Charge-sign dependence of the differential yield of the helium autoionization resonances $(2s2p)^1P$ and $(2p^2)^1D$ excited by fast protons and antiprotons

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The reasons for the appearance of a dependence of the differential yield of the autoionization resonances on the sign of the charge of the incoming particles are analyzed. The influence of the two-step excitation mechanism and the Coulomb interaction in the final state on the differential yield of the helium resonances $(2s2p)^1P$ and $(2p^2)^1D$ is investigated for the energy of the incoming protons and antiprotons equal to 1.84 MeV. The coefficients of the angular anisotropy are used to analyze the effect of the sign of the charge of the scattered particle on the angular distribution of the ejected electrons in the energy range 0.5–5 MeV. © *1996 American Institute of Physics.* [S1063-7761(96)00505-7]

The investigation of the effect of the sign of the charge of the incident particle on the dynamics of various elementary processes is a rapidly developing direction of the physics of atomic collisions. This increased interest in the interaction of atoms with antiparticles is connected with additional possibilities of investigating the dynamics and kinematics of the collisions. For fast collisions, flipping the sign of the charge of the incoming particle (in contrast to changing the speed or absolute value of the charge of the incoming particle) alters only the contribution of the interference terms to the differential cross section, without changing the relations between the various mechanisms of the considered processes.[1,2] Ionization collisions with excitation and subsequent decay of the autoionization states allow us, by changing the sign of the incoming particle, to investigate both the relation between the various mechanisms of excitation and ionization and the effect of peculiarities of the scattering of the particles in the final state on the profile of the auto-ionization resonances.

Experiments on collisions of helium atoms with protons and electrons having the same velocity[3,4] confirm the main conclusion of Ref. 5 that the profile of the autoionization resonances depend on the sign of the charge of the incoming particles. It was experimentally observed that for collisions with protons the differential yield of the resonances $(2s2p)^1P$ and $(2p^2)^1D$ of the helium atom at some ejection angles is substantially higher than for collisions with electrons.[3] At large ejection angles the relation of the differential yields of these resonances, excited by protons and electrons, is just the opposite. Due to the rather large collision velocity, these differences are associated not with the difference in the mass, but with the different sign of the charge of the incoming particle. The angle-integrated ejection yields of the investigated autoionization resonances for electrons and protons coincide within the limits of experimental error.[3] The total cross section of single ionization of the helium atom, which defines the background against which the resonances are manifested, in the energy range under consideration does not depend on the sign of the charge of the incoming particle.

In theoretical studies of the two-electron excitation cross sections the reasons for the appearance of a charge-sign dependence in the differential yields has been analyzed both qualitatively[7] and quantitatively.[8,9] Attention has been focused on the interference of amplitudes describing the different mechanisms of excitation. However, the coupling of the resonances with the continuum and also their population through the continuum were neglected in Refs. 8 and 9. Here it may be noted that the influence of two-step transitions on the excitation cross section of the $(2s2p)^1P$ resonance of the helium atom, which makes the main contribution to the measured yield[3] in the energy region $E_i \gg 1.5-1.84$ MeV, has been found to be insignificant. The available calculations are restricted to a treatment of total cross sections of two-electron excitation and do not touch upon questions of the appearance of a charge-sign dependence of the differential yield of the investigated autoionization resonances.[3] In point of fact, more complete information about the nature of the sign-charge dependence of the autoionization resonances can be obtained by examining their differential characteristics. In this case, the interaction of the ejected electron and the ion-shell of the target with the scattered particle as much as interference of the resonances with the continuum (if these amplitudes have a different charge-sign dependence) are the source of the charge-sign dependence of the profile of the auto-ionization resonances.[10,11] As the velocities of the ejected electron and the scattered charged particle approach each other in magnitude and direction, the conditions of interference of the resonance with the adjacent continuum change. This is manifested, for example, in the marked increase with respect to the background of the intensity of the resonances excited by antiprotons at small ejection angles. The interaction in the final state not only redistributes the knocked-out electrons in energy and ejection angles, but also

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alters the ejection-angle–integrated yield of the resonance.[12] These changes appear not only as a result of disruption of the parametrization of the resonance profile[13] but also because of "switching on" the interference of the resonance with the incoherent background.

