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Polarization mechanism for nonconservation of parity and the effect of a weak neutral interaction in heavy μ-mesic atoms

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An estimate is made of the effect of nonconservation of parity of the $3d$-orbit meson states for mesic atoms with odd nuclei in the $Z=56$–60 and $Z\sim 83$ ranges in which the hyperfine structure terms of the $3p\,^3/_2$–$3d\,^5/_2$ and $3p\,^1/_2$–$3d\,^5/_2$ orbits respectively intersect. The angular distribution coefficient α in $W(\theta)=1+\alpha\cos\theta$ for quanta emitted by polarized mesic atoms in the $3d\to 1s$ transition is determined. The effect of a weak neutral interaction between a muon and a nucleus and the Coulomb polarization mechanism for transferring the nonconservation of parity of nuclear states to mesic atom states in the chain of nonradiative transitions of a meson are considered. In the nuclear ranges indicated above the effect of the weak interaction for $3d$- and $3p$-orbits is smaller by two–three orders of magnitude than the effect of the polarization mechanism. Under optimal conditions for the intersection of $3d\,^5/_2$ and $3p\,^3/_2$ terms of a mesic atom which, as is shown, can be realized in the range $Z=56$–60 this mechanism leads to a value of the coefficient α which is approximately equal to the amplitude for the nonconservation of parity for nuclear states $\beta(I_\nu E_\nu)$ lying at $E_\nu\approx\hbar\omega(3d\to 1s)$ and $E_\nu\approx\hbar\omega(3d\to 2p)$. If dynamic amplification of the nonconservation of parity of nuclear states occurs at $E_\nu\approx 5$–6 MeV, i.e., $\beta(I_\nu E_\nu)\sim 10^{-5}$, then for the quanta arising from a $3d\to 1s$ transition of a meson under real conditions for the transfer of polarization of the meson spin in the cascade of transitions populating the $3d\,^5/_2$-orbit the anisotropy coefficient α can attain a value of $\sim 10^{-5}$.

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INTRODUCTION

1. The weak neutral interaction between a muon and a nucleus leads to a mixing of mesic atom states of opposite parities. For the $2s\frac{1}{2}$ and $2p\frac{1}{2}$ meson orbits this effect has been investigated in a number of papers.[1–7] with the calculation in Ref. 2 being carried out for the range $3\leqslant Z\leqslant 82$. In the case of light ($Z<10$) mesic atoms the observation of the effect of the weak interaction can be significantly impeded by a number of accompanying processes: configuration mixing in the electron shell of the mesic atom, the Stark effect of the electric field of the medium, etc. (cf., Ref. 8), while for heavy mesic atoms the role played by these processes is insignificant. Since the amplitude of the admixture of the state of opposite parity is determined by the ratio of the matrix element of the interaction to the difference between the energies of the states being mixed, it is natural to seek in the spectrum of a mesic atom states of opposite parity close in energy. However, it is necessary that the meson should penetrate the volume of the nucleus sufficiently effectively if our aim is to determine the magnitude of the weak interaction between a muon and a nucleus. A preliminary calculation of the terms of a mesic atom in the Coulomb field of a uniformly charged sphere ($R_0=1.24\times 10^{-13}A^{1/3}$ cm) has picked out three ranges of Z in which the $3dj_1$ and $3pj_3$ terms of a mesic atom intersect:

Range of Z:	$55\leqslant Z\leqslant 60$	$65\leqslant Z\leqslant 70$	$82\leqslant Z\leqslant 85$
Mesic atom terms:	$3d^5/_2-3p^3/_2$	$3d^3/_2-3p^1/_2$	$3d^5/_2-3p^1/_2$
Nuclear spin:	$I\geqslant{}^1/_2$	$I\geqslant{}^1/_2$	$I\geqslant 1$

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As long as the angular momenta of the $3dj_1$ and $3pj_3$ states are different ($j_1 \neq j_3$), the effects of nonconservation of parity can be observed only in mesic atoms whose nuclei possess an angular momentum I different from zero, i.e., practically for odd nuclei. In this case in the hyperfine structure spectra of the $3dj_1$ and $3pj_3$ terms there exist $|3dj_1IF\rangle$ and $|3pj_3IF\rangle$ states close in energy and with the same total angular momenta $\mathbf{F}=\mathbf{j}+\mathbf{I}$ which are mixed when an interaction which does not conserve parity is "switched-on."

2. The magnitudes of the energies $\varepsilon(3dj_1IF)$ and $\varepsilon(3pj_3IF)$ of the intersecting terms are determined by the hyperfine interaction $\hat{H}(hfs)$ of a mseon with the magnetic dipole ($\mu(I)$) and the electric quadrupole ($Q(I)$) moments of the nucleus, and also by the effect of the polarization of the nucleus due to the Coulomb interaction between the nucleus and the meson. The correction to the energy of the mesic atom due to the nuclear polarization depends on the total angular momentum F of the system, particularly in the case when a resonance occurs between the EL-transition of the meson and the EL-transition of the nucleus. A reliable estimate of this correction is impossible to obtain, since it is necessary to have quite detailed information on the nuclear states over a broad range of excitation energy. In subsequent discussion in order to estimate the effects of nonconservation of parity we have utilized the results of the calculations of the energy $\varepsilon(nlj)$ of the levels of meson orbits with a realistic charge distribution published in the tables of Ref. 9. It should be noted that the correcrections given in Ref. 9 due to polarization of the nucleus are practically independent of the angular momentum j for $3dj_1$- and $3pj_3$-meson orbits and do not depend on the magnitude of the total angular momentum of the system F, i.e., these corrections have been obtained within the framework of a highly simplified nuclear model.

