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ABSTRACT

The target chamber of a laser fusion reactor will contain small amounts of background gases. As the beam is focused, it ionizes the gas and Raman scattering is induced. Density limits on the background gas are found in order that the laser beam will not become appreciably decollimated. It is found that laser bandwidth efficiently decreases the scattering effect.

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Recently, Sparks and Sen published an analysis of forward Raman scatter in a laser fusion target chamber[1] that indicated that the incident laser beam could be appreciably decollimated. Unfortunately their analysis was not correct, as pointed out by Eimerl,[2] being appropriate for a turbulent medium and not for a phase-coherent process such as Raman scattering. Here, I analyze this process from a classical standpoint, considering both forward and backscatter, and obtain limits on the background gas density so that only a small percentage of the incoming beam is sufficiently decollimated that it misses the target.

In the following, although the analysis is fairly general, it will be helpful to keep in mind the following typical parameters: light wavelength $\lambda = 1.06$ μm, with corresponding critical density $n_c = 10^{21}$ cm^{-3}; background gas density $n = 10^{15}$ cm^{-3} (corresponding to the vapor pressure of a lithium wall at ~ 1000 K), with temperature T ~ 1-10 eV; the ionizing intensity is $I_i \simeq 10^{12}$ W/cm^2; intensity on target is $I_0 \sim 10^{15}$ W/cm^2; the target diameter is d ~ 500 μm; the lens f-number is f ~ 3.

The differential equations describing the coupling between the incident and scattered waves and the plasma are[3]

$$\left[\frac{\delta^2}{\delta t^2} + \nu\frac{\delta}{\delta t} - 3\,V_T^2\nabla^2 + \omega_p^2\right]\frac{\tilde{n}}{n} = \frac{e^2}{m^2c^2}\,\nabla^2\,(\underset{\sim}{A}_o \cdot \underset{\sim}{A}_s) \tag{1}$$

$$\left[\frac{\delta^2}{\delta t^2} + \nu_A\frac{\delta}{\delta t} - c^2\nabla^2 + \omega_p^2\right]\underset{\sim}{A}_s = -\,\omega_p^2\,\frac{\tilde{n}}{n}\,\underset{\sim}{A}_o \tag{2}$$

where the subscripts o and s refer to the incident and scattered waves, A is the vector potential, $\tilde{n}$ the fluctuating part of the density, ω_p the plasma frequency, and (ν_A) is the plasma (electromagnetic) wave damping decrement. Pump depletion is neglected here, since if it is important, too much laser power has already been lost to scattering.

The dispersion relation is obtained by Fourier-transforming Eqs. (1) and (2) and using $\omega_o^2 = c^2 k_o^2 + \omega_p^2$, where ω_o is the incident beam frequency and k_o is its wave vector.

$$\omega^2 + i\omega\nu - \omega_p^2 - 3k^2V_T^2 = \omega_p^2k^2v_{os}^2\cos\phi\left\{\left[2\omega\omega_o + \omega^2 - c^2(k^2 + 2kk_o\cos\theta) + i(\omega+\omega_o)\nu_A\right]^{-1} + \left[-2\omega\omega_o + \omega^2 - c^2(k^2 - 2kk_o\cos\theta) + i(\omega-\omega_o)\nu_A\right]^{-1}\right\} \tag{3}$$

where V_{os} is the oscillating velocity $\frac{eA_o}{m}$, θ is the angle between $\underset{\sim}{A}_o$ and $\underset{\sim}{A}_s$ and Θ is the angle between $\underset{\sim}{k}$ and $\underset{\sim}{k}_o$. The dependence on θ is the usual polarization dependence of Raman scattering. Henceforth, we stay in the polarization plane, and let $\theta = o$.

