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High-frequency skin effect in a photoionized inert gas plasma

K. Yu. Vagin[1], **T. V. Mamontova**[1,2], **S.A. Uryupin**[1,2]
[1]Lebedev Physical Institute, Russian Academy of Science, Moscow, Russia
[2]National Research Nuclear University MEPhI, Moscow, Russia

E-mail: uryupin@sci.lebedev.ru

Abstract. The penetration of a monochromatic electromagnetic wave into the weakly ionised plasma formed by multiphoton ionization of inert gas atoms has been studied. It is shown how the type of dependence of the photoelectron collision frequency on their velocity determines transverse permettivity in plasma. Explicit expressions for the absorption coefficient in inert gases at different probe radiation frequencies are obtained. In the high-frequency skin effect mode, the possibility of the absorption coefficient increasing due to the Ramsauer-Townsend effect has been demonstrated with argon as an example.

1. Introduction

At multiphoton ionization of inert gas atoms by a short laser pulse, the plasma with an nonequilibrium distribution of electrons is formed (see, for example, [1, 2]). While at times longer than momentum relaxation time, but shorter than energy relaxation time, it is possible to approximate the electron distribution by a function with one narrow energy peak [3, 4]. The form of photoelectron distribution function has significant impact on plasma properties. In particular, in an isotropic photoionized plasma the existence of longitudinal waves with a linear dispersion law and a group velocity comparable to the characteristic velocity of photoelectrons is possible [5, 4]. New effects appear in a photoionized inert gas plasma due to the Ramsauer-Townsend effect [6, 7, 8]. The specific behaviour of the photoelectron cross-section that corresponds to this effect leads to the possibility of amplifying low-frequency pulses [3, 9, 10].

This paper considers features of a monochromatic electromagnetic field penetration into the plasma formed by multiphoton ionization of inert gases. Special attention is given to studying the effects associated with the non-monotonic behaviour of the electron cross-section at neutral inert gas atoms. The expressions for a probe radiation absorption coefficient were obtained in high-frequency skin effect mode. It has been shown that the presence of a positive collision frequency derivative at the point corresponding to the characteristic velocity of photoelectrons leads to an increase in absorption by several times.

2. Penetration of an electromagnetic wave into a photoionized plasma

Let us consider the penetration of a monochromatic electromagnetic wave into the weakly ionised plasma formed by multiphoton ionization of inert gas atoms by a short pulse of laser radiation. We limit ourselves to considering the case when an energy distribution of photoelectrons has the form of one narrow peak corresponding to absorption of a minimum required number of

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K photons to overcome the ionization threshold determined by the ionization potential E_i of the atom. Under the influence of a linearly polarised ionising radiation field, an anisotropic distribution is formed in which the electron velocities are mainly oriented along and against the field strength vector. However, under conditions of weak ionization, the scattering of photoelectrons at neutral atoms leads to rapid relaxation of their momentum. Therefore, at times exceeding the inverse effective frequency of photoelectrons elastic collisions, it is justified to use the isotropic velocity distribution function $f_0\,(v)$. The nonequilibrium distribution $f_0\,(v)$ relaxes to Maxwell's distribution at significantly longer times due to electron-electron collisions and elastic collisions of photoelectrons with neutral atoms, where the amount of transferred energy is determined by a small the electron mass to the half of atom mass ratio. We will consider the influence of radiation, the frequency of which is much higher than the characteristic inverse relaxation time of the function $f_0\,(v)$, at the photoionized plasma.

We consider the interaction of a plane monochromatic electromagnetic wave

$$\mathbf{E}_i(z,t) = \frac{1}{2}\mathbf{E}_0 \exp\left[-i\omega(t - z/c)\right] + c.c., \quad \mathbf{E}_0 = (E_0, 0, 0) \tag{1}$$

