Renormalization of the P- and T-odd nuclear potentials by the strong interaction and enhancement of P-odd effective field

V. V. Flambaum and O. K. Vorov

School of Physics, University of New South Wales, P.O.Box 1, Kensington, New South Wales 2033, Australia

(Received 7 September 1993)

Approximate analytical formulas for the self-consistent renormalization of P,T-odd and P-odd weak nuclear potentials by the residual nucleon-nucleon strong interaction are derived. The contact spin-flip nucleon-nucleon interaction reduces the constant of the P,T-odd potential 1.5 times for the proton and 1.8 times for the neutron. Renormalization of the P-odd potential is caused by the velocity dependent spin-flip component of the strong interaction. In the standard variant of $\pi+\rho$ exchange, the conventional strength values lead to anomalous enhancement of the P-odd potential. Moreover, the π-meson exchange contribution seems to be large enough to generate an instability (pole) in the nuclear response to a weak potential.

PACS number(s): 21.30.+y, 13.75.Cs, 24.80.Dc

Recent measurements of the effects of parity nonconservation (PNC) in nuclear reactions produced several results which still have not been explained: permanent sign PNC effects in neutron capture by ^{232}Th [1] and very large PNC effects in Mössbauer transitions [2]. One can consider these observations as a hint that there could be some new mechanisms to enhance the weak interaction in the nucleus. Therefore, it is time to consider possible corrections which can influence the magnitude of PNC effects. In our work [3] it was pointed out that the residual strong interaction can enhance the two-nucleon PNC interaction $\sim A^{1/3}$ times (A is a nucleon number). We called the residual interaction which combines the action of the weak potential with the residual strong interaction the induced parity nonconserving interaction (IPNCI). However, the dominating part of the two-nucleon IPNCI, which was produced by the velocity-independent contact strong interaction, does not contribute to the single-particle weak potential. In the present work we consider the part of the strong interaction which renormalizes the single-particle weak potential. Renormalization of the P- and T-odd nuclear potentials which contribute to P- and T-odd nuclear moments and to P- and T-odd effects in neutron scattering is also considered.

Let us start from the consideration of P- and T-odd nuclear potentials (see, e.g., Ref. [4]):

$$H_{TP} = \frac{G}{2\sqrt{2}m}\eta(\sigma\nabla)\rho \simeq \theta\sigma\nabla U \ ,$$
$$\theta = \eta\frac{G}{2\sqrt{2}m}\frac{\rho(0)}{U(0)} = -2\times 10^{-8}\eta \text{ fm}, \tag{1}$$

where σ is the doubled nucleon spin, ρ is the nuclear density, m is the proton mass, G is the Fermi constant, η is a dimensionless constant characterizing the strength of T- and P-odd interactions [the limits on these constants for protons (η_p) and neutrons (η_n) were obtained from atomic [5] and molecular [6] electric dipole moment measurements], and U is the strong nuclear potential, 1 fm$=10^{-13}$ cm. The shape of the potential U and the nuclear density ρ is known to be approximately similar. We used this fact in Eq. (1). Correspondingly, the whole potential affecting the nuclear motion is equal to

$$\bar{U} = U + H_{TP} = U(\mathbf{r}) + \theta\sigma\nabla U \simeq U(\mathbf{r}+\theta\sigma). \tag{2}$$

Hence, it is obvious that the nucleon wave function with the H_{TP} taken into account has the form

$$\bar{\psi} = \psi(\mathbf{r}+\theta\sigma) = (1+\theta\sigma\nabla)\psi(\mathbf{r}) = \psi + \delta\psi, \tag{3}$$

where $\psi(\mathbf{r})$ is the nonperturbed wave function. The direct correction to the strong potential induced by a small perturbation can be written as follows:

$$\delta V(1) = \sum_a \int d2[\delta\psi_a^\dagger(2)V(1,2)\psi_a(2) + \psi_a^\dagger(2)V(1,2)\delta\psi_a(2)]. \tag{4}$$

Here the notation $1(2) \equiv [\mathbf{r}_{1(2)}, \sigma_{1(2)}, \tau_{1(2)}]$ stands for the full set of the nucleon variables (coordinate, spin, and isospin) and the summation is carried out over the occupied nucleon states a.