If $\Theta \neq o$, only the second term inside the curly brackets is resonant. This is a 3-wave interaction, photon $\rightarrow$ photon + plasmon. If $\Theta = o$ (i.e. direct forward scatter) both terms may be resonant. This is a 4-wave process, and corresponds to Sparks and Sen's "instantaneous scattering process". This case represents a modulational instability with $k = \omega_p/c$ and a growth rate of $\lambda = \frac{1}{\sqrt{2}} \omega_p^2/\omega_o v_{os}/c$ (hardly infinite, as claimed in Ref. 1). It does not scatter light out of the beam, but does produce density perturbations that may affect the beam propagation through the plasma because of refraction. It also produces a frequency bandwidth in the laser beam.[3]

Now consider $\Theta = 0$ (the resonance width is $\Delta\Theta \simeq \nu_\Lambda/\omega_p < n/n_c = 10^{-6}$). The second term in the curly brackets is resonant if

$$c^2(k^2 - 2kk_o \cos \Theta) = \omega^2 - 2\omega\omega_o$$

or

$$k = k_o \cos \Theta \pm \sqrt{k_o^2 \cos^2 \Theta - 2\frac{\omega\omega_o}{c^2} + \frac{\omega^2}{c^2}} \qquad (4)$$

The plus sign refers to backscatter and the negative sign to forward scatter. We may neglect ω^2 in comparison to $2\omega\omega_o$, since they are in the ratio $\omega_p/\omega_o << 1$. I shall show that, even though back scatter growth rates are higher than those for forward scatter, the latter instability is more efficient in scattering laser energy Thus I consider forward scatter first.

The maximum value of k occurs for $\cos \Theta = \sqrt{2\omega/\omega_o}$ and is

$$k_{max} = \sqrt{2\omega_p\omega_o}\,/c \qquad (5)$$

the corresponding scattering angle is

$$\theta_s \cong \frac{k}{k_o} = \sqrt{2\omega_p/\omega_o} \qquad (6)$$

The threshold is

$$\nu\nu_A = \frac{\omega_p}{\omega_o} k^2 V_{os}^{\,2} \qquad (7)$$

The growth rate well above threshold is

$$\gamma = \frac{1}{2} \sqrt{\frac{\omega_p}{\omega_o}}\; k\, V_{os} \qquad (8)$$

Since $\gamma \sim k$, we find that the maximum growth rate occurs for k_{max} and is

$$\gamma_{max} = \frac{1}{\sqrt{2}}\, \omega_p V_{os}/c \qquad (9)$$

with the threshold

$$\nu\nu_A = 2\omega_p^{\,2} \frac{V_{os}^{\,2}}{c^2} \qquad (10)$$

I now show that this instability is always strongly over driven. The plasma absorption frequency is certainly less than ω_p. Thus

$$V_{os}^2/c^2 \Big|_{thres} \lesssim \frac{\nu_A}{\omega_p} = \frac{Z}{2000}\left(\frac{n}{n_c}\right)^{3/2} \frac{1}{T(keV)^{3/2}} \tag{11}$$

Let the ion charge number $Z = 5$, $n/n_c = 10^{-6}$, $T = 10^{-2}$kev then the threshold $V_{os}/c = 10^{-4}$, corresponding to an intensity of $\sim 10^{10}$ W/cm^2. Since we are interested in $I \gtrsim I_i = 10^{12}$ W/cm^2, the intensity is always well above threshold intensity. This conclusion also applies to the back scattered instability, since its threshold is lower than the forward scattered instability.

Next, I consider whether the instability is absolute or convective. The phase velocities of the scattered photon and plasmon are in the same direction. Thus the instability is absolute if[4]

$$2\gamma = \sqrt{2}\,\omega_p V_{os}/c > \sqrt{\frac{V_T c}{2}}\left(\frac{\nu}{V_T} + \frac{\nu_A}{c}\right) \approx \frac{1}{2}\sqrt{\frac{c}{V_T}}\,\nu$$

or

$$V_{os}/c > \frac{1}{4\sqrt{2}}\sqrt{\frac{c}{V_T}}\left[10^{-9}\frac{n^{1/2}}{T(ev)^{3/2}} + 8\sqrt{2\pi}\left(\frac{\omega_p}{\omega_o}\frac{c^2}{V_T^{\,2}}\right)^{3/2}\exp -\left(\frac{\omega_p}{\omega_o}\frac{c^2}{V_T^{\,2}}\right)\right] \qquad (12)$$

where V_T is the electron thermal velocity. The first term in the square brackets is collisional damping, and the second is Landau damping. For $n = 10^{15}$ and $T = 10$ ev, this condition becomes $V_{os}/c > 2.5 \times 10^{-3}$ or $I \gtrsim 5 \times 10^{12}$ W/cm^2. Thus there may be regions of absolute and of convective instability.

