11-GeV electrons. However, the aim of their work was different, and it is difficult to establish anything about the diurnal variation in the time-of-flight difference between the electrons and photons from their papers as published.

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Orientation of nucleus excited in annihilation of positron on heavy-atom shell

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A calculation is made of the cross sections and orientation spin tensors of a nucleus excited in nonradiative annihilation of a positron by an electron from an (nlj) shell of an atom. Use is made of relativistic wave functions of the electron and positron, obtained by numerical integration of the Dirac equation with a single average Hartree–Fock–Slater atomic potential. The results of calculations of the factors governing the cross sections and orientation of the excited nucleus as a function of the positron energy $E_+ = 0.5$–6.55 MeV are presented graphically for the series of multipole transitions $E0$, $E1$, and $M1$ in annihilation on the K shell of an atom. The results are given for the atoms with $Z = 41,49,60,70,82$, and 92. Estimates are also obtained for the annihilation of a positron on the L shell of lead ($Z = 82$). The results are used to estimate the cross sections of nuclear excitation in positron annihilation and the angular distributions of the decay products (quanta, fission fragments, etc.) of a nucleus excited in this way. A brief discussion is made of possible applications of low-energy positron beams in nuclear spectroscopy.

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§1. INTRODUCTION

When a positron collides with an atom, we can expect not only two- and one-photon annihilation of the positron by an atomic-shell electron but also excitation of the nucleus of the atom as a result of nonradiative EL or ML transition from the ground nuclear state E_1 to a level E_2 is the sum of the positron (E_+) and (nlj)-shell electron (E_{nlj}) energies, agrees—within the width Γ—with the nuclear transition energy $E_2 - E_1$. This nuclear excitation process is of the resonance type and the width Γ is governed by the lifetime of the excited state of the system; in the case of annihilation on the K shell the value of Γ is usually equal to the width of a hole $\Gamma(K)$, but there are some exceptions to this rule.[1)] The resonance nature of the nonradiative excitation of a nucleus as a result of positron annihilation provides a new potential method for investigating the nuclear structure, whose practical implementation required a fairly strong "monochromatic" positron beam of controlled energy E_+; we shall be concerned with the range of low nuclear excitation energies $E^* < 10$ MeV.

Our earlier[1,2] results of calculations of the cross sections of some nuclei in the $40 < Z \leqslant 92$ range show that it is desirable to have positron beams with a cur-

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rent of the order of ~10^{-6} A. It is not our task to consider the technical feasibility of producing such a beam but it seems desirable to discuss possible ways and ranges of applications of such a beam in low-energy nuclear physics, which may stimulate interest in this branch of experimental physics. The proposed applications naturally do not exhaust all possible uses of positron beams.

1. There are obvious applications of a positron beam in generating a similarly "monochromatic" flux of quanta as a result of one-photon positron annihilation (OPPA) on the K and L shells of a heavy atom. Such a photon flux can be used in the nuclear spectroscopy exactly as the photons resulting from an $(n\gamma)$ reaction. Very detailed investigations of the OPPA process involving the K, LI, LII, and LIII shells is reported by Broda and Johnson[3] who (by way of illustration) give a number of specific results of a numerical calculation of the OPPA cross sections for positron energies in the range $E_+ \leqslant 1.75mc^2$. We shall reproduce here some of the results from Ref. 3 for OPPA on the lead ($Z=82$) atomic shells when the positron energy is $E_+=1.5mc^2$:

Shell:	K	LI	LII	LIII
σ, mb:	735.27	109	35.8	3.28

According to Broda and Johnson, the angular distribution of the OPPA photons (in the $Z=82$ case) has a very sharp maximum in the range of small angles of photon emission relative to the positron momentum $\mathbf{p}_+$ and exhibits a very fast fall of the cross section on increase of this angle relative to $\mathbf{p}_+$.

2. We shall consider a different process which is the direct (i.e., without a preliminary conversion of a positron beam into a beam of photons) excitation of a nucleus by positron annihilation on the electrons of an (nlj) shell of an atom. In contrast to OPPA, we shall call this process nonradiative positron annihilation in accordance with the initial stage of the process involving excitation of a nonradiative (EL or ML) nuclear transition $E_1\Pi_1I_1 \to E_2\Pi_2I_2$; here and later, Π_iI_i denote the parity and spin of a nuclear level E_i. The subsequent decay of a nucleus from the E_2I_2 level may occur in various ways: it may evolve emission of a photon cascade, neutron evaporation, or fission. A transition is usually possible to the ground state by the emission of one photon, whose energy $\hbar\omega$ is equal to the energy of the OPPA process.

We reported earlier[1,2] a calculation of the cross sections of the $E0$, $E1$, $E2$, $M1$, and $M2$ nonradiative annihilation process for nuclei in the range $40 < Z \leqslant 92$ with energies 0.55 MeV $\leqslant E_+ \leqslant$ 6.55 MeV. In the case of the strongest $E1$ nonradiative annihilation process on the K shell of a heavy atom the cross section is of the order of 50–100 mb provided the $E1$ nuclear transition $E_1\Pi_1I_1 \to E_2\Pi_2I_2$ is characterized by a reduced probability $B(E1; I_1 \to I_2)$ which is of the same scale as the single-particle (sp) estimate obtained by Weisskopf[4]:

$$[B(E1)]_{sp} = \frac{3}{4\pi}e^2R_0^2\left(\frac{3}{4}\right)^2.$$

Systematics of the experimental data shows that the $E1$ transitions of nuclei in the excitation energy range 1 MeV $\leqslant E^* \leqslant$ 7 MeV are usually hindered very greatly (by a factor of the order of $10^{-3}-10^{-5}$) compared with the single-particle Weisskopf estimate. Therefore, in determining the intensity of the nonradiative annihilation process one has to consider cross sections of the order of $10^{-29}-10^{-30}$ cm^2, which are characteristics of such $E1$ transitions and of the normal unhindered $E0$, $M1$, and $E2$ transitions in a nucleus (see Ref. 2). For example, in the $40 < Z$ range the nonradiative process should have a cross section which is a factor of ~10^{-4} smaller than the OPPA cross section, i.e., the nuclear transitions due to the nonradiative process have to be separated from the background of a very intensive flux of OPPA quanta. Nevertheless, we can identify a number of nonradiative processes which are of considerable interest in nuclear physics because of certain factors favorable for distinguishing these processes against the OPPA background:

a) subbarrier fission in the nonradiative process;

b) excitation of spontaneously fissioning isomers of heavy nuclei with $Z=92$, 94, 95, and 96 when the energy is $E^* \sim 2-3$ MeV and determination of the spin of these isomeric states;

c) evaporation of a neutron from a nucleus excited in the nonradiative positron annihilation process: the interest lies in the range of the dense spectrum of nuclear levels at the binding energy of a neutron;

d) filling of low-lying isomeric states by a cascade of transitions developing after the excitation of a nucleus: in contrast to an $(n\gamma)$ cascade, we can in this case vary in a controlled manner the initial nuclear excitation energy $E^* = E_+ + E_{nlj}$ (attempts to observe this process in the case of $^{115}_{49}$In have been made[5,1] on the basis of the isomer yield);

e) observation of a cascade of characteristic photons emitted by a nucleus excited in the nonradiative process to a fixed level E_2I_2. For example, in the case of excitation energies of even nuclei $2\Delta \leqslant E^* \leqslant 3\Delta$ ($\Delta \sim 0.8$ MeV), usually attributed to the breaking of a nucleon pair, there is usually a group of 1+ levels (for example, in the case of ^{196}Pt as reported in Ref. 6). Excitation in the nonradiative positron annihilation of a system of these 1+ levels and observation of characteristic transitions to collective 2+ levels (vibrational or rotational) can give very important information on the structure of the nuclear energy levels in this range of E^*.

Another similar example is the excitation of even nuclei to the first 3+ levels, which are attributed (in the vibrator model) to three-phonon excitation of a nucleus. It is interesting to measure the probability of the $B(M3; 0 \to 3)$ transition.

3. The above processes can, of course, be produced also by a beam of "monochromatic" photons obtained by conversion of a positron beam; the exception to this rule is only the $E0$ excitation of a nucleus and clearly the excitation of higher multipoles. In each specific variant one can legitimately consider the comparative efficiency of the use of a photon beam after

positron conversion or the direct use of the nonradiative positron annihilation process. The selection between these two approaches can be made on the basis of the results of calculations relating to the nonradiative process and more detailed calculations of the OPPA cross sections (carried out in a wider range of E_+ and for a larger number of nuclei) than those given in Ref. 3. We shall not discuss this problem here and we shall analyze in detail the process of nonradiative positron annihilation.

In processes of the a)–c) type one records a fragment or a neutron, and the contribution of OPPA is proportional to the target thickness, whereas the contribution of the nonradiative process is linear in respect of this thickness, so that these processes can be separated when an intense positron beam is directed onto a thin target. In type e) reactions there is a contribution to the background from two-photon positron annihilation whose characteristics in the case of light atoms ($Ze^2/\hbar c \ll 1$) were considered recently in detail by Gorshkov *et al.*[7]; we are not aware of a similar investigation in the $Ze^2/\hbar c \sim 1$ case.

An important and favorable factor for the observation of the type e) nonradiative process is the radical difference between the angular distribution of the OPPA photons and the photons emitted by a nucleus excited and oriented in the nonradiative process (the angular distribution is measured relative to the momentum of the incident positron $\mathbf{p}_+$). The OPPA photons have a narrow distribution and they deviate through small angles relative to the positron momentum, whereas the angular distribution of photons resulting from the nuclear transitions cascade $E_1I_1 \xrightarrow{\text{nr}} E_2I_2 \xrightarrow{\hbar\omega} E_3I_3$ ("nr" refers to the nonradiative process) is symmetric relative to the plane perpendicular to $\mathbf{p}_+$ and although the anisotropy of the distribution is considerable, it is not as pronounced as in the case of the OPPA photons. The orientation of a nucleus excited by the nonradiative process is naturally manifested also in the anisotropy of the angular distribution of other products of decay of excited nuclei (fission fragments, neutrons, etc.).

4. We shall give the results of a numerical calculation of the charactertistics of the orientation of nuclei excited by the nonradiative annihilation of a positron and an electron from an (nlj) atomic shell. All the quantities necessary for numerical estimates of the effects will be given in graphical form. We shall confine our attention to the most important process of the annihilation on the K shell of a heavy atom with $40 < Z \leq 92$. More detailed numerical results, including the data for the L shell, will be given in a separate communication. Our calculations were carried out only for the orientation of a nucleus described by a set of even spin tensors of the excited nucleus; if a polarized positron is annihilated, a nucleus excited in the nonradiative positron annihilation process becomes polarized, i.e., the odd spin tensors of the nucleus also become finite. Calculations for the latter case can also be made but there is not immediate need for this.

5. The process of excitation of an $I_1 \to I_2$ nuclear transition by nonradiative annihilation of a positron and an electron from an (nlj) atomic shell is the converse of the pair conversion of a nuclear ΛL multipole by an $I_2 \to I_1$ transition which occurs when there is at least one hole in an (nlj) atomic shell. Then, an electron from such a pair is captured by the hole in the (nlj) shell and a "monochromatic" positron is emitted with an energy distributed in an interval of the order of the width Γ near a value $\bar{E}_+$, determined by the law of energy conservation: $\bar{E}_+ = (E_2 - E_1) - E_{nlj}$.

