PHYSICAL REVIEW A VOLUME 44, NUMBER 10 15 NOVEMBER 1991

Surface waves in the regime of the anomalous skin effect

R. Dragila*
Laser Physics Centre, Research School of Physical Sciences, The Australian National University, P.O. Box 4, Canberra, Australian Capital Territory 2601, Australia

E. G. Gamaliy
Lebedev's Physical Institute, Leninskiy prospect 53, Moscow, U.S.S.R.
(Received 6 August 1990)

We have formulated an effective description of surface waves propagating along the boundary between a high-temperature, high-density plasma and vacuum or a high-temperature, high-density plasma and a low-density plasma, within the framework of the anomalous skin effect. Since the anomalous skin effect is essentially collisionless, we have investigated possible collisionless dissipative mechanisms. Two mechanisms have been identified: phase breaking and capacitor heating. The possibility of exciting surface waves by an external electromagnetic wave is also discussed. We propose two types of solid planar targets that after being ionized by an intense ultrashort heating laser beam allow the excitation of surface waves by an obliquely incident *p*-polarized electromagnetic wave.

PACS number(s): 52.40.Db

I. INTRODUCTION

The development of laser facilities [1] generating high-intensity ($>10^{18}$ W/cm^2) ultrashort pulses (<1 psec) allows the creation of laboratory plasmas with solid-state density and very high electron temperatures ($\cong 10-100$ keV). It has been demonstrated [2] that in the presence of a strong electromagnetic field [$E > E_a$, $I > I_a$; where $E_a = e/r_B^2 = 5.1\times10^9$ V/cm (E_a is the atomic field and $I_a = cE_a^2/8\pi = 3.4\times10^{16}$ W/cm^2 is the corresponding laser intensity)] atoms of a solid target are ionized first due to multiphoton ionization in the tunnel limit, and second by electron impact within a time shorter than the period of the laser radiation. One of the important features of plasmas generated on the femtosecond time scale is that, unlike in "classical" laser-produced plasmas involving nanosecond time scales (see e.g., Ref. [3]) (and thus well-developed hydrodynamic motion), there is not enough time to convert the electron energy into the kinetic energy of directed motion of heavy ions and hence no hydrodynamic motion occurs during the pulse. This means that the irradiation by ultrashort laser pulses leads to the creation of a distinct boundary between the high-density warm plasma and the vacuum. Our interest in such a boundary is that it can support surface waves (see, e.g., Ref. [4]) propagating along it. In such a simple geometry these cannot be excited directly by the heating laser radiation, but can be, for example, driven by electron thermal fluctuations [5]. The conditions here are, however, different from those in "standard" surface-wave phenomena [4]—namely, high electron energies can lead to a situation where the electron mean free path substantially exceeds the penetration depth of the electromagnetic field (e.g., due to the laser radiation or due to electron thermal fluctuations) into the plasma. In such a case relationship between the current density and associated electric field is nonlocal and one deals with the so-called anomalous skin effect [6]. Nonlocality is the main feature of such a plasma that allows surface-wave propagation along the plasma boundary. The spatial distribution of the evanescent electromagnetic field within the plasma is not exponential, as in the absence of spatial dispersion, and it can be determined only when the electron distribution function is known.

It is beyond the scope of this paper to analyze the full self-consistent problem in which the distribution function is determined by the presence of the electromagnetic field [7,8]. Instead, we assume that the electron distribution function is known and isotropic. The last assumption permits us to simplify the calculations and is not important for the analysis of surface-wave excitation in this case. We assume also that the amplitude of the field associated with a surface wave is small, causing only a perturbation of the background distribution function. A dispersion relation for surface waves in the regime of the anomalous skin effect is presented and analyzed for some special cases. We also discuss the possibility of exciting a surface wave propagating along plasma surfaces created by an intense ultrashort heating laser beam by an obliquely incident *p*-polarized low-intensity probing beam. This may be of some diagnostic interest because information about the plasma temperature is encoded in the reflection properties for such probing beams.

II. THEORY

Let us consider a simple geometry where the plane $z=0$ separates the half-space $z>0$ occupied by a warm homogeneous plasma and vacuum ($z<0$). The plasma is assumed to be characterized by an isotropic distribution function which, however, is not necessarily Maxwellian. Since our aim is to find solutions which represent surface waves propagating along the boundary $z=0$, it is assumed that an electromagnetic field associated with a surface wave with frequency $\omega < \omega_{pe} = (4\pi e^2 N_e / m)^{1/2}$, where e, N_e, and m are, respectively, the electron charge, the density, and the mass, has the following form: $\mathbf{H}(z) = (H,0,0)$, $\mathbf{E}(z) = (0, E_y, E_z) \propto \exp[i(\omega t - ky)]$, where k is the component of the wave number in the direction of the wave propagation. The spatial distribution of the electromagnetic field is then described by a set of Maxwell's equations,