If the instability is absolute, its saturation level will be determined by particle trapping in the waves[5]. If it is convective, its saturation level will be determined by the minimum of the convective result ($\sim \exp(qIx)$) and the particle trapping level. Thus it is conservative to estimate the saturation level everywhere by the particle trapping limit, in order to estimate the total energy lost from the beam.

The particle trapping limit for the plasma waves is[6]

$$\frac{E_p^{\,2}}{4\pi nmV_p^{\,2}} = 1 + 2\beta^{1/2} - 8/3\beta^{1/4} - \beta/3 \qquad (13)$$

where E_p is the amplitude of the plasma waves, V_p is the phase velocity $\omega_p/k = \sqrt{\frac{\omega_p}{2\omega_o}}c$ and $\beta \equiv 3V_T^2/V_p^2$. evaluate Eq. (12), the temperature must be estimated. Consider heating due to inverse bremsstrahlung. The rate is approximately

$$\nu_{\epsilon} = 10^{12}\left(\frac{n}{10^{21}}\right)^2 \left(\frac{T}{1000}\right)^{-3/2}$$

$$= 5 \times 10^3 \text{ sec}^{-1}$$

So the time required for significant heating is $\tau \sim 2 \times 10^{-4}$ sec, much too long for this problem.

The plasma is locally heated, however, by the plasma waves, at a rate given by $2\gamma \frac{E_p^2}{4\pi}$. The number of photons lost by the main beam is the same as the number of plasmons created,

$$\frac{\Delta E_o^2}{4\pi\ \omega_o} = \frac{\Delta E_p^2}{4\pi\ \omega_p}$$

Thus the energy lost from the main beam per unit time is

$$\nu^* \frac{E_o^2}{8\pi} = 2\gamma\frac{\omega_o}{\omega_p}\frac{E_p^2}{4\pi} \qquad (14)$$

where ν^* is a phenomenological damping rate defined by this equation. The maximum energy flux that the electrons can carry away is 3/2 nm V_t^3. Setting this equal to the absorbed energy flux gives an expression for V_T.

$$\frac{\nu^*}{\omega_o} \cdot \frac{E_o^2}{8\pi} c = 2\frac{\gamma}{\omega_p} \cdot \frac{E_p^2 c}{4\pi} = \frac{3}{2} nmV_T^3 \tag{15}$$

Eq. (15) assumes a near-Maxwellian distribution function, which is often the case for turbulent wave breaking heating of plasmas.

Eq. (1) may be written in terms of β:

$$\frac{E_p^2}{4\pi nm\, V_p^2} = \frac{1}{12\sqrt{3}} \sqrt{\frac{\omega_p}{\omega_o}}\; c/V_{os}\; \beta^{3/2} \tag{16}$$

Substituting Eq. (16) into Eq. (13) gives an expression for β in terms of $\sqrt{\frac{\omega_p}{\omega_o}}\, c/V_{os}$. In the regime of interest, $\beta \sim 1$; i.e. T $\sim$ 10 ev. Eq. (13) then gives

$$\frac{E_p^2}{4\pi} = nmV_p^2 = nmc^2\, \omega_p/2\omega_o \tag{17}$$

Next I determine the total fractional energy lost over a ray path. This is

$$f_E = 1 - \exp\left[\int_\ell^L \frac{\nu^*}{c}\, dx\right]$$

Distances are measured from beam focus. L is the ionization point, I (L) = 10^{12} ω/cm^2; ℓ is the closest distance at which a scattered photon still hits the target.

$$\ell = d\left(f + \frac{1}{2\phi_s}\right)$$

$$= d\left(f + \frac{1}{2\sqrt{2}}\sqrt{\frac{\omega_o}{\omega_p}}\right) \qquad (18)$$

See Fig. (1).

