[3]V. A. Maisheev, A. M. Frolov, E. A. Arakelyan, G. L. Bayatyan, G. S. Vartanyan, N. K. Grigoryan, A. T. Margaryan, and S. S. Stepanyan, Preprint IFVÉ 76-15, Serpukhov, 1976.
[4]V. A. Maisheev, A. M. Frolov, R. O. Avakyan, E. A. Arakelyan, A. A. Armaganyan, G. L. Bayatyan, G. S. Vartanyan, G. A. Vartapetyan, N. K. Grigoryan, A. O. Kechechyan, S. G. Knyazyan, A. T. Margaryan, É. M. Matevosyan, R. M. Mirzoyan, S. S. Stepanyan, L. Ya. Kolesnikov, A. L. Rubashkin, and P. V. Sorokin, Zh. Eksp. Teor. Fiz. **77**, 1708 (1979) [Sov. Phys. JETP **50**, 856 (1979)]; Preprint IFVÉ 79-63, Serpukkov, 1979.
[5]G. L. Bayatyan, G. S. Vartanyan, O. M. Vinnitskiĭ, N. K. Grigoryan, S. G. Knyazyan, V. A. Maisheev, A. T. Margaryan, V. P. Sakharov, Yu. M. Sapunov, and A. M. Frolov, Preprint EFI-64 (74), Erevan, 1974.
[6]M. A. Kumakhov, Phys. Lett. A **57**, 17 (1976).
[7]M. A. Kumakhov, Dokl. Akad. Nauk SSSR **230**, 1077 (1976) [Sov. Phys. Dokl. **21**, 581 (1976)].
[8]M. A. Kumakhov, Zh. Eksp. Teor. Fiz. **72**, 1489 (1977) [Sov. Phys. JETP **45**, 781 (1977)].
[9]G. Diambrini Palazzi, Rev. Mod. Phys. **40**, 611 (1968).
[10]V. N. Baĭer, V. M. Katkov, and V. S. Fadin, Izluchenie relyativistskikh élektronov (Emission of Radiation by Relativistic Electrons), Atomizdat, M., 1973, p. 234.
[11]B. B. Rossi, High-Energy Particles, Prentice-Hall, New York, 1952 (Russ. Transl., Gostekhizdat, 1955, pp. 66–68).
[12]A. I. Akhiezer, V. F. Boldyshev, and N. F. Shul'ga, Fiz. Elem. Chastits At. Yadra **10**, 51 (1979) [Sov. J. Part. Nucl. **10**, 19 (1979)].
[13]M. T. Robinson, Phys. Rev. **179**, 327 (1969).
[14]S. A. Vorob'ev, B. N. Kalinin, V. V. Kaplin, and A. P. Potylitsyn, Pis'ma Zh. Tekh. Fiz. **4**, 1340 (1978) [Sov. Tech. Phys. Lett. **4**, 539 (1978)].
[15]A. O. Agan'yants, Yu. A. Vartanov, G. A. Vartapetyan, M. A. Kumakhov, Kh. Trikalinos, and V. Ya. Yaralov, Pis'ma Zh. Eksp. Teor. Fiz. **29**, 554 (1979) [JETP Lett. **29**, 505 (1979)].
[16]I. I. Miroshnichenko, D. D. Merri, R. O. Avakyan, and T. Kh. Figut, Pis'ma Zh. Eksp. Teor. Fiz. **29**, 786 (1979) [JETP Lett. **29**, 722 (1979)].

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P-odd nuclear forces—a source of parity violation in atoms

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We consider the electromagnetic *P*-odd interaction produced between an electron and a nucleus by parity-violating nuclear forces. New information on these forces can be obtained even now by experimentally investigating the optical activity of heavy atoms and diatomic molecules. In the case of the deuteron, the *P*-odd vector potential is expressed in terms of parameters that characterize parity violation in *np* scattering.

