Drag current in multiphoton ionization of atomic gases

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(Submitted 22 June 1982)
Zh. Eksp. Teor. Fiz. **83**, 2027–2034 (December 1982)

The asymmetry of the photoelectron-emission directions in multiphoton ionization of atoms, required by the law of momentum conservation in the "atom + ionizing radiation" system, is investigated theoretically. The resonance and polarization features of the effect are studied within the framework of the one-electron model-potential technique. The density of the photoelectron-drag current in two-photon ionization of hydrogen is calculated. Estimates are made for alkali-metal vapors.

PACS numbers: 32.80.Fb, 32.80.Kf, 51.70. + f

1. Free electrons produced in photoionization should be asymmetrically distributed relative to the propagation direction of the ionizing radiation. This fact follows already from the momentum conservation law for the "photon + atom" system: the momentum of the absorbed photon is acquired by the electron and by the atomic residue. In the nonrelativistic dipole approximation, however, which is customarily employed in the theory of photoionization of atoms and molecules by radiation in the optical or soft x-ray bands, calculation yields a photoelectron angular distribution that is symmetrical with respect to the wave propagation direction. The asymmetric part of the distribution can be obtained only when account is taken of the spatial inhomogeneity of the electromagnetic wave, and also of the magnetic interaction of the ionized particle with the field.

The asymmetry in the distribution of the electrons in photoionization of hydrogen atoms was first calculated by Sommerfeld (see Ref. 1), who showed that the term following the dipole term in the long-wave approximation determines the preferred emission of the photoelectron along the photon propagation direction. Interest in this phenomenon has increased recently because of experimental observation of currents (drag currents) which are produced in semiconductors under the influence of a rather strong electromagnetic radiation and are due to momentum transfer from the field to the free carriers.[2]

These experiments have stimulated further theoretical investigations of a similar phenomenon in an atomic gas. Calculations were performed of the asymmetry of the angular distribution of photoelectrons in ionization of inert-gas atoms, and a very important feature of the effect was observed, namely, in complex multielectron atoms the direction of the predominant emission of the photoelectrons can be opposite to the photon-momentum direction.[3] This effect is apparently due to the interaction of the outgoing electron with the atomic core, which acquires the excess momentum and by the same token ensures its conservation. Thus, the direction of the drag current depends on the structure of the gas atoms and on the frequency of the incident wave, and can be directed either along or against the wave vector of the incident radiation.

Exactly the same effect can apparently be observed also in other multielectron atoms, if calculations similar to those in Ref. 3 are performed using various methods for determining the matrix elements of the bound-free transitions, particularly the quantum-defect method[4] for alkali-metal atoms or the model-potential method[5] for alkali-metals and inert gases. The solution of this problem is of interest and an experimental verification of these calculations might serve as a check on the applicability of the particular method, but the lact of sufficiently intense vacuum-ultraviolet or x-ray sources to ionize these atoms makes the experiment difficult, since the current densities even at the highest presently available intensities hardly reach the experimentally observed value ($j = 10^{-12}$–10^{-11} A/cm^2, Ref. 3).

It is therefore more important at present to study the asymmetry of the distribution of the photoelectrons in multiphoton ionization by intense sources of optical radiation—pulsed lasers. Multiphoton ionization of rarefied atomic media by intense light fluxes has been thoroughly investigated by now both theoretically[5] and experimentally.[6] However, the asymmetry in the direction of emission of the photoelectrons in multiphoton ionization has not yet been discussed in the literature.[5,6,7] Yet, experimental study of this phenomenon may turn out to be much simpler than in one-photon ionization, owing to the high radiation intensity of modern lasers, which can yield, for example in ionization of alkali atoms, photoelectron emission probabilities $W_N = 10^8 - 10^{12}$ sec^{-1} (Ref. 6), which exceeds by many orders of magnitude the probability of one-photon ionization by existing sources ($W_1 \lesssim 10^{-6}$ sec^{-1}). In addition, in multiphoton ionization it is possible to increase considerably the probability on account of the intermediate resonances with the natural frequencies of the atoms.

