Directed motion of electrons in gases under the action of photon flux

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The phenomenon of directed motion of electrons and ions in gases under the action of ionizing radiation pressure is investigated. It is shown that for photon energies from the thresholds of atomic photoionization to several keV the photoionization process is the main mechanism for the transfer of electromagnetic radiation momentum to an atom. Expressions for the drag currents that appear under the action of ionizing radiation in atomic gases and their mixtures are obtained. The connection between the drag currents and one of the nondipole asymmetry parameters is established. Experimental investigation of the drag currents for use in precision measurement of the asymmetry parameters is discussed, particularly for small photoelectron energies where it is difficult to apply the traditional experimental schemes to measure the differential cross sections for photoionization. Nondipole parameters for the Ne $2s$, Ne $2p$, and Ar $1s$ subshell photoionization are calculated and compared with measurements and other calculations. Partial drag currents for the Ne $2s$, Ne $2p$, and Ar $1s$ subshells are also presented.

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I. INTRODUCTION

Light pressure is found in many natural phenomena and plays a particularly important role in two fields of physics, namely, astrophysics and atomic and subatomic fields. In astrophysics the existence of light pressure explains many processes occurring in interstellar media and in stellar atmospheres [1]. On the atomic level the pressure of resonance radiation effectively acts upon individual atoms forcing them to move with acceleration many times greater than that of gravity [2]. Resonance light scattering is the mechanism for light induced drift of atoms in binary gas mixtures [3–5]. In all these cases, photon scattering either by atomic particles or by unbound protons and electrons is conventionally considered to be the main process of momentum transfer to particles in rarefied gases. The scattering processes occur with particularly high probability for radiation frequencies close to those of discrete atomic transitions. With increase of radiation frequency while going outside the resonance range, the efficiency of photon scattering processes decreases by several orders of magnitude.

For photon energies higher than the ionization potentials of target atoms, the other mechanism of light pressure becomes more important, namely, the photoionization of the atomic particles in which the photon momentum is transferred to the ion and the electron after absorption. In this process, as well as in photon scattering, transfer of the photon momentum to the electron and the ion takes place. This insufficiently studied manifestation of electromagnetic radiation pressure upon matter in the gas phase plays an essential role in many natural phenomena. For instance, ongoing processes in the atmospheres of some stars are explained, to a considerable extent, by photon momentum transfer during atomic photoionization [6–8]. There is no doubt that the role of these processes is great in upper layers of the Earth's atmosphere, where the intensity of the ultraviolet component of solar radiation is sufficient to generate considerable flows of charged particles.

According to conservation laws the photon energy and momentum cannot be given to a single particle, e.g., the atom; the presence of a third body is necessary. This role in the conventional mechanism is played by the quasielastically scattered photon. In photoionization the role of the third body is played by the electron emitted after photon absorption. In this process the photon momentum is shared between the product ion and the ejected photoelectron. The fraction of the momentum given to the photoelectron in the former case was calculated in the well known Sommerfeld work [9,10]. These results were used in the calculation of radiation pressure in stellar atmospheres [6,7], and have been used in many subsequent papers. It was shown [11,12] that the directed motion of photoelectrons due to the photon momentum generates the macroscopic currents in gases.

When the first papers [11,12] dedicated to this phenomenon were published, the possibilities of studying under laboratory conditions the light pressure created in the photoionization process were limited. The limitation resulted from the low intensity of the ultraviolet radiation sources, from x-ray tubes [13–15]. Now, available synchrotrons and storage rings provide intense, tunable, and highly polarized x-ray beams. Experiments that investigate photoelectron angular distributions are carried out by many laboratories (see, for example, [16,17]). In this context it is reasonable and important to discuss the problem of ionizing light pressure and drag currents of electrons by photons not only as a physical phenomenon, but also as a possible method of precision measurement of nondipole terms in the differential cross section for the photoeffect; hence the thrust of this paper.

II. REDISTRIBUTION OF PHOTON MOMENTUM BETWEEN PHOTOELECTRON AND ION

Let us consider the passing of monochromatic electromagnetic radiation with frequency ω through an atomic gas under standard conditions. Suppose that the thermal distribution function $f(\vec{v}_T)$ of the gas atoms is isotropic and normalized by the condition $\int f(\vec{v}_T)d\vec{v}_T=1$. Let us further consider the photoionization of an atom moving in the gas with a velocity $\vec{v}_T$. After a photon is absorbed, its momentum $\hbar\vec{\kappa}=\vec{n}\hbar\omega/c$ is transferred to the center of mass of the electron-ion pair that is formed in the atomic photoionization. As a result, the pair as a whole begins to move in the gas with the velocity $\vec{v}_T+\hbar\vec{\kappa}/(M+m)$. Here m and M are the electron and ion masses, respectively. In the coordinate system connected with the center of mass of this pair, the photoelectron and the ion move in opposite directions with velocities $\vec{v}$ and $\vec{V}$, such that $m\vec{v}=-M\vec{V}$. The velocities of the electron $\vec{v}_e$ and the ion $\vec{v}_i$ in the gas after the photon absorption are respectively equal to

$$\vec{v}_e=\vec{v}_T+\frac{\hbar\vec{\kappa}}{(M+m)}+\vec{v},$$
$$\vec{v}_i=\vec{v}_T+\frac{\hbar\vec{\kappa}}{(M+m)}-\frac{m}{M}\vec{v}. \tag{1}$$