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1. The discovery of the weak electron–nucleon interaction due to neutral currents, which was made in Novosibirsk[1] by observing the optical activity of atomic-bismuth vapor, is undoubtedly only the first positive result in the study of the structure of weak interactions by atomic-spectroscopy methods. The subsequent increase of the accuracy of this experiment[2] makes it quite realistic to obtain qualitatively new results in this field. We have in mind, in particular, the measurement, in atoms[3] and in molecules,[4] of the constant that characterizes the weak interaction of the electron with the nuclear spin. The natural scale of this effect is $\sim 1/Z$ of the already measured one, since it receives contributions not from all the nucleons of the nucleus, but only from the one valence nucleon. It is realistic to expect to measure an effect of this order of magnitude even at present. However, within the framework of the Weinberg-Salam model, at the mixing-parameter value $\sin^2\theta = 0.23$ which follows from the available experimental data, the constant of the interaction in question is numerically very small.

We wish to note in the present article that parity-violation effects that depend on the spin of the nucleus can be produced in atoms not by neutral currents only. They are induced also by *P*-odd nuclear forces. These forces produce in the electromagnetic vector-potential of the nucleus an increment of incorrect *P*-parity, and this increment acts on the electron. The contribution of this mechanism can be noticeably larger than that of the neutral currents. Thus, experiments aimed at observing the effects in question are already realistic at present. They can yield valuable information on parity violation in a nucleus.

Effects peculiar to μ-mesic atoms and due to *P*-odd nuclear forces were considered earlier by Grechukhin and Soldatov.[5] Parity violation in the interaction of an electron with a nucleus, induced by *P*-odd nuclear forces, was discussed from a general point of view in a recent paper by Henley *et al.*[6]

2. The cause of the considered phenomenon is easier to understand by starting from the fact that violation of spatial parity in the nucleus produces in it a spiral spin structure[7,8] and toroidal currents.[7] It is the resultant contribution to the vector-potential of the nucleus which leads to the *P*-odd electron–nucleus interaction of interest to us. We shall show that this electromagnetic interaction must of necessity be of the contact type. This result is contained in essence in a number of papers. According to a private communication from P. Sandars, it was known already to N. Ramsey. None-

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theless, to make the exposition comprehensive, we deem it advisable to present here a proof of this fact.

It is convenient to expand the matrix element J_μ of the electromagnetic-current operator between states with given spin I and momenta $\mathbf{k}$ and $\mathbf{k}'$ in terms of the four independent vectors of the problem:

$$p_\mu=(k'+k)_\mu,\quad q_\mu=(k'-k)_\mu,\quad s_\mu,\quad r_\mu=i\varepsilon_{\mu\nu\varkappa\lambda}p_\nu q_\varkappa s_\lambda, \tag{1}$$

where s_μ is the four-dimensional spin operator which is specified, e.g., for the state with momentum $\mathbf{k}$ and satisfies therefore the condition $k_\mu s_\mu=0$:

$$s_0=\mathbf{Ik}/m,\quad \mathbf{s}=\mathbf{I}+(\mathbf{kI})\mathbf{k}/m(E+m).$$

In this case, of course, $k'_\mu s_\mu \neq 0$.

Taking into account the electromagnetic-current conservation law $q_\mu J_\mu=0$, this matrix element takes the following form[10]:

$$J_\mu=\langle I,\ k'|p_\mu F_1+r_\mu F_2+[q^2 s_\mu-q_\mu(qs)]F_3|I,\ k\rangle. \tag{2}$$

The invariant functions F_i depend on scalars that can be constructed from the vectors (1). Since $pq=0$ and $sk=0$, there are only two such invariants. For the arguments of F_i it is convenient to choose the Hermitian quantity $\tau=iq_\mu s_\mu$ and the usual variable $t=q^2$. Substitution of the operator function $F_i(t,\tau)$ in expression (2), to satisfy the hermiticity condition, must be accompanied by symmetrization of the noncommuting operators. In the expansion of $F_i(t,\tau)$ in powers of τ

$$F_i(t,\tau)=\sum_{n=0}^{N_i} f_{in}(t)\tau^n \tag{3}$$

the maximum power of N_i is obviously determined by the spin I:

$$N_1=2I,\quad N_2=N_3=2I-1.$$

Thus, the total number of electromagnetic form factors f_{in} is $6I+1$.

