Problem of parity nonconservation in ^{57}Fe and ^{119}Sn nuclei

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(Received 19 March 1993)

We consider parity nonconservation effect for the Mossbauer transitions in ^{57}Fe and ^{119}Sn nuclei. To simplify the calculations we use contact approximation for nucleon-nucleon weak interaction. A corresponding parity nonconserving effective Hamiltonian is presented. The result of the calculation of γ-quantum circular polarization degree $P \sim 3\times10^{-8}$ strongly contradicts the experimental value $P \sim 10^{-3}$. One may consider this huge disagreement as an indication of the unknown nuclear phenomena which leads to the enhancement of parity nonconservation. We stress the importance of a new independent measurement of the effect.

PACS number(s): 21.10.Hw, 21.30.+y, 33.45.+x

I. INTRODUCTION

Parity nonconservation (PNC) for low frequency Mossbauer transitions in ^{119}Sn and ^{57}Fe have been observed in Refs. [1–4]. The measured values of the circular polarization of γ quanta are $P = (0.90 \pm 0.1) \times 10^{-3}$ for the 23.9 keV transition in ^{119}Sn and $P = (0.58 \pm 0.12) \times 10^{-3}$ for the 14.4 keV transition in ^{57}Fe. These are transitions between very low-energy levels in the spherical nuclei. Therefore neither the high density of the spectrum [5] nor the accidental closeness of the opposite parity states can enhance the effect. Due to these reasons it is widely believed that the experimental values of P are too large to be explained in the standard picture of the weak interaction in nuclei. It should be mentioned that there is an old theoretical paper [6] which gives the values of P consistent with experiment. However, we do not agree with the results of this work. Actually in Ref. [6] the value of the regular $M1$ amplitude is underestimated almost by 2 orders of magnitude which is evident from a comparison with the experimental value of the level lifetime. On the other hand, in our opinion the weak interaction is substantially overestimated in the above work. Unfortunately only the result of the numerical calculation of the PNC amplitude is presented in [6]. Therefore we cannot point out the reason for this overestimation. Thus in the present paper we perform an accurate calculation of the PNC effect in ^{119}Sn and ^{57}Fe.

II. EFFECTIVE HAMILTONIAN OF THE WEAK INTERACTION

The standard parametrization of the parity nonconserving nucleon-nucleon weak interaction corresponds to the one-boson-exchange approximation with π-, ρ-, and ω-meson exchange taken into account [7]. The ω-meson contribution is small and we neglect it. There are terms of different structure in the contribution corresponding to the ρ-meson exchange. However, the term proportional to the constant h_ρ^0 (see Ref. [7]) is dominating and therefore we neglect other terms. Thus we will consider only the π-meson contribution and h_ρ^0 term from ρ meson contribution. In this approximation the Hamiltonian of the nucleon-nucleon PNC interaction is of the form [7]

$$\begin{aligned} H^{(1b)} &= i\frac{f_\pi g_\pi}{4\sqrt{2}m_p}(\boldsymbol{\tau}_1 \times \boldsymbol{\tau}_2)_z(\boldsymbol{\sigma}_1 + \boldsymbol{\sigma}_2)[\mathbf{p}_{12}, F_\pi] \\ &\quad - \frac{g_\rho h_\rho^0}{2m_p}(\boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2)\big((\boldsymbol{\sigma}_1 - \boldsymbol{\sigma}_2)\{\mathbf{p}_{12}, F_\rho\} + i(1+\mu)(\boldsymbol{\sigma}_1 \times \boldsymbol{\sigma}_2)[\mathbf{p}_{12}, F_\rho]\big), \\ F_a &= \frac{e^{-m_a r}}{4\pi r}, \qquad r = |\mathbf{r}_1 - \mathbf{r}_2|, \qquad \mathbf{p}_{12} = \mathbf{p}_1 - \mathbf{p}_2. \end{aligned} \tag{1}$$

Here $\boldsymbol{\sigma}$ and $\boldsymbol{\tau}$ are spin and isospin Pauli matrices; $\mathbf{p}$ is the momentum operator, m_p is the proton mass; m_π, m_ρ are the π- and ρ-meson masses; [,] is a commutator; and { , } is an anticommutator. $\mu = 3.7$ is an isovector anomalous magnetic moment of the nucleon. The strong nucleon-meson constants are $g_\pi = 13.45$, $g_\rho = 2.79$. The estimates for the weak nucleon-meson constants f_π and h_ρ^0 are presented in Refs. [7,8]. So-called "best" values are $f_\pi = 12 \times (3.8 \times 10^{-8})$, $h_\rho^0 = -30 \times (3.8 \times 10^{-8})$ [7,8]. We would like to stress that "best" does not mean that we really know the values. Information on these constants is still incomplete.

