Laser – Bernstein-mode coupling in a laser-produced plasma

Jetendra Parashar and H. D. Pandey
Center for Energy Studies, Indian Institute of Technology, New Delhi 110016, India

V. K. Tripathi
Physics Department, Indian Institute of Technology, Hauzkhas, New Delhi 110016, India
(Received 16 June 1995)

Laser radiation propagating in a laser-produced plasma, in the presence of a self-generated magnetic field B_S can decay into an electron Bernstein wave and a sideband Langmuir wave. Near the critical density, the laser may be linearly mode converted into a Langmuir wave, which may parametrically excite an electron Bernstein wave and a Langmuir wave. Inhomogeneity in the B_S field determines the convective threshold for resonant and quasimode decay processes. In a plasma with electron temperature $T_e \sim 0.5$ keV and a B_S field scale length of ~ 100 μm the power threshold at a 1-μm wavelength turns out to be 10^{15} W/cm^2. Above the threshold the growth rate is of the order of 3×10^{11} s^{-1} for a resonant decay and 10^{11} s^{-1} for a nonresonant quasimode decay.

PACS number(s): 52.40.Nk, 52.35.Hr, 52.40.Db, 52.35.Mw

I. INTRODUCTION

Strong self-generated magnetic fields have been observed in laser-produced plasmas [1–3]. The generation of these fields has been attributed to a variety of mechanisms, e.g., nonzero $\vec{\nabla}n \times \vec{\nabla}T$, where $\vec{\nabla}n$ and $\vec{\nabla}T$ are density and temperature gradients, respectively [4], thermal Weibel instability [5], pondermotive force [6–8], and electromagnetic two stream instability [9]. The fields are nonuniform and their magnitude increases with the intensity of the laser.

The existence of self-generated magnetic field in a laser-produced plasma allows a variety of collective modes of space charge oscillations, e.g., Bernstein modes, lower hybrid mode, upper hybrid mode, etc. These modes may be driven parametrically unstable by the laser and may play an important role in laser energy absorption and heat transport. Tripathi and Sharma [10] have examined three wave parametric decay instability of laser radiation into lower hybrid and upper hybrid modes, ion acoustic and upper hybrid modes, fast ion wave and plasma wave, in a uniform magnetoplasma. Sharma [11] has studied the resonant decay instability of plasma waves into electron Bernstein waves. Simon [12] has recently studied the parametric excitation of Bernstein modes in a laser-produced plasma where sideband waves are taken as electromagnetic modes. This process has been proposed as a probable scenario for observed high reflectivity of laser light from underdense plasmas. Baton *et al.* [13] have recently observed enhancement of reflectivity of 1 μm, 1-ps pulse from a tenth critical density plasma, from a value of 10^{-4} at 5×10^{14} W/cm^2 to 0.1 at 10^{16} W/cm^2. For intensities between 10^{16} and 10^{17} W/cm^2 a saturation level of 10% is established. The theoretical interpretation of data in terms of stimulated Brillouin scattering (SBS) requires nonthermal noise, orders of magnitude higher than the thermal noise level. Similar difficulties exist in the interpretation of data from experiments with 8–10 ps and nanosecond pulses [14,15]. Further, the Thompson scattering spectra of waves in the ion acoustic range of frequency reveals that high-amplitude waves exist at frequencies ω, up to $\omega \sim 2\omega_a$ where ω_a is the usual ion acoustic frequency generated in a SBS process, $\omega_a \sim 2k_0c_s \approx 2\omega_0 c_s/c$.