$$\nabla \times \mathbf{H} = \frac{4\pi}{c}\mathbf{j} + \frac{1}{c}\frac{\partial \mathbf{E}}{\partial t}\ , \quad \nabla \times \mathbf{E} = -\frac{1}{c}\frac{\partial \mathbf{H}}{\partial t}\ , \quad \nabla \cdot \mathbf{H} = 0\ , \tag{1}$$

where $\mathbf{j}$ is the current density and c is the speed of light in vacuum, or, equivalently, by the following set of coupled wave equations,

$$\frac{d^2 E_y}{dz^2} + k_0^2 E_y + ik\frac{dE_z}{dz} = i\frac{4\pi}{c}k_0 j_y\ , \tag{2}$$

$$(k_0^2 - k^2)E_z + ik\frac{dE_y}{dz} = i\frac{4\pi}{c}k_0 j_z\ . \tag{3}$$

Since we investigate the regime which corresponds to the conditions of the anomalous skin effect, the current density $\mathbf{j}$ in nonlocally related to the electric field $\mathbf{E}$ via the tensor of high-frequency conductivity σ, i.e.,

$$\mathbf{j}(z) = \int_0^\infty \sigma(z,z') \cdot \mathbf{E}(z')dz'\ . \tag{4}$$

To Fourier transform the wave equations (2) and (3), one has to continue the field quantities into the region $z < 0$. This procedure, however, depends on the model used to describe the reflection of electrons from the boundary $z = 0$. In what follows we will adopt the model of specular reflection [6]. Consistently, the field quantities are continued into the region $z < 0$ in accordance with the usual recipe [6],

$$\begin{aligned} &E_y(-z) = E_y(z), \quad H(-z) = -H(z), \\ &E_z(-z) = -E_z(z); \quad z > 0\ . \end{aligned} \tag{5}$$

While E_y is continuous across the boundary $z = 0$, dE_y/dz is an odd function of z and has a discontinuity there. The discontinuities of E_z and dE_y/dz at $z = 0$ can be, for the infinite-medium problem, related to the amplitude of the magnetic field at $z = 0$ as follows:

$$\left[\frac{dE_y}{dz} + ikE_z\right]_{z=0_+} - \left[\frac{dE_y}{dz} + ikE_z\right]_{z=0^-} = 2ikH_0\ , \tag{6}$$

where $H_0 = H(z = 0^+)$. Now, one can Fourier transform the wave equations (2) and (3) over the whole space $-\infty < z < \infty$ being careful of the discontinuities at $z = 0$. Such a transform then results into the following set of coupled equations:

$$(k_0^2 - \kappa^2)E_{y\kappa} - k\kappa E_{z\kappa} = i\frac{4\pi}{c}k_0 j_{y\kappa} - 2ik_0 H_0\ , \tag{7}$$

$$(k_0^2 - k^2)E_{z\kappa} - k\kappa E_{y\kappa} = i\frac{4\pi}{c}k_0 j_{z\kappa}\ , \tag{8}$$

where the Fourier transform ϕ_κ of a function $F(z)$ is defined in a usual manner,

$$\phi_\kappa = \int_{-\infty}^{\infty} F(z)e^{-i\kappa z}dz\ .$$

Here $j_{y,z\kappa} = \sigma_{y,z\kappa}E_{y,z\kappa}$ is the κ component of the Fourier transform of the current density and $\sigma_{y,z\kappa}$ is the Fourier transform of diagonal terms of the plasma high-frequency conductivity. The dependence of $\sigma_{y,z\kappa}$ on κ is the signature of the nonlocal relationship between current and electric field in the regime of the anomalous skin effect. The off-diagonal terms of the tensor σ vanish in accordance with our assumption that the plasma distribution function is isotropic. Solving (7) and (8) for $E_{y\kappa}$ and $E_{z\kappa}$ then results in

$$E_{y\kappa} = \frac{2ik_0 H_0[k_0^2 - k^2 - i(4\pi/c)k_0\sigma_{z\kappa}]}{[k_0^2 - k^2 - i(4\pi/c)k_0\sigma_{z\kappa}][k_0^2 - \kappa^2 - i(4\pi/c)k_0\sigma_{y\kappa}] - k^2\kappa^2} \equiv \xi(\kappa)H_0 \tag{9}$$

and

$$E_{z\kappa} = \frac{k\kappa}{k_0^2 - k^2 - i(4\pi/c)k_0\sigma_{z\kappa}}E_{y\kappa}\ . \tag{10}$$