The angular distribution of photoelectrons is determined by the differential cross section for the photoeffect, $d\sigma_{ph}(\omega)/d\Omega$. Normalized to unity, the distribution function of electrons has the form $\phi(\Omega)=\sigma_{ph}^{-1}(\omega)d\sigma_{ph}(\omega)/d\Omega$. Here $d\Omega$ is the solid angle in which the vector $\vec{v}$ is located after atomic photoionization and $\sigma_{ph}(\omega)$ is the total cross section of the photoeffect. Using Eq. (1) we can obtain the mean values of the projections of the photoelectron $\vec{p}_e$ and ion $\vec{p}_i$ momenta on $\vec{n}$; they are, respectively,

$$\langle\vec{p}_e\cdot\vec{n}\rangle_{T,\Omega}=\int f(\vec{v}_T)d\vec{v}_T\int\phi(\Omega)(m\vec{v}_e\cdot\vec{n})d\Omega$$
$$=\frac{m}{M+m}\hbar\kappa+\frac{m}{\sigma_{ph}(\omega)}\int\frac{d\sigma_{ph}(\omega)}{d\Omega}(\vec{v}\cdot\vec{n})d\Omega,$$
$$\langle\vec{p}_i\cdot\vec{n}\rangle_{T,\Omega}=\int f(\vec{v}_T)d\vec{v}_T\int\phi(\Omega)(M\vec{v}_i\cdot\vec{n})d\Omega$$
$$=\frac{M}{M+m}\hbar\kappa-\frac{m}{\sigma_{ph}(\omega)}\int\frac{d\sigma_{ph}(\omega)}{d\Omega}(\vec{v}\cdot\vec{n})d\Omega. \tag{2}$$

These projections of the momenta are nonzero. Hence the process of atomic photoionization generates electrons and ions moving preferably along or opposite to the photon momentum. The sum of the mean electron and ion momenta, Eq. (2), is equal to the photon momentum for any frequency, as it should be, including that at the photoeffect threshold, rather than zero as implied by Eq. (3) of Ref. [8]. According to Eq. (2), the mean momentum acquired by the ion during photoionization is the sum of two momenta: that connected with the motion of the ion-electron pair as a whole and the momentum due to their relative motion. If the photoelectron is ejected preferably along $\vec{n}$, the integrals in Eq. (2) are positive and the momentum acquired by the ion is less than the photon momentum. When the photoelectron is emitted preferably opposite to the direction of the photon propagation then the ion momentum is greater than $\hbar\vec{\kappa}$.

In multielectron atoms, unlike the hydrogen atom, the direction of preferred ejection of the photoelectron depends on radiation frequency. Therefore, the light pressure force due to atomic photoionization can change its sign for some photon energies, which is impossible in the case of ordinary light pressure due to photon scattering by an atom. In the latter case the light pressure force is always directed along the Poynting's vector. In the dipole approximation the momentum of the relative motion in the ion-electron pair is equal to zero and for the mean momentum acquired by the target atom we have the value $\langle\vec{p}_i\cdot\vec{n}\rangle_{T,\Omega}\approx\hbar\kappa$. The same momentum is transferred to an atom in the process of photon scattering. Let us compare the forces of light pressure acting on the atomic particle in these two processes. If W ($\mathrm{cm}^{-2}\,\mathrm{s}^{-1}$) is the density of the photon flux, then these forces are, respectively, $f_{ph}(\omega)\approx W\sigma_{ph}(\omega)\hbar\kappa$ and $f_{scat}(\omega)\approx W\sigma_{scat}(\omega)\hbar\kappa$, where $\sigma_{scat}(\omega)$ is the cross section for photon scattering. Near the photoionization thresholds of outer atomic subshells typical values of σ_{ph} are of the order of $30\alpha a_0^2$ and of σ_{scat} are of the order of $10\alpha^4 a_0^2$ [18], where α is the fine structure constant and a_0 is the Bohr radius. The cross sections $\sigma_{ph}(\omega)$ and $\sigma_{scat}(\omega)$ maintain comparable values to those quoted above up to several hundreds of eV photon energies. It is seen from these expressions that, for photon energies in the range from outer atomic shell ionization thresholds up to several hundreds of eV, the force $f_{ph}(\omega)$ is several orders of magnitude greater than $f_{scat}(\omega)$. Consequently, in this frequency range the dominant mechanism of light pressure on matter in the gas phase is atomic photoionization.

III. THE KINETICS OF THE GENERATION OF ELECTRON-ION PAIRS AND THE DRAG CURRENT IN ATOMIC GASES

The directed motion of charged particles (photoelectrons and positive ions) results in the formation of macroscopic currents flowing through a medium along or opposite to the direction of electromagnetic wave propagation. Let us derive the formulas for these currents. Suppose that n_a is the density of the gas atoms. Then the rate of generating electron-ion pairs in a unit volume is equal to $W\sigma_{ph}(\omega)n_a$. In the stationary regime the rate of generating the pairs with a nonisotropic momentum distribution is equal to the rate of transformation of this distribution into the Maxwell one that is due to the collisions of electrons and ions with the gas atoms. Thus, one has

$$\frac{n_e}{\tau_e(v_e)}=\frac{n_i}{\tau_i(v_i)}=W\sigma_{ph}(\omega)n_a. \tag{3}$$