Simon suggested that some of these observations may be attributed to the excitation of Bernstein waves. The power threshold for laser decay into a Bernstein wave and an electromagnetic wave sideband is seen to be lower than that for SBS. At 1 MG magnetic field the Bernstein mode frequency (of the order of electron cyclotron frequency) is an order of magnitude higher than that of the ion acoustic wave produced in a SBS process. However, Simon has ignored the inhomogeneity of the magnetic field, which may cause much stronger detuning of the three wave decay process involving a Bernstein wave than a density gradient. This will enhance the power threshold for decay into a Bernstein wave. Parametric decay into a heavily damped Bernstein quasimode may be an important channel in such a situation. Recently Labaune *et al.* [16] have reported exciting experimental results on large-amplitude ion acoustic waves produced via SBS in an underdense laser-produced plasma. At intensities $I > 2\times10^{13}$ W/cm^2 they observed, in addition to an ion acoustic mode, an electrostatic mode with frequency $\omega \sim 10\omega_a$. The frequency ω increases with laser intensity. It is probable that this mode is a Bernstein wave whose frequency depends on electron cyclotron frequency, hence, on plasma temperature and laser intensity.

In this paper we estimate power thresholds for three wave parametric decay of laser radiation into a Bernstein modes and quasimodes due to inhomogeneities in plasma density and ambient magnetic field. In a dense plasma, near the critical layer, we examine the parametric decay of a linearly mode converted Langmuir wave into a Bernstein mode/quasimode and a Langmuir wave sideband. It may also undergo a four wave parametric process.

In Sec. II we study the parametric processes in the underdense region. In Sec. III we study the parametric instability of linearly mode converted Langmuir wave, including finite wave number of the pump wave. A discussion of results is given in Sec. IV.

II. DECAY IN THE UNDERDENSE REGION

Consider the propagation of a laser pump wave through an underdense plasma of electron density n_0^0 and electron temperature T_e, having a dc magnetic field $\vec{B}_S \| \hat{z}$,

$$\vec{E}_0 = \hat{y} E_0 \exp[-i(\omega_0 t - k_0 x)], \quad \vec{B}_0 = \frac{c\vec{k}_0 \times \vec{E}_0}{\omega_0}, \tag{1}$$

where $k_0 = (\omega_0/c)(1-\omega_p^2/\omega_0^2)^{1/2}$, $\omega_p = (4\pi n_0^0 e^2/m)^{1/2}$, $-e$ and m are electronic charge and mass, respectively, and we have assumed $\omega_c \equiv (eB_s/mc) \ll \omega_0$. The pump wave provides an oscillatory velocity to electrons, $\vec{v}_0 = e\vec{E}_0/mi\omega_0$, that couples a low-frequency electron Bernstein mode with electrostatic potential $\phi = \phi \exp[-i(\omega t - \vec{k}\cdot\vec{x})]$, and an electromagnetic wave sideband with electric field $\vec{E}_1 = \vec{E}_1 \exp[-i(\omega_1 t - \vec{k}_1 \vec{x})]$. The linear dispersion relations for the decay waves are

$$\omega = \omega_c \left[1 + \frac{I_1(b)e^{-b}}{1 - I_0(b)e^{-b} + k^2 v_{\text{th}}^2/2\omega_p^2} \right] \approx \omega_c [1 + I_1(b)e^{-b}], \tag{2}$$

and

$$\omega_1^2 = k_1^2 c^2 + \omega_p^2, \tag{3}$$

respectively, where $b = k^2 v_{\text{th}}^2/2\omega_c^2$. The Bernstein mode in the case of quasimode decay is heavily cyclotron damped. The phase-matching conditions $\omega_1 = \omega - \omega_0$, $\vec{k}_1 = \vec{k} - \vec{k}_0$ give

$$k = \frac{1}{c}[(\omega_0^2 - \omega_1^2) - 2\omega_p^2 - 2\sqrt{(\omega_0 - \omega_p)(\omega_1 - \omega_p)} \cos\theta_1]^{1/2}, \tag{4}$$

here θ_1 is the angle between $\vec{k}_1$ and $\vec{k}_0$. The pump and the sideband waves exert a ponderomotive force $\vec{F}_p \equiv e\vec{\nabla}\phi_p$ on the electrons at $(\omega, \vec{k})$, where $\phi_p = -\vec{v}_0 \cdot \vec{E}_1 / 2i\omega_1$. The ponderomotive force and the self-consistent fields at $(\omega, \vec{k})$ produce an electron density perturbation n,

$$n = \frac{k^2}{4\pi e} \chi_e (\phi + \phi_p),$$

where χ_e is the electron susceptibility at $\omega, \vec{k}$:

$$\chi_e = \frac{2\omega_p^2}{k^2 v_{\text{th}}^2} \left[1 + \frac{\omega}{k_z v_{\text{th}}} \sum_s Z \left[\frac{\omega - s\omega_c}{k_z v_{\text{th}}} \right] I_s(b) e^{-b} \right]. \tag{5}$$

