers ($1.05\text{ }\mu\text{m}$, $0.53\text{ }\mu\text{m}$ and $0.35\text{ }\mu\text{m}$ wavelengths). Finally, the thickness of refractive components must be held to an absolute minimum to minimize the B-integral contribution.

Numerous illumination systems were investigated, but the only system which satisfied the above requirements, as well as being within the current state of fabrication art, was the triple bounce system (shown schematically in Figure 4.1-20)

Fig. 4.1-20 The triple bounce illumination system.

which was originally designed in 1977.¹⁹ Unlike many other systems studied, all elements of the triple bounce system are axial and require comparatively mild asphericity. For complete 360° target coverage (180° on a side), the parabola need only have a speed of $f/1.4$. To provide adequate

diagnostics area, the actual design has an $f/1.8$ parabola which results in a target cone angle of 150° on a side. In this system, at least one of the surfaces must have relatively weak higher order aspheric terms. We have chosen to locate this aspheric surface on the target vacuum chamber window.

Because the effective optical path length has been reduced from 117 mm in BK-7 to 25 mm in the quartz vacuum chamber window, deleterious non-linear index effects of high power operation have been reduced by a factor of about 7 over the current lens mirror system. The transmission to the target is increased from 65 to 85 percent since, unlike the current system, all surfaces in the triple bounce system can be effectively coated.

The aspheric surface on the corrector plate was designed to produce focal spots of less than $10\text{ }\mu\text{m}$ in diameter at both $1.05\text{ }\mu\text{m}$ and $0.53\text{ }\mu\text{m}$. The irradiance uniformity over a sphere was determined from 3000 ray traces to be better than 1.2 to 1.¹⁹

Notice from Figure 4.1-20 that the laser beam forms two foci on either side of the target prior to the final implosion. In principle these early foci could disturb the target. However, because of the on-axis entrance holes in the cavity mirrors, paraxial rays are blocked (4 percent of the total area) and a small dead volume surrounds the target through which no rays will pass until the final (target) focus is reached. The dead volume is elliptically shaped with a minor axis of about 6 mm. Experiments at high power with the current mirror lens system and the target displaced from the focal volume to simulate this condition have demonstrated there is no target disturbance from early energy.

As a consequence of the corrector and cavity mirrors having less asphericity than the current aspheric lens and ellipsoidal mirrors, fabrication difficulties are comparable with or even less than those experienced with the current system. Fabrication of the parabola, with its speed of $f/1.8$, is well within the state of the art.

The conic cavity mirrors require only 25 percent as much volume to be removed from the best fitting sphere to achieve the desired conic section as do the current ellipsoids. Although these triple bounce cavity mirrors are conic sections, the eccentricity is so small that the two conjugate foci are only about 5 cm apart. This is advantageous because deep

reflectors are easier to fabricate when they are nearly spheres.

Testing these deep reflectors is more difficult however, because the cone angles required to fill the mirror from either focus are both quite large, so large in fact that there is no known off the shelf lens with the required qualifications. To solve this problem, we have designed a three element, all-spherical, expanding lens which can convert the 28° cone angle from a commercial lens (diffraction limited), into a stigmatic 122° cone which will fill the conic cavity mirror to the required diameter. A cavity mirror test system with the expanding lens is shown schematically in Figure 4.1-21. This diverging lens has a precision test ball permanently affixed to its vertex to permit autocollimation testing. With this arrangement, both cavity mirrors are made identical.

Fig. 4.1-21 Test set up for conic cavity mirror.

The elements of the expanding lens are aplanatic spheres and consequently the outer surface of the lens is concentric with the emerging wavefront. If left uncoated, this surface will produce a spherical reference wave (by reflection) which can be compared with the wave returning from the cavity mirror and test ball in a spherical wave interferometer.

The complete interferometer and cavity mirror test set up is shown in Figure 4.1-22.

Fig. 4.1-22 Spherical wave interferometer

Fabrication and Testing of the TBIS Elements. Fused silica was chosen for the corrector material based on its superior mechanical characteristics. The desired profile is similar to a classical Schmidt corrector plate which departs from the best fit sphere by 732 fringes (at $0.63 \mu\text{m}$).

The aspheric surface of the corrector was generated with a computer directed 3 axis interferometric controlled grinding machine shown in Figure 4.1-23. Fine grinding was performed by

Fig. 4.1-23 Computer directed grinding machine.

hand with subdiameter flexible tools and a spindle. Ring gauges of various diameters with 20 microinch indicating gauges were used through both fine grinding and polishing stages. Polishing was also done by hand on a spindle using flexible laps. Final measurement of the correctors was performed with a 3 axis interferometrically controlled measuring machine which has an accuracy of several microinches. Both correctors, when completed, were within 20 microinches of the theoretical profile. A double pass transmission interferogram (at $0.63 \mu\text{m}$) (Figure 4.1-24) shows the null or inflection point to be uniform to within a few fringes, indicating the part is relatively free of astigmatism.

Fig. 4.1-24 Corrector plate interferogram.

The CER-VIT parabolas were generated, fined and polished out in much the same manner as the correctors with the exception of the flexible tools. These parabolas, even though quite fast at $f/1.8$, were finished to very high accuracy ($\lambda/6$ peak to valley) as shown in Figure 4.1-25.

Fabrication of the cavity mirrors was much more difficult. The major excavation was performed by sand blasting and halted when within a half inch or so of the desired curve. Final rough grinding was

Fig. 4.1-25 Parabola interferogram

performed with the computer controlled grinder described earlier. Many passes were required to reduce the sand blasted curve to the proper profile. Fine grinding was again performed by hand, with frequent checks using the ring gauge and, periodically, the 3 axis measuring machine. Polishing was also largely done by hand and the radial length of any strokes had to be kept short, as the radius of curvature changed rapidly in this direction. In the early stages of figuring, the errors were too large to be effectively evaluated with the spherical wave interferometer and a special knife edge test fixture was used. This test fixture sat directly on the work piece and used a high numerical aperture microscope objective and laser source to illuminate a portion of the ellipsoidal cavity from the outer focus. Since only a portion of the ellipsoid was illuminated (from the outer edge to just past the mirror vertex) the returning beam was passed through the inner focus in such a fashion as to be accessible for both the knife edge and viewing. The knife edge test fixture on the cavity mirror is shown in Figure 4.1-26. With this instrument the cavity was hand figured until a relatively smooth cutoff was achiev-

Fig. 4.1-26 Knife edge test fixture on cavity mirror.

ed at all positions around the mirror. At this point testing was subsequently performed on the spherical wave interferometer. When first put on the interferometer after knife edge testing, errors of up to 20 fringes (double pass) and considerable astigmatism were visible. Weeks of careful hand figuring were required to reduce the errors to a total of approximately 3 fringes, as shown in Figure 4.1-27. Since these fringes are generated in a double pass mode, the actual surface error is about $3/4 \lambda$ at $0.63 \mu\text{m}$.

Fig. 4.1-27 Cavity mirror interferogram.

Mounting and Alignment of the TBIS Optical Elements. Computer studies of alignment sensitivities of the triple bounce system revealed that, while it is no more critical than that experienced with the lens mirror system, the tolerances are nevertheless measured in microns and fractions of an arc minute. Experience with the lens mirror system demonstrated that an adjustable mounting is necessary and must operate inside a vacuum chamber, this led to providing each mirror motion with position encoders. The corrector motions are not motor driven, as they are relatively insensitive and their controls are located on the target chamber exterior.

The cavity mirrors require six degrees of freedom and the parabolas require five. Meeting all these requirements resulted in the rather sophisticated element mounting assembly shown in Figure 4.1-28,

which has a total of 24 motors and encoders. Each motion is controlled by a 0.5 to 5 RPM D.C. motor which is driven with a square wave signal of variable duty cycle. The motor is connected to a 40 thread per inch lead screw. Motions as small as 1 micron are easily made and measured with the encoding system. All motions of the elements within the mount are constrained by either preloaded ball bearings or crossed roller linear bearings.

System Testing. System testing was performed with all mirror surfaces chemically silvered. Additional detailed system testing was performed after all the elements were given their final dielectric coatings. A schematic drawing of this test set-up is shown in Figure 4.1-29 and a photograph of the system is shown in Figure 4.1-30. The entire system is isolated on an air table so that an interferometer

Fig. 4.1-28 Cavity mirror mounting assembly.

Fig. 4.1-29 System test set-up.

Fig. 4.1-30 Photograph of final dielectric coating test set-up.

can be installed in the input beam to provide a reference wave and allow path length changes such that the total path length difference is an even multiple of the laser cavity length.

It was originally planned that localized hand figuring would be performed on the corrector if the system did not meet the encircled energy requirement at the target position. Measurements indicated the system had an 80 percent encircled energy diameter of 10 μm (at $\lambda = 0.063 \mu\text{m}$). The theoretical performance indicates an RMS spot

diameter of 7 μm at this wavelength. Because the relayed image path is very long (> 4 meters) there was some concern that strongly aberrated rays would not properly be relayed and therefore give an overly optimistic estimate of the spot size. After the elements were dielectrically coated, attempts were made to examine the system focus directly. The very large cone angle of the focus made direct measurement difficult.

While no known optical system has sufficient numerical aperture to collect the rays at the system focus (without also obscuring much of the aperture), a microscope objective system was designed to permit some evaluation of this focus by relaying the image of the focus to a small screen. This system uses a 3X microscope objective to relay the spot image and has just barely enough resolution to measure the system focus at 7 μm . However, this system did confirm that there was a "halo" of poorly focused energy around the focus which had not been measured in earlier spot relay testing.

The spot size could be estimated more accurately by measuring the amount of energy transmitted through pinholes of known diameter. However, because of the large cone angle, the pinhole substrate thickness must be less than 25 percent of the hole diameter just to pass rays from a perfect focus. This would require pinholes only a few microns thick. Substituting a knife edge for a pinhole and tabulating the amount of energy blocked versus the knife edge position, the spot size can

Fig. 4.1-31 System focus measurements.

be determined analytically. In the data reduction scheme, the spot is considered to be composed of radially symmetric annuli. Solving for the intensity of each annulus produced the relative intensity versus radius plot in Figure 4.1-31. Integrating this plot (assuming circular symmetry) yields the encircled energy versus radius plot also given in Figure 4.1-31. This technique yields an 80 percent encircled energy diameter of 22 microns at $\lambda = 1.05 \mu\text{m}$, and 35 microns at $0.53 \mu\text{m}$. These diameters are larger than those indicated by the relay method but are acceptable since typical laser fusion targets in use today are the order of $100 \mu\text{m}$.

Dual Frequency Coating Development. A new coating development program was initiated to provide the dual frequency coatings for the TBIS. These coatings will permit operation of the TBIS at the fundamental wavelength of the Nd:glass laser ($1.05 \mu\text{m}$), and at the second harmonic green wavelength ($0.53 \mu\text{m}$). During 1979 the first TBIS was coated with the dual frequency (IR and green) coating.

The following performance goals were chosen for the dual frequency coating:

1. Reflectivity greater than 98 percent at both

wavelengths for mirrors, and anti-reflection coatings with less than 0.5 percent reflectivity per surface at each wavelength.