The operator $\tau=is_\mu q_\mu$ reverse sign not only under space reflection, but also under time reversal. Therefore the T-invariance requirement restricts the summation in (3) to even powers of τ. But then all the functions $F_i(t,\tau)$ turn out to be P-even, and the entire parity nonconservation in the electromagnetic current (2) is due only to the structure of $q^2 s_\mu - q_\mu(qs)$. The term proportional to q_μ corresponds to the gradient of the scalar function in the vector-potential of the particle. This contribution is eliminated by a gauge transformation. The factor q^2 in the remaining term of the current cancels out the propagator $1/q^2$ of the particle field. Thus, a P-odd electromagnetic interaction is indeed of necessity a contact interaction. We note that the term with $n=0$ in the expansion (3) for the function $F_3(t,\tau)$ corresponds to the well known anapole moment first considered by Zel' dovich.[7]

As for the T-odd part of the electromagnetic current, the contact terms here are not only those that conserve the P-invariance. As a result, the contact part of the current is the one with incorrect C-parity.

The total number of invariant form factors f_{in} that conserve (+) and violate (−) the P, T, and C invariances are obtained for a given spin I without difficulty and is listed in the table. If two quantities are contained in one line of the table, the left-hand and right-hand quantities pertain to integer and half-integer spin, respectively. This classification of electromagnetic form factors was first presented by Kobzarev *et al.*[10] using a transition into the annihilation channel.

TABLE I.

P	T	C	Number of form factors
+	+	+	$2I+1$
−	+	−	I | $I+1/2$
+	−	−	I | $I-1/2$
−	−	+	$2I$

In concluding this section we note that in those current matrix elements that are not diagonal in the internal state the T-invariance requirement no longer forbids non-contact P-odd terms. This is attested to by the circular polarization of the radiation in the presence of parity nonconservation. The physical cause of this difference is quite clear. The already mentioned toroidal currents[7] produce a field only inside a system that has no definite P parity. Therefore the diagonal matrix element of the interaction is of the contact type. And only when the system is restructured, i.e., in the case of an off-diagonal matrix element, does this field emerge to the outside.

3. We proceed now to a direct calculation of the parity-violating electromagnetic field of the nucleus. It is natural to regard the nucleus as nonrelativistic, so that the P-odd part of $J_\mu(q)$ takes the form

$$j_0(\mathbf{q})=0,\quad \mathbf{j}(\mathbf{q})=(\mathbf{q}^2\mathbf{I}-(\mathbf{qI})\mathbf{q})F. \tag{4}$$

The first of these equations has a simple physical meaning: the P-odd interaction does not influence the charge distribution in the system. From the second equation it follows, in particular, that

$$\int \mathbf{j}(\mathbf{r})\,d^3r=\mathbf{j}(\mathbf{q}=0)=0. \tag{5}$$

Since the characteristic atomic momenta $\mathbf{q}$ are small compared with the nuclear ones, the dependence of F on q^0 and on $\mathbf{q}\cdot\mathbf{I}$ can be neglected. In other words, we assume the nucleus to be pointlike. In this case the vector potential can be represented in the form

$$\mathbf{A}(\mathbf{r})=\int\frac{d^3r'\,\mathbf{j}(\mathbf{r}')}{|\mathbf{r}-\mathbf{r}'|}=[(\mathbf{b}\nabla)\nabla-\mathbf{b}\Delta]\frac{1}{r}. \tag{6}$$

To find the constant vector $\mathbf{b}$, it suffices to integrate (6) with respect to $\mathbf{r}$, taking (5) into account. As a result we arrive at the following expression for the vector-potential:

$$\mathbf{A}(\mathbf{r})=-\pi\delta(\mathbf{r})\int d^3r'(r')^2\,\mathbf{j}(\mathbf{r}'). \tag{7}$$

We have left out here the term $\nabla(\mathbf{b}\cdot\nabla)r^{-1}$, which is eliminated by a gauge transformation.

