Parity-nonconservation effect in the dielectronic recombination of polarized electrons with heavy He-like ions

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We investigate the parity nonconservation (PNC) effect in the dielectronic recombination of a polarized electron with a heavy He-like ion into doubly excited $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ and $[(1s2s)_0 n\kappa]_{1/2}$ states of a Li-like ion. We determine the nuclear charge number Z for which these opposite-parity levels are almost crossing and therefore the PNC effect will be significantly enhanced. Calculations are performed for quantum numbers $n \geqslant 4$ and $\kappa = \pm 1$.

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I. INTRODUCTION

Investigations of the parity-nonconservation (PNC) effects in atoms play a very important role for tests of the electroweak sector of the standard model (SM) in the low-energy regime [1–3]. In particular, the atomic experiments, being very sensitive to extensions of the standard model that predict an extra Z boson, can set a stronger restriction on the mass of the extra Z boson compared to the high-energy experiments [3]. The unprecedented experimental precision for the PNC amplitude was obtained in ^{133}Cs measurements [4,5] and, together with recent progress in QED and atomic structure calculations (see, e.g., Refs. [6–8] and references therein), provided the most accurate to date test of the SM with atomic systems. From the theoretical side, further progress in studying the PNC effect with neutral atoms is strongly limited by the uncertainties of the electron-correlation contributions. In contrast to that, in heavy, highly charged ions the correlation effects, being suppressed by a factor $1/Z$, can be calculated by perturbation theory up to the required precision. This gives good prospects for studying the PNC effects with highly charged ions.

Parity-nonconservation experiments with few-electron ions were first proposed by Gorshkov and Labzowsky in Ref. [9], where the fact that opposite-parity $2\,^1S_0$ and $2\,^3P_1$ states are almost crossing for He-like ions with $Z \sim 6$ and $Z \sim 29$ was utilized. Since that work, a number of authors considered He-like ions as very promising systems for investigating the PNC effects [10–20]. This is due to the fact that the PNC effects in He-like ions can be significantly enhanced due to the near degeneracy of some opposite-parity states. In a large number of proposals [10,13,15–20] the level crossing between the $2\,^3P_0$ and $2\,^1S_0$ states of He-like ions was exploited. One may expect that the addition of a highly excited electron would not greatly change the energy difference between the corresponding levels in Li-like ions. Indeed, the opposite-parity $[(1s2s)_0 n\kappa]_{1/2}$ and $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ states can be still made to almost cross by choosing the principal quantum number n and the Dirac angular quantum number $\kappa = (-1)^{j+l+1/2}(j+1/2)$. In this work we present such quasidegenerate levels of heavy Li-like ions and propose a scheme for observing the PNC effect in the dielectronic recombination (DR) of free electrons with He-like ions into these doubly excited states of Li-like ions.

In some previous proposals the dielectronic recombination was considered as a convenient probe process, which can be used to measure the parity violation effects. In Ref. [11], Pindzola studied the PNC effect on the Auger-electron emission from He-like uranium. The parity violation in the dielectronic recombination of polarized electrons with H-like ions at $Z < 60$ was discussed by Gribakin *et al.* in Ref. [14]. In our previous work [19] we investigated the PNC effect on the recombination of a polarized electron with unpolarized H-like thorium ($Z = 90$) and gadolinium ($Z = 64$) ions in the case of resonance with a doubly excited state of the corresponding He-like ions. In the present work we investigate the PNC effect in the process of the dielectronic recombination of polarized electrons with heavy He-like ions into the doubly excited $[(1s2s)_0 n\kappa]_{1/2}$ and $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ states of Li-like ions. The energy of the incident electron is considered to be tuned in resonance with one of these levels. The case of a nonmonoenergetic incident electron beam is also studied.