Quite recently, the dependence of the autoionization resonance profile of helium on the sign of the charge of the incoming particle was investigated in the one-center approximation by the method of strong coupling of channels at energies of the incoming protons and antiprotons of 2 MeV (Ref. 14) and 1.5 MeV (Ref. 15). However, because they neglected the effect of the Coulomb interaction in the final state in the resonance ionization amplitude, these theoretical studies are incomplete.

The aim of the present work is a theoretical study of the charge-sign dependence of the differential yield of the $(2s2p)^1P$ and $(2p^2)^1D$ autoionization resonances of the helium atom, excited by fast protons and antiprotons.

The angular distribution of the ejected electrons can be written as[6]

$$\frac{d\sigma}{d\Omega_e}=\int dE_e F(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})+\sum_\mu Y_\mu(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt}). \tag{1}$$

Here θ_c is the ejection angle relative to the incident beam, E_i and Z_p are the energy and charge of the incident ion, E_e is the energy of the electron, $d\Omega_e$ is the solid angle element in the direction of the trajectory of the ejected electron, and ν_{pe} and ν_{pt} are the Coulomb parameters, which take account of the interaction of the scattered ion with the electron and target ion-shell, respectively.[17] The first term in Eq. (1) describes the background created by the direct ionization transitions, and the second term is the sum of the differential yields of the resonances.

Let us consider the region of the kinematics where the quantities describing the resonance profile vary slowly over its width as functions of the energy of the autoionization electron. Then the differential yield of the resonance can be represented in the form[13]

$$Y_\mu(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})=Y_{\mu,\mathrm{res}}(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})+Y_{\mu,\mathrm{int}}(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt}), \tag{2}$$

$$Y_{\mu,\mathrm{res}}(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})=\frac{\pi\Gamma_\mu}{2}4\frac{k_f}{k_i}\int d\Omega_{\mathbf{k}_f}|t_\mu^{\mathrm{PWBA}}K_{\mathrm{res},\mu}(\nu_{pe},\nu_{pt})|^2, \tag{3}$$

$$Y_{\mu,\mathrm{int}}(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})=\frac{\pi\Gamma_\mu}{2}8\frac{k_f}{k_i}\int d\Omega_{\mathbf{k}_f}\,\mathrm{Im}\{t_{\mathrm{dir}}^{\mathrm{PWBA}*}K_{\mathrm{dir}}^*(\nu_{pe},\nu_{pt})t_\mu^{\mathrm{PWBA}}K_{\mathrm{res},\mu}(\nu_{pe},\nu_{pt})\}, \tag{4}$$

where

$$t_{\mathrm{dir}}^{\mathrm{PWBA}}=\sum_L\sqrt{\frac{2L+1}{4\pi}}\,t_{\mathrm{dir}}^{(L)}(Q,k_e)P_L(\hat{Q},\hat{k}_e), \tag{5}$$

$$t_\mu^{\mathrm{PWBA}}=\sqrt{\frac{2}{\pi\Gamma_\mu}}\,t_{ex,\mu}(E_i,Z_p,Q)P_{L_\mu}(\hat{Q},\hat{k}_e) \tag{6}$$

are the amplitudes of direct ionization and ionization via the autoionization resonances in the plane-wave approximation, and the factors $K_{\mathrm{dir}}(\nu_{pe},\nu_{pt})$ and $K_{\mathrm{res},\mu}(\nu_{pe},\nu_{pt})$ allow for the interaction in the final state[17,18]

$$K_{\mathrm{dir}}(\nu_{pe},\nu_{pt})\to 1,\quad K_{\mathrm{res},\mu}(\nu_{pe},\nu_{pt})\to 1\quad \text{for}\quad \nu_{pe}\to 0,\quad \nu_{pt}\to 0.$$

