Parity-nonconservation effect with the laser-induced $2\,^3S_1$–$2\,^1S_0$ transition in heavy heliumlike ions

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The parity-nonconservation (PNC) effect on the laser-induced $2\,^3S_1$–$2\,^1S_0$ transition in heavy heliumlike ions is considered. A simple analytical formula for the PNC correction to the cross section is derived for the case, when the opposite-parity $2\,^1S_0$ and $2\,^3P_0$ states are almost degenerate and, therefore, the PNC effect is strongly enhanced. Numerical results are presented for heliumlike gadolinium and thorium, which seem the most promising candidates for such experiments. In both Gd and Th cases the photon energy required will be anticipated with a high-energy laser built at GSI. Alternatively, it can be gained with ultraviolet lasers utilizing relativistic Doppler tuning at FAIR facilities in Darmstadt.

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I. INTRODUCTION

Measurements of parity-nonconservation (PNC) effects with heavy few-electron ions can provide new opportunities for tests of the Standard Model in the low-energy regime. This is mainly because, in contrast to neutral atoms (see Refs. [1–10]), in highly charged ions the electron-correlation effects, being suppressed by a factor $1/Z$ (where Z is the nuclear charge number), can be accounted for by perturbation theory to a very high accuracy. The simple atomic structure of such ions allows also one to calculate the QED contributions to the required accuracy.

PNC experiments with highly charged ions were first discussed in Ref. [11]. There it was proposed to use close opposite-parity levels $2\,^1S_0$ and 2^3P_1 in He-like ions for $Z \approx 6$ and $Z \approx 29$, where the PNC effect is strongly enhanced. Later, various scenarios for PNC experiments with heavy H- and He-like ions were considered in a number of papers [12–21]. In particular, in Ref. [13] it was proposed to study the induced $2\,^3S_1$–$2\,^1S_0$ transition in He-like ions with $Z \approx 6$ in the presence of electric and magnetic fields. Possibilities of investigating PNC effects in H-like ions at high-energy ion storage rings utilizing relativistic Doppler tuning and laser cooling were considered in Ref. [17]. Most of the works [12,14,16,18,21] exploited, however, the near degeneracy of the $2\,^1S_0$ and $2\,^3P_0$ states in He-like ions at $Z \approx 64$ and $Z \approx 90$. For overviews of the schemes suggested we refer to Refs. [20,21].

In the present paper, we evaluate the PNC effect on the laser-induced $2\,^3S_1$–$2\,^1S_0$ transition in heavy heliumlike ions nearby $Z = 64$ (transition energy of about 114 eV) and $Z = 90$ (transition energy of about 240 eV), where the PNC effect is strongly enhanced. Such experiments seem to be feasible in the near future in view of recent developments in high-energy lasers for heavy-ion experiments (PHELIX project) [22,23]. As an alternative, one may consider employment of relativistic Doppler tuning at FAIR facilities in Darmstadt [24,25]. With ion energies up to 10.7 GeV/u, as anticipated at the FAIR facilities, the Doppler effect can be utilized for tuning ultraviolet laser light with photon energies in the range from 4 to 10 eV in resonance with the transition energies under consideration.

The paper is organized as follows. In Sec. II, the basic formulas for the $2\,^3S_1$–$2\,^1S_0$ transition amplitude are presented. The admixture of the opposite-parity states $2\,^1S_0$ and $2\,^3P_0$ is taken into account and, as a result, the PNC correction to the cross section is derived. It is shown that accounting for the first-order interelectronic-interaction and QED corrections in the velocity gauge can be easily done within the zeroth-order approximation in the length gauge. In Sec. III, numerical results for the PNC correction in heliumlike gadolinium and thorium are presented and possible scenarios for experiments are discussed.

Relativistic units ($\hbar = c = 1$) and the Heaviside charge unit [$\alpha = e^2/(4\pi)$, $e < 0$] are used throughout the paper.

II. BASIC FORMULAS

We consider the absorption of a photon with energy $\omega \approx E_{2\,^1S_0} - E_{2\,^3S_1}$ and circular polarization $\lambda = \pm 1$ by a heavy heliumlike ion being initially prepared in the $2\,^3S_1$ state. If the weak electron-nucleus interaction is ignored, the absorption cross section is completely determined by the magnetic-dipole transition amplitude. For such a transition the interelectronic-interaction effects are suppressed by a factor $1/Z$ and, to zeroth order, we assume that the electrons interact only with the Coulomb field of the nucleus. Then, the wave functions of the initial ($2\,^3S_1$) and the final ($2\,^1S_0$) state are given by

$$u_{JM}(\mathbf{x}_1,\mathbf{x}_2) = \frac{1}{\sqrt{2}} \sum_{m_1,m_2} C^{JM}_{j_1 m_1, j_2 m_2} \big[\psi_{j_1 m_1}(\mathbf{x}_1)\psi_{j_2 m_2}(\mathbf{x}_2) - \psi_{j_1 m_1}(\mathbf{x}_2)\psi_{j_2 m_2}(\mathbf{x}_1)\big], \tag{1}$$

where $\psi_{j_1 m_1}(\mathbf{x})$ is the one-electron $1s$ wave function, $\psi_{j_2 m_2}(\mathbf{x})$ is the one-electron $2s$ wave function, and $C^{JM}_{j_1 m_1, j_2 m_2}$ is the Clebsch-Gordan coefficient. In what follows, we assume that the laser spectral width and the width due to a finite ion-laser interaction time can be neglected. If, for a moment, we further neglect the width of the initial state, the cross section in the resonant approximation is given by (see, e.g., Refs. [26–28])

$$\sigma = (2\pi)^3 \frac{\Gamma_b |\langle b|[R(1)+R(2)]|a\rangle|^2}{(E_a+\omega-E_b)^2+\Gamma_b^2/4}. \quad (2)$$