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

II. BASIC FORMALISM

We consider the process of the dielectronic recombination of an electron having an asymptotic four-momentum $p_i = (\varepsilon_i, \mathbf{p}_i)$ and polarization μ_i with a heavy He-like ion, being originally in the ground $(1s)^2$ state. As a result of this nonradiative capture, one of the nearly degenerate opposite-parity d_1 or d_2 states of the Li-like ion is formed. To simplify the derivation of formulas, we assume that these levels decay via the emission of a photon to some final state f. We suppose that the incoming electron energy ε_i is chosen to get to resonance with one of the doubly excited d_1 or d_2 states. The differential cross section of the process under consideration is defined as

[21,22]

$$\frac{d\sigma_{\mu_i}}{d\Omega} = \frac{(2\pi)^4}{v_i}\omega^2 \times \sum_{\epsilon_f}\left|\sum_{M_{d_1}}\tau_{\gamma_f,f;d_1}\frac{1}{E_i - E_{d_1} + i\Gamma_{d_1}/2}\langle\Psi_{d_1}|I|\Psi_i\rangle. + \sum_{M_{d_2}}\tau_{\gamma_f,f;d_2}\frac{1}{E_i - E_{d_2} + i\Gamma_{d_2}/2}\langle\Psi_{d_2}|I|\Psi_i\rangle\right|^2, \quad (1)$$

where E_{d_k}, Γ_{d_k}, and M_{d_k} are the energy, the total width, and the momentum projection of the d_k state ($k = 1,2$), respectively, $E_i = E_{(1s)^2} + \varepsilon_i$ is the total energy of the initial state of the system, and v_i is the velocity of the incident electron. The outgoing photon γ_f is characterized by the energy ω and the polarization $\boldsymbol{\epsilon}_f$. In addition, $\tau_{\gamma_f,f;d_k}$ is the amplitude of the radiative transition from the d_k state to the f state via the emission of a photon and I is the operator of the interelectronic interaction as defined in Ref. [22].

As mentioned above, for heavy few-electron ions the interelectronic-interaction effects are suppressed by a factor $1/Z$, compared to the interaction of the electrons with the Coulomb field of the nucleus. Therefore, we can generally consider the wave functions of our system in the independent-electron approximation. With this approximation, the initial-state wave function is given by

$$\Psi_{p_i\mu_i,JM}(\mathbf{x}_1,\mathbf{x}_2,\mathbf{x}_3) = A_N\sum_{\mathcal{P}}(-1)^{\mathcal{P}}\mathcal{P}\sum_{m_1m_2}C^{JM}_{j_1m_1,j_2m_2} \times\psi_{n_1\kappa_1m_1}(\mathbf{x}_1)\psi_{n_2\kappa_2m_2}(\mathbf{x}_2)\psi_{p_i\mu_i}(\mathbf{x}_3), \quad (2)$$

where $\psi_{n\kappa m}(\mathbf{x})$ is the one-electron bound-state Dirac wave function, $\psi_{p_i\mu_i}(\mathbf{x})$ is the incident electron wave function, $C^{JM}_{j_1m_1,j_2m_2}$ is the Clebsch-Gordan coefficient, $(-1)^{\mathcal{P}}$ is the parity of the permutation, $\mathcal{P}$ is the permutation operator, and A_N is the normalization factor. From the theoretical viewpoint, it is convenient to formulate the electron capture in the ion rest frame. In this frame we can adopt that the quantization axis (z axis) is directed along the incoming electron momentum $\mathbf{p}_i$. In this case the full expansion of the incoming electron wave function is given by (see, e.g., Refs. [23,24])

$$\psi_{p_i\mu_i}(\mathbf{x}) = \frac{1}{\sqrt{4\pi}}\frac{1}{\sqrt{p_i\varepsilon_i}}\sum_{\kappa} i^l\exp(i\Delta_\kappa) \times\sqrt{2l+1}C^{j\mu_i}_{l0,1/2\mu_i}\psi_{\varepsilon_i\kappa\mu_i}(\mathbf{x}), \quad (3)$$

where Δ_κ is the Coulomb phase shift and $\psi_{\varepsilon_i\kappa\mu_i}(\mathbf{x})$ is the partial electron wave with the Dirac quantum number $\kappa = (-1)^{j+l+1/2}(j+1/2)$, determined by the angular momentum j and the parity l.