In Eqs. (3)–(6) $\mathbf{k}_i$ and $\mathbf{k}_f$ are the momenta of the incident and scattered particle, $\mathbf{Q}=\mathbf{k}_i-\mathbf{k}_f$ is the momentum transfer, Γ_μ is the width of the resonance with orbital angular momentum L_μ, $\hat{k}_e$ is the direction of the knocked-out electron, and $P_L(x)$ are the Legendre polynomials. The amplitude of excitation of the autoionization resonance

$$t_{ex,\mu}(E_i,Z_p,Q)=t_{ex,\mu}^{(1)}(Z_p,Q)+\tilde{t}_{ex,\mu}^{(1)}(Z_p,Q)+t_{ex,\mu}^{(2)}\times(E_i,Z_p,Q) \tag{7}$$

is described within the framework of second-order perturbation theory in the interaction of the target with the incident particle.[16,19] Here $t_{ex,\mu}^{(1)}$ is the amplitude of the two-electron correlated transitions in the first Born approximation, $\tilde{t}_{ex,\mu}^{(1)}$ is the amplitude of the population of states of the two-electron excitation via the continuum, and $\tilde{t}_{ex,\mu}^{(2)}$ is the amplitude of the two-step excitation in the second Born approximation.[20]

The modification of the differential yield $Y_{\mu,\mathrm{res}}$ by the Coulomb interaction in the final state for fast collisions can be neglected:[13]

$$Y_{\mu,\mathrm{res}}(E_i,Z_p,\theta_e,\nu_{pe},\nu_{pt})=Y_{\mu,\mathrm{res}}(E_i,Z_p,\theta_e,\nu_{pe}=0,\nu_{pt}=0)$$

for

$$\text{for}\ \frac{2}{\Gamma_\mu}(k_{pe}V_f+\mathbf{k}_{pe}\mathbf{V}_f)\gg 1, \tag{8}$$

where $\mathbf{k}_{pe}$ is the momentum of the ejected electron relative to the scattered ion. However, the effect of the Coulomb interaction on the interference part of the differential yield (4) cannot be neglected in the kinematic region (8). It should be mentioned that the modification of the relative phase between the amplitudes of the direct and resonance ionization is enhanced to the extent that the resonance width Γ_μ is decreased.[21]

As was already noted, there exists a region of energies E_i, where the effect of the scattered ion on the mechanism of direct ionization can be neglected.[6] Then $t_{\mathrm{dir}}^{\mathrm{PWBA}}$ is described to first order in the interaction of the target with the incident particle and is proportional to Z_p. In this case the charge-sign dependence of the differential yield of the resonance (1) is governed by the following factors:

a) interference between the excitation amplitudes having different charge dependence, and in the pole approximation for $\tilde{t}^{(1)}_{ex,\mu}$ and $t^{(2)}_{ex,\mu}$ in Eq. (7) only the last two terms interfere:[16]

$$\mathrm{Re}\{t^{(2)}_{ex,\mu}(E_i,Z_p,Q)\tilde{t}^{(1)*}_{ex,\mu}(Z_p,Q)\}\propto\frac{Z_p^3}{V_i}; \quad (9)$$

b) interference between the amplitudes of direct and resonance ionization even if only one of them has charge dependence different from Z_p^1:

$$\mathrm{Im}\{t^{(2)}_{ex,\mu}(E_i,Z_p,Q)t_{\mathrm{dir}}^{\mathrm{PWBA}*}\}\propto\frac{Z_p^3}{V_i}; \quad (10)$$

c) the interaction in the final state, taken into account in expressions (3) and (4) by the factors K_{dir} and $K_{\mathrm{res},\mu}$, alters both the absolute values of the amplitudes and their relative phase.