Since the spin-orbit splitting of the $3d\frac{3}{2}$ and $3d\frac{5}{2}$ orbits and of the $3p\frac{1}{2}$ and $3p\frac{3}{2}$ orbits is significantly greater than the magnitude of the hyperfine splitting of the $3dj_1$ and $3pj_3$ terms, then the interaction of the meson with the $\mu(I)$ and $Q(I)$ moments of the nucleus is taken into account by us in an approximation linear in $\hat{H}(hfs)$:

$$\varepsilon(nljIF)=\varepsilon(nlj)+\langle nljIF|\hat{H}(hfs)|nljIF\rangle; \quad (1)$$

here for $\varepsilon(nlj)$ we have adopted values shown in the tables of Ref. 9.

3. Mesons captured by an atom into the initial orbit with $n\approx 10-15$ are polarized. In the subsequent cascade of Auger-processes and radiative transitions of the meson leading to the population of the $3dj_1$-orbit this polarization is passed on sufficiently effectively, as can be inferred from the value of the polarization of mesons (~20%) in the $1s\frac{1}{2}$-orbit in the case of heavy ($Z \geq 48$) mesic atoms (cf., Ref. 10). Under these conditions it is natural to observe the effect of nonconservation of parity in the angular distribution of quanta emitted in the transition of μ-mesons from the $3dj_1$-orbit to the $1s\frac{1}{2}$-orbit. Below we investigate this variant of the mixing of the initial state of a mesic atom with a meson in the $3dj_1$-orbit and the state with a meson in the $3pj_3$-orbit, the observed quantity being the angular distribution of the quanta emitted in the transition of the mesic atom to one of the states $|1s\frac{1}{2}IF\rangle$. In contrast to the variant studied in Refs. 1–7 involving the mixing of the $2s\frac{1}{2}$ and $2p\frac{1}{2}$ states of a mesic atom, in the variant being considered by us of the mixing of the intersecting $|3dj_1IF\rangle$ and $|3pj_3IF\rangle$ states the differences between the energies of the levels of the mesic atom are comparable with the radiative halfwidths $\gamma(3pj_3)$ and $\gamma(3dj_1)$, and this is taken into account in the formula for the angular distribution of the quanta. With this aim in mind we have investigated the development in time of the wave function of the system "mesic atom + radiation field," the wave function for the system being determined at the instant $t=0$ by the amplitudes $C(3dj_1; F'f')$:

$$\Psi|_{t=0}=\sum_{F'f'}C(3dj_1;F'f')|3dj_1IF'f'\rangle|0\rangle, \quad (2)$$

where $|0\rangle$ is the function for the photon vacuum.

The amplitudes $C(3dj_1; F'f')$ fix the values of the spin-tensor describing the population of the states $|3dj_1IF'f'\rangle$

$$\rho_{00}(j_1F')=\sum_{f'}|C(3dj_1;F'f')|^2 \quad (3)$$

and of the spin-tensor describing the polarization of the mesic atoms

$$\rho_{10}(j_1F'F')=\sum_{f'}(F'1f'0|F'1F'f')C^*(3dj_1;F'f')C(3dj_1F'f'), \quad (4)$$

where $(ABab|ABC\gamma)$ is the Clebsch-Gordon coefficient in the notation of Condon and Shortley.[11] It is assumed that all the other orientation spin-tensors of higher ranks ($\varkappa \geq 2$) $\rho_{\varkappa 0}(j_1\tilde{F}'F')$ are equal to zero in the initial state of the mesic atom. For the quantization axis (the z axis) of the angular momentum we have chosen the direction of the polarization of the spin of the meson incident on the atom which corresponds to the direction of the meson momentum.

In the range of nuclei $Z>50$ the magnitude of the hyperfine splitting of the $3dj_1$ term of the mesic atom, as a rule, is significantly greater than the radiative half-width $\gamma(3dj_1)$. In this case the polarization spin-tensor $\rho_{10}(j_1\tilde{F}'F')$ for the mesic atom in the $3dj_1$-orbit is diagonal with respect to the quantity F', i.e.,

$$\rho_{10}(j_1\tilde{F}'F')=\delta_{\tilde{F}'F'}\rho_{10}(j_1F'F'). \quad (5)$$

Below we exhibit the formulas for the effects for this simplified case, i.e., we assume that the condition (5) is satisfied; naturally, situations can exist for which (5) is not satisfied.

4. In the ranges of Z under consideration manifestation of two different processes of nonconservation of parity of the states of the mesic atoms is possible:

A. The effect of a contact weak neutral interaction between a muon and a nucleus. Below we carry out an estimate for two assumed forms for the weak interaction $\hat{H}_1$ and $\hat{H}_2$ given in the notation of Ref. 2:

$$\hat{H}_1=C_1\frac{G}{\sqrt{2}}\bar{\psi}_\mu\gamma_i\gamma_5\psi_\mu\cdot\bar{\Psi}_N\gamma_i\Psi_N, \quad (6)$$

$$\hat{H}_2=C_2\frac{G}{\sqrt{2}}\bar{\psi}_\mu\gamma_i\psi_\mu\cdot\bar{\Psi}_N\gamma_i\gamma_5\Psi_N, \quad (6')$$

where $G=(1.4320\pm 0.0011)\times 10^{-49}$ erg·cm^3. A natural measure of such an interaction is the quantity $GR_0^{-3} \approx 5.2 \times 10^{-2}\ A^{-1}$ eV.

B. The polarization mechanism for the nonconservation of parity for states of mesic atoms. As is well known, parity is not conserved in nuclear states, and therefore the nuclear state of energy E_ν and spin I_ν (component M_ν) is described by a superposition of states of parities Π_ν and $-\Pi_\nu$:

$$|E_\nu I_\nu M_\nu\rangle = |\Pi_\nu I_\nu M_\nu\rangle + i\beta(I_\nu E_\nu)\,|-\Pi_\nu I_\nu M_\nu\rangle, \qquad (7)$$

the amplitude for the admixture $|\beta(I_\nu E_\nu)|$ being usually of the order of $10^{-7}-10^{-6}$, but in the range of nuclear excitation energy ~ 5–7 MeV dynamic amplification of the effect of nonconservation of parity of nuclear states is possible by a factor of the order of 10^1-10^2 (cf., Ref. 12, 13), so that the amplitude $|\beta(I_\nu E_\nu)|$ can attain values of $10^{-5}-10^{-4}$.