The one-boson-exchange Hamiltonian (1) is nonlocal (nonzero range) and therefore it is not convenient for simple nuclear calculations. More convenient is the local (zero range) effective interaction between the nucleons a and b. The zero range (contact) approximation is rather crude, but it allows one to do an analytical calculation. This approximation is very convenient if we do not need high accuracy and therefore we use it. To first order in the momenta the contact interaction is of the form

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$$W_{ab} = \frac{G}{\sqrt{2}}\frac{1}{2m_p}((g_{ab}\boldsymbol{\sigma}_a - g_{ba}\boldsymbol{\sigma}_b)\{\mathbf{p}_{ab}, \delta(\mathbf{r}_{ab})\} + ig'_{ab}(\boldsymbol{\sigma}_a \times \boldsymbol{\sigma}_b)[\mathbf{p}_{ab}, \delta(\mathbf{r}_{ab})]). \tag{2}$$

Here G is the Fermi constant. Any exchange matrix element of the local Hamiltonian (2) can be reduced to the direct one by the Fiertz transformation. We define an effective interaction (2) in such a way that only the direct matrix elements have to be considered. The exchange terms are taken into account by appropriate definition of the dimensionless effective constants g. In the limit $F_a(\mathbf{r}) \to \frac{1}{m_a^2}\delta(\mathbf{r})$ the Hamiltonian (1) can be easily transformed to the form (2)

$$\begin{aligned} H_{nn} = H_{pp} &= -\frac{g_\rho h_\rho^0}{2m_p m_\rho^2}(\mu+2)((\boldsymbol{\sigma}_a - \boldsymbol{\sigma}_b)\{\mathbf{p}_{ab}, \delta(\mathbf{r}_{ab})\} + i(\boldsymbol{\sigma}_a \times \boldsymbol{\sigma}_b)[\mathbf{p}_{ab}, \delta(\mathbf{r}_{ab})]), \\ H_{pn} &= \frac{f_\pi g_\pi}{2\sqrt{2}m_p m_\pi^2}(\boldsymbol{\sigma}_p + \boldsymbol{\sigma}_n)\{\mathbf{p}_{pn}, \delta(\mathbf{r}_{pn})\} \\ &\quad -\frac{g_\rho h_\rho^0}{2m_p m_\rho^2}((2\mu+1)(\boldsymbol{\sigma}_p - \boldsymbol{\sigma}_n)\{\mathbf{p}_{pn}, \delta(\mathbf{r}_{pn})\} - i(\mu-1)(\boldsymbol{\sigma}_p \times \boldsymbol{\sigma}_n)[\mathbf{p}_{pn}, \delta(\mathbf{r}_{pn})]). \end{aligned} \tag{3}$$

Let us stress once more that only the direct matrix elements of (3) have to be considered. This is the standard way of defining the effective interaction in the Fermi liquid theory [9]. Let us discuss now the validity of the zero range approximation. The typical momentum of a nucleon in the nucleus is $p \sim m_\pi$. Therefore the contact limit is justified for the ρ exchange. However, we should take into account short range nucleon-nucleon repulsion. To do it let us introduce into the ρ term in Eq. (3) the factor W_ρ. We also include into this factor the corrections due to the range and exchange character of the ρ exchange force.