Because of the stronger, more energetic dependence of $t^{(2)}_{ex,\mu}$ in comparison with the amplitudes of the correlated transitions $t^{(1)}_{ex,\mu}$ and $\tilde{t}^{(1)}_{ex,\mu}$ the contribution of terms (9) and (10) to the differential yield (2) decreases with growth of the collision velocity V. The effect of the interaction in the final state on the resonance profile also decreases. Therefore, in the limit $V_i\to\infty$ the charge-sign dependence of the resonance profile is absent.

Let us next consider the angular distribution of the autoionization electrons in the case in which we neglect the interaction in the final state $K_{\mathrm{res},\mu}=1$ and the interference of the autoionization resonance with the background $Y_{\mu,\mathrm{int}}=0$. Then itegrating Eq. (4) over the scattering angle we obtain for $L_\mu\geqslant 1$

$$Y_{\mu,\mathrm{res}}(E_i,Z_p,\theta_e,\nu_{pe}=0,\nu_{pt}=0)=Y_0(L_\mu,E_i,Z_p)\times\left\{1+\sum_{n=1}^{L_\mu}\beta_e^{(n)}\times(L_\mu,E_i,Z_p)\cos^{2n}\theta_e\right\}, \quad (11)$$

where the ejection-angle–independent constant $Y_0(L_\mu,E_i,Z_p)$ is related to the excitation cross section $\sigma_{ex,\mu}(E,Z_p)$ by the formula

$$Y_0(L_\mu,E_i,Z_p)=\frac{\sigma_{ex,\mu}}{4\pi}\left(1+\sum_{n=1}^{L_\mu}\frac{\beta_e^{(n)}(L_\mu,E_i,Z_p)}{2n+1}\right)^{-1}. \quad (12)$$

Calculations of the anisotropy coefficients of the angular distributions of the electrons for $L_\mu=1, 2$ give the following results:

$$\beta_e^{(1)}(L_\mu=1,E_i,Z_p)=\frac{3G_1(E_i,Z_p)-1}{1-G_1(E_i,Z_p)}, \quad (13)$$

$$\beta_e^{(1)}(L_\mu=2,E_i,Z_p)=6\frac{42G_1(E_i,Z_p)-45G_2(E_i,Z_p)-5}{11-30G_1(E_i,Z_p)+27G_2(E_i,Z_p)}, \quad (14)$$

$$\beta_e^{(2)}(L_\mu=2,E_i,Z_p)=9\frac{3-30G_1(E_i,Z_p)+35G_2(E_i,Z_p)}{11-30G_1(E_i,Z_p)+27G_2(E_i,Z_p)}, \quad (15)$$

where

$$G_n(E_i,Z_p)=\frac{1}{\sigma_{ex,\mu}}\int d\Omega_{\mathbf{k}_f}\cos^{2n}(\hat{Q},\hat{k}_i)\frac{d\sigma_{ex,\mu}}{d\Omega_{\mathbf{k}_f}} \quad (16)$$

is the mean value of the even powers of the cosine of the angle between the directions of the incident beam $\mathbf{k}_i$ and the momentum transfer $\mathbf{Q}$ $(0\leqslant G_n(E_i,Z_p)\leqslant 1)$. With increase of the collision energy E_i the contribution of the small scattering angles grows in Eq. (16), The ejected-electron distributions (11) are symmetric relative to $\theta_c=90°$. For $L_\mu=1$ the electron angular distribution is in fact determined by only one parameter $\beta_e^{(1)}$. For $\beta_e^{(1)}>0$ the electrons scatter preferentially in the direction of the incident beam:

$$Y_\mu(E_i,Z_p,\theta_e=0,\nu_{pe}=0,\nu_{pt}=0)>Y_\mu(E_i,Z_p,\theta_e=90°,\nu_{pe}=0,\nu_{pt}=0).$$