In the following we shall restrict ourselves to $\omega \sim \omega_c$ hence only $s=0$, $s=1$ terms of the summation need to be retained. The ion density perturbation could be written as $n_i = -(k^2/4\pi e)\chi_i \phi$, where $\chi_i = -\omega_{pi}^2/\omega^2$ is the ion susceptibility. Substituting for n and n_i in Poisson's equation, we obtain

$$\epsilon\phi = -\chi_e \phi_p, \tag{6}$$

where $\epsilon = 1 + \chi_e + \chi_i$. The density perturbation at $(\omega, \vec{k})$ couples with $\vec{v}_0$ to produce nonlinear current density $\vec{J}_1^{\text{NL}} = -\frac{1}{2} ne\vec{v}_0^*$. When $\vec{J}_1^{\text{NL}}$ is used in the wave equation one obtains

$$D_1 E_1 = -\frac{i\omega_1(1+\chi_i)}{2} k^2 v_0^* \phi, \tag{7}$$

where $D_1 = \omega_1^2 - \omega_p^2 - k_1^2 c^2$ and we have assumed that $\vec{k}_1 \cdot \vec{E}_0 \approx 0$. Using Eqs. (6) and (7) we obtain the nonlinear dispersion relation

$$\epsilon D_1 = \chi_e (1+\chi_i) k^2 |v_0|^2 / 4. \tag{8}$$

We solve Eq. (8) in two cases.

A. Resonant decay

Expressing $\omega = \omega_{1r} + i\gamma$ where ω_r is the frequency at which $\epsilon = 0$ and D_1 at $\omega_{1r} = \omega_r - \omega_0$ is simultaneously zero, one obtains the growth rate

$$\gamma^2 = -\frac{(1+\chi_i)^2 k^2 |v_0|^2}{4\partial\epsilon/\partial\omega \partial D_1/\partial\omega_1} \approx \frac{(1+\chi_i)^2 k^2 |v_0|^2}{8\omega_0 \omega_p^2} \omega_c^3 b I_1 e^{-b}. \tag{9}$$

For 1.06 μm, 10^{16} W/cm^6 Nd:Glass laser, $n_0^0 = 10^{20}$ cm^{-3}, $T_e = 0.5$ K eV, $B_s \sim 1$ MG the growth rate turns out to be $\gamma \sim 3 \times 10^{11}$ s^{-1}. Figures 1 and 2 show the variation of growth rate with density and magnetic field, respectively. Growth rate increases with magnetic field but decreases as we move towards critical density.

B. Resistive quasimode decay

If one allows a finite k_z then the Bernstein wave could be cyclotron damped for $\omega \sim \omega_c \simeq k_z v_{\text{th}}$, and it goes over to a quasimode. In that case the growth rate can be ob-

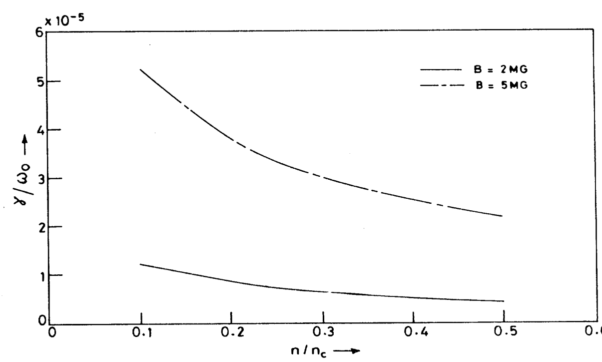

FIG. 1. Normalized growth rate vs plasma density for resonant decay in underdense plasma.