Using Eqs. (17) and (15),

$$\int_\ell^L \frac{\nu^* dx}{c} = \frac{\sqrt{2}}{2}\left(\frac{\omega_p}{\omega_o}\right)^3 \int_\ell^L \frac{k_o dx}{V_{os}/c} \qquad (19)$$

A factor of 1/2 comes from the polarization dependence. Since the beam is focused, $V_{os}(x) \sim 1/x$. Thus relating $V_{os}(x)$ to the target plane x = fd,

$$\int_\ell^L \frac{\nu^* dx}{c} = \frac{1}{2\sqrt{2}}\left(\frac{\omega_p}{\omega_o}\right)^3 \frac{c}{V_{os}(fd)}\, k\, fd \left[\left(\frac{L}{fd}\right)^2 - \left(\frac{\ell}{fd}\right)^2\right] \qquad (20)$$

But $L^2 >> \ell^2$, and $(\ell/fd)^2 = I_o/I(L) = I_o/10^{12}$.
Also, V_{os} $(fd)/c = (\omega_{Nd}/\omega_o)$ $(I_o/10^{18})^{1/2}$, where ω_{Nd} is the Nd laser frequency. Eq. (2) can thus be written

$$\int_{\ell}^{L} \frac{\nu^* dx}{c} = \frac{10^3}{2\sqrt{2}} \left(\frac{n}{n_c}\right)^{3/2} \frac{\omega_o}{\omega_{Nd}} \left(\frac{I_o}{10^{12}}\right)^{1/2} k_o fd \qquad (21)$$

For the parameters of our example, Eq. (21) predicts an scattering of about 10%. This is probably a conservative overestimate, but does show that the instability may be of consequence, and that $n \lesssim 10^{15}$ is an upper limit to the allowable density for these parameters.

Note that the frequency scaling is $\sim \omega_{Nd}/\omega_o$, so that CO_2 lasers are much more susceptible, whereas higher frequency lasers are less susceptible to forward Raman scattering.

Next consider back or side scatter. The absorption fraction $\int \frac{\nu^* dx}{c} \sim \gamma E_p^2 \sim \gamma/k^2$. However, $\gamma \sim k$ giving $\int \frac{\nu^* dx}{c} \sim \frac{1}{k}$. Back and side scatter have $k \gtrsim k_o$. Thus the absorption due to these instabilities is in the ratio

$$\frac{k_{forward}}{k_o} \cong \sqrt{\frac{\omega_p}{\omega_o}} << 1 \qquad (22)$$

So we see that forward scatter is more important in the present case, due to the higher plasma wave saturation level.

Finally, consider the effects of finite bandwidth. Since $\omega_p \ll \omega_o$, it is easy to have a laser bandwidth comparable to the plasma frequency. suppose that $\gamma \ll \Delta\omega \ll \omega_p$. Then the effects of finite bandwidth are well known.[7] The growth rate becomes

$$\gamma = \frac{1}{2} \frac{\omega_p^2}{\Delta\omega} \frac{V_{os}^2}{c^2} \tag{23}$$

In order for the instability to have reasonable growth, it is necessary that $\gamma\tau \gg 1$, where τ is the laser pulse duration. Thus bandwidth will essentially eliminate the instability if

$$\gamma\tau = \frac{1}{2} \frac{\omega_p^2}{\Delta\omega} \tau \frac{V_{os}^2}{c^2} = 1 \quad , \tag{24}$$

requiring a bandwidth of

$$\frac{\Delta\omega}{\omega_o} = \frac{1}{2} \frac{\omega_p^2}{\omega_o^2} \frac{V_{os}^2}{c^2} \omega_o\tau \quad . \tag{25}$$

For I $\sim 10^{15}$ and $\tau = 10^{-9}$ sec this requires $\frac{\Delta\omega}{\omega_0} = 10^{-3}$, which is not unreasonably large.

In conclusion, I find that forward Raman scattering is effective in decollimating a coherent laser beam, requiring that $n/n_c < 10^{-6}$. However, a small amount of laser bandwidth will efficiently reduce the instability.

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FIGURE CAPTION

Figure 1: Geometry of focused beam on target.