Thus, neglecting the dependence of F on $\mathbf{q}$, we consider only the lower of the P-odd multipoles—the anapole moment of the nucleus

$$\mathbf{a}=-\pi\int d^3r\,r^2\mathbf{j}(\mathbf{r}). \tag{8}$$

Its calculation is generally speaking a very complicated task. We confine ourselves here to the simplest shellmodel approximation: one nucleon in excess of the completely closed shells. The anapole moment is the result of the mixing of single-particle levels of different parity with one and the same total angular momentum. We express the resultant state in the form

$$|0\rangle+\sum_k \frac{|k\rangle\langle k|H_w|0\rangle}{E_0-E_k}$$
$$=R_0(r)\Omega_{Ilm}(\mathbf{n})+i\sum_k \eta_k R_k(r)\Omega_{Il'm}(\mathbf{n}), \tag{9}$$

where R_0 and R_k are the radial wave functions, Ω_{Jlm} are spherical spinors, $l'=2I-l$, and $\mathbf{n}=\mathbf{r}/r$. The mixing coefficient

$$i\eta_k=\langle k|H_w|0\rangle/(E_0-E_k) \tag{10}$$

is under the customary definition of spherical function a pure imaginary quantity by virtue of the T-invariance of the weak-interaction operator H_w.

We begin with calculation of the contribution made to the effect by that current term which is connected with the spin magnetic moment of the nucleon. Using the identity

$$\Omega_{Il'm}=-(\boldsymbol{\sigma}\mathbf{n})\Omega_{Ilm},$$

we represent the spin-current axial part of interest to us

$$\mathbf{j}^s(\mathbf{r})=\frac{e\mu}{2m}\operatorname{rot}(\psi^+\boldsymbol{\sigma}\psi) \tag{11}$$

in the form

$$\frac{e\mu}{m}\operatorname{rot}\left\{\sum_k \eta_k R_0 R_k \Omega^+_{Ilm_1}[\mathbf{n},\boldsymbol{\sigma}]\Omega_{Ilm_1}\right\}. \tag{12}$$

Here e and m are the charge and mass of the proton, and μ is the magnetic moment of the nucleon. Substituting (12) in (8) and integrating by parts, we obtain without difficulty

$$\mathbf{a}^s=2\pi\frac{\varkappa\mathbf{I}}{I(I+1)}\frac{e\mu}{m}\sum_k \eta_k r_{0k},$$
$$r_{0k}=\int_0^\infty dr\, r^3 R_0 R_k,\quad \varkappa=(I+{}^1/_2)(-1)^{I+{}^1/_2-l}. \tag{13}$$

We proceed now to calculate the contribution connected with the orbital motion of the proton (there is of course no such contribution for the neutron). The integral $\int d^3r r r^2 \mathbf{j}(r)$ can be conveniently represented with the aid of the substitution $\dot{\mathbf{r}}=i[H,\mathbf{r}]$ ($[a,b]$ is a commutator) in the form

$${}^1/_2 ie\langle r^2[H,\mathbf{r}]+[H,\mathbf{r}]r^2\rangle, \tag{14}$$

where the mean value is taken over the perturbed state

$$|0\rangle+i\sum_k \eta_k|k\rangle. \tag{14a}$$

We break up the total nonrelativistic Hamiltonian of the valent nucleon into three terms

$$H=H_r+H_l+H_w. \tag{15}$$

Here

$$H_r=-\frac{1}{2m}\frac{1}{r}\partial_r{}^2 r+U(r),\quad H_l=\frac{\hat{l}^2}{2mr^2},$$

and H_w is the weak-interaction Hamiltonian, which takes the form

$$H_w=\frac{G}{2^{1/2}m}[\boldsymbol{\sigma}\mathbf{p}f(r)+f(r)\boldsymbol{\sigma}\mathbf{p}]. \tag{16}$$

