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Laser excitation of surface waves over a dense plasma

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An intense short laser pulse incident on a metallic target produces an overdense plasma and excites a plasma surface wave via stimulated Compton scattering. The pump and the surface wave exert a beat-frequency ponderomotive force on the electrons in the skin layer, driving a heavily damped quasimode. The density perturbations due to the quasimode couple with the oscillatory velocity due to the pump to drive the surface wave. The growth period for the process turns out to be about 0.5 ps at a Nd:glass laser intensity of 10^{16} W cm^{-2} and a plasma density of 1.5 times the critical value.

1. Introduction

The interaction of intense short laser pulses with metallic targets has revealed several interesting features, including strong reduction of reflectivity with intensity (Milchberg *et al.* 1988; Chaker *et al.* 1991; Ng *et al.* 1994), excitation of instabilities (Saurbrey *et al.* 1994), emission of higher harmonics (Esarey *et al.* 1993; Rax and Fish 1993; Shen *et al.* 1995), generation of strong DC magnetic fields (Sudan 1993) and X-ray emission (Milchberg and Freeman 1990; Milchberg *et al.* 1991; Eder 1995). Understanding these phenomena involves understanding laser-energy deposition and heat-energy transport schemes. Surface waves could play a significant role in laser absorption and energy transport along the target surface. Early studies (Zoboronkova *et al.* 1976; Yasumoto and Noguchi 1982; Lindgren *et al.* 1982) suggest that surface waves can be excited by parametric decay of a laser into two surface waves.

In this paper we study the excitation of a surface wave via stimulated Compton scattering (Sagdeev and Galeev 1969; Nishikawa and Liu 1976; Aliev *et al.* 1995) of the laser over the surface of a dense plasma created by laser irradiation of a metallic target (Rozmus and Tikhonchuk 1990; Price *et al.* 1995). In this process, the laser decays into a surface wave and a heavily damped quasimode inside the skin layer. The pump and the surface wave exert a beat-frequency ponderomotive force on the electrons, driving a heavily damped quasimode inside the skin layer. The density perturbation due to the quasimode couples with the oscillatory velocity due to the pump to produce a nonlinear current that drives the surface wave. The process is similar to surface Raman scattering from metallic films, where the incident laser couples to a plasma surface wave and a purely growing surface perturbation (ripple). The fields of the laser and the surface wave cause stronger etching in regions of

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surface depressions. The depressions tend to trap the laser and drive the surface wave (Oron and Sorenson 1978; Ehrlich *et al.* 1982). This process has been used for producing grating and dot structures, useful in gate devices. This process is important at moderate laser intensities; however, it occurs on a lower time scale. It may not be important in picosecond laser pulses.

In Sec. 2 we discuss the excitation of plasma surface waves by stimulated Compton scattering. In Sec. 3 we discuss the results.

2. Surface-wave excitation by stimulated Compton scattering

Consider a free-space–dense-plasma (or metal) interface at $x = 0$, with the plasma occupying $x > 0$. A laser pulse is obliquely incident on the interface from the vacuum side. It is polarized in the plane of incidence. The field of the incident laser for $x < 0$ is written as

$$\mathbf{E}_{0i} = \mathbf{A}_0' \exp\left\{-i\left[\omega_0 t - \frac{\omega_0}{c}\sin\theta_0 z - \frac{\omega_0}{c}\cos\theta_0 x\right]\right\}, \tag{1}$$

where θ_0 is the angle of incidence and

$$A_{0x} = -A_{0z}\tan\theta_0.$$

The field inside the plasma can be written as (Jackson 1975)

$$\mathbf{E}_0 = \mathbf{A}_0 e^{-\alpha_0 x} e^{-i(w_0 t - k_{0z} z)}, \tag{2}$$

where

$$\alpha_0 = \frac{(\omega_p^2 - \omega_0^2\epsilon_L + k_{0z}^2 c^2)^{1/2}}{c} = \frac{[\omega_p^2 - \omega_0^2(\epsilon_L - \sin^2\theta_0)]^{1/2}}{c},$$