The spatial distribution of the electromagnetic field within the plasma then becomes

$$E_y(z) = \frac{1}{2\pi}\int_{-\infty}^{\infty}\xi(\kappa)H_0 e^{i\kappa z}d\kappa\ , \quad E_z(z) = \frac{1}{2\pi}\int_{-\infty}^{\infty}\frac{k\kappa\xi(\kappa)}{k_0^2 - k^2 - i(4\pi/c)k_0\sigma_{z\kappa}}E_{y\kappa}e^{-i\kappa z}d\kappa\ . \tag{11}$$

The expression (9) now allows estimation of the penetration depth of the field into the plasma in two limits: (a) when a surface wave is predominantly of electromagnetic character [4], i.e., for $\omega^2 \ll \omega_{pe}^2$, and (b) when $\omega^2 < \omega_{pe}^2$ but $\omega^2 \cong \omega_{pe}^2$. In the electromagnetic limit the refractive index associated with the wave $n = k/k_0$ is very close to 1 (i.e., $n - 1 \ll 1$) and $E_{y\kappa}$ peaks for

$$|(4\pi/c)k_0\,\mathrm{Im}\sigma_{y\kappa}| \approx |(4\pi/c)k_0\,\mathrm{Im}\sigma_{z\kappa}| \approx \kappa^2 \gg k_0^2\ .$$

The penetration depth (skin depth) then becomes

$d_s \approx 1/\kappa$. This depth has different values and is described by different expressions depending on the plasma parameters. Thus for a dense collisional plasma when $\omega \gg \nu_{ei}$ but still $\omega^2 \ll \omega_{pe}^2$, the plasma conductivity becomes [9] $\sigma \cong i\omega_{pe}^2/4\pi\omega$ [the real part of σ is small as $O(\nu_{ei}/\omega)$] and $d_s \cong c/\omega_{pe} \gg d_{ei} = \langle v^2\rangle^{1/2}/\nu_{ei}$, where d_{ei} is the electron mean free path, ν_{ei} the electron-ion collision frequency, and $\langle v^2\rangle^{1/2}$ the average velocity of electron random motion. Such a penetration depth corresponds to the classical case [10] of surface waves propagating along the boundary of an overdense plasma in a regime where the plasma can be described macroscopically in terms of a dielectric permittivity $\epsilon = 1-\omega_{pe}^2/\omega^2$, i.e., when the current density j is a local function of the electric field. In the opposite limit of a strongly collisional plasma when $\omega \ll \nu_{ei}$, the plasma conductivity becomes [9]

$$\sigma \cong i\omega_{pe}^2/4\pi\omega + \omega_{pe}^2\nu_{ei}/4\pi\omega^2 \cong \omega_{pe}^2\nu_{ei}/4\pi\omega^2 ,$$

which results in the penetration depth

$$d_s \cong (c/\omega_{pe})(\omega/\nu_{ei})^{1/2} < c/\omega_{pe}$$

and $d_s \gg d_{ei}$. Finally, in the case of a high-temperature plasma when $d_{ei} > d_s$, which occurs when

$$\left[\frac{\frac{1}{2}m\langle v^2\rangle}{1\ \text{keV}}\right]^2 \left[\frac{N_e}{10^{24}\ \text{cm}^{-3}}\right]^{1/2} > 16 ,$$

one enters the region of the anomalous skin effect and the relation between the current density and corresponding electric field becomes nonlocal. Then,

$$\sigma_{y\kappa} \cong i\ \text{Im}\sigma_\kappa + \text{Re}\sigma_{y\kappa} , \quad \sigma_{z\kappa} \cong i\ \text{Im}\sigma_\kappa + \text{Re}\sigma_{z\kappa} \tag{12}$$

where

$$\text{Im}\sigma_\kappa \cong \frac{1}{4\pi}\frac{\omega_{pe}^2}{\omega}\frac{k_0}{\kappa}\frac{c}{\langle v^2\rangle^{1/2}} . \tag{13}$$

Therefore, as long as $n^2 \ll (\omega_{pe}/\omega)^2$ one can estimate that $E_{y\kappa}$ peaks at

$$\kappa = \kappa_s \cong \frac{\omega_{pe}}{c}\left[\frac{\omega}{\omega_{pe}}\frac{c}{\langle v^2\rangle^{1/2}}\right]^{1/3} , \tag{14}$$

and, consequently, the penetration depth can be estimated as

$$d_s \cong \frac{c}{\omega_{pe}}\left[\frac{\omega_{pe}}{\omega}\frac{\langle v^2\rangle^{1/2}}{c}\right]^{1/3} . \tag{15}$$

Such an estimate of the penetration depth, however, preassumes that $\text{Re}\sigma \ll \text{Im}\sigma$, which holds when the energy gain per electron per transit through the skin layer is small when compared to the initial energy of the electron, $m\langle v^2\rangle/2$, as is also the number of resonant electrons. The first requirement can always be satisfied if the field amplitude H_0 is small enough. We will return to the discussion of these conditions later.

