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Stimulated Raman backward scattering of a laser beam in a magnetized plasma

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ABSTRACT

Stimulated Raman scattering of a laser beam is investigated in the plasma with strong self generated magnetic field. The magnetized plasma supports various localized radial and azimuthal modes of lower hybrid frequencies. The density fluctuations due to lower hybrid modes couple with the oscillating velocity due to the pump, and drive the scattered wave. Equations describing the Raman process are derived and effects of various modes are studied on the growth rate analytically. Self generated magnetic field has a strong localization effect on the Raman process and growth rate is maximum for radial eigen mode number $q=0$ and azimuthal eigen mode number $l=3$. The frequency shift has signatures of self generated magnetic field and could serve as a diagnostic.

1. Introduction

The nonlinear interaction of intense short pulse laser with plasmas has been a subject of much experimental and theoretical study in recent years due to its relevance to laser driven fusion [1,2]. In high power laser plasma interaction, stimulated Raman scattering (SRS) is an important three wave process. In this process, the ponderomotive beatwave between the pump (ω_0) and its stokes/ antistokes sidebands ($\omega_0 \pm \omega_p$) drive a plasma wave (ω_p), which then acts as a grating, scattering the pump and reinforcing the sidebands. The transmission of laser light through a plasma and the coupling of the laser energy into a plasma can be greatly affected by Raman scattering, which consequently can have a large impact on various applications. For example, Raman scattering in the forward direction can be used to drive the self modulated laser wakefield accelerator [3], in which a long laser pulse becomes modulated and produces a large amplitude plasma wave with phase velocity near the speed of light $v_\phi = \omega/k \approx c$. This plasma wave with 10–100 GeV/m accelerating gradients have been demonstrated using present laser technology and can be used to accelerate charged particles to high energies [4–7].

In laser fusion accelerations such as fast ignition [8], the excitation of Raman instabilities can yield poor coupling of laser to the energetic electrons. Thus, SRS plays a destructive role in fast ignition fusion scheme by diverting a significant amount of laser energy away from pellet heating. Growth of SRS can be controlled via Langmuir decay instability [9] and electron trapping [10] by restricting experiments in certain regimes of plasma parameters. The use of finite bandwidth laser pulse has also been considered for enhancement/suppression of Raman instabilities [11,12]. In the fast ignition scheme [13–17], simulations show more than a 30% energy transfer from the laser to generate relativistic electrons. The relativistic electrons give rise to an azimuthal magnetic field around the axis of the laser. Pukhov et al. [5] and Sentoku et al. [18] have observed azimuthal magnetic fields of the order of 100 MG in their particle-in-cell simulations. Borghesi et al. [19] and Mackinnon et al. [20] have experimentally observed multi-megagauss magnetic field generation with picosecond and subpicosecond pulses of intensity $I \sim 10^{19}$ W/cm^2. These magnetic fields play a crucial role in indirect laser acceleration of electrons. Whenever the Doppler shifted laser frequency, as experienced by an accelerating electron, equals the frequency of betatron oscillator in the magnetic field, an efficient energy transfer from laser to electrons occurs [21].

In a uniform magnetic field, the motion of plasma electrons modifies under the influence of the force $-e\left(\vec{v} \times \vec{B}\right)$. It will give rise to change in the dispersion of the laser beam, nonlinear current density and natural plasma modes. The excited plasma modes transfer its energy to the plasma particles via Landau damping and leads to enhanced heating of the plasma [22,23]. The success of laser based particle accelerators, radiation sources and inertial confinement fusion scheme critically depend on the amount of transmitted laser energy through the plasma, and consequently SRS has been studied extensively in both un-magnetized and magnetized plasmas [24–28].

The self generated magnetic field can open up new channels of SRS as well. For example, a magnetized plasma supports lower hybrid waves with frequency ω in the range $\omega_{pi}/(1+\omega_p^2/\omega_c^2) < \omega < \omega_p$, ω_c. Here $\omega_{pi}$is the ion plasma frequency and ω_p, ω_c are the plasma and cyclotron frequencies of the electrons respectively. The frequency and mode structure of the mode could be strongly influenced by the

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inhomogeneous magnetic field. Parametric coupling of the laser with this mode would have strong signatures of the self generated magnetic field.

In this paper, we have studied the SRS of laser beam in plasma with self generated magnetic field. The laser excites an electrostatic lower hybrid mode and a side band electromagnetic wave. The density fluctuations due to lower hybrid modes couple with the oscillatory velocity of the plasma electron due to the pump to produce a nonlinear current driving the sideband waves. The pump and sideband wave exert a ponderomotive force on the electrons driving the lower hybrid wave. We have investigated the effects of various radial and azimuthal lower hybrid modes on the growth of Raman process. In Section 2, we have derived the mode structures of various plasma eigenmodes in a magnetized plasma. In Section 3, we study the three wave parametric instability. Sections 4 and 5 are dedicated to a discussion of the results and conclusions respectively.