The zero range approximation for the π contribution is not so good. In Ref. [10] the π exchange was treated explicitly and the contact limit was considered only for the short range ρ and ω parts of the interaction. Such an approach is suitable for more accurate calculations. In the first approximation, our opinion, the simple contact Hamiltonian (2) is good enough. To take into account the long range and exchange character of the π exchange as well as short range nucleon-nucleon repulsion we introduce into the π term the factor W_π. Following Ref. [11], and based on the calculation of PNC in neutron and proton scattering on ^{4}He [12,13], we set $W_\rho \approx 0.4$, $W_\pi \approx 0.16$. These values agree with that obtained in other works [14,15]. Comparing (3) with (2) we find the dimensionless effective constants g:

$$\begin{aligned} g_{pp} &= -(\mu+2)W_\rho A_\rho h_\rho^0 \approx 1.5, \\ g'_{pp} &= g_{nn} = g'_{nn} = g_{pp}, \\ g_{pn} &= -(2\mu+1)W_\rho A_\rho h_\rho^0 + W_\pi A_\pi f_\pi \approx 6.5, \\ g_{np} &= -(2\mu+1)W_\rho A_\rho h_\rho^0 - W_\pi A_\pi f_\pi \approx -2.2, \\ g'_{pn} &= g'_{np} = (\mu-1)W_\rho A_\rho h_\rho^0 \approx -0.7, \\ A_\rho &= \frac{\sqrt{2}g_\rho}{Gm_\rho^2} \approx 0.57 \times 10^6, \quad A_\pi = \frac{g_\pi}{Gm_\pi^2} \approx 59 \times 10^6. \end{aligned} \tag{4}$$

Numerical values of the constants g_{ab} are presented for the "best" values of f_π and h_ρ^0 [7] (see above). We stress once more that the values of f_π and h_ρ^0 are not known precisely. Therefore one should be rather skeptical of the presented numerical values of g_{ab}.

An effective Hamiltonian of the weak interaction for the unpaired nucleon above the paired nuclear core can be written in the following form:

$$W_a = g_a \frac{G}{\sqrt{2}}\frac{1}{2m_p}\{\boldsymbol{\sigma}\mathbf{p}, \rho(r)\}. \tag{5}$$

Here $\boldsymbol{\sigma}$ and $\mathbf{p}$ are the spin and momentum of an unpaired nucleon. $\rho(r)$ is the core density normalized by the condition $\int \rho d^3r = A$ (we assume that the mass number $A \gg 1$). Due to Eqs. (2) and (4) the dimensionless constant g_a are equal

$$\begin{aligned} g_p &= \frac{Z}{A}g_{pp} + \frac{N}{A}g_{pn} = -W_\rho A_\rho\left(\frac{Z}{A}(\mu+2) + \frac{N}{A}(2\mu+1)\right)h_\rho^0 + W_\pi A_\pi \frac{N}{A}f_\pi \approx 4.5, \\ g_n &= \frac{Z}{A}g_{np} + \frac{N}{A}g_{nn} = -W_\rho A_\rho\left(\frac{Z}{A}(2\mu+1) + \frac{N}{A}(\mu+2)\right)h_\rho^0 - W_\pi A_\pi \frac{Z}{A}f_\pi \approx 0. \end{aligned} \tag{6}$$

Numerical values of g_a correspond to the "best" values of f_π and h_ρ^0, and $\frac{Z}{A} \approx 0.42$. Formulas (6) were obtained in Ref. [11]. An effective single-particle PNC interaction, (5) and (6), was also derived in Ref. [16]. As far as we understand the only important difference is in the factor W_ρ. We include short range nucleon-nucleon correlations into the effective interaction explicitly. It means that our factor $W_\rho \approx 0.4$ accounts for the range and exchange character of the ρ exchange force, as well as short range nucleon-nucleon correlations. The factor W_ρ=0.79–0.88 of Ref. [16] only accounts for the range and exchange character of the ρ exchange force. It does not include any effect due to short range correlations.

III. REGULAR $M1$ AMPLITUDES

We consider the 23.9 keV $|\frac{1}{2}^+\rangle \to |\frac{3}{2}^+\rangle$ transition in ^{119}Sn and the 14.4 keV $|\frac{1}{2}^-\rangle \to |\frac{3}{2}^-\rangle$ transition in ^{57}Fe. The half-life times for $|\frac{3}{2}\rangle$ levels are [17] ^{119}Sn, $t_{1/2} = 17.8$ ns; ^{57}Fe, $t_{1/2} = 98$ ns. Both transitions are predominantly of the $M1$ type [17]. Therefore one can easily calculate the $M1$ amplitudes

$$\begin{aligned} {}^{119}\mathrm{Sn}: &\qquad \langle \tfrac{3}{2}^+, \tfrac{1}{2}|M_z|\tfrac{1}{2}^+, \tfrac{1}{2}\rangle \approx \pm 0.67\mu_N, \\ {}^{57}\mathrm{Fe}: &\qquad \langle \tfrac{3}{2}^-, \tfrac{1}{2}|M_z|\tfrac{1}{2}^-, \tfrac{1}{2}\rangle \approx \pm 0.61\mu_N. \end{aligned} \tag{7}$$