Neglecting the weak interaction, we can write the wave functions of the intermediate d and final f states as

$$\Psi_{J(J')M}(\mathbf{x}_1,\mathbf{x}_2,\mathbf{x}_3) = A_N\sum_{\mathcal{P}}(-1)^{\mathcal{P}}\mathcal{P}\sum_{M'm_3}\sum_{m_1m_2}C^{JM}_{J'M',j_3m_3}C^{J'M'}_{j_1m_1,j_2m_2} \times\psi_{n_1\kappa_1m_1}(\mathbf{x}_1)\psi_{n_2\kappa_2m_2}(\mathbf{x}_2)\psi_{n_3\kappa_3m_3}(\mathbf{x}_3). \quad (4)$$

To account for the weak interaction, the intermediate d_1 and d_2 states should be considered with a small admixture of the closest-lying opposite-parity d_2 and d_1 states, respectively. Then the wave functions of the corresponding doubly excited states modify as

$$|\Psi_{d_1}\rangle \to |\Psi_{d_1}\rangle + i\xi|\Psi_{d_2}\rangle, \quad (5)$$

$$|\Psi_{d_2}\rangle \to |\Psi_{d_2}\rangle + i\xi|\Psi_{d_1}\rangle, \quad (6)$$

where the admixing parameter $i\xi = \langle\Psi_{d_2}|\sum_{i=1}^3 H_W(i)|\Psi_{d_1}\rangle/(E_{d_1} - E_{d_2})$ is determined by the nuclear spin-independent effective Hamiltonian of the weak interaction

$$H_W = -(G_F/\sqrt{8})Q_W\rho_N(r)\gamma_5. \quad (7)$$

Here $Q_W \approx -N + Z(1 - 4\sin^2\theta_W)$ denotes the weak charge of the nucleus, G_F is the Fermi constant, γ_5 is the Dirac matrix, and ρ_N is the nuclear weak-charge density (normalized to unity). After substitution of the modified wave functions into Eq. (1) and summing over all decay channels, one finds

$$\sigma_{\mu_i} = \frac{(2\pi)^3}{v_i}\sum_{M_{d_1}M_{d_2}}\delta_{M_{d_1}M_{d_2}}\Bigg[\frac{\Gamma_{d_1}}{|E_i - E_{d_1} + i\Gamma_{d_1}/2|^2}|\langle\Psi_{d_1}|I|\Psi_i\rangle|^2 + \frac{\Gamma_{d_2}}{|E_i - E_{d_2} + i\Gamma_{d_2}/2|^2}|\langle\Psi_{d_2}|I|\Psi_i\rangle|^2 + 2\left(\frac{\Gamma_{d_1}}{|E_i - E_{d_1} + i\Gamma_{d_1}/2|^2} - \frac{\Gamma_{d_2}}{|E_i - E_{d_2} + i\Gamma_{d_2}/2|^2}\right)\mathrm{Re}\big[i\xi\langle\Psi_{d_1}|I|\Psi_i\rangle(\langle\Psi_{d_2}|I|\Psi_i\rangle)^*\big] + 2(\Gamma_{d_2} - \Gamma_{d_1})\mathrm{Re}\left(i\xi\frac{\langle\Psi_{d_1}|I|\Psi_i\rangle(\langle\Psi_{d_2}|I|\Psi_i\rangle)^*}{(E_i - E_{d_1} + i\Gamma_{d_1}/2)(E_i - E_{d_2} - i\Gamma_{d_2}/2)}\right)\Bigg]. \quad (8)$$

In this expression the terms of order ξ^2 are neglected. The first and second terms are parity conserving, while the third and fourth terms correspond to the parity-violating contributions to the cross section. The third term originates from the weak interaction in the dielectronic recombination process. The parity violation in the decay process is described by the fourth term. In the case of clearly resolved levels ($\Gamma_{d_1},\Gamma_{d_2} \ll |E_{d_1} - E_{d_2}|$), one can consider only the resonant term in Eq. (1). For example, if the energy of the incident electron is tuned to the d_1 state, the total cross section takes the form

$$\sigma_{\mu_i} = \frac{(2\pi)^3}{v_i} \frac{\Gamma_{d_1}}{\left|E_i - E_{d_1} + i\Gamma_{d_1}/2\right|^2} \times \sum_{M_{d_1} M_{d_2}} \delta_{M_{d_1} M_{d_2}} \{ |\langle \Psi_{d_1} | I | \Psi_i \rangle|^2 + 2\,\mathrm{Re}[i\xi \langle \Psi_{d_1} | I | \Psi_i \rangle (\langle \Psi_{d_2} | I | \Psi_i \rangle)^*] \}. \tag{9}$$