Here $|a\rangle \equiv |2\,^3S_1\rangle$ and $|b\rangle \equiv |2\,^1S_0\rangle$ are the initial and final states, respectively, Γ_b is the width of the final state, and $R(i)$ is the transition operator acting on variables of the ith electron. In the transverse gauge, $R = -e\boldsymbol{\alpha}\cdot\mathbf{A}$, where

$$\mathbf{A}(\mathbf{x}) = \frac{\boldsymbol{\epsilon}\exp(i\mathbf{k}\cdot\mathbf{x})}{\sqrt{2\omega(2\pi)^3}} \quad (3)$$

is the wave function of the absorbed photon and $\boldsymbol{\alpha}$ is the vector incorporating the Dirac matrices. In order to account for the width of the initial state a in Eq. (2), we simply replace $E_a \to E$ and

$$|a\rangle\langle a| \to \int dE \frac{\Gamma_a/(2\pi)}{(E-E_a)^2+\Gamma_a^2/4}|a\rangle\langle a|. \quad (4)$$

In a more rigorous approach, one should consider the preparation of the state a as a part of the whole process [27]. With the substitution (4), we get

$$\sigma = (2\pi)^3 \times \int dE \frac{\Gamma_b\Gamma_a|\langle b|[R(1)+R(2)]|a\rangle|^2}{2\pi\left[(E+\omega-E_b)^2+\Gamma_b^2/4\right]\left[(E-E_a)^2+\Gamma_a^2/4\right]}. \quad (5)$$

Integrating over E, we obtain

$$\sigma = (2\pi)^3 \frac{\Gamma_a+\Gamma_b}{[\omega-(E_b-E_a)]^2+(\Gamma_a+\Gamma_b)^2/4} \times |\langle b|[R(1)+R(2)]|a\rangle|^2. \quad (6)$$

In the resonance case, $\omega = E_b - E_a$, we have

$$\sigma = 4(2\pi)^3 \frac{|\langle b|[R(1)+R(2)]|a\rangle|^2}{\Gamma_a+\Gamma_b}. \quad (7)$$

Finally, averaging over the angular momentum projection of the initial state, we obtain

$$\sigma = 4(2\pi)^3 \frac{1}{2J_a+1}\sum_{M_a}\frac{|\langle b|[R(1)+R(2)]|a\rangle|^2}{\Gamma_a+\Gamma_b}. \quad (8)$$

In what follows, because of the smallness of the transition energy, we can write

$$R = -e\boldsymbol{\alpha}\cdot\mathbf{A} = -e(\boldsymbol{\alpha}\cdot\boldsymbol{\epsilon})\frac{\exp(i\mathbf{k}\cdot\mathbf{x})}{\sqrt{2\omega(2\pi)^3}} \approx -e\frac{(\boldsymbol{\alpha}\cdot\boldsymbol{\epsilon})}{\sqrt{2\omega(2\pi)^3}}(1+i\mathbf{k}\cdot\mathbf{x}). \quad (9)$$

For the transition $J_a = 1 \to J_b = 0$ we can restrict ourselves to the dipole approximation and, therefore, represent the transition operator as the sum

$$R = R_{\rm e} + R_{\rm m}, \quad (10)$$

where

$$R_{\rm e} = -e\frac{(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})}{\sqrt{2\omega(2\pi)^3}} \quad (11)$$

is the electric-dipole transition operator in the velocity gauge,

$$R_{\rm m} = i\frac{([\boldsymbol{\epsilon}\times\mathbf{k}]\cdot\boldsymbol{\mu})}{\sqrt{2\omega(2\pi)^3}} \quad (12)$$

is the magnetic-dipole transition operator, and $\boldsymbol{\mu} = (e/2)[\mathbf{x}\times\boldsymbol{\alpha}]$ is the operator of the magnetic moment of electron. If we neglect the weak interaction, the $2\,^3S_1$–$2\,^1S_0$ transition amplitude is the pure magnetic-dipole one. Then, evaluating the matrix elements in Eq. (8), we obtain

$$\sigma_0^{(2\,^3S_1\to 2\,^1S_0)} = \frac{1}{9}\frac{\omega}{\Gamma_{2\,^3S_1}+\Gamma_{2\,^1S_0}}|\langle 2s||\boldsymbol{\mu}||2s\rangle - \langle 1s||\boldsymbol{\mu}||1s\rangle|^2, \quad (13)$$

where $\langle ns||\boldsymbol{\mu}||ns\rangle$ is the reduced matrix element of the magnetic-dipole-moment operator and the subscript "0" stays for the zeroth-order approximation.

To account for the weak interaction we have to first modify the wave function of the $2\,^1S_0$ state due to the admixture of the $2\,^3P_0$ state:

$$|2\,^1S_0\rangle \to |2\,^1S_0\rangle + \frac{\langle 2\,^3P_0|[H_W(1)+H_W(2)]|2\,^1S_0\rangle}{E_{2\,^1S_0}-E_{2\,^3P_0}}|2\,^3P_0\rangle. \quad (14)$$

Here

$$H_W = -(G_F/\sqrt{8})Q_W\rho_N(r)\gamma_5 \quad (15)$$

is the spin-independent part of the effective nuclear weak-interaction Hamiltonian [29]. G_F denotes the Fermi constant, $Q_W \approx -N + Z(1-4\sin^2\theta_W)$ is the weak charge of the nucleus (which is related to the Weinberg angle θ_W), γ_5 is the Dirac matrix, and ρ_N is the effective nuclear weak-charge density normalized to unity. A simple evaluation of the weak-interaction matrix element yields

$$\langle 2\,^3P_0|[H_W(1)+H_W(2)]|2\,^1S_0\rangle = \langle 2p_{1/2}|H_W|2s\rangle = i\frac{G_F}{2\sqrt{2}}Q_W\int_0^\infty dr\, r^2\rho_N(r)\left[g_{2p_{1/2}}f_{2s} - f_{2p_{1/2}}g_{2s}\right], \quad (16)$$