In the present paper we investigate theoretically the asymmetries in the emission directions of photoelectrons in multiphoton ionization of atoms. Within the framework of the one-electron method of the model potential, developed earlier for the calculation of the cross sections for multiphoton ionization,[8] we investigate the resonance and polarization features of this effect, and obtain also estimates for the density of the drag current produced in the medium in multiphoton ionization.

2. The probability of N-photon ionization of an atom by monochromatic radiation of frequency ω, with an electric field intensity vector

 0038-5646/82/121174-05$04.00

$$\vec{\mathcal{F}}(\mathbf{r}, t)=2F \operatorname{Re}\{\mathbf{e}\exp[i(\mathbf{kr}-\omega t)]\}, \quad (1)$$

(e is the unit vector of the polarization accompanied by emission of a photoelectron with momentum $\mathbf{k}_f$ directed into the solid angle $d\Omega_f$) is of the form (the atomic system of units is used)

$$dW_N=2\pi|M_{fi}|^2 d\Omega_f, \quad (2)$$

where M_{fi} is a matrix element of order N, which defines an N-photon transition from an initial state $|i\rangle$ with energy $E_i<0$ into a final state $|f\rangle$ with energy

$$E_f=E_i+N\omega>0.$$

Retaining the terms corresponding to $E1$, $M1$, and $E2$ interactions in the multipole expansion of the operator of the interaction of the electromagnetic wave (1) with the atom

$$\hat{V}(\mathbf{r}, t)=\hat{V}^{\omega}(\mathbf{r})e^{-i\omega t}+\hat{V}^{\omega+}(\mathbf{r})e^{i\omega t},$$

we represent $\hat{V}^{\omega}(\mathbf{r})$ in the form

$$\hat{V}^{\omega}(\mathbf{r})=\hat{V}_d+\hat{V}_\mu+\hat{V}_Q, \quad (3)$$

where

$$\hat{V}_d=-F(\hat{\mathbf{d}}\mathbf{e})$$

is the electric-dipole,

$$\hat{V}_\mu=i2^{1/2}F(\hat{\boldsymbol{\mu}}\{\mathbf{n}\otimes\mathbf{e}\}_1)$$

is the magnetic-dipole, and

$$\hat{V}_Q=-i\alpha\omega 6^{-1/2}F(\hat{Q}_2\{\mathbf{n}\otimes\mathbf{e}\}_2)$$

is the electromagnetic dipole interaction operator ($\mathbf{d}$, $\boldsymbol{\mu}$, and Q_2 are the dipole, magnetic, and quadrupole moments of the atom, respectively, $\alpha=1/137$ is the fine-structure constant, $\{\mathbf{n}\otimes\mathbf{e}\}_j$ is the irreducible tensor product of the vectors $\mathbf{n}=\mathbf{k}/k$ and $\mathbf{e}$, of rank j).

The initial state of the atom $|i\rangle$ will be assumed for simplicity to be a spherically symmetrical S state, and the final state

$$|f\rangle=\sum_{l=0}^{\infty}\sum_{m=-l}^{l} i^l \exp(-i\delta_l) Y_{lm}^*(\mathbf{n}_f)|Elm\rangle, \quad (4)$$

where δ_l is the partial phase shift for the scattering of an optical electron by the atomic core, $\mathbf{n}_f=\mathbf{k}_f/k_f$. The coefficients of expansion in the eigenvectors $|Elm\rangle$ of the continuous spectrum of an isolated atom in (4) were chosen such as to eliminate from (2) the factors corresponding to the final-state density.

