Photon-photon polarization correlations as a tool for studying parity nonconservation in heliumlike uranium

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Due to electron-nucleus weak interaction, atomic bound states with different parities turn out to be mixed. We discuss a prospective method for measuring the mixing parameter between the nearly degenerate metastable states $1s_{1/2}2s_{1/2}\ :\ J=0$ and $1s_{1/2}2p_{1/2}\ :\ J=0$ in heliumlike uranium. Our analysis is based on the polarization properties of the photons emitted in the two-photon decays of such states.

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I. INTRODUCTION

Parity nonconservation (PNC) had been at first theoretically proposed by Lee and Yang in 1956 in order to find a way out of the so-called τ-γ puzzle [1,2]. The next year, Wu and collaborators observed an asymmetry in nuclear β decay ascribed to parity nonconservation in weak processes [3]. Many later experiments in nuclear and high-energy physics confirmed parity violation in weak interactions and precisely recorded weak charge and other related parameters [4–7]. Although, with some initial controversies, the τ-γ puzzle was also solved by understanding that both τ and γ were two decay channels of the same parent particle, known today as the charged kaon K^+ [8,9]. In contrast to nuclear and high-energy physics, fewer experiments have been carried out in atomic physics to measure the properties of weak interaction. In fact, the conflicting results of the early Bismuth experiments in the 1970s [10–13] spread the belief that nothing fundamentally useful could ever have been extracted from atomic physics experiments. Nonetheless, renewed interest on the subject arose in the late 1980s and 1990s and led to the successful measurements of the weak charge Q_w and related parameters in atomic cesium [14–19], thallium [20], lead [21] and yttrium [22]. On the theoretical side, starting from the early work of Curtis-Michel [23], several investigations of PNC have been made in the context of neutral atoms [24], few-electron ions [25,26], and muonic atoms [27,28]. In all the proposed studies, the small role played by PNC effects together with the need of precise measurements have been highlighted.

Parity violation in atomic physics is mainly caused by the exchange of the Z^0 boson between atomic electrons and quarks in the nucleus. All atomic states become mixtures consisting mainly of the state they are usually assigned together with a small percentage of states possessing the opposite parity. The prospective method for measuring the mixing between the states $1s_{1/2}2s_{1/2} : J=0$ and $1s_{1/2}2p_{1/2} : J=0$ in heliumlike uranium by inducing a resonant parity-violating $E1E1$ transition between them has been discussed [29]. The authors concluded that the proposed measurement was not feasible with the technology available at that time, while it is currently under consideration at the GSI facility in Darmstadt (Germany). Some years later, Dunford proposed an analysis with the same goal, based on the circular polarization asymmetry of one of the photons emitted in the two-photon decay of the $1s_{1/2}2p_{1/2} : J=0$ state [30]. The author concluded that the calculations performed were not enough to assess whether or not the polarization asymmetry could lead to useful parity experiments. With the same intent and similar method, we propose another route based on photon polarization properties for the experimental determination of the mixing parameter between the states $1s_{1/2}2s_{1/2} : J=0$ and $1s_{1/2}2p_{1/2} : J=0$ in U^{90+} (in the following, these two states will be briefly called $2\,^1S_0$ and $2\,^3P_0$, respectively). The different polarization properties of the photons emitted in the two-photon decays of such states suggest a way to discriminate the decays and, thereby, to measure the mixing parameter between the states. However, the prospective method presents some technical difficulties, thoroughly discussed in the text, that make the experimental realization a challenge even with current state-of-the-art technologies.

The paper is structured as follows. Section II describes the salient characteristics of the first excited states of U^{90+}. Section III shows the geometry we refer to, while Sec. IV describes the two-photon transition amplitude and the employed atomic model. The polarization-polarization correlation, that is, the function which denotes the probability of detection of photons with certain polarizations, will also be discussed in detail. In Sec. V, results are shown for this correlation function, emphasizing the role played by parity-mixing terms. With results in hand, the experimental setup for the prospective method is then explained and thoroughly discussed. Finally, a brief summary is given in Sec. VI.