where the large and small radial components of the Dirac wave function are defined by

$$\psi_{n\kappa m}(\mathbf{r}) = \begin{pmatrix} g_{n\kappa}(r)\Omega_{\kappa m}(\mathbf{n}) \\ if_{n\kappa}(r)\Omega_{-\kappa m}(\mathbf{n})\end{pmatrix} \quad (17)$$

and $\kappa = (-1)^{j+l+1/2}(j+1/2)$ is the Dirac quantum number. Then formula (14) can be written as

$$|2\,^1S_0\rangle \to |2\,^1S_0\rangle + i\xi|2\,^3P_0\rangle, \quad (18)$$

where

$$\xi = \frac{G_F}{2\sqrt{2}}\frac{Q_W}{E_{2\,^1S_0}-E_{2\,^3P_0}} \times \int_0^\infty dr\, r^2\rho_N(r)\left[g_{2p_{1/2}}f_{2s} - f_{2p_{1/2}}g_{2s}\right]. \quad (19)$$

The admixture of the $2\,^3P_0$ state enables the $2\,^3S_1$–$2\,^3P_0$ transition, which is determined by the electric-dipole amplitude. Since the electric-dipole transition operator depends on the gauge employed, the results may differ in the different gauges,

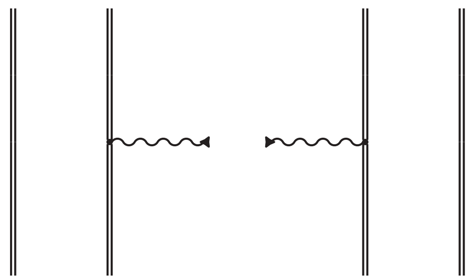

FIG. 1. The absorption of a photon by a heliumlike ion to the zeroth-order approximation (noninteracting electrons).

if the calculations are restricted to a given approximation. The difference can be especially large for a transition between the states having the same (or close) zeroth-order energies, as in the case under consideration. In Ref. [30] it was shown that using the length gauge in the calculation of the zeroth-order $2s$–$2p_{1/2}$ transition amplitude in H-like ions is equivalent to accounting for the one-loop QED corrections in the velocity-gauge calculation, provided the transition energy in the length-gauge calculation includes the corresponding corrections. Let us show that accounting for the one-photon exchange and one-loop QED corrections to the $2\,^3S_1$–$2\,^3P_0$ transition amplitude in the velocity gauge can be performed equivalently within the zeroth-order approximation in the length gauge by employing the transition energy which includes the related corrections.

With this in mind, we consider first the evaluation of the $2\,^3S_1$–$2\,^3P_0$ transition amplitude in the velocity gauge to zeroth and first orders in $1/Z$. The corresponding diagrams are presented in Figs. 1 and 2, respectively. Formal expressions for these diagrams can be derived by using the two-time Green function method [27]. Such a derivation was considered in detail in Ref. [31]. To simplify the analysis, we consider

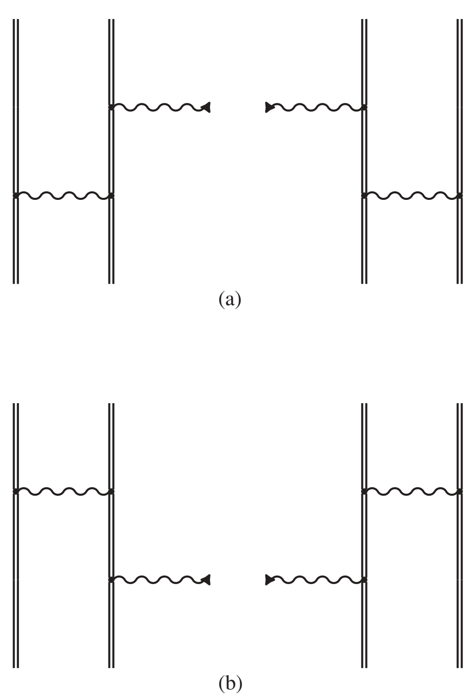

FIG. 2. One-photon exchange corrections to the absorption of a photon by a heliumlike ion.

the matrix elements of the electric-dipole transition operator between the one-determinant wave functions,

$$u_a(\mathbf{x}_1,\mathbf{x}_2) = \frac{1}{\sqrt{2}}\sum_P (-1)^P \psi_{Pa_1}(\mathbf{x}_1)\psi_{Pa_2}(\mathbf{x}_2), \qquad (20)$$

$$u_b(\mathbf{x}_1,\mathbf{x}_2) = \frac{1}{\sqrt{2}}\sum_P (-1)^P \psi_{Pb_1}(\mathbf{x}_1)\psi_{Pb_2}(\mathbf{x}_2), \qquad (21)$$

where it is assumed that $a_1 = b_1 = 1s$, $a_2 = 2s$, $b_2 = 2p_{1/2}$, P is the permutation operator, and $(-1)^P$ is the sign of the permutation. Then, to zeroth order we obtain for the transition amplitude:

$$\begin{aligned}\tau^{(0)} &= -\langle b|[R_{\mathrm{e}}(1) + R_{\mathrm{e}}(2)]|a\rangle \\ &= -\langle b_1|R_{\mathrm{e}}(1)|a_1\rangle\delta_{a_2\,b_2} - \langle b_2|R_{\mathrm{e}}(2)|a_2\rangle\delta_{a_1\,b_1} \\ &= -\langle 2p_{1/2}|R_{\mathrm{e}}|2s\rangle. \end{aligned} \qquad (22)$$

Employing the identity

$$\boldsymbol{\alpha} = i[H,\mathbf{r}], \qquad (23)$$

where H is the one-electron Dirac Hamiltonian, one obtains

$$\tau^{(0)} = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{r})|2s\rangle(\varepsilon_{2p_{1/2}} - \varepsilon_{2s}), \qquad (24)$$

where ε_{2s} and $\varepsilon_{2p_{1/2}}$ are the one-electron Dirac energies of the $2s$ and $2p_{1/2}$ states, respectively. In particular, it follows that for the pure Coulomb field ($\varepsilon_{2s} = \varepsilon_{2p_{1/2}}$) in the velocity gauge the zeroth-order $2\,^3S_1$–$2\,^3P_0$ transition amplitude is equal to zero.