For $\beta_e^{(1)}<0$ this inequality is reversed, and the electrons scatter preferentially perpendicular to the incident beam. The case $\beta_e^{(1)}=0$ corresponds to an isotropic electron distribution. Taking account of the interference with the background created by the direct ionization transitions disrupts dependence (11). The model character of this treatment consists in the fact that both terms in Eq. (2) frequently have comparable magnitudes and possess identical asymptotic limits as $E_i\to\infty$. In general, it is not possible to neglect either one of them. However, in the case of radiative decay of the autoionization resonances of the helium atom the angular distribution of the photons is almost completely determined by the resonance amplitude.[21] The anisotropy coefficients of the photon angular distributions $\beta_\gamma^{(n)}$, defined in terms of $G_n(E_i,Z_p)$, quite rigorously describe the real picture of radiative decay of the autoionization resonances. In the case of radiative decay of the $(2s2p)^1P$ resonance of the helium atom to the ground state the angular anisotropy coefficient is given by the following form:

$$\beta_\gamma^{(1)}(L_\mu=1,E_i,Z_p)=\frac{1-3G_1(E_i,Z_p)}{1+G_1(E_i,Z_p)}. \quad (17)$$

Comparison of the angular anisotropy coefficients of the resonances excited by protons and antiprotons allows us to investigate the effect of the interference of excitation amplitudes having different charge dependences (9) on the angular distributions of the knocked-out electrons.

To investigate the dependence of the differential yield of the resonances on the sign of the charge of the incident particle, we calculated the yield of the $(2s2p)^1P$ and $(2p^2)^1D$ helium resonances excited by protons and antiprotons at $E_i=1.84$ MeV, and also the angular anisotropy coefficients of these resonances in the region $E_i=0.5-5.0$MeV. The non-Coulomb part of the potential of the ion-shell He^+ $(1s)$ was taken into account as in Ref. 22. In the calculations

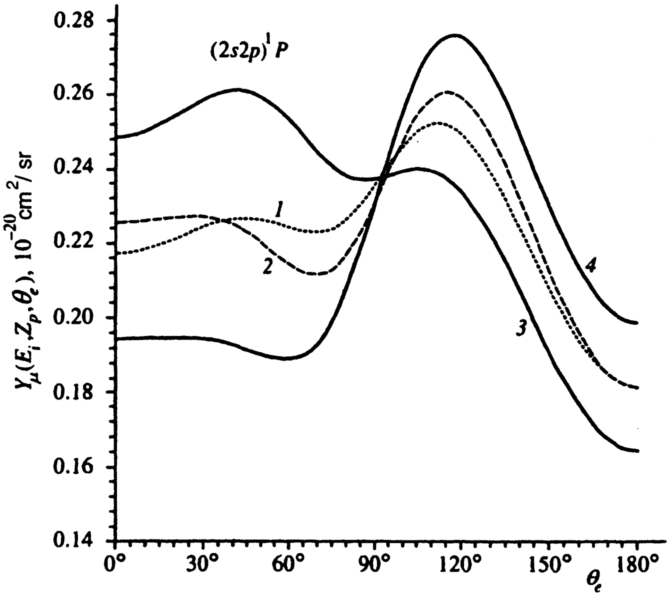

FIG. 1. Differential yield of the helium resonance $(2s2p)^1P$ formed in collisions with protons and antiprotons with energy 1.84 MeV. The dashed lines are the results of the present calculations neglecting the Coulomb interaction in the final state $Y_{\mu,res}(E_i,Z_p,\theta_e,\nu_{pe}=0,\nu_{pt}=0)$ for collisions with protons (curve *1*) and antiprotons (curve *2*). The solid lines show the results of calculations that took account of the interaction in the final state (2) for collisions with protons (curve *3*) and antiprotons (curve *4*).

