Stimulated magnetic scattering of electromagnetic waves in a plasma

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A new mechanism for the stimulated scattering of electromagnetic waves in a plasma by magnetic fluctuations is identified. This process, which is characterized by a side scattering and by the rotation of the radiation polarization vector, occurs without a frequency shift. These systematic features, taken collectively, set this scattering mechanism apart from the previously known mechanisms.

Parametric instabilities in a plasma such as stimulated Brillouin (or Mandel'shtam-Brillouin) scattering (SBS), stimulated Raman scattering (SRS), and stimulated thermal scattering (STS) have been studied extensively. Knowledge of the properties of these processes is the principle upon which the remote diagnostic study

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of the plasma exposed to a strong electromagnetic radiation is based.

In the present letter we report the discovery of a new parametric process: stimulated scattering by magnetic fluctuations, whose properties differ qualitatively from those of SBS, SRS, and STS and whose study shows promise for the development of a more detailed and hence a more reliable plasma diagnostic system. The new parametric instability observed by us corresponds to the onset in the plasma of quasistatic perturbations of the magnetic field and of the transverse rf wave. This essentially means that the electromagnetic radiation to which the plasma is exposed excites a nonpotential aperiodic instability which causes an effective scattering of pump waves on the intensifying fluctuations of the quasistatic magnetic field. We studied the spectral, angular, and polarization characteristics of the scattered light, which have several properties that differ qualitatively from the previously known systematic features of the SBS, SRS, and STS processes.

Let us consider a plasma which is subjected to the effect of an electromagnetic pump wave with a frequency ω_0, wave vector $\mathbf{k}_0$, and electric field strength $\mathbf{E}_0$. To describe the interaction of the pump wave with the plasma, we will use the transport equations for electrons (for the average velocity and viscous stress tensor) in the ten-moment approximation of Grad's method, with allowance for the nonlinear (quadratic in the velocity) viscosity.[3] In this case the characteristic time for the change in the hydrodynamic quantities and the pump wavelength are large in comparison with the mean-free path. Averaging these equations over the pump-wave oscillation period, we find the following relation, which is the analog of the Ohm's law ($|e|B/m_e c \ll \nu_{ei}$):

$$j_k/\sigma = \left(\mathbf{E} - \frac{1}{en_e}\nabla n_e T_e - \frac{e}{4m\omega_0^2}\nabla|\tilde{\mathbf{E}}|^2\right)_k - \frac{e}{5a\omega_0^2 m_e}\frac{\partial}{\partial r_l}$$

$$\times\left(\tilde{E}_k\tilde{E}_l^* + \tilde{E}_k^*\tilde{E}_l - \frac{2}{3}\delta_{kl}|\tilde{\mathbf{E}}|^2\right)$$

$$- \frac{e^2}{5a^2 m_e^2\omega_0^2\nu_{ei}c}\frac{\partial}{\partial r_l}\left(\tilde{E}_l[\tilde{\mathbf{E}}^*,\mathbf{B}]_k + \tilde{E}_k[\tilde{\mathbf{E}}^*,\mathbf{B}]_l + \text{c.c.}\right). \qquad (1)$$

Here $\mathbf{j}$ and $\mathbf{E}(\mathbf{B})$ are the quasistatic current density and the quasistatic electric (magnetic) field; $\sigma = e^2 n_e/m_e\nu_{ei}$ is the electrical conductivity of the plasma; c is the velocity of light; ν_{ei} is the frequency of electron-ion collisions; $a = \frac{6}{5}[1 + (1/\sqrt{2}z)]$, where z is the degree of ionization of the ions; and $\tilde{\mathbf{E}}(\mathbf{r},t)$ is a rapidly alternating (at a frequency close to ω_0) field in the plasma:

$$\tilde{\mathbf{E}}(\mathbf{r},t) = \mathbf{E}_1(\mathbf{r},t)e^{-i\omega_0 t} + \text{c.c.}\ .$$

In the absence of the last term, relation (1) corresponds to the result of Ref. 4 for the ponderomotive force in a strongly collisional plasma.

Supplementing Eq. (1) with equations for the electromagnetic fields

$$\frac{\partial \mathbf{B}}{\partial t} = -c \operatorname{curl} \mathbf{E}, \quad \operatorname{curl} \mathbf{B} = \frac{4\pi}{c}\mathbf{j} ,$$

$$\frac{2i}{\omega_0}\frac{\partial \mathbf{E}_1}{\partial t} + \epsilon(\omega_0)\mathbf{E}_1 - \frac{c^2}{\omega_0^2} \operatorname{curl}\operatorname{curl} \mathbf{E}_1 = i\frac{e\,\omega_{Le}^2}{m_e c\,\omega_0^3}[\mathbf{E}_1, \mathbf{B}] \qquad (2)$$

