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Skin effect and interaction of short laser pulses with dense plasmas

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In this paper we discuss interaction of intense, subpicosecond laser pulses with plasmas. We propose a self-consistent analytical model of the anomalous and normal skin effects in plasmas with steplike density profile. The heat transport is described by classical Spitzer conductivity with new boundary conditions accounting for laser absorption in the thin skin layer. We obtain self-similar solutions for the heat-conduction problem, and the scaling laws for important plasma parameters are also discussed. Our predictions are consistent with the recent experimental results.

I. INTRODUCTION

In recent years the new technology of high-intensity subpicosecond lasers has stimulated exciting developments in physics of laser-matter interactions. These studies have been motivated by new applications for x-ray sources of high brightness,[1,2] x-ray lasers,[3,4] and also by interesting unexplored physics related to behavior of dense plasmas in super-strong electromagnetic fields.[5–7]

Plasmas created by intense subpicosecond laser pulses by irradiation of metal targets have several important differences from conventional laser plasmas studied, for example, in the context of laser fusion. Absorption and electron heating occur within a skin depth, and there is almost no hydrodynamical expansion during the time of interaction. Consequently, the electron density is two to three orders of magnitudes higher and corresponds to solid density. The typical plasma scale lengths involved are much smaller than a laser wavelength. We will analyze processes of absorption and heat conduction in such plasmas, over a wide range of parameters, trying to establish a set of scaling laws and self-similar solutions. Our approach is based on well-known classical physics. Heat conductivity, for example, is given by a Spitzer-type expression and corresponds to the upper bound of the possible heat flux. The interesting question of thermal-transport inhibition, anisotropy in the heated electron distribution function, will be discussed in future publications and compared with the results of a simpler theory derived here.

We will consider an ideal experimental situation, where the short laser pulse interacts with a steplike density profile. Under these conditions, the interaction takes place in the thin skin layer of an overdense plasma. The physics of the classical, stationary skin effect has been studied in great detail for metals and low-intensity electromagnetic fields.[8,9] However, the subpicosecond laser-pulse plasma interaction leads to a time-dependent situation because of the rapid absorption and plasma heating. The theory of anomalous and normal skin effect will be discussed in such plasmas, complementing previous works[10–13] on absorption and thermal conduction.

In recent experimental studies[5–7] on the physics of subpicosecond laser plasmas, intensities on target never exceed 10^{16} W/cm^2 and were achieved by focusing laser pulses of energies below 1 J. In all of these studies, the measurements and theoretical calculations have shown that the hydrodynamical evolution of plasmas did not play a major role during the main pulse duration, but that the density scale length depended on the prepulse energy. The interaction in the thin skin layer took place when the energy of the prepulse was small, and no plasma was produced on the surface of a target before the main pulse arrived. The temperatures achieved in these experiments were of the order of a few hundred electron volts, and the main absorption mechanism was related to electron-ion collisions in the skin layer. This is the physical situation related to normal, collisional skin effect.[8,9]

For higher laser intensities, and therefore higher plasma temperatures, in the range of a few kilo-electron-volts, the physics of the interaction can be quite different. It has been shown in Ref. 11 that the mean-free path of electrons could exceed the skin-layer depth for high-temperature plasmas, leading to the anomalous skin effect and collisionless energy absorption. The characteristic laser parameters correspond to intensities above 10^{17} W/cm^2 and pulse durations of 100 fs. For these fluxes of electromagnetic radiation, the electric field of the laser wave is of the order of an atomic electric field and produces the instantaneous direct ionization of atoms.

Temperatures which can be achieved during rapid plasma heating and absorption of electromagnetic radiation depend on the heat-transport processes. For picosecond and longer laser pulses, the main physical process regulating plasma temperature is an adiabatic cooling related to plasma expansion into vacuum.[3] In the case of shorter laser pulses, the dominant role is played by the thermal transport by electrons into dense cold matter. The role of radiative cooling is negligible in the energy balance because of the small plasma layer thickness.[3,11]

In this paper we will discuss a quasilinear approach to the anomalous and normal skin effects. The self-similar regimes of the laser heating of overdense plasma will be analyzed and compared with experimental results. Our discussion is based on the self-consistent model of energy absorption by the skin effect and classical, collisional heat transport into cold plasma. In Sec. II the theoretical model of an anomalous skin effect will be discussed. Analysis of absorption and heat conduction for the anomalous skin effect will be given in Sec. III. The self-consistent model of laser plasma interaction for the normal skin effect will be described in Sec. IV. Section V gives scaling laws, comparisons with experiments, and a summary.

II. THEORETICAL MODEL

The complete description of laser energy absorption, heating, thermal conduction, and plasma expansion is a complicated problem. In our simple theoretical model, we neglect all hydrodynamical effects related to ion motion, and we also assume instantaneous plasma ionization. The plasma is modeled as a slab with a sharp smooth boundary. The density of electrons is assumed to be constant. The electrons inside the plasma undergo specular reflection from the boundary. We only consider the normal incidence of a laser radiation.

First, we propose a kinetic theory for the anomalous skin effect, based on a quasilinear approximation for the slowly varying part of the electron distribution function. This approach explores the separation of scales between the thickness of the anomalous skin layer and a much longer electron mean free path, which defines the characteristic gradient of the averaged electron kinetic energy. Therefore, absorption in the skin layer will be described as a special kind of boundary condition for the electron distribution function.[14] Later, in Sec. IV, a similar concept will also be introduced for the normal skin effect in the situation where the skin-layer depth is much shorter than characteristic scale of the heat wave.

The kinetic equation will be solved in the region of plasma, where there is no electromagnetic field. The self-consistency of the model is related to the fact that absorption, which is described by the boundary condition, depends on the solution of the kinetic equation outside the absorption region. This is the important difference between anomalous and normal, collisional skin effect.[8,9] In the latter case, the skin depth can be comparable to the characteristic scale of temperature inhomogeneity (see Sec. V for more discussion on the applicability conditions and dominant physical processes for given plasma parameters).

The electron distribution function $f(\mathbf{r},\mathbf{v},t)$ is represented in the skin layer by two parts: a slowly varying part on the scale of the laser wave period $F(\mathbf{r},\mathbf{v},t)$ and a quickly varying part $\delta\mathscr{f}(\mathbf{r},\mathbf{v},t)$ that is proportional to the laser electric field. In Sec. II A, we will find the distribution of electric field and the distribution function f inside the skin layer following the usual approach of the linear theory of the anomalous skin effect.[8,9] Later, the quasilinear equation for F will be derived by our substituting $\delta\mathscr{f}$ into the nonlinear term of kinetic equation and averaging with respect to laser period. The quasilinear equation can be integrated over the skin depth in order to obtain the boundary condition for the distribution function outside the absorption region.

A. Linear theory of the anomalous skin effect

Let us consider a plasma occupying half-infinite space $(z>0)$, with ion density $n_i=n_0/Z$, where $Z\gg1$. The laser electromagnetic wave, with frequency ω_0 (where the critical density is $m\omega_0^2/4\pi e^2\equiv n_c\ll n_0$) propagates in vacuum $(z<0)$ and interacts at a normal angle of incidence with sharp plasma boundary at $z=0$. We assume that an electric field E_0 of the electromagnetic wave penetrates plasmas over the skin-layer depth l_s, which is much smaller than the mean-free path for electron-ion collisions $l_{e\text{-}i}$.

The evolution of electron distribution function $f=F+\delta\mathscr{f}$ is described by the following kinetic equation:

$$\frac{\partial f}{\partial t}+v_z\frac{\partial f}{\partial z}+\frac{e}{m}\left[\mathscr{E}+\frac{1}{c}(\mathbf{v}\times\mathscr{B})\right]\cdot\frac{\partial f}{\partial \mathbf{v}}=C[f]\ , \qquad (1)$$

where the collision operator C is taken in the Landau form. In Eq. (1), e and m stand for electron charge and mass, respectively, and the electromagnetic fields are given by the following expressions:

$$\mathscr{E}(z,t)=\mathscr{R}\mathbf{E}(z,t)\exp(-i\omega_0 t)+\mathbf{E}_a(z,t)\ , \qquad (2)$$

$$\mathscr{B}(z,t)=\mathscr{R}\mathbf{B}(z,t)\exp(-i\omega_0 t)\ , \qquad (3)$$

where the amplitudes $\mathbf{E},\mathbf{B}$ of electromagnetic radiation satisfy the Maxwell equations:

$$c\nabla\times\mathbf{E}=i\omega_0\mathbf{B}\ , \qquad (4)$$

$$c\nabla\times\mathbf{B}=4\pi\mathbf{j}\ . \qquad (5)$$

The displacement current has been neglected in Eq. (5) due to high density of the plasma $n_0\gg n_c$. In the geometry of normal incidence, vectors $\mathbf{E},\mathbf{B}$ are in the (x,y) plane and do not have any component along the z axis. The ambipolar electric field $\mathbf{E}_a$ acts along the z axis. This field is produced in order to maintain electric neutrality disturbed by strong electron heat flux into cold plasma.

