539.141

# *PARITY-NONCONSERVING NUCLEAR FORCES**

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Usp. Fiz. Nauk 95, 647–655 (August, 1968)

## I. WHAT REMAINS OF THE BARE WEAK INTERACTION?

THE hypotheses of universal weak interaction predicts the existence of parity-nonconserving exchange forces between the proton and the neutron in first order in the constant $G = 10^{-5}/m^2$ (m - mass of the nucleon, $\hbar = c = 1$). When T-invariance is violated, these forces can lead to the appearance of electric dipole moments in hadrons.

According to modern-day notions, the weak nucleon current should be the difference between the polar and axial vectors (V – A). The bare form of the current is distorted by strong interaction even in nucleon-lepton processes. In nucleon scattering due to weak forces, the role of strong interactions is even greater, since pion exchange is possible between all particles taking part in the process. Apparently only two properties of the bare "current" Hamiltonian of the weak interaction survive, namely the order of magnitude of the coupling constant and the isotropic structure.

The relative magnitude of weak internucleon forces in nuclei can be reasonably characterized by the dimensionless parameter

$$F=\left(\frac{\mathrm{Sp}\,H_W^2}{\mathrm{Sp}\,H^2}\right)^{1/2}; \tag{1}$$

here $H_W$ and H are the Hamiltonians of the weak and strong interactions of the nucleons, and Sp stands for the trace of the matrix. The internucleon distances in the nuclei have an order of magnitude $1/\mu$ ($\mu$ - pion mass). This gives grounds for expecting, in order of magnitude

$$F\simeq G\mu^2=10^{-5}\left(\frac{\mu}{m}\right)^2\approx 3\cdot 10^{-7}. \tag{2}$$

In some cases, however, the effects of weak nucleon interactions can be greatly enhanced compared with the indicated magnitude. This question is discussed in Sec. II.

Inasmuch as strong interaction is isoscalar, it does not change the isospin structure of the bare current Hamiltonian. The weak nucleon current is an isovector, and the baryon current with unity strangeness change is an isospinor. It follows therefore that the effective Hamiltonian of the weak interaction should be a superposition of an isoscalar, an isovector, and a symmetrical isotensor of second rank with zero trace. It can be expected that the isovector term, obtained by squaring the isospin current, is suppressed by a factor of 15–20 compared with the isoscalar one, owing to the Cabibbo factor. Further, the octet enhancement hypothesis, based on the $SU_3$ symmetry[1], predicts an increase of the isoscalar by an order of magnitude or more compared with the "normal" value of F given in (2).

In the nonrelativistic approximation, the T-invariant Hamiltonian of the parity-nonconserving interaction between nucleons can be written in the form

$$H_W=\boldsymbol{\tau}_1\boldsymbol{\tau}_2U_S+[\boldsymbol{\tau}_1\boldsymbol{\tau}_2I]_zU_V+\left(\tau_{1z}\tau_{2z}-\frac{1}{3}\boldsymbol{\tau}_1\boldsymbol{\tau}_2\right)U_T, \tag{3*}$$

where $\boldsymbol{\tau}_1$ and $\boldsymbol{\tau}_2$ are the nucleon isospin operators, and U are operators acting on the coordinates and the spin variables:

$$U_{S,T}=\frac{\mathbf{r}}{r}[\boldsymbol{\sigma}_1\boldsymbol{\sigma}_2]V_{S,T}(\mathbf{r})+V_{S,T}^{(1)}(\mathbf{r})\mathbf{p}(\boldsymbol{\sigma}_1-\boldsymbol{\sigma}_2), \tag{4}$$

$$U_V=\frac{\mathbf{r}}{r}(\boldsymbol{\sigma}_1+\boldsymbol{\sigma}_2)V_V(\mathbf{r}); \tag{5}$$

here r is the distance between nucleons, $\boldsymbol{\sigma}_1$ and $\boldsymbol{\sigma}_2$ are the spin operators, $\mathbf{p}_1$ is the momentum operator, and V(r) are short-range potentials. The spin part of the isovector term $U_V$ differs from $U_{S,T}$ owing to the requirement of T-invariance (the operator $[\boldsymbol{\tau}\times\boldsymbol{\tau}_2]_z$ reverses sign under time reversal, inasmuch as $\tau_{iy}\rightarrow-\tau_{iy}$).