II. HELIUMLIKE URANIUM ION

Heliumlike uranium ion represents a very suitable candidate for studying PNC due to the fact that the states $2\,^3P_0$ and $2\,^1S_0$ are separated by an energy difference of only a few eV [31,32], out of a total binding energy of order 100 keV. Figure 1 shows the scheme of the first levels of U^{90+} [33]. The $2\,^3P_0$ state has negative parity and a lifetime of about $\sim 10^{-11}$ sec, while the $2\,^1S_0$ state has positive parity and a shorter lifetime of about $\sim 10^{-13}$ sec. Although $2\,^3P_0$ can decay by single photon emission into $2\,^3S_1$, both $2\,^3P_0$ and

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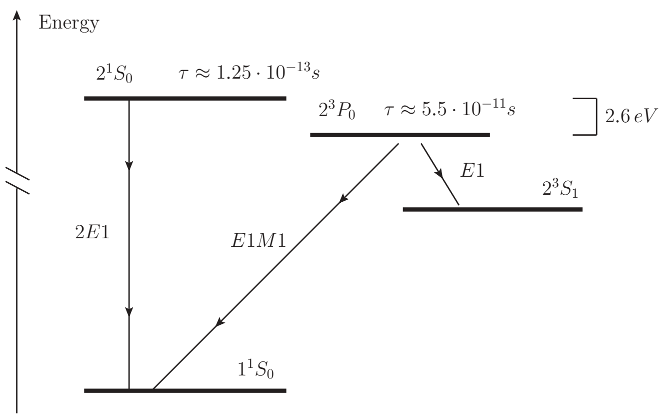

FIG. 1. Level scheme of few low energetic states in U^{90+}.

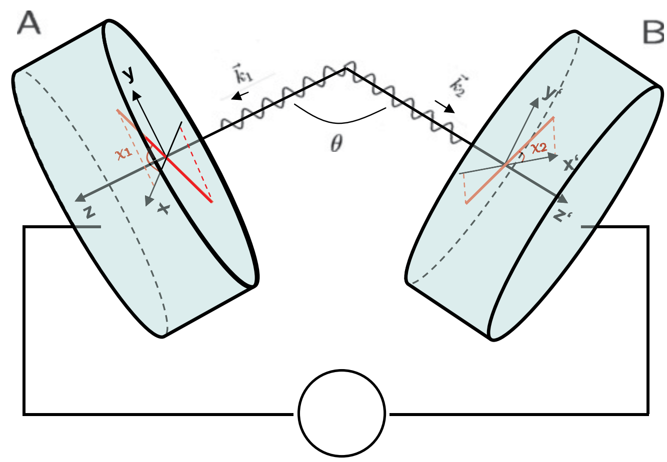

FIG. 2. (Color online) Geometry for the two-photon decay. The z (z') axis is adopted along the propagation direction of the first (second) photon, the x axis is chosen such that x-z is the reaction plane, while the y axis is coincident with y'. The definitions of the opening angle and of the angles which define the detectors' transmission axes are also displayed.

$2\,^1S_0$ decay exclusively by two-photon decay to the ground state, owing to angular momentum conservation. Due to weak interaction between electrons and nucleus, $2\,^3P_0$ acquires a small admixture of $2\,^1S_0$ and vice versa. Since the size of the parity mixing depends inversely on the energy difference between the mixed states [29], both $2\,^3P_0$ and $2\,^1S_0$ do not get any other considerable PNC contribution from any other state. More explicitly, at the first order in perturbation theory, the true $|2\,^3P_0\rangle$ and $|2\,^1S_0\rangle$ states can be written as [30],

$$\begin{aligned} |2\,^3\tilde{P}_0\rangle &\approx |2\,^3P_0\rangle + \eta\,|2\,^1S_0\rangle\,, \\ |2\,^1\tilde{S}_0\rangle &\approx |2\,^1S_0\rangle + \eta\,|2\,^3P_0\rangle\,, \end{aligned} \tag{1}$$

where the tilde notation is here and henceforth used to denote true states, in order to differentiate them from the bare theoretical Dirac states which will be denoted without the tilde. The mixing parameter η in Eq. (1) is given by

$$\eta = \frac{\langle 2\,^3P_0|\,\hat{H}_W\,|2\,^1S_0\rangle}{\Delta E_{\mathrm{PS}}}, \tag{2}$$

where ΔE_{PS} is the energy difference between $2\,^3P_0$ and $2\,^1S_0$, while $\hat{H}_W$ is the operator for the nuclear-spin-independent weak interaction [30]. Up to a very good approximation, we will neglect any parity-mixing effect in any state with the exception of $2\,^3P_0$ and $2\,^1S_0$. Among the low-energetic states in U^{90+}, only these two have, in fact, energies close enough to determine a sizable mixing parameter between them.

