объединенный
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исследований
дубна

E2-80-555

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THE PARITY VIOLATION IN THE FEW-NUCLEON SYSTEMS AS THE QUARK-NUCLEAR EFFECT.

III. Circular Polarization of Photons at Radiative Capture of Thermal Neutrons by Protons

Submitted to "Zeitschrift für Physik".

1980

1. INTRODUCTION

In this paper we continue the development of the quark-nucleus approach to parity violation (P.V.) effects in few-nucleon systems [1,2]. The reaction of the radiative capture of thermal neutrons by protons $np \to d\gamma$, , in which there has been measured the circular polarization of photons ($P_\gamma = (-1.30 \pm 0.45) \cdot 10^{-6}$ [3]), will serve as a concrete example.

Note, within generally accepted approach to the P.V. in NN -system, one considers an exchange by a corresponding meson and at one of the vertices one takes into account the weak interaction [4,5]. The weak long-range NN-correlations in this approach are defined by a π-meson exchange, and such an introduction of the P.V. is apparently consistent with the usual description of the strong NN-interaction by the nuclear potentials, as the behaviour of the latter is well known up to the internucleon distances $R \sim 1$ fm. The only problem is to calculate the effective weak constant $f^W_{\pi NN}$ [5].

In the case, when the selection rules permit an exchange by a vector meson (the NN-interaction at small distances $R \sim \frac{1}{m_\rho} \sim 0.25$ fm), it turn out that the calculated P.V. parameters are much lesser than the experimentally observed ones* . So, one finds, e.g., $P_\gamma^{exp} = (10 \div 100)\, P_\gamma^{th}$ (see Table 1 here and §1 in [7], further I) and the asymmetry in the $\vec{p}p \to pp$ scattering A_{pp}^{exp} (45 MeV) $= (-3.2 \pm 1.1) \cdot 10^{-7} \simeq 10\, A_{pp}^{th}$ [4,5]. Thus, for the ρ, ω-exchanges, the compatibility of the usual approach [4,5] and of the potential description of the NN-interaction, which is extrapolated into the region $R < 0.5$ fm, is doubtful (see, e.g., ref. [9]).

Here we develop an approach, which is more adequate to the problem of describing the weak NN-interaction in nuclei [1,2]. We shall consider nucleons as the 3q-clusters, and np-system at small R as the 6q-object. The radial dependence of

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* However, a calculation of P_γ in $np \to d\gamma$ shows that the essential contribution is given by the 2π-exchanges [8]. The method of this calculation needs further investigation.

Table 1

Theoretical values of $P_\gamma \cdot 10^{+8}$ in $np \to d\gamma$
($P_\gamma^{exp} = (-130 \pm 45) \cdot 10^{-8}$) [1]

strong potential model \ weak theory + ρ,ω-exchange Fig.1 of I	Cabibbo	Weinberg-Salam, Sep.contr.+Non-sep.contr. (Fig.1a) (Fig.1b)	Weinberg-Salam, Sep. and Nonsep. contr. + gluon corrections
Hulthen (R_c = 0.432 fm)	2.2 [2]		
-"- (R_c = 0.561 fm)	1.8 [2]		
Hamada-Johnston	2.9 [3]	1.6+1.0=2.6 [9] 2.34 [10]	0.98 [9] 0.80 [10]
Kishi-Sawada-Watari	1.2 [3]	0.56 [10]	-2.56 [10]
Tamagaki Gaussian soft core	2.0 [3]		
Reid hard core (R_c =0.5 fm)	2.7 [4]	2.01 [10]	1.18 [10]
Reid soft core	2.3 [4-7]	3.2 [10]	0.51 [10]
RSC + PBT [8] $^1S_0, ^3S_1, ^3D_1$	1.2 [8]		
-"- $^1P_1, ^3P_0$	1.6 [8]		
-"- all states	25.8 [8]		