2. Damage thresholds —

- (a) $> 15 \text{ GW}/\text{cm}^2$ average for 250 psec at $1.05 \mu\text{m}$
- (b) $> 12 \text{ GW}/\text{cm}^2$ average for 250 psec at $0.53 \mu\text{m}$
- (c) $> 15 \text{ GW}/\text{cm}^2$ average for 250 psec at $1.05 \mu\text{m}$ and $0.53 \mu\text{m}$ simultaneously, with the power divided equally between both wavelengths.

Coating Development Program. Development contracts were let with two vendors to provide coating designs meeting the performance goals. Working independently, each vendor supplied samples for evaluation of reflectivity and damage performance. From these results, new samples were produced and evaluated. KMSF then selected on the basis of past experience, performance, cost and promised delivery one vendor for production of the TBIS.

Damage Test Facility. A damage test facility was constructed to allow routine testing of optical components, including the dual frequency test coatings and the TBIS. Adequate repeatability and ease of operation resulted from this standardized equipment. The basic facility is outlined in Figure 4.1-32.

Fig. 4.1-32 KMSF damage testing facility.

To damage components successfully, approximately 3 J, delivered in a 250-300 psec FWHM pulse is required. An approximately 300 psec, square wave envelope is produced by the rod amplifier section of the Chroma I laser. Frequency doubled light is provided by a KDP doubling crystal. Doubling efficiencies are typically 40-50 percent.

Results of the Coating Development Program. Parts provided by both contractors were successful in meeting the KMSF reflection and damage performance goals. Figures 4.1-33 and 4.1-34 outline the experimental results. Each datum represents a different coating design.

Fig. 4.1-33 Dual frequency mirror damage thresholds.

Fig. 4.1-34 Dual frequency AR damage thresholds.

Coherent Radiation had two design approaches for mirrors, "conventional" and "optimized." The "optimized" approach was an effort to reduce the maximum E-field present within the coating stack. The results indicate that these designs had a higher damage threshold than the "conventional" designs.

For subnanosecond pulses, the power damage thresholds of various dielectric coatings exhibit an inverse dependence on the pulse width. The most common model for this dependence, at 1.05 μm, is: Damage Threshold (GW/cm²) $\propto t^{-1/2}$.

Four of the dual frequency mirrors were tested at ~70 psec and 300 psec. Figure 4.1-35 shows the results along with the slope of the $t^{-1/2}$ model.

Fig. 4.1-35 Damage threshold vs time (0.53 μm).

Design Optics was chosen to coat the first TBIS. To test the uniformity of reflection over the full aperture, the vendor prepared fixturing which holds fourteen samples aligned to simulate the ellipsoidal surface. Trial coating runs were conducted and the samples spectroscopically analyzed. Variation of the coating parameters produced a 35 nm wide reflection peak at 0.53 μm and a 70 nm wide reflection peak at 1.05 μm. The actual ellipsoidal mirrors were then coated.

One problem was leaching or staining of small areas of the ellipsoidal and flat mirror surfaces prior to coating. This was predominantly a side effect of the silvering process required during the testing and manufacture of the parts. These

stains could not be removed by any common solvent. The flat surfaces on the back of the ellipsoids were coated first so that the second heat cycle, required to coat the ellipsoidal surface, would affect the *flat* surface, not the more critical *ellipsoidal* surface. A second problem arose directly from this dual heat cycle. Crazeing occurred on the flat surfaces when the ellipsoidal surfaces were coated. Except for the stained areas, reflectivity exceeded 98 percent at both wavelengths. This performance was maintained even in the difficult steep section of the ellipsoidal surface. The stained areas reflected about 95 percent of the incident radiation at green, and about 97 percent at IR. The AR coating reflected less than 0.5 percent at both wavelengths. Within the ± 1 percent precision of our measurements, the crazeing present on the flats did not lower the reflection of the dielectric coating.

A worst case damage test was made on the flat, crazeed, mirror surface. Our performance goal was 15 GW/cm^2 . At $1.05 \mu\text{m}$ the damage threshold was $15 \pm 4 \text{ GW/cm}^2$; at $0.53 \mu\text{m}$, $10 \pm 5 \text{ GW/cm}^2$ was measured. The value for $1.05 \mu\text{m}$ meets the performance goal, while the $0.53 \mu\text{m}$ value does not. However, they are sufficiently high for us to operate the TBIS with the Chroma I laser.

4.1.6 Frequency Doubling

L.D. Siebert

Considerable effort has been expended to optimize the frequency doubling efficiency of our large aperture cells. The basic theory was described in the 1977 KMSF Annual Report.²⁰

The experimental results from the two 14 cm aperture Type II KDP crystals (0.9 cm thick) used for two beam green light target experiments are shown in Figure 4.1-36 for each crystal. These data points are for single Gaussian laser pulses with nominal 100 ps widths. The solid curves are the theoretical efficiencies for phase mismatch parameters DK , from 0.0 to 5.0 rad/cm. In Figure 4.1-36 (a) and (b), the laser spectral linewidth was less than $1 \text{ \AA}$ and the laser output passed through a spatial filter with a 1.5 mm diameter pinhole and a focal length of 4.57 m (angular bandpass of $\pm 0.16 \text{ mrad}$). Neither the spectral linewidth nor the spatial angular width produces a significant phase

Fig. 4.1-36 (a) Doubling efficiency for the 14 cm diameter 0.9 cm thick KDP II crystal #1 used in beam A. (b) Doubling efficiency for the 14 cm diameter 0.9 cm thick KDP II crystal #2 used in beam B.

mismatch (i.e. $DK \ll 1$).

The doubling efficiency for the spectrally broadened laser output (from the plasma filter) is shown in Figure 4.1-37. The spectral broadening accounts for a DK of less than 2 rad/cm; however, spatial and temporal amplitude spikes, which cause large power density excursions, probably account for the lower doubling efficiency.

Fig. 4.1-37 Doubling efficiency with the plasma spatial filter.

These results demonstrate that doubling efficiencies of 60 percent to 70 percent can be obtained with a relatively clean laser beam. No damage was observed on either the crystals or the windows in these tests. Hence, higher efficiencies appear possible at higher power densities.

4.1.7 Output Beam Characterization

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The characteristics of the Chroma I laser beam have been measured under four separate operating conditions: normal IR operation, IR with the plasma spatial filter (PSF), normal frequency doubled green operation, and green operation with the PSF.

The near field image for Chroma I shot #4188 is shown in Figure 4.1-38, along with a linearized scan through this image. The power on target was 1.48 TW. One measure of beam quality is the ratio of the peak to the average irradiance, which was 2.81 to 1 for this shot. The dashed lines on Figure 4.1-38 also indicate the computer calculated standard deviation about the mean value. The ratio of standard deviation to the average value was 0.56 for this shot. A plot of this peak to average ratio is shown in Figure 4.1-39 as a function of laser power on target, which

Fig. 4.1-38 Near field photograph for Shot 4188 (1.48 TW).

Fig. 4.1-39 Chroma I laser beam quality (near-field peak to average).

SHOT # 4167 P = 4
0.62 TW
PEAK-TO-AVERAGE = 1.51

SHOT # 4168 P = 5
0.56 TW
PEAK-TO-AVERAGE = 1.57

SHOT # 4169 P = 8
1.9 TW
PEAK-TO-AVERAGE = 1.48

SHOT # 4189 P = 8
1.8 TW
PEAK-TO-AVERAGE = 1.97

Fig. 4.1-40 Near field scans for Chroma I shots.

shows a slight trend toward higher modulation with increased power. The mean value of this near field peak to average ratio taken over 19 shots was 1.87 ± 0.35 .

Several representative near field scans for some early Chroma I shots are shown in Figure 4.1-40. The dashed lines indicate a supergaussian fit to the data of the form $e^{-(r/r_0)^p}$. The average beam diameter (FWHM) for this series was only 10.8 cm,

compared to the 14 cm clear aperture. The effective fill factor was about 0.6, or only 60 percent of the available beam area was effectively utilized.

The equivalent plane modulation was particularly large for these early shots. A far field camera photograph from shot #4188 (1.48 TW on target) is shown in Figure 4.1-41, along with a linearized scan taken at an equivalent longitudinal plane 40 μm from focus. The peak to average modulation is 3.92 to 1.

Fig. 4.1-41 Far field camera data for Shot 4188 (1.48 TW).

Fig. 4.1-42 Initial Chroma I equivalent plane distributions.

Fig. 4.1-43 Chroma I laser beam quality vs laser power (equivalent plane peak-to-average).

The ratio of standard deviation to average was 0.81. Figure 4.1-42 shows several other linearized equivalent plane scans, and Figure 4.1-43 shows the equivalent plane peak to average modulation as a function of power on target. The increased modulation of equivalent plane data relative to near field images is probably due to phase modulation on the laser beam. In particular, a hot outer ring characteristic of longitudinal spherical aberration appeared on most shots. The passive spherical aberration would not account for this ring. Apparently, the beam taper, combined with whole beam self focusing at high power, caused most of this ring structure.

The laser beam for these early Chroma I shots

SHOT # 4239 P = 26 PEAK-TO-AVERAGE = 1.69
3 db DIAMETER = 13.02 cm

SHOT # 4240 P = 29 PEAK-TO-AVERAGE = 1.90
3 db DIAMETER = 13.1 cm

SHOT # 4241 P = 16 PEAK-TO-AVERAGE = 1.75
3 db DIAMETER = 12.4 cm

SHOT # 4242 P ~ 40 PEAK-TO-AVERAGE = 1.56
3 db DIAMETER = 13.56 cm

Fig. 4.1-44 Representative Chroma I near field patterns after beam expansion.

SHOT # 4298 0.57 TW
PEAK-TO-AVERAGE = 4.75

SHOT # 4312 1.1 TW
PEAK-TO-AVERAGE = 4.05

SHOT # 4345 0.9 TW
PEAK-TO-AVERAGE = 2.38

SHOT # 4335 0.66 TW
PEAK-TO-AVERAGE = 1.88

Fig. 4.1-45 Equivalent plane distributions after beam expansion.

was formed by placing a hard aperture on the null of an Airy disc pattern early in the rod amplifier chain. To increase the fill factor, the beam was enlarged in front of this hard aperture. The image relaying property of the system permits this larger fill factor without the growth of deleterious diffraction rings. The fill factor at the end of the laser was increased from 0.60 to about 0.87.

The near field patterns are now much closer to a

flat top, as illustrated in Figure 4.1-44. The equivalent plane modulation decreased from a mean peak to average of 7.11 ± 4.78 down to a mean of 2.5 ± 0.81 . As shown in Figure 4.1-45, the hot outer ring, characteristic of the previous equivalent images (see Figure 4.1-42) has been significantly attenuated by the new laser beam profile.

Conversion to the frequency doubled green

Fig. 4.1-47 Green beam equivalent plane uniformity.

Fig. 4.1-46 Green near field data for shot 4290 (0.84 TW of green).