With regards to the radial dependence of the potentials V(r), we can advance only qualitative considerations. The T-invariance forbids one-pion exchange for the isoscalar and isotensor parts of the Hamiltonian $H_W^2$, and therefore the weak-interaction periphery should be determined essentially by two-pion exchange, i.e., by a radius of the order of $1/2\ \mu$.

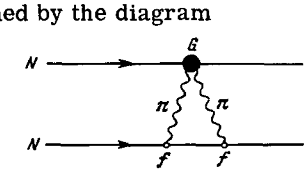
In concrete calculations of probabilities of various kinds of the effects of manifestation of weak interaction of nucleons, use was made of a δ-potential corresponding to the product of the nucleon currents[3]:

$$V_S^{(1)}(\mathbf{r})=-\frac{2}{3}V_T^{(1)}(\mathbf{r})=\frac{G}{4m}\delta(\mathbf{r}),\qquad V_S=V_T=0, \tag{6}$$

and a potential due to two-pion exchange[4]:

$$V_S=-\frac{2}{3}V_T=-\frac{G}{3\pi}f^2\left(\frac{1}{r^4}+\frac{2}{\mu r^5}+\frac{1}{\mu^2r^0}\right)e^{-2\mu r},$$
$$f^2=0.08,\qquad V_S^{(1)}=V_T^{(1)}=0. \tag{7}$$

This potential corresponds to the scattering amplitude determined by the diagram




Formula (3) for $H_W$ was written under the assumption that the bare weak currents are charged. If there exist also neutral currents, then it is necessary to add to the Hamiltonian (3) non-exchange isospin terms (an isoscalar containing no isospin operator, and an isovector of the type $(\tau_{1z}+\tau_{2z})U$). The spin structure of these operators coincides with (4).

*Paper delivered at International Seminar on Problems of CP-violation (Moscow, 22-26 January, 1968).

*$[\tau_1\tau_2I]\equiv[\tau_1\times\tau_2I]$

## II. ENHANCEMENT EFFECTS

We consider the influence of parity nonconservation in nuclear forces on the electromagnetic transitions of nuclei (affirmative experimental results on the observation of weak forces are available so far only for these processes).

We can point to three sources of enhancement of nuclear phenomena due to parity nonconservation in the interaction between nucleons:

1) kinematic enhancement,

2) structural features of the lower states of the nuclei,

3) dynamic enhancement due to the closeness of levels of different parity.

1. Kinematic enhancement. Weak interactions of nucleons are manifest in nuclear forces by the fact that the nuclear states cease to have a definite parity. As a result, in particular, the emission of electric (EL) and magnetic (ML) radiation of identical multiplicity L becomes possible in one and the same transition.

Kinematic enhancement is connected with the fact that, other conditions being equal, the amplitude ; M(ML) is suppressed compared with the amplitude M(EL) by a factor on the order of v/c, where v is the effective velocity of the intranuclear nucleon. The effects of manifestation of weak forces in radiative transitions are determined by the quantity

$$RF=\frac{M(EL)\,M(ML)}{(M(EL))^2+(M(ML))^2} \tag{8}$$

(we assume the radial matrix elements M to be real). If only a magnetic transition is allowed in the absence of weak forces, then $|M(EL)| \ll |M(ML)|$, and therefore

$$RF=\frac{M(EL)}{M(ML)}, \quad M(EL)\sim F. \tag{9}$$

Putting further $M(ML) \cong c/v$, we obtain from (9) $R \cong c/v$, which has an order of magnitude close to 10 for medium and heavy nuclei.

2. Structural enhancement. This effect takes place when a process allowed in the absence of weak forces becomes suppressed as a result of structure singularities in the nuclear states that take part in the transition. The enhancement coefficient R can in general be represented in the form

$$R=\frac{1}{M_{0f}}\sum_n \alpha_{0n}M_{nf}, \tag{10}$$

where $\alpha_{0n}$ is the amplitude of the admixture of the state $|n\rangle$ with opposite parity to the "principal" state $|0\rangle$. In first order of perturbation theory we have

$$\alpha_{0n}=\frac{1}{F}\,\frac{\langle n|H_W|0\rangle}{E_0-E_n} \tag{11}$$

($E_0$, $E_n$ - level energies). The quantities $M_{0f}$, $M_{nf}$ in formula (10) are the radial parts of the matrix elements of the transitions $0 \to f$, $n \to f$ from the principal and impurity states. The structural enhancement takes place if