We use the Landau-Migdal parametrization of the strong interaction,

$$V(\mathbf{r},\mathbf{r}') = C\delta(\mathbf{r}-\mathbf{r}')[f_0 + f_0'\tau\tau' + g\sigma\sigma' + g'\tau\tau'\sigma\sigma'], \tag{5}$$

where $C = 300$ MeV fm^3, $g = 0.575, g' = 0.725$, and only the direct terms are considered (see, e.g., Refs. [7,8]). Using Eq. (3) for $\psi + \delta\psi$ and integration by parts in Eq. (4), we obtain the correction to the T,P-odd potential:

$$\tilde{H}_{TP} = -\sum_a \theta_2 \int d^3r\psi_a^\dagger(\mathbf{r}_2)[\sigma_2\nabla_2, V(\mathbf{r}_1,\mathbf{r}_2)]\psi_a(\mathbf{r}_2) = \sum_a \theta_2(g+g'\tau_1\tau_a)(\sigma_1\nabla)|\psi_a|^2 = \gamma\sigma\nabla\rho \ . \tag{6}$$

$\gamma = C[\theta_p\frac{Z}{A}(g\pm g') + \theta_n\frac{N}{A}(g\mp g')]$ for protons (neutrons). Hereafter, [,] means commutator, and $\rho = \sum_a|\psi_a|^2$. We put the proton density $\rho_p = \frac{Z}{A}\rho$, neutron density $\rho_n = \frac{N}{A}\rho$, and $\langle\sigma_2\rangle = 0$ (we consider the potential created by

 ©1994 The American Physical Society

paired nucleons). Now we should solve the self-consistent equation $H_{TP} = H^0_{TP} + \tilde{H}_{TP}$ for T, P-odd potential:

$$\theta\sigma\nabla U = \theta^0\sigma\nabla U + \gamma\frac{\rho(0)}{U(0)}\sigma\nabla U. \tag{7}$$

Here H^0_{TP} contains the "initial" values of the T, P-odd interaction constants η^0_p and η^0_n (or θ^0_p and θ^0_n), while H_{TP} and $\tilde{H}_{TP}$ contain "final" values of the constants. The solutions for the pair of simple linear algebraic equations for the constants are the following:

$$\eta_p = \frac{1}{D}\{\eta^0_p[1+\tilde{C}(g+g')N/A] - \eta^0_n\tilde{C}(g-g')N/A\} \simeq \frac{\eta^0_p}{1.5},$$
$$\eta_n = \frac{1}{D}\{\eta^0_n[1+\tilde{C}(g+g')Z/A] - \eta^0_p\tilde{C}(g-g')Z/A\} \simeq \frac{\eta^0_n}{1.8},$$
$$D = [1+\tilde{C}(g+g')N/A][1+\tilde{C}(g+g')Z/A] - \tilde{C}^2(g-g')^2ZN/A^2 .$$

Here, $\tilde{C} = C\rho/|U| = \frac{4}{3}\frac{\varepsilon_F}{|U|} = \frac{4}{3}(1+\frac{|\varepsilon|}{\varepsilon_F})^{-1} \simeq 1$ and η^0_p and η^0_n are the initial values of the constants. We used the well known relations :

$$C = \frac{\pi^2}{p_F m}, \quad \rho = \frac{2p_F^3}{3\pi^2}, \quad \varepsilon_F = \frac{p_F^2}{2m}, \quad |U| = \varepsilon_F + |\varepsilon|, \tag{8}$$

where p_F is a Fermi momentum and $|\varepsilon|$ is a nucleon separation energy. We also have taken into account in the numerical estimate that $|g-g'|$ is small.

Thus, the strong residual interaction reduces the values of the T, P-odd potential constants 1.5–1.8 times. Note that the response of the nucleus to the T- and P-odd potentials (1) as a function of the interaction constants has poles ($D = 0$) at $g = \tilde{C}^{-1} \simeq -1$ and $g' \simeq \tilde{C}^{-1} \simeq -1$ (for $N \simeq Z$). The positions of the poles differ from the instability points in an infinite Fermi system $g = g' = -1.5$ (see, e.g., Refs. [9,10]) since the interactions (1) do not exist in the infinite system ($H_{PT} = 0$ at $\rho = \text{const}$).