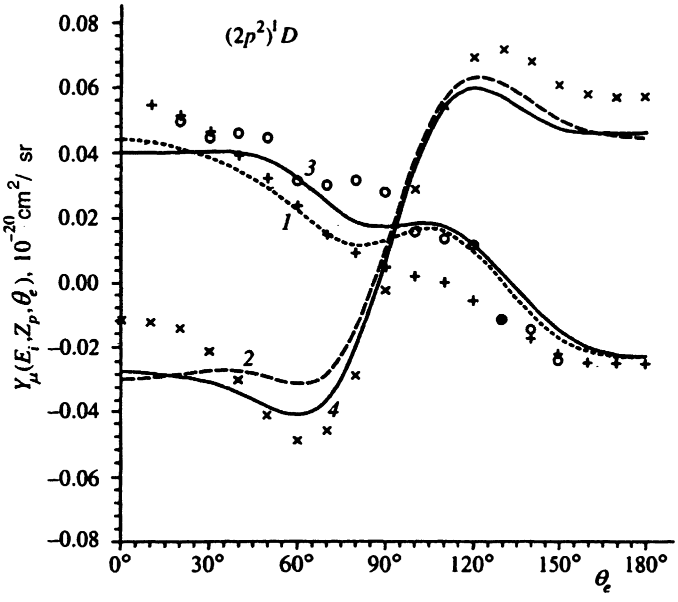

FIG. 2. Differential yield of the helium resonance $(2p^2)^1D$ formed in collisions with protons and antiprotons with energy 1.84 MeV. The results of calculations using the strong-coupling method for the protons (+) and antiprotons (×), and also experimental data for the protons (○) with energy 2 MeV taken from Ref. 14. Remaining notation as in Fig. 1.

of the excitation amplitude (7) we used wave functions obtained by the method of superposition of configurations inthe charge $Z_t=+2$ including the lower states of the one- and two-electron excitations of even parity. In the calculations of the amplitude of the direct ionization the ground state of the helium atom was described by the Hartree–Fock function.[23]

Figure 1 shows the dependence of the differential yield of the $(2s2p)^1P$ resonance excited by protons and antiprotons with energy 1.84 MeV. The asymmetry of the angular distribution relative to $\theta_c=90°$ testifies to the necessity of taking account of interference with the background (the second term in Eq. (2)). Calculations of the differential yield without taking account of the interaction in the final state for the protons and antiprotons differed by not more than 5% over the entire region of ejection angles. This means that terms (9) and (10) alter the differential yield of the $(2s2p)^1P$ resonance (2) only insubstantially. Taking account of the interaction in the final state leads to the result that in the forward hemisphere $\theta_c<90°$ the yield of the resonance excited by the protons is greater than for the antiprotons. For large ejection angles $\theta_c>90°$ the relation is reversed. Thus, the interaction in the final state is the main reason for the charge-sign dependence of the differential yield of the $(2s2p)^1P$ resonance at $E_i=1.84$ MeV.

Figure 2 shows that for all variants of calculation the dependence of the differential yield of the $(2p^2)^1D$ resonance excited by protons with energy $E_i=1.84$ MeV qualitatively differs from the dependence for the collisions with antiprotons. For small ejection angles θ_c for the collisions with the protons the resonance shows up as a peak in the energy spectrum of the ejected electrons ($Y_\mu>o$), and for the antiprotons the differential cross section at the resonance point becomes less than the background ($Y_\mu<0$). At large ejection angles the situation for the protons and antiprotons is reversed. Our results, which neglect the interaction in the final state, are close to those of Ref. 14 for energy of the incident antiprotons $E_i=2$ MeV. Taking account of the interaction in the final state leads only to moderate quantitative changes in the yields of the resonances in the forward hemisphere of ejection angles and improves the agreement of the calculated results with the experimental data of Ref. 14 for $E_i=2$ MeV. Thus, for the $(2p^2)^1D$ resonance the main reason for the appearance of the charge-sign dependence in this region of collision energies is the influence of the two-step excitation both on the interference of excitation amplitudes having a different charge dependence (9) and on the interference of the amplitudes of the transitions through the autoionization levels and the amplitudes of the ionization transitions (10).