In the case when the energy spread of the electron beam exceeds the energy spacing between the quasidegenerate states, one should integrate Eq. (8) over the incident electron energies. It can be performed analytically since the velocity v_i and the DR amplitudes weakly change within the interval of the beam energy distribution. Thus, for the close-lying states, one obtains

$$\overline{\sigma}_{\mu_i} = \frac{(2\pi)^4}{v_i} \sum_{M_{d_1} M_{d_2}} \delta_{M_{d_1} M_{d_2}} \Bigg[|\langle \Psi_{d_1} | I | \Psi_i \rangle|^2 + |\langle \Psi_{d_2} | I | \Psi_i \rangle|^2 + 2\xi (\Gamma_{d_2} - \Gamma_{d_1}) \mathrm{Re}\left(\frac{\langle \Psi_{d_1} | I | \Psi_i \rangle (\langle \Psi_{d_2} | I | \Psi_i \rangle)^*}{(E_{d_1} - E_{d_2}) - i(\Gamma_{d_1} + \Gamma_{d_2})/2} \right) \Bigg], \tag{10}$$

where $\overline{\sigma}_{\mu_i}$ is the integrated cross section. When the energy distribution in the beam exceeds the energy widths but is much less than the energy distance between the quasidegenerate levels, one should integrate Eq. (9). For instance, for a nonmonoenergetic beam tuned to the d_1 state we obtain

$$\overline{\sigma}_{\mu_i} = \frac{(2\pi)^4}{v_i} \sum_{M_{d_1} M_{d_2}} \delta_{M_{d_1} M_{d_2}} \{ |\langle \Psi_{d_1} | I | \Psi_i \rangle|^2 + 2\,\mathrm{Re}[i\xi \langle \Psi_{d_1} | I | \Psi_i \rangle (\langle \Psi_{d_2} | I | \Psi_i \rangle)^*] \}. \tag{11}$$

III. RESULTS AND DISCUSSION

As mentioned above, the enhancement of the PNC effect takes place for close-lying opposite-parity levels. In our previous work [25] we found that for Li-like ions the near degeneracy takes place for several doubly excited opposite-parity $[(1s2s)_0 n\kappa]_{1/2}$ and $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ states with $4 \leqslant n \leqslant 7$, $\kappa = \pm 1$, $Z \sim 60$, and $Z \sim 92$. The energy difference has been evaluated as follows:

$$\Delta E = E_{[(1s2p_{1/2})_0 n\kappa]_{1/2}} - E_{[(1s2s)_0 n\kappa]_{1/2}} = \Delta E^{(\mathrm{He})} + \Delta E^{(\mathrm{ext})}, \tag{12}$$

where $\Delta E^{(\mathrm{He})} = E_{(1s2p_{1/2})_0} - E_{(1s2s)_0}$ is the energy difference of the corresponding levels in the He-like ion and $\Delta E^{(\mathrm{ext})} = E^{(n\kappa)}_{(1s2p_{1/2})_0} - E^{(n\kappa)}_{(1s2s)_0}$ is the difference of the one-photon exchange contributions, describing the interaction between the external $n\kappa$ electron and the inner-shell electrons. The highly accurate values of $\Delta E^{(\mathrm{He})}$, including all second-order two-electron QED contributions, were taken from Ref. [26]. We also have taken into account the mixing of the close-lying $(1s2sns)_{1/2}$ and $(1s2p_{1/2}np_{1/2})_{1/2}$ levels, as well as the $(1s2p_{1/2}ns)_{1/2}$ and $(1s2snp_{1/2})_{1/2}$ levels (see Ref. [25] for details).