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6q wave functions will be restored by two phenomenological nuclear potentials (see Ref. [10], hereafter II). Let us take the weak interaction of quarks to be contact and consider 6q -states in the configurations s^6 and s^4p^2. Their contributions to the wave functions of np-system (deuteron), characterizing a probability of the fluctuon state, are calculated in II by the quark transitional-invariant shell model (TISM)[29] with the oscillator potential[14]. To calculate P_γ, we use two types of the radial wave function, the behaviour of which differs considerably at small R ("soft core"[12] and "node" at $R \sim 0.5$ fm[13], see Figs.1 and 2 in II). Thus we are convinced that our results are rather stable with respect to the magnitude of the effect and to its sign. This is due to the fact, that in the quark-cluster consideration nucleons are really-spread (on the coordinates of their quarks) and the structure of 6q wave functions is strictly fixed by the Pauli principle (see §1 in I and §3 in the paper).

The paper is organized as follows. Paragraph 2 presents a general scheme for calculating the P.V. amplitude of T_{E1}-transition and of P_γ. The obtained results are discussed in §3. The reasons for **increasing** of P_γ in the quark-nuclear approach are also explained in this paragraph. In the conclusive §4 the restrictions to the probability of the admixture of the fluctuon state in deuteron are discussed.

§2. CALCULATIONS OF T_{E1} AND P_γ

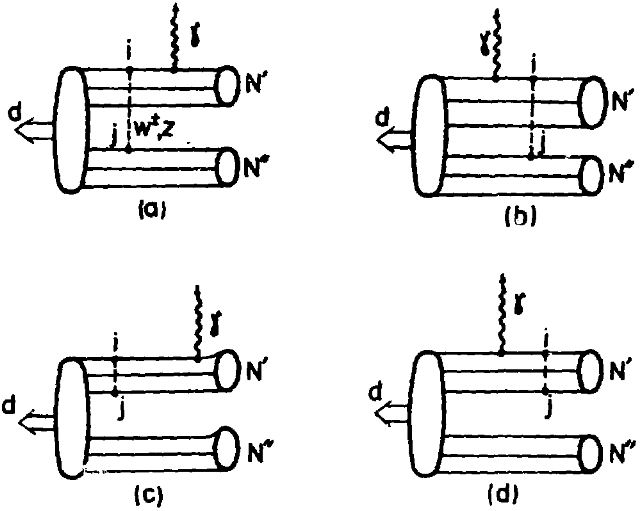
Following papers [1,2] we calculate the amplitude T_{E1} using the SEWM diagram for the direct exchanges by $W^\pm$, Z-bozons between quarks of the 6q-states (Fig.):

$$T^{(\lambda)}_{E1} = \langle \Psi_d(6q) \mid w^{(\lambda)}_{PV} \mid \Psi^{(+)}_{^1S_0}(6q) \rangle . \qquad (1)$$

Here $w^{(\lambda)}_{PV}$ is the operator taking into account the electric transition $^1S_0 - {^3S_1}$ proportional to eG_F; λ is a helicity of an emitted photon; $\Psi^{(+)}_{^1S_0}$ and Ψ_d are the quark wave functions written in the cluster approximation (see II):

$$\Psi_{NN}(6q) = \frac{1}{N_A} A \{ \Phi^{\ell}_{NN}(R)\, \Psi_{N'}(3q)\, \Psi_{N''}(3q) \}_{CST\ell} \, , \qquad (2)$$

where **A** is the antisymmetrizer in the permutations of quarks from various nucleons, N_A is the normalization factor generated by **A**.



<u>Figure</u>. The quark diagrams of the reaction $n + p \to d + \gamma$ with the parity violation.