Sliv was the first to consider this variant of pair conversion[8] and he carried out calculations using wave functions of an electron in the field of a bare point nucleus. We can give a refined estimate of the probability of this process. Following Ref. 8, we shall introduce $w_\mu(\Lambda L[nlj]^{-1}, I_2 \to I_1, \bar{E}_+)$; to denote the probability of emission of a "monochromatic" positron per unit time and per one hole in an (nlj) atomic shell (this probability is integrated over the positron spectrum); Sliv[8] denoted this probability by w_μ^K for the K shell. In this approach it is assumed that a hole has an indefinite lifetime and then the finite lifetime of the hole is allowed for by introducing relative probabilities of the competing processes. The principle of detailed equilibrium relates the value of w_μ to the cross section of the nonradiative positron annihilation on the same shell. Using Eq. (15) of Ref. 1, which gives the cross section for the nonradiative process, we find that w_μ is given by

$$w_\mu(\Lambda L[nlj]^{-1} I_2 \to I_1 : \bar{E}_+) = \frac{1}{\pi^2} \frac{me^4}{\hbar^3} |\langle I_2 E_2 \| \Lambda L \| I_1 E_1 \rangle|^2 p_+^{\,2} a_0^{N(L)} \xi(\Lambda L[nlj]^1 \bar{E}_+),$$

where $\bar{E}_+ = [m^2c^4 + \hbar^2c^2p_+^2]^{1/2}$; p_+ is the wave number of a positron whose energy is $\bar{E}_+$.

Here and later, we shall use the same notation as in Refs. 1 and 2; the formula for the nonradiative annihilation cross section can also be obtained from Eq. (10) for σ_{Q0} with $Q = 0$; the quantities occurring in that equation are explained immediately below it.

In comparison with the probability of the usual pair conversion w_π of a nuclear ΛL multipole ($L \neq 0$) in an $I_2 \to I_1$ nuclear transition we shall introduce the "monochromatic" pair-conversion coefficient $\alpha_\mu(\Lambda L[nlj]^{-1}\bar{E}_+)$ expressing the nuclear matrix elements in terms of the probability of a radiative $I_2 \to I_1$ nuclear transition of the ΛL type. In the specific case of the EL multipoles ($L \neq 0$) and one hole in an (nlj) atomic shell we obtain ($\hbar\omega \equiv E_2 - E_1$):

$$\alpha_\mu(EL[nlj]^{-1}\bar{E}_+) = \left(\frac{1}{2\pi}\right)^3 \frac{L}{L+1} [(2L+1)!!]^2 \times \left(\frac{e}{\hbar c}\right)^{2L-1} \left(\frac{mc^2}{\hbar\omega}\right)^{2L+1} \left[\left(\frac{\bar{E}_+}{mc^2}\right)^2 - 1\right] \xi(EL[nlj]^1\bar{E}_+).$$

The $\xi(\Lambda L[nlj]^1 E_+)$ factors were calculated by us earlier for a number of nuclei (see Refs. 1 and 2) and they are also given in the figures below.

Sliv calculated w_μ^K and the probability of the usual pair conversion w_π for the $E1$ and $E0$ processes in the RaC′ nucleus using the electron positron functions in the field of a point nucleus and ignoring the field screening of the electron shell of the atom, whereas we calculated the $\xi(\Lambda L[nlj]^1 E_+)$ factors allowing for all these effects. We shall not make comparison with the case of the $E0$ transition because it is well known that in this case the effect of the finite dimensions of a nucleus is very large.

In the case of an electric dipole of the RaC′ nucleus ($Z=84$) and the photon energy $\hbar\omega=1400$ keV, the ratio of these two probabilities is $w_{\mu}^{K}/w_{\tau}=\frac{1}{3}$. Clearly, this is not a rational fraction because the ratio varies with the transition energy.[8] Numerical accuracy of this quantity is not known. For a bare point nucleus with $Z=84$ for the photon energy $\hbar\omega=3mc^2$, the usual pair conversion coefficient of a heavy nucleus (hn) is [9]

$$[\alpha_{\pi}(E1)]_{\text{hn}}=1.84\cdot10^{-4}.$$

Hence, we can reconstruct the coefficient of monochromatic pair conversion for a bare point nucleus ($Z=84$) at $\hbar\omega=1400$ keV:

$$[\alpha_{\mu}(E1[K]^{-1})]_{\text{hn}}\approx0.61\cdot10^{-4}.$$

A calculation of α_{μ} in accordance with our formula allowing for the finite size of the nucleus and the screening effect gives

$$\alpha_{\mu}(E1[K]^{-1})=0.74\cdot10^{-4}.$$

This value is somewhat overestimated because linear interpolation in the estimate of the factor ξ for $Z=84$ is used. It seems to us that the agreement between the two values of $\alpha_{\mu}(E1)$ is reasonable; thus, the correction for the screening effect in the OPPA process in a nucleus with $Z=82$ and for $E_{+}=1.75mc^2$ (see Ref. 3) also increases the cross section (by about 3%).

§2. ORIENTATION SPIN TENSORS OF A NUCLEUS EXCITED IN NONRADIATIVE POSITRON ANNIHILATION

1. We shall begin by defining various quantities used and calculated below.

Let us assume that the state of a nucleus with an angular momentum **I** is given by a superposition of the magnetic numbers M along a selected quantization axis:

$$\Psi_I=\sum_M a_M(\mathbf{I})\,|IM\rangle, \tag{1}$$

where the amplitudes $a_M(\mathbf{I})$ are normalized to unity:

$$\sum_M |a_M(\mathbf{I})|^2=1. \tag{2}$$

The polarization and orientation of a nucleus in state (1) is usually described by introducing a system of spin tensors of rank Q:

$$\rho_{Q\nu}(\mathbf{I})=\sum_{MM'} a_{M'}^{*}(I)\,a_M(I)\,(IQM\nu|IQIM'), \tag{3}$$

TABLE I. Nuclear levels of $^{208}_{82}$Pb from Ref. 13 and calculated resonance cross sections for excitation of these levels due to annihilation of positron by filled K shell.

