ISSN 1063-780X, Plasma Physics Reports, 2006, Vol. 32, No. 5, pp. 423–428. © Pleiades Publishing, Inc., 2006.
Original Russian Text © S.G. Bezhanov, S.A. Uryupin, 2006, published in Fizika Plazmy, 2006, Vol. 32, No. 5, pp. 459–464.

LASER PLASMA

Optical Properties of a Plasma Produced by the Tunneling Ionization of Atoms of a Matter in the Field of a Circularly Polarized Wave

S. G. Bezhanov and S. A. Uryupin

Lebedev Physical Institute, Russian Academy of Sciences, Leninskiĭ pr. 53, Moscow, 119991 Russia

Received August 4, 2005

Abstract—General features of the absorption and reflection of a test wave by a nonequilibrium plasma produced in the tunneling ionization of atoms of a matter by a circularly polarized laser pulse are described. Because of the highly anisotropic distribution of photoelectrons, the optical properties of a nonequilibrium plasma differ considerably from those of a plasma with a Maxwellian electron velocity distribution. Physically, an anomalous behavior of the absorption coefficient and of the phase shift stems from the fact that electron kinetics in the skin layer is modified by the alternating magnetic field of the test wave.

PACS numbers: 52.38.–r, 52.38.Dx

DOI: 10.1134/S1063780X06050084

1. INTRODUCTION

The plasma produced during rapid ionization of a matter is characterized by a strongly nonequilibrium velocity distribution of photoelectrons (see, e.g., [1–4]). The characteristic kinetic energy of the photoelectrons substantially exceeds the ion kinetic energy, and their distribution function is highly anisotropic. The regular features of the velocity distribution of photoelectrons depend largely on the mechanism whereby the atoms of a matter are ionized and on the properties of the ionizing radiation. In particular, the distribution of photoelectrons produced in the tunneling ionization of atoms by linearly polarized radiation is close to a bi-Maxwellian anisotropic distribution [1, 2]. The tunneling ionization of atoms by circularly polarized radiation results in a qualitatively different distribution of photoelectrons, one that is characterized by high electron velocities in the polarization plane of the ionizing wave [1, 4].

The physical properties of plasmas with a strongly nonequilibrium distribution of photoelectrons also differ significantly from those of a thermodynamically equilibrium plasma (see, e.g., [5–9]). The unusual properties of a plasma produced in the tunneling ionization by the field of a linearly polarized wave were described in [10–12]. In those papers, it was shown that the physical causes for the anomalous optical plasma properties lie in the influence of the magnetic field of the test wave on the electron kinetics in the skin layer. This influence is great because the electron distribution is anisotropic and the strength of an alternating magnetic field in the skin layer is relatively high in comparison with the electric field strength [10–12].

In the present paper, we consider the features of the absorption and reflection of a test wave that interacts with a plasma produced in the tunneling ionization of atoms of a matter by a circularly polarized wave. We assume that a nonequilibrium plasma has a sharp boundary, from which the electrons are specularly reflected, and that the electron kinetic energy is high enough to ignore rare electron collisions. Under these assumptions, we derive general relationships for the surface impedance of a nonequilibrium plasma, for the absorption coefficient, and for the phase shift of the test wave. We obtain explicit analytical expressions that describe how the absorption coefficient and phase shift depend on the ratio of the characteristic distance traveled by an electron during the period of the test wave to the skin depth in the case of a high-frequency skin effect. We check the analytical expressions against the results of numerical calculations of the phase shift and absorption coefficient and also discuss the conditions under which the optical plasma properties under study can be observed.

