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# Parametric generation of surface deformations in laser interaction with overdense plasmas

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**Abstract**

By analytical modeling and numerical simulation, we show that surface modes in moderately overdense plasmas may be excited parametrically by an intense, ultrashort laser pulse. This process has a feedback effect on fast electron generation and may seed a fast distortion of plasma "moving mirrors."

**Keywords:** Laser beam effects; Parametric instability; Plasma-beam interactions; Plasma production by laser; Plasma simulation

## 1. INTRODUCTION

The interaction of subpicosecond, high-intensity laser pulses with solid targets leads to outstanding nonlinear phenomena such as production of energetic "fast" electrons and generation of high harmonics (HHG) of the incident pulse frequency $\omega$. Both these phenomena may be regarded as "surface effects," since the laser–plasma coupling occurs near the "critical" surface where $n_e = n_c$ (where $n_c = m_e\omega^2/4\pi e^2 \ll n_{solid}$ is the cut-off density for laser propagation) and over a narrow region with a depth of the order of the skin length $d_p < c/\omega$. In turn, the laser–plasma interaction is extremely sensitive to geometry (Ruhl *et al.*, 1999) and, eventually, to surface deformations, both preimposed (e.g., microstructured targets; see Kulcsar *et al.*, 2000, and references therein) or self-generated during the interaction. For instance, absorption in a one-dimensional (1D) geometry is limited by strong constraints (Mulser *et al.*, 2001) that may be removed by the onset of 2D perturbations at the surface. Simulations suggest a spatial correlation between surface rippling and fast electron jets (Macchi *et al.*, 2000). Plasma surface rippling at high laser intensities also leads to the spatial spread of high harmonic emission in experiments (Norreys *et al.*, 1996; Tarasevitch *et al.*, 2000), that is detrimental for applications. Both experimental and simulation results suggest that the surface perturbations grow on short time scales, much faster than ion motion, and are therefore of electronic nature.

In this article, we use particle-in-cell (PIC) simulations (Macchi *et al.*, 2001) to show that electron surface structures growing in steplike density profiles can be due to a parametric three-wave process, namely, the "decay" of a laser-driven oscillation at $\omega$ or $2\omega$ into two surface modes. We also outline an analytical model, based on a cold fluid plasma and nonrelativistic approximation, that shows that the basic process is understood as the excitation of a pair of counterpropagating surface waves (Macchi *et al.*, 2002).

## 2. NUMERICAL EXPERIMENTS

The 2D dynamics of the surface motions was well resolved in 2D PIC simulations. In all simulations, we took $\lambda = 0.25\ \mu$m, that is, $n_c = 1.6 \times 10^{22}$ cm$^{-3}$, normal incidence and uniform intensity of the laser pulse, immobile ions, and a steplike initial density profile. Details are given in Macchi *et al.* (2001). We investigated cases with densities $n_e = 3$–$5n_c$ and irradiances $a_i = 0.85$–$1.7$ where $a_i = 0.85\ (I\lambda^2/10^{18}\ \mathrm{W}\mu\mathrm{m}^2/\mathrm{cm}^2)^{1/2}$.

The temporal evolution of the electron density $n_e(x, y)$ is shown in Figure 1a,b for $n_e = 5n_c$ and $a_i = 1.7$ and Figure 1c