E_2, MeV	$I_2\Pi_2$	ΛL	Γ_N, eV	$B(\Lambda L;\uparrow)$, $10^{-3}\,e^2\cdot$bn	σ_{res}, 10^{-3} bn
5.514	1−	$E1$	22	3.8	6.2
6.720	1−	$E1$	14.5	1.4	2.5
7.061	1+	$M1$	23	1.90	3.5
7.080	1−	$E1$	12	0.98	1.8
7.333	?	{ $M1$	48	3.5	4.9
		{ $E1$	48	3.5	4.2

TABLE II. Comparison of $\xi(\Lambda L[nlj]^1E_+)$ factors for $1s$ and $2s$ electrons of lead atom.

ΛL	$[nlj]$	$\xi(\Lambda L[nlj]E_+)$					
		$E_+=0{,}55$ MeV	1,03 MeV	1,51 MeV	2.23 MeV	3,19 MeV	3.67 MeV
$E0$	$1s$	$5.45\cdot10^3$	$2.04\cdot10^5$	$2.48\cdot10^5$	$2.57\cdot10^5$	$2.49\cdot10^5$	$2.45\cdot10^5$
$E0$	$2s$	$8.44\cdot10^2$	$3.15\cdot10^4$	$3.84\cdot10^4$	$3.97\cdot10^4$	$3.85\cdot10^4$	$3.79\cdot10^4$
$E1$	$1s$	0.311	4.21	4.18	3.98	3.82	3.77
$E1$	$2s$	$4.74\cdot10^{-2}$	0.615	0.596	0.555	0.525	0.516
$E2$	$1s$	78.1	$8.16\cdot10^3$	$1.97\cdot10^4$	$3.34\cdot10^4$	$8.77\cdot10^4$	$1.15\cdot10^5$
$E2$	$2s$	12.1	$1.20\cdot10^3$	$2.83\cdot10^3$	$6.09\cdot10^3$	$1.21\cdot10^4$	$1.58\cdot10^4$
$M1$	$1s$	$3.25\cdot10^{-2}$	1.38	1.82	2.03	2.11	2.12
$M1$	$2s$	$4.48\cdot10^{-2}$	0.189	0.247	0.275	0.284	0.285

where $(ABab|ABDd)$ are the Clebsch–Gordan coefficients in the notation of Condon and Shortley[10] (Condon and Shortley's monograph[10] also has tables). We shall assume that the target is not oriented, i.e., that in the initial state E_1I_1 an ensemble of the target nuclei can be described by a single spin tensor $\rho_{00}(I_1)=1$ of rank zero, whereas all the other spin tensors $\rho_{Q\nu}(I_1)$ with $Q\neq0$ are identically equal to zero. In the nuclear transition $E_1I_1\to E_2I_2$, excited by nonradiative annihilation of a polarized positron beam, there is only one preferred axis in space and that coincides with the direction of the positron momentum $\mathbf{p}_+$. In this case all the odd spin tensors of an excited nucleus $\rho_{Q\nu}(I_2)$ with $Q=1, 3, 5, \ldots$ vanish with precision of the same order as the neglected (by us) effects of parity nonconservation of nuclear states. In this case the nonvanishing even spin tensors ($Q=0, 2, 4, \ldots$), which govern the orientation of an excited nucleus, must obey the axial symmetry condition

$$\rho_{Q\nu}(I_2)=\rho_{Q0}(I_2)\,\delta_{\nu0}. \tag{4}$$

The corresponding angular distributions of the photons emitted from a nucleus (excited by nonradiative positron annihilation) in the subsequent $E_2I_2\to E_3I_3$ transition of multipole order ΛL (EL or ML) is consequently given by the familiar formula

$$W(\theta_k)=\frac{1}{4\pi}\sum_{Q=0,2,4\ldots}(2Q+1)\rho_{Q0}(I_2)P_Q(\cos\theta_k)\,(LQ10|LQL1)\,u(I_2I_3QL;LI_2), \tag{5}$$

where θ_k is the angle between the wave vector of a photon $\mathbf{k}$ and positron momentum $\mathbf{p}_+$; $u(abcd;ef)$ is the normalized Racah function introduced by Jahn[11] (convenient algebraic tables of this function are also given in Ref. 11):

$$u(abcd;\ ef)=[(2f+1)(2e+1)]^{1/2}W(abcd;\ ef). \tag{6}$$

Equation (5) can be generalized in an obvious manner to the case of a mixed $E(L+1)+ML$ nuclear transition $E_2I_2\to E_3I_3$; we shall not reproduce it here.

In the case of fission of a nucleus excited to the E_2I_2 state the spin tensors $\rho_{Q\nu}(I_2)$ defined in this way govern also the angular distribution of the fission fragments relative to the positron momentum $\mathbf{p}_+$; the corresponding formulas can be found, for example, in Ref. 12. In considering a cascade of photon emitted by a nucleus from an initial excited state E_2I_2 we have to consider the transfer of a spin tensor along a chain of transitions.

2. Excitation and orientation of a nucleus in an $E_1 I_1 \to E_2 I_2 M_2$ transition of multipole order ΛL (ML or EL), stimulated by nonradiative annihilation of a positron and an (nlj) electron, in the case of an unoriented target and unpolarized positron beam, will be described by a cross section $\sigma(\Lambda L[nlj]^1 E_+ \mathbf{p}_+; I_1 \to I_2 M_2)$ reduced to one electron; the momentum p_+ defines the quantization axis. It is convenient to introduce a system of spin tensors which are the σ_{Q0} cross sections in accordance with the definition

$$\sigma_{Q0}(\Lambda L[nlj]^1 E_+ \mathbf{p}_+; I_1 \to I_2) = \sum_{M_2} (I_2 Q M_2 0 | I_2 Q I_2 M_2)\, \sigma(\Lambda L[nlj]^1 E_+ \mathbf{p}_+; I_1 \to I_2 M_2) \tag{7}$$

as well as normalized to unity spin tensors of the orientation of a nucleus in an excited state $E_2 I_2$

$$\rho_{Q0}(I_2) = \frac{\sigma_{Q0}(\Lambda L[nlj]^1 E_+ \mathbf{p}_+; I_1 \to I_2)}{\sigma(\Lambda L[nlj]^1 E_+; I_1 \to I_2)}. \tag{8}$$

The spin tensors $\rho_{Q0}(I_2)$ are identical with those described by Eqs. (3) and (4) above.