The interelectronic-interaction corrections, defined by the diagrams depicted in Fig. 2, consist of irreducible and reducible contributions [27,31]. Since, to good accuracy, these corrections can be treated with the pure Coulomb field of the nucleus, in what follows, we restrict our consideration to this approximation. Then, according to Ref. [31] we find that for the $2\,^3S_1$–$2\,^3P_0$ transition the reducible contribution vanishes. As for the irreducible contribution, it can be expressed as the sum [31]

$$\tau_{\mathrm{irr}} = \tau_{\mathrm{irr}}^{(a)} + \tau_{\mathrm{irr}}^{(b)}, \qquad (25)$$

where

$$\begin{aligned}\tau_{\mathrm{irr}}^{(a)} = \frac{e}{\sqrt{2\omega(2\pi)^3}}\sum_P(-1)^P \Bigg\{ &\sum_n^{\varepsilon_{Pb_2}+\varepsilon_n \neq E_a^{(0)}} \langle Pb_1|(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})|n\rangle \\ &\times \frac{1}{E_a^{(0)} - \varepsilon_{Pb_2} - \varepsilon_n}\langle nPb_2|I\big(\varepsilon_{Pb_2} - \varepsilon_{a_2}\big)|a_1a_2\rangle \\ &+ \sum_n^{\varepsilon_{Pb_1}+\varepsilon_n \neq E_a^{(0)}} \langle Pb_2|(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})|n\rangle \frac{1}{E_a^{(0)} - \varepsilon_{Pb_1} - \varepsilon_n} \\ &\times \langle Pb_1n|I\big(\varepsilon_{Pb_1} - \varepsilon_{a_1}\big)|a_1a_2\rangle \Bigg\}, \end{aligned} \qquad (26)$$

$$\tau_{\mathrm{irr}}^{(b)} = \frac{e}{\sqrt{2\omega(2\pi)^3}}$$

$$\times \sum_P(-1)^P \Bigg\{ \sum_n^{\varepsilon_{a_2}+\varepsilon_n \neq E_b^{(0)}} \langle Pb_1Pb_2|I\big(\varepsilon_{Pb_2} - \varepsilon_{a_2}\big)|na_2\rangle$$

$$\times \frac{1}{E_b^{(0)} - \varepsilon_{a_2} - \varepsilon_n} \langle n|(\boldsymbol{\epsilon} \cdot \boldsymbol{\alpha})|a_1\rangle + \sum_n^{\varepsilon_{a_1}+\varepsilon_n \neq E_b^{(0)}} \langle Pb_1 Pb_2|I(\varepsilon_{Pb_1} - \varepsilon_{a_1})|a_1 n\rangle \times \frac{1}{E_b^{(0)} - \varepsilon_{a_1} - \varepsilon_n} \langle n|(\boldsymbol{\epsilon} \cdot \boldsymbol{\alpha})|a_2\rangle \Bigg\}. \quad (27)$$

Here $I(\omega) = e^2\alpha^\mu\alpha^\nu D_{\mu\nu}(\omega)$,

$$D_{\rho\sigma}(\omega,\mathbf{x}-\mathbf{y}) = -g_{\rho\sigma}\int \frac{d\mathbf{k}}{(2\pi)^3}\frac{\exp[i\mathbf{k}\cdot(\mathbf{x}-\mathbf{y})]}{\omega^2-\mathbf{k}^2+i0} \quad (28)$$

is the photon propagator in the Feynman gauge, $\alpha^\rho \equiv \gamma^0\gamma^\rho = (1,\boldsymbol{\alpha})$, $E_a^{(0)} = \varepsilon_{a_1} + \varepsilon_{a_2}$, and $E_b^{(0)} = \varepsilon_{b_1} + \varepsilon_{b_2}$. Taking into account that $E_a^{(0)} = E_b^{(0)}$ and using the identity (23), we get

$$\tau_{\rm irr}^{(a)} = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\sum_P (-1)^P \Bigg\{ \sum_n^{\varepsilon_n \neq \varepsilon_{Pb_1}} \langle Pb_1|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n\rangle \times \langle nPb_2|I(\varepsilon_{Pb_2}-\varepsilon_{a_2})|a_1a_2\rangle + \sum_n^{\varepsilon_n\neq\varepsilon_{Pb_2}} \langle Pb_2|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n\rangle\langle Pb_1 n|I(\varepsilon_{Pb_1}-\varepsilon_{a_1})|a_1a_2\rangle \Bigg\}, \quad (29)$$

$$\tau_{\rm irr}^{(b)} = -i\frac{e}{\sqrt{2\omega(2\pi)^3}}\sum_P (-1)^P \times \Bigg\{ \sum_n^{\varepsilon_n\neq\varepsilon_{a_1}} \langle Pb_1Pb_2|I(\varepsilon_{Pb_2}-\varepsilon_{a_2})|na_2\rangle\langle n|(\boldsymbol{\epsilon}\cdot\mathbf{r})|a_1\rangle + \sum_n^{\varepsilon_n\neq\varepsilon_{a_2}} \langle Pb_1Pb_2|I(\varepsilon_{Pb_1}-\varepsilon_{a_1})|a_1n\rangle\langle n|(\boldsymbol{\epsilon}\cdot\mathbf{r})|a_2\rangle \Bigg\}. \quad (30)$$