Following the familiar theory[8,9] of the anomalous skin effect, we assume the subsequent form of the high-frequency electron distribution function: $\delta\mathscr{f}=\mathscr{R}\delta f\exp(-i\omega_0 t)$, where $\delta f\ll F$. Linearizing Eq. (1), we obtain

$$-i\omega_0\delta f+v_z\frac{\partial\delta f}{\partial z}=-\frac{e}{m}\mathbf{E}\cdot\frac{\partial F}{\partial\mathbf{v}}\ . \qquad (6)$$

We also assume that the plasma heating rate and the collision frequency $\nu_{e\text{-}i}$ are much smaller than ω_0. The Lorentz force term is also neglected in Eq. (6), as we assumed that the anisotropic part of F is small. One can solve Eq. (6) with the boundary conditions: $\delta f(v_z<0,v_x,v_y,z\to\infty,t)\to0$, and with the requirement of the specular reflection of electrons from the plasma boundary: $\delta f(v_z<0,z=0,t)=\delta f(-v_z,z=0,t)$.

One can write the solutions to Eq. (6), for electrons moving to the left and to the right, in the same form using a formal and even extension of the electric field into the region of $z<0$ (cf. Refs. 1 and 2). Thus, by formally taking $\mathbf{E}(z<0,t)=\mathbf{E}(-z,t)$, we have the following solution for the high-frequency part of the electron distribution function:

$$\delta f^{\pm}(|v_z|)=\pm\left[\frac{e}{m|v_z|}\right]\int_z^{\mp\infty}dz'\exp\left[i\omega_0\frac{|z-z'|}{|v_z|}\right]\mathbf{E}(z',t)\cdot\frac{\partial F(z=0)}{\partial \mathbf{v}}\ , \tag{7}$$

where $\delta f^{\pm}$ corresponds to particles with positive v_z (+) and negative v_z (−). The slowly varying part of the distribution function F is approximately constant over the skin-layer depth l_s; therefore, in Eq. (7), we evaluate F at $z=0$.

Using solutions (7) we can write an expression for the high-frequency current induced in the plane of incidence by the electromagnetic field:

$$\mathbf{j}=e\int d\mathbf{v}\,\mathbf{v}_p\delta f=e\int_{v_z>0}d\mathbf{v}\,\mathbf{v}_p(\delta f^{+}+\delta f^{-})\ , \tag{8}$$

where $\mathbf{v}_p$ stands for the velocity vector in (x,y) plane. Substituting Eq. (8) into the Maxwell equations (4) and (5), one obtains an equation for the electron-field amplitude inside the skin layer:

$$\frac{\partial^2\mathbf{E}}{\partial z^2}=\frac{4\pi i e^2\omega_0}{mc^2}\int_0^{\infty}\frac{dv_z}{v_z}\int d\mathbf{v}_pF(\mathbf{v}_p,v_z,z=0)\int_{-\infty}^{+\infty}dz'\,\mathbf{E}(z')\exp\left[i\omega_0\frac{|z-z'|}{v_z}\right]\ . \tag{9}$$

In the usual treatment of the anomalous skin effect,[8,9] Eq. (9) is solved with the assumption that F is known and has the form of the Maxwellian or Fermi distribution function. In our case, we do not know F, and later we will study the equation for the low-frequency distribution function; however, we can at least assume that the anisotropic part of F is small, and use in Eq. (9) only the isotropic part F_0.

In order to solve Eq. (9), we assume the following form of the electric field:

$$\mathbf{E}(z,t)=\frac{\omega_0 l_s}{c}\mathbf{B}(z=0,t)\times\frac{\mathbf{z}}{|\mathbf{z}|}u\left[\frac{z}{l_s(t)}\right]\ ,$$

where $\mathbf{B}(0,t)$ is a value of the magnetic field on the boundary, and $u(\xi)$ is a dimensionless function satisfying the following equation derived from Eq. (9):

$$\frac{d^2u}{d\xi^2}=-\frac{i}{\pi}\int_{-\infty}^{+\infty}d\xi'\,u(\xi')Q(\xi-\xi')\ , \tag{10}$$

where the kernel Q is given by

$$Q(\xi)=\left[\int_0^{\infty}dv\,vF_0(v)\right]^{-1}\times\int_0^{\infty}dv\,vF_0(v)\int_0^1 d\mu\frac{1}{\mu}\exp\left[i\frac{|\xi|}{\mu}\frac{\omega_0 l_s}{v}\right]\ , \tag{11}$$

where $\mu=v_z/v$, and the characteristic depth l_s of the skin layer is given by

$$l_s(t)=(2\pi)^{-1}\left[\left[\frac{e^2\omega_0}{mc^2}\right]\int_0^{\infty}dv\,vF_0(v,z=0,t)\right]^{-1/3}\ . \tag{12}$$

In the regime of the plasma parameters characterizing the anomalous skin effect, the ratio $\omega_0 l_s/v\ll 1$, where v is of the order of the average electron velocity $\langle v\rangle$. Therefore, in Eq. (11) the main contribution to the integral with respect to μ comes from the small values of μ. Physically, this means that only electrons moving almost parallel to the plasma surface can effectively absorb electromagnetic energy. Calculating the Fourier transform of Q (11), we obtain

$$Q(q)=\frac{1}{2\pi}\int_{-\infty}^{+\infty}d\xi\,Q(\xi)e^{-iq\xi}\approx\frac{1}{2|q|}\left[1+iO\left[\frac{\omega_0 l_s}{|q|\langle v\rangle}\right]\right]\ , \tag{13}$$

where the small correction, $\sim\omega_0 l_s/\langle v\rangle$, describes the reactive part of the electric conductivity and does not have considerable influence on the structure of the electric field inside the skin layer. The solution of Eq. (10), with Q given by the leading term of Eq. (13), can be obtained using the Fourier transform method. One should notice, however, that due to our formal extension of the electric field for $z<0$ [cf. Eq. (7)], functions $u(\xi)$ and $E(z)$ have discontinuous derivatives at the boundary. In order to find the values of the derivatives at $z=0$, we use the Maxwell equation (4) and express $\partial\mathbf{E}/\partial z$ in terms of $\mathbf{B}(0,t)$; we find that $du/d\xi|_{\xi=0}=i$. The Fourier transform of the function $u(\xi)$ reads as

$$u(q)=(i/\pi)(i|q|^{-1}-q^2)^{-1}\ . \tag{14}$$

By taking the inverse Fourier transform of Eq. (14), we obtain

$$u(\xi)=\frac{1}{\pi}\int_0^{\infty}dq\,q(1+iq^3)^{-1}\cos(q\xi)\ . \tag{15}$$

Following Ref. 8 we can evaluate integral (15), and we find that

$$u(0)=\frac{1}{3\sqrt{3}}-\frac{i}{3}\ . \tag{16}$$

From the conditions for the parallel components of electric and magnetic fields on the plasma boundary, we have

$$|\mathbf{B}(0,t)|=2E_0\ ,\qquad \mathbf{E}(0,t)=2\frac{\omega_0 l_s(t)}{c}\mathbf{E}_0u(0)\ , \tag{17}$$

where $\mathbf{E}_0$ stands for the electric field of the incoming radiation in vacuum. Note that since $\omega_0 l_s\ll c$, the amplitude of the electric field inside the plasma $\mathbf{E}(0,t)$ is much smaller than the vacuum field $\mathbf{E}_0$.

