Physics Letters B 741 (2015) 61–65

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

State-selective high-energy excitation of nuclei by resonant positron annihilation

Nikolay A. Belov *, Zoltán Harman

Max Planck Institute for Nuclear Physics, Saupfercheckweg 1, 69117 Heidelberg, Germany

ARTICLE INFO

Article history:
Received 23 October 2014
Received in revised form 14 November 2014
Accepted 1 December 2014
Available online 12 December 2014
Editor: J.-P. Blaizot

Keywords:
Positron annihilation
Nuclear excitation
Giant resonances

ABSTRACT

In the annihilation of a positron with a bound atomic electron, the virtual γ photon created may excite the atomic nucleus. We put forward this effect as a spectroscopic tool for an energy-selective excitation of nuclear transitions. This scheme can efficiently populate nuclear levels of arbitrary multipolarities in the MeV regime, including giant resonances and monopole transitions. In certain cases, it may have higher cross sections than the conventionally used Coulomb excitation and it can even occur with high probability when the latter is energetically forbidden.

Positron collisions with atomic matter lead to a number of processes [1–5], among which annihilation with shell electrons is one of the most prominent effects. Typically, annihilation leads to the emission of γ rays. Alternatively, the atomic nucleus may resonantly absorb the whole energy of the annihilating particles and become excited, as first proposed in [6]. We refer to this single-step process as nuclear excitation by resonant positron annihilation (NERPA) [7]. It is represented in Fig. 1(a) by the level schemes of the electron shell and the nucleus. NERPA may be followed by radiative nuclear de-excitation. The leading diagram of this two-step process, labeled as NERPA-γ, is shown in Fig. 1(b).

Attempts to observe NERPA have not been conclusive so far. Only an upper bound of its cross section has been determined in the latest experiment [8]. In this measurement, a monochromatic positron beam was employed, and the results suggest that in previous experiments employing broadband sources [9], some unidentified non-resonant process may have played the dominant role, and, therefore, yielded cross sections well above the theoretical NERPA cross sections [6,10]. Further works [11,12], as well as newer results [13] have improved the theoretical description. However, the disagreement of different measurements is still not understood. In the present work we investigate a series of new cases where NERPA, having higher cross sections, can be observed. Therefore, our studies and future experimental work is anticipated to provide an unambiguous identification of the NERPA process

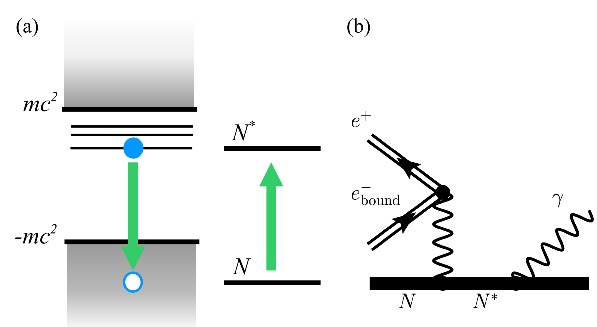

Fig. 1. (Color online.) (a) The illustration of positron annihilation via nuclear excitation by the corresponding fermionic and nuclear level schemes. mc^2 is the rest energy of the e^-, and N (N^*) stands for the nuclear initial (excited) state. (b) The lowest-order Feynman diagram of the NERPA process followed by nuclear de-excitation by emission of a photon. Thick lines denote nuclear states, double lines denote fermions in the Coulomb field of the nucleus, and wave lines represent real or virtual photons.

and a determination of its cross section. Furthermore, we introduce several applications of NERPA for future research in nuclear physics.

NERPA constitutes a way to excite nuclei which is alternative to photo- and Coulomb excitation. Photo-excitation experiments are conventionally done with bremsstrahlung, synchrotron or inverse Compton sources. X-ray free electron lasers, providing the highest photon intensities, are presently limited to the keV photon energy regime [14,15]. The great advantage of photo-excitation is the monochromaticity of the x- or γ-ray beam and the resonant

* Corresponding author.
E-mail address: belov@mpi-hd.mpg.de (N.A. Belov).

http://dx.doi.org/10.1016/j.physletb.2014.12.002

character of the nuclear excitation. The transitions predominantly accessible this way are, however, of electric-dipole ($E1$) type. On the other hand, Coulomb excitation, i.e. excitation by inelastic scattering of massive charged particles may induce transitions of arbitrary multipolarities. As an example, octupole deformations have been recently investigated by Coulomb excitation in collisions with lighter nuclei, and monitoring the subsequent γ decay [16]. Coulomb excitation is however not selective with respect to the nuclear energy levels.