From the expression for the matrix element of the ionization transition under the influence of the perturbation (3) it can be seen that the corrections to the dipole interaction lead to an insignificant change of the ionization probability dW_N, which can become noticeable only in the vicinity of the quadrupole intermediate resonance (at $N\geqslant 2$). Upon integration along the directions of the photoelectron emission, the corrections of odd order in α vanish, so that the contribution of the nondipole terms to the total probability is only of the order of α^2. Therefore, the role of the terms of order α in multiphoton ionization of atoms is negligible—they lead only to a redistribution of the electron emission directions, and only corrections of α^2 were considered in Ref. 5.

The corrections of order odd in powers of α, however, which determine the asymmetry in the photoelectron emission directions, produce in the atomic medium a current that flows along or against the electromagnetic-wave propagation direction determined by the unit vector $\mathbf{n}$. In particular, in the case of one-photon ionization ($N=1$) or a linearly polarized wave

$$dW_1^{\ l}=\frac{F^2}{2}(\mathbf{en}_f)^2\{(\rho^1_{E1,0})^2+\alpha\omega(\mathbf{nn}_f)\cos(\delta_2-\delta_1)\rho^1_{E1,0}\rho^2_{E2,0}\}\,d\Omega_f, \quad (5)$$

where

$$\rho^n_{El,0}\equiv\langle El|r^n|0\rangle$$

is a radial 2^n-pole matrix element. In the case of circular polarization or unpolarized radiation it is necessary to make in (5) the substitution

$$(\mathbf{en}_f)^2\to{}^1/_2[\mathbf{n}\times\mathbf{n}_f]^2.$$

As seen from (5), the asymmetry of the angular distribution of the photoelectrons, due to the quadrupole interaction of the atom with the field $\hat{V}_Q$, is determined by the product of the dipole and quadrupole radial matrix elements and by the phase difference between the D and P states in the continuous spectrum of the atom:

$$R_{21}\equiv\cos(\delta_2-\delta_1)\rho^1_{E1,0}\rho^2_{E2,0}.$$

This expression, generally speaking, can be both positive and negative, corresponding to a preferred emission of the photoelectron along or against the direction of propagation of the electromagnetic wave. Obviously, the sign depends both on the individual properties of the intra-atomic potential and on the frequency of the incident radiation. The direction of the drag current, determined by this sign, can thus serve as an experimental criterion of the applicability of the theory used in the calculation of electromagnetic transitions in atoms. A detailed discussion of the question of the influence of the interelectron interaction on the sign of R_{21} in inert gases was carried out in Ref. 3 on the basis of the random-phase approximation with exchange. In the case of the hydrogen atom $R_{21}>0$ always, i.e., the electrons are dragged only in the wave propagation direction, and the ionization probability takes here the simple form[1]

$$\frac{dW_1^{\ l}}{d\Omega_f}=\frac{4F^2\exp(-4k_f^{-1}\operatorname{arctg}k_f)}{\omega^5[1-\exp(-2\pi/k_f)]}(\mathbf{en}_f)^2\{1+4\alpha k_f(\mathbf{nn}_f)\}.$$

It must be noted that the magnetic interaction determined by the operator $\hat{V}_\mu$ does not contribute to the asymmetric term in (5), since an $M1$ transition between the ground state and the continuous-spectrum state in the atom is forbidden. For the same reason, we can neglect here also the relativistic effects (the electron spin).

3. In the case of N-photon ($N\geqslant 2$) ionization the matrix element in (2) is composite, and in the case of corrections of first order in α it can be represented in the form

$$M_{fi}=(-F)^N\{\langle f|(\hat{\mathbf{ed}})G_{E_i+(N-1)\omega}(\hat{\mathbf{ed}})\ldots G_{E_i+\omega}(\hat{\mathbf{ed}})|i\rangle$$
$$-i2^{1/2}\langle f|(\hat{\boldsymbol{\mu}}\{\mathbf{n}\otimes\mathbf{e}\}_1)G_{E_i+(N-1)\omega}(\hat{\mathbf{ed}})\ldots G_{E_i+\omega}(\hat{\mathbf{ed}})|i\rangle$$
$$-\ldots+i\alpha\omega 6^{-1/2}\langle f|(\hat{Q}_2\{\mathbf{n}\otimes\mathbf{e}\}_2)G_{E_i+(N-1)\omega}(\hat{\mathbf{ed}})\ldots G_{E_i+\omega}(\hat{\mathbf{ed}})|i\rangle$$
$$+\ldots\}, \quad (6)$$

where G_E is the Green's function of the optical electron of the atom. This expression contains one electrodipole and N each magnetodipole and electroquadrupole composite matrix elements that differ only in permutations of the operator $\hat{V}_\mu$ or $\hat{V}_Q$ with the operators $\hat{V}_d$.

