Conf-910123--15

UCRL-JC-106277
PREPRINT

THEORY AND SIMULATION OF RAMAN SCATTERING IN INTENSE, SHORT PULSE LASER PLASMA INTERACTIONS

S. C. Wilks
W. L. Kruer
A. B. Langdon
P. Amendt
D. C. Eder
C. J. Keane

This paper was prepared for the Proceedings of the SPIE--High Power Lasers Meeting held in Los Angeles, CA, January 20-25, 1991.

February 1, 1991

Lawrence Livermore National Laboratory

UCRL-JC--106277

DE91 008734

Theory and simulation of Raman scattering in intense, short pulse laser plasma interactions

S.C. Wilks, W.L. Kruer, A.B. Langdon, P. Amendt, D.C. Eder, C.J. Keane

University of California, Lawrence Livermore National Laboratory
P.O. Box 808, Livermore, CA 94550

ABSTRACT

We report on the stimulated Raman scattering (SRS) that occurs when short ($\sim$ 50-100 femtoseconds), intense (I $\sim 2 \times 10^{17}$ W/cm^2) laser pulses are focused into a plasma, or a neutral gas which quickly becomes a plasma due to multiphoton ionization. Review of the usual SRS growth rates will be given, followed by a discussion on how this theory is modified for short pulses. It is found that for some recombination x-ray laser schemes employing short pulses, the Raman instability can potentially heat the electrons to levels that would greatly reduce the efficiency of such a device. However, there are areas in parameter space where the temperature of heated electrons can be kept to acceptable levels. One dimensional, relativistic, particle-in-cell (PIC) simulations are presented.

1. INTRODUCTION

Some recently proposed recombination x-ray laser (XRL) schemes employ short (< 1 pS), intense ($> 2 \times 10^{17}$ W/cm^2) laser pulses.[1] These schemes usually work as follows. The pulse is focused into a gas, for instance Neon, and the intensity of the pulse is chosen such that multiphoton ionization strips each atom down to a Li or He-like ion. Because of the fact that under many circumstances multiphoton ionization leaves the stripped electrons "cold" ($\sim 10-20$ eV), it is thought that the region just behind the laser pulse would be an ideal environment for population inversion between certain energy levels of the plasma ions. In particular, it may be possible to reach into the x-ray spectrum by exploiting transitions such as the $4d_{5/2}$–$3p_{3/2}$ in Lithium-like Neon. This would result in the possibility of lasing at 98 Å. In order for lasing to occur with reasonable efficiency ($\sim 10^{-5}$), it has been found that the plasma density must be in the range 10^{20} cm^{-3} $< n_e <$ 10^{21} cm^{-3}, and the plasma temperature must be kept below 60 eV. We will show that the Raman backscatter instability will have ample time to develop, causing serious heating of the plasma for many regions in the above-mentioned parameter space.

2. BACKGROUND AND THEORY

There are at least four effects present in this recombination laser method that cause the electron temperature to increase. For instance, the electron will gain energy from the ponderomotive potential which arises due to the transverse profile of the laser.[2] However, since these electrons will gain energy in the transverse direction, they may leave the plasma channel created by the pump laser pulse and leave only cold electrons. An additional source of heat is that the electrons may not give up all of their quiver energy (i.e., its oscillation in the electric field of the laser) to the laser pulse, as the laser passes. This effect can be minimized by choosing an appropriate laser intensity and pulse length.[3] Another systematic method of electron kinetic energy increase arises due to the fact that the electrons may actually be born at a point in the pump laser where the intensity has a large value (above-threshold -ionization). The electron will keep this energy, which can be quite large ($\sim$ 50 eV.) Perhaps the most important heating mechanism arises from collective effects, and it is this topic that we shall address in this paper. It will be shown that for even 50-100 femtosecond pulses, this effect can be substantial and can easily produce a large fraction ($\sim$ 30-50 %) of highly energetic electrons ($\sim$ 1-10 keV.) However, with some modification, some short pulse recombination XRL schemes might avoid this heating.