Knowing the electric field in the plasma [(14) and (15)] and the high-frequency part of the distribution function

(7), one can calculate the absorption efficiency of the laser radiation inside the plasma, which yields

$$A=\left[c\frac{\langle \mathbf{E}_0^2\rangle}{8\pi}\right]^{-1}\frac{1}{2}\mathcal{R}\int_0^\infty dz\,\mathbf{E}^*(z)\cdot\mathbf{j}(z)$$
$$=|E_0|^{-2}\mathcal{R}|\mathbf{E}(0,t)\times\mathbf{B}^*(0,t)|$$
$$=\frac{4}{3\sqrt{3}}\frac{\omega_0 l_s(t)}{c}\ . \tag{18}$$

Thus the solution for the electrodynamical part of the problem depends on the symmetric part of the electron distribution function on the plasma boundary. In Sec. II B, we use the quasilinear approach in order to find F_0 and l_s.

B. Kinetic equation for an electron distribution function in a quasilinear approximation

Let us begin by noting that the high-frequency part of the electron distribution function δf (7) is much smaller than F for the laser electric field satisfying the following relation:

$$v_E=\frac{e|E(0)|}{m\omega_0}\ll\langle v\rangle\ , \tag{19}$$

where v_E is an amplitude of an electron oscillatory velocity in the laser field. Our theory will be valid in the regime of parameters satisfying condition (19).

In order to derive an equation for the slowly varying part of the distribution function, we average Eq. (1) over the period of laser oscillations:

$$\frac{\partial F}{\partial t}+v_z\frac{\partial F}{\partial z}+\frac{e}{m}E_a\frac{\partial F}{\partial v_z}+\frac{e}{2m}\mathcal{R}\mathbf{E}^*\cdot\frac{\partial \delta f}{\partial \mathbf{v}_p}=C[F]\ . \tag{20}$$

The last term on the left-hand side of Eq. (20) describes electron heating inside the skin layer $z\le l_s$. From relation (19) it is evident that the energy gain by electrons in the skin layer is much smaller than the average kinetic energy. As the variation of F in the region of skin layer is small and takes place over a very short distance, we will describe the whole heating process as a boundary condition for the evolution of F inside a cold plasma, i.e., for $z\ge l_s$. Neglecting the time derivative, the ambipolar field term, and the collision operator in Eq. (20) for $z\le l_s$, we can write the following:

$$v_z\frac{\partial F}{\partial z}=-\frac{e}{2m}\mathcal{R}\mathbf{E}^*\cdot\frac{\partial\delta f}{\partial \mathbf{v}_p}\ . \tag{21}$$

We integrate Eq. (21) over the skin-layer depth, and using the condition of specular reflection of electrons from the plasma boundary, we will represent the boundary condition for F in terms of the discontinuity at $z=0$:

$$(F)_{z=0}\equiv F(v_z>0,\mathbf{v}_p,z\gg l_s,t)-F(-v_z,\mathbf{v}_p,z\gg l_s,t)$$
$$=-\frac{e}{2mv_z}\int_0^\infty dz\,\mathcal{R}\mathbf{E}^*(z)\cdot\frac{\partial}{\partial \mathbf{v}_p}(\delta f^++\delta f^-)\ .$$

For simplicity, let us assume that the laser light is unpolarized or circularly polarized, and that we can assume symmetric conditions in the (x,y) plane. Taking these assumptions into account, we can rewrite the boundary condition in the following form:

$$(F)_{z=0}=\frac{v_{E_0}^2}{c^2}\left[\frac{\omega_0 l_s}{v}\right]^4 V\left[\frac{\omega_0 l_s}{\mu v}\right]\left[\frac{v}{\mu^2}\frac{\partial}{\partial v}v-v^3\frac{\partial}{\partial v}\frac{1}{v}\right]$$
$$\times\frac{\partial F_0(v,z=0,t)}{\partial v}\ , \tag{22}$$

where $\mu=v_z/v=\cos\theta>0$, $v_{E_0}=e\langle E_0^2\rangle^{1/2}/m\omega_0$ (the averaged electron velocity in the oscillatory electric field in vacuum), and the function V describes the angular distribution of heated electrons:

$$V(w)=2\mathcal{R}\int_0^\infty d\xi\, u(\xi)\int_{-\infty}^\infty d\xi' u(\xi')\exp(iw|\xi-\xi'|)\ .$$

We can rewrite the expression above for $V(w)$ using Eq. (15) for the spatial structure of the electric field and Fourier representation:

$$V(w)=\frac{4w}{\pi}\mathcal{R}\int_0^\infty dq\frac{iq^2}{(1+iq^3)^2}\frac{1}{w^2-q^2}\ .$$

The integrand above has a pole at $q=w$, which corresponds to the resonance interaction between fields and electrons. We can deform the contour in the vicinity of the real axis, $w\to w+iO$, and close it along the line constructed by rotating the positive real axis by $-\pi/6$ and along an arc for large arguments (cf. Ref. 8). There is no singularity inside the contour, and we can write $V(w)$ in the following form:

$$V(w)=\frac{4w}{\pi}\int_0^\infty d\kappa\frac{\kappa^2}{(1+\kappa^3)^2}\frac{w^2-\kappa^2/2}{w^4+\kappa^4-w^2\kappa^2}\ . \tag{23}$$

The expression (22) will be applied as the boundary condition for the kinetic equation (20) describing an evolution of the slowly varying part of the distribution function. Note that F depends now only on the particle speed v, the angle θ between the velocity vector $\mathbf{v}$ and the axis z, the coordinate z, and the time t. Also, since we have assumed that $Z\gg1$, we can neglect electron-electron collisions in Eq. (20), which now can be written in the following form:

$$\frac{\partial F}{\partial t}+v\cos\theta\frac{\partial F}{\partial z}+\frac{e}{m}E_a\left[\cos\theta\frac{\partial F}{\partial v}-\frac{\sin\theta}{v}\frac{\partial F}{\partial\theta}\right]$$
$$=\nu_{e\text{-}i}\frac{1}{\sin\theta}\frac{\partial}{\partial\theta}\left[\sin\theta\frac{\partial F}{\partial\theta}\right]\ , \tag{24}$$

where $\nu_{e\text{-}i}=2\pi Ze^4\Lambda/m^2v^3\equiv\nu_0/v^3$ stands for the electron-ion collision frequency, and Λ is the Coulomb logarithm. The kinetic equation (24) describes the evolution of the distribution function under the influence of collisions; however, as a result of rare electron-electron collisions, only the effect of isotropization of the velocity distribution function is taken into account in Eq. (24). The ambipolar electric field $\mathbf{E}_a$ arises in Eq. (24) in order to maintain quasineutrality in the presence of an electron thermal flux into cold plasma. This field is defined by the condition that the electric current component along the z axis vanishes, i.e., that

$$\int_0^\infty dv\, v^3 \int_{-1}^{1} d\mu\, \mu F(v,\mu,z=0,t)=0 \ . \tag{25}$$

The system of equations (24) and (25), together with boundary condition (22) and the necessity of particle density conservation,

$$\int_0^\infty dv\, v^2 \int_{-1}^{1} d\mu F(v,\mu,z,t)=\text{const} \ , \tag{26}$$

constitute the main equations in our theory.

It is important to note, that the boundary condition (22) describes the transfer of electromagnetic energy to electrons primarily in the direction perpendicular to z, and results in the energy flux AI_0, where A represents the absorption coefficient (18), and $I_0=c\langle E_0^2\rangle/8\pi$ represents the laser energy flux. The boundary condition (22) does not induce any electric current in plasma; therefore, the ambipolar field in the vicinity of the plasma boundary is zero. The nonzero E_a is related to the isotropization of the electron distribution function due to electron-ion collisions.

III. ABSORPTION AND HEAT CONDUCTION FOR THE ANOMALOUS SKIN EFFECT

A. Equation for the symmetric part of a distribution function

As an example, we first consider the situation where the temperature scale length is much longer than the mean-free path for electron-ion collisions. This somewhat ideal case is nevertheless a good starting point for the discussion of the heat-conduction problem. Our solution will establish the upper bound on the heat-conduction rates.