In the spin-isospin space of n- and p-quarks, the P.V. part of the operator of the effective contact interaction of currents of the i-th and j-th quarks has the form:

$$H_{ij}^{PV} = \frac{G_F}{\sqrt{2}} \cos^2\theta_C \{[\tau_i^+\tau_j^- + \tau_i^-\tau_j^+ + \frac{1}{2\cos^2\theta_C}\tau_i^z\tau_j^z] \times$$

$$\times (\gamma_i^\mu \gamma_{\mu j}\gamma_j^5 + \gamma_i^\mu\gamma_i^5\gamma_{\mu j}) - \frac{2\sin^2\theta_W}{\cos^2\theta_C}(\tau_i^z e_j \gamma_i^\mu\gamma_i^5\gamma_{\mu j} + \tau_j^z e_i \gamma_i^\mu \gamma_{\mu j}\gamma_j^5)\}, \qquad (3)$$

H_{ij}^{PV} includes the contributions both of the charge and neutral currents of SEWM.

Introducing the Jacobi coordinates, which are used for the description of $6q$-states in II, and summing the contributions of the diagrams (<u>Fig.</u>), we may represent the amplitude as follows

$$T_{E1}^{(\lambda)} = \frac{-ih}{2m_q c} \int \Psi_{d(6q)}(\vec{x}_1, \vec{y}_1, \vec{x}_2, \vec{y}_2, \vec{R}) \sum_{i \neq j}^{6} \tilde{w}_{ij}^{(\lambda)} \hat{\delta}_{ij} \times$$

$$\times \Psi_{^1S_0}(\vec{x}_1, \vec{y}_1, \vec{x}_2, \vec{y}_2, \vec{R})\, d^3x_1 d^3y_1 d^3x_2 d^3y_2 d^3R \,. \qquad (4)$$

Here $\delta_{ij} = \delta^3(\vec{r}_i - \vec{r}_j)$ is the operator of the contact interaction of all the quarks in pairs (all in all there are $6(6-1) = 30$ terms in $\sum_{i,j}$, 18 of them when quarks i and j belong to various nucleons ($3q$ -clusters) and 12 when quarks i and j are in one nucleon); $\tilde{w}_{ij}^{(\lambda)} = \frac{1}{2}(w_{ij}^{(\lambda)} + w_{ji}^{(\lambda)})$, $w_{ij}^{(\lambda)} = \{e_q^i \gamma_i^{(\lambda)}(\gamma_i^0 - \gamma_i^3), H_{ij}^{PV}\}_+$, where $\gamma_i^{(\lambda)} = \gamma_i^\mu \epsilon_\mu^{(\lambda)*} = (0, -\frac{1}{\sqrt{2}}\gamma^1, \pm\frac{1}{\sqrt{2}}\gamma^2, 0)$, $\lambda = \pm 1$, the curly brackets $\{\ \}_+$ denote the anticommutator. The nonrelativistic reduction of the matrix elements in expr. (4) allows one to rewrite the operator $w_{ij}^{(\lambda)}$ in the following form:

$$\tilde{w}_{ij}^{(\lambda)} = G_F\sqrt{2}\cos^2\theta_C(\tau_i^- \tau_j^+ - \tau_i^+ \tau_j^-)(\sigma_i^\lambda - \sigma_j^\lambda) +$$

$$+ \sqrt{2}\left(\frac{1}{4} - \frac{5}{9}\sin^2\theta_W\right)(\tau_j^z - \tau_i^z)(\sigma_i^z \sigma_j^\lambda - \sigma_i^\lambda \sigma_j^z), \tag{5}$$

where

$$\sigma^\pm = \frac{1}{2}(\sigma^x \pm i\sigma^y), \quad \tau^\pm = \frac{1}{2}(\tau^x \pm i\tau^y).$$

Now we use the expansion of the $6q$ -state of deuteron over the basis functions of the quark TISM (see II):

$$\Psi_d(6q) = \sum_{n=0,2,\ldots} |n\rangle_{TISM} \langle n|\Psi_d(6q)\rangle, \tag{6}$$

where n is the number of excitation quanta.