III. GEOMETRY

The geometry we want to adopt for the prospective method is displayed in Fig. 2. We define one local system of reference for each emitted photon. The axes definitions are as follows: The propagation direction of the first (second) photon is adopted as z (z') while the angle between the photons' directions (opening angle) is called θ. The x axis is fixed such that the plane defined by the two photons' directions (reaction plane) is the x-z plane. Using standard notation, the photon polarization plane is the plane which is orthogonal to the photon's direction with the origin of coordinate axes located at the position where the photon is detected. As displayed in Fig. 2, the A (B) detector measures the linear polarization of the first (second) photon along the transmission axis defined by the angle $\chi_{1(2)}$ in the polarization plane. The detectors are thought to work as polarizer filters—whenever a photon hits one of them, the detector gives off a click, or no click, which would indicate that the photon has been measured as having its polarization along the direction $\chi_{1,2}$ or $\chi_{1,2} + 90^\circ$, respectively.

Finally, we define the sharing parameter f as the fraction of energy carried away by the first photon:

$$f = \frac{\omega_1}{E_i - E_f} = 1 - \frac{\omega_2}{E_i - E_f}\,, \tag{3}$$

where $E_{i,f}$ are the energies of the initial and final ionic states, while $\omega_{1,2}$ are the recorded energies of the first and second photon, respectively. Energy conservation has been used in the last step of the above equation.

IV. THEORY

For the purpose of measuring the parameter η in Eq. (2), we propose to prepare U^{90+} in the $2\,^3\tilde{P}_0$ state. The efficiency of such preparation is here assumed to be 100%. The prepared $2\,^3\tilde{P}_0$ state will decay either into $2\,^3S_1$ or into the ground state, as extensively displayed in Fig. 1. The two-photon decay channel $2\,^3\tilde{P}_0 \to 1\,^1S_0$, in which we are interested, can be easily selected out in experiments by requiring a two-detector coincidence measurement. The amplitude for this process can be obtained in second-order perturbation theory and reads [34],

$$\begin{aligned} &\mathcal{M}^{\lambda_1\lambda_2}(2\,^3\tilde{P}_0 \to 1\,^1S_0) \\ &= \sum\!\!\!\!\!\!\int_\nu \Bigg(\frac{\langle 1\,^1S_0|\,\vec{\alpha}\cdot\vec{u}^*_{\lambda_1} e^{-i\vec{k}_1\cdot\vec{r}}\,|\nu\rangle\,\langle\nu|\,\vec{\alpha}\cdot\vec{u}^*_{\lambda_2} e^{-i\vec{k}_2\cdot\vec{r}\,'}\,|2\,^3\tilde{P}_0\rangle}{E_\nu - E_i + \omega_2} \\ &\quad + \frac{\langle 1\,^1S_0|\,\vec{\alpha}\cdot\vec{u}^*_{\lambda_2} e^{-i\vec{k}_2\cdot\vec{r}}\,|\nu\rangle\,\langle\nu|\,\vec{\alpha}\cdot\vec{u}^*_{\lambda_1} e^{-i\vec{k}_1\cdot\vec{r}\,'}\,|2\,^3\tilde{P}_0\rangle}{E_\nu - E_i + \omega_1} \Bigg). \end{aligned} \tag{4}$$

Here, $\vec{\alpha}$ is the vector of Dirac matrices, while the symbol $\sum\!\!\!\!\!\!\int_\nu$ stands for both a summation over the discrete and an integration over the continuum part of the ionic spectrum. In addition, E_ν is the energy of the intermediate electronic state $|\nu\rangle$, while $\vec{k}_{1,2}$ and $\vec{u}_{\lambda_{1,2}}$ denote the linear momentum and the polarization vector of the first (second) photon, respectively. The latter directly depends on the photons helicities $\lambda_{1,2} = -1,1$. By introducing Eq. (1) into Eq. (4), the amplitude is split into two terms,

$$\mathcal{M}^{\lambda_1\lambda_2}(2\,^3\tilde{P}_0 \to 1\,^1S_0) \approx \mathcal{M}^{\lambda_1\lambda_2}(2\,^3P_0 \to 1\,^1S_0) + \eta\mathcal{M}^{\lambda_1\lambda_2}(2\,^1S_0 \to 1\,^1S_0). \quad (5)$$