Integrating in (6) with respect to the angle variables of the photoelectron and substituting in (2) we can, particularly for circularly or unpolarized radiation, represent the probability of N-photon ionization in the form

$$\frac{dW_N^c}{d\Omega_f}=\frac{1}{2}\left\{\frac{F^2}{2}[\mathbf{n}\times\mathbf{n}_f]^2\right\}^N\{(\rho^{1,1,\ldots 1}_{EN,N-1,\ldots,1,0})^2+\alpha\omega(\mathbf{n}\mathbf{n}_f)R_{N+1,N}\}, \tag{7}$$

where

$$\rho^{n_{k+1},n_k,\ldots n_2,n_1}_{El,l_k,\ldots l_1,0}\equiv\langle El|r^{n_{k+1}}g^{E_l+(N-1)\omega}_{l_k}r^{n_k}\ldots g^{E_l+\omega}_{l_1}r^{n_1}|0\rangle$$

is the radial component matrix element for the N-th order transition from the ground state $|0\rangle$ into the continuous spectrum $|El\rangle$, g^E_l is the radial Green's function[5] and

$$R_{N+1,N}\equiv\cos(\delta_{N+1}-\delta_N)\rho^{1,1,\ldots 1}_{EN,N-1,\ldots 1,0}[\rho^{2,1,\ldots 1}_{EN+1,N-1,\ldots 1,}$$
$$+\rho^{1,2,1,\ldots 1}_{EN+1,N,N-2,\ldots 1,0}+\ldots+\rho^{1,1,\ldots 1,2}_{EN+1,N,N-1,\ldots 3,2,0}] \tag{8}$$

is the product of the dipole matrix element and the sum of the quadrupole composite matrix elements (the analog of R_{21} in one-photon ionization).

The expression for the probability of N-photon ionization by a linearly polarized wave differs from (7) in having a more complicated dependence on the photoelectron emission direction $\mathbf{n}_f$ of both the symmetric and asymmetric components; this dependence is determined by a combination of composite matrix elements with a different set of orbital angular momenta in the intermediate and final states. Thus, in the case of two-photon ionization

$$\frac{dW_2^l}{d\Omega_f}=\frac{F^4}{18}\left\{(\rho^{1,1}_{E0,1,0})^2+[1-3(\mathbf{e}\mathbf{n}_f)^2]^2(\rho^{1,1}_{E2,1,0})^2\right.$$
$$\left.+\frac{3}{5}\alpha\omega(\mathbf{n}\mathbf{n}_f)[A_1+(1-5(\mathbf{e}\mathbf{n}_f)^2)A_3]\right\}, \tag{9}$$

$$A_i=(\rho^{1,2}_{Ei,2,0}+\rho^{2,1}_{Ei,1,0})[\cos(\delta_i-\delta_0)\rho^{1,1}_{E0,1,0}$$
$$+\cos(\delta_i-\delta_2)[1-3(\mathbf{e}\mathbf{n}_f)^2]\rho^{1,1}_{E2,1,0}]. \tag{10}$$

In contrast to the one-photon case, the magnetic and relativistic effects make a nonzero contribution to the correction terms for M_{fi} in (6), just as the quadrupole interaction. However, the corresponding composite matrix elements are much smaller than the electroquadrupole ones and do not contain at all resonant singularities as in the case of a dipole composite matrix element. At the same time, the sum of the quadrupole matrix elements in (8) and (10) contains, besides all the resonant singularities of the dipole matrix element, its own quadrupole resonances. We can therefore neglect in the calculations the $M1$ interaction of the atom with the field of the wave. Expressions (8) and (10) were written down with account taken of this neglect.