As is explained in II, we may leave in $\sum_n$ the low-lying states $n = 0,2$ only with configurations $s^6[6]^x[2^3]^{CS}$ and $s^4p^2[42]^x[42]^{CS}$. In view of the symmetry of the operator in ij, the antisymmetrizer of the shell function Ψ_n will also act as an antisymmetrizing projector in the initial state. Thus, in $\Psi_{1S_0}^{(+)}$ it is sufficient to take into account the change of the normalization arising due to the antisymmetrization, i.e.:

$$\langle\Psi_n|\tilde{w}_{PV}^{(\lambda)}|\frac{1}{N_A}A\{\Phi_{1S_0}^{(+)}\Psi_N, \Psi_N, \ldots\}_{CST}\rangle = \frac{\langle\Psi_n|\tilde{w}_{PV}^{(\lambda)}|\{\Phi_{1S_0}^{(+)}\Psi_N, \Psi_N, \ldots\}_{CST}\rangle}{\sqrt{\langle\{\Phi_{1S_0}^{(+)}\Psi_N, \Psi_N, \ldots|A|\Phi_{1S_0}^{(+)}\Psi_N, \Psi_N, \ldots\}_{CST}\rangle}} \tag{7}$$

Note, that the operator $\tilde{w}_{PV}^{(\lambda)}$ can be represented in the form

$$\tilde{w}^{(\lambda)}_{PV} = \frac{eG_F}{2\sqrt{2}} [18\delta(\vec{r}_3 - \vec{r}_6)\tilde{w}^{(\lambda)}_{36} + 12\delta(\vec{r}_1 - \vec{r}_2)\tilde{w}^{(\lambda)}_{12}]. \tag{8}$$

Here the numbers "18" and "12" are the combinatorial factors, with which the contributions of the diagrams, of Fig. (a,b) (the pairs ij belong to various nucleons (3q -clusters)) and of the diagrams of Fig. (c,d) enter (the pairs ij belong to one nucleon); $m_q = \frac{1}{3}M_N$, where m_q is the quark mass of the nonrelativistic oscillator model, θ_C is the Cabibbo angle; $(\cos^2\theta_C \simeq 1)$, θ_W is the Weinberg angle ($\sin^2\theta_W =$ (0.2 ÷ 0,3). The matrix elements $\tilde{w}^{(\lambda)}_{ij}$ have opposite signs for the radiation of photons of different helicity $\lambda = \pm 1$; this gives through the interference with the regular amplitude $T^{(\lambda)}_{M1}$ of the magnetic transition the circular polarization, e.g., at $\cos^2\theta_C = 1$:

$$\begin{aligned} &\langle s_{ij} = 1 \;\; m_{ij} = \pm 1 \,|\langle T_{ij} = 1 \;\; T_{zij} = 0 |\, \tilde{w}^{(\mp)}_{ij} \,| s_{ij} = 0\rangle | T_{ij} = 0\rangle = \\ &= \mp \frac{22}{9} \mp \frac{20}{9}(\sin^2\theta_W - \frac{1}{4}), \\ &\langle s_{ij} = 1 \;\; m_{ij} = \pm 1 | \langle T_{ij} = 0 \,|\, \tilde{w}^{(\mp)}_{ij} | s_{ij} = 0\rangle \times \\ &\times | T_{ij} = 1 \;\; T_{zij} = 0\rangle = \pm \frac{14}{9} \mp \frac{20}{9}(\sin^2\theta_W - \frac{1}{4}). \end{aligned} \tag{9}$$

Leaving in the expansion (6) the low-lying states only, we rewrite $T^{(\lambda)}_{E1}$ in the form

$$\begin{aligned} T^{(\lambda)}_{E1} &= -i\frac{eh}{2m_q c}\,\frac{G_F}{\sqrt{2}}[18C_{s^6}\,\tilde{T}^{(\lambda)s^6}_{E1(36)} + \\ &+ C_{s^4p^2}(18\,\tilde{T}^{(\lambda)s^4p^2}_{E1(36)} + 12\tilde{T}^{(\lambda)s^4p^2}_{E1(12)})], \end{aligned} \tag{10}$$

where C_n are the corresponding amplitudes of the fluction states in deuteron (see II), and the amplitudes $\tilde{T}_{E1}$ are defined as follows:

$$\tilde{T}^{(\lambda)s^6}_{E1(36)} = \langle \Psi^{M=1}_{s^6}(6q) | \tilde{w}^{(\lambda)}_{36}\,\hat{\delta}_{36}\,| \frac{1}{N_A}\Phi^{(+)}_{^1S_0}\{\Psi_N, \Psi_{N''}\}_{\substack{S=0,T=1\\T_z=0}}\rangle . \tag{11}$$