Here $\mathbf{M}$ is the magnetic moment operator and $\mu_N = \frac{e}{2m_p}$ is the nuclear magneton. We have to note that the shell model transition in ^{57}Fe is $p_{1/2} \to p_{3/2}$, and in ^{119}Sn it is $s_{1/2} \to d_{3/2}$. Thus in Sn we have an l-forbidden transition. Nevertheless its magnitude is large. We do not discuss this problem in the present work. Nevertheless we would like to stress that this important problem is not actually resolved.

IV. PNC $E1$ AMPLITUDE

Let us consider the 23.9 keV transition in ^{119}Sn. In the naive shell model the ground state $|0\rangle$ corresponds to the $3s_{1/2}$ unpaired neutron above the closed shells: configuration $|\cdots 2d^4_{3/2}3s_{1/2}\rangle$. The first excited state $|1\rangle$ ($E = 23.9$ keV) corresponds to the $2d_{3/2}$ neutron hole in the closed shells: configuration $|\cdots 2d^3_{3/2}3s^2_{1/2}\rangle$. Thus in the independent particle model we have to calculate the PNC $E1$ amplitude $\langle 3s_{1/2}|E1|2d_{3/2}\rangle$. It should be noted that the many-body effects of BCS-type pairing are very important for the PNC $E1$ amplitude. Nevertheless let us consider first the independent particle model.

We will use the oscillator shell model, a very convenient way for calculating the PNC as suggested in Ref. [18]. The same trick has been used in Ref. [19] for calculation of the permanent electric dipole moment, Schiff moment, and anapole moment of a nucleus. The operator of the electric dipole moment

$$\mathbf{d} = \sum_c e_c \mathbf{r}_c \tag{8}$$

with effective charges $e_p = e\frac{N}{A}$, $e_n = -e\frac{Z}{A}$ [20] in the oscillator model can be presented in the form

$$\mathbf{d} = \frac{i}{m_p\omega^2}\left[\sum_c e_c\mathbf{p}_c, H\right], \tag{9}$$

where H is the nuclear oscillator Hamiltonian, and $\omega \approx 40/A^{1/3}$ MeV [21] is the oscillator frequency. The PNC $E1$ amplitude is equal to

$$\mathbf{D} = \sum_n \left(\frac{\langle 1|\mathbf{d}|n\rangle\langle n|W|0\rangle}{E_0 - E_n} + \frac{\langle 1|W|n\rangle\langle n|\mathbf{d}|0\rangle}{E_1 - E_n} \right). \tag{10}$$

Using Eqs. (9) and (2) and the relation $E_1 - E_0 \ll E_n - E_0 \sim \omega$ one can transform $\mathbf{D}$ into the form

$$\begin{aligned} \mathbf{D} &= -\frac{i}{m_p\omega^2}\sum_c \langle 1|[e_c\mathbf{p}_c, W]|0\rangle \\ &= -i\frac{G}{2\sqrt{2}m_p^2\omega^2}\sum_{b,c} e_c g_{ab}\langle 1|\{\boldsymbol{\sigma}_a(\mathbf{p}_a - \mathbf{p}_b), [\mathbf{p}_c, \delta(\mathbf{r}_a - \mathbf{r}_b)]\}|0\rangle. \end{aligned} \tag{11}$$

The index a corresponds to the external unpaired nucleon, and b corresponds to the core nucleon. We recall that the exchange terms should not be calculated in formula (11). They are taken into account by appropriate definition of the constants g_{ab} in Eq. (2). There are two possibilities in Eq. (11): (1) $c = a$, (2) $c = b$. The first corresponds to the virtual excitations of an external nucleon in the weak potential of core (5). Let us denote the corresponding contribution to the PNC $E1$ amplitude as $\mathbf{D}_{\rm ext}$. The second possibility corresponds to the virtual excitations of the core ($\mathbf{D}_{\rm core}$). From Eq. (11) we get