Here n_e and n_i are the mean equilibrium densities of electrons and ions with the momentum projections given by Eq. (2), and $\tau_e(v_e)$ and $\tau_i(v_i)$ are the relaxation times of the electrons and ions in the gas, connected to the transport cross sections of these particles $\sigma_{ea}(v_e)$ and $\sigma_{ia}(v_i)$ by the following relations:

$$\tau_e(v_e)=[n_a\langle v_e\sigma_{ea}(v_e)\rangle_T]^{-1},$$
$$\tau_i(v_i)=[n_a\langle v_i\sigma_{ia}(v_i)\rangle_T]^{-1}. \tag{4}$$

The currents due to dragging of the electrons and ions by photons are equal to, respectively,

$$j_e(\omega)=-|e|n_e\langle\vec{v}_e\cdot\vec{n}\rangle_{T,\Omega},$$
$$j_i(\omega)=Z|e|n_i\langle\vec{v}_i\cdot\vec{n}\rangle_{T,\Omega}. \tag{5}$$

Here e is the electron charge and Z is the charge of the positive ions created in the atomic photoionization. For the ionization of outer atomic shells when the Auger decay of a vacancy is impossible, the ion charge is equal to unity. In the case of a hole in inner or intermediate atomic shells, the Auger cascade or other mechanism can result in the appearance of ions with charge greater than unity. The total drag current flowing through the medium along the direction $\vec{n}$ is the sum of the electron and ion currents,

$$j(\omega)=j_e(\omega)+j_i(\omega)$$
$$=j_0\left\{[Z\tau_i(v_i)-\tau_e(v_e)]n_a\sigma_{ph}(\omega)\frac{\hbar\kappa}{M+m}-\left[\frac{m}{M}Z\tau_i(v_i)+\tau_e(v_e)\right]n_a\int\frac{d\sigma_{ph}(\omega)}{d\Omega}(\vec{v}\cdot\vec{n})d\Omega\right\}, \tag{6}$$

where $j_0=|e|W$. The total drag current given by Eq. (6) can be represented as a current due to the motion of the center of mass of the electron-ion pair,

$$j_C(\omega)=j_0[Z\tau_i(v_i)-\tau_e(v_e)]n_a\sigma_{ph}(\omega)\frac{\hbar\kappa}{M+m}, \tag{7}$$

and a current connected with the relative motion of the fragments of this pair,

$$j_R(\omega)=-j_0\left[\frac{m}{M}Z\tau_i(v_i)+\tau_e(v_e)\right]n_a\int\frac{d\sigma_{ph}(\omega)}{d\Omega}(\vec{v}\cdot\vec{n})d\Omega. \tag{8}$$

Let us compare the contributions of currents Eq. (7) and Eq. (8) to the drag current $j(\omega)$. According to Eq. (8), the drag current $j_R(\omega)$ is equal to zero in the dipole approximation. The situation becomes different if in calculations of the angular distribution of the photoelectrons the interference between the electric dipole ($E1$) and electric quadrupole ($E2$) photoionization amplitudes is taken into account. In this case, the integral in Eq. (8) is nonzero and for the outer shells of atoms its value is of order $v\sigma_{ph}(\omega)\kappa a$, where a is the radius of the atomic shell [19]. By comparing the currents given by Eqs. (7) and (8) we obtain $j_R(\omega)\gg j_C(\omega)$. The transport cross sections of ion-atom scattering $\sigma_{ia}(v_i)$ are much greater than $\sigma_{ea}(v_e)$. Hence in Eq. (8) we can neglect $Z\tau_i(v_i)$ as compared with $\tau_e(v_e)$. Moreover, the relaxation time $\tau_e(v_e)$ [Eq. (4)] can be written in the form $\tau_e(v_e)\cong[n_av_e\sigma_{ea}(v_e)]^{-1}$ for photoelectron velocities v_e that are greater than the thermal ones. Taking into account these points, we obtain from Eqs. (4) and (8) the following expression for the density of the drag current in atomic gases [11,12]:

$$j(\omega)=-j_0\frac{1}{\sigma_{ea}(\epsilon)}\int\frac{d\sigma_{ph}(\omega)}{d\Omega}\cos\vartheta d\Omega, \tag{9}$$

where $\epsilon=mv^2/2\approx\hbar\omega-I_{nl}$. Multiplying both sides of Eq. (9) by the cross section of the photon beam S we have for the total drag current through the surface S

$$J(\omega)=-J_0\frac{1}{\sigma_{ea}(\epsilon)}\int\frac{d\sigma_{ph}(\omega)}{d\Omega}\cos\vartheta d\Omega, \tag{10}$$

where $J_0=|e|WS$.