We will now briefly review the stimulated Raman scatter (SRS) instability. Much theoretical work has already been done on this subject, and we will start out with a review of the work done on the SRS instabilities where a steady state interaction is assumed to have been achieved, and the pump is continuous over all space. Most of this work will follow the discussion in Kruer's book[4]. (More in depth work on these instabilities can be found in papers by Estabrook[5] and Forslund[6].) It is well known that the temporal growth rate for the Raman instability can be written as

$$\gamma_0 = \frac{k v_{osc}}{4}\sqrt{\frac{\omega_p^2}{\omega_{ek}(\omega_0 - \omega_{ek})}}, \tag{1}$$

where ω_0 is the frequency of the incident (pump) laser, v_{osc} is the quiver velocity of an electron in the laser field, and where the wave number $\mathbf{k}$ is given by the dispersion relation

$$(\omega_{ek} - \omega_0)^2 - (\mathbf{k} - \mathbf{k}_0)^2 c^2 - \omega_p^2 = 0. \tag{2}$$

Here, the frequency ω_{ek} of the plasma wave is given by the dispersion relation

$$\omega_{ek}^2 = \omega_p^2 + 3k^2 v_e^2, \tag{3}$$

where v_e is the thermal velocity of the background plasma. Let us concentrate on the case of low density plasmas ($n/n_{crit} < 0.05$). For backscatter, i.e. $k \simeq 2k_0$) it is apparent that the growth rate is a maximum. For direct forward scatter ($k \simeq \omega_p/c$), where the growth rate is a minimum, both the upshifted and downshifted scattered light waves are now nearly resonant. However, as we shall see, the growth rate for forward scatter is, in general, too long to grow to appreciable values for these short pulses. The potential does exist for near-forward scatter to become important,[7] but this will be dealt with in a future study. For now, we concentrate on the backscatter, whose growth rate can be written as

$$\gamma_0 = \frac{1}{2}\frac{v_{osc}}{c}\sqrt{\frac{\omega_p/\omega_0}{(1 - \omega_p/\omega_0)}}\,\omega_0. \tag{4}$$

This can easily allow 10 or more e-foldings of the instability to occur, even for a 50 fSec long pulse.

When the finite length, and the intensity profile, of the pump is taken into account, it is found that the growth rate is somewhat reduced. At the very least, we need to average the growth rate over the intensity profile. We take the pump profile to be $I(t) = I_0 f(t)$, where I_0 is the maximum intensity of the laser, and $f(t)$ is some function describing the temporal profile of the laser pulse. For example, $f(t)$ may be $\mathrm{sech}^2 t$, or $\exp(-t^2/2\tau_{pulse}^2)$ for a Gaussian profile. The intensity is related to the quiver velocity via the formula $v_{osc} = 25.6\sqrt{I}\lambda_0$ [cm/sec], where I is in [W/cm^2] and λ_0, the laser frequency, is given in microns. For our simulations, we chose

$$f(t) = [1 - \frac{4t^2}{\tau_{pulse}^2}]^2 \tag{5}$$

for times between $-\tau_{pulse}/2 < t < \tau_{pulse}/2$. Putting this into the above equation for the effective growth rate and integrating this value over the pulse length, we obtain

$$\gamma_{eff} = \frac{2}{3}\gamma_0 \tag{6}$$

which tells us that the effective growth rate of the instability is $\frac{2}{3}$ of what the growth rate at maximum intensity would be, where the maximum intensity and an infinitely long pulse is assumed.

A more important effect has to do with the finite time any given electron is driven. In particular, a given electron experiences the pump field for a finite amount of time of about $\tau_{pulse} \simeq l/c$, where l is the full width of the pump wave packet. Intuitively, we expect that if $\tau_{pulse} \gg 1/\gamma_0$, the homogeneous theory gives a reasonable approximation for the growth rate. When τ_{pulse} becomes comparable to $1/\gamma_0$, the homogeneous theory over estimates the growth rate. Indeed, consider a simple estimate. The finite pulse length corresponds to an effective bandwidth $\Delta\omega \sim \frac{2\pi}{\tau_{pulse}}$. The growth rate is reduced[9] when $\Delta\omega > \gamma_0$; i.e., when $\gamma_0 \tau_{pulse} \leq 2\pi$. As we will shortly see, our simulations support this qualitative picture and help quantify the reduction in the growth rate as a function of $\gamma_0 l/c$.

