Two-surface wave decay: improved analytical theory and effects on electron acceleration

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Abstract

Two-surface wave decay (TSWD), i.e. the parametric excitation of electron surface waves, was recently proposed as an absorption mechanism in the interaction of ultrashort, intense laser pulses with solid targets. We present an extension of the fluid theory of TSWD to a warm plasma which treats boundary effects consistently. We also present test-particle simulations showing localized enhancement of electron acceleration by TSWD fields; this effect leads to a modulation of the current density entering into the target and may seed current filamentation instabilities.

PACS numbers: 52.38.-r; 52.38.Dx
I. INTRODUCTION

The excitation of electron surface waves (ESWs) is a possible route to collisionless energy absorption in plasmas produced by the interaction of sub-picosecond, high intensity (typically $\geq 10^{16}$ W/cm$^2$) laser pulses with solid targets. This regime is relevant for applications such as generation of ultrashort X-ray pulses, either as uncoherent thermal emission ([1] and references therein) or high laser harmonics [2], or production of energetic electrons and ions [1, 3, 4, 5].

Linear mode conversion of the laser wave into an ESW is not possible at a planar plasma-vacuum interface, because the phase matching between the ESW and the incident wave is not allowed; the process can take place for specially tailored density profiles, e.g. grating targets with a density profile modulated at the surface with a wavevector $k_p$. In this case, the condition for the excitation of a ESW with wavevector $k_s$ (parallel to the plasma surface) and the same frequency $\omega_L$ of the laser pulse is

$$k_L \sin \theta = k_s + k_p,$$

where $k_L = \omega_L/c$, $\theta$ is the angle of incidence of the laser pulse. Experimental investigations of laser absorption and X-ray emission in grating targets are described, e.g., in Refs. [6, 7].

**Nonlinear** mode conversion, i.e. parametric excitation of two ESWs is possible also in a simple “step” profile. We name this process “two–surface wave decay” (TSWD) [8]. Examples of similar TSWD processes were previously considered in regimes other than intense laser–plasma interactions and in the electrostatic limit only [9, 10]. The general phase matching conditions for such a three–wave process are given by

$$k_0 = k_+ + k_- , \quad \omega_0 = \omega_+ + \omega_-,$$

where $k_0, \omega_0$ and $k_\pm, \omega_\pm$ are the wavevector and frequency of the “pump” wave and the two ESWs, respectively. If the electric field of the laser pulse acts as a “pump” for TSWD, then $\omega_0 = \omega_L$, $k_0 = k_L \sin \theta$, and equation (2) implies that two sub–harmonic ESW are generated with frequency around $\omega/2$ (“$\omega \rightarrow \omega/2 + \omega/2$” TSWD). However, also the $v \times B$ term of the Lorentz force may excite TSWD. This was observed in particle-in-cell (PIC) simulations [11] for normal laser incidence; in this case, the $v \times B$ force drives 1D oscillations at the plasma surface with frequency $2\omega_L$; after a few laser cycles the overlap of a standing wave
with frequency $\omega_L$ was observed. This is a clear signature of a “$2\omega \rightarrow \omega + \omega$” process leading to two counterpropagating ESWs both having the frequency of the laser pulse.

According to the theoretical model [8, 12], the maximum growth rate of the $2\omega \rightarrow \omega + \omega$ case is found for normal incidence. As said above, TSWD does not need a structured target; however, it is worth noticing that in a grating target the wavevector $k_p$ of the surface modulation [eq. (1)] equals the wavevector of the two ESWs excited by the $2\omega \rightarrow \omega + \omega$ process at $\theta = 0$. Hence, TSWD is enhanced in such a grating target. It is then interesting to notice that PIC simulations reported in Ref. [7], showing the generation of two countepropagating ESWs in grating targets at $10^{16}$ W/cm$^2$, can now be interpreted as an evidence of TSWD seeded by the surface grating.

At very high intensities ($\geq 10^{18}$ W/cm$^2$), simulations [11] show that TSWD enters a strongly nonlinear regime leading to strong rippling of the plasma surface. This might be relevant to the surface instabilities which have been observed in experiments [13] and play a detrimental role in high harmonic generation from solid targets [14, 15, 16]; such instabilities appear to grow even for pulses of few tens of fs [16] and thus must be of electronic nature, since ions do not move on such a fast scale. Simulations also show that nonlinear TSWD affects fast electron generation. This effect is investigated in section III by means of test-particle simulations. Before that, in section II we present an improvement of the model of Ref. [8] where temperature effects have been included. In both cases, we restrict ourselves to the case of the $2\omega \rightarrow \omega + \omega$ TSWD at normal laser incidence, and we neglect collisions for simplicity and because collective processes dominate absorption at intensities $\geq 10^{16}$ W/cm$^2$.