With the aid of the completeness condition

$$\sum_n |n\rangle\langle n| = 1, \quad (31)$$

we find for the sum of the expressions (29) and (30)

$$\tau_{\rm irr} = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{r})|2s\rangle(\Delta E_b - \Delta E_a), \quad (32)$$

where

$$\Delta E_a = \sum_P (-1)^P \langle Pa_1Pa_2|I(\varepsilon_{Pa_1}-\varepsilon_{a_1})|a_1a_2\rangle, \quad (33)$$

$$\Delta E_b = \sum_P (-1)^P \langle Pb_1Pb_2|I(\varepsilon_{Pb_1}-\varepsilon_{b_1})|b_1b_2\rangle \quad (34)$$

are the first-order interelectronic-interaction corrections to the initial and final states, respectively. The same relation holds if one includes the one-loop QED corrections. The corresponding proof, which was given first in Ref. [30], is presented in the Appendix. Summing the zeroth- and first-order contributions yields

$$\tau = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{r})|2s\rangle(E_b - E_a) = i\frac{\omega}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{d})|2s\rangle, \quad (35)$$

where $\mathbf{d} = e\mathbf{r}$ is the operator of the electric-dipole moment and E_a and E_b are the total binding energies of the initial and final states, respectively. It is evident that similar equations can be derived involving two-electron wave functions (1). Consequently, in what follows, we will use the electric-dipole transition operator in the length gauge

$$R_{\rm e}^{(l)} = -i\frac{\omega(\boldsymbol{\epsilon}\cdot\mathbf{d})}{\sqrt{2\omega(2\pi)^3}}. \quad (36)$$

Substituting the two-electron wave function (18) into Eq. (8) and performing the calculation, we obtain for the PNC contribution to the cross section

$$\sigma_{\rm PNC}^{(2\,^3S_1\to 2\,^1S_0)} = \frac{1}{9}\frac{\omega}{\Gamma_{2\,^3S_1}+\Gamma_{2\,^1S_0}}2\lambda\xi\langle 2p_{1/2}||\mathbf{d}||2s\rangle(\langle 2s||\boldsymbol{\mu}||2s\rangle - \langle 1s||\boldsymbol{\mu}||1s\rangle), \quad (37)$$

where $\langle 2p_{1/2}||\mathbf{d}||2s\rangle$ is the reduced matrix element of the electric-dipole-moment operator and $\lambda = \pm 1$ is the photon polarization. Integrating over the angular variables in the reduced matrix elements yields

$$\langle ns||\boldsymbol{\mu}||ns\rangle = -2e\sqrt{2/3}\int_0^\infty dr\, r^3 g_{ns}(r) f_{ns}(r), \quad (38)$$

$$\langle np_{1/2}||\mathbf{d}||ns\rangle = -e\sqrt{2/3}\int_0^\infty dr\, r^3\big[g_{np_{1/2}}(r)g_{ns}(r) + f_{np_{1/2}}(r)f_{ns}(r)\big]. \quad (39)$$

These integrals are easily evaluated by employing the virial relations for the Dirac equation (see, e.g., Refs. [32–34]). For the case of interest here, one derives

$$2\int_0^\infty dr\, r^3\,[g_{2s}(r)f_{2s}(r) - g_{1s}(r)f_{1s}(r)] = \gamma - \sqrt{(1+\gamma)/2}, \quad (40)$$

$$\int_0^\infty dr\, r^3\big[g_{2p_{1/2}}(r)g_{2s}(r) + f_{2p_{1/2}}(r)f_{2s}(r)\big] = \frac{3(1+\gamma)\sqrt{1+2\gamma}}{2\alpha Z}, \quad (41)$$

where $\gamma = \sqrt{1-(\alpha Z)^2}$. By substituting these expressions into Eq. (13) and Eq. (37) leads to

$$\sigma_0^{(2\,^3S_1\to 2\,^1S_0)} = \frac{2}{27}\frac{\pi\alpha\omega}{\Gamma_{2\,^3S_1}+\Gamma_{2\,^1S_0}}[\sqrt{2(1+\gamma)} - 2\gamma]^2, \quad (42)$$

$$\sigma^{(2\,^3S_1\to 2\,^1S_0)} = \sigma_0^{(2\,^3S_1\to 2\,^1S_0)} + \sigma_{\rm PNC}^{(2\,^3S_1\to 2\,^1S_0)} = (1+\lambda\varepsilon)\sigma_0^{(2\,^3S_1\to 2\,^1S_0)}, \quad (43)$$

TABLE I. The zeroth-order cross section $\sigma_0^{(2\,^3S_1\to 2\,^1S_0)}$, the PNC correction $\sigma_{\rm PNC}^{(2\,^3S_1\to 2\,^1S_0)}$, and the parameter ε, defined by Eq. (43), for the laser-induced $2\,^3S_1 \to 2\,^1S_0$ transition in He-like Gd and Th. The $2\,^3S_1$–$2\,^1S_0$ transition energies are taken from Ref. [35], while the $2\,^3P_0$–$2\,^1S_0$ energy difference is chosen as discussed in the text.

Ion	$(E_{2\,^3P_0}-E_{2\,^1S_0})$ (eV)	$(E_{2\,^1S_0}-E_{2\,^3S_1})$ (eV)	$\sigma_0^{(2\,^3S_1\to 2\,^1S_0)}$ (b)	$\sigma_{\rm PNC}^{(2\,^3S_1\to 2\,^1S_0)}$ (b)	ε
^{158}Gd^{62+}	0.074(74)	114.0	4084.1	±2.1	−0.000 51
^{232}Th^{88+}	0.44(40)	240.1	1217.6	±0.6	−0.000 53

where

$$\varepsilon = 2\xi \frac{\langle 2p_{1/2}||\mathbf{d}||2s\rangle}{\langle 2s||\boldsymbol{\mu}||2s\rangle - \langle 1s||\boldsymbol{\mu}||1s\rangle} = -2\xi \frac{3(1+\gamma)\sqrt{1+2\gamma}}{\alpha Z(\sqrt{2(1+\gamma)}-2\gamma)} \quad (44)$$

is a parameter which characterizes the relative value of the PNC effect. The second term on the right-hand side of Eq. (43) represents the PNC contribution, which changes sign under the replacement $\lambda \to -\lambda$. The PNC parameter can also be represented as

$$\varepsilon = -2\xi \sqrt{\frac{\Gamma_{2\,^3S_1}+\Gamma_{2\,^3P_0}}{\Gamma_{2\,^3S_1}+\Gamma_{2\,^1S_0}} \frac{\sigma_0^{(2\,^3S_1\to 2\,^3P_0)}}{\sigma_0^{(2\,^3S_1\to 2\,^1S_0)}}}, \quad (45)$$

where $\sigma_0^{(2\,^3S_1\to 2\,^3P_0)}$ is the cross section of the resonant absorption into the $2\,^3P_0$ state and $\Gamma_{2\,^3P_0}$ is the total width of this state.