In terms of how much heating will occur for a given pulse length, it is important to know the level the instability will attain before the pump passes this small region in space. For an infinitely long pump, the plasma waves (as well as the scattered light waves) will grow until they nonlinearly saturate, usually by trapping large numbers of electrons which take energy from the wave. This results in the particles gaining kinetic energy, which translates to an increase in temperature of the plasma. When the instability is strongly driven, the temperature of this hot tail of heated electrons, generated during the Raman backscatter instability can be approximated as

$$T_{hot} \sim \frac{1}{2} m_e v_\phi^2 \tag{7}$$

where $m_e = 511$ keV/c^2 is the mass of an electron, and $v_\phi = \omega_p/2k_0$ is the phase velocity of the plasma wave. This is usually around 1-10 keV for our parameters. However, if the pump is finite in extent, the plasma waves may never reach nonlinear saturation, but instead grow up to some lower level and stop. This explains why the temperature of the plasma is dependent on the pulse length times the growth rate, for levels below the nonlinearly saturated values. The shorter and less intense the pulse is, the less the plasma will be heated. This is because the instability can only continue growing if the driver (in this case the pump laser) is present. Since the intensity of the backscattered electromagnetic waves grow like $I_{SRS} = I_{noise} \exp 2\gamma t$, we can get an idea of the level of backscatter attained for a pulse of length τ_{pulse} by noting that the maximum backscatter will be $I_{SRS} = I_{noise} \exp 2\gamma\tau_{pulse}$. Therefore, a good measure of how much heating is obtained for a given set of parameters is the product $\gamma_{eff}\tau_{pulse}$. The goal of this paper is to get some quantitative feeling for when the heating due to backscatter becomes too severe for the recombination X-ray laser schemes to work efficiently. We will now describe some results obtained from the simulations.

3. SIMULATION RESULTS

In this section we will first describe the basic set-up for the simulations, and then discuss the results in detail. The simulations were performed with ZOHAR[10] , a relativistic particle-in-cell code written by A.B. Langdon and B.F. Lasinski. ZOHAR was run, where one spatial dimension (the x, or longitudinal dimension) and two velocities (the longitudinal v_x, and the transverse v_z, which is the direction of the laser electric field) were included. In all cases, the laser was injected from the left boundary, and was absorbed at the right. The laser was a $\frac{1}{4}$ micron wavelength with a parabolic profile (for v_{osc}/c). The pulse length (approximately FWHM) was varied between 50 to 100 fSec. The maximum intensity of the laser was taken to be either 2×10^{17} or 2×10^{18} W/cm^2 during all runs. The plasma density varied between 10^{20} and 1.5×10^{21} cm^{-3}. The plasma was taken to be a Maxwellian at $t = 0$, with an initial temperature of $T_e = 25$ eV. The final temperature of the plasma was measured after the pump wave exited the system. This was because the oscillatory motion of the electrons in the laser electric field (transverse to the propagation direction of the laser) was so large that in the presence of the pulse, it