The presence of resonance in only one of the factors that determine the quantities $R_{N+1,N}$ (8) and A_i (10) presupposes the possibility of reversal of the sign of this quantity when the sign of the detuning from resonance is reversed. Such are, in particular, the one-photon resonance with the D levels, the two-photon resonances with the P and F levels, etc., which are absent from the dipole composite matrix element. Reversal of the signs of the quantities $R_{N+1,N}$ and A_i at these frequencies will reverse the preferred emission of the photoelectrons, and this will make it possible to change the direction of the drag current on going from the red to the violet wing of the dipole-forbidden atomic line. This effect can take place for practically any atom whose dipole- and quadrupole-resonance frequencies are unequal. Only for the hydrogen atoms are the rsonance frequencies of the dipole and quadrupole transitions equal because of a degeneracy of the spectrum. Consequently when hydrogen is ionized the direction of the preferred emission of the photoelectrons always agrees with the direction of propagation of electromagnetic waves. This is obviously due to the fact that the excess momentum produced when the preferred photoelectron direction is reversed can be transferred in a multielectron atom to the electron shell of a core, which hydrogen does not have, so that no such effect can occur in hydrogen.

In the vicinity of the dipole resonances both factors in (8) and (10) are resonant, so that R_{N+1} and $A_i\sim\Delta^{-2}$ and do not depend on the sign of the detuning Δ. These resonances are also of great practical interest because of their strong dependence on Δ, which makes possible a considerable increase of the drag current with decreasing detuning.

4. To find the drag-current density we can use an equation obtained in Ref. 3 on the basis of the transport equation

$$j=-\sigma^{-1}\int\frac{dW_N}{d\Omega_f}(\mathbf{n}\mathbf{n}_f)\,d\Omega_f, \tag{11}$$

where σ is the cross section for the scattering of a photoelectron by an atom. This expression is applicable in the case when the ionizing radiation acts on the medium continuously. If, however, the ionizing flux is applied to the medium in pulses, the expression for the time-averaged current density is obtained by multiplying the right-hand side of (11) by the pulse duration τ and by the number ν of the pulses per unit time. For the transport equation to be valid it is then necessary to satisfy the relation $\tau\gg\tau_{\rm fr}$, where $\tau_{\rm fr}$ is the mean-free path time of the electrons prior to the collision with the atoms of the medium.

At $\tau\ll\tau_{\rm fr}$ we can use the usual expression for the current density

$$dj=-d\rho v_n, \tag{12}$$

where $d\rho$ is the number of photoelectrons, per unit volume, moving in the direction determined by the solid angle $d\Omega_f$, and $v_n=\mathbf{k}_f\cdot\mathbf{n}$ is the projection of the velocity on the direction of the ordered motion. The quantity $d\rho$ can be expressed in terms of the ionization probability $dW/d\Omega_f$, the effective ionizing-pulse duration $\tau\nu$, and the time of directional motion of the photoelectrons, which coincide with the free-path time $\tau_{\rm fr}$:

$$d\rho=N_0\tau\nu\frac{dW}{d\Omega_f}d\Omega_f\tau_{\rm fr},$$

where N_0 is the number of atoms per unit volume of the gas.

Substituting this expression in (12), integrating over the

directions of the emission of the photoelectrons, and using the expression for free-path time $\tau_{fr} = (N_0\sigma k_f)^{-1}$, we obtain

$$j = -\tau\nu\sigma^{-1}\int \frac{dW}{d\Omega_f}(\mathbf{n}\mathbf{n}_f)\,d\Omega_f. \tag{13}$$

It can be seen that this expression coincides fully with (11) if account is taken in the latter of the pulsed character of the ionizing radiation (by multiplying by $\tau\nu$).