According to standard perturbation theory [9], the plasma high-frequency conductivity can be expressed in terms of the unperturbed electron distribution function f as

$$\sigma_{y\kappa} = \frac{ie^2}{m}\int \frac{v_y^2}{\omega - kv_y - \kappa v_z}\frac{1}{v}\frac{\partial f}{\partial v}\,d\mathbf{v} , \tag{16}$$

$$\sigma_{z\kappa} = \frac{ie^2}{m}\int \frac{v_z^2}{\omega - kv_y - \kappa v_z}\frac{1}{v}\frac{\partial f}{\partial v}\,d\mathbf{v} , \tag{17}$$

where $\omega = \omega + i0_+$ and where, as usual, the integration is carried out along a Landau contour. For nonrelativistic plasmas and surface waves with predominantly electromagnetic character ($n \cong 1$), $\omega \gg kv_y$; however, due to the field localization in the anomalous skin layer, the kernels of the integrals (16) and (17) have poles at $\omega \cong \kappa v_z$. These poles determine the real parts of $\sigma_{y,z\kappa}$, i.e., the efficiency of dissipation of the waves within the plasma. One can thus identify two collisionless dissipation mechanisms: (i) "phase breaking" ($\propto \sigma_{y\kappa}$), when an electron with $v_z \geq \omega/\kappa$ undergoes acceleration (driven by E_y) parallel to the boundary and leaves the skin layer before E_y changes its sign (obviously, the efficiency of this mechanism is highest when $\omega \cong \kappa v_z$); and (ii) "capacitor heating," when an electron accelerated by E_z leaves the skin layer before E_z changes its sign. Both mechanisms can operate only when the electromagnetic field is highly localized in the space with the characteristic size $l_s \ll \lambda_0$, where λ_0 is the vacuum wavelength corresponding to the frequency ω.

In contrast, in the electrostatic limit when $n \gg 1$ and $\omega^2 < \omega_{pe}^2$ but $\omega^2 \cong \omega_{pe}^2$, the penetration depth becomes $d_{elst} \cong c/n\omega$, and in the limit $n \to \infty$ the penetration depth $d_{elst} \to 0$, which corresponds to the fact that there is no electrostatic field in a perfect conductor. However, surface waves with predominantly electrostatic character are not of interest to us because these waves are very slow ($\omega/k \ll c$) and, therefore, they (a) cannot be excited by an external electromagnetic field, and (b) can be heavily damped due to the large number of resonant particles (the damping rate can substantially exceed the usual exponentially small Landau damping rate).

The theory which has been developed above provides an understanding of the penetration of a p-polarized electromagnetic field into an overdense warm plasma and of the dissipation mechanisms involved. In what follows an analysis will be given of surface waves propagating along a plasma-vacuum boundary or along a density discontinuity within a stepwise homogeneous plasma in the regime of the anomalous skin effect.

Let us consider the same plasma-vacuum boundary $z=0$ as above. We now look for a solution where a p-polarized electromagnetic wave propagates along such a boundary and is evanescent both left and right from the boundary. Consequently, in vacuum the electromagnetic wave is characterized by

$$E_y = Ae^{\kappa_1 z}e^{i(\omega t - ky)} , \quad E_z = Be^{\kappa_1 z}e^{i(\omega t - ky)} ; \quad z<0 \tag{18}$$

while its z dependence in the plasma ($z>0$) is characterized by (11). Furthermore, since in vacuum $\text{div}\mathbf{E}=0$, the constants A and B are interrelated as $B = ikA/\kappa_1$, where $\kappa_1 = k_0(n^2-1)^{1/2} > 0$. Requiring continuity of E_y and $ikE_z + \partial E_y/\partial z$ at $z=0$, one obtains the following dispersion relation for surface waves:

$$\frac{(\kappa_1^2-k^2)\xi_y}{\kappa_1(\xi_y'-ik\xi_z)}=-1 , \tag{19}$$

where

$$\xi_y=\int_{-\infty}^{\infty}\xi(\kappa)\,d\kappa \tag{20}$$

is the so-called surface impedance [6], and

$$\xi_z=\int_{-\infty}^{\infty}\frac{k\kappa\xi(\kappa)}{k_0^2-k^2-i(4\pi/c)k_0\sigma_{z\kappa}}\,d\kappa$$