KEY: x = 90% ENCLOSED ENERGY DIAMETER
o = 80% " " "
Δ = 50% " " "

Fig. 4.1-48 Green focal spot diameters.

wavelength ($0.53 \mu\text{m}$) does not significantly change the near field data for shot #4290, which delivered 0.84 TW to the target. The peak to average ripple was 1.87. (The fringes on Figure 4.1-46 were generated by a $0.53 \mu\text{m}$ bandpass interference filter placed in front of the near field camera.)

The green equivalent plane and focal spot characteristics are shown in Figures 4.1-47 and 4.1-48. There is no indication of decreased quality with increased green power on target suggesting that small scale self focusing is not dominant at these power levels.

The green beam characteristics are compared with the average IR beam characteristics in Table 4.1-1. The near field and equivalent plane beam qualities are quite similar. The green focal spot sizes are significantly larger than the measured IR enclosed energy diameters. The passive aberrations in the doubling crystal cell do not account for this increase in focal spot size.

Table 4.1-1
Comparison of Green
and IR Beam Characteristics

	IR ($1.05 \mu\text{m}$)	Green ($0.53 \mu\text{m}$)
Near-Field Ripple*	$1.57 \pm .46$	$1.67 \pm .22$
Equivalent-Plane Ripple*	$2.50 \pm .81$	$2.51 \pm .47$
Far-Field Diameter, 90% enclosed energy	$9.16 \mu\text{m} \pm 2.9 \mu\text{m}$ ($161 \mu\text{rad}$)	$15.5 \mu\text{m} \pm 2.8 \mu\text{m}$ ($272 \mu\text{rad}$)

*Peak-to-average

The addition of the plasma spatial filter (PSF) smooths the time averaged, low frequency rings in laser beam profiles, as shown in Figure 4.1-49,

(a)

(b)

Fig. 4.1-49 Green near field image with the plasma spatial filter for Shot 4311.13 (0.176 TW of green).

Fig. 4.1-50 Far field camera data with the plasma spatial filter for Shot 4311.13 (0.52 TW of IR).

which is a green near field image with the PSF. (Compare with Figure 4.1-46, a non PSF green near field.) The hole in the center of the near field PSF image was generated by a central obscuration on the PSF lens; this obscuration is required to block the on axis pre-pulse reflection from the PSF lens. No extra energy is lost, however, since this central area is also blocked by the target illumination system. In

general, the time averaged peak to average spatial noise away from this hole is slightly higher than for non-PSF shots (1.79 vs. 1.67). However, the instantaneous patterns are certainly much noisier since the far field patterns show considerably more wide angle flare (see Figure 4.1-16), which translates to high spatial frequency noise in the near field.

A far field camera photograph for shot #4311.13 with the PSF is shown in Figure 4.1-50, along with a linearized scan through the equivalent plane (40 μm from focus in target coordinates). Again, the low spatial frequency ripple in the time averaged pattern is smoother with the PSF.

The average characteristics of the PSF beam are compared with the non-PSF features in Table 4.1-2.

Table 4.1-2
PSF and non-PSF Beam Characteristics

	Normal Operation	PSF Operation
IR Data:		
Near Field Ripple*	$1.57 \pm .46$	NA
Equivalent Plane Ripple*	$2.50 \pm .81$	$1.62 \pm .19$
90% Far Field Diameter	$9.16 \mu\text{m} \pm 2.9 \mu\text{m}$	$22.1 \mu\text{m} \pm 7.3 \mu\text{m}$
Green Data:		
Near Field Ripple*	$1.67 \pm .22$	$1.79 \pm .20$
Equivalent Plane Ripple*	$2.51 \pm .47$	
90% Far Field Diameter	$15.5 \mu\text{m} \pm 2.8 \mu\text{m}$	$42.7 \mu\text{m} \pm 19 \mu\text{m}$

*Peak-to-average

4.2

Advanced Laser Projects

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Theoretical investigation of two advanced laser driver systems continued during 1979. As described in the KMSF 1978 Annual Report, the chemically pumped iodine laser (CPIL) is unique in that it derives all of its energy from a chemical reaction.¹ Pulsed power supplies are not required, and the wall plug energy efficiency of a reactor type CPIL should exceed 6 percent, with potential for efficiencies between 20 and 30 percent with more efficient chemical regeneration techniques.²

One potential difficulty remained with regard to the excited iodine lifetime, that is, if the excited iodine atoms are quenched rapidly compared to the time between reactor pulses, then the excited oxygen/iodine mixture must be injected into the CPIL amplifier at very high velocity. To maintain high system efficiency, this gas flow would have to be pulsed immediately before each laser pulse and repeated periodically at the laser repetition rate.

During 1979 we modified the CPIL kinetic codes to include an analysis of thermal effects. Our results indicate that, based on our limited knowledge of the appropriate reaction rates at low temperatures, the excited iodine state can be maintained for more than 0.1 second after mixing. With these long lifetimes and for laser pulse rates of 10 pps or higher, the excited oxygen and iodine gases can be mixed in a CW gas flow at the edge of the amplifier as indicated in Figure 4.2-1. The flow velocity is adjusted so that one aperture width of excited iodine traverses the amplifier during each interpulse period, thereby eliminating the need for pulsed high speed mixing valves.

The wavelength of the CPIL ($1.315 \mu\text{m}$) may be too long for optimum target coupling, however its projected efficiency is sufficiently high that we could consider frequency multiplication. For example, frequency quadrupling the iodine output to $0.33 \mu\text{m}$ would yield a UV wavelength sufficiently short for good target coupling, but sufficiently long to avoid significant UV material problems such as low dielectric film damage thresholds. Assuming a

Fig. 4.2-1 Chemically pumped iodine laser with CW gas flow.

40 percent quadrupling efficiency, a CPIL system efficiency of 20 to 30 percent would translate into a total UV laser efficiency of 8 to 12 percent, which is higher than the predicted efficiencies of other short wavelength driver candidates.

Our second advanced laser candidate, the free electron laser (FEL), deviates further from conventional laser technology. It can operate in a vacuum at any (pre-selected) wavelength, and requires accelerators (and possibly storage rings) to generate a relativistic electron beam. Since high energy accelerator technology is relatively well developed, the scaling of the FEL to reactor driver sizes may proceed faster than other advanced laser candidates. One EPRI sponsored study indicates that the cost of a 1 MJ light ion beam driver based on accelerator technology would be a \$75M, while a 1 MJ CO_2 or KrF laser might cost \$200M.³ Therefore, FEL reactor drivers might be less expensive than other, more conventional lasers. Finally, the system efficiency of the FEL, like the CPIL, may exceed 20 percent.

As reported in the KMSF 1978 Annual Report, we developed a single particle computer code to simulate the interaction of a light wave with a copropagating relativistic electron beam traveling through a rippled magnetic field.⁴ Preliminary performance results were obtained for a constant-period FEL amplifier. During 1979, this simulation code was extended to include variation of the magnet period and field and was improved by including the optical phase shift, electrostatic effects,

and energy spread. To increase energy extraction efficiency above 1 percent, the period and/or strength of the magnet field must be changed as a function of distance along the amplifier. We have analyzed variable period FEL amplifier designs which convert up to 40 percent of the electron energy into laser energy in one pass while minimizing the electron energy spread, a prerequisite for recirculating the electron beam in a multiple pass system. We have also analyzed the effects of a finite energy spread on the incoming electron beam, and these FEL amplifier designs can tolerate an electron energy spread which is up to 0.8 percent of its nominal energy.

The detailed results of our research on the chemically pumped iodine laser are presented in Section 4.2.1 below, followed by the details of our FEL research in Section 4.2.2. FEL theory is presented in Section 4.2.3.

4.2.1 Kinetic Analyses of Energy Storage in a Chemically Pumped Iodine Laser

Geo. Busch

Photolytically pumped iodine lasers appear to be scalable to meet the power and energy requirements for laser fusion.⁵⁻⁶ Unfortunately, the extrapolated efficiencies of this laser remain close to 1 percent, such that interest in the laser for fusion applications has been rather limited. However, the wall plug energy efficiency of iodine lasers will be substantially improved if the highly efficient CW chemically pumped iodine laser⁷⁻⁸ (CPIL) can be adapted to high power, short pulse operation. Pulsed fusion lasers will probably require much higher energy storage densities than have been used to date in the CW lasers. The question of whether quenching reactions at high concentrations will limit energy storage times to unacceptable ranges must be addressed. This Section describes the results of kinetic analyses of the chemical and physical processes in CPIL amplifying media which include predictions of energy storage times. The possibility of increasing energy storage times by cooling the reaction medium is also evaluated.

In the CPIL the upper laser level $I^*(^2P_{1/2})$ of atomic iodine is pumped by near resonant transfer of excitation from the metastable species $O_2^*(^1\Delta)$ in

the Reaction

The ground state atomic iodine species $I(^2P_{3/2})$ is produced by reaction of another metastable species, $O_2^{**}(^1\Sigma)$, with molecular iodine:

Reaction (1) and its reverse, as well as Reaction (2), occur at nearly gas kinetic collision rates, such that a steady state is established in $\leq 10\mu$ sec at 300°K, for reactant partial pressures in the torr range. However, even if flow speeds are supersonic, uniform mixing of reactants throughout amplifier volumes of meter dimensions will require minimum mixing times ≥ 1 ms, and thus the steady state concentrations of the four species in Reaction (1) will track the mixing process. Unfortunately, present information on quenching processes such as

indicates that the quenching rates may be comparable to the maximum achievable mixing rates for pressures of reactants which give energy storage densities sufficiently high to be of interest.

Unlike the severe conditions occurring in gas lasers pumped by e-beams, electrical discharges, or ultraviolet photolysis, relatively mild conditions exist in the CPIL because electronic excitation is introduced into the amplifying medium via the low energy species $O_2^*(^1\Delta)$, which is produced in a prior chemical reaction and transported to the amplifier. Also very little heat is released in the near resonant transfer of Reaction (1). Thus primarily low energy reactions will be important in the CPIL. Although this greatly reduces the number of reactions that need to be considered to model the system, a significant number still remain (see Table 4.2-1). These have been incorporated into a computer model in which the corresponding kinetic rate equations are numerically integrated to yield the concentrations of the important species in a CPIL as a function of time. In the present model, complete mixing of reactants is assumed at $t = 0$.

Table 4.2-1

Reactions and Rate Constants for CPIL Computer Model

Reaction	Rate Constant ^a (cm ³ Molecule ⁻¹) ^N sec ⁻¹
1. $I + O_2(^1\Delta) \longrightarrow I^* + O_2(^3\Sigma)$	7.6×10^{-11}
2. $I^* + O_2(^1\Delta) \longrightarrow I + O_2(^3\Sigma)$	2.6×10^{-14}
3. $O_2(^1\Delta) + M \longrightarrow O_2(^3\Sigma) + M$	8.0×10^{-21} b
4. $O_2(^3\Sigma) + M \longrightarrow O_2(^1\Delta) + M$	1.0×10^{-17} b
5. $I^* + O_2(^3\Sigma) \longrightarrow I + O_2(^1\Delta)$	2.6×10^{-11}
6. $I^* + O_2(^3\Sigma) \longrightarrow I + O_2(^3\Sigma)$	5.7×10^{-14}
7. $I + O_2(^1\Delta) \longrightarrow I + O_2(^3\Sigma)$	1.3×10^{-13}
8. $I^* + M \longrightarrow I + M$	5.0×10^{-18} b
9. $O_2(^1\Delta) + O_2(^1\Delta) \longrightarrow O_2(^3\Sigma) + O_2(^3\Sigma)$	2.0×10^{-17}
10. $O_2(^3\Sigma) + I_2(^1\Sigma) \longrightarrow O_2(^3\Sigma) + 2I$	2.3×10^{-10}
11. $I^* + I_2 \longrightarrow I + I_2$	3.1×10^{-11}
12. $I^* \longrightarrow I + h\nu$	6.0
13. $O_2(^3\Sigma) \longrightarrow O_2(^3\Sigma) + h\nu$	7.7×10^{-2}
14. $2I + I_2 \longrightarrow 2I_2$	3.7×10^{-30}
15. $2I + O_2 \longrightarrow I_2 + O_2$	4.2×10^{-32}
16. $2I + M \longrightarrow I_2 + M$	3.8×10^{-33} b

a. N = 0, 1, 2 for the first, second, and third order reactions respectively at 300° K.
b. M = helium.