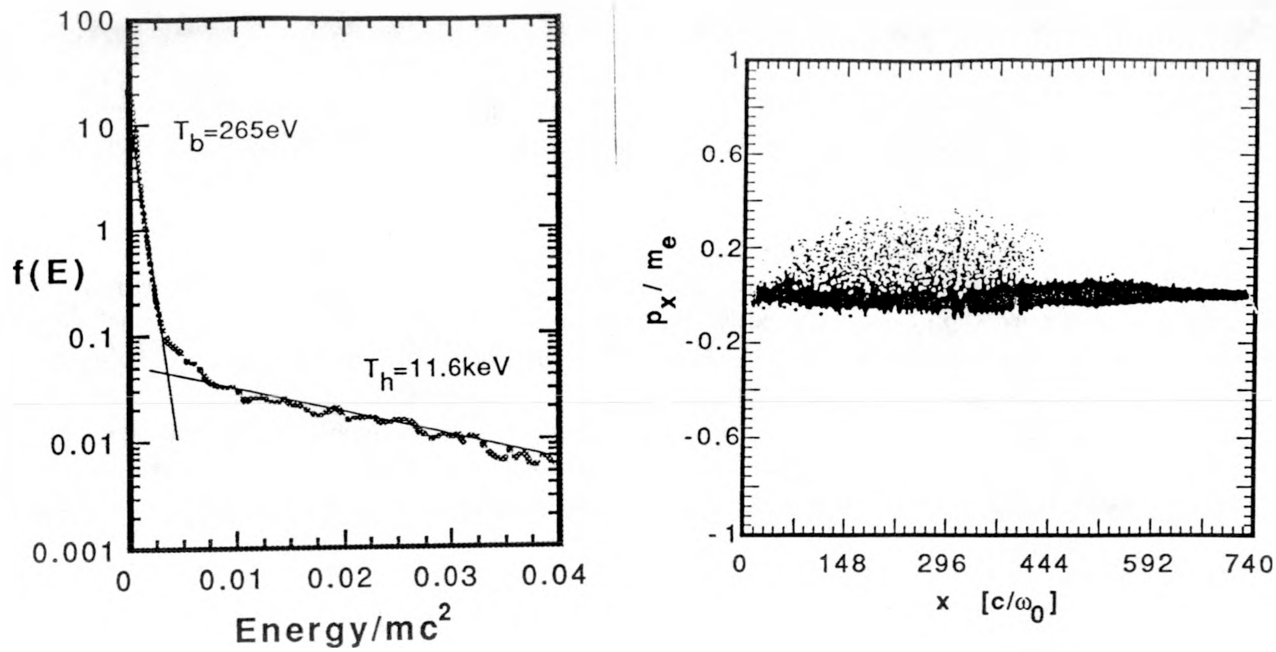

Figure 1: (a) Energy distribution plot of plasma after pump laser has passed through the system. The plasma density is 5×10^{20} cm^{-3}, the maximum intensity of the pump laser is I=2×10^{17} W/cm^2, and the pulse length is τ_{pulse} = 100 fSec. (b) The longitudinal phase space of the simulation described in (a).

dominated the energy distribution. However, the majority of this energy was given back to the laser as it passed through the system. The heating associated with this transverse motion (self-consistently included in the simulation) is negligible, as usual, for these schemes.

Fig. 1a shows a typical electron energy distribution plot that is obtained for every run. The parameters are given in the figure caption. Note that in this case, the final background electron temperature reached 265 eV, while a hot ($\sim$ 10 keV) tail formed as well. Fig. 1b shows the actual longitudinal phase space for this distribution. Notice the gradual growth of the instability (in space) which starts from the right boundary and increases as one travels left. This substantiates the idea that the final temperature depends on the product $\gamma_{eff}\tau_{pulse}$, since a much shorter pulse would have produced a lower temperature. In fact, this pulse was 100 fSec long. When an identical run was made, only this time the pulse length was 50 fSec, a final plasma temperature of 30 eV was measured.

Further evidence of this type of behavior can be seen in Fig. 2, which shows plasma temperature (in eV), normalized backscattered electromagnetic wave intensity, and energy gained by the plasma as functions of the product $\gamma_{eff}\tau_{pulse}$. Note that as the product $\gamma_{eff}\tau_{pulse}$ increases, so do the plasma parameters. This correlation, the fact that the various parameters increase in similar ways, points out that the heating is due to low levels of stimulated backscatter and that a good measure of this is the $\gamma_{eff}\tau_{pulse}$ product. Fig. 3 shows the final plasma temperature, as a function of $\gamma_{eff}\tau_{pulse}$, in detail for the temperature range of interest for recombination x-ray lasers. From this graph, we conclude that if the $\gamma_{eff}\tau_{pulse}$ product can be kept less than about 12, the plasma temperature may stay low enough for lasing to occur. This is probably a conservative estimate, since the noise level will be lower in the experiments. However, the acceptable $\gamma_{eff}\tau_{pulse}$ is only logarithmically sensitive to the noise level, so the difference between simulation and experiment should not be too great.

Fig. 4 shows a graph of the energy efficiency for a Li-like Ne laser versus plasma density (see P. Amendt *et al.*, these proceedings for details on how this plot was obtained). The 3 points on the lines are simulation results that were obtained by choosing the plasma densities, running the code for

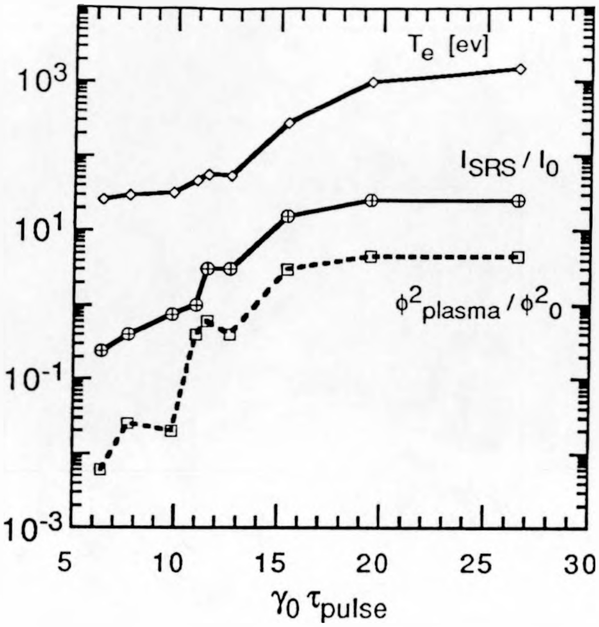

Figure 2: Various plasma parameters measured and plotted versus the product of the growth rate times the pulse width. The correlation between the amount of backscatter and the temperature of the plasma confirms that the heating mechanism is Raman backscatter.