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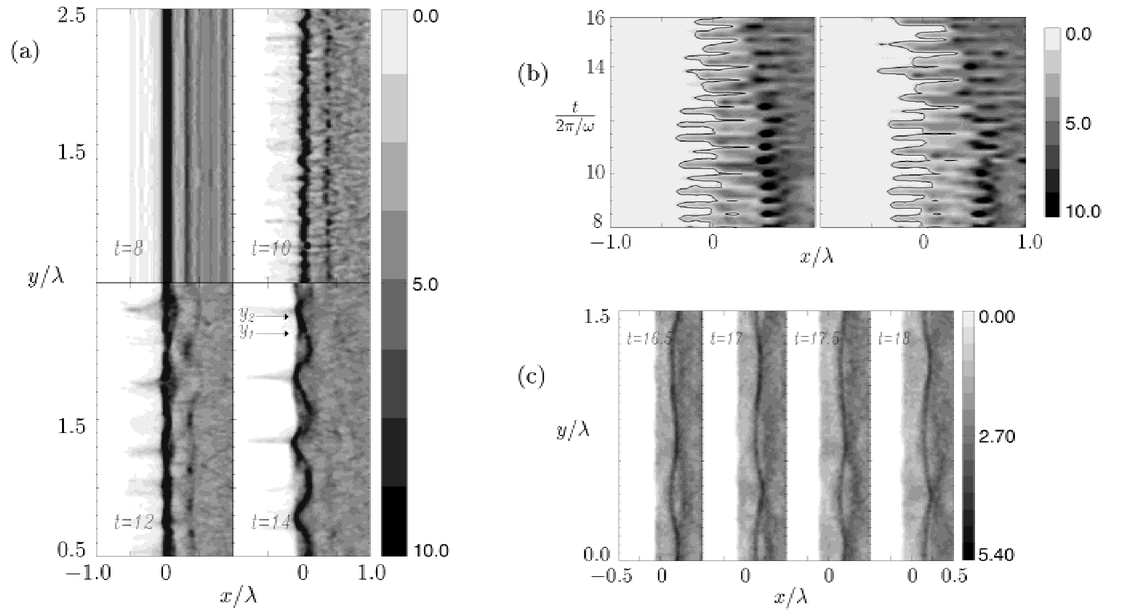


**Fig. 1.** a, c: Contours of $n_e/n_c$ for $a_i = 1.7$, $n_0 = 5n_c$, at various times (see labels) in laser cycle units, for $a_i = 1.7$, $n_o/n_c = 5$ (a) and $a_i = 0.85$, $n_o/n_c = 3$ (c). b: Space-time plot of $n_e(x, y_i, t)/n_c$ at $y = y_1$ (left) and $y = y_2$ (right), for the case (a).

for $n_e = 3n_c$ and $a_i = 0.85$. In all these runs, the laser was $s$-polarized. Initially, the surface oscillation is planar, that is, uniform along $y$ (first plot of Fig. 1a) and has a frequency $2\omega$ (Fig. 1b), where $\omega = 2\pi c/\lambda$ is the laser frequency. This 1D motion is driven by the longitudinal $\mathbf{j} \times \mathbf{B}$ force at $2\omega$. Later, we observe the growth of surface structures having small wavelengths ($\simeq 0.1\lambda$) in a first stage and then evolving into a standing (i.e., not propagating in $y$) oscillation with wavelength $\lambda_s \simeq 0.5\lambda$ and frequency $\simeq \omega$, which is superimposed to the oscillation at $2\omega$; see plots (a) and (b) of Figure 1 for $a_i = 1.7$ and $n_0/n_c = 5$. Figure 1c shows the "snaking" within one laser cycle at frequency $\omega$ of the surface layer for $a_i = 0.85$ and $n_0/n_c = 3$; in this case, the oscillation amplitude and the density compression are lower and the deformation wavelength is larger ($\lambda_s \simeq 0.75\lambda$). In the following we will discuss only the long-wavelength structures oscillating at $\omega$.

The generation of surface structures oscillating at $\omega$ may at first be surprising, since there is no longitudinal component of the laser force at such a frequency. Our simulation results suggest that the 1D, electrostatic oscillation ($k_0 = 0$, $\omega_0 = 2\omega$) "decays" into two surface modes $(k_1, \omega_1)$ and $(k_2, \omega_2)$. The "matching conditions" for this process give $k_1 = -k_2$ and $\omega_1 = \omega_2 = \omega$. The two overlapping modes thus form a *standing* oscillation with frequency $\omega$ and wave vector $k = k_1 = -k_2$. Note that this reasoning is *simply based on the existence of transverse surface modes* with some dispersion relation.