Thermal effects resulting from energy release as chemical reactions proceed have been incorporated into the model. The temperature of the reaction mixture is followed by integrating the rate expression

$$\frac{dT}{dt} = -(C_p)^{-1} \sum_i \Delta H_i \frac{dN_i}{dt} \quad (I)$$

in which C_p is the heat capacity of the system, and ΔH_i and dN_i/dt are the enthalpy of formation and the rate of change of concentration for species i . The temperature is then used to compute the (temperature dependent) rate constants at each stage of the reaction process. The first order rate constants Reactions (12) and (13) are assumed to be independent of temperature while the empirical functional forms for the temperature dependences of third order iodine atom recombination reactions (Reactions 14-16) are obtained from the work of Ip and Burns.⁹ Sparse data exist^{10,11,12} for the temperature dependence of second order rate constants for some of the reactions in Table 4.2-1; however, a single approach which assumed the usual gas kinetic Arrhenius function form

$$k_i = Z_i e^{-\theta_i/T} \quad (II)$$

was used for all binary rate constants. Here Z_i is a

reactive collision frequency factor, or high temperature rate constant, and θ_i is an activation temperature. Since all reactions involve simple atomic or diatomic species, values of Z_i equal to one tenth the appropriate gas kinetic collision frequency factors were used. Measured rate constants at 300°K then give values for θ_i 's, which in turn are used in the numerical modeling program to determine the instantaneous rate constants as the reactions proceed and temperatures change. The results of a computer run for a CPIL amplifying mixture starting at 300°K are shown in Figure 4.2-2. The initial partial pressures of $O_2(^1\Delta_g)$ and I_2 are 5 torr and 0.25 torr, respectively. The initial partial pressure of $O_2(^3\Sigma_g)$ is chosen to be 1.5 times that of I_2 , or 0.38 torr, to be sure I_2 is completely dissociated. Since $O_2(^3\Sigma_g)$ is produced from $O_2(^1\Delta_g)$ by Reaction 9, the initial concentration of ground state O_2 is taken equal to that of $O_2(^3\Sigma)$.

A steady state condition is seen to develop in the system of Figure 4.2-2 in less than 1 μsec and is seen to last for the order of several hundred μsec before quenching reactions take over. In a laser, energy extraction would cease when the gain coefficient, which is proportional to $I^*(^2P_{1/2}) - 1/2 I(^2P_{3/2})$, drops to zero. Combining the latter constraint with a steady state treatment, it is possible to show that the extractable energy density in a particular mixture is

$$E_x = h\nu \left\{ \frac{2}{3} [I^*(^2P_{1/2}) - \frac{1}{2} I(^2P_{3/2})] + (2K+1)^{-1} [2K(O_2(^1\Delta_g)) - O_2] \right\} \quad (III)$$

Fig. 4.2-2 CPIL kinetic results at 300°K.

in which $h\nu$ is the photon energy and K is the steady state equilibrium constant given by the ratio of the rate constants for Reactions 1 and 5 ($K = k_1/k_5$). At early times, E_x for the mixture of Figure 4.2-2 is about 21 J/l.

Defining the time at which E_x drops to half its initial value (after the steady state is established) as the energy storage time $\tau_{1/2}$, it is found that $\tau_{1/2}$ for the mixture modeled in Figure 4.2-2 is about 500 μsec . Whereas energy storage times this short may be usable for pulsed fusion drivers, longer times would place far less severe constraints on mixing and flow systems in order to uniformly fill amplifier cross sections of meter dimensions while maintaining high efficiency. The question then arises as to what pulsed amplifier conditions might increase $\tau_{1/2}$ to more useful ranges.

Storage times may be lengthened either by decreasing concentrations of quenchers or by decreasing the rate coefficients for quenching processes. Both techniques appear possible and we evaluate each, beginning with the former.

Since quenching occurs via alternate reactions between primary oxygen and iodine species, as opposed to being due to impurities,¹⁵ a reduction in concentrations of oxygen and/or iodine is required. Decreasing $O_2^*(^1\Delta_g)$ lowers energy storage density and thus necessitates increasing amplifier lengths to achieve high laser energy fluences. Alternatively, decreasing the amount of iodine requires larger numbers of passes of the laser beam through the amplifiers to efficiently extract the stored energy. Multiple passes of the amplifier are already envisioned for most applications requiring pulsed energy extraction of a CPIL because the energy transfer is reversible such that significant energy remains stored in $O_2^*(^1\Delta_g)$. Short laser pulses ($\ll 1 \mu\text{sec}$) extract only the energy stored in $I^*(^2P_{1/2})$, because the transfer rate (at $O_2^*(^1\Delta_g)$ pressures near 1 torr) is only $\sim 10^6 - 10^7 \text{ sec}^{-1}$. However, if the laser pulse can be repassed through the amplifier after a time delay $\sim 10^{-6} - 10^{-7} \text{ sec}$, then more energy will again have been transferred to the iodine, and this energy can be extracted. The majority of E_x can be extracted when the number of passes is $\sim R = O_2^*(^1\Delta_g)/I_2^0$, where I_2^0 is the initial iodine molecule concentration, assuming initial uniform distribution prior to any reaction.

The computer model has been used to determine

Fig. 4.2-3 Energy storage time $\tau_{1/2}$ at 300°K.

the dependence of $\tau_{1/2}$ on I_2^0 and $O_2^*(^1\Delta_g)$. The results, for gas mixtures at 300°K, are shown in Figure 4.2-3, where $\tau_{1/2}$ is plotted vs. $O_2^*(^1\Delta_g)$ for various values of R . The cross hatched area in the Figure marks the likely experimental regions of interest for a laser amplifier, i.e.: $O_2^* > 0.1 \text{ torr}$, $\tau_{1/2} > 3 \mu\text{sec}$, and $R < 40$. The region is sufficiently large to permit some flexibility in amplifier design. However, since all of the above limits are almost certainly extrema, the answer to whether a CPIL amplifier at 300°K would be useful as a fusion driver awaits both experimental verification that the model utilized is sufficiently accurate and refined estimates of the requirements for quantities such as $\tau_{1/2}$ and R .

Low Temperature CPIL Operation. If the quenching reactions limiting energy storage times are characterized by moderate activation temperatures, θ (see Equation II), then cooling may be effective in increasing $\tau_{1/2}$. Figure 4.2-4 shows kinetic model results for a system starting at initial temperature of 100°K. The value of $\tau_{1/2}$ for this case is 0.2 sec, which is about 400 times longer than that for operation at 300°K. Maximum $\tau_{1/2}$ values were obtained near 100°K. Below this temperature storage times are determined by the recombination of iodine atoms, which accelerates at lower temperatures producing I_2 , which is an efficient quencher (Reaction eleven, Table 4.2-1).

Fig. 4.2-4 CPIL kinetic results at 100°K.

The dependence of $\tau_{1/2}$ on O_2^* ($\triangle$) and R at 100°K is shown in Figure 4.2-5. Comparison with Figure 4.2-3 reveals that much larger ranges of concentrations of O_2^* ($\triangle$) and I_2^0 give useful energy storage times, with a significant region of the cross hatched area lying above 0.1 sec. Storage times greater than 0.1 sec open the possibility of employ-

Fig. 4.2-5 Energy storage time $\tau_{1/2}$ at 100°K.

ing continuous gas flows in a CPIL pulse amplifier at very moderate flow rates. It is recognized however that low temperature operation would introduce its own set of problems, such as condensation, even if the crude assumptions of Arrhenius type scaling of rate constants employed here do adequately approximate the behavior of the chemical system in a low temperature CPIL amplifier.

In conclusion, operation of a pulsed CPIL with the high efficiencies demonstrated in CW operation and with acceptable energy storage densities and times appears to be feasible. Furthermore, if low temperature operation can be realized, the highly attractive possibility of energy storage over the entire time interval (~ 0.1 sec) between pulses of a fusion driver may also be achievable. In this case the CPIL may become a very attractive reactor driver candidate, combining high efficiency, short wavelength (through frequency conversion), and a continuous, closed loop chemical flow in place of pulsed electrical power supplies.

4.2.2 The Application of the FEL to Laser Fusion

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KMSF has been investigating the free electron laser (FEL) as a fusion driver for the past two years. The FEL has several advantages; it is tuneable and can therefore be designed to operate at any wavelength found to be optimum for the laser target interaction, and overall laser efficiencies greater than 20 percent appear feasible. The FEL can operate totally in vacuum with no transparent optical components in the beam path. Self focusing and thermal distortion in transparent optical components will therefore not occur. The only optical elements would be mirrors and beam diameter at the mirrors could be made large enough to prevent damage from high intensity laser light. Also, the beam could be passed through the same amplifier an arbitrary number of times without degradation due to thermal distortion. Examples of multipass

FEL systems and electron storage rings to recirculate the electron beams have been described in previous reports.¹⁶⁻¹⁷

Multipass and Single Pass Concepts. In one design for a driver system¹⁷ (Figure 4.2-6), we considered a scheme with two optical cavities and four laser amplifiers. Each amplifier had associated with it a storage ring loop with a booster accelerator to replace energy loss. An electron pulse is circulated in each loop. Using 500 MeV electron beams with 5 percent fractional energy conversion per amplifier, 0.2 percent optical losses per pass, and 500 passes per target shot, we could obtain 100 TW of optical power in 4 optical beams which intersect at the target focal volume.

Fig. 4.2-6 Design for a fusion driver system with two optical cavities and four laser amplifiers.

An alternative fusion driver concept employing single pass FEL amplifiers has been proposed by A. Szöke of The Lawrence Livermore National Laboratory.¹⁸ In this concept, two large induction linacs produce a train of 50,000 amp 500 MeV short

pulses. The energy of each pulse is slightly different and the pulses are deflected along different paths to an array of FEL amplifiers. An auxiliary laser is needed to provide a high power pulse to each of the FEL amplifiers. It is assumed that 25 percent of the electron energy is converted to laser energy in each amplifier, producing a gain of 1000. An output power of 100 TW would require 16 amplifiers.

