CONVECTIVE AMPLIFICATION OF STIMULATED RAMAN BACKSCATTERING IN INHOMOGENEOUS PLASMAS

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الخلاصة :

في هذا البحث نشتق تحليليّاً مجموعة من المعادلات التي تصف التولد البطيء لتشتت (رامان) الخلفي في البلازما غير المتجانسة حيث تتبع الكثافة نموذجاً خطياً . ثم ندرس التكبير الانتقالي لظاهرة عدم الاستقرار في البلازما غير المتجانسة والمنخفضة الكثافة وتأثير التصادمات المثبطة في حالة عدم التجانس .

ABSTRACT

In this work we derive analytically a set of equations that describe the slow coupling process of Raman backscattering in an inhomogeneous medium in which the plasma density is a linear ramp. Then we consider the convective amplification of the Raman instability in the underdense inhomogeneous plasma taking into consideration the collision damping in addition to the background inhomogeneity.

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1. INTRODUCTION

The complex problem of laser–plasma interaction has attracted a great deal of attention. This is due to interest in laser pellet fusion, as a possible source of energy with no external confinement apart from the inertia of the plasma itself (inertial confinement) [1]. This approach is completely different from the relatively old magnetic field confinement approach.

Since energy absorption is necessary for the operation of laser pellet fusion [1–3], it is necessary to know the amount of incident laser energy scattered and the conditions governing this scattering. Forslund *et al.* [4] addressed this question and showed that for large systems, *i.e.* a long region of underdense plasma, the ratio of backscattered to incident laser energy flux can be the ratio of their frequency.

Guzdar *et al.* [5] investigated the effect of bandwidth on convective amplification of the Raman instability in an underdense inhomogeneous plasma and showed that there is no effect of bandwidth on the convective amplification when the growth rate $\gamma_0 \ll \Delta\omega$, where $\Delta\omega$ is the bandwidth. Also they concluded that for $\gamma_0 \geq \Delta\omega$ there is a statistical enhancement in the amplification factor.

In this paper we consider convective amplification of the Raman instability in the underdense inhomogeneous plasma where we examine the effect of collision damping on the growth rate, the amplification factor, and the instability threshold.

Consider an electromagnetic wave normally incident onto an inhomogeneous plasma slab with linearly increasing density, *i.e.* $n_e(x) = n_i(x) = n_0(1 + x/L)$, where n_e is the electron density, n_i is the ion density, L is the inhomogeneity scale length of the plasma and x is the distance along the density gradient, so that the electromagnetic wave propagates in the underdense region, where the local plasma frequency is less than the incident electromagnetic wave frequency. At the critical density $(m\omega_0^2 / 4\pi e^2)$, the incident electromagnetic wave is reflected. In the overdense region, the incident wave will be exponentially attenuated. This picture is essentially correct for a low intensity electromagnetic wave. However, when the incident radiation exceeds a certain intensity level, it becomes capable of exciting collective modes parametrically inside the plasma.

Near the critical density a strong external electromagnetic field, such as a laser, can decay into a plasma wave (plasmon) and an ion wave (phonon) [6] leading to the anomalous absorption of the incident radiation instead of reflection. In the underdense region the incident laser can decay into a plasmon or a phonon and a scattered electromagnetic wave, where the former case is called the stimulated Raman scattering (SRS) and the latter is called the stimulated Brillouin scattering (SBS) [7]. Since the growth rates for these scattering processes are maximum for backscattering, the incident wave is primarily backscattered once its intensity exceeds a certain threshold value. Sidescattering is also possible and usually has a lower threshold than that of backscattering [8]. Thus the underdense region is a scatterer rather than transparent.

SRS is basically a high-frequency nonlinear phenomenon that can lead to anomalous reflection of the incident laser light. However, the most important feature of SRS is its ability to generate high energy electrons causing preheating of the pellet core and thus preventing the efficient compression which is necessary for the production of a useful amount of fusion energy [9].

SRS has been observed experimentally [10] and predicted theoretically as well as in computer simulation [11].