To understand how our method works, it is useful to analyze these matrix elements by reducing the $6q$ -states into $3q$ -subsystems. Then separating the pairs of quarks (ij) by means of the two particle fractional parentage coefficients (see Tables in I) and using the results (9), we obtain for the contact interaction of two quarks from the same $3q$ -cluster:

$$<\Psi_N|\tilde{w}^{(-)}_{12}\hat{\delta}_{12}|\Psi_N> = 0,$$

$$<\Psi^{S_z=1/2}_{\tilde{N}**}|\tilde{w}^{(-)}_{12}\hat{\delta}_{12}|\Psi^{S_z=-1/2}_{p}> =$$

$$= -<\Psi^{S_z=1/2}_{\tilde{N}**}|\tilde{w}^{(-)}_{12}\hat{\delta}_{12}|\Psi^{S_z=-1/2}_{n}> = \frac{1}{18}[-\frac{22}{9}-\frac{20}{9}(\sin^2\theta_W-\frac{1}{4})](\frac{3\Omega}{2\pi})^{3/2}, \quad (12)$$

$$<\Psi_{N**}|\tilde{w}^{(-)}_{12}\hat{\delta}_{12}|\Psi_N> = 0,$$

Thus, we see that the diagonal $E1$ -transitions are impossible. The nonzero matrix elements correspond to the transitions to the twice p -excited nucleon-like states of the Ropper-resonance type belonging to the 70-plet of the old SU_6 -symmetry. Such states $\Psi_{\tilde{N}**}$ appear in the fractional parentage expansion of the configuration s^4p^2 only, but they are absent in the configuration s^6. Therefore, in paper /1/ the matrix element (12) was not taken into account. As we have found, the contribution of (12) gives almost one half of P_γ.

Finally, we have the following axpressions:

$$\tilde{T}^{(-)s^4p^2}_{E1(12)} = \frac{1}{N_A}\frac{\sqrt{2}}{360}[\frac{22}{9}+\frac{20}{9}(\sin^2\theta_W-\frac{1}{4})](\frac{3\Omega}{2\pi})^{3/2} \times$$

$$\times\int\Phi_{00}(R)\,\Phi^{(+)}_{{}^1S_0}(R)\,d^3R\,;$$

$$\tilde{T}^{(-)s^6}_{E1(12)} = 0;$$

$$N_A \cong \sqrt{\frac{6!}{3!\,3!2}} = \sqrt{10}$$

$$\tilde{T}^{(-)s^6}_{E1(36)} = \frac{1}{N_A}\frac{35\sqrt{2}}{1458}[1-\frac{10}{7}(\sin^2\theta_W-\frac{1}{4})](\frac{3\Omega}{2\pi})^{3/4}\int\Phi^2_{00}(R)\,\Phi^{(+)}_{{}^1S_0}(R)\,d^3R\,;$$

$$\tilde{T}^{(-)s^4p^2}_{E1(36)} = \frac{1}{N_A}\ \frac{23\sqrt{2}}{14580}\left[1-\frac{20}{69}\left(\sin^2\theta_W-\frac{1}{4}\right)\right]\left(\frac{3\Omega}{2\pi}\right)^{3/4} \times$$

$$\times \int \Phi_{00}(R)\,[\Phi_{20}(R)-\sqrt{6}\,\Phi_{00}(R)]\,\Phi^{(+)}_{^1S_0}(R)\,d^3R\ +$$

$$+\frac{1}{N_A}\ \frac{35\sqrt{2}}{1458}\left[\frac{17}{70}\sqrt{\frac{2}{3}}\left(\sin^2\theta_W-\frac{1}{4}\right)\right]\left(\frac{3\Omega}{2\pi}\right)^{3/4}\ \int \Phi^2_{00}(R)\,\Phi^{(+)}_{^1S_0}(R)\,d^3R. \qquad (13)$$