The advantage of the single pass approach is that it is not necessary to correct for electron energy spread induced in the FEL amplifier since each electron beam is only used once. Obtaining sufficiently low beam emittance from the induction linac appears to be as formidable a problem as controlling energy spread in the storage ring. We believe that if multipass systems can be made to work, they would be preferable to single pass systems in commercial fusion reactors. We have, therefore, concentrated our effort on the multipass approach.

Several problems must be addressed to determine whether a multipass FEL could be made to operate in a high power high efficiency pulsed mode, and whether such a laser could meet the specific requirements of a fusion driver.

Energy Conversion and Energy Spread. Correction or compensation for electron energy spread, is the fundamental problem associated with recirculating systems, and most of the effort this year has been directed toward this problem. In the interaction of a laser beam with an electron beam in a free electron laser amplifier, electrons are both accelerated and decelerated. In order to obtain net gain, more electrons must be in a decelerating region with respect to the wave than are in an accelerating region. Bunching of the electron beam must therefore occur on the scale of the wavelength of the laser light. Assuming conservation of phase space area, when the spatial extent of the beam is contracted in phase space, there must be a corresponding increase in the momentum or energy spread. Therefore, laser gain cannot take place without increasing the electron energy spread.

Soon after the principle of the FEL was first demonstrated at Stanford University, it was shown¹⁹⁻²⁰ that, for a wiggler magnet with constant periodicity, the energy spread produced in the amplifier was of the same order as the mean electron energy loss. Calculations performed at

KMSF²¹ also indicated that the maximum fractional energy conversion that could be obtained with a constant period magnet with constant field amplitude was of the order of 1 percent. Therefore, both the mean energy and energy spread from such devices would be unsuitable for a fusion system.

To obtain high fractional energy conversion with reduced energy spread, it is necessary to use magnet designs in which either the magnet period or field strength or both are varied along the length of the magnet. Calculations at KMSF²² and other laboratories²³ have shown that using such tapered wigglers, up to several tens of percent of the electron energy can be converted to photon energy on a single pass of an electron beam through an FEL amplifier. Most of these calculations were carried out for proposed single pass systems and no attempt was made to minimize energy spread. In the KMSF calculations, however, energy conversion was carried out in a way that minimized energy spread, and a reduction of the ratio of energy spread to mean energy loss of about an order of magnitude was achieved compared with constant period systems.

Different approaches for dealing with energy spread are being pursued at other laboratories. Madey²⁴ has proposed design modifications that permit the FEL amplifier to operate efficiently with a larger energy spread. Kroll, Morton and Rosenbluth²⁵ have proposed a scheme that permits energy to be extracted from a beam with a relatively large energy spread without significantly increasing the existing spread. More study will be needed to determine the most appropriate design for a fusion system with a recirculating electron beam, but sufficient progress has already been made to indicate that solutions to this problem probably can be found.

Energy Spread and Storage Ring Design. Although the primary approach to reducing energy spread is to prevent a large spread from developing in the amplifier, some reduction in spread can be achieved in the storage ring. Synchrotron radiation is the primary means of reducing energy spread in all conventional storage rings. In continuous operation, energy spread in a storage ring is reduced to a fundamental limit determined by the spread in direction and energy of the photons of synchrotron radiation. This minimum energy spread determines

the electron beam diameter in the ring. Collisions which scatter electrons out of the trapped beam region then determine the maximum current density and hence the maximum current that can be carried in the ring under equilibrium conditions. Typical operating parameters for a conventional storage ring designed to operate at 500 MeV would be energy spread $\sim 10^{-4}$ and maximum current density $\sim 10^4$ A/cm² in a beam of the order of 100 mm cross section. The limiting current in such a ring would therefore be about 100 A, which is much less than that required for a fusion system.

To obtain higher currents in a storage ring it is necessary to increase beam energy spread. This increases the electron beam diameter and, for a given limiting current density, permits a higher total current in the beam. The FEL amplifier provides this increase in electron energy spread. The problem is matching the spread that is produced with the maximum spread that the amplifier can tolerate without serious degradation of laser gain. In calculations carried out at KMSF we have been able to obtain high fractional energy conversion without loss of particles from the closed orbit region of phase space with electron energy spreads of up to 0.8 percent. We believe that an energy spread of a few tenths of a percent can provide the spread needed for high beam current without significant degradation of amplifier performance.

Amplifier Design Considerations. Although we have not yet fully optimized our amplifier designs, it appears that rather long amplifiers may be required in order to convert several percent of the electron energy to photon energy while still controlling energy spread. Present calculations indicate that for a 500 MeV electron beam, the amplifier may be of the order of 100 m long. The most straightforward way to reduce the length of the amplifier is by lowering the electron energy. It is doubtful that the electron energy could be reduced much below 200 MeV, since the periodicity and hence the diameter of the wiggler magnet must decrease as the electron energy decreases to maintain the desired laser wavelength.

Amplifier simulation studies indicate that amplifier performance on any given pass is dependent on the initial conditions for that pass. Input laser intensity is one of those initial conditions,

and in a pulsed fusion system initial laser intensity will vary on every pass. To accommodate the changing laser intensity, magnet parameters probably will need to be varied as a function of time. Pulsing of the amplifier current windings would take place over a period of about a millisecond.

Reactor Design Constraints. In the baseline FEL design that we have proposed the fusion reactor is located inside the FEL optical cavity and the optical beam passes many times through the target focal volume before it is intercepted by an injected target. A potential problem with this approach is scattering of the optical beam by the residual background gas in the target chamber. A background gas $< 10^{-3}$ torr is necessary in the target chamber to prevent such scattering. Lithium wetted wall or waterfall reactors would probably have too high a background pressure for use with this type of system. One would prefer to operate a fusion reactor at a much higher temperature of about 1000-2000°C in order to achieve high thermodynamic efficiency. For the FEL one would also like a reactor in which the background pressure at the time of the pulse buildup would be limited to residual target gases that had not been removed between pulses, rather than vapor from the walls of the chamber.

Two possible reactor structures could accommodate these demands. The first would be a reactor with a sacrificial first wall made from a material like carbon, which would be replaced periodically. The second possibility is to use a waterfall reactor but with ceramic pellets containing lithium rather than a liquid metal stream.

Input Energy System. Possibly the largest single component of the cost of a laser fusion power plant using a conventional gas laser would be the pulsed power supply. In a multipass FEL some pulsed power may be needed to operate the amplifiers, but most of the power will be supplied by RF accelerating cavities. During the period of pulse buildup, several hundred megawatts of pulsed RF power must be supplied to the electron beams. If several hundred megawatts of pulsed RF capacity must be installed in order to operate the laser, then the cost of the power supply might be greater for a multipass FEL system than for other types of lasers.

However, since the system is only expected to pulse about 10 times a second with RF, power being supplied about 1 to 2 percent of the time, a dramatic reduction in supply requirements could be achieved if the required RF power were produced continuously and dumped into the accelerating cavities only during the laser pulse buildup.

System Efficiency. The overall efficiency of a multipass FEL system depends on energy losses incurred correcting for energy spread, the efficiency with which the electrons are accelerated, optical losses, and the duty cycle for a pulsed system. Laser efficiency for a recirculating system is not directly dependent on fractional energy conversion per pass.

Depending on the mechanisms used to correct or compensate for energy spread, losses can range from many times the energy in the emitted laser light to a fraction thereof. In order for the energy loss to be less than the laser energy, energy spread in the amplifier must be limited to a small fraction of the mean electron energy loss.

The efficiency with which electrons can be accelerated in a conventional RF cavity accelerator is presently about 50 percent. Using superconducting accelerating cavities and more efficient methods of producing the accelerating fields this value may be significantly increased, efficiencies as high as 90 percent may be possible.²⁶

Losses in the optical resonator both limit the peak intracavity power and reduce efficiency. To obtain high power without greatly sacrificing efficiency, we assume the number of optical passes will be limited to $1/\delta$ where δ is the fractional loss per pass.

Table 4.2-2 gives multipass FEL efficiencies for a range of assumed conditions. In calculating these values, it was assumed that 37 percent of the emitted laser light had been lost in the optical system. Under the most optimistic conditions an overall laser efficiency of about 50 percent is obtained.

The efficiency values calculated in Table 4.2-2 do not include fixed energy drains on the system such as vacuum pumps, control equipment, energy required to produce and maintain the high current electron beams at the desired current level, and energy required to produce the magnetic fields. It is assumed that in a system designed to operate at maximum energy efficiency, all constant magnetic

Table 4.2-2
Multipass FEL Efficiency

Energy Lost Correcting For Energy Spread	Efficiency of Electron Acceleration	
	50 %	90 %
$E_{loss} = 10 E_{laser}$	2.9	5.2
$E_{loss} = E_{laser}$	15.8	28.4
$E_{loss} = 0.1 E_{laser}$	28.4	51.0

fields would be produced using permanent magnets or superconducting magnets, and superconductors would be used in accelerating structures and variable field magnets when major energy savings would result from their use.

Since the fixed energy costs are for the most part independent of the rate at which the laser is fired, their effect on overall system efficiency becomes smaller as the repetition rate increases. It is estimated that about 5 MW of power will be required to operate auxiliary equipment for a 100 TW pulsed FEL laser system. Figure 4.2-7 is a power flow diagram for an FEL inertial confinement fusion (ICF) power plant. It is conservatively assumed in this diagram that the booster accelerators are 50

percent efficient. Based on expected improvements in amplifier design, it is estimated that $E_{Loss} = 0.1 E_{Laser}$, and the optical resonator is assumed 63 percent efficient. Net driver efficiency is then calculated to be 28.4 percent. For a repetition rate of ten one-megajoule pulses per second, 35.2 MW of input power would be needed, which would be much greater than the 5 MW of fixed power required to operate the auxiliary equipment.

The high overall efficiency of the FEL presents many advantages for reactor design. Part of the output of the power plant is needed to operate the laser. It is desirable to have this recirculating fraction of the electrical energy as low as possible. The higher the efficiency of the laser, the lower the target gain required to keep the recirculating power fraction low. For the example in Figure 4.2-7 a target gain of 100 would result in a recirculating power fraction of 10.6 percent assuming an electrical generation efficiency of 38 percent.

An advantage of being able to operate with low target gain is that it would be possible to build smaller power plants. Since the FEL could be tuned to operate at the optimum frequency for target implosion, the power output required to achieve a given target gain should be lower for the FEL than for other lasers. For the system shown in Figure 4.2-7 net output power to the utility grid would be about 340 MW.

Fig. 4.2-7 Power flow in an FEL/ICF power plant.

4.2.3 Free Electron Laser Theory

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Constant Period FEL Amplifiers. At the beginning of this year we had completed writing a computer code which could simulate the interaction of an initially monoenergetic electron beam with a monochromatic laser beam in a helical magnetic amplifier of constant periodicity and with a constant magnetic field amplitude. Optical phase shift and electrostatic forces were not included. This computer simulation code was used to determine the performance of constant period helical amplifiers over a wide range of parameters.²¹ Optimum amplifier performance was obtained by determining the off resonance laser frequency that gave the highest gain and determining the amplifier length providing maximum amplification for that frequency.

Fig. 4.2-8 Fraction of the electron energy converted to laser energy for a constant period, constant field, helical wiggler.