For high intensities, density "spikes" are generated at the surface of the plasma. Figure 2a shows a detail of the density plot of Figure 1a at late times, with the electron fluid velocity field $\mathbf{v}_e$ superimposed. Notice the radial structure of $\mathbf{v}_e$ inside the spike, which suggests a sort of Coulomb explosion. Figure 2b shows density contours for a simulation

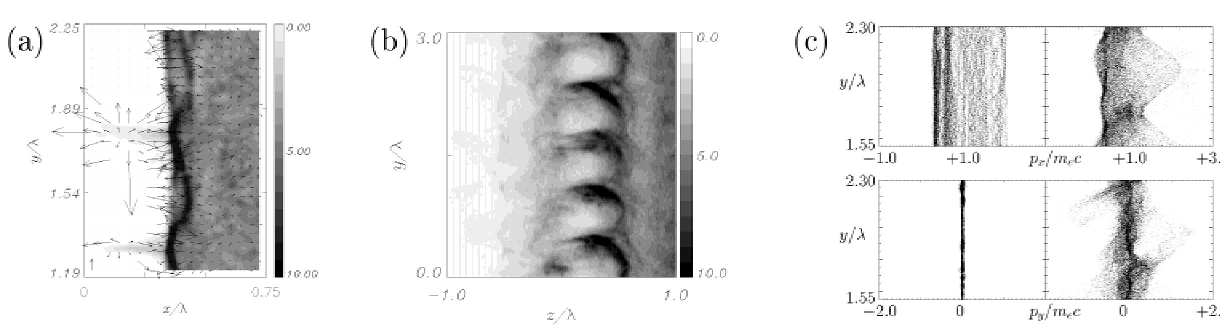


**Fig. 2.** a: Detail of the density plot of Figure 1 at $t = 14$ with the electron fluid velocity field $\mathbf{v}_e$ superimposed. b: Density contours at $t = 15$ laser cycles for $a_i = 1.2$, $n_e/n_c = 5$, and $p$-polarization. c: Phase space projections in the $(y, p_x)$ and $(y, p_y)$ at $t = 8.5$ (left) and $t = 13.5$ (right) laser cycles, for $a_i = 1.7$ and $n_e/n_c = 5$.

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with $n_e = 5n_c$, $a_i = 1.2$ and $p$-polarization. In this case, one observes a transverse bending of the density spikes by the electric field lying in the simulation plane.

The onset of 2D surface oscillations has a substantial impact on fast electron generation. Phase space projections (Fig. 2c) show that at early times the momentum distribution is uniform in $y$, with no accelerated particles in $p_y$ and most energetic electrons having $p_x \simeq 2m_e c$. At later times, stronger forward acceleration occurs near the maxima of the 2D oscillation, showing that most oscillatory energy has been transferred to the 2D modes; strong transverse acceleration in $p_y$ also occurs. One consequence of this correlation is that surface oscillations give an "imprint" on the transverse structure of the fast electron currents, that may seed the growth of current filamentation instabilities (Califano *et al.*, 1998) observed in simulations (Lasinski *et al.*, 1999; Sentoku *et al.*, 2000).

## 3. PARAMETRIC EXCITATION OF SURFACE WAVES

The driving mechanism of the 2D surface oscillations is a parametric, three-wave process involving the excitation by the laser of a pair of electron surface waves (ESWs). Examples of similar two-surface wave decay (TSWD) processes were considered so far in very different regimes and in the electrostatic limit only (Gradov & Stenflo, 1981; Stenflo, 1996). We recently showed that TSWD allows the excitation of fully *electromagnetic* SWs in a planar, steplike density profile (while linear mode conversion of the laser wave into a single ESW requires in general tailored density profiles) and it can be driven either by the electric or magnetic components of the Lorentz force (Macchi *et al.*, 2002). In the following, we give a brief sketch of the analytical modeling and its results, while details are reported in Macchi *et al.* (2002).