In the present work we consider the PNC effect in the process of the dielectronic recombination into $d_1 = [(1s2p_{1/2})_0 n\kappa]_{1/2}$ and $d_2 = [(1s2s)_0 n\kappa]_{1/2}$ states of Li-like ions. First, let us denote the cross sections for positive and negative helicities (spin projection onto the electron momentum direction) of the incident electron by σ_+ and σ_-, respectively. We also introduce designations for the cross section without the PNC effect $\sigma_0 = (\sigma_+ + \sigma_-)/2$ and the PNC contribution $\sigma_{\mathrm{PNC}} = (\sigma_+ - \sigma_-)/2$. Deviation of σ_{PNC} from zero indicates the parity-violation effect. Finally, one should determine the requirements on the luminosity L, provided the PNC effect is measured to a relative accuracy η [14,15],

$$L > L_0 = \frac{\sigma_+ + \sigma_- + 2\sigma_b}{(\sigma_+ - \sigma_-)^2 \eta^2 T}. \tag{13}$$

Here σ_b is the background magnitude and T is the acquisition time. In our calculations we neglect the background signal, set $T = 2$ weeks, and $\eta = 0.01$. In the case of the nonmonoenergetic incident electron beam the integrated cross sections $\overline{\sigma}_0 = (\overline{\sigma}_+ + \overline{\sigma}_-)/2$ and $\overline{\sigma}_{\mathrm{PNC}} = (\overline{\sigma}_+ - \overline{\sigma}_-)/2$ should be used instead of σ_0 and σ_{PNC}. Here we denote the integrated cross sections for positive and negative helicities of the incident electron by $\overline{\sigma}_+$ and $\overline{\sigma}_-$, respectively.

In order to investigate whether or not the levels mixed by the weak interaction are distinguished we introduce the coefficient $R = |E_{d_1} - E_{d_2}|/(\Gamma_{d_1} + \Gamma_{d_2})$. Evaluating the cross section according to Eqs. (8) and (9), it was found that the results became similar at $R \geqslant 5$. Thus, levels with $R \geqslant 5$ are regarded as distinguishable.

In Tables I and II we present numerical results for the most promising case of the resonant DR into the $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ state at n, κ, and Z, which provide the minimum values of

TABLE I. Cross section of the dielectronic recombination of a polarized electron with a He-like ion in the case of resolved levels ($R \geqslant 5$). The electron energy is tuned in resonance with the $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ state. The parameters n, κ, and Z correspond to the minimal values of luminosity L_0, $\Delta E = E_{[(1s2p_{1/2})_0 n\kappa]_{1/2}} - E_{[(1s2s)_0 n\kappa]_{1/2}}$ is the energy difference, R is the coefficient indicating whether or not the states are resolved, σ_0 is the cross section without the PNC effect, σ_{PNC} is the parity-violating contribution, and $\Delta\sigma_0$ indicates the increase of the process cross section related to the usage of Eq. (8) instead of Eq. (9).

Z	$n\kappa$	ΔE (eV)	R	ε_i (keV)	L_0 (cm^{-2} s^{-1})	σ_0 (b)	$\Delta\sigma_0$ (%)	σ_{PNC} (b)
88	7s	3.17(29)	19.3	84.76	1.1×10^{30}	3.8×10^2	0.1	1.2×10^{-3}
90	5s	4.13(47)	7.7	86.91	1.4×10^{30}	2.8×10^2	0.9	9.1×10^{-4}
	6s	2.51(47)	7.9	88.36	5.3×10^{29}	2.7×10^2	0.9	1.5×10^{-3}
	7s	1.75(47)	8.5	89.22	2.6×10^{29}	2.6×10^2	0.8	2.0×10^{-3}
92	5s	2.97(28)	5.0	91.43	5.1×10^{29}	2.5×10^2	2.1	1.4×10^{-3}
	7s	−1.60(28)	7.2	93.86	1.5×10^{29}	2.2×10^2	1.1	-2.4×10^{-3}

TABLE II. Cross section of the dielectronic recombination of a polarized electron with He-like ion in the case of unresolved levels ($R < 5$). The electron energy is tuned in resonance with the $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ state. The parameters n, κ, and Z correspond to the minimal value of luminosity L_0, $\Delta E = E_{[(1s2p_{1/2})_0 n\kappa]_{1/2}} - E_{[(1s2s)_0 n\kappa]_{1/2}}$ is the energy difference, R is the coefficient indicating whether or not the states are resolved, σ and $\overline{\sigma}$ are the cross sections corresponding to the monoenergetic and nonmonoenergetic energy distribution of the incident electron beam, respectively, σ_0 is the cross section without the PNC effect, and $\sigma_{\rm PNC}$ is the parity violating contribution.