Here the numerical factors after $1/N_A$ are the result of the averaging of the matrix elements (9) over the fractional parentage expansion; $\Phi_{00}(R)$ and $\Phi_{20}(R)$ are the oscillator functions of the quark TISM for the 0s - and 2s -states on the relative coordinate R (their forms are given in II). Using the usual approximate expression for P_γ (see, e.g.,[14]) and the standard nonrelativistic amplitude of the magnetic transition $^1S_0 - {}^3S_1$, we can write following formula

Table 2

The fluctuon amplitudes and P_γ in two models of NN-interaction at small R : "core" and "node".

	core(RSC)			node(FSM)		
Ω (fm^{-2})	6	7,56	9	6	7,56	9
$r_0=\Omega^{-1/2}$ (fm)	0,408	0,364	0,333	0,408	0,364	0,333
$C(s^6)$	0,084	0,063	0,052	0.052	≈ 0	-0,008
$C(s^4p^2)$	.098	0,075	0,060	0,122	0,141	0,116
$\hat{T}^{(-)\,s^6}_{E1(36)}$ ($fm^{-3/2}$)	[illegible]76	1,696	1,788	-1,965	-3,613	-5,288
$\hat{T}^{(-)\,s^4p^2}_{E1(36)}$ ($fm^{-3/2}$)	-0,290	-0,313	-0,331	0,078	0,273	0,483
$\hat{T}^{(-)\,s^4p^2}_{E1(42)}$ ($fm^{-3/2}$)	1,660	1,812	1,928	0,492	≈ 0	-0,833
$\frac{i}{\hbar c\,\mu_N}\,T^{(-)}_{E1}$ ($fm^{1/2}$)	$3,58\cdot 10^{-6}$	$2,93\ 10^{-6}$	$2,53\ 10^{-6}$		$0,65\ 10^{-6}$	
$\frac{1}{\hbar c\,\mu_N}\,T^{(-)}_{M1}$ ($fm^{1/2}$)	15,26	15,26	15,26		15,26	
P_γ	$-4,70\ 10^{-7}$	$-3,84\ 10^{-7}$	$-3,32\ 10^{-7}$		$-0,85\ 10^{-7}$	

$$
\frac{\dfrac{d\sigma^{(+)}}{d\Omega}-\dfrac{d\sigma^{(-)}}{d\Omega}}{\dfrac{d\sigma^{(+)}}{d\Omega}+\dfrac{d\sigma^{(-)}}{d\Omega}} \cong -2\,\frac{T^{(-)}_{E1}}{T^{(-)}_{M1}} = -\,\frac{2G_F}{\sqrt{2}}\,\frac{\dfrac{eh}{2m_q c}}{\dfrac{eh}{2M_N c}}\;\times
$$

$$
\times\;\frac{18C_{s^6}\tilde{T}^{(-)s^6}_{E1(36)}+18C_{s^4p^2}\tilde{T}^{(-)s^4p^2}_{E1(36)}+12C_{s^4p^2}\tilde{T}^{(-)s^4p^2}_{E1(12)}}{E_d(\mu_p-\mu_n)\int\Phi_d(R)\,\Phi^{(+)}_{^1S_0}(R)\,d^3R}, \tag{14}
$$

where $m_q = \frac{1}{3}M_N$, , $E_d = 2.22$ MeV, $\mu_p = 2.79$, , $\mu_n = -1.91$. The orbital integrals in (13) have been evaluated for two models of NN-interaction: "core" [12] and "node" [13] (the expressions of the corresponding wave functions are given in Figs.1 and 2 in II). The final results are collected in

Table 2. The stability of the results is checked with respect to variations of values $\sin^2\theta_W$ and Ω. . The latter parameter is the single free dynamical one of our approach; it defines the space spreading of the basis functions of the quark TISM (the radius of the quark oscillator).