Following the standard procedure for the solution of the kinetic equations, we expand the distribution function F into a series of eigenfunctions of the collision operator (24). In our case, these are Legendre polynomials $P_n(\mu)$. Substituting

$$F(v,\mu,z,t)=\sum_{n=0}^{\infty} F_n(v,z,t)P_n(\mu) \tag{27}$$

into the kinetic equation (24) and keeping only the first two Legendre polynomials, we obtain

$$\frac{\partial F_0}{\partial t}+\frac{v}{3}\frac{\partial F_1}{\partial z}+\frac{eE_a}{3mv^2}\frac{\partial(v^2F_1)}{\partial v}=0 \ , \tag{28}$$

$$F_1=-\frac{v}{2\nu_{e\text{-}i}}\frac{\partial F_0}{\partial z}-\frac{eE_a}{2m\nu_{e\text{-}i}}\frac{\partial F_0}{\partial v} \ . \tag{29}$$

In the equation (29) we have dropped $\partial F_1/\partial t$ as compared with $\nu_{e\text{-}i}F_1$, and we have also neglected terms with F_2. Both of these assumptions are the standard elements of classical-transport theory and are related to the dominant role played by collisions and to the Eq. (19). This theory is only valid for the characteristic distances longer than the mean free path. For the anomalous skin effect, the transition range within a collisional mean free path of the skin layer is especially poorly approximated by taking only zero- and first-order Legendre polynomials in Eq. (27). Beyond this transition region, however, the local temperature is a well-defined quantity, and our solution to the kinetic equation (24), based on only two Legendre polynomials, will give reasonable approximation to the thermal-transport problem.

Using Eqs. (29) and (25), we can find the explicit form for the ambipolar electric field:

$$E_a(z,t)=\frac{m}{6e}\left[\int_0^\infty dv\, v^5F_0\right]^{-1}\frac{\partial}{\partial z}\int_0^\infty dv\, v^7F_0 \ . \tag{30}$$

In order to obtain the boundary conditions for F_1, we take the first μ moment of Eq. (22):

$$F_1(v,z=0,t)=\frac{v_{E_0}^2}{c^2}\left[\frac{\omega_0 l_s}{v}\right]^4\left[s_1\left[\frac{\omega_0 l_s}{v}\right]v\frac{\partial}{\partial v}\left[v\frac{\partial F_0}{\partial v}\right]-s_2\left[\frac{\omega_0 l_s}{v}\right]v^3\frac{\partial}{\partial v}\left[\frac{1}{v}\frac{\partial F_0}{\partial v}\right]\right] , \tag{31}$$

where

$$s_1(u)=\tfrac{3}{2}\int_0^1 d\mu\,\mu^{-1}V\left[\frac{u}{\mu}\right] ,$$
$$s_2(u)=\tfrac{3}{2}\int_0^1 d\mu\,\mu V\left[\frac{u}{\mu}\right] . \tag{32}$$

Equations (28)–(30) and the boundary condition (31) constitute the complete description of the evolution of the symmetric part of the distribution function. We can further simplify the boundary condition (31) by taking the asymptotic values of the functions $s_1(u)$ and $s_2(u)$ for $u=\omega_0 l_s/v \ll 1$, which is a proper region of parameters for the anomalous skin effect. In this limit, $s_2\approx -\frac{4}{3}\sqrt{3}u \ll 1$ [and the corresponding term in Eq. (31) may be neglected], the function s_1 is of order unity:

$$s_1(u\ll 1)\approx s_0=\frac{\pi}{2\sqrt{3}} \ . \tag{33}$$

As a result, we obtain the following kinetic equation for the isotropic part of the slowly varying distribution function:

$$\frac{\partial F_0}{\partial t}=\frac{v^2}{6\nu_{e\text{-}i}}\frac{\partial^2 F_0}{\partial z^2}+\frac{ev}{6m\nu_{e\text{-}i}}\frac{\partial E_a}{\partial z}\frac{\partial F_0}{\partial v}+\frac{eE_a}{3m\nu_{e\text{-}i}v^2}\frac{\partial^2}{\partial z\partial v}(v^3F_0)+\frac{e^2E_a^2}{6m^2\nu_{e\text{-}i}v^5}\frac{\partial}{\partial v}\left[v^5\frac{\partial F_0}{\partial v}\right] . \tag{34}$$

This equation should be complemented by the relation (30) for the ambipolar electric field and by the boundary condition (31), which can be now written in the following form

$$\left.\frac{\partial F_0}{\partial z}\right|_{z=0}=-\frac{eE_a(0,t)}{mv}\frac{\partial F_0(v,0,t)}{\partial v}-2\nu_{e\text{-}i}s_0\frac{v_{E_0}^2}{c^2}\left[\frac{\omega_0 l_s}{v}\right]^4\frac{\partial}{\partial v}\left[v\frac{\partial F_0(v,0,t)}{\partial v}\right] . \tag{35}$$

In Sec. III B we will derive from Eq. (34) a simple theory of heat conduction, based on the assumption about the existence of the local Maxwellian distribution function.

B. Solution to the thermal-transport problem

Let us propose the electron distribution function in the form of a local Maxwellian, with the space- and time-dependent temperature:

$$F_0(v,z,t)=\frac{n_0 m^{3/2}}{[2\pi T(z,t)]^{3/2}}\exp\left[-\frac{mv^2}{2T(z,t)}\right] . \qquad (36)$$

One should stress that the form (36) of the distribution function is not a solution to the kinetic equation (34), and is only used here in order to obtain the approximate behavior of the averaged electron kinetic energy as a function of space and time. Our simple approach will elucidate the important role of the new boundary condition (35) characterizing the anomalous skin effect.

Taking the v^4 moment of the kinetic equation (34), with F_0 given by (36), and using relation (30), we obtain the following expression for the ambipolar electric field:

$$E_a=\frac{5}{2e}\frac{\partial T}{\partial z} , \qquad (37)$$

and the equation for thermal conductivity,

$$\frac{\partial T}{\partial t}=-\frac{2}{3n_0}\frac{\partial q}{\partial z} , \qquad (38)$$

where the thermal flux q is given by

$$q=-32\left[\frac{2}{\pi}\right]^{1/2}\frac{n_0}{\nu_0 m^{5/2}}T^{5/2}\frac{\partial T}{\partial z} . \qquad (39)$$

Taking the moment of Eq. (35) with v^9, we obtain the boundary condition for the thermal heat flux:

$$\begin{aligned} q(z=0,t) &= \frac{2}{3}\left[\frac{2}{\pi}\right]^{7/6} s_0 n_0 m c^3\left[\frac{v_{E_0}}{c}\right]^2\left[\frac{\omega_0}{\omega_p}\right]^{8/3}\left[\frac{T}{mc^2}\right]^{1/6} \\ &\equiv A I_0 , \end{aligned} \qquad (40)$$

where A is the absorption coefficient of Eq. (18), and $\omega_p=(4\pi e^2 n_0/m)^{1/2}$.

The problem of thermal conduction defined in Eqs. (38)–(40) is very similar to the usual formulation of the heat-conduction problem,[15] with one important difference, however, in the boundary conditions (40), which is now given in terms of a relation between temperature and heat flux.

We now introduce the dimensionless units of

$$\begin{aligned} & t=t_0\tau, \quad v=v_{T0}u , \\ & z=v_{T0}t_0\xi, \quad T=mv_{T0}^2\Theta , \end{aligned} \qquad (41)$$

to bring Eqs. (38)–(40) to normalized forms:

$$\begin{aligned} & \frac{\partial\Theta}{\partial\tau}=\frac{\partial}{\partial\xi}\Theta^{5/2}\frac{\partial\Theta}{\partial\xi} , \\ & \Theta^{5/2}\frac{\partial\Theta}{\partial\xi}\bigg|_{\xi=0}=-\Theta^{1/6} , \\ & \Theta(\xi\to\infty)=0 . \end{aligned} \qquad (42)$$

Normalizing Eqs. (38) and (39) requires that $v_{T0}^3 t_0^{-1}=\frac{3}{64}(\pi/2)^{1/2}\nu_0$. Then v_{T0} is determined from Eq. (40), and hence t_0, as

$$\begin{aligned} t_0 &= \frac{64}{3}\left[\frac{2}{\pi}\right]^{1/2}\frac{v_{T0}^3}{\nu_0} \\ &= 32\left[\frac{2}{3}\right]^{13/4}\left[\frac{2}{\pi}\right]^{29/16}\frac{c^3}{\nu_0}s_0^{9/8}\left[\frac{v_{E_0}}{c}\right]^{9/4}\left[\frac{\omega_0}{\omega_p}\right]^3 , \\ v_{T0} &= c\left[\frac{2}{3}\right]^{3/4}\left[\frac{2}{\pi}\right]^{7/16}s_0^{3/8}\left[\frac{v_{E_0}}{c}\right]^{3/4}\frac{\omega_0}{\omega_p} . \end{aligned} \qquad (43)$$