In Section 2 of this paper we derive analytically the slow coupling equations in an inhomogeneous plasma. In Section 3 we investigate the convective amplification in such plasmas and derive an expression to determine its value. Finally in Section 4 we present our conclusions.

2. THE SLOW COUPLING EQUATIONS

We start our analysis by the definition of electron and ion densities in an inhomogeneous plasma, the total electron density n_e is:

$$n_e = n_e(x) + \bar{n}_e(x,t), \tag{1}$$

where $n_e(x)$ is the slowly varying density due to inhomogeneity and is given by

$$n_e(x) = n_i(x) = n_0\left(1 + \frac{x}{L}\right) \qquad (2)$$

and $\bar{n}_e(x,t)$ is the fast, harmonically varying, density due to the electrons' fast response to the high frequency fields.

The fundamental set of equations for the study of SRS in the underdense plasma are the momentum equation:

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v}\cdot\nabla)\mathbf{v} = -\frac{e}{m}\mathbf{E} - \frac{e}{mc}\mathbf{v}\times\mathbf{B} - \frac{3T}{mn_e(x)}\nabla n_e - \nu\mathbf{v}; \qquad (3)$$

the continuity equation

$$\frac{\partial n_e}{\partial t} + \nabla\cdot(n\mathbf{v}) = 0; \qquad (4)$$

Poisson's equation

$$\nabla\cdot\mathbf{E} = -4\pi e\,\bar{n}_e(x,t); \qquad (5)$$

Faraday's law

$$\nabla\times\mathbf{E} = -\frac{1}{c}\frac{\partial B}{\partial t}; \qquad (6)$$

and Ampere's law

$$\nabla\times\mathbf{B} = -\frac{4\pi}{c}en_e\mathbf{v} + \frac{1}{c}\frac{\partial \mathbf{E}}{\partial t}, \qquad (7)$$

where $\mathbf{v}$ is the electron velocity, T is the electron temperature in energy units, ν is the electron ion collision frequency, and $\mathbf{E}$ and $\mathbf{B}$ are the electric and magnetic fields of the waves involved in the three wave process.

As usual, the pump wave (laser) and the two daughter waves must satisfy the frequency and wave-vector matching conditions

$$\omega_0 = \omega_1 + \omega_2, \quad \mathbf{k}_0 = \mathbf{k}_1 + \mathbf{k}_2, \qquad (8)$$

where ω_0, ω_1, ω_2 are the frequencies of the laser wave, the scattered electromagnetic wave, and the plasma wave respectively, and $\mathbf{k}_0$, $\mathbf{k}_1$, $\mathbf{k}_2$ are the propagation vectors of the laser wave, the scattered electromagnetic wave, and the plasma wave respectively.

Each wave of the three interacting waves satisfies its own dispersion relation

$$\omega_0^2 = \omega_p^2 + c^2k_0^2 \quad \text{pump wave (laser)} \qquad (9)$$

$$\omega_1^2 = \omega_p^2 + c^2k_1^2 \quad \text{scattered electromagnetic wave} \qquad (10)$$

$$\omega_2^2 = \omega_p^2 + \tfrac{3}{2}v_{th}^2k_1^2 \quad \text{Plasma wave} \qquad (11)$$

where c is the speed of light, $v_{th} = \sqrt{\frac{2T}{m}}$ is the electron thermal velocity, $\omega_p = \left(\frac{4\pi n_0e^2}{m}\right)^{1/2}$ is the plasma frequency, $n = n_e(x)$ is the electron density, and e, m, and T are the electron charge, mass, and temperature respectively.