II. TEMPERATURE AND BOUNDARY EFFECTS ON TSWD

In Ref. [8] the TSWD growth rate was calculated for a step-like density profile [$n_i = n_0 \Theta(x)$, being $\Theta(x)$ the Heaviside step function] using an Eulerian, fluid model with immobile ions and using the cold plasma approximation. We adopted the following expansion for all fields

$$f(x, y, t) = f_i(x, t - y \sin \theta/c) + \epsilon f_0(x)e^{i\omega_0(t - y \sin \theta/c)} + \epsilon^2 [f_+(x)e^{ik_yy - i\omega_+t} + f_-(x)e^{ik_yy - i\omega_-t}],$$

(3)
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. The first term ($f_i$) of eq.(3) includes zero-order, unperturbed fields or oscillating fields that are non-resonant with the excited modes; the term $f_0$ represents the “pump” field at the frequency $\omega_0$; the last term is the sum of two counterpropagating surface modes. For the $2\omega \rightarrow \omega + \omega$ process at $\theta = 0$, $\omega_0 = 2\omega_L$ and $k_+ = -k_-$, and the zero- and first-order fields do not depend on $y$ (i.e. they are “1D” fields). The coupling between the pump and the surface modes (of order $\epsilon^3$) originates from the nonlinear terms in the Euler equation ($-en_e v \times B$, $-m_e n_e \nabla v$) and the current density $J = -en_e v$.

The calculation for a warm, isothermal plasma proceeds very similarly [17] and thus only the differences from the “cold” case and their consequences are discussed below, while details of the calculation will be published elsewhere. The only difference in the starting Maxwell-Euler systems of equations comes from the pressure term in the Euler equation for electrons:

$$m_e d\mathbf{v}/dt = -en_e (\mathbf{E} + \mathbf{v} \times \mathbf{B}) - k_B T_e \nabla n_e,$$

where $n_e$, $\mathbf{v}$ and $T_e$ are the density, fluid velocity and temperature of electrons.

Assuming reflective conditions at the plasma boundary ($x = 0$), the charge density $\rho = e(n_i - n_e) \neq 0$ only in a surface layer with a thickness on the order of the Debye length. At $T_e = 0$ this corresponds to a surface charge layer [i.e. $\rho \sim \delta(x)$, where $\delta(x)$ is the Dirac delta function], while the longitudinal velocity ($v_x$) is discontinous at the surface. This is relevant because in the $T_e = 0$ case nonlinear coupling terms involving the product of fields that are singular at $x = 0$ occur. The pressure term removes such singularities and enables to evaluate all such terms correctly. Therefore, by calculating the TSWD growth rate and then taking the $T_e \rightarrow 0$ limit one achieves an improvement of the result obtained in the cold plasma case.

The final result of the calculation is the $2\omega \rightarrow \omega + \omega$ growth rate $\Gamma$ shown in Fig.1 (the complete analytical expression is very lengthy [17] and is not reported here). With respect to the $T_e = 0$ case, the two divergences of $\gamma$ are both quenched for growing $T_e$, but for different reasons. The divergence for $\omega \rightarrow \omega_p/\sqrt{2}$ is quenched since in this limit the ESWs have shorter wavelengths and are thus more affected by the thermal pressure that inhibits the formation of small-scale structures. The “pump” resonance at $\omega \simeq \omega_p/2$, due to the excitation of plasmons by the $\mathbf{v} \times \mathbf{B}$ force, is quenched by energy transport out of the skin.
layer because for $T_e \neq 0$ the resonant plasmon propagates into the overdense plasma.