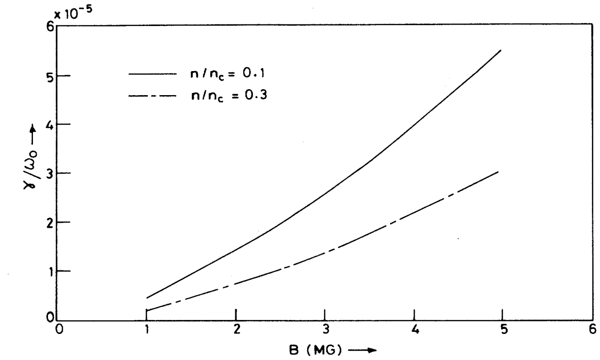

FIG. 2. Normalized growth rate vs magnetic field for resonant decay in underdense plasma.

tained from Eq. (8) by expanding $\omega=\omega+i\gamma$, $\omega_1=\omega_{1r}+i\gamma$ where $\omega_{1r}=\omega-\omega_0$ is the zero of D_1,

$$\gamma=\left[\frac{k^2|v_0|^2(1+\chi_i)^2}{4\partial D_1/\partial\omega_1}\right]\mathrm{Im}\left[-\frac{1}{\epsilon}\right]=\frac{k^2|v_0|^2(1+\chi_i)^2}{16k_0c}\frac{k^2v_{\mathrm{th}}^2F}{2\omega_p^2}, \tag{10}$$

where

$$F=\frac{2\omega_p^2}{k^2v_{\mathrm{th}}^2}\frac{\epsilon_i}{\epsilon_r^2+\epsilon_i^2},$$

$$\epsilon_i=\frac{2\omega_p^2}{k^2v_{\mathrm{th}}^2}\frac{\omega I_1e^{-b}}{k_zv_{\mathrm{th}}}2\sqrt{\pi}\exp\left[-\frac{\omega-\omega_c}{k_zv_{\mathrm{th}}}\right]^2$$

and

$$\epsilon_r=1+\frac{2\omega_p^2}{k^2v_{\mathrm{th}}^2}\left[1-I_0e^{-b}+\frac{\sqrt{\pi}\omega I_1e^{-b}}{k_zv_{\mathrm{th}}}\times\exp\left(-\frac{\omega-\omega_c}{k_zv_{\mathrm{th}}}\right)^2\right].$$

For $\omega-\omega_c\approx k_zv_{\mathrm{th}}$,

$$F\approx\frac{0.6\omega I_1e^{-b}/(\omega-\omega_c)}{[(1-0.3\omega I_1e^{-b}/\omega_1-\omega_c)^2+(0.6\omega I_1e^{-b}/\omega-\omega_c)^2]}.$$

For the parameters similar to the resonant decay case we get growth rate $\gamma\sim10^{11}\ \mathrm{s}^{-1}$. The growth rate scales as $\gamma\sim T_e^{1/4}$, with plasma temperature T_e.

C. Effect of plasma inhomogeneity

The plasma inhomogeneity localizes the region of parametric interaction and introduces convection losses. Let the density n_0^0 and B_S field profiles in the plasma be $n_0^0=N_0(1+x/L_n)$, $B_S=B_{S0}(1+x/L_B)$. The threshold condition for resonant decay instability can be written as

$$\frac{\gamma^2}{|v_{gx}v_{g1x}(\partial/\partial x)(k_x-k_{1x}-k_{0x})|}\geq1,$$

which gives

$$\frac{|v_0|_{\mathrm{TH}}^2}{c^2}\geq\frac{2}{k_0L_BbI_1e^{-b}}.$$

For 1.06 μm, Nd:glass laser, $L_B\sim100\ \mu$m, $T_e=0.5$ keV, $B_S\sim1$ MG, $n_0^0\sim10^{20}\ \mathrm{cm}^{-3}$ the threshold power P_{TH} is $1.5\times10^{15}\ \mathrm{W/cm^2}$. The threshold power $P_{\mathrm{TH}}\sim\omega_c$ and $P_{\mathrm{TH}}\sim1/T_e^{1/2}$.