We pick out the value $\Omega^{-1/2} = R =$ 0.36 fm since it has been fitted by the electromagnetic nucleon form-factors in a large interval $0 \le q^2 \le 25$ GeV^2/c^2 [11], the proper electromagnetic form-factor of quark $F_q(q^2)$ being added, which takes into account effectively the contribution of ρ-meson to the nucleon form factor:

$$
F_q(q^2) = \frac{1}{1+q^2/m_\rho^2}, \quad m_\rho = 0.78\ \mathrm{GeV}. \tag{15}
$$

In particular, a half of the nucleon radius is contributed by the ρ-meson cloud. We consider the valent quarks of the nucleon rest with $R = 0.36$ fm only. Just this $6q$-state of TISM is used by us as the most physical one.

We have seen also that the dependence of our results on the Weinberg angle is very weak. The variation of $\sin^2\theta_W$ from 0.2 to 0.3 changes P_γ from $3.88\cdot10^{-7}$ to $3.80\cdot10^{-7}$ in the "core" model and from $0.86\cdot10^{-7}$ to $0.84\cdot10^{-7}$ in the "node" model. Such slight changes are caused by the cancellations of various contributions, each changing consuderably.

We may point out two reasons of some more increasing in P_γ within the framework of our approach. First, we may take into account $F_q(q^2)$ that will obviously be equivalent to the introduction of a direct interaction of ρ-meson with quarks ([4], V.D. and V.Z.). This has to enlarge the values of P_γ due to the larger overlapping of the wave functions. Second, if there existed a dibaryon with $M_B \approx 2M_N$, the expression for P_γ would have contained the increasing resonance factor. Then, the dynamical calculations of the $6q$-component would be necessary within the method of resonating group (see discussion in II). This would have established the relation of the parameters Ω and R with the spectrum of baryons and dibaryons. In particular, it allows one to find relative weights of s^6 and s^4p^2 configurations. The calculations of such a type have been undertaken already[19,28].

§3. DISCUSSIONS OF THE RESULTS

First of all let us note the differences of our method of calculation of the P.V. effect from the generally accepted approach[4]: i) we have taken into account the contribution of the direct $W^\pm, Z$-exchange between the nucleons, which is not reduced to any weak correction in the vertex of the meson-exchange[4]; ii) the direct $W^\pm, Z$-exchange is evaluated within the quark model taking into account the identity of quarks from different nucleons of the system. The latter clarifies the difference of our results from the calculation of [15], in which the direct $W^\pm, Z$-exchange is considered in the framework of NN-approach with the phenomenological nucleon form factors. In our approach the region of small NN-distances is "filled" by the really spread in the space (on the Jacobi coordinates of quarks) $6q$ Fermi system, the structure of which is essentially determined by the generalized Pauli principle. Therefore, the dependence of our results on the input data as the form of the NN-wave function $\Phi_{NN}(R)$ is weakened. Actually, the wave function $\Psi_{NN}(6q)$ in the region $R < R_c$ depends not only on the form of $\Phi_{NN}(R)$, but on the functions $\Psi_N(3q)$; due to the antisymmetrization of quarks from different nucleons, there arises additional dependence on the relative coordinate R. Hence, our results for P_γ in the "core" and "node" models are similar ($-4\cdot10^{-7}$ and $-1\cdot10^{-7}$ under the optimal choice of the parameter $\Omega = 7.56\ fm^{-2}$), in spite of the strong difference in the behaviour of $\Phi_{NN}(R)$ at $R < R_c$ (in the "node" model nucleons penetrate freely into each other and a fixed position of the node at the point $R_n \cong 0.4$ fm, is used in the NN-scattering problem as a boundary condition[13]). Converse as has been mentioned in I, in the approach [16] the amplitude T_{E1} is determined by the values of NN-wave functions $\Phi_{NN}(R)$ and their first derivatives at the point $R = 0$, where they are not defined from a physical point of view and depend badly on the choice of the model.