In this Letter we show that NERPA has an attractive combination of the above advantages: the resonant character of the excitation and a significant cross section regardless of the multipolarity. We find that in certain cases its cross section can be comparable to or an order of magnitude higher than that of Coulomb excitation, leading to several applications. One of examples is the possibility to resonantly induce monopole transitions, that is particularly important for the studies of deformed nuclei. Another application is a collective nuclear excitation, the giant monopole resonance [17,18], also termed as "breathing" mode as it involves an oscillation of the nuclear volume, and provides the only way known for the experimental study of nuclear compressibility. In the case of giant nuclear resonances of any multipolarity, NERPA bears all the above-mentioned advantages of Coulomb excitation. A further property of NERPA is that experiments may be performed with neutral atoms, e.g. solid targets, without the necessity of stripping off the atomic electrons or accelerating the nuclei. However, the shift of the resonance energy due to change of the electron structure after possible annihilation of incoming positrons with higher-shell electrons has to be taken into account when planning such an experiment.

Very recently, intense positron jets with MeV energies or above have been generated in laser-plasma interactions [19–22], with a quasi-monoenergetic spectrum, and with positron numbers reaching 10^{10} per shot [19]. In this Letter we therefore put forward the indirect laser excitation of nuclei via NERPA, utilizing positrons produced by strong laser pulses. At present, such positron beams [19] have an energy width W much larger than the NERPA linewidth $\Gamma_{\rm nucl} + \Gamma_K$, being the sum of the nuclear and the atomic K-shell width. Thus, only an approximate fraction $\frac{\Gamma_{\rm nucl}+\Gamma_K}{W}$ of positrons are utilized for NERPA. However, broadband beams can be used for the excitation of broad nuclear lines, including giant resonances, or, furthermore, some certain energy regions within these resonances. This scheme is complementary to excitation by γ photons generated via Compton backscattering as planned, e.g., for the ELI facility [23,24], and to other direct or indirect nuclear excitation mechanisms utilizing optical [25–29] or x-ray light sources [30–34].

The NERPA process is the time reverse of e^-e^+ internal pair conversion accompanying the decay of an excited nucleus. A concise model of that process is given in [35,36]. We extend the formalism of these works for the case of NERPA. Its transition rate (probability per unit time) is given by

$$P_{\rm N}^{j',j} = \frac{\delta(E+E'-\omega)}{[J_i, j', j]} \sum_{M_i=-J_i}^{J_i} \sum_{M_f=-J_f^*}^{J_f^*} \sum_{\mu=-j}^{j} \sum_{\mu'=-j'}^{j'} (2\pi)^2 \times \left| \alpha \int d\mathbf{r}_n \int d\mathbf{r}_f \, j_n(\mathbf{r}_n) j_f(\mathbf{r}_f) \frac{e^{i\omega|\mathbf{r}_n-\mathbf{r}_f|}}{|\mathbf{r}_n-\mathbf{r}_f|} \right|^2, \quad (1)$$

where j and j' are the total angular momentum quantum numbers of the e^+ and e^- states, respectively, J_i and J_f^* are angular momenta of the nuclear initial and final states, and the M-s and μ-s are the associated magnetic quantum numbers. Furthermore, α is the fine-structure constant, $\mathbf{r}_n$ and $\mathbf{r}_f$ denote nuclear and fermionic coordinates, and ω stands for the virtual γ energy.