In the case of quasimode decay ω_c is a function of x, the condition for heavy cyclotron damping $\omega-\omega_c(x)=k_zv_{\mathrm{th}}$ is satisfied only locally. As one moves away from this region $\mathrm{Im}(-1/\epsilon)$ decreases rapidly, with typical half-width $\sim x_0$, and the instability threshold may be written as

$$\gamma x_0/v_{g1x}\geq1$$

or

$$\frac{|v_0|_{\mathrm{TH}}^2}{c^2}\geq\frac{1}{4}\frac{\omega_{c0}}{(\omega-\omega_c)k_0L_B}.$$

The threshold power minimized for $\omega-\omega_c\sim\omega_cI_1e^{-b}$. For 1.06 μm, Nd:glass laser, $L_B\sim100\ \mu$m, $B_S\sim1$ MG, $n_0^0\sim10^{20}\ \mathrm{cm}^{-3}$ threshold power $P_{\mathrm{TH}}\sim10^{16}\ \mathrm{W/cm^2}$.

III. DECAY NEAR THE CRITICAL LAYER

Now we examine the decay of a linearly mode converted Langmuir wave,

$$\phi_0=\phi_0e^{-i(\omega_0t-k_0x)}, \tag{11}$$

where $\omega_0^2=\omega_p^2+k_0^2v_{\mathrm{th}}^2$ and $v_{\mathrm{th}}=\sqrt{T_e/m}$, T_e is electron temperature. The pump wave provides an oscillatory velocity to electrons, $\vec{v}_0=e\vec{E}_0/mi\omega_0$, where $\vec{E}_0=-\vec{\nabla}\phi_0$, and couples a low-frequency electron Bernstein wave of potential $\phi=\phi\exp[-i(\omega t-\vec{k}\cdot\vec{x})]$, and a sideband Langmuir wave of potential $\phi_1=\phi_1\exp[-i(\omega_1t-\vec{k}_1\cdot\vec{x})]$ where $\omega_1=\omega_0-\omega$, $\vec{k}_1=\vec{k}_0-\vec{k}_1$, and

$$k=[(\omega_0^2-\omega_p^2)/c^2+(\omega_1^2-\omega_p^2)/v_{\mathrm{th}}^2-2\sqrt{(\omega_0-\omega_p)(\omega_1-\omega_p)}/(v_{\mathrm{th}}c)\cos\theta_1]^{1/2}.$$

The pump and the sideband waves exert a ponderomotive force $\vec{F}_p\equiv e\vec{\nabla}\phi_p$ on the electrons at $(\omega,\vec{k})$ where $\phi_p=(\vec{v}_0\cdot\vec{k}_1/2\omega_1)\phi_1$. The ponderomotive force and the self-consistent field at $(\omega,\vec{k})$ produce an electron density perturbation n and ion density perturbation n_i: $n=[k^2/(4\pi)]\chi_e(\phi+\phi_p)$, $n_i=-[k^2/(4\pi e)]\chi_i\phi$, respectively. Substituting for n in the Poisson equation we obtain

$$\epsilon\phi=-\chi_e\phi_p, \tag{12}$$

where $\epsilon=1+\chi_e+\chi_i$. Nonlinear electron density n_1^{NL} at $(\omega_1,\vec{k}_1)$ can be obtained from the equation of continuity

$$\frac{\partial n_1^{\rm NL}}{\partial t}+\vec{\nabla}\cdot\left[\frac{1}{2}n\vec{v}_0^*\right]=0\ ,$$

$$n_1^{\rm NL}=\frac{\vec{k}_1\cdot\vec{v}_0^*}{2\omega_1}\frac{k^2}{4\pi e}(\phi+\phi_p)\ .$$