$$\left|\frac{M_{nf}}{M_{0f}}\right| \gg 1. \tag{12}$$

The accuracy of modern calculations of nuclear matrix elements of radiative transitions that are not suppressed by special (model) selection rules does not exceed 10–30% (in amplitude). This means, in particular, that if a certain model of the nucleus gives $M_{0f} = 0$ and $M_{nf} \neq 0$, then we can actually expect only that

$$\left|\frac{M_{nf}}{M_{0f}}\right| \geqslant 10. \tag{13}$$

The estimate (13) can be improved if $M_{0f}$ of the suppressed "principal" transition $0 \to f$ is known from experiment. Precisely such a situation is known for the electromagnetic transition $5/2^+ \to 7/2^+$ with energy 482 keV in the $Ta^{181}$ nucleus. The lifetime of the $5/2^+$ level and the multipole composition of the radiation (97%(E2) + 3%(M1) are known from experiment. If follows therefore that the absolute probability of the (M1) transition is $3\times10^6$ sec$^{-1}$, whereas the normal value for this region of energies and nuclei is $5\times10^2$ sec$^{-1}$. Thus, the considered (M1) transition is suppressed by a factor of approximately $10^6$. The physical reason for the suppression is the fact that, according to the shell model, the $5/2^+ \to 7/2^+$ transition in the $Ta^{181}$ nucleus is coupled with the change of the orbital momentum of the nucleon by two units, as a result of which the emission of the magnetic dipole quantum by one nucleon is impossible. If the impurity transition (E1) $n \to f\,(5/2^{\pm} \to 7/2^{\pm})$ is not suppressed, then it should have a probability of the order of $10^{15}$–$10^{14}$ sec$^{-1}$. This means that in the given case

$$\left|\frac{M_{nf}}{M_{0f}}\right| \cong 10^4-10^3. \tag{14}$$

The factor (14) contains also the kinematic enhancement ($\approx 10$). Thus the structural enhancement in this example is of the order of $10^3$ - $10^2$.

In order to proceed from (14) to the estimate of the coefficient R, it is necessary to know the amplitude $\alpha_{0n}$. At the present state of our knowledge, it is hardly possible to indicate a definite order of magnitude of $\alpha_{0n}$. The point is that the nuclear wave functions used in the concrete calculations describe the states not of nucleons but of quasiparticles (elementary excitations of the nuclear Fermi liquid). The latter, when moving in the self-consistent field, are scattered by one another. The quasiparticle scattering amplitude differs greatly from the scattering amplitude of free nucleons[5]. Both the parameters of the self-consistent field, and the scattering amplitude of the quasiparticles are introduced into nuclear theory as empirical data (the calculation of these quantities for real nucleus is among the unsolved problems). In connection with the foregoing, it should be clear that even if the Hamiltonian (3) of the weak interaction of the nucleons were known, its use for the calculation of nuclear effects would be an exceedingly complicated problem. Should one assume that the parity-nonconserving forces enter in the self-consistent nuclear field? Do they appear in the quasiparticle scattering amplitude? We still have no answers to these questions.

Inasmuch as quasiparticle wave functions are used for the description of the nuclear states, the operator $H_W$ in formula (11) should be replaced by the effective Hamiltonian $\overline{H}_W$, which consists of two parts:

$$\overline{H}_W=H_W^{(S)}+H_W^{(Q)}, \tag{15}$$

where the first term $H_W^{(S)}$ is the parity-nonconserving self-consistent field, and the second $H_W^{(Q)}$ is the weak interaction of the quasiparticles with one another. The general expression for $H_W^{(Q)}$ coincides with (3)—(5). The only difference lies in the potentials $V(r)$. As to $H_W^{(S)}$, the only possible form of this operator is

$$H_W^{(S)} = V^{(S)}(\mathbf{r})\,\sigma\mathbf{p}, \tag{16}$$

where $\sigma$ and $\mathbf{p}$ are the spin and momentum operators of the quasiparticles in the self-consistent field. The amplitudes $\alpha_{0n}$ will now consist of two terms:

$$\alpha_{0n} = \frac{1}{F}\frac{1}{E_0 - E_n}(\langle n | H_W^{(S)} | 0\rangle + \langle n | H_W^{(Q)} |\rangle). \tag{17}$$

Both these terms have, generally speaking, the same order of magnitude and can therefore cancel each other. This is why it is difficult to obtain a reliable estimate of the order of magnitude of the coefficient R, even in the case of $Ta^{181}$, when the structural enhancement is large and is relatively well known.