The calculated results shown in Fig. 3 for the yields of the $(2s2p)^1P$ and $(2p^2)^1D$ resonances demonstrate qualitative agreement with the experimental results of Ref. 3. Although the differential yield of the $(2s2p)^1P$ resonance exceeds the yield of the $(2p^2)^1D$ resonance by an order of magnitude (see Figs. 1 and 2), the charge-sign dependence of the total yield is determined by both the interaction in the final state and the effect of the two-step excitation on the resonance profile.

The results of the calculations of the angular anisotropy coefficients of the $(2s2p)^1P$ and $(2p^2)^1D$ helium resonances in the collisions with protons and antiprotons are given in Tables I and II. The calculated coefficient $\beta_e^{(1)}$ of the $(2s2p)^1P$ helium resonance at $E_i=3$ MeV corresponds within the limits of error to the experimental value

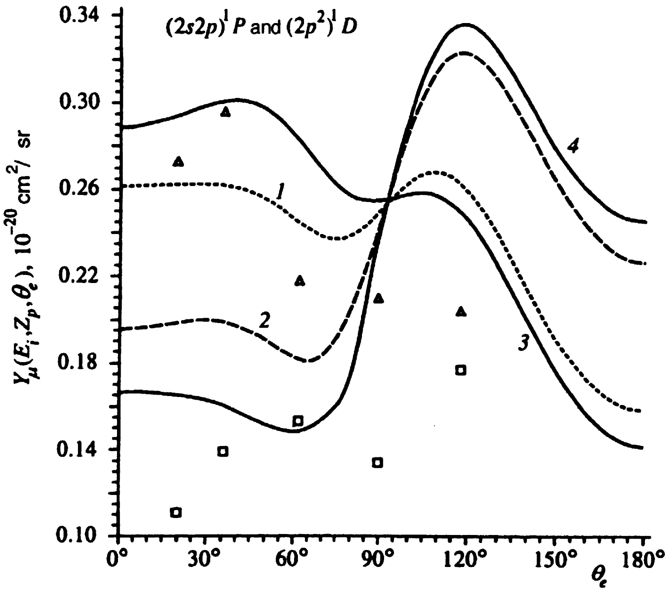

FIG. 3. Sum of the differential yields of the helium resonances $(2s2p)^1P$ and $(2p^2)^1D$ formed in collisions with protons and antiprotons with energy 1.84 MeV. The experimental points are taken from Ref. 3 with allowance for the correction given in Ref. 4: Δ—for collisions of protons with energy 1.84 MeV, □—for collisions of electrons with energy 1 keV. Remaining notation as in Fig. 1.

$$\beta_e^{(1)}(L_\mu=1,E_i,Z_p)=-(0.3\pm0.1),$$

obtained by fitting the resonance yield of Ref. 24 by formula (11). With increase of the collision energy the effect of two-step processes on the excitation amplitude (7) decreases, which leads to a weaker charge-sign dependence of the angular anisotropy coefficients. In the energy region $E_i \geqslant 1.84$ MeV for the $(2s2p)^1P$ resonance and $E_i \geqslant 3$ MeV for the $(2p^2)^1D$ resonance the angular anisotropy coefficients in the collisions with protons and antiprotons differ insignificantly. In these energy regions the change in the charge dependence of the excitation amplitude due to two-step transitions has no effect on the angular distributions of the ejected electrons.