Substituting $n_1^{\rm NL}$ in the Poisson's equation we get

$$\epsilon_1\phi_1=\frac{k^2}{k_1^2}\frac{\vec{k}_1\cdot\vec{v}_0^*}{2\omega_1}\phi\ , \qquad (13)$$

where $\epsilon_1=1-(\omega_p^2+k_1^2v_{\rm th}^2)/\omega_1^2$. Equations (12) and (13) yield the nonlinear dispersion relation

$$\epsilon=-\frac{\chi_e(1+\chi_i)k^2|v_0|^2}{4}\frac{1}{D_1}\ , \qquad (14)$$

where $D_1=\omega_1^2-(\omega_p^2+k_1^2v_{\rm th}^2)$. Had we considered the upper sideband $\omega_2,\vec{k}_2$ also (where $\omega_2=\omega+\omega_0$, $\vec{k}_2=\vec{k}+\vec{k}_0$) then the above equation would have $1/D_1$ on the right-hand side replaced by $1/D_1+1/D_2$ where $D_2=\omega_2^2-(\omega_p^2+k_2^2v_{\rm th}^2)$.

A. Resonant decay

Solving Eq. (14) around the simultaneous zeros of ε and D_1, one obtains the growth rate for resonant three wave decay,

$$\gamma^2=-\frac{\chi_e k^2|v_0|^2}{8\partial\epsilon/\partial\omega\partial\epsilon_1/\partial\omega_1}$$

$$=\frac{(1+\chi_i)^2k^2|v_0|^2}{8\omega_0\omega_p^2}\omega_c^3 bI_1e^{-b}\ . \qquad (15)$$

For $\omega_0=2\times10^{15}$ rad/s, $k_0\sim3\omega_0/c$, $v_0/c\sim0.1$ (this may correspond to laser power density $\sim10^{15}$ W/cm^2 if one assumes 10% of laser power conversion into the Langmuir wave), $n_0^0\sim10^{21}$ cm^{-3}, $T_e=0.5$ keV, $B_S\sim1$ MG, $\gamma\sim5\times10^{11}$ s^{-1}.

B. Resistive quasimode decay

In the case of a quasimode decay, the growth rate can be obtained from Eq. (14) by expanding $\omega=\omega+i\gamma$, $\omega_1=\omega_{1r}+i\gamma$ where $\omega_{1r}=\omega-\omega_0$ is the zero of D_1,

$$\gamma=\left[\frac{k^2|v_0|^2\chi_e^2}{4\partial D_1/\partial\omega_1}\right]{\rm Im}\left[-\frac{1}{\epsilon}\right]$$

$$=\frac{k^2|v_0|^2}{16\omega_p^2}\frac{k^2v_{\rm th}^2}{\omega_0}\chi_e^2F\ . \qquad (16)$$

For 1.06 μm, Nd:glass laser, $n_0^0\sim10^{21}$ cm^{-3}, $T_e=0.5$ keV, $B_S\sim1$ MG the growth rate turns out $\gamma\sim2\times10^{11}$ s^{-1}.

C. Effect of plasma inhomogeneity

Assuming the density n_0^0 and B_S field profile of the plasma be $n_0^0=N_0(1+x/L_n)$, $B_S=B_{S0}(1+x/L_B)$. Following a similar analysis to the preceding section, the threshold condition for resonant decay instability can be written as

$$\frac{|v_0|^2}{c^2}\geq\frac{4v_{\rm th}^3\omega_0}{\omega_c^2bL_Bc^2}\ . \qquad (17)$$

For 1.06 μm, Nd:glass laser, $n_0^0\sim10^{21}$ cm^{-3}, $T_e\sim0.5$ keV, $B_S\sim1$ MG, $L_B\sim100$ μm, the threshold power is $P_{\rm TH}=7\times10^{14}$ W/cm^2. $P_{\rm TH}\sim\omega_c$ and $P_{\rm TH}\sim1/T_e^{1/2}$. The threshold condition for the quasimode decay is

$$\frac{|v_0|_{\rm TH}^2}{c^2}\geq\frac{16\omega_{c0}k_0}{k^2k_zv_{\rm th}L_B}\ . \qquad (18)$$

For parameters similar to the case of resonant decay the threshold power is $P_{\rm TH}\sim10^{15}$ W/cm^2.