The $5/2^+ \to 7/2^+$ transition in the $Ta^{181}$ nucleus is of particular interest because experimental measurement data are available for the experimental circular polarization of the $\gamma$ radiation emitted in this transition. The magnitude of the circular polarization is determined by the formula

$$P = \frac{2}{1+q^2}RF, \tag{18}$$

where $q^2 = 41 \pm 10$ is the probability ratio of the transitions E2 and M1.

The Leningrad group (Leningrad Physico-technical Institute) obtained for P the following result[6]:

$$P = -(6 \pm 1)\cdot 10^{-6}. \tag{19}$$

If it is assumed that the amplitudes $\alpha_{0n} \cong 1$, then, in accordance with the foregoing, $|R| \cong 10^4 - 10^3$.

We then obtain from (18) and (19)

$$F \simeq 10^{-7} - 10^{-8}. \tag{20}$$

A detailed calculation of R for the considered transition in $Ta^{181}$ was performed by Wahlborn[7]. He assumed that $\overline{H}_W = H_W^{(S)}$, i.e., he took into account only the self-consistent weak field and used a single-particle model of the nonspherical nucleus (one nucleon in a nonspherical ellipsoidal well). For the radial dependence of the potential $V^{(S)}(r)$ it was assumed that

$$V^{(S)}_{(r)} = \begin{cases} V_0, & r < r_0, \\ 0, & r > r_0. \end{cases} \tag{21}$$

The natural dimensionless parameter determining the relative intensity of the weak forces is, for the given model, the quantity

$$\mathscr{F} = m r_0 V_0. \tag{22}$$

If the considered model is realistic, then we should have obviously $|\mathscr{F}| = F$. For RF we obtain the following expression:

$$RF = \mathscr{R}\mathscr{F}, \qquad \mathscr{R} = K\frac{\langle f | (\sigma\mathbf{r})\,\mathbf{r} | 0\rangle}{M_{0f}}, \tag{23}$$

where K is a known positive numerical factor*. To estimate R, experimental data were used on the internal-conversion coefficient of the M1 radiation on the K-shell of the $Ta^{181}$ atom. It is well known that the amplitude of the probability of the internal conversion is proportional to $|M_{0f}|^2$, if we confine ourselves to the first nonvanishing term of the expansion in powers of $r_0/\lambda\!\!\!^{-}$ ($\lambda\!\!\!^{-}$ - radiation wavelength). If the transition is suppressed, as in the present case, then an important role is played by terms of order $(r_0/\lambda\!\!\!^{-})^2$, which take into account the change of the wave function of the electron over the volume of the nucleus. The expression for the conversion coefficient $\beta_\mu$ has in this case the form

$$\beta_\mu = \beta_\mu^{(1)} + b\frac{\langle f | \mathbf{r}\,(\sigma\mathbf{r}) | 0\rangle}{M_{0f}}; \tag{24}$$

here $\beta_\mu^{(1)}$ (tabulated conversion coefficient) and b are known functions of the transition energy.

Substituting (24) and (23), we can obtain

$$\mathscr{R} = K\frac{\Delta\beta_\mu}{b},$$

where

$$\Delta\beta_\mu = \beta_\mu - \beta_\mu^{(1)}. \tag{25}$$

Substitution of the experimental data in (24) yields

$$\mathscr{R} = +(3.4 \pm 0.8)\cdot 10^3. \tag{26}$$

The uncertainty in $\mathscr{R}$, calculated in this manner, is determined by the experimental errors in the measurements of the conversion coefficient. If we take into account, however, the possible limits of variation of the factor K in (24) on going from the single-particle model to an allowance for the collective effects, then, according to the estimate of the author of[7], we should write

$$\mathscr{R} = +(1.5 + 6.9)\cdot 10^3. \tag{27}$$

From (18), (19), (23), and (27) we then get

$$\mathscr{F} = -(0.3 + 4)\cdot 10^{-7}. \tag{28}$$

Concluding this analysis, we emphasize that the estimate (28) is no more reliable than (20), since all the derivations were essentially based on the choice of the effective weak-interaction Hamiltonian in the form of a self-consistent field and the simplest model of the nucleus. These are precisely the circumstances which made it possible to obtain formula (24) and to use for the estimate of $\mathscr{R}$ the experimental data on the conversion coefficient.