2. Eigen modes

Consider a plasma of electron density n_0. When a laser pulse of spot size r_0 propagate through it along z-axis with normalized Gaussian amplitude profile $a^2=A_L^2 e^{-x^2/r_0^2} e^{-\psi^2}$, where $\psi=(t-z/\eta c-t_0)/\tau$, $\eta=(1-\omega_p^2/\omega_0^2)^{1/2}$ and $\omega_p=(n_0 e^2/\varepsilon_0 m)^{1/2}$, it produces a self generated azimuthal magnetic field $B=B_0(r/r_0)\exp\left[-r^2/2r_0{}^2\right]\hat{\theta}$, as suggested by Gorbunov et al. [29]. Here –e and m are electronic charge and rest mass respectively, $\eta c$is the group velocity, ω_p is electron plasma frequency and ε_0 is free space permittivity.

In equilibrium, the drift velocity of plasma electrons is considered to be zero. We perturb this equilibrium by an electrostatic perturbation $\vec{E}=-\nabla\phi$, where $\phi=\phi(r)e^{-i(\omega t-k_z z-l\theta)}$ where k_z is the wave number along z-axis and l is the azimuthal mode number. Using equation of motion for plasma electrons, we obtain the electron velocity components parallel and perpendicular to the magnetic field respectively. The electron density perturbation is obtained by solving equation of continuity as follows:

$$n_1=+\frac{en_0}{m\omega^2}\left[\frac{1}{r}\nabla_{\text{II}}^2\phi+\frac{1}{(1-\omega_c^2/\omega^2)}\nabla_\perp^2\phi\right] \tag{1}$$

where $\omega_{c0}=eB_0/m$, $\nabla_{\text{II}}^2=\frac{1}{r^2}\frac{\partial^2}{\partial\theta^2}$ and $\nabla_\perp^2=\frac{\partial^2}{\partial r^2}+\frac{1}{r}\frac{\partial}{\partial r}-k_z{}^2$. Using this value of density perturbation in the Poisson equation $\nabla^2\phi=ne/\varepsilon_0$, we obtain the mode structure equation

$$\frac{\partial^2\phi}{\partial r^2}+\frac{1}{r}\frac{\partial\phi}{\partial r}+\left[-k_z^2-\frac{l^2}{r^2}+\frac{2\omega_p^2\ \omega_{c0}^2}{\omega^2\left(\omega_p^2-\omega^2\right)^2}\frac{l^2}{r_0^2}+\frac{\omega_{c0}^2\left(4\omega_p^2+\omega_{c0}^2+4\omega^2\right)}{\left(\omega_p^2-\omega^2\right)^2}\frac{l^2}{r_0^4}r^2\right]\phi=0 \tag{2}$$

Here, we have neglected the higher terms of r, that is valid only for $\omega^2\ll\omega_{c0}^2$. Eq. (2) has associated Laguerre polynomial solutions

$$\phi=A\psi(r) \tag{3}$$

where $\psi(r)=\frac{1}{2\alpha R}\left(\frac{r}{R}\right)^l\exp\left(r^2/2R^2\right)L_P^l\left(r^2/R^2\right)$ and q is an integer characterizing the eigen values

$$\alpha^2R^2=2q+l;q=0,1,2\ldots \tag{4}$$

where $\alpha=\sqrt{\frac{2\omega_p{}^2\omega_{c0}{}^2}{(\omega_p{}^2-\omega^2)^2}\frac{l^2}{r_0{}^2}-k_z{}^2}$ and $R=\sqrt{\frac{r_0{}^2(\omega_p{}^2-\omega^2)}{2\omega_{c0}\quad l\sqrt{4\omega_p{}^2+\omega_{c0}{}^2+4\omega^2}}}$

ϕ vanishes on the axis and peaks at $r\approx 1/\alpha$ because the lower hybrid mode exists only for ω_{c0}, $\omega_P>\omega$ and since ω_c vanishes at $r=0$, the mode has vanishing amplitude there. It exists at values of r where $\omega_c>\omega$. The pump parametrically couple with modes given by Eq. (3) and undergo Raman scattering in the lower hybrid range. The scattered electromagnetic wave has strong signature of self generated azimuthal magnetic field.