Equation (8) includes the positron annihilation cross section $\sigma(\Lambda L[nlj]^1 E_+; I_1 \to I_2)$, defined by us earlier in Eq. (15) of Ref. 1 or in Eq. (9) of Ref. 2; this cross section is a spin tensor of zero rank ($Q=0$) but for simplicity we shall omit the index of this tensor.

In considering reactions caused by a strongly nonmonochromatic positron beam with a given distribution $S(E_+)$ in the energy scatter range ΔE_+

$$\int_{\Delta E_+} S(E_+)\, dE_+ = 1, \tag{9}$$

when the nucleus has a dense energy level spectrum so that the interval ΔE_+ covers several nuclear levels $E_2 I_2 \Pi_2$, we can calculate the angular distribution of the reaction products using naturally the spin tensors representing the cross sections σ_{Q0} integrated over the distribution $S(E_+)$ allowing for the contribution of each level to the observed decay channel. The general formula for this variant will not be reproduced here: it is self-evident in each specific case.

For the spin-tensor cross section σ_{Q0} of the annihilation of a monochromatic positron accompanied by the excitation of a ΛL nuclear transition $E_1 I_1 \to E_2 I_2$, we shall use the notation adopted earlier in Refs. 1 and 2:

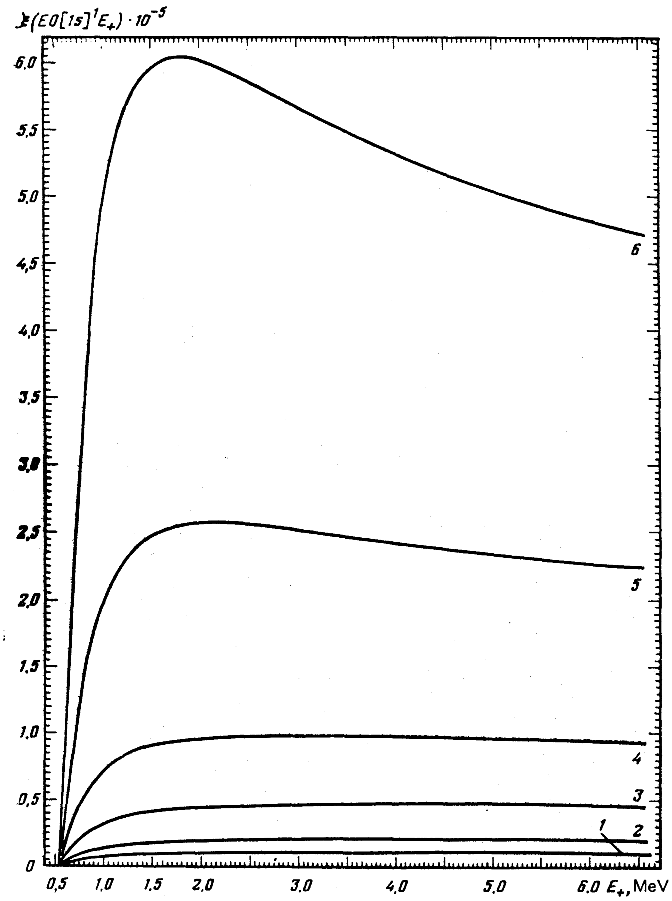

FIG. 1. Cross section factor ξ of the $E0$ transition plotted as a function of the energy of a positron annihilated by a K-shell electron of atoms with the following values of Z: 1) 41; 2) 49; 3) 60; 4) 70; 5) 82; 6) 92.

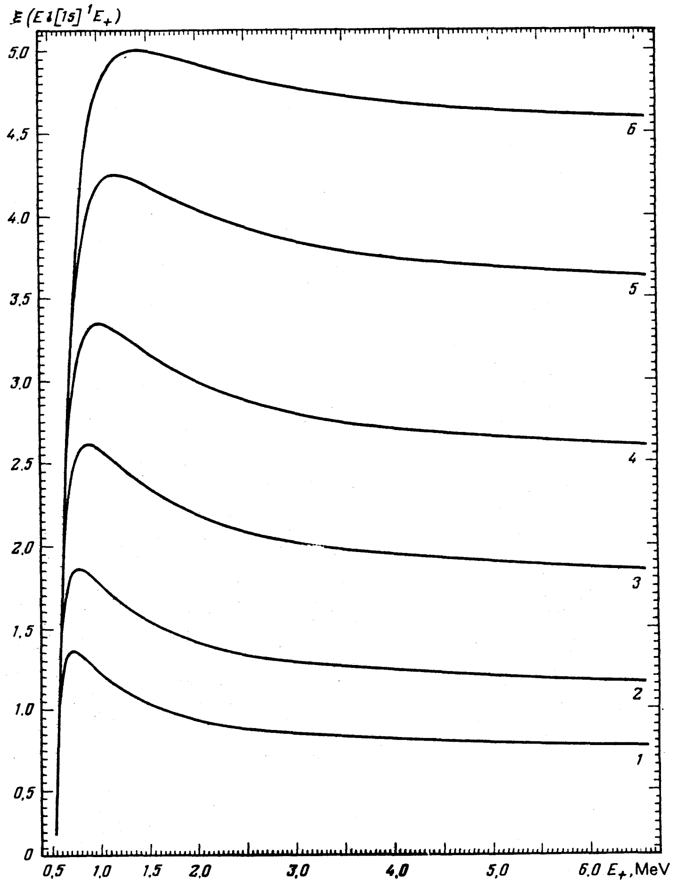

FIG. 2. Cross section factor ξ of the $E1$ transition plotted as a function of the energy of a positron annihilated by a K-shell electron of atoms with the following values of Z: 1) 41; 2) 49; 3) 60; 4) 70; 5) 82; 6) 92.