Fig. 4.2-9 Amplifier length required for maximum laser gain for a constant period, constant field, helical wiggler.

It was found that all optimized solutions could be represented by a single universal curve (obtained by fitting the data points in Figure 4.2-8), and that for low gain the optimum length of the amplifier could be expressed as a function of electron energy, amplifier magnetic field, and laser intensity (Figure 4.2-9). From the simulation data of Figure 4.2-8, the fraction of the electron energy converted to laser energy for a constant period, constant field helical wiggler under optimum conditions is given by

$$\frac{\eta}{s} = .32 \frac{K}{\sqrt{s}} \left(1 + \frac{1}{4} \left[\frac{K}{\sqrt{s}} \right]^{2/3} \right) \quad (1)$$

where η is the fractional energy conversion, s is the ratio of energy density in the laser beam to energy density in the electron beam at the input end of the amplifier and

$$K = \frac{e^2 B_m \lambda_m^2 N \rho^{1/2}}{\pi^{1/2} (\gamma_0 m c^2)^{3/2}} \quad (2)$$

In Equation (2) B_m and λ_m are the magnetic field strength and periodicity, respectively, e and m are the electronic charge and rest mass, γ_0 is the initial ratio of electron energy to rest mass energy, N is the number of periods in the amplifier, ρ is the electron density and c is the speed of light.

From the solid line in Figure 4.2-9, the amplifier length required for maximum laser gain is given by

$$L = 5.9 \frac{\gamma_0}{B_m^{1/2} I_0^{1/4}} \quad (3)$$

where L is in meters when B_m is in kG, and I_0 , the input laser intensity, is in W/cm². Equation (3) is accurate when the laser gain is low and the input laser intensity is only slightly different than the output intensity. However, because of the very weak dependence of amplifier length on laser gain, an error of only 5 percent in the value of L would result when the output intensity is a factor of 50 greater than the input intensity. Points shown in Figure 4.2-9 which deviate significantly from the straight line which defines Equation (3) are for cases with laser gains between 10^2 and 10^4 .

Variable Period FEL Amplifiers. From our study of constant period FEL amplifiers, we concluded that no more than about 1 percent of the energy in the electron beam could be converted to photon energy for attainable values of electron beam and magnet parameters. For high efficiency, high power lasers it is therefore necessary to consider magnet designs in which the magnet periodicity and field strength are not held constant.

During this year, the computer simulation code has been upgraded to investigate amplifier designs in which the magnet periodicity and field amplitude are varied with position. Graphics have been added to assist in determining when the electrons remain trapped in the closed orbit region of phase space, where their energy spread is limited. Optical phase shift, electrostatic forces and initial energy spread have also been added to the computer simulation. Using this upgraded version of the code, we have shown that all of the electrons in the beam can be trapped in the closed orbit region of phase space, and that the electron energy spread can be made much smaller than the mean energy loss of the electrons.

In the variable period computer simulation code, particle trajectories are determined by forward integration of the single particle differential equations of motion for the electrons in the combined magnet and laser fields. The electrons are assumed to enter the amplifier as a spatially uniform beam, and the laser light is taken to be monochromatic with self consistently determined amplitude. It is assumed that at each time step the amplifier can be treated as

if it were a constant period helix, but the magnet period is permitted to vary from time step to time step. Typical variations of the magnet period between time steps are on the order of 10^{-4} of the magnet period.

In the computer simulation a segment of an electron beam one radiation wavelength long was followed through the amplifier. The electron beam was represented by a linear array of 100 particles initially equally spaced over one radiation wavelength. Electrons entering the amplifier passed through a constant period pre-buncher with the periodicity chosen so that the mean electron energy was the resonance energy. After about 70 percent of the electrons had been bunched to within a quarter of an optical wavelength, the magnet wavelength was slowly decreased to trap the particles. The product of magnetic field amplitude and magnet period was held constant in the simulation runs as the length of the magnet period was decreased.

Variation of magnet period with distance was determined by the deceleration of a test particle moving with the electron beam. The magnet period was varied so that the test particle always remained at resonance. The relative phase of the test particle was changed between time steps to vary the rate of deceleration. A small linear shift in optical phase was also found to be necessary to trap all the electrons in the closed orbit region of phase space.

Pendulum Phase Space Diagrams and Particle Trapping. For a simple constant period helical amplifier with negligible gain, the motion of electrons in the beam can be described exactly by the simple pendulum equation²⁰

$$\ddot{z} = \Omega^2 \sin \psi \quad (4)$$

where

$$\Omega^2 = \frac{(1 + \alpha^2) \alpha}{\gamma^4} \frac{e}{m} E_r \quad (5)$$

and

$$\psi = (k_r + k_m) z - \omega_r t + \phi. \quad (6)$$

Here ψ is the angle between the laser magnetic field vector $\vec{B}$, and the perpendicular component of the electron velocity, and

$$\alpha = \gamma \beta_{\perp} = \frac{|e| B_m}{m c k_m}$$

in mks units. E_r is the amplitude of the laser electric

field, k_m and k_r are the wave numbers for the magnet period and laser radiation, respectively, z is the electron position, ω is the laser frequency, and ϕ is a phase factor for the electromagnetic wave. Several of these orbits are shown in Figure 4.2-10.

Fig. 4.2-10 Phase space orbits for electrons in a constant period helical amplifier. The closed orbits represent the motion of particles bound in the ponderomotive potential and the open orbits correspond to unbound particles.

Electrons which move in closed orbits in phase space are described as trapped. The largest permissible closed orbit is called the separatrix. The equation of the separatrix is given by

$$\dot{z} = \pm \sqrt{\frac{2\Omega^2}{k_r} (1 - \cos \psi)} \quad (7)$$

when $\phi = \text{constant}$.

Including the separatrix curve on phase space diagrams showing electron velocity as a function of position immediately indicates whether there are untrapped electrons in the beam.

The electron orbits produced by Equation (4) accurately describe the motion of electrons in the prebunching section of an amplifier where laser gain is negligible. After the electrons are prebunched in our computer simulations, the magnet period is varied to reduce the mean energy of the electrons by keeping the electron bunches in a position where most of the electrons are decelerated by the electromagnetic wave. It is desirable to vary the

magnet periods in such a way that the electrons remain inside the closed orbit region of phase space, where the range of energy values they can have is bounded, while they are decelerating. Electrons located in the open orbit region of phase space will pass through all phase values of the axial ponderomotive force and will not, on average, contribute energy to the laser beam.

It has been shown in a number of papers^{23,25} that the electron orbits given by Equation (4) do not accurately describe the motion of the electrons in a variable period device. A more accurate set of orbit equations can be derived assuming a slow variation of the magnet period with position along the amplifier.

To derive the orbit equations for a variable period device we start by differentiating Equation (6) with respect to time, assuming $\phi = \text{constant}$:

$$\begin{aligned} \dot{\psi} &= (k_m + k_r)\dot{z} - \omega_r \\ &= (k_m + k_r)\beta_z c - k_r c \\ &= k_m c - c(k_m + k_r)(1 - \beta_z) \\ &\approx k_m c - ck_r \left(\frac{1 + \alpha^2}{2\gamma^2} \right) \end{aligned} \quad (8)$$

where

$$\beta_z = v_z/c, \quad \omega_r = k_r c, \quad 1 + \beta_z \approx 2, \quad k_m + k_r \approx k_r$$

and

$$1 - \beta_z^2 = \frac{1 + \alpha^2}{\gamma^2} \quad (9)$$

which can be derived from the definitions of γ and α .

Using $\frac{d\psi}{dt} = \frac{d\psi}{dz} \frac{dz}{dt} \approx c \frac{d\psi}{dz}$, we obtain

$$\frac{d\psi}{dz} = k_m - \frac{k_r}{2\gamma^2} (1 + \alpha^2) \quad (10)$$

Since only the electric field can do work on the electrons, we have

$$\frac{d\gamma}{dt} = \frac{e}{mc} E_r \cdot \beta_{\perp} = \frac{e}{m} \beta_{\perp} \times B_r = \frac{\alpha}{\gamma} \frac{e}{m} B_r \sin \psi$$

and

$$\frac{d\gamma}{dz} = \frac{\alpha}{\gamma} \frac{e B_r}{mc} \sin \psi \quad (11)$$

in mks units.

Equations (10) and (11) hold for any electron

and, in particular, for the test or synchronous particle which is used to determine the variation of magnet period with position. The synchronous particle is assumed to be moving at the resonance velocity where one wavelength of light passes over it as it moves one magnet period. The rate at which the velocity of the synchronous particle changes due to the ponderomotive force acting on it determines the rate at which the magnet period and field must be varied to maintain the resonance condition. For the synchronous particle we have

$$\frac{d\psi_s}{dz} = k_m - \frac{k_r}{2\gamma_s^2} (1 + \alpha^2) \quad (12)$$

and

$$\frac{d\gamma_s}{dz} = \frac{\alpha}{\gamma_s} \frac{eB_r}{mc} \sin \psi_s \quad (13)$$

Subtracting Equation (12) from (10) and (13) from (11), we obtain

$$\frac{d}{dz}(\psi - \psi_s) = \left(\frac{1}{\gamma_s^2} - \frac{1}{\gamma^2} \right) \frac{k_r}{2} (1 + \alpha^2) \quad (14)$$

and

$$\frac{d}{dz}(\gamma^2 - \gamma_s^2) = 2\alpha \left(\frac{eB_r}{mc} \right) (\sin \psi - \sin \psi_s) \quad (15)$$

Assuming $\gamma^2 - \gamma_s^2 \sim (\gamma - \gamma_s) 2\gamma_s$ and defining $P = \gamma - \gamma_s$, we obtain

$$\psi' = \psi'_s + \frac{k_r}{\gamma_s^3} (1 + \alpha^2) P \quad (16)$$

and

$$P' = \frac{\alpha}{\gamma_s} \left(\frac{eB_r}{mc} \right) (\sin \psi - \sin \psi_s) \quad (17)$$

In deriving the above equations, we have assumed that $\phi' = 0$ so that there is no optical phase shift, and that the resonance condition is

$$1 + \alpha^2 = \frac{2\gamma_s^2 k_m}{k_r} \quad (18)$$

We will now also assume that the phase of the synchronous particle varies slowly (adiabatic approximation) so that $\psi'_s \cong 0$. Defining

$$A(z) = \frac{2k_m(z)}{\gamma_s(z)} \quad \text{and} \quad C(z) = \frac{\alpha(z)}{\gamma_s(z)} \frac{eB_r(z)}{mc}$$

we obtain

$$\psi' = A(z) P \quad (19)$$

$$P' = C(z) (\sin \psi - \sin \psi_s) \quad (20)$$

$$\frac{dP}{d\psi} = \frac{C(\sin \psi - \sin \psi_s)}{AP} \quad (21)$$

If we now assume that C/A is slowly varying and to first order can be assumed to be constant (low gain), we can integrate Equation (21) to obtain the orbit equations for a decelerating electron beam. From Equation (20), extrema for the electron orbits will occur when $\sin \psi = \sin \psi_s$. This will be the case for

$$\left. \begin{aligned} \psi &= \psi_s \pm 2\pi n \\ \psi &= \pi - \psi_s \pm 2\pi n \end{aligned} \right\} n = 0, 1, 2, \dots$$

