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2.B Filamentation of Short Wavelength Laser Beams in Plasmas

A plane electromagnetic beam propagating in a plasma has been shown[1,2,3] to be unstable against modulation in the direction perpendicular to the direction of propagation. This instability causes the beam to break up into narrow filaments of high intensity and low plasma density and is known as the filamentation instability. Such an effect may be important in laser fusion since non-uniformities in illumination at the critical surface may produce non-uniform ablation pressures which will prevent the symmetric implosion of the target. The increased intensity in a filament can drive other plasma instabilities such as stimulated Raman scattering and the two-plasmon instability, these instabilities produce fast electrons which may preheat the fuel.

In a laser produced plasma two mechanisms that can cause the laser beam to form filaments are the ponderomotive force[1] and thermal forces.[4] If the plasma is collisionless the mechanism responsible for filamentation is the ponderomotive force on the electron fluid. This force is proportional to $-\nabla <E^2>/\omega^2$, where E is the electric field of the wave of frequency ω and <> denotes time average. The net effect of the ponderomotive force is to push electrons away from regions of highest intensity. However, if the plasma is collisional then inverse bremsstrahlung heating becomes important and filamentation due to thermal effects[4] can occur if the mean free path of an electron, λ_{mfp}, is smaller than the filament width. In regions of higher intensity inverse bremsstrahlung leads to a temperature increase which in turn causes a hydrodynamic expansion. In both cases there results a reduction in plasma density where the intensity is greatest. Since the index of refraction, $N = (1-4\pi e^2 n_o / m_e \omega_0^2)^{1/2}$, of a plasma depends on the local electron density, n_o, a decrease in plasma density increases the index of refraction and decreases the phase velocity, $v_p = c/N$, causing the wave fronts to curve in such a way that the beam is focused toward the original region of enhanced intensity becoming even more intense.

As the wavelength of the beam decreases the radiation can penetrate to higher densities where the plasma is more collisional and for a given intensity the ponderomotive force decreases.

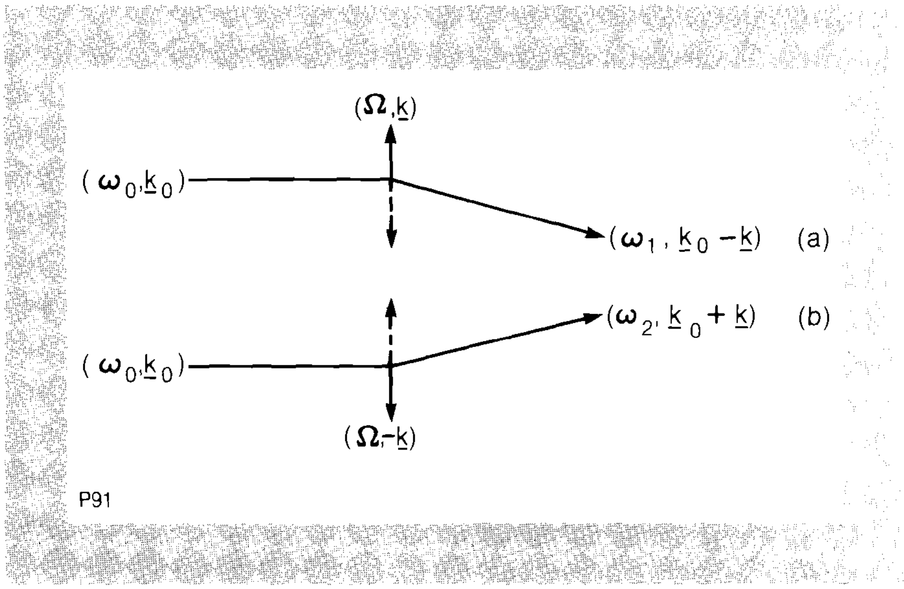

Fig. 15
Configuration of wave vectors for filamentation (a) (ω_o,$\underline{k}_o$) the incident electromagnetic wave, (ω_1, $\underline{k}_o$-$\underline{k}$) the Stokes electromagnetic wave and (Ω, $\underline{k}$) a density perturbation; (b) (w_2, $\underline{k}_o$+$\underline{k}$) the anti-Stokes electromagnetic wave and (Ω,$-\underline{k}$) a density perturbation. The beating of an incident electromagnetic wave with a density perturbation leads to the generation of Stokes and anti-Stokes electromagnetic waves which grow exponentially if the incident wave intensity is above a threshold value.