2. BACKGROUND STATE

We consider a half-space $z > 0$ that is filled with a plasma produced by the ionization of atoms of a matter by a circularly polarized high-power ultrashort laser pulse. The characteristic plasma production time and the duration of the ionizing pulse are assumed to be short in comparison with the characteristic time scale on which the nonequilibrium distribution of photoelectrons varies because of collisions and because of the possible onset of electromagnetic instabilities. We also ignore the change in the photoelectron density due to the plasma expansion into a vacuum ($z < 0$) during the pulse. If the plasma is produced by tunneling ionization, then the characteristic time scale on which the nonequilibrium photoelectron distribution forms is shorter than or on the order of the reciprocal of the frequency of the ionizing radiation. If, in this case, the ionizing radiation flux density I satisfies the inequalities

$$\left(\frac{4E_i}{3\hbar\omega_L}\right)^2 cnE_i \gg I \gg \frac{\omega_0^2}{\omega_L^2} cnE_i, \quad (1)$$

where E_i is the ionization potential of the atoms of a matter, n is the photoelectron density, c is the speed of light, $\hbar$ is Planck's constant, ω_0 is the ionizing radiation frequency, and ω_L is the electron plasma frequency, then the resulting nonequilibrium photoelectron distribution can be approximated by the function [1]

$$F(\mathbf{v}) \simeq \frac{n}{4\pi^2 v_E v_T^2} \exp\left[-\frac{v_z^2}{2v_T^2} - \frac{1}{2v_T^2}(v_\perp - v_E)^2\right]. \quad (2)$$

In expression (2), the characteristic velocities v_E and v_T depend on the ionizing radiation intensity:

$$v_E^2 = \frac{2I}{nmc}\frac{\omega_L^2}{\omega_0^2}, \quad (3)$$

$$v_T^2 = \frac{\hbar\omega_L}{2m}\sqrt{\frac{I}{cnE_i}}, \quad (4)$$

where m is the mass of an electron. In accordance with inequalities (1), we have $v_E \gg v_T$.

3. ABSORPTION AND REFLECTION OF A TEST WAVE

We consider the interaction of a test electromagnetic wave with a tunneling-ionization-produced plasma that occupies the half-space $z > 0$ and in which the nonequilibrium electron distribution is given by formulas (2)–(4). The field of a linearly polarized incident test wave can be represented as

$$\frac{1}{2}\mathbf{E}_L \exp(-i\omega t + ikz) + \text{c.c.}, \quad z < 0, \quad (5)$$

where $\mathbf{E}_L = (E_L, 0, 0)$ and $\omega = ck$, with k and E_L being the wavenumber and electric field strength, respectively. The frequency ω of the test wave is assumed to be much lower than ω_L. A wave with field (5) penetrates into a nonequilibrium plasma over a distance on the order of the skin depth and is reflected. The field of the reflected wave can be represented as

$$\frac{1}{2}\mathbf{E}_r \exp(-i\omega t + ikz) + \text{c.c.}, \quad z < 0. \quad (6)$$

Here, $\mathbf{E}_r = (E_r, 0, 0)$ with $E_r = RE_L$, where R is the complex reflection coefficient.

The electric field in the plasma has the form

$$\frac{1}{2}\mathbf{E}(z)\exp(-i\omega t) + \text{c.c.}, \quad z > 0,$$
$$\mathbf{E}(z) = (E(z), 0, 0). \quad (7)$$

The magnetic field of the incident and reflected waves is determined from the equation

$$\frac{\partial \mathbf{B}}{\partial t} = -c(\nabla \times \mathbf{E}). \quad (8)$$

Taking into account Eq. (8) and representations (5)–(7), we obtain from the continuity conditions for the tangential components of the magnetic and electric fields at the plasma boundary ($z = 0$) the following expression for the complex reflection coefficient:

$$R = \frac{Z-1}{Z+1} = |R|\exp(i\psi), \quad (9)$$

where Z is the surface plasma impedance,

$$Z = \frac{E(z=0)}{B(z=0)} = \frac{i\omega}{c}\frac{E(0)}{E'(0)} = \operatorname{Re} Z + i\operatorname{Im} Z, \quad (10)$$

and ψ is the reflection-induced phase shift,

$$\psi = \arctan\left[\frac{2\operatorname{Im} Z}{|Z|^2 - 1}\right]. \quad (11)$$

By definition, the absorption coefficient A has the form

$$A = 1 - |R|^2 = \frac{4\operatorname{Re} Z}{|Z|^2 + 1 - 2\operatorname{Re} Z}. \quad (12)$$

From expressions (9)–(12) we see that, in order to determine the absorption and reflection coefficients, it is sufficient to find surface impedance (10) of a nonequilibrium plasma.