$$-\frac{ice}{m_e\omega_0^3}\{[\operatorname{curl}\mathbf{B}, \operatorname{curl}\mathbf{E}_1] + \nabla_k(E_{1i}\operatorname{curl}_k\mathbf{B} + E_{1k}\operatorname{curl}_i\mathbf{B})\}$$

and perturbing these fields, we see, after linearizing (1) and (2), that the stability of the plasma described by Eqs. (2) can be solved by analyzing the dispersion equation ($|\omega| \ll \nu_{ei}$, $kv_{Te} \ll \nu_{ei}$)

$$\Gamma^2(\omega\quad \mathbf{k}) - \frac{\Gamma(\omega,\mathbf{k})}{5a}\left\{\frac{2[\mathbf{k},\mathbf{v}_E]^2}{a\nu_{ei}} + i\frac{k^2 v_E^2(\omega_{Le}^2 + c^2k^2)\left(1 - \frac{k^2}{4k_0^2}\right)}{\omega_0 D_+}\right\}$$

$$+\frac{(\mathbf{k}\mathbf{v}_E)^2(\mathbf{v}_E[\mathbf{k},\mathbf{k}_0])^2(\omega_{Le}^2 + c^2k^2)}{25a^2\omega_0 k_0^2 D_+}\left\{\frac{2}{a\nu_{ei}} + i\frac{\omega_{Le}^2 + c^2k^2}{\omega_0 D_+}\right\} = 0, \qquad (3)$$

where $D_\pm \equiv (\omega \pm \omega_0)^2\epsilon(\omega \pm \omega_0) - c^2(\mathbf{k} + \mathbf{k}_0)^2$, $\epsilon(\omega_0)$ is the dielectric constant, $\mathbf{v}_E = e\mathbf{E}_0/m_e\omega_0$, E_0 is the amplitude of the pump wave, and

$$\Gamma(\omega,\mathbf{k}) = -i\omega + \frac{c^2k^2}{4\pi\sigma} + \frac{2(\mathbf{k}\mathbf{v}_E)^2}{5a^2\nu_{ei}} + i\frac{(\mathbf{k}\mathbf{v}_E)^2(\omega_{Le}^2 + c^2k^2)}{5a\omega_0 D_+}$$

or by analyzing an equation which differs from (3) in the replacement of D_+ by D_-.

In a low-density plasma, $\omega_0^2 \gg \omega_{Le}^2$, the maximum γ_m of the instability growth rate is reached when

$$\mathbf{k}\cdot\mathbf{v}_E = 0, \quad k^2 = \mp 2\mathbf{k}\cdot\mathbf{k}_0 = \pm 2kk_0\cos\varphi ,$$

which corresponds to $\mathrm{Re}\omega = 0$ and $\mathrm{Re}D_+ = 0$, i.e., to the growth of the nonpotential aperiodic instability with the excitation of the transverse electromagnetic wave (the scattered wave). The threshold strength of the electric pump field is given by

$$\frac{v_{E,\,\mathrm{thr}}}{c} = \sqrt{5a}\,\frac{\nu_{ei}}{\omega_0} . \qquad (4)$$

The threshold in (4) is realized at $\varphi = 90° \pm 45°$. Near the instability threshold for γ_m we have

$$\gamma_m = \frac{\omega_{Le}^2}{2\omega_0^2}\nu_{ei}\left(\frac{v_E^2}{v_{E,\,\mathrm{thr}}^2} - 1\right). \qquad (5)$$

If the threshold is exceeded by a fair amount, with $v_E/c \gg \nu_{ei}/\omega_0$, the growth rate will aslo peak at $\varphi_m = 90° \pm 45°$. This instability growth rate maximum γ_m is

$$\gamma_m = \frac{1}{10a} \frac{v_E^2}{c^2} \frac{\omega_{Le}^2}{\nu_{ei}} \quad . \tag{6}$$

The wave vector of the scattered waves $\mathbf{k} \pm \mathbf{k}_0$, which forms a 90° angle with the wave vector $\mathbf{k}_0$, lies in the plane perpendicular to the vector $\mathbf{E}_0$. The scattered wave in this case is polarized in the same plane, which is a direct consequence of the scattering by magnetic fluctuations. Equations (4)–(6), strictly speaking, correspond to a case in which the plasma is completely ionized. The results obtained can, however, be easily extended to a slightly ionized plasma. The quantity ν_{ei} is then taken to mean the frequency of the electron-neutral particle collisions and a should be replaced by ≈ 1.

The parametric instability which we have detected corresponds to the scattering of an electromagnetic pump wave by quasistatic magnetic-field perturbations. This process, which is characterized by a side scattering and a 90° rotation of the electromagnetic-wave polarization vector, occurs without a frequency shift. These systematic features, taken collectively, set this scattering mechanism apart from the previously known mechanisms (SBS, SRS, STS).

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