We consider in such a situation the electric interaction between a meson and a nucleus. Suppose that we have populated the $|n_1 l_1 j_1 IF'f\rangle$ state of a mesic atom with a meson in the $n_1 l_1 j_1$-orbit. In a subsequent nonradiative EL-transition of the meson to the $n_2 l_2 j_2$-orbit excitation of the nucleus occurs to the level $|E_\nu I_\nu\rangle$. Since parity is not conserved in nuclear states, the inverse nonradiative $E\tilde{L}$-transition of a nucleus to the ground state ($\tilde{L} = L+1$) is possible with the meson making a transition to the $n_3 l_3 j_3$-orbit, where $l_3 = l_1 + 1$. The matrix element for the $E\tilde{L}$-transition is proportional to the amplitude $\beta(I_\nu E_\nu)$. As a result of this process superposition of mesic atom states $|n_1 l_1 j_1 IF'f'\rangle$ and $|n_3 l_3 j_3 IF'f'\rangle$ occurs where $l_3 = l_1 + 1$. Schematically the process is represented by the diagram in Fig. 1. The superposition of the states of the mesic atom with the meson in the $n_1 l_1 j_1$- and $n_3 l_3 j_3$-orbits that arises manifests itself in the usual manner in radiative transitions to lower lying states $|n_4 l_4 j_4 IFf\rangle$. If the mesic atom was polarized in the initial state, then the emitted quanta will have an angular distribution of the form

$$W(\theta) = 1 + \alpha\cos\theta, \qquad (8)$$

where θ is the angle between the propagation vector of the quantum **k** and the polarization vector of the total angular momentum **F′** of the mesic atom in the initial state.

5. For a more compact presentation of the formulas we utilize below the abbreviated notation for the nlj meson orbits:

$$n_1l_1j_1 \equiv j_1 \equiv 3dj_1,\quad j_1 = 3/2 \text{ or } 5/2;$$
$$n_3l_3j_3 \equiv j_3 \equiv 3pj_3,\quad j_3 = 1/2 \text{ or } 3/2;$$
$$n_4l_4j_4 \equiv j_4 \equiv 1s1/2,\quad j_4 = 1/2.$$

In the case of the polarization mechanism the formula contains intermediate meson orbits $1s\frac{1}{2}$, $2s\frac{1}{2}$, $2p\frac{1}{2}$, $2p\frac{3}{2}$, for which we utilize the notation $n_2 l_2 j_2 \equiv j_2$.

§1. ANGULAR DISTRIBUTION OF QUANTA EMITTED BY A POLARIZED MESIC ATOM

1. In the general case which, in particular, includes resonance between a meson EL-transition and a nuclear EL-transition, i.e., which takes into account the possibility of formation of a mixed nuclear-meson atom state, the formula for the angular distribution of quanta is fairly awkward. This case was discussed by us in Ref.

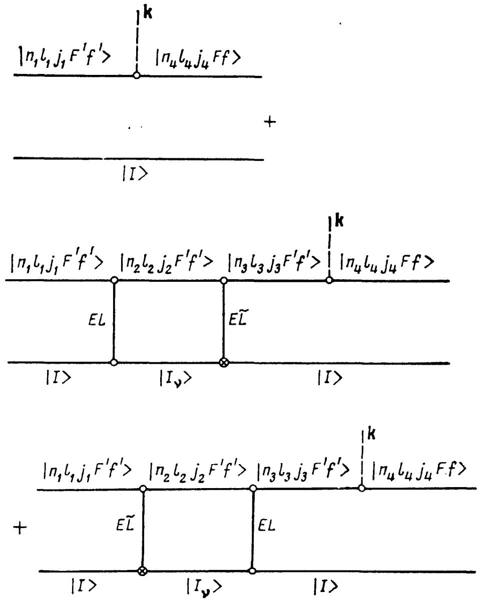

FIG. 1. Schematic diagram of the polarization mechanism for the nonconservation of parity of mesic atom states. The vertex proportional to the amplitude $\beta(I_\nu E_\nu)$ is denoted by a cross in a circle. Conservation of the total angular momentum F' of the system is arbitrarily noted in meson states.

14. In order to provide an idea of the order of magnitude of quantities we limit ourselves here to a real, but a simpler, variant, viz.: we assume below that the polarization Coulomb interaction between the meson and the nucleus can be included as a correction to the mesic atom energy terms and that moreover, conditions (5) are satisfied, i.e., the hyperfine splitting for the $3dj_1$-orbit is significantly greater than the radiative halfwidth $\gamma(3dj_1)$. Within the framework of the restrictions indicated above and in the linear approximation with respect to the weak interaction between a meson and a nucleus (correspondingly, in the linear approximation in terms of the amplitude $\beta(I_\nu E_\nu)$), we have in the case of intersecting terms $3dj_1 IF'$ and $3pj_3 IF'$ of a mesic atom, for the probability of the emission of a quantum ($\hbar\omega$; $d\hbar\omega$) into the solid angle $d\Omega$ at an angle θ to the quantization axis (the z axis) corresponding to the transition of a mesic atom into the group of states $|n_4 l_4 j_4 IFf\rangle$ with a meson in the $1s\frac{1}{2}$-orbit

$$dW(\hbar\omega;\theta) = d\hbar\omega\,\frac{d\Omega}{4\pi}\,[w_0(\hbar\omega) + w_1(\hbar\omega)\cos\theta], \qquad (9)$$

where $w_0(\hbar\omega)$ gives the spectral distribution of the emitted quanta:

$$w_0(\hbar\omega) = \frac{1}{2\pi}\sum_{FF'} \hbar W(E2; j_1 \to j_4)\,|\Delta(j_1F'j_4F\omega)|^{-2}[u(F'I2j_4; j_1F)]^2; \qquad (10)$$