and

$$\xi_y'=i\int_{-\infty}^{\infty}\kappa\xi(\kappa)\,d\kappa .$$

The quantities ξ_y, ξ_z, and ξ_y' have their analogs in the classical (local and macroscopic) theory of surface waves, where the plasma is described in terms of the dielectric permittivity $\epsilon_2=1-\omega_{\rm pe}^2/\omega^2$, and, in particular

$$\xi_y\leftrightarrow\frac{i\kappa_2}{k_0\epsilon_2} ,\quad \xi_z\leftrightarrow\frac{k}{k_0\epsilon_2} ,\quad \xi_y'\leftrightarrow\frac{i\kappa_2^2}{k_0\epsilon_2} , \tag{21}$$

where $\kappa_2=k_0(n^2-\epsilon_2)^{1/2}$. Inserting these analogs into the dispersion relationship (19), one obtains the well-known classical dispersion relation for surface waves propagating along a plasma-vacuum $[\epsilon(z<0)=\epsilon_1=1]$ boundary [10],

$$\frac{\epsilon_1\kappa_2}{\kappa_1\epsilon_2}=-1 . \tag{22}$$

Generalization to the case of surface waves propagating along a plasma-plasma boundary is straightforward. To do so we will assume that the plasma occupying the region $z>0$ is characterized by the same set of parameters as in the previous case while the plasma occupying the region $z<0$ is rare, being characterized by the electron plasma frequency $\omega_{\rm pe1}\ll\omega_{\rm pe2}$. Since the linear classical theory of surface waves [4] requires that the dielectric permittivity must change sign across the boundary, then the necessary condition for a surface-wave solution to exist is that $\omega\geq\omega_{\rm pe1}$ [see the dispersion relation (22)], the condition for the anomalous skin effect to occur,

$$\frac{\omega_{\rm pe1}}{\omega}\frac{\langle v^2\rangle^{1/2}}{c}>1 ,$$

is not satisfied in the underdense plasma ($z<0$) and, consequently, it can be described classically via $\epsilon_1=1-\omega_{\rm pe1}^2/\omega^2$. The dispersion relation (19) remains formally the same where, however, now $\kappa_1=k_0(n^2-\epsilon_1)^{1/2}$. Obviously, the surface-wave-type solutions exist only for $n^2>\epsilon_1$.

The dispersion relation (19) is in general very complex and represents an implicit relation between ω and k. The details of this dependence are affected by the particular form of the electron distribution function, f, corresponding to the overdense plasma. Nevertheless, a qualitative analysis of (19) is possible in the electromagnetic limit $\omega_{\rm pe1}<\omega\ll\omega_{\rm pe}$ when n^2 exceeds, but is close to, ϵ_1. In such a case one can approximate the surface impedance ξ_y as $\xi_y\approx ik_0d_s$, where d_s is the penetration depth and, consequently, $\xi_y'\approx ik_0$ and $\xi_z\approx(kd_s)(k_0d_s)$. Then, neglecting dissipation in accordance with our assumption that ${\rm Re}\sigma\ll{\rm Im}\sigma$, the dispersion relation (19) becomes

$$\frac{k_0\epsilon_1k_0d_s}{\kappa_1[(kd_s)^2-1]}=-1 . \tag{23}$$

Since $1>\epsilon_1>0$, the dispersion relation (23) can have a solution corresponding to a surface wave only if $kd_s\approx k_0d_s<1$, which, in given conditions, is always well satisfied for both the anomalous skin effect [see (15)] and the classical case. The refractive index n can now be expressed in the following explicit form:

$$n^2\cong\epsilon_1\left[1+\left[\frac{\omega}{\omega_{\rm pe}}\right]^{4/3}\left[\frac{\langle v^2\rangle}{c}\right]^{2/3}\epsilon_1\right] , \tag{24}$$

while in the classical case,

$$n^2\cong\epsilon_1\left|1+\epsilon_1\left[\frac{\omega}{\omega_{\rm pe}}\right]^2\right| .$$

Therefore, in the electromagnetic limit the dispersion curve (24) lies only slightly below the classical one, which is not surprising when one realizes that in both cases the electromagnetic field associated with surface waves is spatially distributed mainly in the region $z<0$, characterized by $\epsilon=\epsilon_1$.

Since all previous considerations fall within the framework of the anomalous skin effect when $d_s<d_{ei}$, the energy of the surface waves is dissipated in the collisionless regimes of "phase breaking" and "capacitor heating." We will now estimate the efficiency of these mechanisms.