$$\sigma_{Q0}(\Lambda L[nlj]^1E_+\mathbf{p}_+;I_1\to I_2)$$
$$=a_0^{N(L)}\left(\frac{2I_2+1}{2I_1+1}\right)|\langle I_2E_2\|\Lambda L\|I_1E_1\rangle|^2D(E_2-E_1-E_{nlj}-E_+)$$
$$\times u(I_1LI_2Q;I_2L)\xi_Q(\Lambda L[nlj]^1E_+), \quad (10)$$

where

$$D(x)=\frac{1}{2\pi}\frac{I_0\Gamma}{x^2+(\Gamma/2)^2}, \quad (11)$$

$$\left.\begin{array}{l} a_0=\hbar^2/me^2\approx 0.529\cdot 10^{-8}\,\text{cm},\ I_0=me^4/\hbar^2=27.2\,\text{eV}, \\ N(L)=\begin{cases}-2 & \text{for } E0\\ -2L+2 & \text{for } ML \text{ and } EL \text{ for } L\neq 0\end{cases}\end{array}\right\} \quad (12)$$

The matrix elements $\langle I_2E_2\|\Lambda L\|I_1E_1\rangle$ of EL and ML nuclear $I_1E_1\to I_2E_2$ transitions are related to the reduced probabilitiies $B(\Lambda L;I_1\to I_2)$ adopted in the literature[4] by

$$\left.\begin{array}{l} e^2\left(\frac{2I_2+1}{2I_1+1}\right)|\langle I_2E_2\|EL\|I_1E_1\rangle|^2=B(EL;I_1\to I_2), \\ e^2\frac{L}{L+1}\left(\frac{2I_2+1}{2I_1+1}\right)|\langle I_2E_2\|ML\|I_1E_1\rangle|^2=B(ML;I_1\to I_2). \end{array}\right\} \quad (13)$$

The atomic-positron factors $\xi_Q(\Lambda L[nlj]^1E_+)$, per one electron in an (nlj) shell do not contain any unknown (nuclear) quantities and they can be calculated and tabulated. In the case of a spin tensor of rank zero representing a cross section (i.e., in the case when $Q=0$), the factor ξ_0 is naturally identical with the factor $\xi(\Lambda L[nlj]^1E_+)$ for the nonradiative positron annihilation cross section introduced earlier by Eqs. (9)–(11) of Ref. 1. The formulas for the orientation factors ξ_Q with $Q\neq 0$ $(Q=2,4,\dots)$ are more cumbersome than for the ξ factor; naturally, they contain all the same radial integrals and phases for the scattering of a positron by the screened Coulomb field of the nucleus. These formulas are the opposite of clear and we shall not give them because only the values of the factors ξ_Q are of practical interest.

In estimating the effects of the orientation of a nucleus in nuclear reactions caused by nonradiative positron annihilation it is convenient to use the coefficients $A_Q(\Lambda L[nlj]^1E_+)$ in accordance with the equation

$$\xi_Q(\Lambda L[nlj]^1E_+)=A_Q(\Lambda L[nlj]^1E_+)\xi(\Lambda L[nlj]^1E_+). \quad (14)$$

The results of calculations of the factor ξ and of the coefficients $A_Q(\Lambda L[nlj]^1E_+)$, together with Eqs. (10)–(14), represent thus the solution of the problem formulated above.

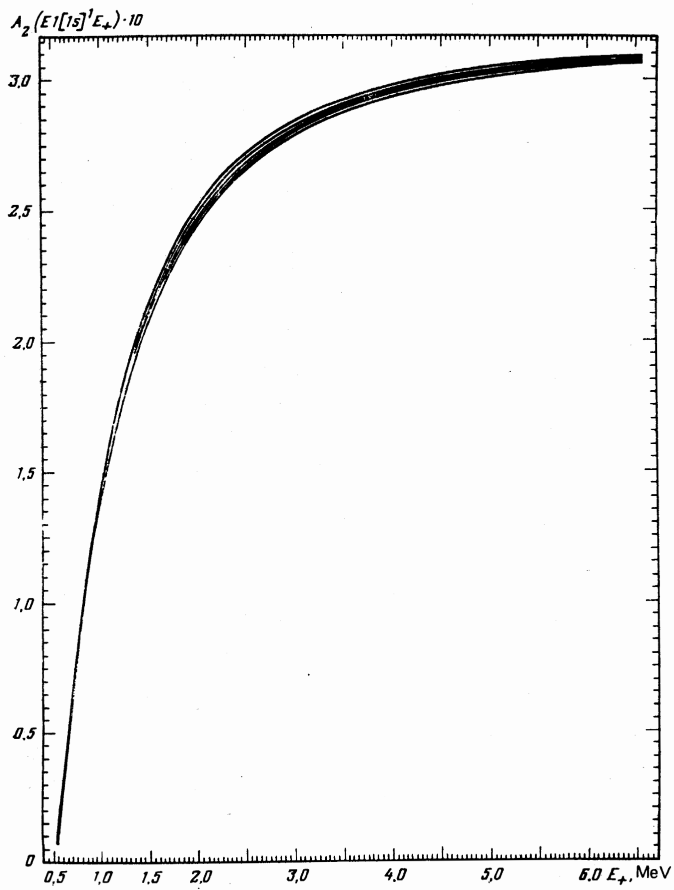

FIG. 3. Orientation coefficient A_2 plotted as a function of the energy of a positron annihilated by a K-shell electron of atoms with $Z=41$–92. The dependence on the charge Z is weak. The upper limit corresponds to $Z=92$ and the lower limit to $Z=41$.