In order to suggest any experiment whose goal is the measurement of the mixing parameter η, we should be first able to theoretically discriminate the two amplitudes of the right-hand side of Eq. (5). The key point of the prospective method is that such discrimination can be obtained by studying the photons' polarization properties contained in those amplitudes. It has recently been shown that, in the case that nearly equal energy is shared between the photons, the two-photon decay $2\,^3P_0 \to 1\,^1S_0$ is characterized by photon linear polarizations which are exclusively orthogonal to each other [linear polarizations of the first (second) photon are detected, correspondingly, along the x,y' or y,x' axes], while the two-photon decay $2\,^1S_0 \to 1\,^1S_0$ is characterized by photon linear polarizations which are exclusively parallel to each other (linear polarizations of the first (second) photon are detected, correspondingly, along the x,x' or y,y' axes) [35,36]. While the first assertion is true independently of the opening angle θ, the second one holds only in case the photons are recorded either collinearly or back to back ($\theta = 0^\circ, 180^\circ$). However, as it will be evident in the following, the linear polarizations of photons emitted in $2\,^1S_0 \to 1\,^1S_0$ decay can be considered parallel in the whole intervals $0^\circ \leqslant \theta \lesssim 2^\circ$ and $178^\circ \lesssim \theta \leqslant 180^\circ$, due to the fact that the orthogonal corrections to the polarization state are negligible in that region, even for the delicate problem under consideration. As a matter of fact, for the case $0^\circ \leqslant \theta \lesssim 2^\circ$ (or $178^\circ \lesssim \theta \leqslant 180^\circ$) and $f = 0.5$, it can be demonstrated that the polarization state of the two photons emitted in consequence of the decay of the prepared $2\,^3\tilde{P}_0$ state can be simply described by the ket vector [36]

$$|\Psi\rangle = O^{\rm PS}_{f,Z,\theta}\,(|xy\rangle + |yx\rangle) + \eta\, O^{\rm SS}_{f,Z,\theta}\,(|xx\rangle + |yy\rangle)\,, \quad (6)$$

where $|O^{\rm PS}_{f,Z,\theta}|^2$ is the probability of detecting the emitted photons with polarizations along $\chi_1 = 0^\circ$, $\chi_2 = 90^\circ$, or $\chi_1 = 90^\circ$, $\chi_2 = 0^\circ$ while $|O^{\rm SS}_{f,Z,\theta}|^2$ is the probability of detecting the photons with polarizations along $\chi_1 = 0^\circ$, $\chi_2 = 0^\circ$, or $\chi_1 = 90^\circ$, $\chi_2 = 90^\circ$. Both $O^{\rm PS}_{f,Z,\theta}$ and $O^{\rm SS}_{f,Z,\theta}$ contain the dependence on the energy-sharing parameter f, the atomic number Z, and the opening angle θ given, respectively, by the amplitudes $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^3P_0 \to 1\,^1S_0)$ and $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^1S_0 \to 1\,^1S_0)$.

In order to inspect the polarization properties of the photons emitted in the two-photon decay of the $2\,^3\tilde{P}_0$ state, we define the polarization-polarization correlation function, which is the physical quantity we mean to investigate. This function is given by [35],

$$\Phi^f_{\chi_1,\chi_2}(\theta) = \frac{\mathcal{N}^2}{4(2J_i+1)} \sum_{\substack{\lambda_1\lambda_1' \\ \lambda_2\lambda_2'}} e^{i(\lambda_1-\lambda_1')\chi_1}\, e^{i(\lambda_2-\lambda_2')\chi_2}\, \mathcal{M}^{\lambda_1\lambda_2} \times (i \to f)\mathcal{M}^{\lambda_1'\lambda_2'*}(i \to f)\,, \quad (7)$$

where J_i is the total angular momentum of the initial ionic state. Thus, Φ represents the normalized probability density of measuring, in coincidence, two photons with well-defined wave vectors $\vec{k}_1$, $\vec{k}_2$ and with certain linear polarizations which are characterized by the angles χ_1, χ_2 with respect to the reaction plane (see Fig. 2 for details concerning the notation). The normalization constant $\mathcal{N}$ is chosen such that $1/\mathcal{N}^2$ is the sum of the probability densities of the four independent polarization outcomes $\chi_{1,2} = 0^\circ, 90^\circ$.