$$k_{0z} = \frac{\omega_0}{c}\sin\theta_0,$$

$$A_{0z} = A_{0z}' T, \qquad E_{0x} = \frac{ik_{0z}}{\alpha_0} E_{0z},$$

$$T = 2\left(1 + \frac{\epsilon_0\omega_0\cos\theta}{i\alpha_0 c}\right),$$

$\epsilon_0 = e_L - \omega_p^2/\omega_0^2, \epsilon_L$ is the metal lattice permittivity (for a plasma $\epsilon_L = 1$), $\omega_p^2 = 4\pi n_0^0 e^2/m$, n_0^0 is the plasma density, and $-e$ and m are the electronic charge and mass respectively. Here we have assumed that the collision frequency $\nu \ll \omega$. The laser field vanishes beyond a few skin depths; the skin depth $\delta \equiv \alpha_0^{-1}$.

Inside the skin layer the electrons acquire a large oscillatory velocity

$$\mathbf{v}_0 = \frac{e\mathbf{E}_0}{mi\omega_0}. \tag{3}$$

The laser decays into a surface wave (ω_1, k_1) (Liu and Tripathi 1994) and a lower-frequency space-charge quasimode (ω, k), whose fields, for $x > 0$, are $\mathbf{E}_1$ and $\mathbf{E} = -\nabla\phi$;

$$\mathbf{E}_1 = \mathbf{E}_1 e^{-\alpha_1 x} e^{-i(\omega_1 t - k_{1z} z)}, \tag{4}$$

$$\phi = \phi\, e^{-i(\omega t - k_z z)}. \tag{5}$$

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Here

$$\alpha_1^2 = k_{1z}^2 - \frac{\omega_1^2 \epsilon_L - \omega_p^2}{c^2}.$$

The phase-matching conditions demand that

$$\omega_1 = \omega_0 - \omega, \tag{6a}$$

$$k_{1z} = k_z - k_{0z}. \tag{6b}$$

The surface wave induces an oscillatory velocity on the electrons,

$$\mathbf{v}_1 = \frac{e\mathbf{E}_1}{mi\omega_1}. \tag{7}$$

The pump (ω_0, k_{0z}) and the surface wave (ω_1, k_{1z}) exert a ponderomotive force on electrons at (ω, k_z):

$$\mathbf{F}_p = e\boldsymbol{\nabla}\phi_p = -\tfrac{1}{2}m\boldsymbol{\nabla}(\mathbf{v}_0 \cdot \mathbf{v}_1)$$

where

$$\phi_p = -\frac{m}{2e}(\mathbf{v}_0 \cdot \mathbf{v}_1). \tag{8}$$

ϕ and ϕ_p produce an electron-density perturbation n at $(\omega, \mathbf{k})$. Following Liu and Tripathi (1994), n can be written as

$$n = \frac{k^2}{4\pi e}\chi_e(\phi + \phi_p), \tag{9}$$

where

$$\chi_e = \frac{2\omega_p^2}{k^2 v_{\text{th}}^2}\left[1 + \frac{\omega}{k_z v_{\text{th}}} Z\left(\frac{\omega}{k_z v_{\text{th}}}\right)\right]$$

is the electron susceptibility at $(\omega, \mathbf{k})$, v_{th} is the electron thermal speed and $Z(\xi)$ is the plasma dispersion function.

The ion response can be written as

$$n_i = -\frac{k^2}{4\pi e}\chi_i \phi, \tag{10}$$

where $\chi_i = -\omega_{pi}^2/\omega^2$ is the ion susceptibility and ω_{pi} is the ion plasma frequency. In the case of a metal $\chi_i = 0$.

Using n and n_i in Poisson's equation, we get

$$\phi + \phi_p = \frac{1 + \chi_i}{\epsilon}\phi_p, \tag{11}$$

where $\epsilon = 1 + \chi_e + \chi_i$. Using (11) in (9), n can be rewritten as

$$n = \frac{k^2}{4\pi e}\frac{\chi_e(1 + \chi_i)}{\epsilon}\phi_p. \tag{12}$$

The nonlinear current density at (ω_1, k_{1z}) is

$$\mathbf{J}_1^{NL} = -\tfrac{1}{2}ne v_0^*. \tag{13}$$

Here we have neglected the contribution due to the surface charges, since it is ω/ω_0 times smaller than the nonlinear term in (13).