III. ELECTRON ACCELERATION

At high intensity of the laser pulse, it is well known that most of the absorbed energy goes into “fast” electrons injected into the overdense plasma region. At normal incidence, due to the leading action of the $\mathbf{v} \times \mathbf{B}$ force, fast electron bunches are produced twice per laser cycle. The PIC simulations of Ref.[11] showed that after the growth of nonlinear TSWD the generation of fast electrons was enhanced near spatial maxima of the standing surface oscillation (Fig.2 top). Thus, this effect may give an “imprint” for the formation of electron filaments, whose size and spacing would be close to (and scaling with) the laser wavelength as observed in other simulations [18, 19], and affect energy transport by fast electrons into the plasma.

To investigate this effect further, we performed “test particle” simulations of electron motion in the overlapping “pump” and TSWD fields. In other words, we solve the equation of motion for electrons into a force field of the form given by eq.(3), i.e. the sum of an one-dimensional (1D) force of frequency $2\omega$ (whose analytical expression is obtained from the theoretical model) and the force field from a standing 2D surface wave of frequency $\omega$. A similar study, focused on the acceleration of electrons by a single ESW, was recently reported [20].

The 1D and 2D fields vary in time as $\cos 2\omega t$ and $\sin \omega t$, so that the temporal maxima and minima of the ESW field are always coincident in time with maxima of the 1D field, as it was found in theory and simulation [8, 11]. In what follows, $x$ is the direction normal to the plasma surface located at $x = 0$ (the plasma occupies the $x > 0$ region) and $y$ is the direction of propagation of ESW. The field amplitudes are chosen so that the system remains far from relativistic conditions and the ESW field can be considered as a perturbation. The ratio between the amplitudes of standing SW and the $2\omega$ pump was varied as a parameter. For the simulations reported in this paper, in terms of normalized amplitudes, $a = eE/m_e\omega$ where $E$ is the amplitude of the electric field, the values are $a_L = 0.2$ for the laser field and $a_{ESW} = 0.019$ for the ESW field. The wavevectors of the laser wave and of the ESW are given by the known expressions for a cold plasma with density $n_e/n_c = \omega_p^2/\omega^2 = 5$. To initialize the simulations, we gave the particles an initial position $x(0) > 0$ such that the
evanescent fields are negligible at that point, and an initial velocity \( v_x(0) < 0 \) (on average \( v_x = -0.1c \)) so that the particles move towards the surface region, where the ESW fields are localized. The particles are distributed in \( y \) uniformly over a region of width \( \lambda_s \), where \( \lambda_s \) is the ESW wavelength.

The \((y, p_x)\) projection of the test particle phase space (Fig. 2, bottom) looks similar to the one from the fully self-consistent PIC simulation (Fig. 2, top), showing that electron acceleration is indeed enhanced near spatial maxima of the standing wave. The overlap of the SW fields with the \( \mathbf{v} \times \mathbf{B} \)-driven fields (homogeneous along the \( y \)-direction) thus leads to a modulation along \( y \) of the longitudinal momenta.

The enhancement (or quenching) of the accelerating field occurs once per laser cycle (because the ESW frequency equals the laser frequency), but with opposite phase between contiguous maxima of the standing ESW. This is shown by Fig. 3 where the complete \((x, p_x)\) phase space projection, showing electron bunches generated at \( 2\omega \) rate, is compared with the same plot but including now only the electrons whose starting position lies within an interval of \( \lambda/4 \) width around the \( y = \lambda/4 \) spatial maximum of the ESW. It is thus evident how the most energetic electrons in the bunches penetrating in the \( x > 0 \) region are generated near the ESW maximum and at \( \omega \) rate, i.e. with the ESW frequency. Taking only the electrons around the \( y = 3\lambda/4 \) spatial maximum gives an almost identical picture, except that the most energetic bunches are now out of phase by an angle of \( 2\pi/\omega \) with respect to those coming from \( y = \lambda/4 \).