Therefore, there will be a regime where both filamentation mechanisms must be considered together.

The filamentation instability is basically a four wave process and is illustrated in Fig. 15. (ω_0, $\underline{k}_0$) is the frequency and wavenumber of the initial plane elecromagnetic pump wave which beats with a density perturbation with frequency Ω and wavenumber $\underline{k}$ to produce two sidebands, the Stokes wave with frequency ω_1 $\omega_0-\Omega$ and wavenumber $\underline{k}_1 = \underline{k}_0-\underline{k}$ and the anti-Stokes wave with frequency $\omega_2 = \omega_0+\Omega$ and wavenumber $\underline{k}_2 = \underline{k}_0+\underline{k}$.

The configuration of the problem is taken to be the following. The initial electromagnetic wave propagates in the x-direction and is polarized in the z-direction and the density perturbation is in the y-direction. The Stokes and anti-Stokes waves are also polarized in the z-direction.

Using a two fluid model together with Maxwell's equations the following equations for the electromagnetic fields and density perturbation can be obtained:

$$c^2\nabla^2\underline{E} - \frac{\partial^2 \underline{E}}{\partial t^2} - \omega_{pe}{}^2\underline{E} = i4\pi e\omega\tilde{n}\underline{v} + 4\pi n_o e\nu_e\underline{v} \quad (1)$$

$$\tilde{n} = i n_o e < (\underline{v}\mathrm{x}\underline{B})_y > / k_y k_B T_e - n_o \tilde{T}_e/T_e \quad (2)$$

where c is the velocity of light, ω_{pe} is the plasma frequency, e the charge of a proton, v the electron velocity field due to all the electromagnetic waves, ω the frequency of the electromagnetic field, n_o the equilibrium density, k_y the wavenumber of the density perturbation in the y direction, k_B the Boltzmann constant, T_e the electron temperature, B the magnetic field of the waves, υ_e the electron ion collision frequency and $\tilde{T}_e$ the perturbation of the electron temperature. The first term on the R.H.S. of Eq. (2) represents the ponderomotive force while the second term is due to thermal effects. We have treated the density perturbation as a driven response and neglected ion inertia. Balancing the wave heating with the thermal conduction term in the energy equation for the electrons and neglecting all other terms allows us to obtain an expression for the electron temperature perturbation $\tilde{T}_e$ in terms of the electromagnetic fields

$$\tilde{T}_e = -e\mathrm{Re} < \underline{E} \cdot \underline{v} > / \kappa k_y^2 \quad (3)$$

where Re means real part and κ is the electron thermal conductivity.

We now write the pump, Stokes and anti-Stokes waves as

$$E_j(x,t) = \tfrac{1}{2} \epsilon_j(x,t)\, e^{i(k_j \cdot x - \omega_j t)} + c \cdot c,\ j = 0,1,2$$

where the slow amplitude variation $\epsilon_j(x,t)$ is determined by the non-linear interaction and linear dissipation and the phase factor is due to linear dispersion. Using a perturbation procedure on equations (1)—(3) we obtain the following equations for the Stokes and anti-Stokes waves