In our model, we consider a steplike plasma density profile $n_i(x) = n_i \Theta(x)$, where $\Theta(x)$ is the Heaviside step function. The laser pulse has linear polarization and wavevector $\mathbf{k} = (\omega_c)(\hat{\mathbf{x}} \cos\theta + \hat{\mathbf{y}} \sin\theta)$. Translational invariance along $z$ is assumed. In the Maxwell–Euler cold plasma equations, we adopt the following expansion for all fields:

$$F(x,y,t) = F_i(x) + \epsilon F_0(x, t - y\sin\theta/c) + \epsilon^2[f_+(x,y,t) + f_-(x,y,t)], \qquad (1)$$

where $\epsilon$ is a small expansion parameter and $f$ stands for either the electron density or velocity or for the EM fields in the $(x,y)$ plane. In this expansion, $F_i$ represents unperturbed fields of zero order (e.g., $n_i$). The term $F_0$ represents the "pump" field at the frequency $\omega_0$ (yet to be specified), that is taken to be of order $\epsilon$ and is written as

$$F_0 = \tilde{F}^{(\omega_0)}(x) e^{i\omega_0(t - y\sin\theta/c)} + \text{c.c.} \qquad (2)$$

The last term is the sum of counterpropagating surface modes, which are assumed to be of order $\epsilon^2$ and are written as

$$f_\pm = \Re[\tilde{f}_\pm^{(\omega_\pm)}(x) e^{ik_\pm y - i\omega_\pm t}]. \qquad (3)$$

Using expansion (1), the coupling between the pump and the surface modes is of order $\epsilon^3$. Thus, from the Maxwell–Euler equations, we obtain, to order $\epsilon$, the pump fields and, to order $\epsilon^2$, the dispersion relations for ESWs in a cold plasma:

$$k_\pm^2 = \frac{\omega_\pm^2}{c^2} \frac{\omega_p^2 - \omega_\pm^2}{\omega_p^2 - 2\omega_\pm^2}. \qquad (4)$$

The dispersion relation (4) is shown in Figure 3a. All feedback effects of the nonlinear coupling on the pump fields are neglected. The coupling terms that lead to the excitation of the ESW with frequency $\omega_\pm$ may be represented by the nonlinear force

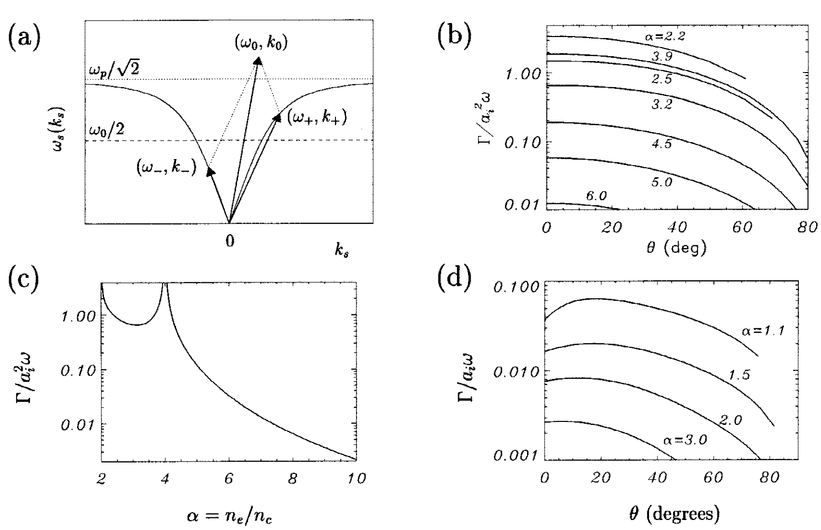


**Fig. 3.** a: Dispersion relation of electron surface waves (thick lines), Eq. (4), and the matching conditions for TSWD represented by arrows. b–c: The TSWD growth rate $\Gamma$ (normalized to $a_i^2\omega$) for the $2\omega \to \omega + \omega$ process as a function of $\alpha = n_e/n_c$ and $\theta$ (b), Eq. (6), and as a function of $\alpha$ for $\theta = 0$ (c). d: The TSWD growth rate $\Gamma$ (normalized to $a_i\omega$) for the $\omega \to \omega/2 + \omega/2$ process at oblique incidence, as a function of $\theta$ and $\alpha$.