$$\begin{aligned} D^z &= D^z_{\rm ext} + D^z_{\rm core}, \\ D^z_{\rm ext} &= e_n g_n \xi = -\frac{Z}{A}\left(\frac{Z}{A}g_{np} + \frac{N}{A}g_{nn}\right)e\xi, \\ D^z_{\rm core} &= -\left(e_n\frac{N}{A}g_{nn} + e_p\frac{Z}{A}g_{np}\right)\xi = -\frac{ZN}{A^2}(g_{np} - g_{nn})e\xi, \\ \xi &= i\frac{G}{2\sqrt{2}m_p^2\omega^2}\langle 2d_{3/2}, \tfrac{1}{2}|\{(\boldsymbol{\sigma}\cdot\mathbf{p}), [p_z, \rho]\}|3s_{1/2}, \tfrac{1}{2}\rangle, \end{aligned} \tag{12}$$

where $\rho(r)$ is the nuclear density. Let us use the constant density approximation

$$\rho(r) = \rho_0\theta(R - r), \qquad \rho_0 = \frac{3}{4\pi r_0^3}, \qquad R = r_0 A^{1/3}, \qquad r_0 = 1.1 \text{ fm}. \tag{13}$$

Then

$$\xi = i\frac{G\rho_0}{6m_p\omega}\left(\frac{d\chi_{2d}}{dx}\chi_{3s} - \chi_{2d}\frac{d\chi_{3s}}{dx} + \frac{3}{x}\chi_{2d}\chi_{3s}\right)_{x=R\sqrt{m_p\omega}} = i\frac{0.16}{6}\frac{G\rho_0}{m_p\omega}. \tag{14}$$

The radial wave functions are normalized by the condition $\int \chi^2 dx = 1$. We have used the values of wave functions at the nuclear surface calculated for ^{119}Sn in the oscillator model: $\chi_{3s} = 0.60$, $\frac{d\chi_{3s}}{dx} = 1.11$, $\chi_{2d} = 0.73$, $\frac{d\chi_{2d}}{dx} = 0.69$.

Now we can discuss the influence of BCS-type pairing on to the PNC $E1$ amplitude. It is well known that due to the pairing the matrix element of the single-particle operator between the quasiparticle states $|0\rangle$, $|1\rangle$ should be multiplied by a factor $u_1u_0 - v_1v_0$ for a T-even operator and by a factor $u_1u_0 + v_1v_0$ for a T-odd operator (see, e.g., [20,22]). u and v are the parameters of Bogoljubov transformation. The levels we consider are very close to the Fermi level. Therefore we get by direct calculation $u_0 \approx u_1 \approx v_0 \approx v_1 \approx 1/\sqrt{2}$, $u_1u_0 + v_1v_0 \approx 1$, $u_1u_0 - v_1v_0 \leq 1/10$. Comparison of a quasiparticle-quadrupole vibration coupling constant with experiment confirms the estimation $u_1u_0 - v_1v_0 \leq 1/10$.

Due to Eq. (12) the PNC $E1$ amplitude is proportional to the matrix of effective operator $i\{(\boldsymbol{\sigma}\cdot\mathbf{p}), [p_z, \rho]\} = \{(\boldsymbol{\sigma}\cdot\mathbf{p}), \frac{d\rho}{dz}\}$. It is T even. Surely it is quite obvious without any calculation: The effective operator corresponds to second order in $E1$ and weak interaction (10), but both the $E1$ amplitude (8) and weak interaction (2) are T even. Thus to take into account the pairing the result of the independent particle shell model calculation (12) should be multiplied by a factor $u_1u_0 - v_1v_0$ [18]. Using Eqs. (12) and (14) we get the value of the PNC $E1$ amplitude:

$$E1_{\rm PNC} = -\frac{Z}{A} g_{np} e\xi (u_1u_0 - v_1v_0) = -i\frac{0.16}{6}\frac{Z}{A} g_{np}\frac{eG\rho_0}{m_p\omega}(u_1u_0 - v_1v_0). \tag{15}$$

Due to Eqs. (4), (7), and (15) the value of the γ-quantum circular polarization for the 23.9 keV transition in ^{119}Sn is equal to

$$P = 2\,{\rm Im}\frac{E1_{\rm PNC}}{M1} = \pm 0.14\frac{G\rho_0}{\omega}(u_1u_0 - v_1v_0) = \pm 3\times 10^{-7}(u_1u_0 - v_1v_0) \sim \pm 3\times 10^{-8}. \tag{16}$$