4. SURFACE IMPEDANCE

The surface impedance can be determined using the kinetic equation for a small correction δF to the photoelectron distribution function $F(\mathbf{v})$ given by expression (2). Assuming that photoelectrons have a sufficiently high kinetic energy, we ignore collisions between them, as well as their collisions with ions, in calculating the correction δF introduced by the field of the test wave. This assumption is justified when the effective electron collision frequency is much lower than both the test wave frequency and the ratio of the characteristic electron velocity to the skin depth. We represent the correction δF in the form

$$\delta F = \frac{1}{2}\delta f(z, \mathbf{v})\exp(-i\omega t) + \text{c.c.} \quad (13)$$

to see that the function $\delta f = \delta f(z, \mathbf{v})$ satisfies the equation

$$-i\omega\delta f + v_z\frac{\partial}{\partial z}\delta f = Q(z, v_x, v_y, v_z) \equiv Q(z, v_z), \quad (14)$$

where we have introduced the notation

$$Q(z, v_z) = -\frac{e}{m}\left[E(z)\frac{\partial F}{\partial v_x} + \frac{i}{\omega}E'(z)\left(v_z\frac{\partial F}{\partial v_x} - v_x\frac{\partial F}{\partial v_z}\right)\right], \quad (15)$$

with e being the charge of an electron. We present the solution to Eq. (14) with the simplest boundary conditions, which are often used in the skin effect theory (see, e.g., [10–15]). Assuming that, far from the plasma boundary, the distribution of electrons with velocities $v_z < 0$ is unperturbed, we obtain from Eq. (14) the expression

$$\delta f^{(-)} = -\frac{1}{v_z}\int_z^{\infty} dz' \exp\left[\frac{i\omega}{v_z}(z - z')\right]Q(z', v_z), \quad v_z < 0, \quad (16)$$

where the subscript (–) indicates that $\delta f^{(-)}$ is the distribution function of the electrons that move toward the plasma surface. The convergence of the integral with respect to z' at the upper limit is ensured by the fact that the frequency ω has a small positive imaginary part $i\epsilon$, which corresponds to the switching-on of the field in the infinite past, $t \longrightarrow -\infty$. Using the condition of specular reflection of the electrons from the plasma boundary, $\delta f^{(-)}(z = 0, -v_z) = \delta f^{(+)}(z = 0, v_z)$, we obtain from Eq. (14) the following expression for the distribution function of the electrons that fly away from the plasma surface:

$$\delta f^{(+)} = \frac{1}{v_z}\int_0^{z} dz' \exp\left[\frac{i\omega}{v_z}(z - z')\right]Q(z', v_z) + \frac{1}{v_z}\int_0^{\infty} dz' \exp\left[\frac{i\omega}{v_z}(z + z')\right]Q(z', -v_z), \quad v_z > 0. \quad (17)$$

Using expressions (13), (16), and (17) for the correction to the distribution function, we can evaluate the density of the current generated in the plasma. Representing the current density as

$$\frac{1}{2}\mathbf{j}(z)\exp(-i\omega t) + \text{c.c.}, \quad \mathbf{j}(z) = (j(z), 0, 0), \quad (18)$$

we find the density of the current flowing along the x axis:

$$j(z) = e\int d\mathbf{v}\, v_x \delta f(z, \mathbf{v}). \quad (19)$$

Here, for $v_z < 0$ and $v_z > 0$, we must integrate the function $\delta f^{(-)}$ (16) and the function $\delta f^{(+)}$ (17), respectively. Substituting current density (18) into the wave equation yields the following equation for determining the field in the plasma:

$$E''(z) + \frac{\omega^2}{c^2}E(z) = -\frac{4\pi i\omega}{c^2}j(z), \quad z > 0. \quad (20)$$