According to Eq. (10) the drag current does not depend on the density of gas atoms. This is connected with the fact that the effective electromotive force $f_{ph}(\omega)n_a$ is proportional to the concentration of atoms, while the electron mobility $|e|\tau_e(v_e)/m$ is inversely proportional to n_a, and according to Ohm's law [Eq. (5)], the current in the gas is the product of these magnitudes. The situation is quite different in atomic gas mixtures. There, the drag current depends on atomic gas densities. Let us consider the photoionization of type A atoms which are in the atmosphere of a buffer gas of type B. Let us suppose that the buffer gas density n_a^B is greater than n_a^A. If the frequency of radiation is such that the photoionization of type B atoms is impossible ($\hbar\omega<I_{nl}^B$), or their cross sections are small compared with $\sigma_{ph}^A(\omega)$, then the rate of generating the photoelectrons, the right side of Eq. (3), is $W\sigma_{ph}^A(\omega)n_a^A$. Since $n_a^B\gg n_a^A$, the relaxation time in this binary gas mixture is determined by the electron scattering by B atoms; therefore $\tau_e(v_e)\cong[n_a^Bv_e\sigma_{ea}(v_e)]^{-1}$. Now, using Eqs. (3), (4), and (8) we obtain the following modified expression for the drag current in the gas mixture [20]:

$$J(\omega)=-J_0\frac{n_a^A}{n_a^B}\frac{1}{\sigma_{ea}^B(\epsilon)}\int\frac{d\sigma_{ph}^A(\omega)}{d\Omega}\cos\vartheta d\Omega. \tag{11}$$

Here $d\sigma_{ph}^A(\omega)/d\Omega$ and $\sigma_{ea}^B(\epsilon)$ are, respectively, the cross sections for photoionization of A atoms and elastic scattering of electrons by B atoms.

IV. RELATIONSHIP BETWEEN DRAG CURRENTS AND NONDIPOLE ASYMMETRY PARAMETERS

According to Eqs. (10) and (11), the drag current is defined by the angular differential cross section for photoionization. In the case of unpolarized radiation it has the form [11,12,21]

$$\frac{d\sigma_{ph}(\omega)}{d\Omega}=\frac{\sigma_{ph}(\omega)}{4\pi}\Big[1-\frac{\beta(\omega)}{2}P_2(\cos\vartheta)+\kappa\gamma(\omega)P_1(\cos\vartheta)+\kappa\eta(\omega)P_3(\cos\vartheta)\Big], \tag{12}$$

where the coefficients of the Legendre polynomials $P_l(\cos\vartheta)$ are functions of the dipole and quadrupole matrix elements and phase shifts of the photoelectron wave functions. For linearly polarized radiation, instead of Eq. (12) we have [22]

$$\frac{d\sigma_{ph}(\omega)}{d\Omega}=\frac{\sigma_{ph}(\omega)}{4\pi}\{1+\beta(\omega)P_2(\cos\Theta)+[\delta^C(\omega)+\gamma^C(\omega)\cos^2\Theta]\sin\Theta\cos\Phi\}, \tag{13}$$

where Θ is the polar angle of the photoelectron velocity with respect to the photon polarization vector $\vec{e}$, and Φ is the azimuthal angle defined by the projection of $\vec{v}$ in the plane perpendicular to $\vec{e}$ and containing the photon propagation vector $\vec{\kappa}$. The nondipole asymmetry parameters $\gamma^C(\omega)$ and $\delta^C(\omega)$ are connected to $\gamma(\omega)$ and $\eta(\omega)$ by the relations

$$\frac{\gamma^C}{5}+\delta^C=\kappa\gamma \quad \text{and} \quad \frac{\gamma^C}{5}=-\kappa\eta. \tag{14}$$

Substituting Eqs. (12) and (13) in Eq. (10) and integrating over the solid angle we obtain the following formula for the drag current:

$$J(\omega)=-J_0\frac{1}{3}\frac{\sigma_{ph}(\omega)}{\sigma_{ea}(\epsilon)}\kappa\gamma(\omega). \tag{15}$$

Thus, the drag current is directly connected with one of the nondipole asymmetry parameters $\gamma(\omega)$ in the expansion of $d\sigma_{ph}(\omega)/d\Omega$ in terms of the Legendre polynomials [Eq. (12)]. Therefore, the experimental investigation of the drag current can be used for precision measurement of this parameter. This is very important for small photoelectron energies, where it is difficult to apply the traditional experimental schemes to measure the differential photoionization cross section.

V. NONDIPOLE ASYMMETRY PARAMETERS

We first calculate the nondipole asymmetry parameters for some subshells of Ne and Ar atoms, and then use them to obtain the drag currents in these gases. For arbitrary orbital angular momentum l of an atomic subshell the nondipole asymmetry parameters are defined by rather complicated formulas presented in Refs. [11,12,22]. These formulas simplify considerably for the atomic s subshells. For this case $\delta^C=0$ and we have a very simple expression [19],

$$\gamma^C(\omega)=5\kappa\gamma(\omega)=6\kappa\frac{Q_2}{D_1}\cos(\delta_2-\delta_1), \tag{16}$$

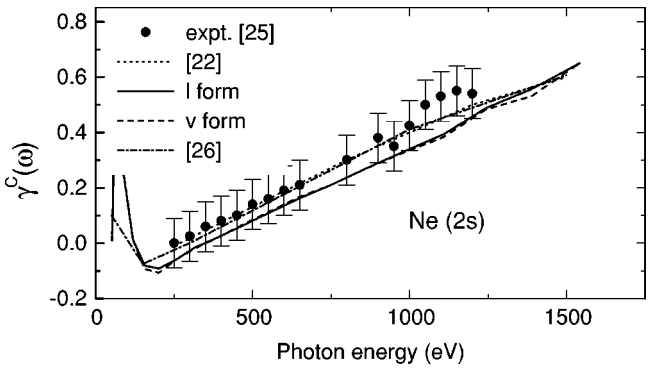

FIG. 1. Nondipole asymmetry parameter γ^C, as a function of photon energy, for photoionization out of the $2s$ subshell of Ne.