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Using (13), the wave equation for $\mathbf{E}_1$ can be written as

$$\nabla^2\mathbf{E}_1-\boldsymbol{\nabla}(\boldsymbol{\nabla}\cdot\mathbf{E}_1)+\frac{\omega_1^2-\omega_p^2}{c^2}\mathbf{E}_1=-\frac{\mathbf{v}_0^*}{4c^2}\frac{\chi_e(1+\chi_i)}{4\epsilon}k^2\mathbf{v}_0\cdot\mathbf{E}_1. \tag{14}$$

Taking the divergence of this equation, we get

$$\boldsymbol{\nabla}\cdot\mathbf{E}_1=-\frac{k^2}{\omega_1^2-\omega_p^2}\frac{\chi_e(1+\chi_i)}{4\epsilon}\boldsymbol{\nabla}\cdot(\mathbf{v}_0^*\,\mathbf{v}_0\cdot\mathbf{E}_1). \tag{15}$$

Using (15) in (14), we get

$$\nabla^2\mathbf{E}_1+\frac{\omega_1^2-\omega_p^2}{c^2}\mathbf{E}_1=-\frac{k^2}{c^2}\frac{\chi_e(1+\chi_i)}{4\epsilon}\left\{\mathbf{v}_0^*\,\mathbf{v}_0\cdot\mathbf{E}_1+\frac{c^2}{\omega_1^2-\omega_p^2}[\boldsymbol{\nabla}\cdot(\mathbf{v}_0^*\,\mathbf{v}_0\cdot\mathbf{E}_1)]\right\}. \tag{16}$$

In the absence of the right-hand side, the solution of (16) can be written as

$$\mathbf{E}_1=A_1\,\boldsymbol{\psi}_1, \tag{17}$$

where

$$\boldsymbol{\psi}_1=\begin{cases}\left(\hat{\mathbf{z}}+\dfrac{ik_{1z}}{\alpha_1}\hat{\mathbf{x}}\right)e^{-\alpha_1 x} & \text{for } x>0,\\[2ex] \left(\hat{\mathbf{z}}-\dfrac{ik_{1z}}{\alpha_1^1}\hat{\mathbf{x}}\right)e^{\alpha_1^1 x} & \text{for } x<0,\end{cases}$$

where

$$\alpha_1^2=k_{1z}^2-\frac{\omega_1^2\epsilon_L-\omega_p^2}{c^2},\quad \alpha_1^{12}=k_{1z}^2-\frac{\omega_1^2}{c^2}.$$

The dispersion relation for the surface wave is

$$k_{1z}^2=\frac{\omega_1^2}{c^2}\frac{\epsilon_1}{1+\epsilon_1}, \tag{18}$$

where $\epsilon_1=\epsilon_L-\omega_p^2/\omega_1^2$. We denote the solution of (18), for a given k_{1z}, by ω_{1r}, the eigenfrequency in the linear case. With this value of ω_1, in the linear case α_1 can be written as α_{1r}. In the presence of the right-hand side, we assume that the eigenfunction $\boldsymbol{\psi}_1$ remains the same – only the eigenvalue changes. Using (17) in (16), multiplying the resulting equation by $\boldsymbol{\psi}_1^*$ and integrating over x, we get

$$\omega_1^2-\omega_{1r}^2=-\frac{1}{N}\frac{k^2\chi_e(1+\chi_i)}{4\epsilon}I, \tag{19}$$

where

$$I=|v_{0z}|^2\left(1-\frac{k_{0z}k_{1z}}{\alpha_0\alpha_1}\right)^2\frac{1}{2(\alpha_0+\alpha_1)}\left[1-\frac{c^2k_{1z}}{\omega_1^2-\omega_p^2}\frac{2\alpha_0}{\alpha_1}\frac{k_z+k_{0z}-\dfrac{\alpha_1}{\alpha_0}k_{0z}}{1-\dfrac{k_{0z}k_{1z}}{\alpha_0\alpha_1}}\right],$$

$$N=\left[\left(1+\frac{k_{1z}^2}{\alpha_1^{12}}\right)\frac{1}{2\alpha_1^2}+\left(1+\frac{k_{1z}^2}{\alpha_1^2}\right)\frac{1}{2\alpha_1}\right].$$