III. RESULTS AND DISCUSSION

The formulas (42)–(44) allow one to evaluate the cross section and the corresponding PNC effect. The most promising situation for observing the PNC effect occurs in cases where the levels $2\,^1S_0$ and $2\,^3P_0$ are almost degenerate. According to the most to-date elaborated calculations [35–38] (see also the related table in Ref. [21]) such cases are gadolinium ($Z = 64$) and thorium ($Z = 90$), where the levels $2\,^1S_0$ and $2\,^3P_0$ are nearly crossing. For Gd the energy interval amounts to −0.023(74) eV, while in case of Th it is 0.44(40) eV [38]. Since the uncertainties are comparable to the energy differences, to estimate the PNC effect we take the values 0.074 eV and 0.44 eV for the $2\,^3P_0$–$2\,^1S_0$ energy difference in Gd and Th, respectively. The widths of the $2\,^3S_1$ and $2\,^1S_0$ states, which enter formula (42), are mainly defined by the one-photon $M1$ and two-photon $E1E1$ transitions, respectively. We evaluate the decay rate of the $M1$ transition $2\,^3S_1 \to 1\,^1S_0$ employing the transition energy taken from Ref. [35]. The interelectronic-interaction corrections to the transition amplitude are calculated to first order in $1/Z$ within a systematic QED approach (for details see Ref. [31]). As the result, we obtain decay rates $w_{M1}^{(2\,^3S_1\to 1\,^1S_0)} = 2.301 \times 10^{12}$ s^{-1} for Gd and $w_{M1}^{(2\,^3S_1\to 1\,^1S_0)} = 9.470 \times 10^{13}$ s^{-1} for Th. These values are in fair agreement with those from Refs. [31,39,40]. The two-photon decays $2\,^3S_1 \to 1\,^1S_0$ and $2\,^1S_0 \to 1\,^1S_0$ are calculated in the length gauge with the transition energies taken from Ref. [35]. The interelectronic-interaction effects are approximately accounted for by means of a Kohn-Sham potential. The calculated transition rates are $w_{2\gamma}^{(2\,^3S_1\to 1\,^1S_0)} = 8.74 \times 10^8$ s^{-1} and $w_{2\gamma}^{(2\,^1S_0\to 1\,^1S_0)} = 9.04 \times 10^{11}$ s^{-1} for Gd and $w_{2\gamma}^{(2\,^3S_1\to 1\,^1S_0)} = 2.07 \times 10^{10}$ s^{-1} and $w_{2\gamma}^{(2\,^1S_0\to 1\,^1S_0)} = 6.25 \times 10^{12}$ s^{-1} for Th. These values together with the $E1E1$ channel also include higher multipole contributions, such as $M1M1$ etc. For Th the dominant $E1E1$ decay channel yields $w_{E1E1}^{(2\,^3S_1\to 1\,^1S_0)} = 1.62 \times 10^{10}$ s^{-1} and $w_{E1E1}^{(2\,^1S_0\to 1\,^1S_0)} = 6.25 \times 10^{12}$ s^{-1}. It is worth noticing that for the $2\,^3S_1$ state the higher multipoles contribute up to 20% to the total two-photon decay rate. Comparing the $E1E1$ decay rates with the results of Ref. [41], we find excellent agreement for the $2\,^3S_1$ state and a slight deviation for the $2\,^1S_0$ state, which is mainly due to employing the more accurate transition energies in our calculations. Finally, the total widths are $\Gamma_{2\,^3S_1} = 1.515$ meV and $\Gamma_{2\,^1S_0} = 0.595$ meV for Gd and $\Gamma_{2\,^3S_1} = 62.35$ meV and $\Gamma_{2\,^1S_0} = 4.11$ meV for Th. The results of the calculations of the PNC effect by formulas (42)–(44) for Gd and Th are presented in Table I.

As one can see from the table, in both Gd and Th cases the PNC effect amounts to about 0.05%, which is a rather large value for parity-violation experiments. Moreover, one may expect that the PNC effect can be further increased, at least, by an order of magnitude by choosing proper isotopes, provided the $2\,^1S_0$–$2\,^3P_0$ energy difference is known to a higher accuracy. With the current experimental techniques [42], accurate measurements of the difference considered seem feasible.

Because of a large transition energy (>100 eV), until recently experimental scenarios with laser-induced $2\,^3S_1$–$2\,^1S_0$ transition in heavy He-like ions were far from being possible. However, the situation has changed in view of the very significant progress in x-ray laser development. Such lasers will be available in the near future with a high repetition rate [43]. Already now, there is such an x-ray laser available at the heavy-ion facility GSI (PHELIX facility) where photon energies of up to 200 eV have been reached [22,23]. As an alternative scenario, the excitation energy can be obtained by counter-propagating the ultraviolet laser beam with the photon energy in the range from 4 to 10 eV and the He-like ion beam with energy up to 10.7 GeV/u, which will be available at the FAIR facility in Darmstadt [24,25]. The population of the $2\,^1S_0$ level can be measured by observing the $E1E1$ decay to the ground state. In the second scenario, due to the strong Lorentz boosting, the decay photons are emitted at the forward direction, which considerably simplifies their detection.