where $\delta_{1,2}(v)$ are the phases of photoelectron scattering in the atomic field with a vacancy in the ns subshell. Dipole D_1 and quadrupole Q_2 matrix elements are determined by the integrals

$$D_1=\int P_{ns}(r)rP_{\epsilon,1}dr \quad \text{and} \quad Q_2=\tfrac{1}{2}\int P_{ns}(r)r^2P_{\epsilon,2}dr. \tag{17}$$

$P_{ns}(r)/r$ and $P_{\epsilon,1,2}(r)/r$ are the radial parts of the electron wave functions of the atomic ns subshell and the continuous spectrum, respectively. For multielectron atoms the one-electron approximation often describes atomic photoprocesses unsatisfactorily . Much more accurate results can be obtained if the multielectron correlations are taken into account as is done in the random phase approximation with exchange (RPAE) [23]. In the RPAE, $\gamma(\omega)$ is determined by Eq. (16) in which the following substitutions are made [11,12]:

$$\frac{Q_2}{D_1}\cos(\delta_2-\delta_1)\to\{(D_1'Q_2'+D_1''Q_2'')\cos(\delta_1-\delta_2)+(D_1'Q_2''-D_1''Q_2')\sin(\delta_1-\delta_2)\}\times[D_1'^2+D_1''^2]^{-1}, \tag{18}$$

where D_1', Q_2' and D_1'', Q_2'' are the real and imaginary parts of the matrix elements, respectively.

The calculations of the dipole and quadrupole matrix elements and phase shifts of the photoelectron were performed using standard codes [24]. Figure 1 compares our calculated $\gamma^C(\omega)$ for the $2s$ subshell of the Ne atom with other results. The calculations of the matrix elements were performed in the so-called length and velocity forms. They had to give the same results not only for precise initial and final state wave functions, but also in the framework of the RPAE as well. The agreement between the two forms demonstrates the high quality of our calculations. The curves obtained are in good agreement with the available experimental data [25]. Also included in Fig. 1 are calculations using the Herman-Skillman potential [22] and those of a recent paper [26]. In Fig. 2 our calculated linear combination of parameters $\gamma^C+3\delta^C$ for the $2p$ subshell of the Ne atom are presented

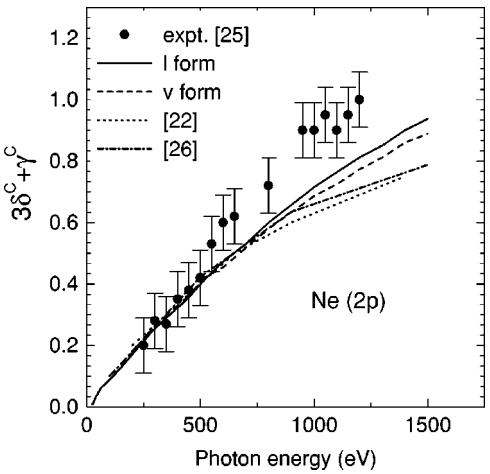

FIG. 2. Nondipole asymmetry parameter $3\delta^C+\gamma^C$, as a function of photon energy, for photoionization out of the $2p$ subshell of Ne.

in both the length and velocity forms. Measurement [25] and calculation [22,26] results are also included for comparison.

The photoionization of the outer atomic subshells, as a rule, has a collective character. However, for the Ne atom the role of multielectron correlations is not important, so that similar curves are also obtained in the one-electron Hartree-Fock approximation. In Fig. 3 our calculated (in the length form) parameter $\gamma^C(\omega)$ for the $1s$ subshell of the Ar atom is compared with measurement [16] and calculation [22]. It is well known that the photoionization of deep atomic shells is accompanied by rearrangement of the residual positive ion. The rearrangement consists of the variation of the wave functions of all atomic electrons due to the creation of a deep vacancy and its subsequent decay. This leads to changes in the mean field in which the knocked out photoelectron moves, which affects the probability of electron ejection [27]. Therefore, in calculations of the parameter $\gamma^C(\omega)$ for the deep $1s$ hole we take into account its Auger decay. As

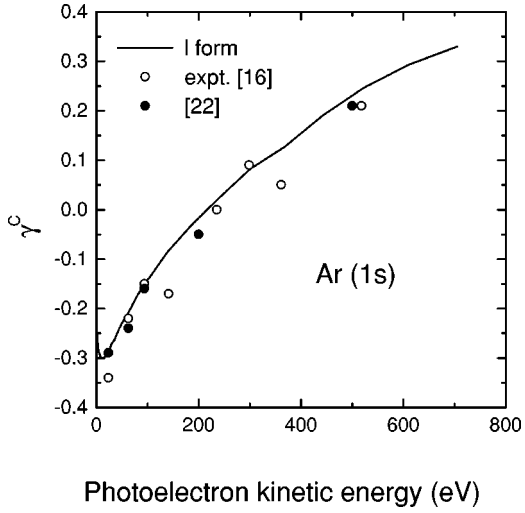

FIG. 3. Nondipole asymmetry parameter γ^C, as a function of photon energy, for photoionization out of the $1s$ subshell of Ar.