It is interesting that the T- and P-odd interactions induce a spin hedgehog ($\sigma \sim \mathbf{r}$) in the nucleon spin distribution within a spherical nucleus. A simple calculation with the wave function (3) gives the following proton and neutron spin distributions:

$$\sigma_p(\mathbf{r}) = \theta_p\nabla\rho_p(\mathbf{r}), \quad \sigma_n(\mathbf{r}) = \theta_n\nabla\rho_n(\mathbf{r}) . \tag{9}$$

The interaction $\tilde{H}_{TP}$ in Eq. (6) is, in fact, a strong interaction of the nucleon with the spin hedgehog $[Cg\sigma\sigma(\mathbf{r})]$.

Now we turn to considering corrections to the weak P-odd and T-even potentials,

$$W = \frac{G}{2\sqrt{2}m}g^W(\sigma\mathbf{p}\rho + \rho\sigma\mathbf{p}). \tag{10}$$

Here p is the nucleon momentum; the dimensionless constants g^W_p for the proton and g^W_n (for the neutron) are of order of unity (the notation $\varepsilon \simeq 1.0 \times 10^{-8}g$ is also adopted in the current literature). In a simple model of a constant nuclear density it is easy to find the result of the action of the perturbation W (see Ref. [11] and the first paper of Ref. [12]):

$$\tilde{\psi} = \exp(-i\xi\sigma\mathbf{r})\psi(\mathbf{r}) \simeq (1 - i\xi\sigma\mathbf{r})\psi(\mathbf{r}) ,$$
$$\xi = \frac{G}{\sqrt{2}}g^W\rho = \varepsilon m . \tag{11}$$

In the general case (real density shape and spin-orbit interaction taken into account) the correction to the wave function contains an extra spherically symmetric function $\varphi_a(r)$ (see, e.g., Ref. [13]):

$$\delta\tilde{\psi}_a = -i(\sigma\mathbf{r})\varphi_a(r)\psi_a(\mathbf{r}). \tag{12}$$

The P-odd weak interaction (10) also changes the spin distribution. It rotates the spin around the vector $\mathbf{r}$ [see Eq. (11)] by the angle ξr and creates a spin spiral [13]. However, after summation over paired nucleons this spin structure disappears. As a result, the contact spin-dependent strong interaction (5) does not contribute to the renormalization of the weak potential [because of the factor i in Eqs.(11) and (12) the contribution of $(\delta\psi^\dagger)$ compensates the contribution from $(\delta\psi)$ in Eq. (4) for the correction to the potential]. This result looks natural since the only possible orientation of the spin in the spherical nucleus $\sigma \sim \mathbf{r}$ violates both P and T invariances and cannot be produced by a T-even weak interaction (10).

The correlation which is actually produced by the weak interaction is $\sigma\mathbf{p}$. To reveal such structures the strong interaction must be spin and momentum dependent as well (another possibility is related to a finite range exchange interaction; it will be considered below). Within Landau-Migdal theory, the momentum dependence is usually described [9] by the following extra term in (5):

$$V_1 = \frac{1}{4}Cp_F^{-2}(g_1 + g_1'\tau_1\tau_2)(\sigma_1\sigma_2)[\mathbf{p}_1\mathbf{p}_2\delta(\mathbf{r}_1-\mathbf{r}_2) + \mathbf{p}_1\delta(\mathbf{r}_1-\mathbf{r}_2)\mathbf{p}_2 + \mathbf{p}_2\delta(\mathbf{r}_1-\mathbf{r}_2)\mathbf{p}_1 + \delta(\mathbf{r}_1-\mathbf{r}_2)\mathbf{p}_1\mathbf{p}_2]. \tag{13}$$