Structural enhancement apparently takes place also for the transition $9/2^- \to 7/2^+$ with energy 396 keV in the $Lu^{175}$ nucleus (multipole composition (E1) + (M2)). For the circular polarization of the $\gamma$ radiation, the Leningrad group obtained[8]

$$P = +(4 \pm 1)\cdot 10^{-5}. \tag{29}$$

The authors propose that

$$|\mathscr{F}| = (2 + 8)\cdot 10^{-7}. \tag{30}$$

Here, as in (28), the lower limit is theoretically less reliable.

3. Dynamic enhancement. In the preceding analysis it was assumed that the impurity amplitudes are $\alpha_{0n} \lesssim 1$. This is probable for the lower states of the nuclei, but may be incorrect at high excitations: the $\alpha_{0n}$ may increase as a result of the decrease of the energy de-

*Formula (23) can be readily obtained by neglecting the spin-orbit coupling in the nuclear Hamiltonian H of the shell model [3] (to this end it is necessary to use the relationship $\mathbf{p} = i\frac{m}{\hbar}[H, \mathbf{r}]$, $M_{nf} \sim \langle f | \mathbf{r} | n\rangle$ and the complete system of functions). In [7] formula (23) was obtained without this assumption, but it was assumed that the spin-orbit part affects only the values of the energy levels, withoug influencing greatly the radial wave functions. An oscillator potential is also used.

nominators in (11). The associated enhancement of the manifestation of weak forces will be called dynamic.

Dynamic enhancement may turn out to be significant in the excitation-energy region on the order of 8--10 MeV, realized in capture of slow neutrons by nuclei. A reaction of this type is $Cd^{113}(n,\gamma)Cd^{114}$ in a beam of polarized thermal neutrons, and was investigated in a number of papers[9-11] In these investigations, they measured the angular asymmetry of the $\gamma$ radiation relative to the neutron polarization direction. The angular correlation function is

(31)

where $\theta$ is the angle between the quantum momentum and the neutron polarization direction. At 100% polarization of the neutron beam, the coefficient a for the transition $1^+ \rightarrow 0^+$ is given by the formula

$$a = 2RF, \tag{32}$$

where R is the coefficient defined in (10). Thus, in the absence of enhancement one would expect values $a \cong 10^{-7}$. In the experiments of the Moscow Group (ITEF) they obtained[9-11] the following results for the coefficient a:

$$a = \begin{cases} -(3.7 \pm 0.9)\cdot 10^{-4} & (1964), \\ -(3.5 \pm 1.2)\cdot 10^{-4} & (1967). \end{cases} \tag{33}$$

Similar investigations by Danish workers[12-14] yield

$$a = \begin{Bmatrix} -(8.4 \pm 2.8)\cdot 10^{-4} & (1965), \\ -(3.8 \pm 2.4)\cdot 10^{-4} & (1966), \\ -(2.5 \pm 2.2)\cdot 10^{-4} & (1967). \end{Bmatrix} \tag{34}$$

There are also known results of Italian experimenters[15]:

$$a = (0.216 \pm 1.13)\cdot 10^{-4} \quad (1965.) \tag{35}$$

and the older data by R. Haas and co-workers[16]:

$$a = (1.2 \pm 7.8)\cdot 10^{-4} \quad (1959.) \tag{36}$$

It can be concluded from the foregoing data at least that values $a = 10^{-4}$ cannot be excluded at the present. The enhancement coefficient $\cong 10^3$ cannot be structural in this case, since the main transition (M1) is not suppressed. The kinematic enhancement (impurity transition (E)), as always, is a factor of the order of 10. Consequently, the dynamic enhancement should yield a coefficient on the order of $10^2$. It is important to ascertain whether such values of the dynamic enhancement are possible for real nuclei.