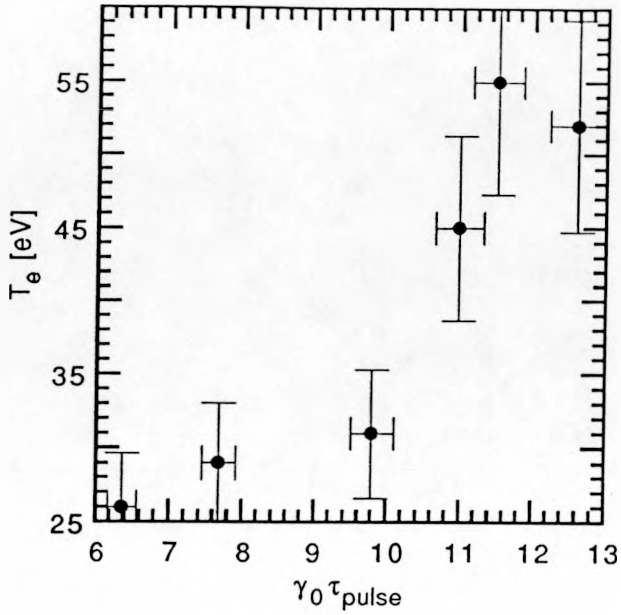

Figure 3: Plot of the plasma temperature versus $\gamma_{eff}\tau_{pulse}$, where the temperature is in the range of interest for the recombination x-ray laser, showing that some simulations gave acceptable levels of heating.

10,000 time steps, then measuring the plasma temperature for the 3 cases. The laser had the following parameters: I=2 $\times$ 10^{17} and a pulse length of 50 fSec. These runs show that even when the collective effects of the plasma are present (i.e., parametric instabilities are included, such as Raman backscatter) x-ray efficiencies on the order of 10^{-5} are theoretically possible.

Finally, Fig. 5 shows a graph of the growth rates measured from the simulations (normalized by γ_{eff}) versus $\gamma_{eff}\tau_{pulse}$. These values were obtained by measuring the increase of the electrostatic energy ($\propto (\nabla\phi)^2$) as a function of time. For the shortest pulses studied, measuring the growth of the electrostatic energy was complicated by the fact that it takes some time for the growth, due to the instability, to well up out of the noise initially present in the simulation. Since the growth rate may not be constant for the front of the pulse, (in fact, it is probably lower than the homogeneous growth rate) we are measuring only that growth rate that is associated with the entire length of the pulse. As expected, for large values of $\gamma_{eff}\tau_{pulse}$, γ_{sim} approaches γ_{eff}. However, for small $\gamma_{eff}\tau_{pulse}$, we find a smaller growth rate, which seems to be linear up to about a $\gamma_{eff}\tau_{pulse}$ of 15. This is qualitatively as expected, as discussed in the previous section.

4. CONCLUSIONS

We have reported on a series of 1-D computer simulations that explored the longitudinal heating of electrons during the passage of ultrashort ($\sim 50-100$ fSec), intense ($\sim 10^{17}-10^{18}$ W/cm^2) laser pulses due to Raman backscatter. However, there seems to be a region of plasma density-laser intensity parameter space that may allow for these schemes to work, once the driver laser technology has progressed to a sufficient point. The simulations have stressed the need for shorter wavelength pumps ($\leq 1/4$ micron) with pulse widths of ≤ 100 fSec, in order for the plasma temperature to remain low. This is an important consideration because for recombination x-ray lasers employing short pulse