The constants of this interaction are found to be $g_1 = -0.5$, $g_1' = -0.26$ [9]. Using Eqs. (4), (11), and (13) we can calculate the corresponding correction to the weak potential:

$$\tilde{W} = \sum_a \int d^3r_2 \psi_a^\dagger(\mathbf{r}_2)[i\xi_a\sigma_2\mathbf{r}_2, V_1]\psi_a(\mathbf{r}_2)$$
$$= -\sum_a \xi_a(g_1 + g_1'\tau_1\tau_a)\frac{C}{2p_F^2}\{\sigma_1\mathbf{p}_1, |\psi_a|^2\} = Q[(\sigma\mathbf{p})\rho + \rho(\sigma\mathbf{p})], \tag{14}$$

where $Q = -\frac{C}{2p_F^2}[\frac{Z}{A}(g_1 \pm g_1')\xi_p + \frac{N}{A}(g_1 \mp g_1')\xi_n]$ for protons (neutrons) correspondingly. A self-consistent solution of the equation for the total P-odd nuclear potential $W = W^0 + \tilde{W}$ gives the following values of the potential constants:

$$g_p^W = \frac{1}{D}\left\{g_p^{W0}\left[1 + \frac{2N}{3A}(g_1 + g_1')\right] - \frac{2N}{3A}g_n^{W0}(g_1 - g_1')\right\},$$
$$g_n^W = \frac{1}{D}\left\{g_n^{W0}\left[1 + \frac{2Z}{3A}(g_1 + g_1')\right] - \frac{2Z}{3A}g_p^{W0}(g_1 - g_1')\right\},$$
$$D = \left[1 + \frac{2N}{3A}(g_1 + g_1')\right]\left[1 + \frac{2Z}{3A}(g_1 + g_1')\right] - \frac{4NZ}{9A^2}(g_1 - g_1')^2. \tag{15}$$

We have taken into account here that $C\rho m p_F^{-2} = 2/3$. It is interesting that the poles ($D = 0$) in the response of a nucleus to the weak potential $W \sim \sigma\mathbf{p}$ coincide with the boundary of stability for a Fermi liquid with the interaction (13): $g_1 = g_1' = -1.5$ at $N = Z$ (see, e.g., Refs. [9,10]). This is not too surprising since we used the approximation $\rho =$ const to obtain the wave function (11).[1]

The interaction V_1 with the constants $g_1 = -0.5$, $g_1' = -0.26$ does not cause instability. However, it acts in the direction of the poles and increases the P-odd potential:

$$g_p = 1.3g_p^0 + 0.18g_n^0, \quad g_n = 1.4g_n^0 + 0.12g_p^0. \tag{16}$$

Therefore, the Landau-Migdal interaction $V + V_1$ [Eqs. (5) and (13)] does not produce crucial changes in the values of interaction constants for P,T-odd and P-odd potentials. The corrections are of the same size as, say, corrections to the Schmidt values of the magnetic moments. In fact, the Landau-Migdal interaction originates from the underlying ($\pi + \rho$)-exchange interaction [14] which generates also tensor components. To account for the latter destroys this "idyllic" picture at least for the P-odd potential. The rest of the paper is devoted to the calculation of the $\pi + \rho$ contribution.

In the present random-phase-approximation-like calculations, the correction to the nucleon P-odd potential $\tilde{W} = \sum_{\mu\nu}\delta w_{\mu\nu}\mu^\dagger\nu$ due to the strong interaction $\hat{V} = \frac{1}{2}\sum_{abcd}a^\dagger c^\dagger V_{abcd}db$ is given by the expression

$$\delta w_{\mu\nu} = \sum_{ab}(A_{ab}V_{ba\mu\nu} - V_{ab\mu\nu}A_{ba})n_a - \sum_{ab}(A_{ab}V_{b\nu\mu a} - V_{a\nu\mu b}A_{ba})n_a, \tag{17}$$

where the first sum is the direct contribution and the second is exchange one ($\mu^\dagger, \mu, a^\dagger, a$, ... are creators and destructors of nucleons in the corresponding single-particle states), n_a are the occupation numbers, $n_a \equiv \langle a^\dagger a\rangle$, and $A_{ab} = \langle\psi_a|i\xi(\sigma\mathbf{r})|\psi_b\rangle$ are single-particle matrix elements of the P-odd mixing operator [see Eq. (11)]. For the $\pi + \rho$ interaction [14] $\hat{V}_{abcd}$ is given by

$$V_{abcd} = \int d1d2\psi_a^\dagger(1)\psi_c^\dagger(2)V^{\pi+\rho}(1,2)\psi_b(1)\psi_d(2), \tag{18}$$