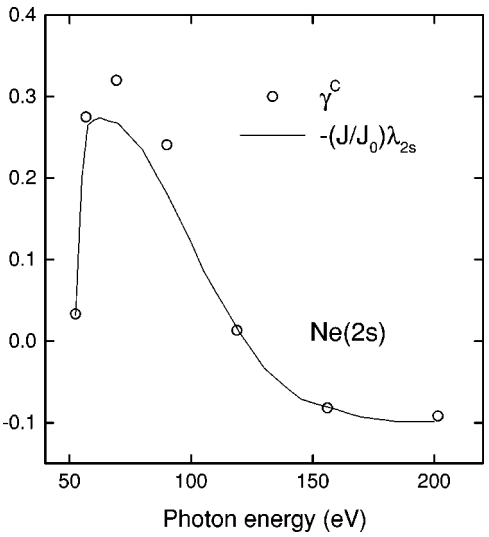

FIG. 4. Partial drag current for Ne $2s$ photoionization.

seen from Fig. 3 our calculated values are in good agreement with the experimental data [16].

VI. DRAG CURRENTS

In Figs. 4 and 5 the ω dependence of the partial drag currents in Ne gas are presented. The ratio of the cross sections $\sigma_{ph}(\omega)/\sigma_{ea}(\epsilon)$ was calculated for frequencies ω near the $2s$ and $2p$ thresholds of the Ne atom ($I_{2s}=52.51$ eV, $I_{2p}=23.13$ eV). Here the photoionization cross sections obtained in [28,29] were used. The elastic scattering cross sections for the Ne atom in the energy range 0–200 eV are taken from [30]. In both figures the nondipole asymmetry parameters $\gamma^C(\omega)$ for the $2s$ subshell and $\kappa\gamma(\omega)$ for the $2p$ subshell are represented by open circles; the constants λ_{2s} and λ_{2p} in these figures are equal to 10^4 and 0.7×10^3, respectively.

As seen from these figures, the qualitative behavior of the ω dependence of the asymmetry parameters and partial drag

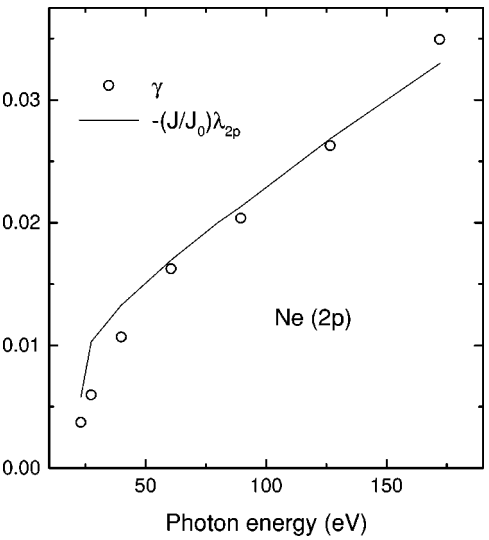

FIG. 5. Partial drag current for Ne $2p$ photoionization.

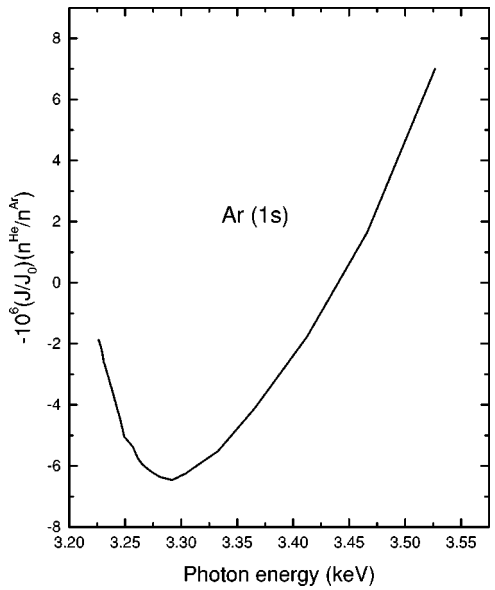

FIG. 6. Frequency dependence of the drag current from the Ar $1s$ subshell.

currents is similar. In particular, the positions of the zeros on the energy scale for the current j_{2s} and the parameter $\gamma^C(\omega)$ coincide. Near the thresholds of photoionization the ratio $\sigma_{ph}(\omega)/\sigma_{ea}(\epsilon)$ is in most cases a slowly varying function, almost a constant. Therefore, the experimental investigation of the drag current would permit the frequency dependence $\gamma^C(\omega)$ and $\gamma(\omega)$ to be reconstructed near the photoionization thresholds. The measurement of the drag current as a method of investigating the nondipole asymmetry parameter is particularly effective, provided that the ratio $\sigma_{ph}(\omega)/\sigma_{ea}(\epsilon)$ varies slowly with the radiation frequency. In cases where this ratio varies rapidly, the frequency dependence of the $\gamma^C(\omega)$ parameter can be masked. The noble gases Ar, Kr, and Xe, owing to the presence of the Ramsauer minimum, provide a good example. In such a case it is convenient to study the drag effect in mixtures of atomic gases by choosing the density of the gas under study n_a^A and that of the buffer one n_a^B so that the photoionization of the A atoms is accompanied by relaxation due to the elastic scattering of electrons by the B atoms. In a binary mixture of Ar and He gases near the ionization threshold of the $1s$ subshell of Ar ($I_{1s}=3226.19$ eV), the photoionization cross section of He atoms can be neglected as compared to $\sigma_{ph}^{\mathrm{Ar}}(\omega)$. Therefore, for $n_a^{\mathrm{He}} \gg n_a^{\mathrm{Ar}}$ we have, according to Eq. (11),

$$J(\omega) = -J_0 \frac{1}{15} \frac{n_a^{\mathrm{Ar}}}{n_a^{\mathrm{He}}} \frac{\sigma_{ph}^{\mathrm{Ar}}(\omega)}{\sigma_{ea}^{\mathrm{He}}(\epsilon)} \gamma_{\mathrm{Ar}}^C(\omega). \tag{19}$$