Thus, the drag-current density j is practically independent of the density of the atomic medium (at $\rho \ll N_0$) and is determined only by the integral cross section σ for the scattering of the electron by the atom. This cross section, however, depends on the characteristics of the ionizing radiation—on the power, duration, number of pulses per unit time and on the relation between the radiation frequency and natural frequencies of the atoms, a relation on which the manifestation of the resonant singularities in composite matrix elements of multiphoton ionization depends.

Substituting (7) in (13) we obtain an expression for a drag-current density in N-photon ionization of atoms by a circularly polarized wave:

$$j_N{}^c = -\frac{N!}{(2N+3)!!}2\pi\nu\tau\sigma^{-1}\alpha\omega F^{2N}R_{N+1,N}. \tag{14}$$

In the case of linear polarization of electromagnetic waves, as shown by calculations, the largest contribution to the transition amplitude is made by the matrix elements with the maximum possible value of the orbital quantum number of the photoelectron in the intermediate and final states,[5] i.e., by those which determine the quantity $R_{N+1,N}$ in (14). If we neglect the contribution of the remaining matrix elements, we can obtain an approximate relation between the current densities in N-photon ionization by a circularly (j_N^c) and linearly polarized wave (J_n^l):

$$j_N{}^c/j_N{}^l \approx 2(2N+1)!!/(N+2)!. \tag{15}$$

We note that (15) is the upper bound of the ratio j_N^c/j_N^l, which can vary significantly depending on the relation between the frequency of the electromagnetic wave and the natural frequencies of the atoms. Thus, for a circularly polarized wave we can indicate the frequencies at which the dipole or the sum of all the quadrupole composite matrix elements in (8) vanishes. The vanishing of the entire sum of the radial matrix elements with different (not the largest possible) quantum numbers that determine j_N^l has low probability, since at these frequencies j_N^c/j_N^l is zero. At the same time expression (15) shows that $j_N^c/j_N^l \geqslant 1$ and increases rapidly with increasing N. In addition, when N increases the number of the radial matrix elements with nonmaximal orbital number and their contribution to the amplitude of ionization by a linearly polarized wave increase, so that actually the ratio (15) decreases at large N. Finally, the composite matrix elements that determine the current due to dragging of photoelectrons by a linearly polarized wave have more resonant singularities at $N \geqslant 3$ than in the case of a circularly polarized wave. In particular, we have here two-photon resonances with S states in the dipole matrix element and with P states in the quadrupole matrix elements, resonances which are not contained in the quantity $R_{4,3}$ that determines the drag current in three-photon ionization by a circularly polarized wave ($R_{4,3}$ contains two-photon resonances only with the D and F states, resonances present also in a linearly polarized field). In the vicinity of these resonances j_N^l increases sharply, so that $j_N^c/j_N^l \to 0$. By virtue of all the foregoing factors, the applicability of expression (15) is restricted to small values of N and to frequencies that are far from resonances and are absent in $R_{N+1,N}$.

To estimate the absolute value of the drag-current density and to assess the possibility of experimentally observing this effect in multipliphoton ionization, we consider the simplest example, two-photon ionization of hydrogen atoms by radiation of frequency $\omega = \omega_{2P,1S} - \Delta$, where $\Delta \ll \omega$ is the detuning of resonance with the frequency of the leading line of the Lyman series ($\omega_{2P,1S} = 3/8$ a.u.). In the vicinity of the resonance we can retain in the composite matrix elements in (7) and (9) only the terms that make the principal contribution. All the ratial matrix elements can then be calculated in analytic form. As a result we obtain

$$j_2{}^c = -(^{16}/_3)^{10}\frac{\pi\alpha\omega k_f(1+k_f^2)\exp(-4k_f^{-1}\operatorname{arctg}2k_f)}{35(1+4k_f^2)^8[1-\exp(-2\pi/k_f)]} \times F^4\nu\tau\sigma^{-1}\Delta^{-2},$$