One can find a self-similar solution to Eqs. (42). By introducing the new variable $\eta=\xi/\tau^{4/5}$, and the self-similar function $\Theta(\xi,\tau)=\tau^{6/25}\Phi(\eta)$, we rewrite Eqs. (42) in the following form:

$$\begin{aligned} & \frac{5}{4}\frac{d}{d\eta}\Phi^{5/2}\frac{d\Phi}{d\eta}+\eta\frac{d\Phi}{d\eta}-\frac{3}{10}\Phi=0 , \\ & \Phi^{7/3}\frac{d\Phi}{d\eta}\bigg|_{\eta=0}=-1, \quad \Phi(\infty)=0 . \end{aligned} \qquad (44)$$

Equations (42) satisfy the physical condition[15] that the total energy contents of the heat wave, $W=\int_0^\infty dz\,\Theta(z)$, changes in time with the rate given by the energy flux at the left boundary:

$$\frac{dW}{d\tau}=\Theta^{1/6}(\xi=0,\tau) . \qquad (45)$$

Relation (45), in terms of the self-similar variables and Eq. (44), leads to the following expression:

$$\int_0^\infty d\eta\,\Phi(\eta)=\frac{25}{26}\Phi^{1/6}(0) . \qquad (46)$$

We solved Eq. (44) numerically for several boundary values of $\Phi(0)$, and chose the proper value satisfying condition (46). The shape of the heat wave is shown in Fig. 1. The obtained thermal wave front is similar to the front

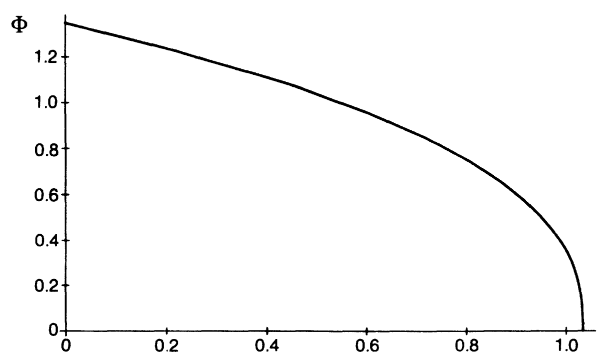

FIG. 1. Self-similar solution to the heat-transport equation (44) for the anomalous skin effect. Function Φ is proportional to dimensional temperature, and $\eta=\xi/\tau^{4/5}$ in the self-similar variable.

in the case of the constant heat flux;[15] however, because of the uncommon boundary condition (40), the time dependence of the temperature and the position of the heat front edge are different:

$$T(0,t)=mv_{T0}^2\Phi(0)\tau^{6/25} ,$$
$$z_f=v_{T0}t_0\eta_f\tau^{4/5} , \tag{47}$$

where $\Phi(0)=1.35$ and $\eta_f=1.07$ are found from the numerical solutions of Eq. (44).

Using the scaling coefficients (43) and the power laws obtained from the self-similar solution, one can obtain scaling laws for all of the important plasma parameters which characterize the efficiency of laser-plasma coupling:

$$A(t)=\frac{4}{3^{3/2}}\left[\frac{2}{\pi}\right]^{1/6}\left[\frac{\omega_0}{\omega_p}\right]^{2/3}\left[\frac{T(0,t)}{mc^2}\right]^{1/6} ,$$
$$l_s(t)=\left[\frac{2}{\pi}\right]^{1/6}\frac{c}{\omega_0^{1/3}\omega_p^{2/3}}\left[\frac{T(0,t)}{mc^2}\right]^{1/6} , \tag{48}$$
$$l_{e\text{-}i}=v_{T0}t_0\tau^{12/25} ,$$

where $l_{e\text{-}i}$ is a mean-free path for electron-ion collisions. Our results (47) and (48) display a weak dependence on the time of the absorption coefficient and the skin-layer thickness, $\sim t^{0.04}$. On the other hand, the electron mean-free path increases rapidly with time, $\sim t^{0.48}$, and so does the penetration depth of the heat front $\sim t^{0.8}$.

The weak time dependence of the absorption coefficient $A(t)$ results in an almost constant value of the heat flux on the boundary. Therefore, our power laws are close to the classical result[15] of the heat wave penetration with a constant heat flux on the boundary.

The following plasma conditions have to be satisfied for the solutions (47) and (48) to be valid:

$$\tau \gg 1 , \tag{49}$$
$$l_{e\text{-}i} > l_s , \tag{50}$$
$$\left[\frac{ZT_e}{m_i}\right]^{1/2} t \le l_s . \tag{51}$$

Equation (49) corresponds to the validity condition for the self-similar description of the classical-thermal conductivity, which should be valid for the characteristic temperature scale length being much longer than electron mean free path, i.e., $z_f \gg l_{e\text{-}i}$. The second relation above [Eq. (50)] describes the region of applicability of the anomalous skin effect.

Finally, Eq. (51) defines the time period of the interaction, when we can neglect the hydrodynamical effects. It is an important restriction for our theory, which assumes a sharp plasma boundary and a homogeneous plasma density. The ion sound speed in Eq. (51) approximates the expansion velocity in the isothermal rarefaction wave model. For longer times, violating Eq. (51), plasma corona will be formed, and part of the absorption process will take place in the critical region leading to the overall drop in absorption efficiency. This suggests that the maximum temperature in this model will be achieved for $t^* \sim l_s/(ZT_e/m_i)^{1/2}$. We will discuss this applicability condition in more detail in Sec. V.

IV. ABSORPTION AND HEAT CONDUCTION FOR THE NORMAL SKIN EFFECT

Our theory has been devoted thus far to the description of an anomalous skin effect. The derived system of equations is based on the conjecture that the depth of the skin layer is much smaller than the electron mean-free path. Therefore, the absorption in the skin layer is modeled as a special kind of boundary condition for the heat wave propagation into cold plasma. Note, however, that similar theoretical concepts can be applied to the description of the opposite physical situation of the normal skin effect. When $l_{e\text{-}i} < l_s$, and $\nu_{e\text{-}i} > \omega_0$, but the characteristic penetration depth of the heat wave is much deeper than the skin-layer depth, we can still assume that the absorption process takes place in an infinitesimally thin region on the surface of the plasma.

Let us start, as before, from the kinetic equation (1), with a collision operator in the Landau form. Equation (6), for the high-frequency part of the distribution function, will now take the following form:

$$\frac{e}{m}\mathbf{E}\cdot\frac{\partial F}{\partial \mathbf{v}}=C[\delta f] . \tag{52}$$

The solution to Eq. (52) reads as

$$\delta f=-\frac{e}{2m\nu_{e\text{-}i}(v)}\mathbf{E}\cdot\frac{\partial F_0}{\partial \mathbf{v}} , \tag{53}$$

where $\nu_{e\text{-}i}(v)=\nu_0/v^3$ is a collision frequency, and F_0 stands for an isotropic part of the distribution function. Taking F_0 to be a Maxwellian distribution function, we can calculate the electric current, $\mathbf{j}=\sigma\mathbf{E}$, where plasma dc conductivity is given by

$$\sigma=\frac{e^2n_0}{m\nu_{\text{eff}}} , \tag{54}$$

and the effective collision frequency is given in the terms of the electron thermal velocity, $v_T=(T_e/m)^{1/2}$, as

$$\nu_{\text{eff}}=\frac{1}{8}\left[\frac{\pi}{2}\right]^{1/2}\nu_0 v_T^{-3} . \tag{55}$$

By solving the Maxwell equations with the appropriate boundary conditions[8,9] we obtain an electric-field amplitude in the plasma:

$$E(z,t)=E_b\exp[(i-1)z/l_s(t)] ,$$
$$E_b=(1-i)\frac{\omega_0 l_s(t)}{c}E_0 , \tag{56}$$

where $l_s=(c/\omega_p)(2\nu_{\text{eff}}/\omega_0)^{1/2}$ is the skin depth related to the normal skin effect. Using the above expressions, we obtain for the absorption coefficient the following:

$$A=2\frac{\omega_0 l_s}{c} . \tag{57}$$

The temperature distribution and heating rates are given by the solutions to Eq. (38), for the thermal conduction and boundary condition, as

$$q(z=0,t)=q_{\rm abs}=A(t)I_0 \; . \tag{58}$$

Equation (58) simply expresses the energy balance, where the flux at $z=0$ corresponds to the absorbed energy flux. The difference between the normal skin effect [(38) and (58)] and the anomalous skin effect [(38) and (40)] follows from different boundary conditions, in particular, from a different dependence of the absorption coefficient on temperature.