Figure 6 shows the calculated drag current in this mixture of gases as a function of ω. The cross section $\sigma_{ea}(\epsilon)$ was taken from [31] and the photoionization cross section $\sigma_{ph}^{\mathrm{Ar}}(\omega)$ near the $1s$ subshell threshold came from [27,32]. We note that in this case also the frequency dependence of $\gamma_{\mathrm{Ar}}^C(\omega)$ (not shown in Fig. 6) is similar to that of $J(\omega)$. This makes it possible to reconstruct the behavior of the nondipole asymmetry parameter near the threshold of the deep $1s$ subshell of the Ar atom from the data on the current. The absolute value of the drag current density is determined by the photon flux W. The total flux of synchrotron radiation reaches the value of $WS \approx 10^{12}$ s^{-1}. For this photon flux, the scaling parameter in our formulas is $J_0 = 1.6 \times 10^{-7}$ A. According to Figs. 4–6, the partial drag currents in Ne gas at photon energy $\hbar\omega = 100$ eV have values $J_{2s} = 0.2 \times 10^{-11}$ A and $J_{2p} = 0.5 \times 10^{-11}$ A for the $2s$ and $2p$ subshells, respectively. The drag current from the $1s$ subshell of Ar gas has a maximum absolute value $J_{1s} = 1.0 \times 10^{-13}$ A for $n_a^{\mathrm{He}}/n_a^{\mathrm{Ar}} = 10$. These currents are readily available for experimental observation. Free electron lasers have now been developed that can provide photon fluxes many orders of magnitude higher than these values. They are therefore suitable for studying drag currents.

The experimental study of drag currents is essentially a method of direct measurement of one of the nondipole parameters, namely, the parameter $\gamma(\omega)$ in the differential cross section in Eq. (12). This method is completely different from the traditional ones used for measuring the angular distribution of photoelectrons. In those methods the measurement of nondipole parameters is performed by recording the number of photoelectrons ejected at a definite angle ϑ relative to $\vec{\kappa}$. The differential cross section, Eq. (12), at this angle $d\sigma_{ph}(\omega, \vartheta_0) \sim [1 - \beta P_2(\cos\vartheta_0)/2 + \kappa\gamma P_1(\cos\vartheta_0) + \kappa\eta P_3(\cos\vartheta_0)] d\Omega$ is defined by the number of electrons emitted within the elementary solid angle $d\Omega$. Consequently, in this type of experiment one obtains information only on a linear combination of the coefficients in the expansion of the differential cross section in Eq. (12) in terms of the Legendre polynomials. In contrast, the drag current is defined only by $\gamma(\omega)$ [Eq. (15)]; this gives the possibility of directly measuring this parameter.

VII. POSSIBLE SCHEME OF THE EXPERIMENT

Let us imagine the following experimental scheme designed to observe and study the effect of photoelectron dragging. The ionizing radiation traverses the atomic gas, placed between two metallic grids playing the role of electrodes, with the help of which the current can be measured. Since the macroscopic drag current $J(\omega)$ is formed by the processes of photoelectron collisions with gas atoms, this experiment could be performed at normal gas densities. This is contrary to the traditional experimental scheme where the collisions of photoelectrons with the atoms of the target gas are the obstacle for precision measurements of the differential cross section.

For the elimination from the measured current of the electrons directly ionized from the grids another experimental scheme is suggested, similar to that for the measurement of the Hall effect. Imagine that the electrodes are placed parallel to the photon flux. The ionizing gas is placed in a magnetic field oriented perpendicular to the photon flux direction and parallel to the electrode surfaces. Under the action of this field, a Hall current $\vec{J}_H(\omega)$ is created in the gas. The total current through the gas target in the magnetic field is given by [33]

$$\vec{J}_H(\omega)=\frac{1}{1+\xi^2H^2}\vec{J}(\omega)+\frac{\xi}{1+\xi^2H^2}[\vec{J}(\omega)\times\vec{H}], \quad (20)$$

where $\xi=e\mu\tau_e/mc$, $J(\omega)$ is the drag current without the magnetic field given by Eq. (10), and μ is the gas magnetic permittivity. For the magnetic induction $\mu H=10^4$ G and the atomic concentration $n_a\approx10^{19}$ cm^{-3}, the product ξH is of the order of 10^{-2}. Under these conditions the current between the electrodes [the second term in Eq. (20)] is about $J_\perp(\omega)\approx10^{-2}J(\omega)$. As mentioned above, in this experimental scheme the electrode surfaces are not affected directly by x-ray radiation. Therefore, the problem of taking into account the surface photoeffect is eliminated.