where $V^{\pi+\rho}(1,2)$, in p representation, is

$$V^{\pi+\rho}(1,2) = -4\pi(\tau_1\tau_2)\left[\frac{f_\pi^2}{m_\pi^2}\frac{(\sigma_1\mathbf{q})(\sigma_2\mathbf{q})}{q^2 + m_\pi^2} + \frac{f_\rho^2}{m_\rho^2}\frac{[\sigma_1 \times \mathbf{q}][\sigma_2 \times \mathbf{q}]}{q^2 + m_\rho^2}\right], \tag{19}$$

where $\mathbf{q}$ is the momentum transfer, $m_{\pi(\rho)}$ is the pion (rho meson) mass, $f_\pi^2 = 0.08$ is a pion coupling constant, and f_ρ^2 is corresponding ρ-meson coupling ranging from 1.86 to 4.86 ("weak" and "strong" couplings correspondingly [14]). In the coordinate representation, the last expression becomes a potential depending on $|\mathbf{r}_1 - \mathbf{r}_2|$. Thus its commutator with $A = i\xi(\sigma\mathbf{r})$ in (17) (direct terms) is zero, while the exchange terms contain (due to nonlocality of the potential) an effective velocity dependence and yield a nonzero contribution to $\tilde{W}$. To calculate the latter, we should reduce

[1]We had known from a private communication with V.G. Zelevinsky that he independently obtained a similar result: The correction to the effective field $\sigma\mathbf{p}$ diverges at the same point where the first harmonic of the Landau interaction $g_1(\sigma_1\sigma_2)(\mathbf{p}_1\mathbf{p}_2)$ leads to the instability of the Fermi liquid.

the exchange terms in (17) to a direct form which requires the change $\mathbf{q} \to \mathbf{p}_1 - \mathbf{p}_2$ (the nucleons are on the Fermi surface) and Fierz transformation of the spin and isospin tensor structures [8]. After performing that, we obtain, for $V_{b\nu\mu a}$,

$$V_{b\nu\mu a} = \int d1d2\psi_b^\dagger(2)\psi_\mu^\dagger(1)V'(1,2)\psi_\nu(1)\psi_a(2),$$

with $V'(1,2)$ being equal to

$$\begin{aligned} V'(1,2) = & -2\pi\left(\frac{3}{2} - \frac{1}{2}(\tau_1\tau_2)\right)\sum_{\alpha\beta}[2\sigma_{1\alpha}\sigma_{2\beta} + \{1-(\sigma_1\sigma_2)\}\delta_{\alpha\beta}] \\ & \times\left[\frac{f_\pi^2}{m_\pi^2}\frac{(\mathbf{p}_1-\mathbf{p}_2)_\alpha(\mathbf{p}_1-\mathbf{p}_2)_\beta}{(\mathbf{p}_1-\mathbf{p}_2)^2+m_\pi^2} + \frac{f_\rho^2}{m_\rho^2}\frac{\delta_{\alpha\beta}(\mathbf{p}_1-\mathbf{p}_2)^2 - (\mathbf{p}_1-\mathbf{p}_2)_\alpha(\mathbf{p}_1-\mathbf{p}_2)_\beta}{(\mathbf{p}_1-\mathbf{p}_2)^2+m_\rho^2}\right] . \end{aligned} \tag{20}$$

By use of that, the second sum in Eq. (17) is reduced to the expectation value of the commutator $[i\xi_2(\sigma_2\mathbf{r}_2), V'(1,2)]$ and we obtain the meson exchange correction $\tilde{W}$ to the P-odd potential acting on the first nucleon:

$$\tilde{W}^{\pi+\rho} = -\langle\psi_a^\dagger(2)|[i\xi(\sigma_2\mathbf{r}_2), V'(1,2)]|\psi_a(2)\rangle . \tag{21}$$