In order to complete the theoretical background needed for the prospective method, we conclude this section by explaining the model we use for the calculations. The description of two-electron ions is indeed a theoretical challenge of the current state of research in atomic physics. The method of relativistic finite-basis sets, for instance, has been shown to be valid and efficient in order to obtain highly accurate calculations of the two-photon $E1M1$ decay rate from the $2\,^3P_0$ state [37]. Alternatively, the salient characteristics of heavy heliumlike ions can be described by the independent particle model (IPM), which is the model used here for the calculations. Although this model treats the electrons as independent particles bound to the nucleus (the nuclear Coulomb attraction is assumed to be much stronger than the electron-electron repulsion), it takes the Pauli principle into account. Moreover, this model allows a drastic simplification of the two-electron amplitude which appears in Eq. (5), allowing it to be reduced to a summation over one-electron amplitudes [38]. Although the calculation of the latter quantity is itself a challenging theoretical problem, several methods have been successfully proposed in the past decades to precisely perform it [39,40]. The tool we adopt here for its calculation is the relativistic Dirac-Coulomb-Green function. For details regarding this approach, useful information can be found in Refs. [41,42].

The results shown in Sec. V are obtained by taking into account the full multipole contribution of the photons' fields. Finally, the effective nuclear charge used for the computation is $Z = 91.275$. This accounts for the electromagnetic screening that one electron makes on the other one, allowing for a basic electron-electron interaction.

V. PROSPECTIVE METHOD: RESULTS AND DISCUSSION

After explaining the theory at the basis of our prospective method as well as the model we used, we are now ready to concretely present the proposal. In order to measure the mixing parameter η in Eq. (2), we propose, as previously mentioned, to prepare the uranium ion U^{90+} in the $2\,^3\tilde{P}_0$ state, to place two polarization detectors at a fixed position in the reaction plane, and to use them as polarizer filters. While one of the two detectors will be kept at a fixed orientation (fixed transmission axis), the transmission axis of the other will be continuously rotated to record the correlation function in Eq. (7) for different photons' polarization configurations. In order to suggest a workable experimental scenario, we must inevitably look for opening angle and energy values which enable $\eta\, O^{\rm SS}_{f,Z,\theta}$ to be comparable with $O^{\rm PS}_{f,Z,\theta}$, in Eq. (6). In other words, since η is considerably small, we must find a configuration where the amplitude $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^3P_0 \to 1\,^1S_0)$ is small in comparison with $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^1S_0 \to 1\,^1S_0)$. For this purpose, it has been shown that the decay rate for the $2\,^3P_0 \to 1\,^1S_0$ transition is strongly suppressed for photons' opening angle $0 \leqslant \theta \lesssim 2^\circ$ and equal energy sharing, whereas, for the same configuration, the decay rate $2\,^1S_0 \to 1\,^1S_0$ gets almost its maximum value [38,43]. Choosing small values of θ and equal energy sharing will also ensure that the different amplitudes in Eq. (5) will determine different photons' polarization outcomes (as remarked in Sec. IV), which is decisive for the scope of the prospective method. An optimal configuration for our purposes can be found, for instance, at $\theta = 1^\circ$ and $f = 0.5$. For such a configuration, the coefficients $O^{\rm PS}_{f,Z,\theta}$ and $O^{\rm SS}_{f,Z,\theta}$, which compose the ket vector in Eq. (6), assume the values -8.49×10^{-11} and 4.43×10^{-5}, respectively. The correlation function Φ related to this polarization state can be easily calculated:

$$\begin{aligned}\Phi^{f=0.5}_{\chi_1,\chi_2}(\theta = 1^\circ) &= \mathcal{N}^2\,[-8.49 \times 10^{-11}\,(\cos\chi_1 \sin\chi_2 + \sin\chi_1 \cos\chi_2) \\ &\quad + \eta\, 4.43 \times 10^{-5}\,(\cos\chi_1 \cos\chi_2 + \sin\chi_1 \sin\chi_2)]^2. \qquad (8)\end{aligned}$$

We draw the above function in Fig. 3, where χ_1 has been arbitrarily set to 90° for a better visualization, while η has been fixed to the predicted theoretical value 1.75×10^{-6}, which can be obtained by correcting the value obtained in Ref. [29] with