We continue the field $E(z)$ and Eq. (20) into the half-space $z < 0$ in such a way that the resulting field and equation are even. The derivative $E'(z)$ of the field so continued has a jump at $z = 0$. The jump, which is equal to $2E'(0)$, is determined by the surface plasma impedance:

$$2E'(0) = 2ikB(0) = 4ikE_L/(Z + 1). \quad (21)$$

Applying the Fourier transform to Eq. (20) continued in an even fashion, we obtain

$$E(z) = \frac{2ikE_L}{\pi(Z + 1)}\int_{-\infty}^{\infty}\frac{dq}{k^2\varepsilon(\omega, |q|) - q^2}\exp(iqz), \quad (22)$$

where the surface impedance Z (10) has the form

$$Z = \frac{ik}{\pi}\int_{-\infty}^{\infty}\frac{dq}{k^2\varepsilon(\omega, |q|) - q^2}. \quad (23)$$

The dielectric function, which determines the wave field $E(z)$ and surface impedance Z, is an even function of $\mathbf{q}$. In accordance with relationships (15)–(17) and (19), it can be written as $\varepsilon(\omega, |q|) = \varepsilon_{xx}(\omega, |q|) \equiv \varepsilon_{xx}(\omega, \mathbf{q})$,

$$\varepsilon_{ij}(\omega, \mathbf{q}) = \delta_{ij} + \frac{\omega_L^2}{n\omega}\int d\mathbf{v}\frac{1}{\omega - \mathbf{q}\cdot\mathbf{v} + i\epsilon} \times\left[\delta_{sj}\left(1 - \frac{\mathbf{q}\cdot\mathbf{v}}{\omega}\right) + \frac{v_j q_s}{\omega}\right]v_i\frac{\partial F}{\partial v_s}, \quad (24)$$

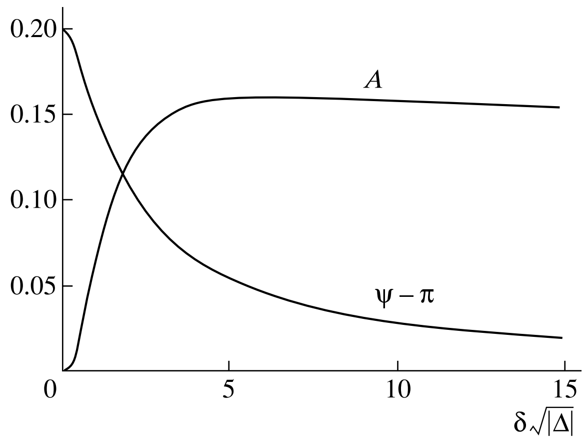
where δ_{ij} is the Kronecker delta and $\epsilon > 0$. For a plasma with a nonequilibrium photoelectron distribution given by function (2), expression (24) simplifies to the approximate form

$$\varepsilon(\omega, |q|) = 1 - \frac{\omega_L^2}{\omega^2}J_+\left(\frac{\omega}{|q|v_T}\right) - \Delta\frac{\omega_L^2}{\omega^2}\left[1 - J_+\left(\frac{\omega}{|q|v_T}\right)\right], \quad (25)$$

where we have used the following notation for Δ and $J_+(\beta)$ [16]:

$$-\Delta = \frac{v_E^2}{2v_T^2} + \frac{1}{2} \simeq \frac{v_E^2}{2v_T^2}, \quad (26)$$

$$J_+(\beta) = \frac{\beta}{\sqrt{2\pi}}\int_{-\infty}^{\infty}\frac{dx}{\beta - x + i\epsilon}\exp\left(-\frac{x^2}{2}\right), \quad \beta > 0. \quad (27)$$



Absorption coefficient A (28) and phase shift ψ (36) of the reflected wave vs. parameter $\delta\sqrt{|\Delta|}$ (32), characterizing the extent to which the skin effect is anomalous. The calculations were carried out for $\Omega = 0.1$ and $|\Delta| = 9$.