Although the amplitude of ESW fields is 0.1 times the pump field amplitude, the relative modulation of the longitudinal momentum \( p_x \) is about 30%. This can be qualitatively explained as follows. The electrons acquire energy from an evanescent, oscillating field if their transit time across the region where the field has non-vanishing amplitude is shorter than an oscillation period; this is the condition for non-adiabaticity of electron motion with respect to the field. The ratio between the transit time and the oscillation period is approximately given by the parameter \( \eta = L/v_0T \) where \( L \) is the evanescence length, \( v_0 \) is the average velocity of the electron and \( T \) is the oscillation period. For the ESW, \( T = 2\pi/\omega \) and \( l_{ESW} = (c/\omega)\sqrt{\alpha - 2}/(\alpha - 1) \) (where \( \alpha = \omega_p^2/\omega^2 \)) while for the 1D field at \( T = \pi/\omega \) and \( l_{2\omega} = (c/2\omega)(1/\sqrt{\alpha - 1}) \) and \( \text{[8]} \). Thus, \( \eta_{ESW}/\eta_{2\omega} = \sqrt{(\alpha - 2)/(\alpha - 1)} < 1 \) which means that the electron motion (for a given \( v_0 \)) is more non-adiabatic with respect to the ESW field rather than to the \( 2\omega \) oscillation. Thus, the contribution of the ESW in accelerating
(or decelerating) electrons is enhanced by the relatively low frequency and short evanescence length.

The “imprint” effect of the standing ESW is also noticeable in the contourplot of the fluid velocity of electrons, shown in Fig.4. The velocity field has been computed by averaging the velocity over a spatial grid, as in PIC codes, and over time. As a consequence the electric current in the overdense plasma region has a spatial transverse modulation with the same wavelength of ESWs. The simulations of Ref.7 for “grating” targets also show a modulation of the fast electron current correlated with same period of the grating, that may be due to the local enhancement of the longitudinal field by TSWD or to the geometrical “funnel” effect of the surface deformation [21]. Our test particle results indeed show that due to TSWD a modulated fast electron current can be produced even by a plane–wave pulse on a flat surface.

From figures 2 and 3 it is also found that only near ESW spatial maxima a few electrons are ejected into vacuum, i.e. in the $x < 0$ region (in Fig.3 for these electrons the oscillation of $p_x$ vs. $x$ is due to the effect of the $v \times B$ force in vacuum). Their origin is likely to be due to the longitudinal field component directed into vacuum that is associated to the ESWs. In Ref.20, the features of electrons accelerated in vacuum by the field of a single ESW are investigated. Note, however, that electrostatic back-holding fields are not self-consistently included in test particle simulations and thus the number of electrons escaping in vacuum is likely to be overestimated. Nevertheless, it is interesting to notice that in the PIC simulations at very high intensities 11,12 “plumes” of electrons extending into the vacuum region are generated near the maxima of the standing ESW. A “plume” structure in the vacuum region is also evident in the velocity field shown in Fig.4.

Acknowledgments.

This work was partly supported by the Italian Ministry of University and Research (MIUR) through the project “Generation of fast electron and ion beams by superintense laser irradiation”. Discussions with Serena Bastiani-Ceccotti, Caterina Riconda and Francesco
Pegoraro are gratefully acknowledged.

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FIG. 1: Growth rate of the $2\omega \rightarrow \omega + \omega$ process vs. normalized electron density $n_e/n_c = \omega_p^2/\omega_L^2$ (x-axis) and temperature $\sqrt{k_B T_e/m_e c^2}$ (labels). The dashed line is the “cold” result previously obtained [8]. The growth rate is normalized to $a_L^2 \omega$ where $a_L = eE_L/m_e c\omega_L$ is the dimensionless field amplitude of the laser pulse.

FIG. 2: Top: $(y, p_x)$ phase space projections from PIC simulations [11] at two subsequent times, showing electron acceleration localized near spatial maxima of the growing surface oscillations. Bottom: $(y, p_x)$ phase space projection from test particle simulations, showing similar features. The “black stripe” around $v_x = -0.1$ represents electrons that have not been reached yet the surface region at the time shown.

FIG. 4: Contours of time-averaged fluid velocity from test particle simulations, showing a transverse modulation due to the enhanced electron acceleration near ESW maxima ($y \simeq \lambda_s/4$, $y \simeq 3\lambda_s/4$).

FIG. 3: Left: phase space $(p_x,x)$ projection (integrated over $y$) from test particle simulations, showing fast electron jets penetrating into the plasma region ($x > 0$) where they move ballistically ($p_x \propto x$). The jets are produced twice per laser cycle. All particles in the simulation are included in the plot. Right: same as left side, but restricted to particles around $y = \lambda_s/4$ (spatial maximum of the standing surface wave), showing enhanced acceleration around this position once per laser cycle. In both figures electrons propagating into vacuum ($x < 0$) are also evident. They are found only near ESW maxima ($y \simeq \lambda_s/4$, $y \simeq 3\lambda_s/4$).
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