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$$\mathbf{f}_{\pm}^{(NL)} = \tilde{\mathbf{f}}_{\pm}^{(NL)}(x) e^{ik_{\pm} y - i\omega_{\pm} t}$$
$$= -\left[ m_e(\mathbf{v}_{\mp}\cdot\nabla\mathbf{v}_0 + \mathbf{v}_0\cdot\nabla\mathbf{v}_{\mp}) + \frac{e}{c}(\mathbf{v}_0 \times \mathbf{B}_{\mp} + \mathbf{v} \times \mathbf{B}_0)\right]_{res}. \tag{5}$$

On the r.h.s. of Eq. (5) only "resonant" terms, having the same phase of $\mathbf{f}_{\pm}^{(NL)}$, are included. Obviously, these terms exist if the matching conditions $k_0 = k_+ + k_-$ and $\omega_0 = \omega_+ + \omega_-$ hold. These relationships are shown in Figure 3a. The "pump" mode may be driven by the laser pulse either by the electric or the magnetic ($\mathbf{v} \times \mathbf{B}$) force terms, which have frequencies $\omega_0 = \omega$ and $\omega_0 = 2\omega$, respectively. Thus, from the matching conditions, one obtains that these processes correspond to a "$\omega \to \omega/2 + \omega/2$" and to a "$2\omega \to \omega + \omega$" decay, respectively, since one may write $\omega_{\pm} = \omega_0/2 \pm \delta\omega$ where the shift $\delta\omega$ vanishes at normal incidence for symmetry reasons.

Figure 3b shows the analytically calculated growth rate $\Gamma$ for the "$2\omega \to \omega + \omega$" process at oblique laser incidence as a function of $\theta$ and $n_e/n_c$. In this case, $\Gamma$ is proportional to $a_i^2\omega$, that is, to the laser intensity. Figure 3c shows $\Gamma$ as a function of $n_e/n_c$ for $\theta = 0$, that is given approximately by

$$\Gamma \simeq 4\omega a_i^2 \frac{(\alpha - 1)^{3/2}}{|\alpha - 4|[(\alpha - 1)^2 + 1](\alpha - 2)^{1/2}}. \tag{6}$$

The leading term in the nonlinear pump force is found to be $(m_e(\mathbf{v}_{\mp}\cdot\nabla v_x^{(2\omega)}\hat{\mathbf{x}} + v_x^{(2\omega)}\hat{\mathbf{x}}\cdot\nabla\mathbf{v}_{\mp}))$, which is most efficient at $\theta = 0$. The divergence at $n_e = 4n_c$ is due to the resonant excitation of plasmons by the longitudinal force at $2\omega$ making $v_x^{(2\omega)}$ very large, while that for $n_e \to 2n_c$ is due to the fact that ESWs have zero energy in this limit. This also corresponds to $k_{\pm} \to \pm\infty$, thus in this limit, one expects a strong thermal damping of ESWs (neglected in our cold approximation), smoothing the resonance out.

Figure 3c gives $\Gamma$ for the "$\omega \to \omega/2 + \omega/2$" process driven by a $p$-polarized laser pulse. The growth rate is now proportional to the laser field amplitude $a_i$. Again, the TSWD is driven mostly by the longitudinal electron motion resulting in a maximum of $\Gamma$ at some "optimal" angle. However, the pump force does not vanish for $\theta = 0$, resulting in $\Gamma \neq 0$ even at normal incidence.

The analytical modeling suggests that the driving mechanism of the 2D oscillations observed in the simulations is due to a $2\omega \to \omega/ + \omega$ process, and predicts spatial and temporal scales which reasonably agree with the simulations in regimes where our approximations are valid (i.e., for nonrelativistic intensities $a_i < 1$). For higher intensities, the observed wavelengths of the surface modes are smaller than predicted by the simple cold fluid TSWD model, suggesting that kinetic and relativistic effects play an important part.

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