$u(abcd; ef)$ is the normalized Racah function, tables for which are given in Ref. 15; $W(EL; j \to j_4)$ is the rate of the radiative EL-transition of a mesic atom $nlj \to n_4 l_4 j_4$;

$$\Delta(jF'j_4F\omega) = \varepsilon(nljIF') - \varepsilon(n_4l_4j_4IF) - \hbar\omega - i\gamma(nlj), \qquad (11)$$

where $\varepsilon(nlj)$ contains corrections for the polarization of the nucleus; in a specific calculation $\varepsilon(nlj)$ are taken from Ref. 9. The effect of nonconservation of parity manifests itself in the second term of formula (9);

$$w_1(\hbar\omega)=\frac{1}{2\pi}\sum_{FF'}\frac{\hbar W(E2;j_1\to j_4)}{|\Delta(j_1F'j_4F\omega)|^2}\rho_{10}(j_1F'F')\frac{3}{\sqrt{2}}\eta\left(\frac{E1;j_3\to j_4}{E2;j_1\to j_4}\right)\cdot$$

$$\times[u(1F'2F;F'1)u(F'I2j_4;j_1F)u(F'I1j_4;j_3F)]i\left\{\frac{\langle 3pj_3IF'|\hat{H}'|3dj_1IF'\rangle}{\Delta(j_3F'j_4F\omega)}-\text{c.c.}\right\}. \tag{12}$$

Here we have introduced the quantities

$$\eta\left(\frac{E1;j_3\to j_4}{E2;j_1\to j_4}\right)=\binom{+}{-}\left\{\frac{W(E1;3pj_3\to 1s^1/_2)}{W(E2;3dj_1\to 1s^1/_2)}\right\}^{1/2}; \tag{13}$$

the sign of η is determined by the radial matrix elements of the EL- and $E2$-multipoles of the mesic atom. In the variants discussed below of the "intersection" of the $3d\frac{5}{2}$–$3p\frac{3}{2}$ and the $3d\frac{5}{2}$-$3p\frac{1}{2}$ meson orbits the quantity η is negative.

By $\langle 3pj_3IF'|H'|3dj_1IF'\rangle$ we have denoted the matrix element of the effective operator for the weak interaction between the meson and the nucleus. In the case of a contact weak interaction the operator $\hat{H}'$ is generated by the densities $\hat{H}_1'$ or $\hat{H}_2'$ (cf., formuls (6) and (6′)), while in the polarization mechanism we have for the matrix element of the operator in accordance with the scheme for the process (cf., Fig. 1)

$$\langle 3pj_3IF'|\hat{H}'|3dj_1IF'\rangle=-i\sum_{L\tilde{L}}\sum_{E_\nu I_\nu}\sum_{(n_2l_2j_2)}\frac{\beta(I_\nu E_\nu)}{\Delta_{2\nu}(F'F\omega)}\cdot$$

$$\times\langle 3pj_3IF'|\hat{H}_c(EL)|n_2l_2j_2I_\nu F'\rangle\langle n_2l_2j_2I_\nu F'|\hat{H}_c(E\tilde{L})|3dj_1IF'\rangle. \tag{14}$$

The calculation of the matrix elements of the Coulomb EL-interaction between a meson and a nucleus is possible only if a concrete nuclear model is used. But an estimate of the value of the elements (apparently with an error of $\lesssim 20\%$) can be given on the basis of experimental data if we utilize an approximate formula for the operator $\hat{H}_c(EL)$:

$$\hat{H}_c(EL)\approx -e^2\frac{4\pi}{(2L+1)}\sum_m Y_{Lm}(\mathbf{r}_\mu)\left(\frac{1}{r}\right)^{L+1}\int d^3r r^L Y_{Lm}^*(\mathbf{r})[\hat{\rho}(\mathbf{r})-\langle\rho(\mathbf{r})\rangle], \tag{15}$$

where $\mathbf{r}_\mu$ is the meson coordinate, $\hat{\rho}(\mathbf{r})$ is the operator for the density of the protons of the nucleus, $\langle\rho(r)\rangle$ is the average distribution of the protons over the volume of the nucleus. In formula (15) the region $r_\mu < r$ is neglected, so that the greatest error occurs in the estimation of the matrix elements for the $1s\frac{1}{2}$ meson orbit.

The energy denominator of the polarization element $\Delta_{2\nu}(F'F\omega)$ has the form

$$\Delta_{2\nu}(F'F\omega)=(E_\nu-E)+\tilde{\varepsilon}(n_2l_2j_2I_\nu F')-\varepsilon(n_1l_1j_1IF)-\hbar\omega-i\gamma_{2\nu}, \tag{16}$$

where E is the energy of the ground state of the nucleus. The energy of the intermediate meson state $\tilde{\varepsilon}(n_2l_2j_2I_\nu F')$ does not include corrections for the Coulomb polarization of the nucleus. The quantity $\gamma_{2\nu}$ is given by

$$\gamma_{2\nu}=\gamma(n_2l_2j_2)+\gamma_\nu, \tag{17}$$

where $\gamma(nlj)$ are the radiative halfwidths of the nlj-meson orbits; γ_ν is the halfwidth of the E_ν-level of the nucleus. Since the decay time for a meson in the $1s\frac{1}{2}$ orbit is significantly greater than the times for the radiative transitions of a meson ($\tau(1s\frac{1}{2})\sim 10^{-7}$ sec for $Z\geq 50$), we have assumed for the $1s\frac{1}{2}$ orbit $\gamma(n_4\rho_4j_4)=0$.