Z	$n\kappa$	ΔE (eV)	R	ε_i (keV)	L_0 (cm^{-2} s^{-1})	σ_0 (b)	$\sigma_{\rm PNC}$ (b)	$\overline{\sigma}_0$ (b eV)	$\overline{\sigma}_{\rm PNC}$ (b eV)
62	$7s$	$-0.103(64)$	2.0	39.56	3.6×10^{29}	1.4×10^{3}	-4.0×10^{-3}	4.8×10^{2}	5.5×10^{-5}
88	$7p_{1/2}$	$-2.46(29)$	4.4	84.76	1.7×10^{30}	2.8×10^{1}	-2.6×10^{-4}	9.6×10^{2}	-7.0×10^{-6}
90	$6p_{1/2}$	$-1.26(47)$	1.1	88.37	1.0×10^{30}	9.6×10^{1}	-6.2×10^{-4}	1.6×10^{3}	-2.6×10^{-5}
92	$6s$	$-1.07(28)$	3.0	92.96	7.3×10^{28}	2.5×10^{2}	-3.8×10^{-3}	6.8×10^{2}	1.7×10^{-5}
	$6p_{1/2}$	2.38(27)	2.0	92.96	1.3×10^{30}	4.3×10^{1}	3.6×10^{-4}	1.6×10^{3}	-1.8×10^{-5}
	$7p_{1/2}$	2.38(28)	3.2	93.86	8.1×10^{29}	2.8×10^{1}	3.8×10^{-4}	1.0×10^{3}	-4.0×10^{-6}

the luminosity L_0. Table I corresponds to the case of resolved opposite-parity states, whereas the case of unresolved states is presented in Table II.

It is clearly seen from Table I that the R coefficient can be applied in order to distinguish cases of resolved and unresolved states. Indeed, at the border value ($R = 5$), σ_0 increases only by about 2% for the calculations utilizing Eq. (8) instead of Eq. (9). For the other parameters n, κ, and Z listed in Table I the growth of the cross section amounts to 1% or less.

According to Tables I and II the PNC effect seems to be most promising for the dielectronic recombination of a polarized electron with He-like uranium ($Z = 92$). When the energy of the incident electron is tuned in resonance with the $[(1s2p_{1/2})_0 6s]_{1/2}$ state, the ratio $\sigma_{\rm PNC}/\sigma_0$ equals -1.5×10^{-5}. After integration over ε_i it turns into $\overline{\sigma}_{\rm PNC}/\overline{\sigma}_0 = 2.5 \times 10^{-8}$. Let us compare the obtained results with similar calculations presented in Ref. [14]. In that work, the authors considered the process of the dielectronic recombination into the $(2s)^2$ and $(2s2p_{1/2})_0$ states for $Z = 48$, where the enhancement of the parity-violating effect takes place due to the quasidegeneracy of these levels. The PNC asymmetry of the process considered in Ref. [14] amounted to 5×10^{-9}, while for the process considered in the present work it reaches 1.5×10^{-5}. The increase of the effect by more than three orders of magnitude is caused by the fact that the admixing parameter ξ for $Z = 48$, obtained in Ref. [14], equals 6.0×10^{-9}, whereas for Li-like uranium we get $\xi = 4.0 \times 10^{-6}$.

In Fig. 1, $\sigma_{\rm PNC}$ is displayed as a function of the energy of the incident electron in the case of unresolved levels. As one can see from this figure, the PNC cross section is mainly formed by the parity-violation effect in the dielectronic recombination process [third term in Eq. (8)]. Nevertheless, the contribution from the subsequent radiative decay [fourth term in Eq. (8)] slightly enhances $\sigma_{\rm PNC}$ for the energy of the incident electron tuned in resonance with d_2 state. Conversely, for the energy tuned in resonance with the d_1 state a small decrease of the PNC contribution is observed. One can observe the energy of the incident electron at which $\sigma_{\rm PNC}$ turns to zero. It approximately corresponds to the energy just in between the quasidegenerate d_1 and d_2 states.