To obtain the generalized wave equation, we take the curl of Equation (6); thus we get

$$\nabla^2\mathbf{E} - \frac{1}{c^2}\frac{\partial^2\mathbf{E}}{\partial t^2} - \nabla(\nabla\cdot\mathbf{E}) = -\frac{4\pi en_e(x)}{c^2}\frac{\partial\mathbf{v}}{\partial t} - \frac{4\pi e}{c^2}\frac{\partial}{\partial t}(\bar{n}_e\mathbf{v}) \qquad (12)$$

Using Equations (3) and (5), this equation can be written in the form

$$\nabla^2\mathbf{E} - \frac{1}{c^2}\frac{\partial^2\mathbf{E}}{\partial t^2} - \nabla(\nabla\cdot\mathbf{E}) - \frac{4\pi e^2n_e(x)}{mc^2}\mathbf{E} - \frac{4\pi en_e(x)}{c^2}\nu\mathbf{v} + \frac{3T}{mc^2}\nabla(\nabla\cdot\mathbf{E}) = \frac{1}{c^2}\frac{\partial}{\partial t}(\mathbf{v}\nabla\cdot\mathbf{E}) + \frac{4\pi en_e(x)}{c^2}\left[\tfrac{1}{2}\nabla(\mathbf{v}\cdot\mathbf{v}) - \mathbf{v}\times\left(\nabla\times\mathbf{v} - \frac{e\mathbf{B}}{mc}\right)\right] \qquad (13)$$

where we have made use of the identity

$$\mathbf{v}\times(\nabla\times\mathbf{v}) = \tfrac{1}{2}\nabla(\mathbf{v}\cdot\mathbf{v}) - (\mathbf{v}\cdot\nabla)\mathbf{v}.$$

Equation (13) is the generalized inhomogeneous, nonlinear wave equation, where the R.H.S. represents the coupling terms.

For our present purpose it is enough to expand the coupling terms about the linear values of $\mathbf{v}$, $\mathbf{B}$, and $\mathbf{E}$, so using the linear value of the electron velocity (known as the quiver velocity) $\mathbf{v} = \frac{e\mathbf{E}}{im\omega}$. It is easy to show that $\nabla\times\mathbf{v} = \frac{e\mathbf{B}}{mc}$ and Equation (13) becomes:

$$\nabla^2\mathbf{E} - \frac{1}{c^2}\frac{\partial^2\mathbf{E}}{\partial t^2} - \nabla(\nabla\cdot\mathbf{E}) - \frac{4\pi e^2n_e(x)}{mc^2}\mathbf{E} - \frac{4\pi en_e(x)}{c^2}\nu\mathbf{v} + \frac{3T}{mc^2}\nabla(\nabla\cdot\mathbf{E}) = \frac{1}{c^2}\frac{\partial}{\partial t}(\mathbf{v}\nabla\cdot\mathbf{E}) + \frac{2\pi en_e(x)}{c^2}\nabla(\mathbf{v}\cdot\mathbf{v}) \qquad (14)$$

Taking the x-component of this equation and ignoring the nonresonant term $\frac{1}{c^2}\frac{\partial}{\partial t}(v_x \nabla \cdot \mathbf{E})$, since its frequency response ($2\omega_2$) does not conform with the matching conditions, we easily get:

$$\frac{3T}{mc^2}\frac{\partial^2 E_x}{\partial x^2} - \frac{1}{c^2}\frac{\partial^2 E_x}{\partial t^2} - \frac{4\pi e^2 n_e(x)}{mc^2} E_x - \frac{v 4\pi e n_e(x)}{c^2} v_x = \frac{2\pi e n_e(x)}{c^2} \nabla(\mathbf{v}\cdot\mathbf{v}). \quad (15)$$

Taking the y-component of Equations (14), we get

$$\nabla^2 E_y - \frac{1}{c^2}\frac{\partial^2 E_y}{\partial t^2} - \frac{4\pi e^2 n_e(x)}{mc^2} E_y - \frac{v 4\pi e n_e(x)}{c^2} v_y = \frac{1}{c^2}\frac{\partial}{\partial t}(v_y \nabla \cdot \mathbf{E}) \quad (16)$$

where ∇ in the above equations is given by the relation $\nabla = \hat{x}\frac{\partial}{\partial x}$.