The next problem to be addressed is the preparation of ions in the $2\,^3S_1$ state that is required in both scenarios considered.

As follows from the study presented in Refs. [44,45], in collisions with gas atoms one can produce selectively both the $2\,^1S_0$ state and the $2\,^3S_1$ one [46]. However, it would be of great importance to populate exclusively only the $2\,^3S_1$ state. The only way to accomplish this is to first form the doubly excited $(2s2p_{1/2})_0$ state via dielectronic recombination of an electron with a H-like ion. Since the main decay channel of the $(2s2p_{1/2})_0$ state is the transition into the $2\,^3S_1$ state, this enables selective production of ions in the $2\,^3S_1$ state.

The PNC effect is to be measured by counting the intensity difference in the $E1E1$ decay of the $2\,^1S_0$ state for polarizations $\lambda = \pm 1$. The background emission can be separated by switching off the laser light. Changing the photon energy allows one to eliminate the interference with a nonresonant transition via the $2\,^3P_0$ state, which could also be evaluated to good accuracy if necessary. Moreover, since the $E1E1$ emission can be measured relative to the intensity of the $M1$ x-ray line (decay of the $2\,^3S_1$ state), such an experimental scenario appears to be quite realistic.

IV. CONCLUSION

In this paper we have studied the PNC effect with laser-induced $2\,^3S_1$–$2\,^1S_0$ transition in heavy heliumlike ions. A simple analytical formula for the photon-absorption cross section derived enables easy evaluation of the PNC effect for ions near $Z = 64$ and $Z = 90$, where the effect is strongly enhanced due to the near crossing of the opposite-parity $2\,^1S_0$ and $2\,^3P_0$ levels. The calculations performed showed that the effect can amount to about 0.05% and even more for the ions of interest. Prospects for the corresponding PNC experiments have been discussed. It is found that the desired photon energy can be achieved either by x-ray lasers that are presently getting developed at GSI (PHELIX project) as well as at the Helmholtz-Institute in Jena [47] or by counter-propagating the ultraviolet laser beam and the He-like ion beam at the FAIR facility in Darmstadt.

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APPENDIX: QED CORRECTIONS TO THE TRANSITION AMPLITUDE

The one-electron QED corrections to the $2\,^3S_1$–$2\,^3P_0$ transition amplitude are determined by the corresponding contributions to the $2s$–$2p_{1/2}$ amplitude in one-electron ions as defined by the diagrams shown in Fig. 3. Formal expressions for these corrections are almost the same as for the corresponding corrections to the emission amplitude [27]. Let us consider the one-loop self-energy correction. According to formulas provided in Ref. [27], it is given by the sum of the irreducible, reducible, and vertex contributions. For an electron interacting with a pure Coulomb field together with the dipole approximation $\exp(i\mathbf{k}\cdot\mathbf{x}) \to 1$, the reducible contribution vanishes. The irreducible contribution is given by

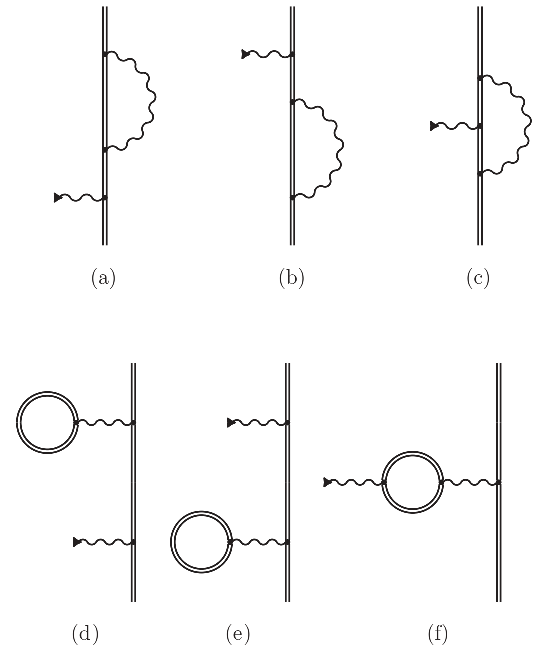

FIG. 3. One-loop QED corrections to the photon absorption.

$$\begin{aligned}\tau_{\rm irr}^{\rm (SE)} &= -\langle 2p_{1/2}|R|\xi_{2s}\rangle - \langle \xi_{2p_{1/2}}|R|2s\rangle \\ &= \frac{e}{\sqrt{2\omega(2\pi)^3}}\big[\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})|\xi_{2s}\rangle + \langle \xi_{2p_{1/2}}|(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})|2s\rangle\big],\end{aligned} \tag{A1}$$

where

$$|\xi_a\rangle = \sum_n^{n\neq a} \frac{|n\rangle\langle n|\Sigma(\varepsilon_a)|a\rangle}{\varepsilon_a - \varepsilon_n}, \tag{A2}$$

$$\langle \xi_b| = \sum_n^{n\neq b} \frac{\langle b|\Sigma(\varepsilon_b)|n\rangle\langle n|}{\varepsilon_b - \varepsilon_n}, \tag{A3}$$

and

$$\langle a|\Sigma(\varepsilon)|b\rangle = \frac{i}{2\pi}\int_{-\infty}^{\infty} d\omega \sum_n \frac{\langle an|e^2\alpha^\rho\alpha^\sigma D_{\rho\sigma}(\omega)|nb\rangle}{\varepsilon - \omega - \varepsilon_n(1-i0)}. \tag{A4}$$

By means of the identity (23) and the completeness relation (31), we obtain

$$\begin{aligned}\tau_{\rm irr}^{\rm (SE)} = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\{&\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{r})|2s\rangle(\langle 2p_{1/2}|\Sigma(\varepsilon_{2p_{1/2}})|2p_{1/2}\rangle \\ &- \langle 2s|\Sigma(\varepsilon_{2s})|2s\rangle) + \langle 2p_{1/2}|[(\boldsymbol{\epsilon}\cdot\mathbf{r}),\Sigma(\varepsilon_{2s})]|2s\rangle\}.\end{aligned} \tag{A5}$$