A more or less reasonable estimate of the upper limit of the dynamic enhancement was first published by R. Blin-Stoyle[4]. In somewhat similar form, a similar result was obtained by the present author[17]. In the foregoing arguments, the initial point will be the equation (1). We assume, first, that this equality takes place not only for the entire spectrum of the states, but also for any sufficiently numerous group of levels. In other words, we assume that Eq. (1) can be replaced by

$$F \cong \left[\frac{(\operatorname{Sp} H_W^2)_{\Delta E}}{(\operatorname{Sp} H^2)_{\Delta E}}\right]^{1/2}; \tag{37}$$

here

$$(\operatorname{Sp} A)_{\Delta E} \equiv \sum_{(\Delta E)} \langle n | A | n \rangle \tag{38}$$

differs from

$$\operatorname{Sp} A = \sum_n \langle n | A | n \rangle \tag{39}$$

in that the summation is carried out not over all the levels n, but also over those levels which are located in the energy interval $\Delta E$. If the latter contains N levels, then

$$(\operatorname{Sp} H_W^2)_{\Delta E} = \sum_{n=1}^{N} \sum_m |\langle n | H_W | m \rangle|^2 \tag{40}$$

(in the internal sum, the summation extends over all the levels). If the mean distance between the levels in the interval $\Delta E$ under consideration is equal to D, and we are interested in the region near the nuclear stability boundary, then we can write

$$(\operatorname{Sp} H^2)_{\Delta E} \cong \frac{(\Delta E)^3}{3D}. \tag{41}$$

We now introduce a second assumption, namely that only one limited group of states with different parity, which belongs so to speak to a single family (e.g., an aggregate of states which is almost degenerate as a result of some dynamic symmetry), become particularly well "mixed" with one another. Families of this kind are known in nuclear physics (their existence leads to the appearance of the so-called "giant resonances". If the interval $\Delta E$ spans the indicated group of states, then (4) can be replaced by

$$(\operatorname{Sp} H_W^2)_{\Delta E} \leqslant \left(\frac{\Delta E}{D}\right)^2 \overline{|\langle n | H_W | m \rangle|^2}, \tag{42}$$

where the bar denotes averaging. We now substitute (41) and (42) in (37), and obtain

$$\bar{\alpha}_{mn} \leqslant \left(\frac{\Delta E}{D}\right)^{1/2}, \tag{43}$$

where

$$\bar{\alpha}_{mn} = \frac{1}{F} \frac{\{\overline{|\langle n | H_W | m \rangle|^2}\}^{1/2}}{D}. \tag{44}$$

For nuclei with $A \cong 100$ we have $D = 20$ eV in the excitation-energy region close to the neutron binding energy. If we assume that $\Delta E \simeq 0.1 - 1$ MeV (this is precisely the extent of the known "giant resonances"), then it follows from (43) that

$$\bar{\alpha}_{mn} \leqslant 10^2. \tag{45}$$

This very rough estimate shows that one cannot exclude the possibility of a dynamic enhancement on the order of $10^2$. Comparing (32), (33), and (45), and taking into account the kinematic enhancement (a factor on the order of 10), we find that the data of the Moscow group imply

$$F \gtrsim 10^{-7}. \tag{46}$$

## III. CONCLUSION

It follows from all the foregoing that at present it is possible to indicate more or less definitely the upper limit of the values of the nuclear factors of the enhancement of the effects of parity nonconservation in nucleon interaction.

From the theoretical point of view, we can therefore conclude from the experimental data on circular polarization of the $\gamma$ radiation of $Ta^{181}$ (Eq. (19))

$$F \gtrsim 10^{-8}. \tag{47}$$

From analogous data on $\gamma$ radiation of $Lu^{175}$ (Eq. (29)) and the angular asymmetry of the $\gamma$ rays in the reaction $Cd^{113}(n,\gamma)Cd^{114}$ on polarized neutrons (Eq. 33)) it follows that

$$F \geqslant 10^{-7}. \tag{48}$$

It must be emphasized that the experimental data at our disposal are not yet as reliable as desired: each of the three positive results (19), (29), and (33) were obtained only by one experimental group.

In conclusion we note that of particular interest for the theory would be the observation of effects of parity nonconservation in the reaction $p(n, d)\gamma$ (or in the inverse process $d(\gamma, n)p$). This reaction is of interest, first, because experimental data concerning it would yield direct information on the two-nucleon interaction $H_W$ (formula (3)) and not concerning the effective Hamiltonian $\overline{H}_W$ (formula 15)). Second, in this case it becomes possible to distinguish the contribution of the isoscalar and isotensor parts of the Hamiltonian $H_W$ from the isovector interaction (e.g., the isovector term contributes only to the angular asymmetry of the $\gamma$ radiation, while the isoscalar and isotensor terms only to the circular polarization (see[18-20])).

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Translated by J. G. Adashko