In the case of phase breaking an electron is accelerated by E_y along the boundary during the period $t_s=d_s/v_z<1/\omega$ until it leaves the skin layer (here $v_z\approx\langle v^2\rangle^{1/2}$). Therefore, the energy gain ΔE_e per electron per transit through the skin layer can be estimated as

$$\Delta E_e\approx eE_yv_{E_y}t_s ,$$

where $v_{Ey}=eE_y/m\omega$ is the electron quiver velocity. Substituting for d_s and normalizing ΔE_e to the electron average energy $\langle E_e\rangle=m\langle v^2\rangle/2$, one finally obtains

$$\frac{\Delta E_e}{\langle E_e\rangle}\approx\left[\frac{v_{E_y}}{c}\right]^2\left[\frac{mc^2}{\langle E_e\rangle}\right]^{4/3}\left[\frac{\omega}{\omega_{\rm pe}}\right]^{2/3} . \tag{25}$$

For example, for the frequency of a Nd glass laser, $\omega\cong2\times10^{15}$ rad s^{-1}, the solid-state electron density $N_e=10^{24}$ cm^{-3}, and the average electron energy $\langle E_e\rangle=10$ keV, one obtains $\Delta E_e/\langle E_e\rangle\approx(v_{E_y}/c)^2$. Consequently, as long as the quiver velocity of electrons within the skin layer is relatively low, $(v_{E_y}/c)^2\ll1$, the relative energy gain per electron per transit through the skin layer, $\Delta E_e/\langle E_e\rangle$, also remains small. Once $\Delta E_e/\langle E_e\rangle$ is known, one can estimate ${\rm Re}\sigma_y$ from the following simple energy-balance relationship:

$$\beta N_e\frac{d\langle E_e\rangle}{dt}\approx\beta N_e\frac{\Delta E_e}{t_s}\approx{\rm Re}\sigma_yE_y^2 ,$$

which then results in

$$\mathrm{Re}\sigma_y \approx \beta \frac{\omega_{\rm pe}^2}{4\pi\omega} , \quad (26)$$

where the magnitude of the numerical factor $\beta \ll 1$ depends on details of the distribution function f and βN_e is the effective density of electrons involved in the dissipation process.

In the case of capacitor heating an electron with average velocity $v_z > \omega d_s$ is accelerated in the direction perpendicular to the boundary over a distance d_s and gains the energy $\Delta E_e \approx eE_z d_s$. Repeating the energy-balance consideration of the previous case, one again obtains

$$\mathrm{Re}\sigma_z \approx \beta \frac{\omega_{\rm pe}^2}{4\pi\omega} . \quad (27)$$

However, since $(E_z/E_y)_{z=0^+} = \xi_z/\xi_y \approx k\delta$, the energy of surface waves with predominantly electromagnetic character ($n \approx \epsilon_1$, and therefore $k\delta \ll 1$) is dissipated mainly due to the phase breaking.

Finally, if now, by analyzing (16) and (17), one factorizes the coefficient β into two parts, one due to the finite size of the skin depth d_s and the other, α, purely due to the details of the shape of the electron distribution function,

$$\beta = \alpha \frac{\omega d_s}{\langle v^2 \rangle^{1/2}} ,$$

one can relate the real and the imaginary parts of the high-frequency conductivity as

$$\frac{\mathrm{Re}\sigma}{\mathrm{Im}\sigma} = \alpha \frac{\omega}{\omega_{\rm pe}} \frac{c}{\langle v^2 \rangle^{1/2}} , \quad (28)$$

where typically $\alpha < 1$, but not necessarily $\alpha \ll 1$, as is the case for Landau damping of waves with phase velocity substantially exceeding $\langle v^2 \rangle$; for example, for a Maxwellian electron distribution,

$$\alpha = \exp \left[-\frac{\omega^2 d_s^2}{\langle v^2 \rangle} \right] .$$

At the same time, in the regime of the anomalous skin effect,

$$\frac{\omega}{\omega_{\rm pe}} \frac{c}{\langle v^2 \rangle^{1/2}} < 1$$

and, therefore, our earlier assumption $\mathrm{Re}\sigma \ll \mathrm{Im}\sigma$ is justified. If one considers a numerical example of the plasma with $m\langle v^2 \rangle/2 = 10$ keV and $N_e = 10^{24}$ cm^{-3}, one obtains

$$\frac{\omega}{\omega_{\rm pe}} \frac{c}{\langle v^2 \rangle^{1/2}} \cong 0.1$$

for $\omega = 2\times 10^{15}$ rad s^{-1}, which corresponds to the frequency of the Nd glass laser, and

$$\frac{\omega}{\omega_{\rm pe}} \frac{c}{\langle v^2 \rangle^{1/2}} \cong 0.01$$