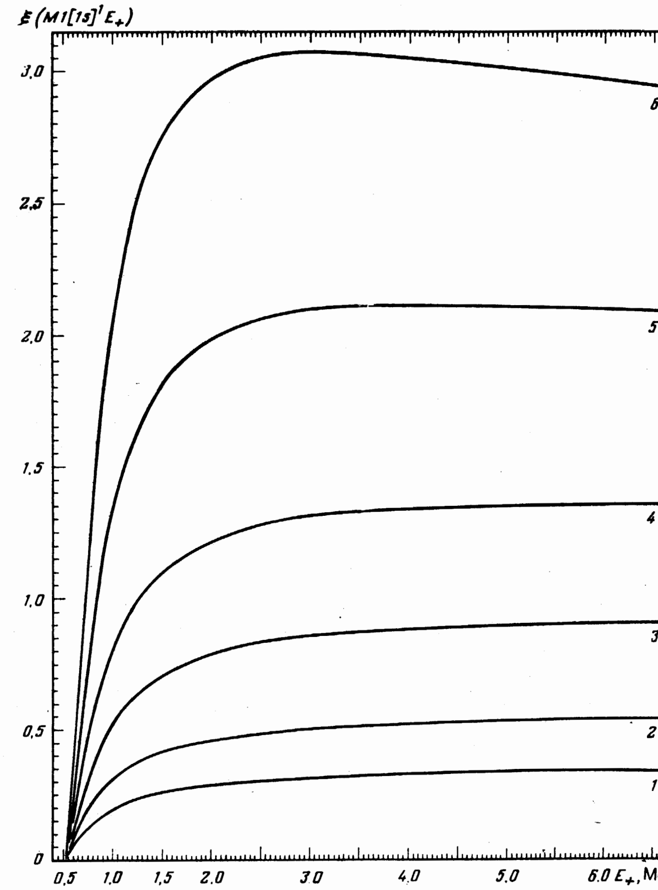

FIG. 4. Cross section factor ξ of the $M1$ transition plotted as a function of the energy of a positron annihilated by a K-shell electron of atoms with the following values of Z: 1) 41; 2) 49; 3) 60; 4) 70; 5) 82; 6) 92.

3. The resonance nature of the cross section (10) is given by the function $D(x)$, which includes the width Γ of the final state that the system assumes as a result of nonradiative positron annihilation. As a rule, Γ is governed by the lifetime of a hole in an (nlj) electron shell; for example, in the case of a K hole and $Z=80$–90 the width $\Gamma(K)$ is of the order of ~100 eV, which is considerably greater than the width of nuclear levels Γ_N. However, there are exceptions: for example, a group of 1+ and 1− levels with a large width Γ_N, due to radiative transition to the ground state 0+ of the $^{208}_{82}$Pb nucleus, is reported in Ref. 13. We calculated the cross sections for the excitation of these levels by nonradiative positron annihilation processes ($E1$ and $M1$) at a point E_+ corresponding to a resonance. The results are given in Table I. The cross sections of these nonradiative positron annihilation processes reach ~2–6 mb, which is naturally much smaller than the OPPA cross sections but the angular distribution of the $M1$ and $E1$ photons resulting from transitions to the ground state from the 1+ and 1− levels, relative to the direction of the positron momentum $\mathbf{p}_+$,

$$W(\theta_k)=\frac{1}{4\pi}\left\{1+5\sqrt{\frac{1}{10}}A_2(\Lambda L[1s]^1E_+)P_2(\cos\theta_k)\right\} \quad (15)$$

is very different from the distribution of OPPA photons which falls rapidly in the rear hemisphere, which makes it possible to identify these nuclear states against the background of smooth variation of the OPPA cross section considered as a functions of E_+.

§3. RESULTS OF A NUMERICAL CALCULATION OF THE NUCLEAR ORIENTATION FACTORS

Our earlier calculations[1] were carried out for the $^{115}_{49}$In and $^{235}_{92}$U nuclei. Bearing in mind, the above-mentioned possible applications of positron beams in nuclear spectroscopy, we calculated the factors $\xi_Q(\Lambda L[nlj]^1E_+)$ in the nuclear charge range $40<Z\leq 92$ for positron energies 0.55 MeV $\leq E_+\leq$ 6.66 MeV.

All the calculations were carried out using relativistic wave functions of an atomic-shell electron and a positron, which were obtained by numerical integration of the Dirac equation with a single average Hartree–Fock–Slater atomic potential, which was found employing a program written by I. M. Band and M. B. Trzhaskovskaya.[14]

The most detailed calculations of the factor $\xi(\Lambda L[nlj]^1E_+)$ and the orientation coefficients $A_Q(\Lambda L[nlj]^1E_+)$ were carried out for the series of multipoles $E0$, $E1$, and $M1$ of the process of annihilation of a positron and a K electron. Selective calculations for the LI, LII, and $LIII$ shells of the lead atom indicated that the factor $\xi(\Lambda L[nlj]^1E_+)$ is considerably less than for a K electron and that gradation of the value of the factor ξ in the LI, LII, and $LIII$ subshells is approximately the same as for the OPPA process, i.e., the largest factor $\xi(\Lambda L)$ is for an LI electron but it is approximately an order of magnitude small than the corresponding factor for a K electron. This is illustrated for a number of energies E_+ in Table II, calculated for the lead atom.

It seems to us that the main task at present is to estimate the feasibility of using positron beams in nuclear spectroscopy and, therefore, the data for the strongest multipole transitions ($E0$, $E1$, $M1$, and $E2$) in nonradiative annihilation of a positron by a K-shell atomic electron are of practical interest. Consequently, Figs. 1–5 in the present paper give the results of numerical calculations of the factors $\xi(\Lambda L[1s]^1E_+)$ and of the nuclear orientation coefficients $A_Q(\Lambda L[1s]^1E_+)$ for the nonradiative positron annihilation processes of the lowest multipole order.