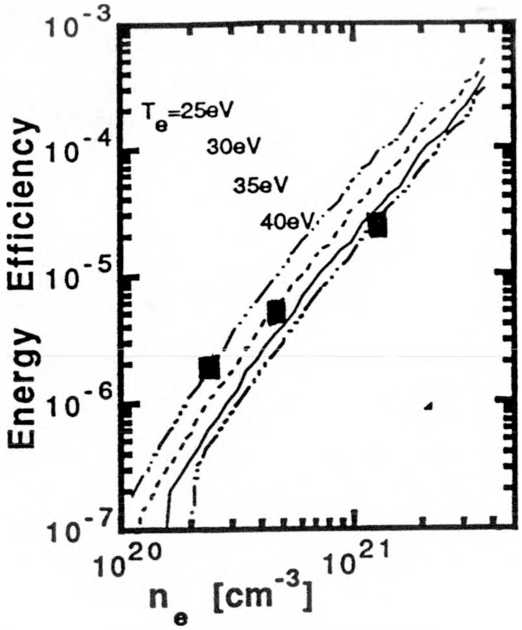

Figure 4: Energy efficiency as a function of density for the Li-like Ne laser discussed in Ref. 1. Here, the energy efficiency is defined as the energy deposited into the short x-ray pulse divided by the energy in the incident pump laser pulse.

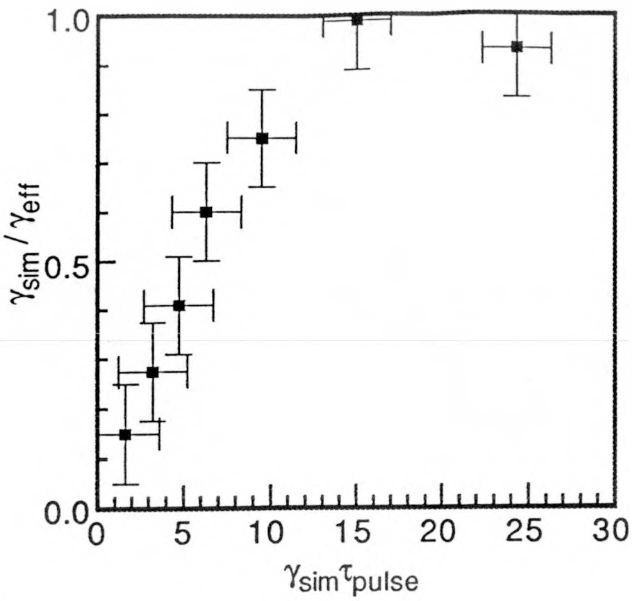

Figure 5: Plot of growth rate measured from the simulations (γ_{sim}, normalized by γ_{eff}) versus product $\gamma_{eff}\tau_{pulse}$.

drivers, heating greatly decreases the efficiency of the lasing medium. In fact, we have found from simulation results that the product of the growth rate and the pulse length should not exceed about 10-15. This will ensure reasonable levels for the plasma temperature ($\leq$ 60 eV) once the pump laser has propagated through the plasma.

5. ACKNOWLEDGMENTS

We wish to thank E.A. Williams and K. Estabrook for useful discussions on this subject. This work was performed under the auspices of the U.S. Department of Energy by the Lawrence Livermore National Laboratory under contract number W-7405-ENG-48.

6. REFERENCES

1. P. Amendt *et al.*, "Optically ionized plasma recombination x-ray lasers," these proceedings.

2. P.B. Corkum, N.H. Burnett and F. Brunel, "Multiphoton Ionization in Large Ponderomotive Potentials," to be published in **Atoms in Intense Radiation Fields**, edited by M. Gavrila (Academic Press Inc., Harcourt Brace Janovich, Orlando, Florida, 1991).

3. B. Penetrante and J.N. Bardsely, submitted to *Phys. Rev. A* (1990).

4. W.L. Kruer, **The Physics of Laser Plasma Interactions**, (Addison-Wesley, New York, 1988).

5. K. Estabrook and W.L. Kruer, *Phys. Fluids* **26**(7), 1892 (1983); K. Estabrook, W.L. Kruer, and B.F. Lasinski, *Phys. Rev. Lett.*, **45**(17), 1399 (1980).

6. D.W. Forslund, J.M. Kindel, and E.L. Lindman, *Phys. Rev. Lett.*, **30**, 739 (1973).

7. S.C. Wilks, *et al.*, "Theory and Simulation of Raman Scatter at Near Forward Angles," in preparation.

8. K. Nishikawa and C.S. Liu, "General Formalism of Parametric Excitation"; in **Advances in Plasma Physics**, Vol. 6, (A. Simon and W. Thomson, eds.), p. 3-81. Wiley, New York, 1976.

9. J.J. Thomson and J.I. Karush, *Phys. Fluids*, **17**, 1608 (1974).

10. A.B. Langdon and B.F. Lasinski, in **Methods in Computational Physics**, edited by J. Killeen *et al.* (Academic, New York) Vol 16, 327, (1976).

Technical Information Department • Lawrence Livermore National Laboratory
University of California • Livermore, California 94550