We now introduce a set of dimensionless variables:

$$\begin{aligned} &t=\hat{t}_0\chi, \quad v=w\hat{v}_{T0} \; , \\ &z=\hat{v}_{T0}\hat{t}_0\zeta, \quad T=m\hat{v}\,^2_{T0}\Xi \; , \end{aligned} \tag{59}$$

so that

$$\frac{\partial \Xi}{\partial \chi}=-\frac{\partial \Pi}{\partial \zeta}, \quad \Pi=-\Xi^{5/2}\frac{\partial \Xi}{\partial \zeta}, \quad \Pi(\zeta=0,\chi)=\Xi^{-3/4} \; . \tag{60}$$

Again, as for Eqs. (38) and (39), we require that $v_{T0}^3/t_0=\frac{3}{64}(\pi/2)^{1/2}\nu_0$; however, the boundary flux is now given by Eq. (58) instead of Eq. (40), and therefore we have a different normalization for v_{T0}, and hence for t_0, namely,

$$\hat{v}_{T0}=\left[\frac{1}{3}\left[\frac{\pi}{2}\right]^{1/4}\right]^{2/9} c \left[\frac{v_{E_0}}{c}\right]^{4/9}\left[\frac{n_c}{n_0}\right]^{2/9}\left[\frac{\nu_0\omega_0}{\omega_p^2 c^3}\right]^{1/9} ,$$

$$\hat{t}_0=\frac{64}{3^{5/3}}\left[\frac{2}{\pi}\right]^{1/3}\left[\frac{n_c}{n_0}\right]^{2/3}\frac{\omega_0}{\omega_p^2}\left[\frac{\nu_0\omega_0}{\omega_p^2 c^3}\right]^{-2/3}\left[\frac{v_{E_0}}{c}\right]^{4/3} . \tag{61}$$

The dimensionless temperature Ξ and heat flux Π satisfy equations very similar to Eqs. (42) for the anomalous skin effect, but with different boundary conditions.

Equations (60) admit self-similar transformation, $\phi=\zeta\chi^{-17/24}$, $\Xi=\chi^{1/6}\Psi(\phi)$, leading to the following set of ordinary differential equations:

$$\begin{aligned} &\frac{24}{17}\frac{d}{d\phi}\Psi^{5/2}\frac{d\Psi}{d\phi}+\phi\frac{d\Psi}{d\phi}-\frac{4}{17}\Psi=0 \; , \\ &\Psi^{13/4}\frac{d\Psi}{d\phi}\bigg|_{\phi=0}=-1, \quad \Psi(\infty)=0 \; . \end{aligned} \tag{62}$$

The numerical solution to Eqs. (62) is shown in Fig. 2. It corresponds to $\Psi(0)=1.46$, and $\phi_f=1.27$, which were found following the same procedure as for the Fig. 1. The shape of the heat wave in Fig. 2 is very similar to the solution found for the anomalous skin effect, Fig. 1. The small differences follow from the different boundary conditions.

From the solution of Eqs. (62) and the self-similar transformation, we derive scaling laws for important physical parameters: The temperature for the constant laser intensity grows with time as $\sim t^{1/6}$, which is a

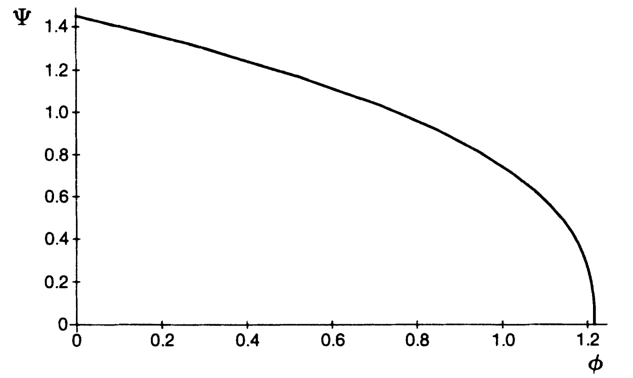

FIG. 2. Self-similar solution to the heat-transport equations (62) for the normal skin effect. Function Ψ is proportional to dimensional temperature, and $\phi=\zeta\chi^{-17/24}$ is the self-similar variable.

slower rate as compared with the anomalous skin effect; the absorption and the skin-layer thickness decrease with the same rate of $\sim t^{-1/8}$ (while, in the case of the anomalous skin effect, these values were slowly increasing with time); finally, the mean-free path for electron-ion collisions and the position of the heat front increase with time as $\sim t^{1/3}$ and $\sim t^{17/24}$, respectively. Our theoretical model of the normal-skin-effect heating is valid for

$$l_{e\text{-}i}<l_s \; , \tag{63}$$

$$z_f>l_s \; . \tag{64}$$

Condition (63) defines the separation between the anomalous and normal skin effect. Equation (64) lets us describe the absorption process as a boundary condition for the heat wave. However, in the cases of shorter laser pulses and/or smaller intensities, Eq. (64) may not be satisfied, and the heat front will not penetrate the cold material beyond the skin-layer depth. In such cases the heat-conduction losses are negligible, and we obtain from the energy balance equation,

$$\tfrac{3}{2}n_0\frac{dT}{dt}=\tfrac{1}{2}\sigma|E|^2\approx\omega_0\frac{E_0^2}{4\pi} \; ,$$

the linear growth in time of temperature in the absorption region:

$$T\approx\omega_0 t\frac{E_0^2}{6\pi n_0} \; . \tag{65}$$

This result was obtained from our simple model, which neglects the energy losses related to ionization, and also nonideal plasma effects that can increase specific heat and, therefore, decrease temperature. In Sec. V we will discuss the conditions of the applicability of various theoretical models.

V. SUMMARY AND CONCLUSIONS

We studied absorption and heat-conduction processes during the interaction of subpicosecond laser pulses with metal targets. In particular, we discussed theories of anomalous and collisional skin effects. In order to summarize our results, we will first describe the scaling laws of various physical parameters obtained from self-similar solutions to the heat-transport equations (44) and (62). We will use the following set of units: laser intensity, $I=I_0/10^{18}$ W/cm^2; laser wavelength, $\lambda=\lambda_0/1$ μm; plasma temperature will be given in kilo-electron-volts; distances will be expressed in nanometers; time duration will be given in femtoseconds.

For the anomalous skin effect, the characteristic velocity and time (43) have the following form:

$$v_{T0}=(1.5\times 10^9)I^{3/8}\lambda^{-1/4}Z^{-1/2}\ \text{cm/s}\ ,\qquad t_0=(4.9\times 10^2)Z^{-7/2}I^{9/8}\lambda^{-3/4}\ \text{fs}\ , \tag{66}$$

where we took for the ion density, $n_i=6\times 10^{22}$ cm^{-3}, and $n_0=Zn_i$. In estimates for the anomalous skin effect, the Coulomb logarithm is approximately $\Lambda=5$. Using Eqs. (66), we can write the scaling laws for the important physical parameters (47) and (48) as

$$\begin{aligned} T&=0.4Z^{-4/25}I^{12/25}\lambda^{-8/25}t^{6/25}\ ,\\ A&=(5.5\times 10^{-2})Z^{-9/25}I^{2/25}\lambda^{-18/25}t^{1/25}\ ,\\ l_s&=11.4Z^{-9/25}I^{2/25}\lambda^{7/25}t^{1/25}\ ,\\ z_f&=56.7Z^{-6/5}I^{3/5}\lambda^{-2/5}t^{4/5}\ . \end{aligned} \tag{67}$$