Only open orbits will have both solutions. For positive values of P , the first set of solutions will be maxima and the second set minima. The limiting open orbit equation is the separatrix. Looking at a section of P, ψ space between 0 and 2π , the equation for the separatrix can be written

$$\frac{P^2}{2} = \frac{\alpha e B_r}{2mck_m} (-\cos \psi - \psi \sin \psi_s) \Big|_{3\pi - \psi_s}^{\psi}$$

where $\pi < \psi_s < 3/2\pi$. This can be written

$$P = \pm \left(\frac{\alpha e B_r}{mck_m} [-\cos \psi - \cos \psi_s + (3\pi - \psi - \psi_s) \sin \psi_s] \right)^{1/2} \quad (22)$$

The value of P at $\psi = \psi_s$ is

$$P_{max} = \pm \left(\frac{2\alpha e B_r}{mck_m} [-\cos \psi_s + (3/2\pi - \psi_s) \sin \psi_s] \right)^{1/2} \quad (23)$$

The closed orbit region in a reference frame decelerating at a rate equal to that of an electron located at ψ_s is shown in Figure 4.2-11. The closed orbit region for a constant period helix in which the resonant electron velocity is not changing is also shown. It can be seen that the closed orbit region of phase space is smaller for a tapered wiggler than for a constant period device. Therefore, in an amplifier with an initial constant period prebuncher followed by a tapered wiggler, some electrons that start out in the closed orbit region of phase space may move into unbounded orbits as the closed orbit region contracts. This effect would be particularly detrimental in systems with recirculating electron

Fig. 4.2-11 Separatrix for closed orbit region of electrons decelerating at same rate as electron located at ψ_s . When $\psi_s = \pi$, there is no deceleration, and the separatrix becomes that of the constant period case.

beams where one wishes to minimize energy spread by maintaining the electrons in the bounded closed orbit region of phase space.

The equations obtained from Equation (21) are a better approximation to the actual electron orbits than Equation (4), but one should remember that Equation (21) is only valid when the assumptions made in deriving it hold. It accurately describes the motion of particles only if the amplitude of the orbits in phase space remain constant, the phase of the synchronous particle varies slowly, and the phase ϕ of the optical wave remains constant. These assumptions may not hold in a high gain tapered wiggler. Nevertheless, the phase space diagrams obtained from Equation (21) can be helpful in predicting system performance and determining the proper variation of magnet period needed to optimize amplifier performance.

Inclusion of Optical Phase Shift in the Simulation Code. Using an approach similar to that of Colson and Ride,²⁷ we derive a linearized form of the wave equation for the magnetic vector potential of the light waves from which both the amplitude and phase of the electromagnetic wave can be calculated.

The wave equation for the magnetic vector potential of an infinite plane electromagnetic wave

can be derived from Maxwell's equation and is given by

$$\frac{\partial^2 \vec{A}}{\partial z^2} - \frac{1}{c^2} \frac{\partial^2 \vec{A}}{\partial t^2} = -\mu_0 \vec{J} \quad (24)$$

where $\vec{A}(z, t)$ is the magnetic vector potential which is assumed to have the form

$$\vec{A}(z, t) = -\frac{B_r(z, t)}{k_r} (\cos \alpha, \sin \alpha, 0) \quad (25)$$

with

$$\alpha = k_r z - \omega_r t + \phi(z, t). \quad (26)$$

$\vec{J}$ is the electron current.

Substituting Equation (25) in Equation (24) and neglecting all terms with second derivatives or products of first derivatives we obtain two linear first order differential equations for the field amplitude B_r and phase ϕ .

$$\frac{\partial B_r}{\partial z} + \frac{1}{c} \frac{\partial B_r}{\partial t} = \frac{\mu_0}{2} \vec{J}_\perp \cdot \hat{e}_2 \quad (27)$$

and

$$B_r \left(\frac{\partial \phi}{\partial z} + \frac{1}{c} \frac{\partial \phi}{\partial t} \right) = -\frac{\mu_0}{2} \vec{J}_\perp \cdot \hat{e}_1 \quad (28)$$

where we define

$$\hat{e}_1 = (\cos \alpha, \sin \alpha, 0) \text{ and } \hat{e}_2 = (-\sin \alpha, \cos \alpha, 0)$$

with $\hat{e}_1 \cdot \hat{e}_2 = 0$.

The perpendicular component of the electron current, $J_\perp$, is made up of the contributions from all of the charges in the distribution. In the simulation program

$$\vec{J}_\perp = \sum_{i=1}^N \frac{e\rho}{N} \vec{v}_{\perp i} = -\frac{|e|\rho}{N} \sum_i \vec{v}_{\perp i} \quad (29)$$

Here $\hat{e}_1$ is a unit vector in the direction of the laser magnetic field, $\hat{e}_2$ is a unit vector in the direction opposite the laser electric field and $\vec{J}_\perp \cdot \hat{e}_2$ has the same sign as the work done by the electrons.

$$\vec{J}_\perp \cdot \hat{e}_2 = -\frac{|e|\rho}{N} \sum v_{\perp} \sin \psi \quad (30)$$

where ψ is the angle between B_r and $v_\perp$.

The phase factor for $\vec{J}_\perp \cdot \hat{e}_2$ is given by

$$-\sin \psi = \cos(\pi/2 + \psi).$$

The phase angle between $\vec{J}_\perp$ and $\hat{e}_1$ will be $\pi/2$ less than the phase angle between $\vec{J}_\perp$ and $\hat{e}_2$, so that the phase factor for $\vec{J}_\perp \cdot \hat{e}_1$ will be $\cos \psi$.

By taking the sum of the contributions of the electrons to a total electron current we are averaging over space. Also by using a reference frame moving with the electron beam we essentially only look at the time variation of the electron beam and not the spatial variation. We can, therefore, simplify Equations (27) and (28) to

$$\frac{\partial B_r}{\partial t} = \frac{\mu_0 c}{2} \langle \vec{J}_\perp \cdot \hat{e}_2 \rangle \quad (31)$$

and

$$\frac{\partial \phi}{\partial t} = \frac{-\mu_0 c}{2 B_r} \langle \vec{J}_\perp \cdot \hat{e}_1 \rangle \quad (32)$$

where the brackets around the scalar products represent a spatial average. Substituting the expression for $J_\perp$ and using finite differentials we obtain

$$\Delta B_r = -\frac{\mu_0 c}{2} Q \Delta t \sum v_\perp \sin \psi \quad (33)$$

and

$$\Delta \phi = \frac{\mu_0 c}{2 B_r} Q \Delta t \sum v_\perp \cos \psi \quad (34)$$

where $Q \equiv \frac{|e| \rho}{N}$ as defined in the computer program.

We calculate the value of the laser magnetic field B_r at every time step in the computer simulation from conservation of energy, and we can compare the expression used in the computer simulation with the expression for the field obtained using the linearized wave equation.

In the program

$$B_r = \left(\frac{1}{\epsilon_0 c^3} \right)^{1/2} P^{1/2} = \left(\frac{\mu_0}{c} \right)^{1/2} P^{1/2} \quad (35)$$

where P is the laser intensity in watts per meter squared.

$$\Delta B_r = \left(\frac{\mu_0}{c} \right)^{1/2} (P_n^{1/2} - P_{n-1}^{1/2}) \quad (36)$$

where the subscripts refer to time steps.

$$P_n = P_{n-1} - c^2 \Delta t Q B_{n-1} \sum v_\perp \sin \psi \quad (37)$$

Therefore,

$$\Delta B_r = \left(\frac{\mu_0}{c} \right)^{1/2} \left[(P_{n-1}^{1/2} - c^2 \Delta t B_{n-1} Q \sum v_\perp \sin \psi)^{1/2} - P_{n-1}^{1/2} \right]$$

$$= \left(\frac{\mu_0}{c} \right)^{1/2} P_{n-1}^{1/2} \left\{ \left[1 - \frac{1}{2} \left(\frac{\Delta P}{P_{n-1}} \right) - \frac{1}{8} \left(\frac{\Delta P}{P_{n-1}} \right)^2 - \frac{1}{16} \left(\frac{\Delta P}{P_{n-1}} \right)^3 \dots \right] - 1 \right\}$$

$$\approx - \left(\frac{\mu_0}{c} \right)^{1/2} P_{n-1}^{1/2} \frac{c^2 \Delta t Q \sum v_\perp \sin \psi B_{n-1}}{2 P_{n-1}} \quad (38)$$

Let $K \equiv c \Delta t Q \sum v_\perp \sin \psi$

Then,

$$\Delta B_r = - \left(\frac{\mu_0}{c} \right)^{1/2} \frac{c B_{n-1} K}{2 P_{n-1}^{1/2}} \quad (39)$$

But

$$P_{n-1}^{1/2} = \left(\frac{c}{\mu_0} \right)^{1/2} B_{n-1}$$

so

$$\Delta B_r = - \frac{\mu_0 K}{2} \quad (40)$$

Equation (40) is a linearized form of the expression used in the computer simulation to determine B_r . Equation (40) has exactly the same form as (33) which shows that the linearized wave equation gives exactly the same value for the laser magnetic field as a linearized version of our present method of computation. The difference between ΔB_r obtained from (36) and (40) is a few parts in 10^4 so the linear approximation appears to be adequate. In the simulation code we use Equation (35) to compute B_r , and Equation (34) to calculate the optical phase shift.

Inclusion of optical phase shift in the computer simulation can dramatically affect the dynamics of the interaction of the laser beam with the electrons, especially in high gain amplifiers. One can achieve high gain and high fractional energy conversion in an FEL only by carefully maintaining the electron bunches at the proper phase with respect to the optical wave. If an amplifier is designed to maintain the proper phase relations neglecting phase shift, large phase shifts can produce entirely different phase relationships which dramatically lower laser gain.

It is, however, possible to totally compensate for even large optical phase shifts by proper amplifier design, permitting one to use all of the techniques developed for amplifier optimization neglecting

phase shift. The resonance condition occurs when one wavelength of light passes over an electron in the time it moves one magnet period. With optical phase shift the electron energy for which this condition is satisfied is different than when optical phase is constant. The location of the closed orbit region of phase space will also be different with optical phase shift than without. Figure 4.2-12 is a "snapshot" of the phase space of an electron beam in the prebunching section of an amplifier that has been designed to correct for optical phase shift. The frame of reference in Figure 4.2-12 is traveling with the velocity of an electron which satisfies the constant phase resonance condition of Equation (18). It is seen from Figure 4.2-12 that the closed orbit region is centered about an electron velocity which is considerably lower than the constant phase resonant velocity.

Fig. 4.2-12 Displacement of the closed orbit region of phase space as a result of optical phase shift. The reference frame is moving with the velocity of an electron which satisfies the resonance condition with zero optical phase shift.