5. ABSORPTION COEFFICIENT

In accordance with definition (10), the absolute value of the surface impedance is much less than unity when the characteristic penetration depth of the field into the plasma, $|E(0)/E'(0)|$, is small in comparison with the reciprocal of the wavenumber of the test wave, $c/\omega = \lambda/2\pi$ (where λ is the wavelength). Expression (12) then implies that, for $|Z| \ll 1$, the absorption coefficient A is substantially smaller than unity and is contributed mainly by the real part of the surface impedance,

$$A \simeq 4\,\mathrm{Re}\,Z. \tag{28}$$

Taking into account relationships (23), (25), and (27), we represent absorption coefficient (28) as

$$A = \frac{8}{\pi}\delta\Omega\int_0^\infty \frac{\mathrm{Im}(x)dx}{[\mathrm{Im}(x)]^2 + [\mathrm{Re}(x)]^2}, \tag{29}$$

where we have introduced the notation $\Omega = \omega/\omega_L$ and $\delta = v_T\omega_L/\omega c$, and the real and imaginary parts of the impedance have the form

$$\mathrm{Im}(x) = \sqrt{\frac{\pi}{2}}(1-\Delta)\delta^2 x^3 \exp\left(-\frac{x^2}{2}\right), \tag{30}$$

$$\mathrm{Re}(x) = 1 + \delta^2 x^2\left[\Delta - \Omega^2 + (1-\Delta)x\exp\left(-\frac{x^2}{2}\right)\int_0^x \exp\left(\frac{t^2}{2}\right)dt\right]. \tag{31}$$

Since, under the conditions in question, we have $\delta\Omega = v_T/c \ll 1$, the dependence of $\mathrm{Re}(x)$ on Ω can be ignored. According to definitions (3) and (4) and inequalities (1), we also have $|\Delta| \simeq 1 - \Delta \simeq v_E^2/2v_T^2 \gg 1$. Ignoring the difference between the quantities $-\Delta$ and $1 - \Delta$ in formulas (30) and (31), we see that the integral with respect to x in expression (29) depends on a single parameter,

$$\delta\sqrt{|\Delta|} = \frac{v_E\omega_L}{c\omega\sqrt{2}}. \tag{32}$$

For large and small values of the parameter $\delta\sqrt{|\Delta|}$, this integral is described by simple asymptotic formulas. For $\delta\sqrt{|\Delta|} \ll 1$, the small function $[\mathrm{Im}(x)]^2$ in the denominator of the integral in expression (29), as well as the difference of the real part $\mathrm{Re}(x)$ (31) of the surface impedance from unity, can be ignored. In this case, expression (29) reduces to the approximate form

$$A \simeq 8\sqrt{\frac{2}{\pi}}\Omega|\Delta|\delta^3 = 4\sqrt{\frac{2}{\pi}}\frac{v_T}{c}\left(\frac{v_E\omega_L}{\omega c}\right)^2, \quad \delta\sqrt{|\Delta|} \ll 1. \tag{33}$$

For $\delta\sqrt{|\Delta|} \gg 1$, the main contribution to the integral in expression (29) comes from the values $x \simeq 1/\delta\sqrt{|\Delta|} \ll 1$. Using the corresponding approximate formula for the integrand in expression (29) in the limit $x \ll 1$, we can write the absorption coefficient as

$$A \simeq 4\sqrt{\frac{2}{\pi}}\Omega\delta^3|\Delta|\int_0^\infty \frac{x^3dx}{(1+\Delta x^2\delta^2)^2 + (\pi/2)\Delta^2\delta^4x^6} \simeq 4\frac{\Omega}{\sqrt{|\Delta|}} = 4\sqrt{2}\frac{v_T}{c}\left(\frac{\omega c}{v_E\omega_L}\right), \quad \delta\sqrt{|\Delta|} \gg 1. \tag{34}$$

In order to gain insight into a case intermediate between the above two limits, we show in the figure how the absorption coefficient depends on the parameter $\delta\sqrt{|\Delta|}$. The dependence was calculated for $\Omega = 0.1$ and $|\Delta| = 9$, which correspond to $\omega = 0.1\omega_L$ and $v_E \simeq 4v_T$.