2. Here we do not consider the resonance situation with the formation of a mixed mesic atom-nuclear state. In the absence of such resonance the energy of the quanta $\hbar\omega$ is determined by the denominator $|\Delta(j_1F'j_4F\omega)|^2$ in a band of width $\sim\gamma(3dj_1)$ near the difference between terms of the mesic atom

$$\hbar\omega\approx\varepsilon(3dj_1IF')-\varepsilon(1s^1/_2IF). \tag{18}$$

The maximum effect of nonconservation of parity manifests itself in the case when the difference between the "intersecting" terms $|3dj_1IF'\rangle$ and $|3pj_3IF'\rangle$ is equal to the halfwidth $\gamma(3pj_3)$, as can be seen from formula (12). Since for real mesic atoms it is practially impossible to calculate the position of the terms with an accuracy of the order of $\gamma(3pj_3)$ due to the indefiniteness of the estimate of the effect of the Coulomb polarization of the nucleus, therefore below we estimate the possible magnitude of the effect under the condition that the following equality is satisfied

$$|\varepsilon(3dj_1IF')-\varepsilon(3pj_3IF')|\approx\gamma(3pj_3). \tag{19}$$

The "intersection" of the terms for the $3d\frac{5}{2}$- and 3 $3p\frac{3}{2}$-orbits of a meson in the nuclear range $Z=56-60$ is demonstrated in Tables I and II, where values are given of the characteristics of nuclei and of mesic atom states needed for subsequent numerical estimates. As was noted already, the values of $\varepsilon(3d\frac{5}{2})$ and $\varepsilon(3p\frac{3}{2})$ were taken from the tables of Ref. 9. For comparison it is natural to take as the reference level the magnitude of $\varepsilon(3d\frac{5}{2})$ and to introduce the deviations

$$a(3d^5/_2IF')=\langle 3d^5/_2IF'|\hat{H}(hfs)|3d^5/_2IF'\rangle,$$
$$b(3p^3/_2IF')=\varepsilon(3p^3/_2)-\varepsilon(3d^5/_2)+\langle 3p^3/_2IF'|\hat{H}(hfs)|3p^3/_2IF'\rangle, \tag{20}$$

and also the difference between terms with angular momentum F'

$$\Delta_0(F')=b(3p^3/_2IF')-a(3d^5/_2IF'). \tag{21}$$

The hyperfine interaction $\hat{H}(hfs)$ of the meson with the magnetic dipole $\mu(I)$ and the electric quadrupole $Q(I)$ moments of the nucleus in the ground state was taken into account by us in the linear approximation. For the calculation we used the relativistic meson wave func-

TABLE I. Characteristics of mesic atoms and nuclei in the region $Z=56-60$.

Characteristic	Nucleus			
	Ba_{56}^{135}	La_{57}^{139}	Pr_{59}^{141}	Nd_{60}^{145}
$I\Pi$ (ground state)	$^3/_2+$	$^7/_2+$	$^5/_2+$	$^7/_2-$
$\mu(I)$, μ_B	0,8365	2,778	4.3	−0,654
$Q(I)$, b	0.18±0.02	0.22±0.03	−0.07	−0.25
$\varepsilon(3p^3/_2)-\varepsilon(3d^5/_2)$, keV*	−0,72	−0.42	+0.30	+1,0
$\gamma(3d^5/_2)$ **	43.5	46	52	55.5
$\gamma(3p^3/_2)$	63.5	66	73	76.5
$\gamma(2p^3/_2)$	240	255	280	290
$\gamma(2p^1/_2)$	223	230	255	265
$\gamma(2s^1/_2)$	3.35	3.70	4.75	5.45
$\eta\left(\frac{E1;\ 3p^3/_2\to 1s^1/_2}{E2;\ 3d^5/_2\to 1s^1/_2}\right)$	−6.7	−6.5	−6.15	−5.95

*According to the data from the tables of Ref. 9.
**The values of γ are given in electron volts.

TABLE II. Hyperfine splitting of the $3d^{5/2}$- and $3p^{3/2}$-orbits of a mesic atom in the region of nuclei $Z=56-60$.

F'	$a(j_1IF')$, keV	$b(j_3IF')$, keV	$\Delta_0(F')$, keV	Transition $F'\to F$	$[u_2]^2$ for $F'\to F$	$[u_1u_2u_3]$ for $F'\to F$
			Nucleus Ba_{56}^{135}			
0		+1,018				
1	+0.316	−0,352	−0,668	$1\to 1$	$^9/_{10}$	$-\sqrt{3/16}$
2	−0.033	−1,762	−1,729	$2\to 2$	$^3/_{10}$	$-\sqrt{21/400}$
3	−0.234	−0.403	−0,169	$3\to 2$	$^3/_5$	$-\sqrt{3/25}$
4	+0.097					
			Nucleus La_{57}^{139}			
1	+0.340					
2	+0.176	+0.649	+0.473	$2\to 3$	$^9/_{10}$	$-3/5$
3	−0.012	−0.670	−0.658	$3\to 3$	$^3/_4$	$-\sqrt{27/128}$
4	−0.153	−1.219	−1.066	$4\to 3$	$^{11}/_{20}$	$-\sqrt{33/640}$
5	−0,156	−0.068	+0.088	$5\to 4$	$^7/_{10}$	$-\sqrt{21/125}$
6	+0.098					
			Nucleus Pr_{59}^{141}			
0	−0.040					
1	−0.006	+0.073	+0.079	$1\to 2$	$^{14}/_{15}$	$-\sqrt{42/100}$
2	+0.047	+0.471	+0.424	$2\to 3$	$^1/_5$	$-2/15$
3	+0.084	+0.629	+0,545	$3\to 3$	$^2/_5$	$-\sqrt{1/12}$
4	+0.052	+0.023	−0.029	$4\to 3$	$^2/_3$	$-\sqrt{3/20}$
5	−0.112	−	−	−	−	−
			Nucleus Nd_{60}^{145}			
1	−0.397	−	−	−	−	−
2	−0.189	−0.227	−0.038	$2\to 3$	$^9/_{10}$	$-3/5$
3	+0.043	+1.342	+1.299	$3\to 4$	$^1/_4$	$\sqrt{3/128}$
4	+0.210	+1.936	+1.726	$4\to 4$	$^9/_{20}$	$\sqrt{1617/16000}$
5	+0.191	+0.446	+0.255	$5\to 4$	$^7/_{10}$	$-\sqrt{21/125}$
6	−0.167	−	−	−	−	−

tions for the Coulomb field of a uniformly charged sphere of radius $R_0=1.24\times 10^{-13}A^{1/3}$ cm.