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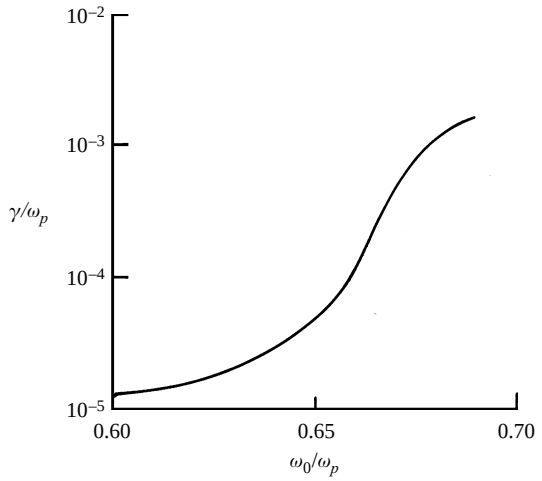


Figure 1. Variation of growth rate with laser frequency for a laser power of 10^{16} W cm^{-2}, a wavelength of 1 μm and a plasma temperature $T_e = 300$ eV.

We choose $\omega_1 = \omega_{1r} + i\gamma$, and, on substituting into (19), we obtain the growth rate

$$\gamma = \frac{k^2}{N}\frac{I}{8\omega_0}(1+\chi_i)^2\frac{\epsilon_i}{\epsilon_r^2+\epsilon_i^2}, \tag{20}$$

where ϵ_r and ϵ_i are the real and imaginary parts of ϵ respectively. For the validity of the perturbation theory, $\gamma < \omega$.

For a plasma (Price *et al.* 1995) with $T_e \approx 300$ eV and a Nd:glass laser with $\lambda \approx 1\,\mu$m and $P_0 = 10^{16}$ W cm^{-2}, the growth rate turns out to be about 2×10^{12} s^{-1} at $\omega_0 = 0.7\omega_p$. In Fig. 1 we show the variation of γ with ω_0. The growth rate rises sharply from $\gamma \approx 10^{-5}\omega_p$ at $\omega_0 = 0.6\omega_p$ to $10^{-3}\omega_p$ at $\omega_0 = 0.65\omega_p$. At larger values of the pump frequency, $0.65\omega_p < \omega_0 < \omega_p/\sqrt{2}$, the growth rate reaches saturation.

3. Discussion

Excitation of surface waves via stimulated Compton scattering is sensitive to laser frequency. For a Nd:glass laser with power density $P_0 = 10^{16}$ W cm^{-2} at 1 μm wavelength incident on a plasma with $T_e \approx 300$ eV, the growth rate $\gamma \approx 2\times10^{12}$ s^{-1} at $\omega_0 \approx \omega_p/\sqrt{2}$. We see a sharp variation in growth rate with ω_0. As ω_0 increases beyond $\omega_0 = 0.6\omega_p$ to $\omega_0 = 0.65\omega_p$, the growth rate increases by two orders of magnitude. Near $\omega_0 = \omega_p/\sqrt{2}$, the growth rate saturates. For $\omega_0 > \omega_p/\sqrt{2}$, nonlinear Landau damping is kinematically forbidden. Because of the finite extent of the region of interaction, the instability could suffer convective losses. The convective threshold may be given as $\gamma r_0/v_{g1} \geqslant 1$ or $\gamma \geqslant (c/r_0)\cos\theta_i$, where r_0 is the laser spot size. For $\gamma \approx 2\times10^{12}$ s^{-1}, one needs $r_0 \geqslant 100\,\mu$m. The Compton scattering process is liable to produce hot electrons travelling along the target surface. These electrons may create a dense plasma of greater dimensions than the laser spot size, and may give rise to stronger X-ray emission.

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The present treatment is restricted to planar surfaces, and may not apply to cylindrical surfaces, where parametric instability could be significantly different (Gradov and Stenflo 1983). However, the study could be relevant to ionospheric backscattering experiments when steep density gradients are realized (Stenflo *et al*. 1978).

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