In Sec. IV we presented a theory of the normal skin effect, which was based on the idea of introducing special boundary conditions into the heat-transport equations, in order to describe absorption in a manner similar to the anomalous skin effect. From the self-similar transformation, we obtained for the characteristic velocity and time (61) the following:

$$v_{T0}=(1.3\times 10^9)I^{2/9}\lambda^{-1/9}Z^{-1/9}\ \text{cm/s}\ ,\qquad t_0=(5.3\times 10^2)I^{2/3}\lambda^{-1/3}Z^{-7/3}\ \text{fs}\ , \tag{68}$$

where the Coulomb logarithm was taken as $\Lambda=3$. From Eqs. (68) we can derive scaling laws for

$$\begin{aligned} T&=0.53Z^{1/6}I^{1/3}\lambda^{-1/6}t^{1/6}\ ,\\ A&=(2.2\times 10^{-2})Z^{3/8}I^{-1/4}\lambda^{-3/8}t^{-1/8}\ ,\\ l_s&=1.75Z^{3/8}I^{-1/4}\lambda^{5/8}t^{-1/8}\ ,\\ z_f&=10^2Z^{-19/24}I^{5/12}\lambda^{-5/24}t^{17/24}\ ,\\ l_{e\text{-}i}&=(3.2\times 10^2)Z^{-5/3}I^{2/3}\lambda^{-1/3}t^{1/3}\ , \end{aligned} \tag{69}$$

where the collision mean-free path is given by $l_{e\text{-}i}=v_T/\nu_{\text{eff}}=\frac{3}{8}v_{T0}t_0(t/t_0)^{1/3}$. Our results [(68) and (69)] for the normal skin effect are valid only when the penetration depth of the heating wave, given by the position of the heat front edge z_f, is larger than the skin depth l_s, i.e., when

$$I>(5\times 10^{-3})Z^{3/4}\lambda^{5/4}t^{-5/4}\ . \tag{70}$$

For low laser intensities, when Eq. (70) cannot be satisfied, the heat losses by conduction from the skin-layer region are negligible. For such bulk heating of plasma, within skin-layer depth, the temperature (65) changes linearly with time:

$$T(0,t)=2.6It\lambda^{-1}Z^{-1}\ . \tag{71}$$

For small temperatures, the plasma parameter becomes of order unity, and the nonideal effects become important. By requiring that the Coulomb logarithm $\Lambda=2$, we can write using (71) an additional restriction for the laser intensity:

$$I>(5\times 10^{-3})Z^{4/3}\lambda t^{-1}\ . \tag{72}$$

For intensities lower than predicted by Eq. (72), one has to deal with strongly coupled plasma physics[16] in order to properly describe collisional absorption and transport. Interesting measurements by Milchberg *et al.*[5] are partly obtained in this regime of parameters. Note, also, that the low-temperature nonideal effects will play an important role near the heat front for the wide range of intensities.

In addition, we have to make sure that the oscillatory velocity of the electrons is smaller than the thermal velocity (19). It is one of the required assumptions for the quasilinear approach to the solution of the kinetic equation for the anomalous skin effect. It results in the following relation for the laser flux:

$$I<0.18Z^{14/17}\lambda^{-22/17}t^{4/17}\ . \tag{73}$$

As we will see, condition (73) is easily satisfied for realistic plasma parameters. Finally, we neglected in our theory all effects related to hydrodynamical expansion (51). For the anomalous skin effect, this restriction, together with scaling laws (67), leads to

$$I<\mathcal{A}^{25/8}Z^{-39/8}\lambda^{11/4}(t/40)^{-27/4}\ . \tag{74}$$

Similarly, for the normal skin effect, Eqs. (51) and (69) imply that

$$I<\mathcal{A}^{6/5}Z^{-1/2}\lambda^{1.7}(t/7)^{-2.9}\ , \tag{75}$$

where $\mathcal{A}$ is the atomic mass number. Our results (67) can be applied to the situation of anomalous skin effect, defined by the Eq. (50), and to the opposite case of normal skin effect (69), which corresponds to condition (64). Using results (69) for the normal skin effect, we can define a separation line between these two regimes:

$$I=(3.2\times 10^{-3})Z^{49/22}\lambda^{23/22}t^{-1/2}\ . \tag{76}$$

In order to illustrate our results, we will first plot the curves defining different regimes of laser plasma interaction for the two representative cases of aluminum ($\mathcal{A}=27$, $Z=6-10$) and gold ($\mathcal{A}=200$, $Z=10$). In Fig. 3, for the case of heavy metal and for the laser wavelength $\lambda=0.25$ μm, we plot logarithms of intensity with respect to pulse duration: curve 1 separates the regions of the anomalous and normal skin effects (76); curve 2 distinguishes between the normal skin effect, with heat conductivity, and the case of bulk heating, without thermal-conduction losses (70); below line 3 [Eq. (72)], plasma becomes strongly coupled; dotted line 4 [Eq. (73)] bounds the region where oscillatory velocity of electrons is smaller than thermal velocity; the dashed curve 5 [Eq. (74)]

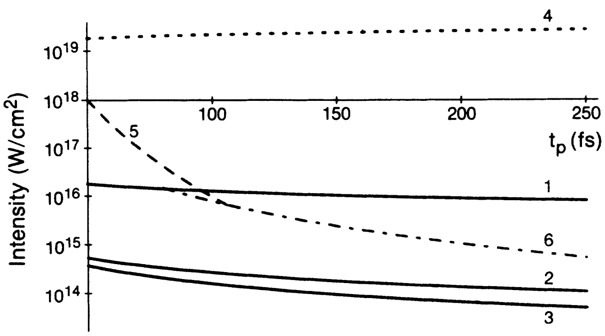

FIG. 3. Logarithm of laser-pulse intensity as a function of pulse duration for the gold target $\mathcal{A}=200$, $Z=10$, and laser wavelength $\lambda=0.25$ μm. Curve 1 separates the regions of anomalous and normal skin effects (76); curve 2 distinguishes between the normal skin effect with heat conductivity and the case of bulk heating without thermal conduction losses (70); below line 3 (72), the plasma becomes strongly coupled; dotted line 4 (73) bounds the region, where the oscillatory velocity of the electrons is smaller than the thermal velocity; dashed curve 5 (74) defines the lower boundary of the regime, where hydrodynamic expansion creates plasma corona and the anomalous skin effect does not occur; dot-dashed curve 6 (75) gives a similar restriction, related to hydrodynamical expansion, for the normal skin effect.

defines the lower boundary of the regime, where hydrodynamic expansion creates the plasma corona, and the anomalous skin effect does not occur; the dot-dashed curve 6 [Eq. (75)] gives a similar restriction, related to hydrodynamical expansion, for the normal skin effect. One can see from Fig. 3 that the anomalous skin effect can dominate the absorption for intensities $I>2\times10^{16}$ W/cm^2, and pulse duration $t_p<100$ fs. Normal skin effect can take place in a much larger region of parameters. The region of bulk heating, between curves 2 and 3, is narrow, which is a general result for all cases considered here.

Figure 4 presents the same information as Fig. 3 for the light-metal case. The anomalous skin effect cannot happen in this regime of parameters, and its occurrence is prevented by hydrodynamical expansion. The plasma expansion also limits the normal skin effect to intensities below $I<10^{15}$ W/cm^2.

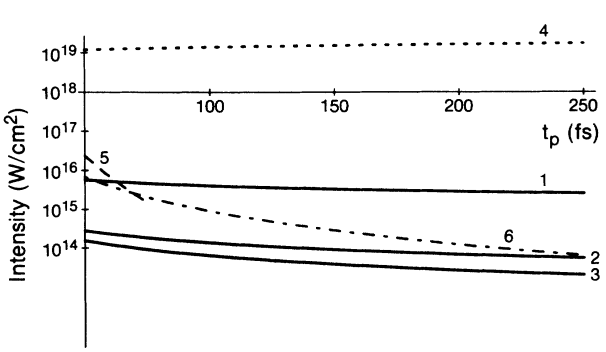

FIG. 4. Logarithm of laser-pulse intensity as a function of pulse duration for the aluminum target $\mathcal{A}=27$, $Z=6$. The curves on the plot have the same meaning as on Fig. 3.

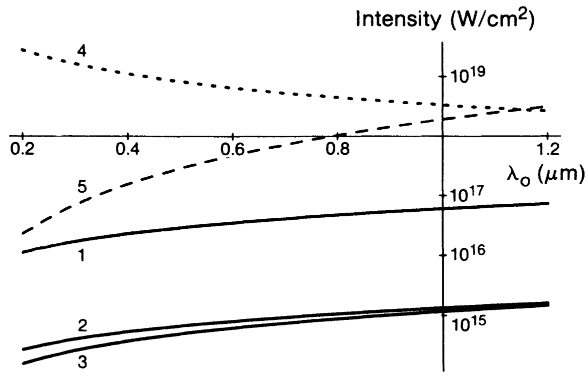

FIG. 5. Logarithm of laser-pulse intensities as a function of laser wavelength for the parameters of Fig. 3 and for the pulse duration $t_p=80$ fs.

Figures 5 and 6 illustrate various physical regimes, in terms of wavelength dependence, for the pulse duration $t_p=80$ fs. One can see from Fig. 5 that the region of applicability of the anomalous skin effect increases with the laser wavelength; also, for light targets (Fig. 6), the anomalous skin effect can occur for $\lambda>0.8$ μm.