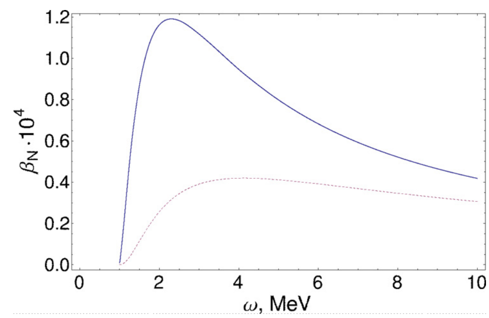

Fig. 2. (Color online.) The NERPA coefficient $\beta_{\rm N}^{j',j}$ for $E2$ nuclear transition between $1/2+$ and $5/2+$ nuclear states vs. the photon energy ω for $Z = 50$ (^{115}Sn). Blue solid line: the case of $j = L - 1/2$, brown dashed line: $j = L + 1/2$.

E and E' are the e^+ and e^- energies, respectively, and j_n and j_f are nuclear and fermionic 4-currents [37]. Also, the notation $[j_1, \ldots, j_n] = \prod_{i=1}^{n}(2j_i+1)$ has been introduced.

The further derivation involves a splitting of the double integral over nuclear and fermionic coordinates into a product of nuclear and fermionic parts. This can be performed in the non-penetration approximation, where the overlap integral of the bound-electron wave function within the nuclear volume is neglected. This simplification is generally used in the calculation of quantities related to the interaction of shell electrons with the nucleus [35,36,38,39]. Then it is possible to factorize the NERPA transition rate as $P_{\rm N}^{j',j} = \beta_{\rm N}^{j',j} P_\gamma$, where P_γ is the rate for the state reached by NERPA to decay by γ emission into the initial state.

For the coefficient $\beta_{\rm N}^{j',j}$ one may obtain expressions for any electric ($\lambda = E$) or magnetic ($\lambda = M$) nuclear transition multipolarity, i.e. for any (non-zero) values of the angular momentum L of the virtual photon:

$$\beta_{\rm N}^{j',j}(\lambda L) = \frac{[J_f]}{[J_i, j', j]} \sum_{\kappa\kappa'} \frac{4\pi\alpha\omega}{L(L+1)} s \left|\kappa\kappa'\right| \rho_\lambda, \quad (2)$$

where the radial part for the $\lambda = E$ case is $\rho_E = |(\kappa-\kappa')(R_3+R_4) + L(R_1+R_2+R_3-R_4)|^2$, and for the $\lambda = M$ case $\rho_M = |(\kappa+\kappa')(R_5+R_6)|^2$. We have also introduced the following notation in terms of a 3-j symbol [40]: $s = \begin{pmatrix} j & j' & L \\ \frac{1}{2} & -\frac{1}{2} & 0 \end{pmatrix}^2$. For the Dirac angular momentum quantum number κ (κ') of the positron (electron) holds: $||\kappa| - |\kappa'|| \le L \le |\kappa| + |\kappa'| - 1$, $j = |\kappa| - 1/2$, $j' = |\kappa'| - 1/2$. The radial integrals R_a, $a \in \{1, \ldots, 6\}$ are defined as in [35] with the analytical form of the radial Coulomb–Dirac wave functions for bound and free particles. All results are obtained here for the $1s_{1/2}$ electron orbital, yielding the highest rate of NERPA.

Fig. 2 shows the calculated dependence of the NERPA coefficient $\beta_{\rm N}^{j',j}$ for an electric-quadrupole ($E2$) nuclear transition between $1/2+$ and $5/2+$ states on the virtual photon energy ω in ^{115}Sn. Curves are given for two different possible positron angular momenta $j = L \pm 1/2$; both values $\beta_{\rm N}^{j',L\pm 1/2}$ have to be taken into account in the calculation of the cross section. One can see that nuclear transitions about 2 MeV, corresponding to the maximum of these curves, are preferred. This behaviour is analogous to the energy dependence of the bound-free pair creation cross section in heavy ion collisions [41].

A promising approach to observe the NERPA process is to measure the photon emitted by the decaying nucleus, but it is a comparably weak effect superimposed on the background of other photon emission processes in the studied system. To circumvent this problem, as in Refs. [6,10], we consider nuclei with a long-living state, which can be populated by some radiative transition from the state excited by NERPA. The decay of this metastable