Noting that:

$$\mathbf{E}_x(x,t) = \hat{x} E_2(x,t) \exp i(k_2 x - \omega_2 t) + c.c. \quad (17)$$

$$\mathbf{E}_y(x,t) = \hat{y}[E_1(x,t) \exp -i(k_1 x + \omega_1 t) + E_0(x,t) \exp i(k_0 x - \omega_0 t) + c.c.] \quad (18)$$

$$v_j = \frac{eE_j}{im\omega_j} \text{ where } j = 0,\ 1,\ 2 \quad (19)$$

$$\mathbf{v}_x = \hat{x}\, v_2 \exp i(k_2 x - \omega_2 t) + c.c. \quad (20)$$

$$\mathbf{v}_y = \hat{y}[v_1 \exp -i(k_1 x + \omega_1 t) + v_0 \exp i(k_0 x + \omega_0 t) + c.c.] \quad (21)$$

and using the slow coupling limit

$$\frac{\partial^2}{\partial x^2}(E_0, E_1, E_2) = 0 \quad (22)$$

$$\frac{\partial^2}{\partial t^2}(E_0, E_1, E_2) = 0$$

we get by comparing phases on both sides of Equations (15) and (16), the following set of equations

$$\left[\frac{\partial}{\partial t} + V_0\frac{\partial}{\partial x} + \frac{i\omega_p^2(0)x}{2\omega_0 L} + \frac{v}{2}\frac{\omega_p^2(0)}{\omega_0^2}\left(1+\frac{x}{L}\right)\right]E_0 = -\frac{ik_2 v_1}{2} E_2 \exp(i\Psi) \quad (23)$$

$$\left[\frac{\partial}{\partial t} - V_1\frac{\partial}{\partial x} + \frac{i\omega_p^2(0)x}{2\omega_1 L} + \frac{v}{2}\frac{\omega_p^2(0)}{\omega_1^2}\left(1+\frac{x}{L}\right)\right]E_1 = \frac{ik_2 v_0}{2} E_2^* \exp(-i\Psi) \quad (24)$$

$$\left[\frac{\partial}{\partial t} + V_2\frac{\partial}{\partial x} + \frac{i\omega_p^2(0)x}{2\omega_2 L} + \frac{v}{2}\frac{\omega_p^2(0)}{\omega_2^2}\left(1+\frac{x}{L}\right)\right]E_2 = \frac{ik_2 v_0}{2}\frac{\omega_{pe}^2(0)}{\omega_1\omega_2} E_1^* \exp(-i\Psi), \quad (25)$$

where

$$V_0 = \frac{k_0 c^2}{\omega_0}$$

is the group velocity of the pump wave,

$$V_1 = \frac{k_1 c^2}{\omega_1}$$

is the group velocity of the scattered electromagnetic wave,

$$V_2 = \frac{3v_{th}^2 k_2}{2\omega_2}$$

is the group velocity of the electron plasma wave, and $\Psi = (\omega_0 - \omega_1 - \omega_2)t - (k_0 + k_1 - k_2)x$ is a phase mismatch introduced for generality purposes.

Equations (23), (24), (25) and their complex conjugates are the slow coupled equations for Raman Backscattering.

3. CONVECTIVE AMPLIFICATION IN AN INHOMOGENEOUS PLASMA

In an inhomogeneous plasma an absolute (temporally growing) instability occurs near $n_c/4$.

Below $n_c/4$ the instability is known to be convective (spatially growing) [12].