For the vertex contribution one derives

$$\tau_{\rm ver}^{(\rm SE)} = -e^2 \frac{i}{2\pi}\int_{-\infty}^{\infty} d\omega \int \frac{d\mathbf{k}}{(2\pi)^3}\, \frac{1}{\omega^2-\mathbf{k}^2+i0} \times \sum_{n_1,n_2} \langle 2p_{1/2}|\alpha^\rho \exp{(i\mathbf{k}\cdot\mathbf{y})}|n_1\rangle \times \frac{1}{\varepsilon_{2p_{1/2}}-\omega-\varepsilon_{n_1}(1-i0)}\langle n_1|\frac{e(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})}{\sqrt{2\omega(2\pi)^3}}|n_2\rangle \times \frac{1}{\varepsilon_{2s}-\omega-\varepsilon_{n_2}(1-i0)}\langle n_2|\alpha_\rho \exp{(-i\mathbf{k}\cdot\mathbf{x})}|2s\rangle. \tag{A6}$$

Transforming

$$\frac{1}{\varepsilon_{2p_{1/2}}-\omega-\varepsilon_{n_1}(1-i0)}\frac{1}{\varepsilon_{2s}-\omega-\varepsilon_{n_2}(1-i0)} = \frac{1}{\varepsilon_{n_1}-\varepsilon_{n_2}}\left(\frac{1}{\varepsilon_{2p_{1/2}}-\omega-\varepsilon_{n_1}(1-i0)} - \frac{1}{\varepsilon_{2s}-\omega-\varepsilon_{n_2}(1-i0)}\right), \tag{A7}$$

$$\langle n_1|(\boldsymbol{\epsilon}\cdot\boldsymbol{\alpha})|n_2\rangle = i\langle n_1|[H,(\boldsymbol{\epsilon}\cdot\mathbf{r})]|n_2\rangle = i(\varepsilon_{n_1}-\varepsilon_{n_2})\langle n_1|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n_2\rangle, \tag{A8}$$

we get

$$\tau_{\rm ver}^{(\rm SE)} = -\frac{e}{\sqrt{2\omega(2\pi)^3}}e^2\frac{i}{2\pi}\int_{-\infty}^{\infty} d\omega \int \frac{d\mathbf{k}}{(2\pi)^3}\,\frac{1}{\omega^2-\mathbf{k}^2+i0} \times \sum_{n_1,n_2}^{\varepsilon_{n_1}\neq\varepsilon_{n_2}} \langle 2p_{1/2}|\alpha^\rho \exp{(i\mathbf{k}\cdot\mathbf{y})}|n_1\rangle \times i\left(\frac{1}{\varepsilon_{2p_{1/2}}-\omega-\varepsilon_{n_1}(1-i0)} - \frac{1}{\varepsilon_{2s}-\omega-\varepsilon_{n_2}(1-i0)}\right)\langle n_1|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n_2\rangle \times \langle n_2|\alpha_\rho \exp{(-i\mathbf{k}\cdot\mathbf{x})}|2s\rangle$$
$$= -i\frac{e}{\sqrt{2\omega(2\pi)^3}}e^2\frac{i}{2\pi}\int_{-\infty}^{\infty} d\omega \int \frac{d\mathbf{k}}{(2\pi)^3}\,\frac{1}{\omega^2-\mathbf{k}^2+i0} \times \Bigg\{\sum_{n_1}\langle 2p_{1/2}|\alpha^\rho \exp{(i\mathbf{k}\cdot\mathbf{y})}|n_1\rangle \times \frac{1}{\varepsilon_{2p_{1/2}}-\omega-\varepsilon_{n_1}(1-i0)} \times \Bigg[\langle n_1|(\boldsymbol{\epsilon}\cdot\mathbf{x})\alpha_\rho \exp{(-i\mathbf{k}\cdot\mathbf{x})}|2s\rangle - \sum_{n_2}^{\varepsilon_{n_1}=\varepsilon_{n_2}}\langle n_1|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n_2\rangle\langle n_2|\alpha_\rho \exp{(-i\mathbf{k}\cdot\mathbf{x})}|2s\rangle\Bigg] - \sum_{n_2}\Bigg[\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{y})\alpha^\rho \exp{(i\mathbf{k}\cdot\mathbf{y})}|n_2\rangle - \sum_{n_1}^{\varepsilon_{n_1}=\varepsilon_{n_2}}\langle 2p_{1/2}|\alpha^\rho \exp{(i\mathbf{k}\cdot\mathbf{y})}|n_1\rangle\langle n_1|(\boldsymbol{\epsilon}\cdot\mathbf{r})|n_2\rangle\Bigg] \times \frac{1}{\varepsilon_{2s}-\omega-\varepsilon_{n_2}(1-i0)}\langle n_2|\alpha_\rho \exp{(-i\mathbf{k}\cdot\mathbf{x})}|2s\rangle\Bigg\}$$
$$= -i\frac{e}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|[(\boldsymbol{\epsilon}\cdot\mathbf{r}),\Sigma(\varepsilon_{2s})]|2s\rangle. \tag{A9}$$

The sum of both irreducible and vertex contributions yields [30]

$$\tau_{\rm tot}^{(\rm SE)} = i\frac{e}{\sqrt{2\omega(2\pi)^3}}\langle 2p_{1/2}|(\boldsymbol{\epsilon}\cdot\mathbf{r})|2s\rangle(\langle 2p_{1/2}|\Sigma(\varepsilon_{2p_{1/2}})|2p_{1/2}\rangle - \langle 2s|\Sigma(\varepsilon_{2s})|2s\rangle). \tag{A10}$$

A similar equation can be derived for the vacuum-polarization contribution.

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