In Table II are given these values of $a(3d_{\frac{5}{2}}IF')$, $b(3p_{\frac{3}{2}}IF')$, $\Delta_0(F')$, and also the values of the products of the Racah coefficients appearing in formula (10): $|u_2|^2=[u(F'I2j_4;\,j_1F)]^2$, and in formula (12):

$$[u_1u_2u_3]=[u(1F'2F;F'1)\,u(F'I2{}^1/_2;{}^5/_2F)\,u(F'I1{}^1/_2;{}^3/_2F)].$$

§2. THE EFFECT OF A WEAK NEUTRAL INTERACTION BETWEEN A MUON AND A NUCLEUS

We estimate for the case of mixing of $3dj_1$- and $3pj_3$-states of a meson the effect of a weak interaction between the meson and the nucleus represented by the densities $\hat{H}_1$ and $\hat{H}_2$ (cf., formulas (6) and (6′)). For the sake of simplicity we shall specify the meson functions $\psi_{nlj\mu}$ by two subscripts: the angular momentum j and the component of the angular momentum μ. Utilizing in explicit form the relativistic meson functions $\psi_{nlj\mu}$

$$\psi_{nlj\mu}(\mathbf{r})=\begin{cases} ig_{nlj}(r)\,\Omega_{lj\mu}(\mathbf{n}) \\ f_{nl'j}(r)\,(\sigma\mathbf{n})\,\Omega_{l'j\mu}(\mathbf{n}) \end{cases}, \tag{22}$$

where $\mathbf{n}=\mathbf{r}/r$, $\Omega_{lj\mu}$ is a spherical spinor (cf., Ref. 16), we obtain an expansion for the meson currents; thus, for the operator $\hat{H}_1$ we have

$$\psi^{+}_{j_3\mu_3}\gamma_5\psi_{j_1\mu_1}=i[g_{j_3}(r)f_{j_1}(r)-g_{j_1}(r)f_{j_3}(r)]\frac{1}{\sqrt{4\pi}}\cdot$$

$$\times\sum_L\{j_3l_3[A(L)]j_1l_1\}(j_3L\mu_3m|j_3Lj_1\mu_1)\,Y_{Lm}(\mathbf{r}). \tag{23}$$

In the case of the operatore $\hat{H}_2$ the expansion formula is more awkward and contains both combinations of the radial meson functions ($\xi=r/R_0$):

$$(1/R_0)^3\varphi_\pm(\xi)=[g_{j_3}(r)f_{j_1}(r)\pm g_{j_1}(r)f_{j_3}(r)]. \tag{24}$$

As calculations have shown, the coefficients $\{j_3l_3[A(L)]\times j_1l_1\}$ (and analogous ones for the operator $\hat{H}_2$) are all comparable to unity, and, therefore, a real measure of the weak neutral interaction between a meson and a nucleus is the quantity

$$\frac{1}{4\pi}GR_0^{-3}\varphi_\pm(1)\approx\frac{5{,}2}{4\pi}\cdot 10^{-2}A^{-1}\varphi_\pm(1)\ \ [\mathrm{eV}]. \tag{25}$$

The functions $\varphi_\pm(\xi)$ were calculated using relativistic radial functions for meson orbits in the Coulomb field of a uniformly charged sphere of radius $R_0=1.24\times 10^{-13}A^{1/3}$ cm.

In the region of nuclei $Z=56-60$, where the meson terms $3p_{\frac{3}{2}}$ and $3d_{\frac{5}{2}}$ "intersect" the functions $\varphi_\pm(\xi)$ within the nuclear volume ($\xi\leq 1$) are well approximated by the relation

$$\varphi_\pm(\xi)\approx\varphi_\pm(1)\xi^3. \tag{26}$$

For the group of nuclei under consideration the values of $\varphi_\pm(1)$ for the intersecting $3p_{\frac{3}{2}}$ and $3d_{\frac{5}{2}}$ meson orbits are given below:

Nucleus:	Ba_{56}^{135}	La_{57}^{139}	Pr_{59}^{141}	Nd_{60}^{145}
$10^3\varphi_+(1)$:	−4.06	−4.71	−5.86	−6.74
$10^3\varphi_-(1)$:	−0.85	−0.99	−1.24	−1.43

while for the $3p_{\frac{1}{2}}$ and $3d_{\frac{5}{2}}$ orbits of the Pb_{82}^{208} nucleus they are as follows:

$$\varphi_+(1)=2.14\cdot 10^{-3},\quad \varphi_-(1)=-6.95\cdot 10^{-3}.$$

Estimating for the mesic atom states $|3d_{\frac{5}{2}}IF'\rangle$ the initial value of the polarization spin-tensor ρ_{10} to have the value of the order of $0.2\rho_{00}$ and taking into account the amplification factor η in formula (12) we obtain for the expected value of the coefficient α in the expression for the angular distribution of quanta $W(\theta)=1+\alpha\cos\theta$ in the optimal variant of the "intersection" of the $3dj_1$ and $3pj_3$ terms

$$\alpha\lesssim(4\pi)^{-1}GR_0^{-3}\varphi_\pm(1)[\gamma(3pj_3)]^{-1}. \tag{27}$$

Thus, in the regions $Z=56-60$ and $Z=82-85$ the effect of neutral currents for the case of mixing of the $3pj_3$ and $3dj_1$ states leads to values of the coefficient $\alpha\lesssim 10^{-9}$. Under these conditions there is no sense in carrying out a more detailed estimate of α utilizing some specific nuclear models. As will be seen below, the polarization mechanism for the nonconservation of parity leads to significantly greater values of the coefficient α.