with the photoionised plasma, which has the isotropic distribution of electrons $f_0\,(v)$ and occupies half-space $z > 0$. An electromagnetic wave generates in the plasma an electric field $(1/2)\mathbf{E}\,(z) \cdot \exp\left(-i\omega t\right) + c.c.$ directed along the $0x$ axis and causes a little perturbation to the photoelectron velocity distribution function $(1/2)\,\delta f(\mathbf{v}, z) \exp\left(-i\omega t\right) + c.c.$ We are interested in the conditions under which the electron displacement over the period of the incident wave change is small compared to the characteristic depth of the field penetration. The effects connected with electron movement are significant in anomalous skin effect mode, which has been described by us in [11]. Taking above mentioned conditions into account, when we find $\delta f\,(\mathbf{v}, z)$, we can disregard the term with a spatial derivative in the kinetic equation. We will use the linearised kinetic equation with the collision integral that describes photoelectrons velocity directions relaxation without their energy changing

$$-i\omega\delta f(\mathbf{v}, z) + \frac{ev_x E\,(z)}{mv}\frac{\partial f_0(v)}{\partial v} = -\nu\,(v)\left[\delta f(\mathbf{v}, z) - \int \frac{d\Omega}{4\pi}\delta f(\mathbf{v}, z)\right], \tag{2}$$

where e is the electron charge, $d\Omega$ is a solid angle element in the velocity space, $\nu\,(v)$ is the characteristic collision frequency of electrons. In a rarefied plasma, photoelectron collision frequencies are relatively low and the conditions are of interest when $\omega \gg \nu$. In this case, a small perturbation of the distribution function corresponding to the equation (2) solution can be presented as follows

$$\delta f(\mathbf{v}, z) = -i\frac{ev_x E\,(z)}{mv\omega}\frac{\partial f_0(v)}{\partial v}\left(1 - i\frac{\nu\,(v)}{\omega}\right). \tag{3}$$

Computing the electric current density from (3) and using Maxwell's equations, we obtain the following equation for the field in plasma

$$\frac{d^2 E\,(z)}{dz^2} + \frac{\omega^2}{c^2}\varepsilon_{tr}\,(\omega)\,E\,(z) = 0, \tag{4}$$

where transverse permittivity has the form

$$\varepsilon_{tr}(\omega) = 1 - \frac{\omega_L^2}{\omega^2}\left[1 - \frac{i}{n}\int d\mathbf{v}\, f_0\,(v)\left(1 + \frac{1}{3}\frac{\partial \ln[\nu\,(v)]}{\partial \ln v}\right)\frac{\nu\,(v)}{\omega}\right]. \tag{5}$$

In (5) n is a photoelectron density, $\omega_L = \sqrt{4\pi n e^2/m}$ is the Langmuir frequency of electrons.

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In the range of frequencies not exceeding the Langmur electron frequency $\omega < \omega_L$, the real part of the transverse permittivity $\mathrm{Re}[\varepsilon_{tr}(\omega)] < 0$. Then, using the continuity property of the electric and magnetic fields tangential components at the plasma boundary, for the electric field in plasma, which corresponds for the solution of the equation (4) decaying deep into plasma, we get

$$E(z) = \frac{2E_0}{1 + i\sqrt{-\varepsilon_{tr}(\omega)}} \exp\left[-\sqrt{-\varepsilon_{tr}(\omega)}\,\frac{\omega}{c}\,z\right], \tag{6}$$

with $\mathrm{Re}[\sqrt{-\varepsilon_{tr}(\omega)}\,] > 0$ and, according to (5), $\mathrm{Im}[\sqrt{-\varepsilon_{tr}(\omega)}\,] < 0$. The value $c/\omega\,\mathrm{Re}[\sqrt{-\varepsilon_{tr}(\omega)}\,]$ determines the characteristic depth of the electric field (6) penetration into a plasma.

Let us define the reflection coefficient at the plasma boundary as $R = E_r/E_0$, where E_r is the amplitude of the wave moving from the plasma boundary into the $z < 0$ region. Then the absorption coefficient $A(\omega)$, characterising the fraction of the energy transmitted to a plasma by the incident wave (1), is as follows

$$A(\omega) = 1 - |R|^2 = \frac{-4\mathrm{Im}\left[\sqrt{-\varepsilon_{tr}(\omega)}\right]}{\left(\mathrm{Im}\left[\sqrt{-\varepsilon_{tr}(\omega)}\right] - 1\right)^2 + \left(\mathrm{Re}\left[\sqrt{-\varepsilon_{tr}(\omega)}\right]\right)^2}. \tag{7}$$