$$\left(\frac{\partial}{\partial t} + v_1 \cdot \frac{\partial}{\partial x} + \gamma_T \right) \epsilon_1(x,t) = i\Gamma \left(|\epsilon_0|^2 \epsilon_1 + \epsilon_0^2 \epsilon_2^* e^{-i2\delta t} \right)$$
$$\left(\frac{\partial}{\partial t} + v_2 \cdot \frac{\partial}{\partial x} + \gamma_T \right) \epsilon_2(x,t) = i\Gamma \left(|\epsilon_0|^2 \epsilon_2 + \epsilon_0^2 \epsilon_1^* e^{-i2\delta t} \right) \quad (4)$$

where v_1 and v_2 are the group velocities, γ_T is the damping,

$$\delta = \omega_o - \omega_1,\ \Gamma = \frac{e^2 \omega_{pe}^2}{8 m_e^2 \omega_0^2 \omega_1 v_{Te}^2} \left(1 + \frac{1}{2.5 k_y^2 \lambda_{mfp}^2}\right)$$

and v_{Te} is the electron thermal velocity. The first term in the coupling coefficient comes from the ponderomotive force while the second is due to thermal effects. The ratio between them is a comparison between the strength of the ponderomotive force and thermal force, this ratio is $2.5 k_y^2 \lambda_{mfp}^2$. Thermal effects can dominate over ponderomotive force effects whenever λ_{mfp} is less than the filament width.

The spatial growth rate in the direction of propagation $k_{||}$ and threshold can be obtained by solving equation (4) for ϵ_0 = constant. The maximum spatial growth rate being given by

$$\left(\frac{k_{||}}{k_0} \right)_{max} = (1/2 \left[\beta (\beta + 4c^2 / 5\omega_0^2 \lambda_{mfp}^2) \right]^{1/2} - \gamma_T/\omega_o)/(1 - \omega_{pe}^2/\omega_0^2)^{1/2} \quad (5)$$

where

$$\beta = \omega_{pe}^2 v_0^2 / 4\omega_0^2 v_{Te}^2 \text{ and } v_0, = \left| \frac{eE_o}{m_e \omega_o} \right| ,$$

is the quiver velocity in the pump field, maximum growth occurs for $k_y = \omega_o \beta^{1/2} c$. The threshold in terms of v_0 is given by

$$\left(\frac{v_o}{v_{Te}} \right)^2_{threshold} = \frac{8\omega_0^2}{\omega_{pe}^2} \left(\left[(0.4/k_y^2 \lambda_{mfp}^2)^2 + (\gamma_T/\omega_0)^2 \right]^{1/2} - 0.4/k_y^2 \lambda_{mfp}^2 \right) \quad (6)$$

Following the treatment of Rosenbluth,[5] Eq. (4) can be modified by the addition of a spatial mismatch, to include the effects of plasma inhomogeneity. We find that the effect of a density gradient on filamentation is just to limit the size of the growth

