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High-frequency waves in plasma formed as a result of tunnel ionization of atoms by circularly polarized radiation

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ABSTRACT

New dependencies of frequency and damping decrement of high-frequency longitudinal waves on the wave vector in photoionized plasma formed by tunnel ionization of atoms in the field of circularly polarized radiation are found. Weakly damped longitudinal waves with a frequency much higher than the electron Langmuir frequency are predicted.

1. Introduction

The ionization of substance atoms by the field of intense laser radiation leads to creation of nonequilibrium plasma with anisotropic velocity distribution of photoelectrons (see, e.g., [1–9]). The difference of the photoelectron distribution from the Maxwell distribution results in the fact that many properties of photoionized plasma differ qualitatively from those of weakly nonequilibrium plasma. Due to the effect of an alternating magnetic field on the photoelectron kinetics the absorption and reflection of test electromagnetic radiation have new unusual properties [10, 11]. The effect of tunnel and multiphoton ionization on the generation of terahertz radiation in the field of femtosecond laser pulses is described [12]. Polarization dependent terahertz response from plasma with anisotropic photoelectron distribution is predicted [13]. In the paper [14] the possibility of terahertz electromagnetic waves amplification in photoionized xenon plasma is shown. The anisotropy of photoelectrons leads to the possibility of aperiodic instability development [15]. Under the conditions of such instability development the opportunity to amplify the low-frequency pulses of electromagnetic radiation by photoionized plasma is predicted [16–18]. Photoionization two-stream instability is studied in the plasma created during the interaction of ultrashort X-ray pulse with gas target [19]. In the present paper it is shown that the dispersion law and the collisionless damping of high-frequency longitudinal electron waves in photoionized plasma may significantly differ from those inherent in the equilibrium plasma. The dispersion equation for longitudinal electron waves in photoionized plasma formed by tunnel ionization of substance atoms in the field of intense circularly polarized laser radiation is studied below. It is found that the dispersion law and the damping of the waves propagating along the anisotropy axis of the photoelectron distribution are the same as in equilibrium plasma. In contrast, the waves propagating across the anisotropy axis have a different dispersion law and significantly greater damping. At the same values of wave numbers the group velocity of these waves is much larger than the group velocity of the waves propagating along the anisotropy axis, and in the short-wave region it is close to the phase velocity. In the region of small wave numbers the frequency of the waves is close to the Langmuir frequency of photoelectrons. At large wave numbers a linear dispersion law with characteristic frequency much higher than the Langmuir frequency is obtained.

2. Photoelectron velocity distribution

Let us consider fully ionized plasma formed as a result of substance atom ionization in the field of ultrashort pulse of circularly polarized laser radiation with the field of the form

$$\mathbf{E}_{pump} = E_{pump}\left\{-\sin(\omega_{pump}t), \cos(\omega_{pump}t), 0\right\}, \tag{1}$$

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where ω_{pump} is the carrier frequency of the ionizing pulse, E_{pump} is the strength of its electric field. The pulse duration will be considered more than the time of atoms ionization, but less than the time of nonequilibrium photoelectrons distribution variations. We assume that frequency and field strength (1) satisfy the conditions

$$mv_E^2 \gg 2E_i \gg \frac{3}{2}\hbar\omega_{pump}\sqrt{\frac{mv_E^2}{2E_i}}, \tag{2}$$

where $v_E = |eE_{pump}/m\omega_{pump}|$ corresponds to velocity of photoelectron circular motion determined by the field of ionizing radiation (1), e and m are the charge and the mass of electron, $\hbar$ is Planck's constant, E_i is the ionization potential of substance atoms. Under these conditions the regime of tunnel ionization of atoms in the electric field (1) is realized, and photoelectrons distribution over velocities $\mathbf{v}$ corresponds to the ionization probability $W(\mathbf{v})$, derived in the paper [1]

$$W(\mathbf{v}) \propto \exp\left\{-\frac{2}{3\hbar\omega_{pump}\sqrt{m}v_E}\left[2E_i + mv_z^2 + m(v_\perp - v_E)\right]^{3/2}\right\}. \tag{3}$$

From this formula and from the inequalities (2), it follows that the distribution function of bulk photoelectrons with velocities v_z^2, $(v_\perp - v_E)^2 \ll 2E_i/m$ can be approximately represented by function of the form [1,20]

$$f(\mathbf{v}) \simeq \frac{n}{4\pi^2 v_E v_T^2}\exp\left[-\frac{(v_\perp - v_E)^2}{2v_T^2} - \frac{v_z^2}{2v_T^2}\right], \tag{4}$$

where $v_\perp = \sqrt{v_x^2 + v_y^2}$, n is the photoelectron density, and

$$v_T = \sqrt{\frac{\hbar\omega_{pump}v_E}{2\sqrt{2mE_i}}} \ll v_E \tag{5}$$

characterizes the spread of the photoelectron velocity distribution. The expression (4) is applicable under the conditions of inequality $v_E \gg v_T$, when the mean kinetic energy of photoelectron motion across the anisotropy axis is greater than along it. Further we will understand the relationship v_E/v_T as the degree of anisotropy of photoelectrons distribution. The anisotropic photoelectron distribution of the form (4) could be implemented, for example, when hydrogen atoms are ionized by a pulse of circularly polarized radiation with the energy flux density of value $I_{pump} \approx 2 \cdot 10^{16}$ W/cm^2, the frequency $\omega_{pump} \approx 2 \cdot 10^{15}$ s^{-1} and the duration of a few tens of femtoseconds. Under these conditions we have $mv_E^2 \simeq 3.2$ keV, $mv_T^2 \simeq 7.2$ eV. Note that in this case it is difficult to talk about the realization of the tunnel ionization regime (see [4]), but it can be assumed that the photoelectron distribution is close to that described by the function (4).