The matrix element

$${}^1/_2 ie\langle r^2[H_r,r]\mathbf{n}+[H_r,r]r^2\mathbf{n}+2r^3[H_l,\mathbf{n}]+2r^2[H_w,\mathbf{r}]\rangle$$

obtained by substituting (15) in (14) is transformed with the aid of the identity

$$r^2[H_r,r]+[H_r,r]r^2={}^2/_3[H_r,r^3]$$

in the following manner:

$$ie\langle{}^1/_3[H-H_l-H_w,r^2\mathbf{r}]+[H_l,r^2\mathbf{r}]+r^2[H_w,\mathbf{r}]\rangle$$
$$=\frac{2}{3}ie\left\langle\frac{1}{2m}[\hat{l}^2,\mathbf{r}]+r^2[H_w,\mathbf{r}]-\mathbf{r}(\mathbf{r}[H_w,\mathbf{r}])\right\rangle. \tag{17}$$

The mean value of the commutator $[\hat{l}^2,\mathbf{r}]$ over the state (14a) is easy to calculate. It is equal to

$$-i\frac{l(l+1)-l'(l'+1)}{I(I+1)}\mathbf{I}\sum_k \eta_k r_{0k}=-2i\frac{\varkappa\mathbf{I}}{I(I+1)}\sum_k \eta_k r_{0k}. \tag{18}$$

The terms in (17) that contain the operator of the so-called contact current

$$\mathbf{j}_c=ie[H_w,\mathbf{r}]=\frac{G}{2^{1/2}}\frac{e}{m}\boldsymbol{\sigma}f(r), \tag{19}$$

reduce to the form

$$-\frac{2}{3}\frac{G}{2^{1/2}}\frac{e}{m}\frac{\varkappa\mathbf{I}}{I(I+1)}\langle 0|r^2 f(r)|0\rangle. \tag{20}$$

We ultimately obtain the following orbital contribution to the anapole:

$$\mathbf{a}^{orb}=-\frac{2\pi}{3}\frac{e}{m}\frac{\varkappa\mathbf{I}}{I(I+1)}\left\{\sum_k \eta_k r_{0k}-\frac{G}{2^{1/2}}\langle 0|r^2 f(r)|0\rangle\right\}. \tag{21}$$

Direct calculation, not based beforehand on the contact character of the interaction, is incomparably more laborious but leads, naturally, to the same results for the vector potentials A^s and A^{orb}.

The resultant P-odd electron–nucleus interaction Hamiltonian is

$$H^{em}=e\boldsymbol{\alpha}(\mathbf{A}^s+\mathbf{A}^{orb})$$
$$=\delta(\mathbf{r})\frac{2\pi\alpha}{m}\frac{\varkappa\mathbf{I}\boldsymbol{\alpha}}{I(I+1)}\left\{\left(\mu-\frac{q}{3}\right)\sum_k \eta_k r_{0k}+\frac{q}{3}\frac{G}{2^{1/2}}\langle 0|r^2 f(r)|0\rangle\right\}, \tag{22}$$

where $\boldsymbol{\alpha}$ is the Dirac matrix for the electron and q is the charge of the external nucleon in units of e.

We present for comparison the Hamiltonian of the weak interaction due to the nucleon-axial and electron-vector neutral currents,

$$H^{nc}=-\delta(\mathbf{r})\frac{G}{2^{1/2}}K_2\frac{(\varkappa-{}^1/_2)\mathbf{I}\boldsymbol{\alpha}}{I(I+1)}. \tag{23}$$

Here K_2 is a dimensionless constant that characterizes this interaction. In the Weinberg-Salam model

$$K_{2p}=-K_{2n}=-{}^1/_2(1-4\sin^2\theta)\cdot 1.25\approx -0.05. \tag{24}$$

The Hamiltonian (22) can be rewritten in similar form

$$H^{em}=\delta(\mathbf{r})\frac{G}{2^{1/2}}K\frac{\varkappa\mathbf{I}\boldsymbol{\alpha}}{I(I+1)}. \tag{25}$$

To estimate the dimensionless constant K, we retain in it only the term containing μ, recognizing that $\mu_p \gg \frac{1}{3}$. Then

$$K\approx 2^{1/2}\pi\alpha\mu m\cdot 10^5\sum_k \eta_k r_{0k}. \tag{26}$$

Putting $r_{0k} \sim A^{1/3}/m_\pi$ and $\eta_k \sim 10^{-7}-10^{-6}$ we get

$$K \sim 0.1-1. \quad (27)$$

This quantity can thus turn out to be much larger than the constant K_2 in the Weinberg-Salam model.