This value is very small. However we would like to stress that there is no special suppression except the factor $\sim 1/5$ due to compensation in formula (14) and factor $(u_1u_0 - v_1v_0) \sim 1/10$ due to pairing. Excluding these factors Eq. (16) is the maximal possible estimation for spherical nuclei: $P \sim G\rho_0/\omega$. A similar situation exists for the 14.4 keV transition in ^{57}Fe. We do not present the calculations for ^{57}Fe, but it is obvious that $E1_{\rm PNC}$ is of the same order of magnitude as for ^{119}Sn. Due to (7) it means that estimation (16) is valid for ^{57}Fe as well.

It is interesting to notice that due to Eq. (15) $E1_{\rm PNC}$ depends only on g_{np}. There is no dependence on g_{nn} due to compensation between $D^z_{\rm ext}$ and $D^z_{\rm core}$ in Eq. (12). However we would not like to consider this compensation very seriously. It is obtained in a very simplified model of the nucleus. For example, the account of $E1$ giant resonance probably renormalizes $D^z_{\rm core}$ by a factor 1/2 and destroys the compensation in the g_{nn} contribution. However it does not influence the magnitude of the effect.

Finally we would like to compare our estimation (16) with well-known results for the $\frac{5}{2}^+ \to \frac{7}{2}^+$ 483 keV transition in ^{181}Ta [23,18,10]. This transition is predominantly of $E2$ type: 97% $E2$+3% $M1$ [17]. Therefore

$$\begin{aligned} P_{\rm Ta} &= \frac{M1^2}{E2^2 + M1^2} 2\,{\rm Im}\frac{E1_{\rm PNC}}{M1} \\ &= (3\times 10^{-2}) 2\,{\rm Im}\frac{E1_{\rm PNC}}{M1}. \end{aligned} \tag{17}$$

Comparing half-life times for the levels of ^{181}Ta and ^{119}Sn we conclude that $M1_{\rm Ta}/M1_{\rm Sn} \sim 2.4\times 10^{-3}$. If we suppose that $E1_{\rm PNC}$ in Ta is the same as in Sn, we get $P_{\rm Ta} \sim 3.8\times 10^{-7}$. However let us recall that in the Sn frequency $\omega \approx 23.9$ keV is much smaller than the pairing gap Δ, therefore $(u_1u_0 - v_1v_0) \sim 0.1$. It is not the case in ^{181}Ta where $\omega \approx 482$ keV and $\Delta \approx 700$ keV, and therefore $(u_1u_0 - v_1v_0) \sim 0.5$. Thus we have to multiply the above estimation by a factor 5:

$$P_{\rm Ta} \sim \pm 2\times 10^{-6}. \tag{18}$$

The nucleus ^{181}Ta is deformed, but we base on the calculation for spherical ^{119}Sn. Therefore the estimation (18) is very crude. Nevertheless it quite agrees with the experimental value $P_{\rm Ta} = -(5.2\pm 0.5)\times 10^{-6}$ [23] and the results of accurate calculations (see, e.g., [18,10]). This agreement confirms our estimation (16) for ^{119}Sn and ^{57}Fe.

V. CONCLUSION

In the framework of the standard description of a weak interaction in nuclei we get the estimation of γ-quantum degree of circular polarization $P \sim 3\times 10^{-8}$ for Mossbauer transitions in ^{119}Sn and ^{57}Fe nuclei. This result is obtained in an oscillator shell model which is rather crude. Besides that the factor $(u_1u_0 - v_1v_0)$, which takes into account the pairing, was really only estimated. Nevertheless we hardly believe that theoretical value of P could exceed 10^{-7}. Thus the result of calculation strongly contradicts the experimental value $P \sim 10^{-3}$. One may consider this huge disagreement as an indication of the unknown nuclear phenomenon which leads to the enhancement of parity nonconservation. In this connection we would like to mention another phenomenon which presently is not understood. This is the correlation of the signs of PNC effects in the neutron scattering on Th and U nuclei [24]. We stress the importance of new independent measurements of parity nonconservation in Mossbauer transitions.

We are grateful to P. N. Isaev, I. B. Khriplovich, and V. V. Sokolov for their interest in the work and very helpful discussions.

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