Here notation $\langle\cdots\rangle$ stands for the expectation value taken in the subspace of the wave functions of the core nucleons (label 2), and the summation is assumed over the states a occupied by the nucleon 2. Calculating the commutator in (21) using (20) and the relations $[\mathbf{r}_{2\alpha}, \mathbf{p}_{2\beta}] = i\delta_{\alpha\beta}$, $\langle\sigma_{2\alpha}\sigma_{2\beta}\rangle = \delta_{\alpha\beta}$, we obtain

$$\begin{aligned} \tilde{W}^{\pi+\rho} = & K_\pi\langle\psi_a^\dagger(2)|[\sigma_1(\mathbf{p}_1-\mathbf{p}_2)]\left[\frac{1}{(\mathbf{p}_1-\mathbf{p}_2)^2+m_\pi^2} - \frac{1}{3}\frac{(\mathbf{p}_1-\mathbf{p}_2)^2}{[(\mathbf{p}_1-\mathbf{p}_2)^2+m_\pi^2]^2}\right]|\psi_a(2)\rangle \\ & -K_\rho\langle\psi_a^\dagger(2)|[\sigma_1(\mathbf{p}_1-\mathbf{p}_2)]\left[\frac{1}{(\mathbf{p}_1-\mathbf{p}_2)^2+m_\rho^2} - \frac{1}{2}\frac{(\mathbf{p}_1-\mathbf{p}_2)^2}{[(\mathbf{p}_1-\mathbf{p}_2)^2+m_\rho^2]^2}\right]|\psi_a(2)\rangle , \end{aligned} \tag{22}$$

where the right hand side remains to be an operator acting on the wave functions of nucleon 1, and the constants are $K_\pi = 6\pi(f_\pi^2/m_\pi^2)[3-(\tau_1\tau_2)]\xi_2$, $K_\rho = 8\pi(f_\rho^2/m_\rho^2)[3-(\tau_1\tau_2)]\xi_2$. To evaluate the expression (22), one can employ, e.g., the Fermi-gas approximation to parametrize the density of core nucleons $\sum_a \psi_a^\dagger(2)\psi_a(2)$ as has been widely used in such calculations (e.g., in obtaining the "bare" nucleon P-odd potential [15]). We obtain from Eq. (22)

$$\tilde{W}^{\pi+\rho} = 2Q^{\pi+\rho}\rho_0(\sigma_1\mathbf{p}_1) , \tag{23}$$

where the constant $Q^{\pi+\rho}$ for the proton and neutron has the following form:

$$Q_p^{\pi+\rho} = q\left(\xi_p\frac{Z}{A} + 2\xi_n\frac{N}{A}\right) , \quad Q_n^{\pi+\rho} = q\left(\xi_n\frac{N}{A} + 2\xi_p\frac{Z}{A}\right) , \quad q = 6\pi\left[\frac{f_\pi^2}{m_\pi^4}W_\pi\left(\frac{p_F}{m_\pi}\right) - \frac{4}{3}\frac{f_\rho^2}{m_\rho^4}W_\rho\left(\frac{p_F}{m_\rho}\right)\right] , \tag{24}$$

and the nonlocality factors W ($W_{\pi,\rho} \to 1$ for $m_{\pi,\rho} \to \infty$) are $W_\pi(\frac{p_F}{m_\pi}) = 0.11$, $W_\rho(\frac{p_F}{m_\rho}) = 0.69$ for the pion and ρ meson correspondingly. The nonlocality effect is greater for the pion due to its smaller mass ($m_\pi = 0.7\,\text{fm}^{-1}$ compared to $p_F \simeq 1.3$ fm^{-1}, while $m_\rho = 3.7$ fm^{-1}). The above value W_π is quite close to the result $W_\pi = 0.16$ for the nonlocality factor for the "bare" weak potential obtained in α-cluster calculations [12].

To obtain the renormalization of the P-odd weak potential with account for $\pi + \rho$ exchange, one should use $\tilde{W}^{\pi+\rho}$ instead of $\tilde{W}$ in Eq. (15) in the self-consistent determination of W. With account for that, the renormalization equations of the potential constants $g_{p,n}$ (15) take the form

$$\begin{aligned} g_p^W &= \frac{1}{D}\left\{g_p^{W0}\left[1-\frac{N}{A}k\right] + 2\frac{N}{A}g_n^{W0}k\right\} , \\ g_n^W &= \frac{1}{D}\left\{g_n^{W0}\left[1-\frac{Z}{A}k\right] + 2\frac{Z}{A}g_p^{W0}k\right\} , \end{aligned} \tag{25}$$

where $k = 2q\rho m$ and, in that case, the determinant D is equal to

$$D = \left(1-\frac{N}{A}k\right)\left(1-\frac{Z}{A}k\right) - \frac{NZ}{A^2}4k^2 . \tag{26}$$