VIII. CONCLUSIONS

In this paper we have shown that the photoionization process is the main mechanism of transferring electromagnetic radiation momentum to atomic particles in gaseous media for photon energies from the photoionization thresholds up to several hundreds of eV. Also, in the photoionization process, the photon momentum is absorbed by the ion produced and the ejected electron. The absorption of photons results in the generation of electron-ion pairs, and the directed motion of these charges in media creates macroscopic currents. The redistribution of photon momentum between the photoelectron and ion has been analyzed. We have also shown that the mean momenta of the directed motion of these particles are defined by nondipole parameters of the photoeffect. The values and directions of these momenta depend on the atomic structure and dynamics of the interaction of the atomic particles with electromagnetic radiation, namely, on the dipole and quadrupole amplitudes, as well as on phase shifts of the photoelectron wave functions. General expressions for the current generated through dragging of electrons by photons have been derived. The drag currents in Ne gas and a binary gaseous Ar-He mixture have been calculated. The possibility of using this effect to study nondipole corrections in the angular distributions of photoelectrons near the photoionization threshold has been analyzed.

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[1] L. Spitzer, Jr., *Physical Processes in Interstellar Medium* (John Wiley & Sons, New York, 1978).
[2] A. Ashkin, Sci. Am. **226**, 63 (1972).
[3] I. Kuščer, L. J. F. Hermans, P. L. Chapovsky, J. J. Beenakker, and G. J. van der Meer, J. Phys. B **26**, 2837 (1993).
[4] F. Kh. Gel'mukhanov and A. M. Shalagin, Pis'ma Zh. Éksp. Teor. Fiz. **29**, 773 (1979) [JETP Lett. **29**, 711 (1979)].
[5] E. R. Eliel, Adv. At., Mol., Opt. Phys. **30**, 199 (1992).
[6] G. Michaud, Astrophys. J. **160**, 641 (1970).
[7] G. Michaud, Phys. Scr., T **47**, 143 (1993).
[8] M. J. Seaton, J. Phys. B **28**, 3185 (1995).
[9] A. Sommerfeld, *Atombau und Spectrallinien: Wellenmechanische Ergänzungsband* (Vieweg, Braunschweig, 1928).
[10] A. Sommerfeld, *Wave Mechanics* (Methuen, London 1930).
[11] M. Ya. Amusia, P. U. Arifov, A. S. Baltenkov, A. A. Grinberg, and S. G. Shapiro, Phys. Lett. **47A**, 66 (1974).
[12] M. Ya. Amusia, A. S. Baltenkov, A. A. Grinberg, and S. G. Shapiro, Zh. Éksp. Teor. Fiz. **68**, 28 (1975) [Sov. Phys. JETP **41**, 14 (1975)].
[13] G. Rakavy and A. Ron, Phys. Rev. **159**, 50 (1967).
[14] M. O. Krause, Phys. Rev. **177**, 151 (1969).
[15] F. Wuilleumier and M. O. Krause, Phys. Rev. A **10**, 242 (1974).
[16] B. Krassig, M. Jung, D. S. Gemmell, E. P. Kanter, T. LeBrun, S. H. Southworth, and L. Young, Phys. Rev. Lett. **75**, 4736 (1995).
[17] O. Hemmers, S. B. Whitfield, P. Glans, H. Wang, D. W. Lindle, R. Wehlitz, and I. A. Sellin, Rev. Sci. Instrum. **69**, 3809 (1998).
[18] V. B. Berestetsky, E. M. Lifchitz, and L. P. Pitayevski, *Quantum Electrodynamics* (Nauka, Moscow, 1980).
[19] M. Ya. Amusia, A. S. Baltenkov, Z. Felfli, and A. Msezane, Phys. Rev. A **59**, R2544 (1999).
[20] M. Ya. Amusia and A. S. Baltenkov (unpublished).
[21] A. Bechler and R. H. Pratt, Phys. Rev. A **39**, 1774 (1989).
[22] J. W. Cooper, Phys. Rev. A **47**, 1841 (1993).
[23] M. Ya. Amusia, *Atomic Photoeffect* (Plenum Press, New York, 1990).
[24] M. Ya. Amusia and L. V. Chernysheva, *Computation of Atomic Processes* (IOP Publishing Ltd., Bristol, 1997).
[25] O. Hemmers, G. Fischer, P. Glans, D. L. Hansen, H. Wang, S. B. Whitfield, R. Wehlitz, J. C. Levin, I. A. Sellin, R. C. C. Perera, E. W. B. Dias, H. S. Chakraborty, P. C. Deshmukh, S. T. Manson, and D. W. Lindle, J. Phys. B **30**, L727 (1997).
[26] W. R. Johnson, A. Derevianko, K. T. Cheng, V. K. Dolmatov, and S. T. Manson, Phys. Rev. A **59**, 3609 (1999).
[27] M. Ya. Amusia, V. K. Ivanov, and V. A. Kupchenko, J. Phys. B **14**, L667 (1981).
[28] M. Ya. Amusia, V. K. Ivanov, N. A. Cherepkov, and L. V. Chernysheva, Phys. Lett. **40A**, 361 (1972).
[29] D. K. Kennedy and S. T. Manson, Phys. Rev. A **5**, 227 (1972).
[30] W. C. Fon and K. A. Berrington, J. Phys. B **14**, 323 (1981).
[31] F. J. de Heer, R. W. Wagenaar, H. J. Blaauw, and A. Tip, J. Phys. B **9**, L269 (1976).
[32] H. W. Schnopper, Phys. Rev. **131**, 2558 (1963).
[33] T. G. Cowling, *Magnetohydrodynamics* (Interscience, New York, 1957).