§3. THE POLARIZATION MECHANISM FOR THE NONCONSERVATION OF PARITY OF STATES OF MESIC ATOMS

1. For the evaluation of the polarization mechanism we use the approximate formula (15) for the Coulomb interaction between a meson and a nucleus, and this enables us to express the matrix element of the effective operator $\hat{H}'$ (formula (14)) in terms of observable quantities:

$$\langle 3pj_3IF'|\hat{H}'|3dj_1IF'\rangle\approx -i\left(\frac{e^2}{a_0}\right)^2\sum_{L\tilde{L}}\sum_{(n_2l_2j_2)}g(3pj_3(EL)n_2l_2j_2(E\tilde{L})3dj_1)\times$$

$$\times\sum_{E_\nu I_\nu}\frac{\exp\{i\pi n_\nu(L\tilde{L})\}\beta(I_\nu E_\nu)}{\Delta_{2\nu}(F'F\omega)}|\langle I\|EL\|I_\nu\rangle\langle I\|E\tilde{L}\|I_\nu\rangle^*|$$
$$\times u(IF'Lj_2;j_3I_\nu)\,u(IF\tilde{L}j_2;j_1I_\nu) \tag{28}$$

Here we have separated out the mesic atom factor being calculated

$$g(3pj_3(EL)n_2l_2j_2(E\tilde{L})3dj_1)=(-)^{j_1-j_3}\left[\frac{(2l_2+1)(2l_2+1)}{3\cdot5}\right]^{1/2}$$
$$\times(Ll_200|Ll_210)(\tilde{L}l_200|\tilde{L}l_220)\,u(j_3{}^1/_2Ll_2;1j_2)\,u(j_1{}^1/_2\tilde{L}l_2;2j_2) \tag{29}$$
$$\times\left(\frac{R_0}{a_0}\right)^{L+\tilde{L}}\left\langle 3pj_3\left|\left(\frac{1}{x}\right)^{L+1}\right|n_2l_2j_2\right\rangle\left\langle 3dj_1\left|\left(\frac{1}{x}\right)^{\tilde{L}+1}\right|n_2l_2j_2\right\rangle.$$

The radial mesic-atom elements of EL-multipoles required for the calculation

$$\left\langle nlj\left|\left(\frac{1}{x}\right)^{L+1}\right|n_2l_2j_2\right\rangle=\int_0^\infty(g_lg_{l_2}+f_lf_{l_2})\left(\frac{1}{x}\right)^{L+1}x^2\,dx \tag{30}$$

are given in Table III. Here in (28)–(30) we use mesic-atom units: $a_0=\hbar^2/e^2m_\mu$, $e^2/a_0\approx5.63$ keV; x is the meson coordinate; $x=r_\mu/a_0$.

The nuclear matrix elements of EL-multipoles exhibited there which are defined by the equation

$$\left\langle IM\left|\hat{\rho}(\mathbf{r})\left(\frac{r}{R_0}\right)^L\left(\frac{4\pi}{(2L+1)}\right)^{1/2}Y_{Lm}{}^*(\mathbf{r})\right|I_\nu M_\nu\right\rangle$$
$$=(ILMm|ILI_\nu M_\nu)\langle I\|EL\|I_\nu\rangle, \tag{31}$$

can be expressed either in terms of the reduced probabilities $B(EL;I\to I_\nu)$ measured in the Coulomb excitation of a nucleus (cf., Ref. 17):

$$|\langle I\|EL\|I_\nu|^2=\frac{4\pi}{(2L+1)}\left(\frac{2I+1}{2I_\nu+1}\right)\frac{B(EL;I\to I_\nu)}{e^2R_0{}^{2L}}, \tag{32}$$

or in terms of the partial radiative halfwidths $\gamma(EL;I_\nu\to I)$ for transitions to the ground state of the nucleus:

$$|\langle I\|EL\|I_\nu\rangle|^2=\gamma(EL;I_\nu\to I)\left[\frac{e_2}{\hbar c}\frac{(E_\nu-E)^{2L+1}}{[(2L+1)!!]^2}\left(\frac{R_0}{\hbar c}\right)^{2L}\frac{(2L+1)(L+1)}{L}\right]^- \tag{33}$$

2. Further we start from the natural assumption that the sign of the amplitude $\beta(I_\nu E_\nu)$ and the sign of the product of the reduced matrix elements of EL- and $E\tilde{L}$-multipoles, represented in (28) by the factor $\exp[i\pi n_\nu(L\tilde{L})]$ (n_ν is an integer), vary randomly from level to level, although, of course, a situation can not be excluded when several close nuclear levels give an in-phase contribution to (28). Within the framework of the assumption made above the magnitude of the element (28) is determined in practice by one nuclear level, for which in the case of a fixed frequency of the quantum $\hbar\omega$ the denominator $\Delta_{2\nu}(F'F\omega)$ is minimal.

Two regions of nuclear excitation energy make a contribution to the effective matrix element (28):

1) The region E_ν, corresponding to the energies of meson transitions $3d\to2p$ and $3d\to2s$, iee., for $Z=56-60$ the region $E_\nu=1-1.5$ MeV.

2) The region E_ν, corresponding to the meson transition $3d\to1s$, i.e., for $Z=56-60$ the region $E_\nu=5.2-5.7$ MeV.

As an example, for the region $E_\nu=1-1.5$ MeV we consider the chain of meson transitions

TABLE III. Coulomb radial elements of EL-multipoles of a mesic atom for the region $Z=56-60$.