It is seen from the last term in Eq. (24) that the contribution from ρ-meson exchange tends to compensate the effect of the π meson, whereas the latter strongly pushes the solution [Eq. (28)] in the direction of the pole ($D = 0$). The equation $D = 0$ determines a curve (function of N/A) corresponding to the border of stability of the nuclear response to the P-odd field. For real nuclei ($N/A \simeq 0.5$–0.6) the position of the pole corresponds to the critical value of $k = k_c = 0.67$. The π meson alone (with no ρ exchange) gives $k = k_\pi \simeq 1$ and produces instability in the "shell-model" nucleus. The ρ-meson exchange reduces the value of k: "strong" ρ-meson coupling ($f_\rho^2 = 4.86$) gives $k = 0.4$ which corresponds to enhancement factors $g_p^W/g_p^{W0} = 1.6$, $g_n^W/g_p^{W0} = 0.7$ (for $g_p^{W0} = 4$

and $g_n^{W0} \simeq 0$, see, e.g., [12]). Thus $g_n^W \sim g_p^W$ even for very small initial values of g_n^{W0}. "Weak" coupling ($0.4f_\rho^2$) gives $k \simeq 0.7 \simeq k_c$ ("infinite" enhancement). Of course, the accuracy of the present consideration is not sufficient to give a definite answer in this situation (note that we have considered the linear response only and neglected fine effects like smoothing of the pion in nuclear matter [7,16]), besides the uncertainty in π and ρ coupling constants in a nucleus. At least one can say that $D \simeq 0$ means a possibility of strongly enhanced P-odd effects.

Interpretation of this fact, resulting mostly from the strong π-meson exchange contribution, is not straightforward: Definitely, it is related to the question of the stability of a nucleus under the tensor π-exchange interaction which has been already widely discussed in the literature [8], in particular, in relation to the problem of π condensation in nuclei (see, e.g., [7,16]). Also, a large enhancement factor is naturally associated with the low-lying 0^- excitation (a pole in D for nonzero frequency of the PNC field). The influence of the 0^- resonance on the PNC effects was discussed in Refs. [17,18]. On the other hand, some effects decreasing the role of the π-exchange interaction in a real nucleus may exist and the stability of the nucleus may be restored (e.g., due to a particular shell structure). At least, our consideration proves that a possible mechanism exists leading to the enhancement of the nuclear P-odd weak potential which is caused by the velocity-dependent spin-flip component of the conventional residual strong interaction with the standard values of its constants. Thus, new reliable experimental information on the P-odd nuclear effects would be desirable. In view of the present considerations, it might be important not only for the weak interaction theory, but also for the study of strong interaction effects and nuclear structure.

To conclude, we have considered the renormalization of the nuclear T,P-odd and P-odd potentials due to the residual strong interaction in the Landau-Migdal parametrization and with account for the tensor component of the one π,ρ-meson exchange. The T,P-odd potential is found to be renormalized moderately, while the renormalization of the P-odd potential proves to be greater and the tensor velocity-dependent interaction, with standard values of the parameters, turns out to be able to produce a substantial enhancement, "driving" the solution for the self-consistent P-odd field towards the region of instability.

---

[1] C.M. Frankle *et al.*, Phys. Rev. Lett. **67**, 564 (1991).

[2] L.V. Inzhechik *et al.*, Yad. Fiz. **44**, 1370 (1986) [Sov. J. Nucl. Phys. **44**, 890 (1986)]; Zh. Eksp. Teor. Fiz. **93**, 800 (1987) [Sov. Phys. JETP **66**, 450 (1987)]; **93**, 1569 (1987) [**66**, 897 (1987)]; V.G. Tsinoev, talk at the International Symposium on Weak and Electromagnetic interactions in Nuclei, Dubna, 1992; O.P. Sushkov and V.B.Telitsin, Phys. Rev. C (to be published).