$$j_2{}^c/j_2{}^l = \frac{96(1+k_f^2)}{67+60k_f^2}.$$

Substituting in these expressions the numerical value of the photoelectron momentum

$$k_f = [2(2\omega+E_{1S})]^{1/2} = 2^{-1/2},$$

as well as the cross section $\sigma = 10\pi$ a.u. for the scattering of an electron with this momentum by hydrogen atoms,[9] we obtain

$$j_2{}^c = -2.54\cdot 10^{10}\,\nu\tau F^4/\Delta^2\,[\mathrm{A/cm^2}],\quad j_2{}^c/j_2{}^l = {}^{144}/_{97}.$$

At $\nu\tau = 10^{-8}$, $F = 5\times 10^5$ V/cm and $\Delta = 100\,\mathrm{cm}^{-1}$ the current density is $j_2^c = -1.2\times 10^{-7}$ A/cm^2, much larger (by approximately five orders of magnitude) than the estimate obtained in Ref. 3 for the photoelectron drag-current density in one-photon ionization of atoms.

Obviously, in multiphoton ionization of alkali-metal vapor one can attain a much higher photoelectron drag-current density by using modern high-power pulsed lasers. For alkali atoms it is then quite easy to choose the resonant frequency, including those that are dipole-forbidden, in the vicinity of which the reversal of the drag-current direction can be observed.

In particular, for two-photon ionization of sodium atoms by the second harmonic of a ruby laser ($\omega = 2\omega_R = 28\,800\ \mathrm{cm}^{-1}$) the detuning of the quadrupole resonance from the 3^2D level is $\Delta_Q = 373\ \mathrm{cm}^{-1}$, and that of the dipole resonance from the 4^2P level is $\Delta_d \approx 1467\ \mathrm{cm}^{-1}$. Using the previously calculated values of the cross sections for two-photon ionization[8] we obtain at $F = 5\times 10^6$ V/cm, $\nu\tau = 10^{-8}$, $j_2^c \approx -10^7$ A/cm^2. With increasing frequency the current increases, since Δ_Q and Δ_d decrease, and at a frequency $\omega > 29\,173\ \mathrm{cm}^{-1}$ the current direction should be reversed, since Δ_Q reverses sign.

A similar effect can take place in two-photon ionization

of rubidium by the second harmonic of a neodymium-laser ($\omega = 2\omega_N = 18\,880$ cm^{-1}). The detuning from resonance with the frequency of the quadrupole transition into the 4^2D state is here $\Delta_Q \approx 475$ cm^{-1}, and the frequencies of the dipole transitions are far from resonance. Owing to the much larger cross section for two-photon ionization than in sodium,[8] the drag current at the same laser-pulse parameters is larger in this case by approximately two orders of magnitude, $j_2^c \approx 10^{-5}$ A/cm^2. The direction of the current is reversed at the frequency $\omega > 19\,355$ cm^{-1}.

In the examples considered, the relation (15) is satisfied with good accuracy, i.e., $j_2^c/j_2^l \approx 5/4$.

5. Thus, the results of our calculations show that the photoelectron drag current in multiphoton ionization of atoms by the existing high-power optical-radiation sources can exceed by many orders of magnitude the drag current in one-photon ionization by available vacuum-ultraviolet or x-ray sources. The two principal factors that cause this increase are the appreciable radiation intensities of the optical-band lasers and the possibility of a resonant enhancement of the effect in question, a possibility that does not exist in one-photon ionization.

These factors make the dragging of photoelectron by optical radiation useful not only in experiment but also in practice. In particular, the dragged current can serve as an autonomous indicator of multiphoton ionization of atoms and molecules. Another indicator can be also the drag-current alternating magnetic field produced when the laser pulse is turned on and off. In addition, converters of high-power light pulses into electric current can be produced on the basis of cells with atomic gases that have high effective multiphoton ionization.

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Translated by J. G. Adashko.