Assuming that there is perfect matching ($\Psi = 0$), $\omega_p(0) = \omega_2$, and defining:

$$a_j = \frac{E_j}{\sqrt{8\pi\omega_j}} \quad j = 0,\ 1,\ 2, \tag{26}$$

we can write Equations (24) and (25) as

$$\left[\frac{\partial}{\partial t} - V_1 \frac{\partial}{\partial x} + \frac{i\omega_p^2(0)x}{2\omega_1 L} + \frac{v}{2}\frac{\omega_p^2(0)}{\omega_1^2}\left(1 + \frac{x}{L}\right)\right]a_1 = \frac{ik_2 v_0}{2}\sqrt{\frac{\omega_p(0)}{\omega_1}}a_2^* \tag{27}$$

$$\left[\frac{\partial}{\partial t} + V_2 \frac{\partial}{\partial x} + \frac{i\omega_p(0)x}{2L} + \frac{v}{2}\left(1 + \frac{x}{L}\right)\right]a_2 = \frac{ik_2 v_0}{2}\sqrt{\frac{\omega_p(0)}{\omega_1}}\ a_1^*. \tag{28}$$

Noting that the homogeneous growth rate is

$$\gamma_0 = \frac{k_2 v_0}{2}\sqrt{\frac{\omega_p(0)}{\omega_1}}, \tag{29}$$

and normalizing the spatial displacement and time by the transformation,

$$x_n = \frac{x}{x_0}, \quad t_n = \frac{t}{t_0}, \tag{30}$$

where we have chosen

$$V_1 = \frac{x_0}{t_0} = \frac{k_1 c^2}{\omega_1} \tag{31}$$

and

$$x_0 t_0 = \frac{2\omega_1 L}{\omega_p^2(0)} \tag{32}$$

we get

$$\left[\frac{\partial}{\partial t_n} - \frac{\partial}{\partial x_n} + i x_n + z\left(1 + \frac{V_1 t_0 x_n}{L}\right)\right]a_1 = i\gamma a_2^* \tag{33}$$

$$\left[\frac{\partial}{\partial t_n} + \beta\frac{\partial}{\partial x_n} + i\alpha x_n + \frac{v t_0}{2}\left(1 + \frac{V_1 t_0}{L}x_n\right)\right]a_2 = i\bar{\gamma}a_1^* \tag{34}$$

where

$$\alpha = \frac{\omega_1}{\omega_p(0)},\ \beta = \frac{V_2}{V_1},\ \bar{\gamma} = \gamma_0 t_0, \text{ and}$$

$$z = \frac{v}{2}\frac{\omega_p^2(0)}{\omega_1^2}t_0.$$

Equations (33) and (34) are the fundamental equations that describe the slow coupling Raman backscattering instability in the underdense plasma. These equations are valid when the growth rate γ_0 and the bandwidth $\Delta\omega$ are much smaller than $\omega_p(0)$ and also if the interaction length x_0 and coherence length $L_c = \dfrac{V_0}{\Delta\omega}$ are less than the plasma size [5]. These conditions are readily met for recent experiments with bandwidth [13, 14].

It is known [15, 16] that such equations give only convective amplifications (spatially growing) and no absolute instability (temporally growing).

Since we are interested in the convective instability, *i.e.* the spatial variation, we can put $\dfrac{\partial}{\partial t_n} = 0$, and assuming $\beta = 0$ ($V_2 \ll V_1$), Equation (34) gives

$$a_2 = \frac{i\bar{\gamma}a_1^*}{\frac{v t_0}{2}\left(1 + \frac{V_1 t_0 x_n}{L}\right) + i\alpha x_n}. \tag{35}$$

Taking the complex conjugate of this equation and substituting for a_2^* in Equation (33), we obtain

$$\left[\frac{\partial}{\partial x_n} - i x_n - z\left(1 + \frac{V_1 t_0 x_n}{L}\right)\right]a_1 = \frac{-|\bar{\gamma}|^2 a_1}{\frac{v_e t_0}{2}\left(1 + \frac{V_1 t_0 x_n}{L}\right) - i\alpha x_n}. \tag{36}$$

Again, taking the complex conjugate of this equation, we get

$$\left[\frac{\partial}{\partial x_n} + i x_n - z\left(1 + \frac{V_1 t_0 x_n}{L}\right)\right]a_1^* = \frac{-|\bar{\gamma}|^2 a_1^*}{\frac{v t_0}{2}\left(1 + \frac{V_1 t_0 x_n}{L}\right) + i\alpha x_n}. \tag{37}$$