We do not take into account here the contribution made to the anapole moment of the nucleus to the intrinsic anapole moment of the nucleon, due to the radiative corrections for weak interactions. This contribution to the constant K amounts to approximately 10^{-2}.

Atomic calculations with the electromagnetic Hamiltonian (25) do not differ in any way with the earlier[3] calculations with the Hamiltonian (23).[1)] For the transition in bismuth with wavelength $\lambda = 648$ nm the difference in the degree of circular polarization of the radiation for different hyperfine components of the line reaches 4% at $K=1$, a value sufficiently close to the accuracy 14% already attained at the present time.[2] A unique possibility of measuring anapole moments of different nuclei can be afforded in principle by the stay of the optical activity of a gas of diatomic molecules, an activity due precisely to the interaction (22) in (23).

Thus, atomic and molecular experiments are a possible source of new information on the parity violation in nuclei.

4. We would like to dwell separately on the calculation of the anapole moment of the deuteron, inasmuch as in the approximation of a zero radius of action of nuclear forces it can be expressed in terms of parameters that characterize parity violation in np scattering.

In this approximation, the wave function of the deuteron, with allowance for the P-odd effects, was obtained in the brilliant paper of Danilov[15]:

$$\psi(\mathbf{r}) = \left(\frac{\varkappa}{2\pi}\right)^{1/2} \{1 + [c(\sigma_p+\sigma_n) + \lambda_t(\sigma_p-\sigma_n)](-i\nabla)\} \frac{e^{-\varkappa r}}{r}\chi. \quad (28)$$

Here $\mathbf{r} = \mathbf{r}_p - \mathbf{r}_n$, $\varkappa = (m\varepsilon_d)^{1/2}$, ε_d is the deuteron binding energy, and χ is the triplet spin function of the two nucleons. The constant c determines the P-odd part of the np-scattering amplitude, which changes the isotopic spin T of the system, but conserves the usual spin S. The constant λ_t corresponds, on the contrary, to conservation of T and to a change of S. The densities of the proton and neutron currents are calculated in accordance with the usual formulas with the aid of the wave function (28). Substituting the expressions obtained in this manner in (8), we arrive to the following results for the anapole moment of the deuteron:

$$\mathbf{a} = \frac{\pi e}{m}\mathbf{I}\left\{c\left(\mu_p - \mu_n - \frac{1}{3}\right) + \lambda_t(\mu_p + \mu_n)\right\}. \quad (29)$$

The contribution of the contact current to the anapole, as well as the contribution of the structural electromagnetic vertex, can be neglected in the zero-radius approximation. In fact, since the weak-interaction radius certainly does not exceed the radius r_0 of the nuclear forces, all these terms should be proportional to the small factor $|(r_0)|^2 r_0^3 \sim \varkappa r_0$. On the other hand, the absence of the indicated quantity from (29) is due to the fact that the corresponding vector potential is described in terms of perturbation theory by pole diagrams with small energy denominators. Thus, the result (29) is essentially a low-energy theorem of sorts. Unfortunately, the real accuracy of the employed approximation is low. It is known that for the deuteron the parameter $\varkappa r_0 \sim \frac{1}{3}$.