The experiment suggested in our paper involves a stored heavy ion beam intersecting with a beam of polarized electrons in an electron target or cooler. The polarized electrons can be produced with a semiconductor photocathode with circularly polarized laser light [27]. They are electrostatically accelerated to the energies of tens of keV that are required for the experiment. The cooler with a photocathode was constructed, for instance, for the TSR storage ring at the MPI-K Heidelberg. Such coolers can in principle be made to produce polarized electrons and they are now under consideration for the FAIR facility and for the CRYRING at GSI, Darmstadt.

The high electron energy definition in the rest frame of the ion is required to achieve the DR resonance. The electron beam energy spread depends on the collision energy ε and the transverse $kT_\perp$ and the longitudinal $kT_{||}$ temperatures of the electron beam: $\Delta\varepsilon = \sqrt{(\ln 2kT_\perp)^2 + 16\ln 2\varepsilon kT_{||}}$ [28]. The laser-produced beams of electrons are intrinsically cold and can be further cooled using an adiabatic beam expansion technique. Beams with a transverse temperature of 3.6 meV and a longitudinal temperature of 38 μeV were produced in electron cooler devices [29]. Accordingly, the energy spread of a few eV at 90 keV can be experimentally achieved at present. This means that the DR resonance structure will be integrated out. To take full advantage of the enhancement of PNC in dielectronic recombination the electron energy spread must be

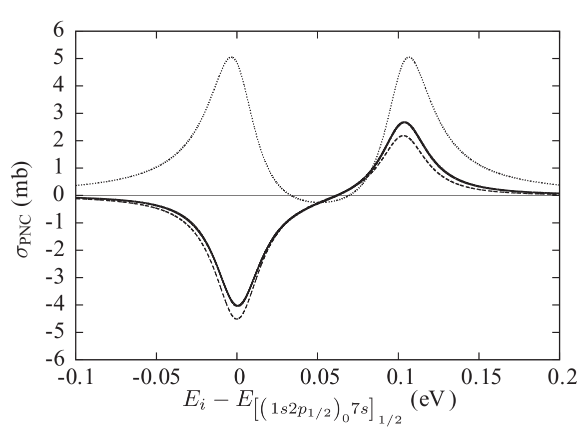

FIG. 1. The PNC cross sections of the dielectronic recombination into the $[(1s2p_{1/2})_0 7s]_{1/2}$ and $[(1s2s)_0 7s]_{1/2}$ states of Li-like samarium ($Z = 62$). The difference $E_i - E_{[(1s2p_{1/2})_0 7s]_{1/2}}$ determines uniquely the energy of the incident electron. The solid line corresponds to $\sigma_{\rm PNC}$, the dashed line is the parity-violating contribution from the dielectronic recombination, and the dotted line is the PNC contribution from the decay process multiplied by a factor of 10.

made smaller than 0.1 eV at 90 keV, which is nowadays not possible. Therefore, further developments will be required to produce an electron beam that is cold enough. In addition, the ion beam momentum spread should be reduced below 10^{-6}. This, however, was demonstrated at the storage ring ESR, albeit with a significant reduction of the beam intensity [30].

IV. CONCLUSION

In the present work we have considered the PNC effect on the cross section of the dielectronic recombination into the $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ and $[(1s2s)_0 n\kappa]_{1/2}$ states of heavy Li-like ions. The calculations have been performed for the parameters n, κ, and Z, which provide the enhancement of the parity-violation effect due to quasidegeneracy of the corresponding levels. It has been found that at energies of incident electron tuned in resonance with the $[(1s2p_{1/2})_0 n\kappa]_{1/2}$ state the PNC effect becomes most pronounced. The estimation of the PNC asymmetry for the most promising case of $[(1s2p_{1/2})_0 6s]_{1/2}$ and $[(1s2s)_0 6s]_{1/2}$ states for $Z = 92$ has given -1.5×10^{-5}, which is by several orders of magnitude bigger than the result obtained for a similar process in Ref. [14].

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