IV. DISCUSSION

In the presence of a 1 MG self-generated magnetic field of scale length $L_B\sim100$ μm a Nd:glass laser of power density $\geq7\times10^{14}$ W/cm^2 is likely to cause the excitation of Bernstein waves via three wave parametric process. The power thresholds for resonant and nonresonant decay processes are primarily determined by the inhomogeneity in the magnetic field and are comparable to each other. Above the threshold, the growth rate at 10^{16} W/cm^2 is of the order of 3×10^{11} s^{-1} for a resonant decay and 10^{11} s^{-1} for a nonresonant decay in the underdense region. Near the critical layer the growth rate is of the order of 5×10^{11} s^{-1} for a resonant decay and 10^{11} s^{-1} for a nonresonant quasimode decay at power density of 10^{17} W/cm^2. In the presence of an inhomogeneity in the magnetic field, the power threshold for resonant decay in underdense region is 6 times more as compared to critical density region whereas for a quasimode decay it is 10^{15} W/cm^2 near the critical layer.

For the parameters of Labaune *et al.*'s experiment, $n_0^0\sim0.1n_{\rm cr}$, $\lambda_0=1.05$ μm, $L_n\sim1000$ μm, $T_e\sim0.4$ keV, $c_s\sim10^7$ cm/s, $\omega_a\equiv2k_0c_s\sim1.2\times10^{12}$ if one assumes $B_S\sim0.7$ MG, $L_B\sim L_n$ then one obtains $\omega\sim10\omega_a$ and threshold power turns out to be 5×10^{14} W/cm^2. This power is an order of magnitude larger than the one reported in the experiment. If there are turning points for the electron Bernstein wave then the instability could have a much lower threshold. This could happen when B field has a parabolic profile with r. The experiment has indicated that the frequency ω of the low-frequency mode increases with the intensity of the laser. This may be due to the fact the frequency of the Bernstein wave increases with the magnetic field hence with the laser intensity.

ACKNOWLEDGMENTS

This work is supported by D.A.E. and D.S.T., Government of India.

[1] V. K. Tripathi and C. S. Liu, Phys. Plasmas **1**, 990 (1994).
[2] V. Yu Bychenkov, Yu. S. Kas'yanov, G. S. Sarkisov, and V. T. Tikhochuk [JETP Lett. **58**, 184 (1993)].
[3] W. L. Kruer, *The Physics of Laser Plasma Interactions* (Addison-Wesley, New York, 1988).
[4] J. A. Stamper, K. Papadopoulos, R. N. Sudan, S. O. Dean, E. A. Mclean, and J. M. Dawson, Phys. Rev. Lett. **26**, 1012 (1972).
[5] D. A. Tidman and R. A. Shanny, Phys. Fluids **17**, 1207 (1994).
[6] J. A. Stamper and D. A. Tidman, Phys. Fluids **16**, 2024 (1973).
[7] J. A. Stamper, Phys. Fluids **18**, 735 (1975).
[8] W. Woo and J. S. Degroot, Phys. Fluids **21**, 124 (1978).
[9] K. Nishihara, Y. Oshawa, K. Mima, and T. Tange, Phys. Fluids **19**, 1833 (1976).
[10] V. K. Tripathi, and R. R. Sharma, Phys. Fluids **22**, 1799 (1979).
[11] O. P. Sharma, Phys. Rev. A **26**, 2854 (1982).
[12] A. Simon (unpublished).
[13] S. D. Baton *et al.*, Phys. Rev. E **49**, R3602 (1994).
[14] H. A. Baldis *et al.*, Phys. Fluids B **5**, 3319 (1993).
[15] A. Chirokilch (unpublished).
[16] C. Labaune *et al.*, Phys. Rev. Lett. **75**, 248 (1995).