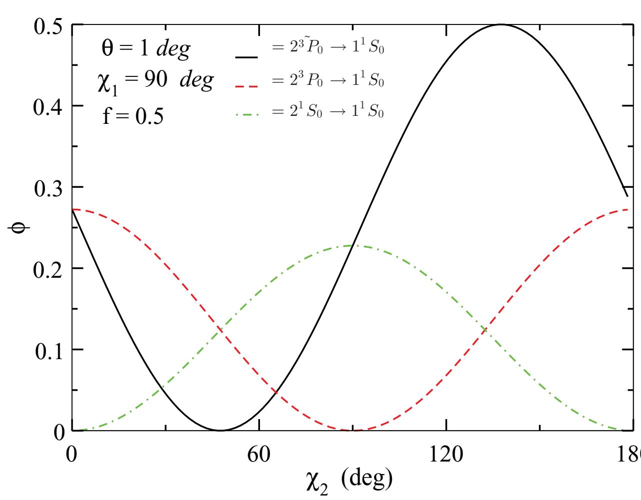

FIG. 3. (Color online) Polarization-polarization correlation function in Eq. (7) for the $2\,^3\tilde{P}_0 \to 1\,^1S_0$ two-photon decay of heliumlike uranium ion. The contribution of the different amplitudes in Eq. (5) are separately displayed. The dashed (red) line and the dash-dotted (green) line represent, respectively, the $P \to S$ and $S \to S$ contribution to the correlation function, while the solid (black) line denotes the total $\tilde{P} \to S$ correlation function. Here, θ is the opening angle, f is the photons' sharing energy, and $\chi_{1,2}$ are the linear polarization angles at which the first (second) detectors' transmission axes are set (see Fig. 2 for details).

the precise calculation of the $2\,^1S_0$-$2\,^3P_0$ energy gap shown in Ref. [32]. The different contributions of the two addends in Eq. (5) are separately displayed, as well as the total correlation function. We can easily notice that the parity-allowed $(|xy\rangle + |yx\rangle)$ and parity-forbidden $(|xx\rangle + |yy\rangle)$ components of the photon polarization state have approximately the same magnitude. In concordance with Ref. [35], it can be seen in the figure, as well as from Eq. (8), that the amplitudes $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^1S_0 \to 1\,^1S_0)$ and $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^3P_0 \to 1\,^1S_0)$ determine the probability of detecting parallel and orthogonal linearly polarized photons, respectively. In an ideal experiment, we could then scan the function Φ over the whole or part of the domain $\chi_{1,2} \in [0,180^\circ]$, in order to be able to determine the parameter η by fitting the measured polarization correlation with the η-dependent function (8).

The proposal is based on the fact that, for $f \to 0.5$ and $\theta \to 0$, the transition $2\,^3P_0 \to 1\,^1S_0$ model-independently vanishes. If we consider the two-photon transition $2\,^1\tilde{S}_0 \to 1\,^1S_0$, it can be easily seen from Eqs. (1) and (4) that the amplitude for that process would turn out to be equal to Eq. (5), with the replacement $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^1S_0 \to 1\,^1S_0) \leftrightarrow \mathcal{M}^{\lambda_1,\lambda_2}(2\,^3P_0 \to 1\,^1S_0)$. Since, unfortunately, there is no geometry for which the transition $2\,^1S_0 \to 1\,^1S_0$ is suppressed, the polarization of the emitted photons would be completely dominated by the parity-allowed component that, in that case, would be $(|xx\rangle + |yy\rangle)$. An initial preparation of the $2\,^1\tilde{S}_0$ state, therefore, although easier from an experimental point of view [44,45], would not give rise to the interference pattern shown in Fig. 3, for any given geometry.

Moreover, the amplitude $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^1S_0 \to 1\,^1S_0)$ is approximately one order of magnitude larger than $\mathcal{M}^{\lambda_1,\lambda_2}(2\,^3P_0 \to 1\,^1S_0)$, as can be seen from the lifetimes of the states displayed in Fig. 1. This fact represents an advantage for studying PNC effects in $2\,^3\tilde{P}_0 \to 1\,^1S_0$ rather than in $2\,^1\tilde{S}_0 \to 1\,^1S_0$, since this difference partially compensates for the small value of the mixing parameter η in Eq. (5) and so helps the two addends in the same equation to be comparable.