3. Field absorption in inert gases

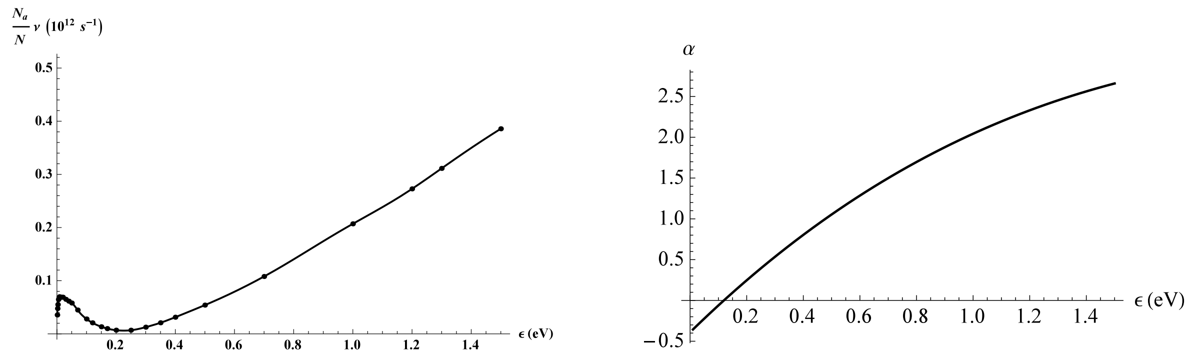
Let us consider the features of probe radiation absorption by the plasma formed during multiphoton ionization of inert gas atoms. In the case of multiphoton ionization threshold process, the energy ϵ_0 that corresponds to the maximum of photoelectron distribution function does not exceed a few eV. In inert gases, the scattering of electrons in this energy range is mainly determined by their elastic collisions with neutral atoms. Then, in the equation (2) the characteristic collision frequency of electrons is $\nu(v) = N\sigma_{tr}(v)\,v$, where N is the concentration of neutral atoms, $\sigma_{tr}(v)$ is the transport cross-section of electrons elastic collisions with neutral atoms, which depends on the photoelectron velocity. A feature of inert gases is the minimum of function $\sigma_{tr}(v)$ in the energy range slightly less than 1eV corresponding to the Ramsauer-Townsend effect [6, 7, 8]. The dependency $\nu(\epsilon)$ for argon is shown in the figure 1(a). Points on the curve of the figure. 1(a) correspond to the experimental data for the transport cross-section of the elastic collisions of electrons with neutral Ar atoms, given at [12]. Generally speaking, in argon the elastic collision frequency increases in the energy range $\epsilon \sim (0.2 \div 15)$ eV. However, starting with $\epsilon = 1.5$ eV, the cross-section of inelastic electron collisions becomes non-zero. Let us therefore limit our consideration by the energies from 0.2 eV to 1.5 eV. Similar dependencies for the xenon and argon atoms are given in [13].

To describe the velocity distribution of photoelectrons, we will use a relatively simple model of a 'cold' plasma that does not take into account the velocity spread in the energy peak. This model corresponds to an approximation of the photoelectron distribution by an isotropic function $f_0(v) = \left(n/4\pi v_0^2\right)\delta(v - v_0)$, where $v_0 = \sqrt{2\epsilon_0/m}$, $\delta(x)$ is Dirac delta function. Then the transverse permittivity (5) will be written as follows

$$\varepsilon_{tr}(\omega) = 1 - \frac{\omega_L^2}{\omega^2}\left[1 - i\left(\frac{\alpha}{3} + 1\right)\frac{\nu_0}{\omega}\right], \tag{8}$$

where $\nu_0 \equiv \nu(v_0)$, $\alpha = \partial\ln[\nu(v_0)]/\partial\ln v_0$ is a parameter determined by the dependence of collision frequency on the photoelectron velocity. The dependence $\alpha(\epsilon)$ is shown in the figure 1(b). When we obtain the curve 1(b) we use the collision frequency dependence on energy shown in the figure 1(a).