for a CO_2-laser frequency $\omega = 2\times 10^{14}$ rad s^{-1}. Also, under these conditions the electron mean free path is $d_{ei} \cong 3\times 10^{-5}$ cm and substantially exceeds the penetration depth of electromagnetic field into the plasma, $d_s \cong 2\times 10^{-6}$ cm (for $\omega = 2\times 10^{15}$ s^{-1}) and $d_s \cong (4-5)\times 10^{-6}$ cm (for $\omega = 2\times 10^{14}$ s^{-1}), consistent with the model of the anomalous skin effect. Finally, under the conditions when $\mathrm{Re}\sigma \ll \mathrm{Im}\sigma$, one can introduce the effective collision frequency representing the essentially collisionless dissipative processes as

$$\frac{\nu_{\rm eff}}{\omega} = \frac{\mathrm{Re}\sigma}{\mathrm{Im}\sigma} = \alpha \frac{\omega}{\omega_{\rm pe}} \frac{c}{\langle v^2 \rangle^{1/2}} . \quad (29)$$

Since the energy of surface waves investigated above is dissipated, these waves can exist only if they are excited by some energy source, e.g., an incident electromagnetic wave. However, as already mentioned, surface waves propagating along a plane plasma-vacuum boundary propagate with a phase velocity less than the speed of light in vacuum, c, and therefore cannot be excited by an external electromagnetic wave. If, however, a p-polarized electromagnetic wave is incident obliquely at the angle of incidence θ onto a structure like vacuum–low-density plasma–high-density plasma, such that $\omega > \omega_{\rm pe1}$, where $\omega_{\rm pe1}$ is the electron plasma frequency corresponding to the low-density plasma, it can excite a surface wave with phase velocity $\omega/k > c$ and wave number $k = k_0 \sin\theta = nk_0$. Experimentally, such a situation would correspond to the situation when a high-intensity ultrashort (heating) pulse impinges upon a solid high-density (high atomic number) target overcoated with a low-density (foam with low atomic number) film. The film and a part of the substrate are then "instantaneously" ionized and no mass motion develops during the heating pulse if it was short enough. A p-polarized low-entensity probing beam, which does not substantially affect the electron distribution function formed by the interaction of the heating beam with the target, is then launched obliquely at the angle of incidence θ onto the plasma structure to excite a surface wave. This can be diagnosed by measuring the amplitude coefficient of reflectivity R of the probing beam,

$$R = -\frac{\epsilon_1 \kappa_0 \kappa_1 \gamma + i\kappa_1^2 \gamma \sinh\kappa_1 d + i\kappa_1 \epsilon_1 k_0^2 \xi_y \cosh\kappa_1 d}{\epsilon_1 \kappa_0 \kappa_1 \gamma - i\kappa_1^2 \gamma \sinh\kappa_1 d - i\kappa_1 \epsilon_1 k_0^2 \xi_y \cosh\kappa_1 d} , \quad (30)$$

where $\kappa_0 = (k_0^2 - k^2)^{1/2}$ and $\gamma = ik\xi_z - \xi_y'$. It is a well-established fact in the theory of linear surface waves that multilayered plasma structures that allow external excitation of surface waves can act as total absorbers of the pump wave [11]. The coefficient of reflectivity R as a function of the angle of incidence θ has a characteristic dip $R = 0$ for some angle θ_0 (which is larger than the angle $\theta_{\rm TIR}$ for total internal reflection at the plasma-vacuum boundary: $\sin^2\theta_{\rm TIR} = \epsilon_1$). Formally, the coefficient of reflectivity vanishes when the numerator of (30) becomes zero. This, however, in fact represents two conditions. First, that the real part of the numerator becomes zero, which happens when the dispersion relation for surface waves [which in the special case when $d \rightarrow \infty$ reduces to

(19)] is satisfied. Second, that the imaginary part vanishes, which occurs when the rate of the energy tunneling through the evanescent layer of the plasma with lower density [$\propto \exp(-2\kappa_1 d)$] is exactly compensated by the rate of energy dissipation within the skin layer ($\propto \text{Im}\sigma$). Since vanishing reflectivity is a resonant effect occurring when the pump wave is totally linearly converted into a damped surface wave, the width of the resonant curve $R = R(\theta)$ is proportional to $\text{Im}\sigma$, but is also sensitive to d, because for $R \neq 0$ there is another kind of dissipation due to re-emission of the energy of surface wave back into the vacuum. Therefore, the width of the resonant curve $R = R(\theta)$ is a function of the average electron energy $m\langle v^2\rangle/2$, which may possibly form a basis for alternative diagnostics of the high-density plasma "temperature."

Finally, we propose an alternative method for exciting the surface wave by a probing beam incident on a high-density, high-temperature plasma created by an ultrashort intense laser beam. This time, instead of coating the solid target by a low-density film, one can corrugate the surface of the target to form a grating [4]. An ultrashort heating pulse ionizes the target to some depth and thus creates a static (during the laser pulse) plasma grating on the surface. Bragg reflection from such a periodic structure now provides matching of the phase velocities of the probe beam and the surface wave.