Figure 7 illustrates the time dependence of the plasma temperature for the set parameters as in Fig. 3. The curves 1 and 2 correspond to the anomalous skin effect (67) for intensities $I=10^{18}$ and 2×10^{17} W/cm^2, respectively. Curve 3 is defined by the normal skin effect for $I=10^{16}$ W/cm^2. For the high-intensity case, the temperature can reach values above 1 keV, which, in reality can be even further increased because of thermal-inhibition.

We will compare our predictions with the results of some recent experimental studies.[1,5,7] Figure 8 shows our results for the aluminum target ($\mathcal{A}=27$, $Z=6$) and the laser wavelength $\lambda=0.31$ μm. All the curves have the same meaning as those in Figs. 3 and 4. For intensity $I=10^{15}$ W/cm^2, we obtain from Fig. 8 an expansion time of 110 fs in the regime of the normal skin effect. At this moment in time, Fig. 9 shows the plasma temperature $T_{\max}=190$ eV. This is the maximum temperature that

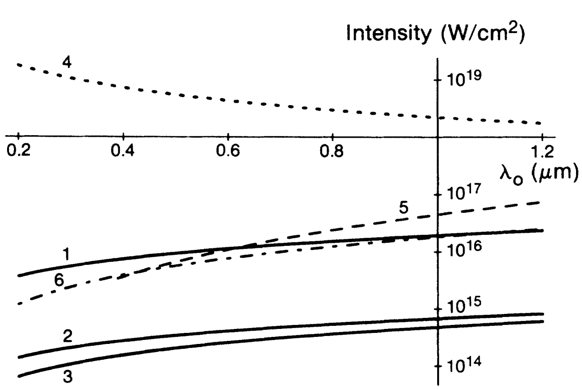

FIG. 6. Logarithm of laser-pulse intensities as a function of laser wavelength for the parameters of Fig. 4 and for the pulse duration $t_p=80$ fs.

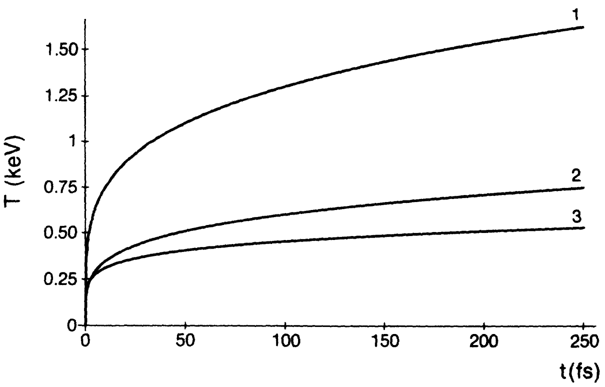

FIG. 7. Temperature as a function of pulse durations for the parameters of Fig. 1. Curves 1 and 2 correspond to the anomalous skin effect (67) for laser-pulse intensities $I=10^{18}$ and 2×10^{17} W/cm^2, respectively. Curve 3 is defined by the normal skin effect (69) and the laser intensity of 10^{16} W/cm^2.

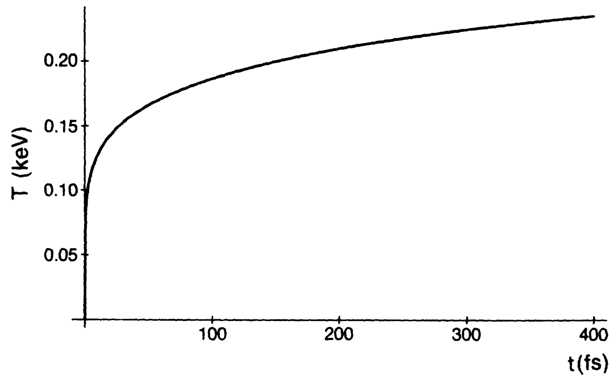

FIG. 9. Temperature as a function of pulse duration for the parameters of Fig. 8 and laser intensity $I=10^{15}$ W/cm^2.

can be achieved in the solid density region during the interaction, because, for longer times, plasma corona prevents the effective laser absorption in the dense plasma. For the same parameters, Milchberg *et al.*[5] reported a temperature of $T\approx105$ eV. The discrepancy of a factor of 2 could be attributed to the effect of hydrodynamical expansion, which occurs during the pulse duration of 400 fs. Only for intensities 10^{14} W/cm^2 does the expansion time become comparable to the laser-pulse duration. Therefore, the effective laser-plasma coupling takes place for intensities lower than 10^{14} W/cm^2, which can explain the change in the slopes of reflectivities and resistivity curves in Ref. 5.

For the parameters of the experiment by Fedosejevs *et al.*[7] ($\mathcal{A}=27$, $Z=6$ $\lambda=0.25$ μm, $I=3\times10^{15}$ W/cm^2), we obtain from (75) the expansion time of 75 fs and the maximum temperature (69) of 260 eV, in reasonable agreement with the reported temperature of 300 eV. Similarly, for the data of Murnane *et al.*,[1] the calculated temperature is in reasonable agreement with the observations.

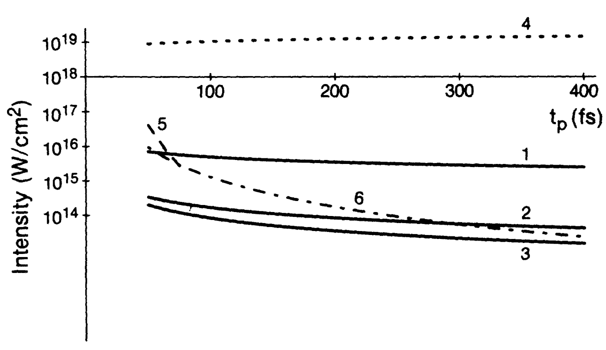

FIG. 8. Logarithm of laser-pulse intensities as a function of pulse duration for the case of aluminum target $\mathcal{A}=27$, $Z=6$, and laser wavelength $\lambda=0.31$ μm. All the curves are defined as in Fig. 3.

At last, we would like to comment on the predictions of our model for the absorption coefficient A. For both cases of anomalous (67) and normal (69) skin effects, absorption coefficients depend very weakly on pulse duration t and ionic charge Z. Figure 10 shows two absorption curves as functions of laser intensity given for the times of expansion, which are found by taking equalities in Eqs. (74) and (75), and for $\lambda=0.25$ μm. Our theory predicts an absorption of 35%–40% for the normal skin effect at 10^{14} W/cm^2, which drops with intensity until the anomalous skin effect takes over, given a 10% absorption at 10^{19} W/cm^2. In both cases, our results describe the absorption process in the solid density region.

In summary, we have presented the analytical model of the subpicosecond laser-pulse interaction with plasmas. Our analysis is based on the classical, Spitzer-type approach to the thermal transport and electron collisions. It is shown that the heat-conduction losses become important for laser intensities above 10^{15} W/cm^2, and the

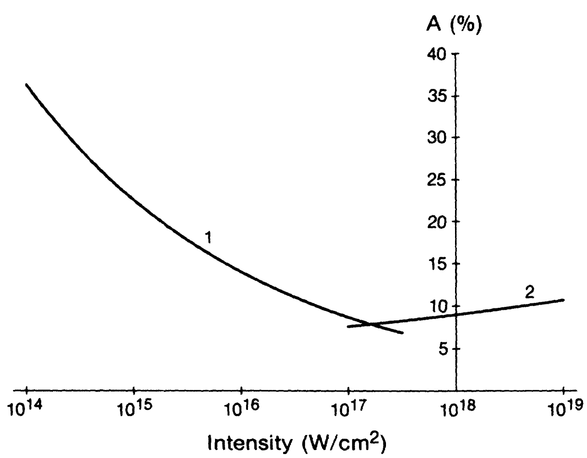

FIG. 10. Absorption as a function of laser-pulse intensity. Curve 1 corresponds to the normal skin effect (69), curve 2 is related to the anomalous skin effect (67). For both curves, times of interaction are evaluated from the conditions (74) and (75), and they correspond to the moment when hydrodynamical expansion becomes important. Additional parameters are $\mathcal{A}=27$, $Z=10$, and $\lambda=0.25$ μm.

regime of the anomalous skin effects can be achieved for heavy target materials and pulse durations $t_p < 100$ fs. Our results predict heat-flux values, which are comparable to the free-streaming limit. This may indicate the importance of thermal-flux inhibition for the proper description of interaction processes at high intensities, which may lead to higher plasma temperatures.

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