Element	Z: 56	57	58	59	60
$\langle 3p^3/_2\,|\,x^{-2}\,|\,1s^1/_2\rangle$	$2.62\cdot10^2$	$2.70\cdot10^2$	$2.77\cdot10^2$	$2.85\cdot10^2$	$2.90\cdot10^2$
$\langle 3p^3/_2\,|\,x^{-2}\,|\,2s^1/_2\rangle$	$7.53\cdot10$	$7.84\cdot10$	$8.15\cdot10$	$8.46\cdot10$	$8.79\cdot10$
$\langle 3d^5/_2\,|\,x^{-2}\,|\,2p^3/_2\rangle$	$7.33\cdot10$	$7.60\cdot10$	$7.87\cdot10$	$8.15\cdot10$	$8.42\cdot10$
$\langle 3d^5/_2\,|\,x^{-3}\,|\,1s^1/_2\rangle$	$1.72\cdot10^3$	$1.82\cdot10^3$	$1.91\cdot10^3$	$2.01\cdot10^3$	$2.10\cdot10^3$
$\langle 3d^5/_2\,|\,x^{-3}\,|\,2s^1/_2\rangle$	$-7.43\cdot10$	$-7.30\cdot10$	$-7.05\cdot10$	$-6.89\cdot10^2$	$-6.07\cdot10$
$\langle 3p^3/_2\,|\,x^{-3}\,|\,2p^3/_2\rangle$	$3.47\cdot10^3$	$3.64\cdot10^3$	$3.81\cdot10^3$	$4.00\cdot10^3$	$4.14\cdot10^3$

$$3d^5/_2[E1]2p^3/_2[E2]3p^3/_2.$$

For the nuclei $Z=56-60$ with $E_\nu=1-1.5$ MeV the $E2$ transitions are weakly collectivized, while the $E1$-transitions are generally delayed by a factor of the order of 10^{-4}. Therefore, in making the estimate we take for the values of the nuclear elements

$$|\langle I\|E1\|I_\nu\rangle|=10^{-2},\qquad|\langle I\|E2\|I_\nu\rangle|=1 \tag{34}$$

and for the denominator we take the value $\Delta_{2\nu}(F'F\omega)=50$ keV. As a result, utilizing the data of Table III, in the optimum variant of the intersection of the $3d\frac{5}{2}$ and $3p\frac{3}{2}$ terms, i.e., for $\Delta(j_3F'j_4F\omega)=0.1$ keV we obtain the factor determining the magnitude of the coefficient of the angular distribution of the quanta in the transition of the mesic atom $3d\frac{5}{2}\to1s\frac{1}{2}$:

$$\left|\mathrm{Im}\frac{\langle 3p^3/_2IF'|\hat{H}'|3d^5/_2IF'\rangle}{\Delta(j_3F'j_4F\omega)}\right|\leqslant0.2|\beta(I_\nu E_\nu)|. \tag{35}$$

For estimating the contribution of the region $E_\nu\approx\hbar\omega\times(3d\to1s)$ for the nuclei $Z=56-60$ we utilize the data obtained from investigating the reaction of radiative capture of thermal neutrons.[18–20] We estimate the values of the reduced nuclear $E1$- and $E2$-elements using the radiative widths of the resonances ($\Gamma_\gamma=0.15-0.05$ eV) and the observed intensity $I_\gamma(E1)$ of the hard lines for the transitions to the ground state of the nucleus or to states close to the ground state. In accordance with the data of Ref. 18–20 for nuclei $Z=55-60$, the values of $I_\gamma(E1)$ vary within the limits from 1 to 50%. For making the estimate we take $I_\gamma(E1)=10\%$, the half-widths $\gamma(E1;I_\nu\to I)=10^{-2}$ eV and $\gamma(E2;I_\nu\to I)=10^{-3}$ eV, while for $\Delta_{2\nu}$ we take half the distance $D(I_\nu)$ between the levels $I_\nu E_\nu$ at an energy $E_\nu=5.2-5.7$ MeV. The magnitude of $D(I_\nu)$ for the corresponding nuclear spins I_ν is estimated from the distance between neutron s-resonances according to the well-known formula for the density of nuclear levels (cf., for example Ref. 21, 22). Thus, taking $\Delta_{2\nu}=0.3$ keV in the optimum variant of the "intersection" of the $3d\frac{5}{2}$ and $3p\frac{3}{2}$ meson terms ($\Delta_0(F')=0.1$ keV), we obtain

$$\left|\mathrm{Im}\frac{\langle 3p^3/_2IF'|\hat{H}'|3d^5/_2IF'\rangle}{\Delta(j_3F'j_4F\omega)}\right|\leqslant0.4|\beta(I_\nu E_\nu)|. \tag{36}$$

As has been noted above, for the region of the spectrum of the nucleus $E_\nu=5-6$ MeV a dynamic amplification is possible of the effect of nonconservation of parity in nuclear states by a factor of $\sim10^1-10^2$, and, therefore, the contribution of this region must be greater by an order of magnitude than the contribution of the region of $E_\nu = 1-1.5$ MeV. Naturally, for specific nuclei the values of (35) and (36) can be obtained more precisely, and, in particular, the special case of mesic atom-nuclear resonance which was not considered here is of interest.

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Vibrational autoionization of a molecule and recombination of asymmetrical molecular ions

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The properties of highly excited electronic states of asymmetrical diatomic molecules are investigated. The rate of vibrational autoionization, a process in which the highly excited electron is ionized by transfer of vibrational energy from the ions, is determined. The inverse process, the attachment of an electron to a molecular ion followed by excitation of vibrations of the latter, is considered. Actual calculations are performed for the ion HeH^+. An expression is obtained for the recombination flux due to the vibrational attachment and the attachment due to triple collision of the ion with two electrons (triple recombination). It is shown that the investigated processes should play a substantial role in the kinetics of low-temperature plasma.

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INTRODUCTION

We investigate here the properties of highly excited ($n>5$) bound electronic states of asymmetrical diatomic molecules. The dimensions of the orbit of the excited electron are much larger than the distance between the nuclei, so that molecule can be regarded as a combination of two weakly-interacting systems: the molecular-ion core and the excited electron. As a result of the weak but finite interaction, these systems can exchange energy. In particular, the energy of the vibrational–rotational motion of the ion can be transferred to the excited electron. If the magnitude of the vibrational–rotational quantum is larger than the binding energy of the excited electron, then such an exchange will cause detachment of the electron. This effect can be naturally called vibrational autoionization. We have found no mention of this effect in the literature,[1–4] although an analysis shows that it should play a substantial role in the kinetics of a low-temperature plasma.