III. CONCLUSIONS

An effective description has been formulated of surface waves propagating along a high-temperature, high-density plasma-vacuum boundary or a high-temperature, high-density plasma-low-density plasma boundary within the framework of the anomalous skin effect. Such plasma density discontinuities can be created in short-pulse (<1 psec high-intensity laser-produced plasma experiments. Assuming that the plasma electron distribution function is known and is isotropic, we have found the dispersion relation for these waves. The dispersion properties of surface waves are only weakly affected by the nonlocality associated with the anomalous skin effect. This is easy to understand because most of the electromagnetic field of a surface wave is spatially distributed outside the high-density plasma. Furthermore, since the anomalous skin effect is essentially collisionless, we have investigated possible collisionless dissipative mechanisms. Two mechanisms have been identified: phase breaking and capacitor heating. It has been found that under the conditions considered here it is the first mechanism that dominates. This is because the component of the electric field associated with a surface wave (which in the linear regime is always p polarized) that is parallel to the boundary, E_y, substantially exceeds the amplitude of the component perpendicular to the boundary, E_z. Possibilities for exciting surface waves by an external electromagnetic wave have also been discussed. We have proposed two types of solid planar targets, which, after being ionized by an intense ultrashort heating laser beam, allow the excitation of surface waves by an obliquely incident p-polarized electromagnetic wave. The amplitude of this wave was assumed to be small (we called it a probing beam) relative to that of the heating wave since we are restricted by the use of linear theory.

The next obvious step is to consider the nonlinear self-consistent problem in which the surface wave is excited by the heating beam itself. In that case the electron distribution function is affected by the presence of the surface wave, which, in turn affects the dispersion and attenuation properties of the surface wave. If the heating wave is now p polarized and obliquely incident onto a target at an angle of incidence such that it provides total absorption of the heating wave at the instant when its amplitude peaks, the net portion of the pulse energy absorbed by the plasma can possibly be substantially increased. Since, to our knowledge, all existing theoretical models of the interaction of ultrashort pulses with plasmas are very pessimistic and predict the absorption rate at a level of $\cong 10\%$, the proposed method of surface-wave-assisted absorption has, in our opinion, some potential.

ACKNOWLEDGMENTS

The authors wish to acknowledge useful discussions with B. Luther-Davies and the assistance of Yanjie Wang.

*Deceased.

[1] P. Maine, D. Strickland, P. Bado, M. Pessot, and G. Morou, IEEE J. Quantum Electron. **QE - 24**, 398 (1988); M. D. Perry, F. G. Patterson, and J. Weston, Opt. Lett. **15**, 381 (1990); M. Ferray, L. A. Lompré, O. Gobert, A. L'Huillier, G. Mainfray, C. Manus, A. Sanchez, and A. S. Gomes, Opt. Commun. **75**, 278 (1990).

[2] E. G. Gamaliy and V. T. Tikhonchuk, Pis'ma Zh. Eksp. Teor. Fiz. **48**, 413 (1988) [JETP Lett. **48**, 453 (1989)].

[3] N. G. Basov, Yu. A. Zakharenkov, N. N. Zorev, G. V. Slizkov, A. A. Rupasov, and A. S. Shikanov, *Heating and Compression of Thermonuclear Targets by Laser Beam* (Cambridge University Press, Cambridge, United Kingdom, 1986).

[4] *Electromagnetic Surface Waves,* edited by A. D. Boardman (Wiley, New York, 1982).

[5] S. Vukovic, R. Dragila, and N. Aleksic, Phys. Rev. A **32**, 606 (1985).

[6] G. E. H. Reuter and E. H. Sondheimer, Proc. R. Soc. London, Ser. A **195**, 336 (1949); E. M. Lifshitz and L. P. Pitaevskiy, *Physical Kinetics* (Pergamon, Oxford, 1981).

[7] E. G. Gamaliy and R. Dragila, Phys. Rev. A **42**, 929 (1990).

[8] E. G. Gamaliy, A. E. Kiselev, and V. T. Tikhonchuk (unpublished); P. N. Lebedev Report No. N22, 1990 (unpublished).

[9] N. A. Krall and A. W. Trivelpiece, *Principles of Plasma Physics* (McGraw-Hill, New York, 1973).

[10] Yu. A. Romanov, Izv. Vyssh. Uchebn. Zaved. Radiofiz. **7**, 242 (1964).

[11] Yu. M. Aliev, S. Vukovic, O. M. Gradov, A. Yu. Kyrie, and V. M. Cadez, Pis'ma Zh. Eksp. Teor. Fiz. **25**, 351 (1977) [JETP Lett. **25**, 326 (1977)].