Electrostatic Effects. The primary force acting on the electrons in a magnetic wiggler is the ponderomotive force produced by the interaction of the amplifier magnetic and laser electromagnetic fields. The axial component of the ponderomotive force is, to a good approximation, given by

$$F_z = e v_{\perp} \times B_r \quad (41)$$

where e is the charge on an electron. In addition to the ponderomotive force, there are also the forces of electrostatic repulsion of one electron by another. These space charge forces, while usually negligible, may affect the interaction of the electron beam with the laser beam in cases where the electron energy is low and the current density high. The perturbing effect of these electrostatic forces on the axial motion of the electrons is included in the computer simulation code by determining the electric field E_z acting on each electron due to the presence of other electrons in the beam. The axial force on each electron is then given by

$$F_z = e(v_{\perp} \times B_r + E_z) \quad (42)$$

To calculate E_z , we start with Maxwell's Equation

$$\nabla \cdot \vec{E} = \rho / \epsilon_0 \quad (43)$$

Since the laser field is assumed to be transverse and a function of z only, Equation (43) reduces to

$$\frac{\partial E_z}{\partial z} = \frac{\rho}{\epsilon_0} \quad (44)$$

We assume that both the electron density and electrostatic field are periodic functions, so that $\rho(z) = \rho(z + n\lambda_r)$ and $E_z(z) = E_z(z + n\lambda_r)$, where λ_r is the wavelength of the laser radiation. It is therefore only necessary to solve for the electrostatic field in a segment λ_r in length. Integrating Equation (44) we obtain

$$\int_0^z E_z dz' = \frac{1}{\epsilon_0} \int_0^z \rho dz'$$

or

$$E_z(z) - E_z(0) = \frac{1}{\epsilon_0} \int_0^z \rho dz' \quad (45)$$

If we assume a beam of infinite extent, the electrostatic field will equal zero when the charge distribution is uniform. The charge density we want in Equation (45) is, therefore, not the total charge density, but only the difference between the actual

charge density $\rho(z)$ and the average charge density ρ_0 .

$$E_z(z) - E_z(0) = \frac{\rho_0}{\epsilon_0} \int_0^z \left(\frac{\rho}{\rho_0} - 1 \right) dz' \quad (46)$$

By adding up all the charges between 0 and z ($0 \leq z \leq \lambda_r$), we can determine the net charge in the electric field in this interval, but do not know the value of $E_z(0)$. We do know, however, from Equation (46) that the net charge in electric field over a beam segment one radiation wavelength long is zero. We divide the interval λ_r into m equal segments Δz , determine the electric field at the end-point of each segment, and from these values, determine the average value of the electric field in the interval.

$$\sum_{z=\Delta z}^{m\Delta z} [E_z(z) - E_z(0)] = \frac{\rho_0}{\epsilon_0} \sum_{z=\Delta z}^{m\Delta z} \int_0^z \left(\frac{\rho}{\rho_0} - 1 \right) dz' \quad (47)$$

The first term on the left hand side of (7) averages out to zero so that

$$E_z(0) = -\frac{\rho_0}{\epsilon_0 m} \sum_{z=\Delta z}^{m\Delta z} \int_0^z \left(\frac{\rho}{\rho_0} - 1 \right) dz' \quad (48)$$

Substituting (48) into (46) gives an expression for $E_z(z)$.

In the computer program, the electric field is found at equally spaced intervals along a beam segment one radiation wavelength long, and the electric field E_z acting on a particle at any given position z is then found by interpolation.

When determining the electric field at any location z it is assumed that each particle in the simulation has its charge uniformly distributed over a region $0.01\lambda_r$ in length. A particle closer than $0.005\lambda_r$ to the point z will only contribute a fraction of its charge to the integral in Equation (46). In the computer code we compare the magnitude of the axial forces produced by the laser field and space charge to determine when space charge effects become important.

In optimized amplifier designs, the electrons may be tightly bunched so that the deviation of the electron density from its average value may be of the same order as the electron density itself. If $\rho = 2\rho_0$, the integrand in Equation (46) equals 1 and the value of the integral is z . The largest value z can have is λ_r , so an order of magnitude value for the peak electrostatic field in the interval is given by

$$E_z \approx \frac{\rho_0 \lambda_r}{\epsilon_0} \quad (49)$$

To neglect electrostatic forces we must, therefore, have

$$v_{\perp} B_r > \frac{\rho_0 \lambda_r}{\epsilon_0} \quad (50)$$

Since $v_{\perp} = \frac{\alpha c}{\gamma}$, $\alpha = \frac{e B_m \lambda_r}{2 \pi m c}$, and $B_r = \left(\frac{\mu_0}{c} \right)^{1/2} I^{1/2}$

$$n < \left(\frac{\mu_0}{c} \right)^{1/2} \frac{\epsilon_0}{\pi m} \frac{\gamma B_m I^{1/2}}{(1 + \alpha^2)}$$

or

$$n < 2.00 \times 10^{11} \frac{\gamma B_m I^{1/2}}{(1 + \alpha^2)} \quad (51)$$

where $n = \rho_0 / e$

and we have used the relation

$$\lambda_r = \frac{\lambda_m}{2\gamma^2} (1 + \alpha^2) \quad (52)$$

For a laser amplifier, the smallest value of the laser intensity I and the smallest value of B_m will occur at the input end of the amplifier. In general, bunching at the input end of the amplifier will also be weak so that the relation (51) using the initial values of B_m and I at the input end of the amplifier would place too low a limit on the electron density required to significantly perturb the system due to space charge effects. From computer simulations we have found that space charge effects produce only a few percent change in laser gain when the density is an order of magnitude higher than the limit given by (51) using initial values for the parameters on the right hand side of the relation. Space charge effects start to have a significant effect only when

$$n > 2 \times 10^{12} \frac{\gamma_0 B_{m0} I_0^{1/2}}{(1 + \alpha^2)} \quad (53)$$

For example, for an electron energy of 500 MeV ($\gamma = 979$), initial magnetic field of 0.1 Tesla, input laser intensity of 10^{12} W/m², and amplifier period $\lambda_m = 0.205$ m ($\alpha = 1.9$), the right hand side of relation (53) equals 4.25×10^{19} m⁻³. Space charge effects due to microscopic charge fluctuations would be small or negligible for electron densities below this value.

Energy Spread. We have made some preliminary runs including energy spread in the beam in a simplified manner. For an initial γ and a total energy spread $\Delta\gamma$, the particles were divided into 3

equal groups with γ 's symmetric about the initial γ . The percent full energy spread is defined as

$$PC = 100 \frac{\Delta\gamma}{\gamma} \quad (54)$$

Three of our previous optimizations with a monoenergetic beam were rerun to include values of PC up to 1 percent. The parameters for these runs were $\lambda = 0.5 \mu\text{m}$, $B_0 = 1 \text{ kG}$, $E_0 = 500 \text{ MeV}$ and sets (I_0, n) of $(10^{11}, 10^{17})$, $(10^{12}, 10^{18})$ and $(10^{13}, 10^{19})$ with I_0 in W/M^2 and n in m^{-3} . For the monoenergetic case, all the particles remained trapped throughout the entire length of the amplifier, except for $(I_0, n) = (10^{11}, 10^{17})$. In the latter case, a few particles moved outside the bucket but were subsequently trapped as the bucket grew.

With the introduction of energy spread, a fraction of the particles was initially outside the bucket. These particles were all eventually trapped if the initial energy spread was sufficiently small. It took about half the amplifier length to trap all the electrons. Under conditions where all electrons were

eventually trapped, the gain and the final fractional energy conversion (FEC) for each optimization were unchanged from the cases with initially monoenergetic electron distributions (Figure 4.2-13).

A summary of these preliminary results is given in Table 4.2-3. We define t/T as the ratio of the distance necessary to trap the particles to the total amplifier length. For $I_0 = 10^{13} \text{ W/m}^2$, $n_0 = 10^{19} \text{ elec/m}^3$, an initial energy spread of 0.8 percent was tolerable. An example of an optimization run showing particles outside the bucket which are then retrapped is shown in Figure 4.2-14. The initial energy spread for the run shown is 0.5 percent.

Table 4.2-3
Energy Spread for Trapped Particles

$I_0, n \text{ (W/m}^2, \text{elec/m}^3\text{)}$	$PC \text{ (\%)}$	t/T
$10^{11}, 10^{17}$	0.1	0.34
	0.3	0.69
$10^{12}, 10^{18}$	0.1	0.39
	0.4	0.51
	0.6	Not Trapped
$10^{13}, 10^{19}$	0.1	0.15
	0.5	0.37
	0.8	0.48
	1.0	Not Trapped

Fig. 4.2-13 Gain - FEC curves for FEL amplifier optimization runs without energy spread and with an energy spread of 0.8 percent. The final values for length, gain and fractional energy conversion are the same. The laser parameters are the same as those for Figure 4.2-15

Graphics Capabilities. The FEL simulation code results can be presented using computer-generated plots. "Snapshot" plots of the evolving system are produced at desired intervals as the equations are solved and particles traced through the amplifier. Several additional plots are generated after completion of each amplifier run.

Figure 4.2-15 is an example of a snapshot plot. It shows the axial force on the electrons (including the space charge force), the density distribution, the

Fig. 4.2-14 Phase space plots for an FEL amplifier simulation run with initial energy spread of 0.5 percent. The particles are finally trapped at about 85 m. The final FEC is 31 percent at 231 m, the same as for our optimization without energy spread.

Fig. 4.2-15 Sample of computer plot generated at desired points throughout an optimization run. This run contained an initial energy spread of 0.5 percent.

phase space particle distribution, and the energy distribution. Superimposed on the phase space distribution is the separatrix, or line separating the bound and unbound regions of phase space. The separatrix is calculated for the average γ of the particle ensemble. Also listed are values for some of the parameters at the position along the amplifier for which the plot is generated. This combination of plots has proved to be extremely useful in determining quickly the effects of parameter variation on the

optimizations.

Provisions also exist for storing on disc the velocity and position coordinates for all the particles at selected intervals as well as the gain, FEC, magnetic field, and magnet period and amplifier length. Auxiliary plotting programs can then generate plots of the phase space distributions (complete with the separatrices) (Figure 4.2-16a & b) or can generate the velocity position trajectory for any particular particle (Figure 4.2-17). One may

Fig. 4.2-16a Phase space evolution near the beginning of the FEL amplifier simulation. The initial beam, prebunching and early trapping are shown. This simulation has 100 particles assumed initially monoenergetic and spatially uniform.

Fig. 4.2-16b Phase-space evolution near the end of the simulation. This illustrates the point at which particles begin to escape (or leak) from the bucket.

Fig. 4.2-17 Plot of the trajectories of selected particles in the distribution. The initial position for the N^{th} particles is given by $Z(N) = N\lambda_r/100$ where λ_r is the laser wavelength. The trajectories are shown for only about half the amplifier length and the stationary bucket is drawn for the last time step.

also plot laser power gain, fractional energy conversion and the optical phase shift, ϕ , vs distance along the amplifier (Figure 4.2-18). Finally, Figure 4.2-19

gives the distance variation of the wiggler magnet field, B , the magnet period, LB , and the vector potential, α .

Fig. 4.2-18 Gain, FEC and ϕ vs length for a sample FEL amplifier simulation run.

Fig. 4.2-19 Wiggler magnet B-field, magnet period (LB) and the vector potential α vs length for a sample FEL amplifier simulation run.

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