[3] V.V. Flambaum and O.K. Vorov (unpublished); see also Phys. Rev. Lett. **70**, 4051 (1993).

[4] V.V. Flambaum, I.B. Khriplovich, and O.P. Sushkov, Nucl. Phys. **A449**, 750 (1986).

[5] S.K. Lamoreax *et al.*, Phys. Rev. Lett. **59**, 2275 (1987).

[6] D. Cho, K. Sangster, and E.A. Hinds, Phys. Rev. A **44**, 1783 (1991).

[7] A.B. Migdal, *Theory of Finite Fermi Systems and Applications to Atomic Nuclei* (John Wiley & Sons, New York, 1967).

[8] G.E. Brown, Rev. Mod. Phys. **43**, 1 (1971); V. Klemt, S.A. Moszkowski, and J. Speth, Phys. Rev. C **14**, 302 (1976); J. Speth, E. Werner, and W. Wild, Phys. Rep. **33**, 127 (1977), and references therein; G. Bertsch, D. Cha, and H. Toki, Phys. Rev. C **24**, 533 (1981); V.A. Khodel and E.E. Sapershtein, Phys. Rep. **92**, 183 (1982), and references therein; J.W. Negele, Rev. Mod. Phys. **54**, 913 (1982); R. De Haro, S. Krewald, and J. Speth, Nucl. Phys. **A388**, 265 (1982); K. Goeke and J. Speth, Annu. Rev. Nucl. Part. Sci. **32**, 65 (1982); F. Osterfeld, Rev. Mod. Phys. **64**, 491 (1992), and references therein.

[9] V.A. Khodel and E.E. Sapershtein, Phys. Rep. **92**, 183 (1982).

[10] D. Pines and P. Nozieres, *The Theory of Quantum Liquids* (Benjamin, New York, 1966).

[11] F.C. Michel, Phys. Rev. **133**, 329 (1964).

[12] V.V. Flambaum, I.B. Khriplovich, and O.P. Sushkov, Phys. Lett. **146B**, 367 (1984); V.F. Dmitriev, V.V. Flambaum, O.P. Sushkov, and V.B. Telitsin, *ibid.* **125B**, 1 (1983); V.V. Flambaum, V.B. Telitsin, and O.P. Sushkov, Nucl. Phys. **A444**, 611 (1985).

[13] I.B. Khriplovich, *Parity Nonconservation in Atomic Phenomena* (Gordon and Breach, New York, 1991); Ya.B. Zeldovich, Zh. Eksp. Teor. Fiz. **33**, 1531 (1957).

[14] G.E. Brown, S.-O. Backman, E. Oset, and W. Weise, Nucl. Phys. **A286**, 191 (1976); M.R. Anastasio and G.E. Brown, *ibid.* **A285**, 516 (1977); J. Speth, E. Werner, and W. Wild, Phys. Rep. **33**, 127 (1976); J. Speth, V. Klemt, J. Wambach, and G.E. Brown, Nucl.Phys. **A343**, 382 (1980); S. Krewald, K. Nakayama, and J. Speth, Phys. Rep. **161**, 103 (1988); S. Krewald, F. Osterfeld, J. Speth, and G.E. Brown, Phys. Rev. Lett. **46**, 103 (1981); S. Krewald and J. Speth, *ibid.* **45**, 417 (1980); K. Nakayama and W.G. Love, Phys. Rev. C **38**, 51 (1988).

[15] E.G. Adelberger and W.C. Haxton, Annu. Rev. Nucl. Part. Sci. **35**, 501 (1985).

[16] S.A. Fayans, E.E. Sapershtein, and S.V. Tolokonnikov, Nucl. Phys. **A326**, 463 (1979); D.V. Aleksandrov *et al.*, *ibid.* **A436**, 338 (1985); J.Bang *et al.*, *ibid.* **A440**, 445 (1985).

[17] S. G. Kadmensky, V. P. Markushev, and V.I. Furman, Yad. Fiz. **37**, 581 (1983) [Sov. J. Nucl. Phys. **37**, 345 (1983)].

[18] N. Auerbach and J.D. Bowman, Phys. Rev. C **46**, 2582 (1992).