Let us return to the study of the coefficient of electromagnetic radiation absorption by a photoionized plasma. First, we consider the case, when the incident wave frequency does not



(a) (b)

Figure 1. Dependencies of the frequency of photoelectron elastic collisions with neutral argon atoms ν and the parameter α on the energy of photoelectrons ϵ. $N_a = 2.5 \cdot 10^{19} cm^{-3}$ is the concentration of atoms at atmospheric pressure.

exceed the electron Langmuir frequency and is not too close to it

$$\omega_L - \omega \gg \frac{\nu_0}{4}\left(\frac{\alpha}{3}+1\right). \tag{9}$$

For the function $\sqrt{-\varepsilon_{tr}(\omega)}$, that defines reflection and absorption coefficients in such conditions, we have

$$\sqrt{-\varepsilon_{tr}(\omega)} \equiv \mathrm{Re}\sqrt{-\varepsilon_{tr}(\omega)} + i\mathrm{Im}\sqrt{-\varepsilon_{tr}(\omega)} \approx \sqrt{\frac{\omega_L^2}{\omega^2}-1} - i\,\frac{\omega_L^2}{2\omega^2}\left(\frac{\alpha}{3}+1\right)\frac{\nu_0}{\sqrt{\omega_L^2-\omega^2}}. \tag{10}$$

By substituting (10) in (7), for the frequencies (9) not too close to plasma frequency we find the absorption coefficient

$$A(\omega) \approx \left(\frac{\alpha}{3}+1\right)\frac{2\nu_0}{\sqrt{\omega_L^2-\omega^2}}. \tag{11}$$

The absorption coefficient (11) is proportional to the collision frequency of electrons ν_0 and increases with the probe wave frequency ω. This expression differs from the known [14] by the presence of the parameter α which takes into account the dependence of collision frequency on velocity. As can be seen in the figure 1(b), due to the Ramsauer-Townsend effect, the α parameter changes between 0 and 2.7, which is accompanied by an almost doubling of the absorption.

In the limit, when the frequency of the wave falling on plasma is close to electron Langmuir one and inequality

$$\omega_L - \omega \ll \frac{\nu_0}{4}\left(\frac{\alpha}{3}+1\right) \tag{12}$$

is fair, from the expression for dielectric permittivity (8) we get

$$\sqrt{-\varepsilon_{tr}(\omega)} = \frac{1-i}{\sqrt{2}}\frac{\sqrt{\omega_L\,\nu_0\,(\alpha/3+1)}}{\omega}. \tag{13}$$

Under such conditions, the absorption coefficient (7) is as follows

$$A(\omega \approx \omega_L) \approx \sqrt{\frac{8\,\nu_0}{\omega_L}\left(\frac{\alpha}{3}+1\right)}. \tag{14}$$

A comparison of the expressions (11) and (14) shows that in the high-frequency skin effect as the frequency ω approaches the Langmuir frequency ω_L the wave absorption increases.

4. Conclusion

In this paper we study the features of monochromatic wave penetration into the semi-bounded plasma obtained as a result of multiphoton ionization of inert gas atoms in the high-frequency skin effect mode. Expressions for transverse permittivity and absorption coefficient taking into account the dependence of photoelectrons collision frequency on their velocity are obtained. In the limit that does not take into account the spread of photoelectron velocities within the peak of distribution function, explicit expressions for the absorption coefficient in different frequency ranges of the probe wave are obtained. For inert gases, the main difference is the presence of the parameter α - the value determined by the average energy of photoelectrons and the type of the transport cross-section dependence on energy. For frequencies noticeably smaller than the Langmuir frequency, the field attenuates at $\sim c/\sqrt{\omega_L^2 - \omega^2}$ distance into the plasma and the absorption coefficient is linearly dependent on the collision frequency and its derivative. When the field frequency approaches the Langmuir frequency, the effective penetration depth is inversely proportional to $\sqrt{\nu + (1/3)v\,\partial\nu/\partial v}$ and turns out to be much greater than at lower frequencies. The penetration also increases and depends on the collision frequency and its derivative as $\sqrt{\nu + (1/3)v\,\partial\nu/\partial v}$. The presence of the Ramsauer-Townsend effect leads to an increase in absorption over the whole range of the frequencies under consideration.

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