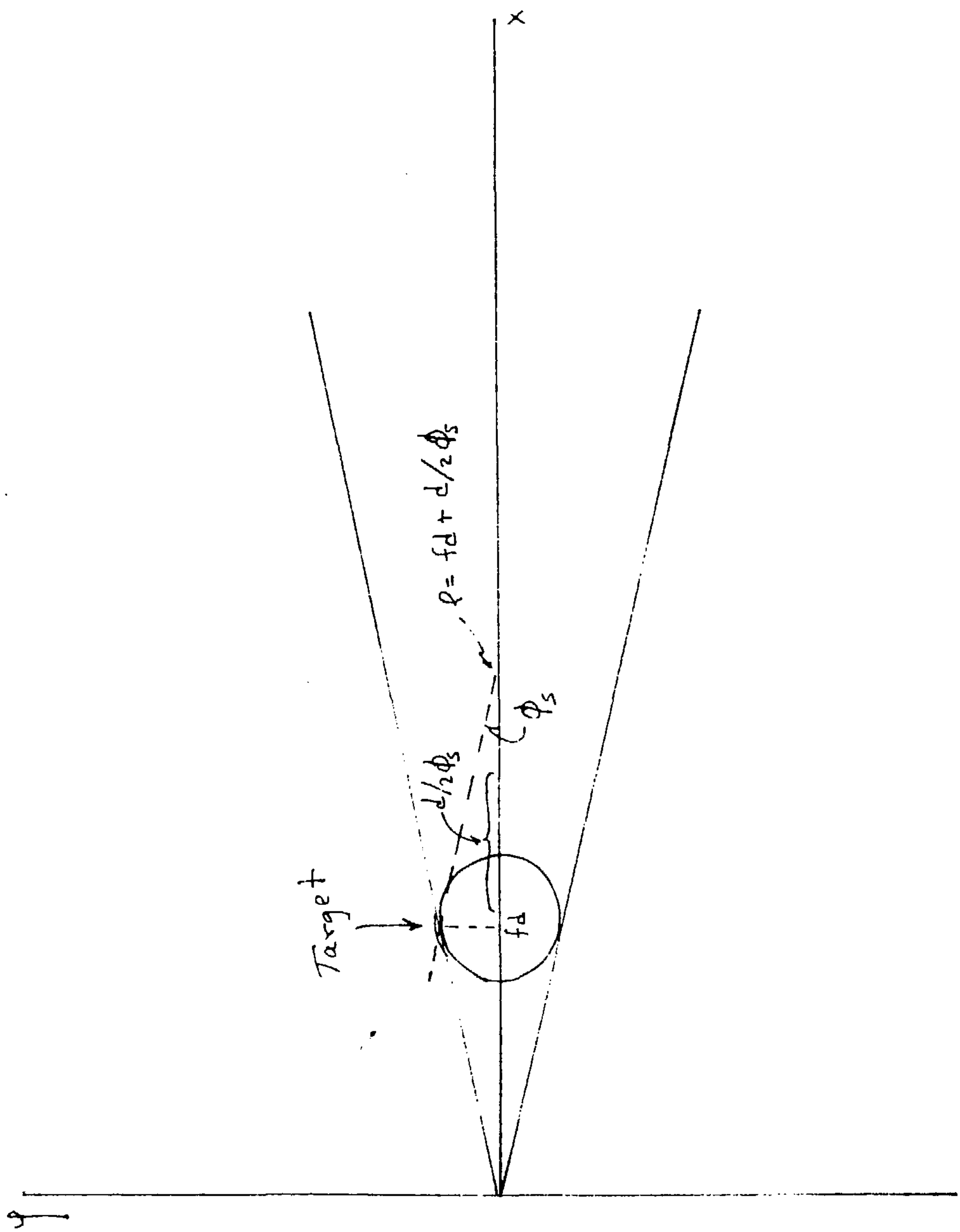

Fig. 1

REFERENCES

1. M. Sparks and P. Sen, Phys. Rev. Lett., 39, 751 (1977).
2. D. Eimerl, to be published.
3. D. W. Forslund, J.M. Kindel, and E.L. Lindman, Phys, Fluids, 18, 1002 (1975).
4. K. Nishikawa, in Adv. In Plasma Physics, Vol. 6, A. Simon and W.B. Thompson, ed., John Wiley and Sons (New York, 1976).
5. W.L. Kruer, ibid.
6. T. Coffey, Phys. Fluids 14, 1402 (1971).
7. J.J. Thomson and Jack I. Karush, Phys. Fluids, 17, 1608 (1974).