From the analysis in Sec. 2, it follows that the deuteron has only one P-odd form factor even in the general case (i.e., without assuming it to be pointlike). We take the opportunity to note that the two P-odd structures cited by Henley *et al.*[6] for the electromagnetic vertex of the deuteron (see also Ref. 16), can be reduced to each other by means of the identity

$$\xi_\mu^* \xi_\nu - \xi_\nu^* \xi_\mu = -{}^1/{}_2 e_{\mu\nu\alpha\beta} e_{\alpha\beta\varkappa\lambda} \xi_\varkappa^* \xi_\lambda$$

We write down the P-odd electromagnetic interaction of the electron with the deuteron in the form

$$H^{em} = \delta(r) \cdot 2^{-1/2} GK\mathbf{I}\boldsymbol{\alpha} \quad (30)$$

(of course, the nonrelativistic approximation for the operator $\boldsymbol{\alpha}$ is perfectly sufficient in the case of the deuteron). Even at the optimistic estimate of the np-scattering parameters $c \sim \lambda \sim 10^{-6} m^{-1}$, the dimensionless constant is here $K \sim 10^{-2}$. Nonetheless, the contribution of the anapole to the P-odd effects that depend on the spin of the nucleus may turn out to be substantial in the deuterium atom. The point is that in the Weinberg-Salam model the axial current of the–adrons is an isovector and is therefore absent from the isoscalar deuteron. The discussed effect in deuterium are therefore due, apart from the calculated anapole moment of the deuteron, only to the isoscalar part of the radiative corrections.

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[1)]We note only that an accurate allowance for the finite dimensions of the nucleus lead in the atomic problem to corrections that do not exceed $Z^2\alpha^2/2$.

[1]L. M. Barkov and M. S. Zolotorev, Pis'ma Zh. Eksp. Teor. Fiz. 27, 379 (1978) [JETP Lett. 27, 357 (1978)].
[2]L. M. Barkov and M. S. Zolotorev, Phys. Lett., 85B, 308 (1979).
[3]V. N. Novikov, O. P. Sushkov, V. V. Flambaum, and I. B. Khriplovich, Zh. Eksp. Teor. Fiz. 73, 802 (1977) [Sov. Phys. JETP 46, 420 (1977)].
[4]O. P. Sushkov and V. V. Flambaum, Zh. Eksp. Teor. Fiz. 75, 1208 (1978) [Sov. Phys. JETP 48, 608 (1978)].
[5]D. P. Grechukhin and A. A. Soldatov, Zh. Eksp. Teor. Fiz. 73, 31 (1977) [Sov. Phys. JETP 46, 15 (1977)].
[6]E. M. Henley, W.-Y. P. Hwang, and G. N. Epstein, Phys. Lett. 88B, 349 (1979).
[7]Ya. B. Zel'dovich, Zh. Eksp. Teor. Fiz. 33, 1531 (1957) [Sov. Phys. JETP 6, 1184 (1957)].
[8]L. M. Barkov, I. B. Khriplovich, and M. S. Zolotorev, Comments Atom. Mol. Phys. 8, 79 (1979).
[9]L. Durand III, P. C. De Celles, and R. B. Marr, Phys. Rev. 126, 1883 (1962).
[10]A. A. Cheshkov and Yu. M. Shirokov, Zh. Eksp. Teor. Fiz. 44, 1982 (1963) [Sov. Phys. JETP 17, 1333 (1963)].
[11]A. A. Cheshkov, Zh. Eksp. Teor. Fiz. 50, 144 (1966) [Sov. Phys. JETP 23, 97 (1966)].
[12]I. Yu. Kobzarev, L. B. Okun', and M. V. Terent'ev, Pis'-

ma Zh. Eksp. Teor. Fiz. 2, 466 (1965) [JETP Lett. 2, 289 (1965)].
[13]A. D. Dolgov, Pis'ma Zh. Eksp. Teor. Fiz. 2, 494 (1965) [JETP Lett. 2, 308 (1965)].
[14]V. M. Dubovik and A. A. Cheshkov, Zh. Eksp. Teor. Fiz. 51, 1369 (1966) [Sov. Phys. JETP 24, 924 (1966)].
[15]G. S. Danilov, Phys. Lett. 18, 40 (1965).
[16]W.-Y. P. Hwang and E. M. Henley, Preprint RLO-1388-804, Univ. Washington (1979).

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