The anisotropic distribution (4) is formed during the time comparable to the ω_{pump}^{-1} and exists within a limited time interval. Firstly, the electron collisions lead to isotropization of the distribution function (4). The characteristic time of isotropization t_c is determined by the inverse frequency of the electron-ion collisions $t_c \sim m^2 v_E^3 (4\pi Ze^4 n\Lambda)^{-1}$ [21], where Z is the degree of ion ionization, Λ is the Coulomb logarithm. For the specified above values of energy flux density and frequency of the ionizing radiation in a plasma with the photoelectron number density $n = 10^{18}$ cm^{-3} and typical values of the parameters $Z = 1$, $\Lambda = 7$ we obtain estimate $t_c \approx 70$ psec. We note that under the considered conditions the collision damping decrement of high-frequency waves is of the order of the electron-ion collision frequency and is two orders of magnitude smaller than the frequency of the analyzed waves. Secondly, plasma with anisotropic electron velocity distribution (4) is unstable with respect to the development of aperiodic instability [15,16,22]. In numerical modeling of aperiodic instability development it is shown [23] that conversion of the instability to the nonlinear stage, when an essential change in the initial anisotropic electron distribution function occurs, is realized when the energy density of the generated quasi-stationary magnetic field in plasma reaches ten percent of the electron kinetic energy density. According to [16] it happens after the time interval $t_{NL} \sim 10 \div 15\gamma_i^{-1}$, where $\gamma_i \approx (v_E/\sqrt{2}c)\omega_L \ll \omega_L$ is the maximum possible growth rate of this instability (see, e.g., [20,25]), and for the parameters of ionizing radiation and formed plasma considered above we have $t_{NL} \approx 5$ psec. Here $\omega_L = \sqrt{4\pi e^2 n/m}$ is the Langmuir electron frequency. Both time intervals t_c and t_{NL}, characterizing the existence time interval of the original distribution of photoelectrons (4), are much longer than the period of the Langmuir oscillations $2\pi\omega_L^{-1} \approx 0.1$ psec. Therefore, being interested in the properties of high-frequency electron plasma waves with frequency close to or higher than ω_L, we neglect time variations of the original distribution function (4).

In connection with the possible effect of instabilities on the lifetime of distribution of the form (4), we note that in plasma with such distribution there is no purely aperiodic instability, which is an analog of the two-stream instability [24]. Concluding the discussion of the non-equilibrium distribution (4) existence conditions, we note that using it we confined ourselves to the range of radiation intensities at which relativistic effects can be neglected. The transition to high intensities is accompanied by a significant modification in the shape of anisotropic photoelectron distribution [3]. The study of high-frequency properties of plasma produced in this case requires a separate, original consideration.

3. Dispersion law and collisionless damping of longitudinal electron waves

It is known that in anisotropic plasma at arbitrary orientation of the wave vector of the perturbation $\mathbf{k}$ relative to the anisotropy axis the waves are not purely longitudinal or transverse ones (see, e.g., [26]). However, in special cases of wave vector orientation exactly along the anisotropy axis or across it both longitudinal and transverse waves may propagate in plasma. The dispersion equation for longitudinal waves in these limit cases has the form

$$\varepsilon(\omega, \mathbf{k}) = 1 + \frac{\omega_L^2}{nk^2}\int d\mathbf{v}\,\frac{1}{\omega - \mathbf{kv}}\left(\mathbf{k}\frac{\partial f(\mathbf{v})}{\partial \mathbf{v}}\right)$$
$$\equiv 1 + \frac{\omega_L^2}{n}\frac{\partial}{\partial\omega}\int d\mathbf{v}\,\frac{f(\mathbf{v})}{\omega - \mathbf{kv}} = 0, \tag{6}$$

where ω is the complex frequency of the considered waves. Below we consider weakly damping waves. In this case the real frequency of these waves $\omega'(\mathbf{k}) \equiv \mathrm{Re}[\omega(\mathbf{k})]$ is determined by the equation $\mathrm{Re}[\varepsilon(\omega', \mathbf{k})] = 0$. The damping decrement, that is small compared with the frequency, is determined by the relation

$$\gamma(\mathbf{k}) \equiv \mathrm{Im}[\omega(\mathbf{k})] \simeq -\left.\frac{\mathrm{Im}\left[\varepsilon(\omega', \mathbf{k})\right]}{\partial\,\mathrm{Re}\left[\varepsilon(\omega', \mathbf{k})\right]/\partial\omega'}\right|_{\omega'=\omega'(\mathbf{k})}. \tag{7}$$

Thus, in order to find the frequency and the damping decrement it is necessary to calculate the real and the imaginary part of $\varepsilon(\omega, \mathbf{k})$, as a function of a real variable ω'.

Taking into account the dependence of the initial photoelectron distribution function (4) on v_z it is easy to see that in the limit case when the wave vector $\mathbf{k} \equiv \mathbf{k}_\parallel = \{0, 0, k_z\}$ is directed along the anisotropy axis, the expressions for frequency and damping decrement of high-frequency longitudinal waves have the form similar to their analogues for the Langmuir waves in the plasma with