Since, we are interested in the intensity amplification we obtain from Equations (36) and (37), the equation

$$\left[\frac{\partial}{\partial x_n} + \frac{|\bar{\gamma}|^2 \upsilon t_0\left(1 + \frac{V_1 t_0 x_n}{L}\right)}{\left[\frac{\upsilon t_0}{2}\left(1 + \frac{V_1 t_0 x_n}{L}\right)\right]^2 + \alpha^2 x_n^2} - 2z\left(1 + \frac{V_1 t_0 x_n}{L}\right)\right]|a_1|^2 = 0. \quad (38)$$

Defining

$$y_n = 1 + \frac{V_1 t_0}{L} x_n$$

$$f_1 = \frac{\upsilon t_0^2 V_1}{\alpha^2 L}|\bar{\gamma}|^2$$

and

$$f_2 = \frac{4\alpha^2 L^2}{4\alpha^2 L^2 + \gamma^2 t_0^4 V_1^2},$$

we get

$$\frac{d|a_1|^2}{dy_n} = \frac{2Lz}{V_1 t_0} y_n |a_1|^2 - \frac{f_1 f_2 y_n |a_1|^2}{y_n^2 - 2f_2 y_n + f_2}. \quad (39)$$

Assuming that interaction region is symmetric and extends from $-y_n$ to y_n, then, by integration, we obtain:

$$\ln\left[\frac{|a_1(u_+)|^2}{|a_1(-u_-)|^2}\right] = -\frac{f_1 f_2^2}{B}\left[\tan^{-1}\left(\frac{u_+}{B}\right) - \tan^{-1}\left(\frac{u_-}{B}\right)\right], \quad (40)$$

where $u_\pm = \pm(y_n - f_2)$ and $B = \sqrt{f_2(1 - f_2)}$.

Extending the interaction region from $-\infty$ to $+\infty$, we get

$$\ln\frac{|a_1(\infty)|^2}{|a_1(-\infty)|^2} = \frac{-f_1 f_2^2 \pi}{\sqrt{f_2(1 - f_2)}}$$

hence

$$|a_1(-\infty)|^2 = |a_1(\infty)|^2 \exp\left(\frac{2\pi|\bar{\gamma}|^2}{\alpha} \times \frac{1}{1 + \frac{\upsilon^2 t_0^4 V_1^2}{4\alpha^2 L^2}}\right). \quad (41)$$

From this equation we conclude that there is an increase in the scattering as the density scale length increases and collisional suppression of the instability, a result that agrees with recent experiments [10, 17].

As $\upsilon \to 0$, we get

$$|a_1(-\infty)|^2 = |a_1(\infty)|^2 \exp(2\pi|\bar{\gamma}|^2/\alpha),$$

which is the same result obtained by Guzdar *et al.* [5].

The amplification factor A can be obtained from Equation (41), where

$$A = \frac{2\pi|\bar{\gamma}|^2}{\alpha}\frac{1}{\left(1 + \frac{\upsilon^2 t_0^4 V_1^2}{4\alpha^2 L^2}\right)}. \quad (43)$$

So, in order to have a growth rate, we must have $A > 1$, then $A = 1$ effectively represents a threshold for this instability, then $|\bar{\gamma}|^2$ given by

$$|\bar{\gamma}|^2 = \frac{\alpha}{2\pi}\left(1 + \frac{\upsilon^2 t_0^4 V_1^2}{4\alpha^2 L^2}\right) \quad (44)$$

is proportional to the pump wave intensity.

4. CONCLUSIONS

We have investigated analytically the Raman backscattering instability in an inhomogeneous medium, and obtained a set of equations that describe this instability in the slow-coupling limit. We derived an expression for the scattered density which shows that there is an increase in the scattering as the density scale length increases, and collisional suppression of the instability. We find from theory that taking the collision damping into consideration reduces the growth rate of the instability and increases the threshold intensity, a result can be explained on the basis that collision provides another mechanism of energy absorption.

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