Although the suggested settings $f = 0.5$ and $\theta = 1^\circ$ ensure, as needed, that the $2\,^3P_0 \to 1\,^1S_0$ channel is strongly suppressed, they determine, at the same time, a challenging arrangement for the experimental investigation of the prospective method. Specifically, because of the required small opening angle θ, the two x-ray photon detectors would have to be placed at a relatively long distance from the source of radiation and thus the detection efficiency would be substantially decreased. An additional hindrance lies in the fact that the polarizations of both photons have to be measured at equal energy sharing ($f = 0.5$). For the case of uranium, this fact would imply that each photon has about 50 keV rest-frame energy. The polarization-resolved experiments in this x-ray energy regime are currently normally performed by using Compton polarimeters [46–51]. By selecting events recorded in coincidence, which have the desired (Compton) scattering angle, such polarimeters can be used to measure the polarization state of the photon pair. The selection of events can, however, increase the statistical uncertainty considerably.

Further experimental difficulties for the realization of the prospective method might arise from the angle-energy resolution needed to record the interference pattern shown in Fig 3. The $P \to S$ channel rises fast, glossing over the other $S \to S$ channel in which we are interested, as soon as we depart from the exact theoretical proposed configuration $f = 0.5$, $\theta = 1^\circ$. In other words, slightly different angle-energy settings would bring about a completely different polarization-polarization correlation function with respect to Eq. (8). As a matter of fact, the opening angle and energy resolutions needed in order to select events for which the correlation function does not change its approximate shape, would be, according to our calculations, 0.5° and 5 eV, respectively. Even though the required angle resolution may be achieved, the energy resolution needed is approximately three orders of magnitude higher than the available resolution in current Compton polarimeters. A possible way to overcome the energy resolution limitation would be to use the so-called absorption edge technique [52]. In this technique the photons pass through an absorption foil. The *K*-shell absorption edge of the foil atoms serves as a photon-energy filter. The photons with energy below the *K*-shell photoionization energy will have a significantly higher transmission probability than the photons with the higher energies. Since, in the proposed experimental scheme, both of the entangled photons have the same energy, one foil can be used as the energy filter for both of the photons. By adjusting the ion-beam velocity, the photon energy can be Doppler tuned such that it is less than 5 eV below the *K*-edge. A Compton scattering polarimeter behind the absorption foil can then be used for the polarization analysis of the transmitted photons. Another possible experimental approach would involve high-energy resolution calorimeters and a Rayleigh scattering polarimetry technique [53]. Here, the energy of the Rayleigh-scattered photon and its scattering direction could be measured with high resolution by an array of x-ray calorimeters. Such arrays are currently being developed [54,55] and likely to reach the required energy resolution at the energy of 50 keV in the near future.

The small expected value of the mixing coefficient η is certainly at the basis of the technical difficulties explained above. A way to ease such difficulties might be represented, for instance, by selecting a suitable isotope of U^{90+}. By virtue of the fact that the energy gap between the $2\,^1S_0$ and $2\,^3P_0$ states varies slightly by changing the mass number of the ion [32], the mixing of the two states itself would depend on the considered isotope [see Eq. (2)]. In particular, by choosing an isotope of uranium whose mass number is smaller than 238, we would be able to increase the mixing parameter of the two states up to a factor of $\approx$1.6. However, besides the technical difficulties related to the radioactive properties that the chosen isotope might show, such an improvement would not be enough to bring considerable advantages to the prospective method.

VI. SUMMARY

In summary, a prospective method for measuring the parity mixing parameter between the states $1s_{1/2}2s_{1/2} : J = 0$ and $1s_{1/2}2p_{1/2} : J = 0$ in heliumlike uranium has been presented. The core of the prospective method lies in the discernment of the two-photon decay of such states by using the polarization properties of the emitted photons. Within relativistic second-order perturbation theory and the independent-particle model, we explored the polarization-polarization correlation function of the photon pair for a chosen angle-energy configuration in which the role played by parity-mixing terms is highlighted. Within the suggested settings, the presence of parity-mixing contributions changes quantitatively and qualitatively the shape of the correlation function in the overall domain $0 \leqslant \chi_{1,2} \leqslant 180^\circ$. Such changes could, in principle, be measured in a polarization-angle-resolved experiment. However, the prospective method presents some technical difficulties, discussed in the text, which currently hamper its experimental investigation. The theoretical analysis which has been carried out on the polarization properties of the emitted photons may be also used as a side study for any other experimental investigations of PNC effects in atoms or ions.

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