We must mention a characteristic feature of the nonradiative positron annihilation processes whose multipole orders are $E1$ and $E2$, discovered only as a result of numerical calculation: the cross section factors $\xi(EL[1s]^1E_+)$ have a definite dependence on the nuclear charge Z, and the coefficients $A_Q(EL[1s]^1E_+)$, can be seen from Fig. 3 to depend very weakly on Z throughout the range $40<Z\leq 92$, whereas in the case of a magnetic multipole ML a strong dependence on Z is exhibited by the factor $\xi(ML[1s]^1E_+)$, and by the coefficients $A_Q(ML[1s]^1E_+)$. These results allow us interpolation throughout the range of Z.

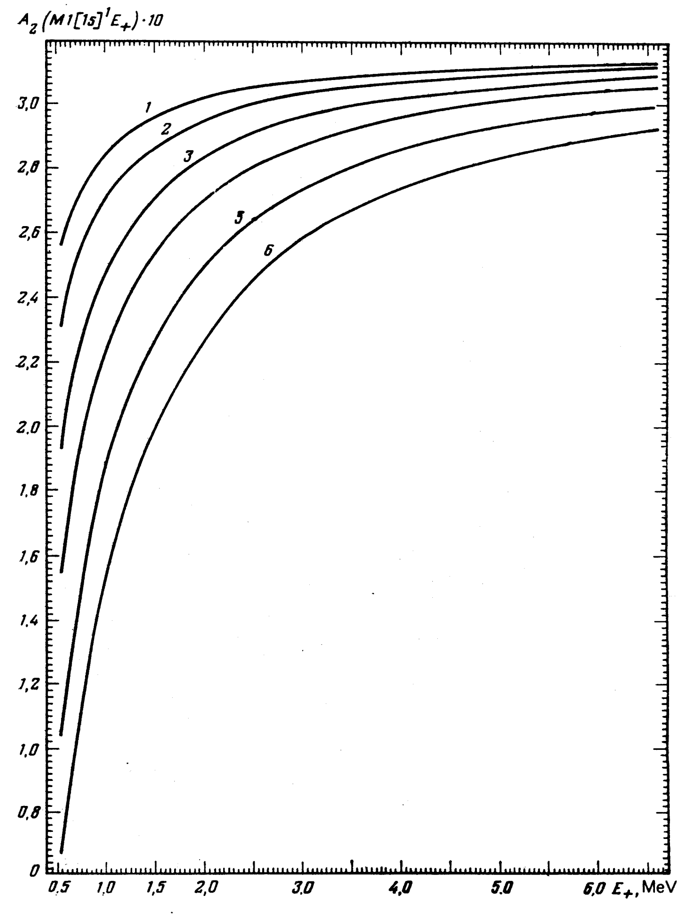

FIG. 5. Orientation coefficient A_2 plotted as a function of the energy of a positron annihilated by a K-shell electron of atoms with the following values of Z: 1) 41; 2) 49; 3) 60; 4) 70; 5) 82; 6) 92.

[1)]In the course of a nonradiative annihilation of a positron and a K electron the whole system (i.e., the nucleus and the electron shell of the atom) go over to an intermediate state which is in resonance with the initial state; in this intermediate state the nucleus is excited to a level E_2 and a hole appears in the K shell of an atom. The amplitude of this intermediate state of the whole system decays as a result of subsequent nuclear (radiative, conversion, nucleon evaporation, etc.) transitions from the level E_2 and also as a result of transitions in the electron shell (radiative or Auger) resulting in the filling of the K hole. Only in the case of hydrogen- and helium-like ions can we regard the K-hole lifetime as infinitely long. The width of a resonance state of the system is thus governed by the total contribution of the nuclear and electron processes: $\Gamma=\hbar/\tau_N+\hbar/\tau_K$, where τ_N is the lifetime of the nucleus at the level E_2 and τ_K is the lifetime of a K hole. As a rule, for $Z>40$ and $E_2<6$ MeV, we have $\tau_K \ll \tau_N$ but there may be exceptions (for example, the $^{208}_{82}$Pb nucleus).

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Effects of nonconservation of spatial and temporal parities in spectra of diatomic molecules

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It is shown that, after allowance for the Λ doubling, the degree of nonconservation of the spatial parity in the scattering of light by homonuclear diatomic molecules can reach values of the order of 10^{-4} for the hydrogen molecule (H_2) and of unity for the iodine molecule (I_2). The effects of nonconservation of the temporal parity (T-invariance violation) in the spectra of heteronuclear diatomic molecules with the Λ doubling of the ground states are enhanced by five orders of magnitude, compared with the atomic spectra, in electric fields of a few kilovolts per centimeter.

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Much theoretical work has been done recently on the effects of nonconservation of the spatial parity in atoms associated with the existence of weak neutral currents (see, for example, the reviews in Refs. 1–4). The results of the first experimental investigations have been published.[5-7] Similar effects in molecules have also been considered.[8,9] Quite a few papers have been devoted also to the search of the effects of nonconservation of the temporal parity in atoms and molecules, namely to the discovery of the influence of the dipole moments of an electron and a proton and of the constants of the scalar and tensor interactions violating the temporal parity on the dipole moments of atoms and molecules.[10-14]

We shall show that the effects of nonconservation of the spatial and temporal parity in diatomic molecules may be enhanced considerably in the case of the levels exhibiting the Λ doubling effect.

1. NONCONSERVATION OF THE SPATIAL PARITY

We shall consider resonance fluorescence emitted by the hydrogen molecule as a result of its transition from the ground para state $^1\Sigma_g^+(I=0,K=0;+)$ to the ground ortho state $^1\Sigma_g^+(I=1,K=1;-)$. Here, I is the total nuclear spin and K is the rotational quantum number; the sign in parentheses is the sign of the state (total spatial