Thus, the source of the charge-sign dependence of the asymmetry of the differential yield is different for different resonances. The amplitude of the two-step excitation of the $(2p^2)^1D$ resonance is determined by two dipole transitions while for the $(2s2p)^1P$ resonance this amplitude is determined by one dipole and one monopole transition. Consequently, the amplitude of the two-step excitation for the $(2p^2)^1D$ resonance is substantially greater than for the $(2s2p)^1P$ resonance at the same collision energy. Further, the contribution of the two-step transitions to the total excitation cross section of the $(2p^2)^1D$ resonance is greater than for the $(2s2p)^1P$ resonance[19] since the excitation of the $(2p^2)^1D$ resonance in first-order perturbation theory is very weak (the amplitude $t^{(1)}_{ex,\mu}$ is nonzero only if we take into account the angular correlations of the electrons). Thus, the influence of the two-step excitation on the total cross sections and the angular anisotropy coefficients is greater for the $(2p^2)^1D$ than for the $(2s2p)^1P$ resonance. The effect of the interaction in the final state on the differential yields of the $(2s2p)^1P$ and $(2p^2)^1D$ resonances is also different. At the incident energy $E_i=1.84$ MeV the velocity of the scattered ion is greater than that of the ejected electron and the profile of the autoionization resonance is weakly distorted. However, the conditions of interference of the amplitudes of direct and resonance ionization are very sensitive even to moderate changes in the relative phase. Since the width Γ_μ of the $(2s2p)^1P$ resonance is almost 3.5 times smaller than the width of the $(2p^2)^1D$ resonance, the effect of the Coulomb interaction in the final state on the interference with the background for the $(2s2p)^1P$ resonance turns out to be stronger.

In conclusion, we note that in the present paper we have carried out a complete analysis of all the factors leading to the appearance of a dependence of the profile of the autoionization resonance on the sign of the charge of the fast incident particle. Our investigations have revealed the existence of several reasons for the appearance of a dependence of the differential yield of the $(2s2p)^1P$ and $(2p^2)^1D$ helium autoionization resonances on the sign of the charge of the fast incident particle. The calculated angular anisotropy coefficients allow one to analyze the effect of two-step excitation on the charge-sign dependence of the angular distributions of the ejected electrons over a wide energy range. Our results show that for the energy $E_i=1.84$ MeV the charge-sign dependence of the differential yield of the $(2s2p)^1P$ resonance is determined primarily by the interaction in the final state, and for the $(2p^2)^1D$ resonance the main reason for the appearance of the charge-sign dependence consists in the change due to two-step excitation in the conditions of amplitude interference corresponding both to different mechanisms for populating the autoionization level and to direct ionization transitions into the resonance region.

TABLE I. Energy dependence of the coefficients of angular anisotropy of the electrons $\beta_e^{(1)}(L_\mu=1,E_i,Z_p)$ and photons $\beta_\gamma^{(1)}(L_\mu=1,E_i,Z_p)$ for the $(2s2p)^1P$ helium resonance, formed in collisions with protons and antiprotons.

	$\beta_e^{(1)}(L_\mu=1, E_i,Z_p)$		$\beta_\gamma^{(1)}(L_\mu=1,E_i,Z_p)$	
Energy, MeV	Protons	Antiprotons	Protons	Antiprotons
0.5	0.1057	−0.0183	−0.0502	0.0092
1.0	−0.1196	−0.2008	0.0636	0.1116
1.84	−0.2449	−0.2994	0.1396	0.1761
3.0	−0.3415	−0.3633	0.2059	0.2220
5.0	−0.3988	−0.4471	0.2491	0.2879

TABLE II. Energy dependence of the coefficients of angular anisotropy of the electrons for the $(2p^2)^1D$ helium resonance, excited in collisions with protons and antiprotons.

	$\beta_e^1(L_\mu=2,E_i,Z_p)$		$\beta_e^{(2)}(L_\mu=2,E_i,Z_p)$	
Energy, MeV	Protons	Antiprotons	Protons	Antiprotons
0.5	1.7920	0.8745	−2.0500	−1.1850
1.0	0.0978	0.3552	−0.5426	−0.0395
1.84	−1.0570	−1.2010	0.6266	0.8103
3.0	−1.6570	−1.6970	1.2660	1.3300
5.0	−2.0560	−2.0580	1.7020	1.7170

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