Let us compare expressions (33) and (34) for the absorption coefficient with the corresponding expressions used in the theory of an equilibrium plasma with a Maxwellian electron distribution. For the normal high-frequency skin effect, $v_T/\omega \ll c/\omega_L$, in an equilibrium plasma, we have $A \simeq 8\sqrt{2/\pi}\Omega\delta^3$ with $\delta \ll 1$, and, for the anomalous skin effect, $v_T/\omega \gg c/\omega_L$, we have $A \simeq (8/3\sqrt{3})(2/\pi)^{1/6}\Omega\delta^{1/3}$ with $\delta \gg 1$. These approximate formulas for the absorption coefficient differ from expressions (33) and (34). According to the latter two expressions, the normal high-frequency skin effect changes into the anomalous one at $v_E/\omega \sim c/\omega_L$ rather than at $v_T/\omega \sim c/\omega_L$. For $v_E/\omega \ll c/\omega_L$, or for $\delta\sqrt{|\Delta|} \ll 1$, absorption coefficient (33) is greater than that for an equilibrium plasma by a factor of $v_E^2/2v_T^2 \gg 1$. For $v_E/\omega \gg c/\omega_L$, the absorption coefficient A (34) decreases with increasing parameter $\delta\sqrt{|\Delta|} \gg 1$, so, at $\delta\sqrt{|\Delta|} > \delta^{-2}$, it is less than that for an equilibrium plasma. Such a large discrepancy is attributed to the influence of the magnetic field of the test wave on the electron kinetics in the skin layer. Because of the effect of an alternating magnetic field, the absorption coefficient not only depends on the relatively low energy of thermal motion across the plasma boundary but is also determined by the substantially higher electron kinetic energy in the polarization plane of the test wave. Note that the expressions obtained for the absorption coefficient in [10] are formally similar to expressions (33) and (34), but they refer to a plasma produced in the tunneling ionization of atoms by a linearly polarized radiation that forms a bi-Maxwellian anisotropic photoelectron distribution greatly extended along the polarization vector of the ionizing wave.

6. PHASE SHIFT

For a small absolute value of the surface plasma impedance, the absolute value of the reflection coefficient is close to unity,

$$|R| \simeq 1 - 2\,\mathrm{Re}\,Z, \tag{35}$$

and the phase shift of the reflected wave is close to π,

$$\psi \simeq \pi - 2\,\mathrm{Im}\,Z. \tag{36}$$

The difference between ψ and π is largely determined by the small imaginary part of the surface impedance. Using notations (30) and (31), we can represent the imaginary part of the impedance Z (23) as

$$\mathrm{Im}\,Z = -\frac{2}{\pi}\delta\Omega\int_0^{\infty}\frac{\mathrm{Re}(x)dx}{[\mathrm{Im}(x)]^2 + [\mathrm{Re}(x)]^2}. \tag{37}$$

We present asymptotic formulas for $\mathrm{Im}Z$ that were derived for conditions such that the dependence of $\mathrm{Re}(x)$ on Ω is unimportant. In this case, the imaginary part of the impedance, $\mathrm{Im}Z$, depends on the parameter $\delta\sqrt{|\Delta|}$ (32) with $|\Delta| \gg 1$.

For $\delta\sqrt{|\Delta|} \ll 1$, the main contribution to the integral in expression (37) comes from the values $x \simeq 1/\delta \gg 1$. For such x values, the function $[\mathrm{Im}(x)]^2$ in the denominator of the integral in expression (37) can be ignored and the real part of the impedance can be described by the approximate formula $\mathrm{Re}(x) \simeq 1 + \delta^2x^2$. From relationship (36) and expression (37) we then obtain

$$\mathrm{Im}\,Z \simeq -\frac{2}{\pi}\delta\Omega\int_0^{\infty}\frac{dx}{1+\delta^2x^2} = -\Omega, \tag{38}$$

$$\psi - \pi = 2\Omega, \quad \delta\sqrt{|\Delta|} \ll 1. \tag{39}$$

For $\delta\sqrt{|\Delta|} \gg 1$, the main contribution to the integral in expression (37) comes from the values $x \simeq \sqrt{|\Delta|} \gg 1$. For such x values, the function $[\mathrm{Im}(x)]^2$ in expression (37) can be omitted and the real part of the impedance, $\mathrm{Re}(x)$, can be represented as $\mathrm{Re}(x) \simeq 1 + x^2\delta^2 + \delta^2|\Delta|$. Having made these simplifications, from relationship (36) and expression (37) we find

$$\mathrm{Im}\,Z \simeq -\frac{2}{\pi}\delta\Omega\int_0^{\infty}\frac{dx}{1+\delta^2|\Delta|+\delta^2x^2} \simeq -\frac{\Omega}{\delta\sqrt{|\Delta|}}, \tag{40}$$

$$\psi - \pi \simeq \frac{2\Omega}{\delta\sqrt{|\Delta|}}, \quad \delta\sqrt{|\Delta|} \gg 1. \tag{41}$$