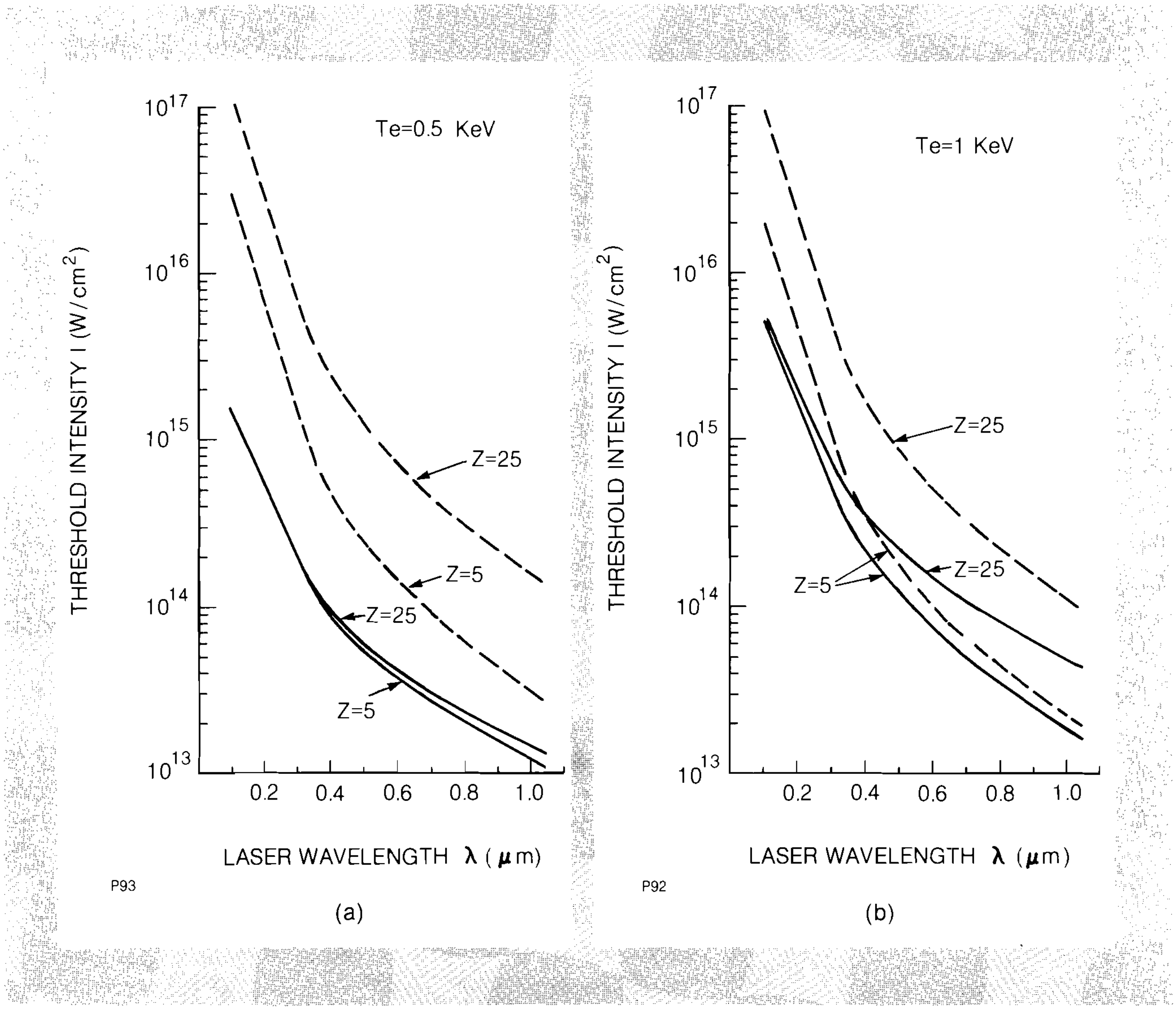

Fig. 16
Laser intensity I(Watts/cm²) for the onset of the filamentation instability as a function of laser wavelength in plasmas at quarter critical with different charge states Z and different temperatures. The solid lines represent the case when both the ponderomotive force and thermal effects are included in the theory. The dashed lines represent the case for the ponderomotive force only.

region. This is in contrast to three wave resonant processes such as stimulated Raman scattering and stimulated Brillouin scattering where a detuning of the resonance occurs. To fully understand the effects of a density gradient on the filamentation instability other effects such as the dependence of the coupling coefficient and damping on x must be taken into account.

The laser threshold intensities for the onset of the filamentation instability as a function of laser wavelength for different charge states Z are plotted in Fig. 16a, 16b. The solid lines represent the case when both the ponderomotive force and thermal effect are included in the equation for the threshold, the dashed lines represent the case when the thermal effect is neglected. For low temperature plasmas the thermal effect is dominant especially for short wavelengths and high Z materials. At higher temperatures, $T_e \sim 1$ KeV, the thermal effect is still dominant at short wavelengths while the ponderomotive force becomes more important at longer wavelengths and low Z materials.

In summary we have shown that for short wavelength lasers and plasma temperatures of about 1 KeV the filamentation mechanism is dominated by the thermal force rather than the ponderomotive force.

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