3. Stimulated Raman process

Consider the propagation of a laser beam with Gaussian intensity profile in a plasma with electric field profile $E=E_0(r)e^{-i(\omega_0 t-k_0 z)}$, where $E_0^2(r)=A_0^2\exp[-r^2/r_0^2]$. The laser imparts oscillatory velocity to electrons given by $\bar{v}_0=e\bar{E}_0/mi\omega_0$. We have assumed $v_0^2/2c^2<<1$ and neglected relativistic laser guiding effects. Liu and Tripathi [30] have studied the relativistic laser guiding in a strong azimuthal magnetic field in plasma. When $\omega_0\approx\omega_{c0}$, the laser undergoes anisotropic self focusing acquiring an elliptic cross-section. However for $\omega_0>>\omega_c$, the ellipticity is weak. The laser couples to a lower hybrid mode of potential $\phi=\phi(r)e^{-i(\omega t-k_z z-l\theta)}$ and a backscattered electromagnetic wave of field $E_1=E_1(r)e^{-i(\omega_1 t-k_1 z)}$, where $\omega_1=\omega-\omega_0$ and $k_{1Z}=k_Z-k_0$. Since $\omega\ll\omega_0$, we have $|\omega_1|\approx\omega_0$ and $|k_{1Z}|\approx 2k_0$.

The scattered wave imparts an oscillatory velocity $\bar{v}_1=e\bar{E}_1/(mi\omega_1)$ to electrons. The pump and the scattered wave exert a ponderomotive force $\bar{F}_p=e\nabla\phi_p$ on plasma electrons, where $\phi_p=(m/2c)\left(\vec{v}_0.\vec{v}_1\right)$. The electron density perturbation at (ω, k_Z) due to ϕ and ponderomotive potential ϕ_p, can be obtained from Eq. (1) by replacing ϕ by $\phi+\phi_p$. The electron density perturbation contains both linear and nonlinear parts. Ignoring ion motion and using perturbed electron density in the Poisson equation, we obtain

$$\frac{\partial^2\phi}{\partial r^2}+\frac{1}{r}\frac{\partial\phi}{\partial r}+\left[-k_z^2-\frac{l^2}{r^2}+\frac{2\omega_p^2\omega_{c0}^2}{\left(\omega_p^2-\omega^2\right)^2}\frac{l^2}{r_0^2}+\frac{\omega_{c0}^2\left(4\omega_p^2+\omega_{c0}^2+4\omega^2\right)}{\left(\omega_p^2-\omega^2\right)^2}\frac{l^2}{r_0^4}r^2\right]\phi$$

$$=\frac{\omega_p^2}{\left(\omega^2-\omega_{c0}^2-\omega_p^2\right)}\left[\frac{\left(\omega^2-\omega_{c0}^2\right)}{\omega^2}\frac{l^2}{r^2}\phi_p-k_z^2\phi_p+\frac{\partial^2\phi_p}{\partial r^2}+\frac{1}{r}\frac{\partial\phi_p}{\partial r}\right](5) \tag{5}$$

If one ignores nonlinear coupling, Eq. (5) gives an associated Laguerre polynomial solution given by Eqs. (3) and (4). On retaining the nonlinear coupling, we assume that the mode structures of lower hybrid waves are not modified significantly, however their eigen values are modified.

The density perturbation couples with $\vec{v}_0$ to produce a nonlinear current density at ω_1 given by $J_1^{NL}=(e/2)nv_0{}^*=-(1/8\pi)\nabla^2\phi v_0{}^*$. Using this value of nonlinear current density in the electromagnetic wave equation, we obtain

$$\frac{\partial^2 E_1}{\partial r^2}+\frac{1}{r}\frac{\partial E_1}{\partial r}+\left[\alpha_z^2+\frac{\omega_p^2}{c^2}\frac{v_0^2}{4c^2}\frac{r^2}{r_0^2}-\frac{l^2}{r^2}\right]E_1=\frac{i\omega_1}{2c^2}\nabla^2\phi\bar{v}_0^* \tag{6}$$

where $\alpha_1^2=\frac{\omega_1^2-\omega_p^2}{c^2}-k_{1z}^2$. If nonlinear coupling is ignored, Eq. (6) reduces to an associated Laguerre equation with

$$E_1=A_1\psi_1(r)$$
$$\psi_1=\frac{1}{\sqrt{2\alpha_1 r_1}}\left(\frac{r}{r_1}\right)^l\exp\left(-r^2/2r_1{}^2\right)L_P^L\left(r^2/r_1{}^2\right) \tag{7}$$

Using Eq. (3) into Eq. (5) and multiplying by $\psi_1(r)$ and integrating over $r\,dr$ from 0 to ∞, we obtain

$$2c^2 I_1\left(\alpha_1{}^2-\alpha_{1p}{}^2\right)A_1=i\omega_1 v_0^* A\beta_1 \tag{8}$$