The dependence of the phase shift on $\delta\sqrt{|\Delta|}$, calculated for $\Omega = 0.1$ and $|\Delta| = 9$, also is shown in the figure.

In the case of reflection of a test wave from a thermodynamically equilibrium plasma, the phase shift is described by the following asymptotic formulas: $\psi - \pi \simeq 2\Omega$ for $\delta \ll 1$ and $\psi - \pi \simeq (4/3)(2/\pi)^{1/6}\Omega\delta^{1/3}$ for $\delta \gg 1$. Expression (39) is similar to the corresponding expression in the case of equilibrium plasma but is applicable over a narrower parameter range, such that $\delta\sqrt{|\Delta|} \ll 1$ with $|\Delta| \gg 1$. According to relationships (41), at $\delta\sqrt{|\Delta|} \gg 1$, the difference between ψ and π decreases with increasing $\delta\sqrt{|\Delta|}$ and turns out to be less than that for an equilibrium plasma. Relationships (38) and (40) are similar to those derived in [11] for a plasma in which photoelectrons obey a strongly anisotropic axisymmetric bi-Maxwellian distribution and which interacts with a test wave polarized along the anisotropy axis of the photoelectron distribution.

7. DISCUSSION

A photoionized plasma exhibits the above optical properties only when its state is close to the initial nonequilibrium state. There are several factors that destroy the nonequilibrium state and thereby limit the time interval during which the anomalies in the absorption and reflection of a test wave can be observed. One such factor is the hydrodynamic expansion of a nonequilibrium hot plasma into a vacuum, accompanied by the destruction of the sharp plasma boundary. In order of magnitude, the characteristic expansion time $\tau_{\exp}$ is determined by the ratio of the effective skin depth d to the speed of sound v_s: $\tau_{\exp} \sim d/v_s$. Consequently, the duration of the test wave pulse, τ_p, should be less than $\tau_{\exp}$. On the other hand, the pulse duration τ_p should not be too short. The reason is that, because of the spectral broadening of the test pulse, the absorption coefficient depends on the pulse duration τ_p when the latter is less than $1/\omega$ (see [17] for details). For $\omega\tau_p \gtrsim 1$, this effect is unimportant, however. The inequalities $\omega\tau_p \gtrsim 1$ and $\tau_{\rm exp} > \tau_p$ yield the inequality $\omega\tau_{\rm exp} > 1$.

Another factor that destroys the nonequilibrium state is electron collisions, due to which nonequilibrium photoelectrons relax to a Maxwellian distribution. According to [18], the characteristic time scale τ_c on which the initial anisotropic photoelectron distribution is relaxed by electron–ion collisions is on the order of $\tau_c \sim v_E^2/\nu v_T^2$, where ν is the electron–ion collision frequency dependent on the velocity v_E. This yields the following restrictions on the duration and frequency of the test wave pulse: $\omega\tau_c > \omega\tau_p \gtrsim 1$.

The third factor is the onset of the Weibel instability [5], accompanied by the generation of a nonuniform quasi-stationary magnetic field, which in turn exerts an inverse effect on the shape of the anisotropic photoelectron distribution. When the photoelectron distribution is described by expression (2), the maximum growth rate γ_m of the Weibel instability is equal to $\gamma_m = \omega_L v_E/\sqrt{2}c$ [18, 19]. In the initial (linear) stage of instability, the instability-generated magnetic field increases exponentially, $B_q \sim B_{\rm sp}\exp(\gamma_m t)$, where $B_{\rm sp}$ is the strength of the spontaneous magnetic field. The magnetic field B_q may produce a significant effect when it is comparable in strength to the magnetic field of the test wave, $B_L = E_L$, or when the electromagnetic energy density $B_q^2/4\pi$ is not too low in comparison with the kinetic energy density of the photoelectrons, $\sim nmv_E^2$. Under the assumption that the spontaneous magnetic field $B_{\rm sp}$ is much weaker than $B_m = \min(B_L, \sqrt{4\pi nm}v_E)$, the effect of the Weibel instability on the optical plasma properties can be considered inessential if

$$\omega\tau_w \sim \frac{\omega}{\gamma_m}\ln\frac{B_m}{B_{\rm sp}} = \left(\frac{\omega c\sqrt{2}}{\omega_L v_E}\right)\ln\frac{B_m}{B_{\rm sp}} > \omega\tau_p \gtrsim 1. \quad (42)$$

Note that, for large values of the parameter $\delta\sqrt{|\Delta|}$ (32), the last of conditions (42) can be satisfied at the expense of a large value of the ratio $B_m/B_{\rm sp} \gg 1$.

ACKNOWLEDGMENTS

This work was supported in part by the RF Program for State Support of Leading Scientific Schools, project no. NSh-1385.2003.2.

REFERENCES

1. N. B. Delone and V. P. Krainov, J. Opt. Soc. Am. B **8**, 1207 (1991).
2. N. B. Delone and V. P. Kraĭnov, Usp. Fiz. Nauk **168**, 531 (1998) [Phys. Usp. **41**, 469 (1998)].
3. V. P. Silin, Zh. Éksp. Teor. Fiz. **117**, 926 (2000) [JETP **90**, 805 (2000)].
4. V. D. Mur, S. V. Popruzhenko, and V. S. Popov, Zh. Éksp. Teor. Fiz. **119**, 893 (2001) [JETP **92**, 777 (2001)].
5. V. I. Arefyev, V. P. Silin, and S. A. Uryupin, Phys. Lett. A **255**, 307 (1999).
6. W. P. Leemans, C. E. Clayton, W. B. Mori, et al., Phys. Rev. Lett. **68**, 321 (1992).
7. P. B. Corkum, N. H. Barnett, and F. Brunel, Phys. Rev. Lett. **62**, 1259 (1989).
8. D. Giulietti, L. A. Gizzi, A. Giulietti, et al., Phys. Rev. Lett. **79**, 3194 (1997).
9. S. J. McNaught, J. P. Knauer, and D. D. Meyerhofer, Phys. Rev. Lett. **78**, 626 (1997).
10. G. Ferrante, M. Zarcone, and S. A. Uryupin, Eur. Phys. J. D **19**, 349 (2002).
11. G. Ferrante, M. Zarcone, and S. A. Uryupin, Eur. Phys. J. D **22**, 109 (2003).
12. G. Ferrante, M. Zarcone, and S. A. Uryupin, Phys. Rev. Lett. **91**, 085 005 (2003).
13. A. A. Andreev, E. G. Gamaliĭ, V. N. Novikov, et al., Zh. Éksp. Teor. Fiz. **101**, 1808 (1992) [Sov. Phys. JETP **74**, 963 (1992)].
14. W. Rozmus and V. T. Tikhonchuk, Phys. Rev. A **42**, 7401 (1990).
15. T.-Y. Brian Yang, W. L. Kruer, A. B. Langdon, and T. W. Johnston, Phys. Plasmas **3**, 2702 (1996).
16. V. P. Silin and A. A. Rukhadze, *Electromagnetic Properties of Plasmas and Plasmalike Media* (Gosatomizdat, Moscow, 1961) [in Russian].
17. V. A. Isakov, A. P. Kanavin, and S. A. Uryupin, J. Russ. Laser Res. **25**, 401 (2004).
18. A. Yu. Romanov, V. P. Silin, and S. A. Uryupin, Zh. Éksp. Teor. Fiz. **126**, 843 (2004) [JETP **99**, 727 (2004)].
19. V. P. Krainov, Zh. Éksp. Teor. Fiz. **123**, 487 (2003) [JETP **96**, 430 (2003)].

Translated by I.A. Kalabalyk