The calculations with ρ, ω-exchanges [4] contain a less rough but analogical arbitrariness.

Now we consider the machinery of our calculations. In the quark-nuclear approach the magnitude of P_γ is a sum of various terms with the different signs, and therefore the differences in the spin-isospin structure of different 6q configurations may influence the final value of P_γ. In particular, the configuration s^4p^2 participates in the weak interaction in a different manner in the "core" and "node" models. Let us consider, for example, the contribution of $W^\pm$, Z-exchanges in side one of the 3q-clusters (the Figure (c,d)). As we are convinced (see (12)), in this case the matrix element is not diagonal with respect to nucleon wave functions and is different from zero for the transit in the orbital excited 3q-cluster with the configuration $sp^2[21]^x$ only. Such a 3q-cluster may be separated of the wave function with the configuration s^4p^2 in the form of a term $\Gamma\Phi_{00}(R)\Psi_{sp^2}(3q)\Psi_{s^3}(3q)$, , where Γ is a fractional parentage coefficient of the expansion $q^6 \to q^3 \times q^3$ (see Table 1 in I). Just the function $\Phi_{00}(R)$ determines the magnitude of the matrix elements (12), since the overlapping integral $\int \Phi_{00}(R)\,\Phi^{(+)}_{NN}(R)\,d^3R$ enters as a factor into its expression. Hence, it follows immediately, that in the model with the node function $\Phi^{(+)}_{NN}(R)$, , which is orthogonal to the nodeless oscillator 0s-state by the definition [13], the given contribution is equal to zero. However, it is this contribution that is essential in the "core" model and it defines almost one half of P_γ. For the same reason the overlapping coefficient with the configuration s^6 in the "node" model vanishes ($C^{node}_{s^6} = 0$). Really, in the fractional parentage expansion $q^6 \to q^3 \times q^3$ of the s^6 configuration all terms are proportional to the function $\Phi_{00}(R)$, the overlapping of which with $\Phi^{node}_{d}(R)$ is equal to zero.

The requirement of the orthogonality to the oscillator 0s-state in the "node" model has been formulated in paper [13] under the assumption, that the contribution of the configuration s^4p^2 dominates at small distances in the NN-system, whereas the contribution of the configuration s^6 is neglected. Now we see, the P.V. effect is sensitive to keen details of the wave function, whereas the data on the NN-scattering may be described equally well in the "core" and "node" models. Unfortunately, the data on P_γ in the "node" model are unstable with respect to variations of the oscillator parameter Ω (in the last row of Table 2, we see $P_\gamma > 0$ for $\Omega = 6\ fm^{-2}$!). On the contrary, the "core" model results are rather stable with respect to variations Ω (with changing from 9 fm^{-2} to 6 fm^{-2} the value of P_γ changes from $-3.3 \cdot 10^{-7}$ to $-4.7 \cdot 10^{-7}$ only!).

§4. ABOUT WEIGHT OF FLUCTUON IN DEUTERON

It is interesting to compare our values of the admixture of the 6q-component in the deuteron with other estimations [16-20]. Let us write conventionally

$$|d> = \alpha|NN> + \beta|6>, \tag{16}$$

where according to our work II, the weight of the 6q-component β^2 is assumed to be defined by the expression

$$\beta|6q> = C_{s^6}\Psi_{s^6}(6q) + C_{s^4p^2}\Psi_{s^4p^2}(6q) + \ldots$$
$$\alpha = 1, \quad \beta^2 = |C_{s^6}|^2 + |C_{s^4p^2}|^2 + \ldots \tag{17}$$

Table 2 shows that $\beta^2 \simeq 1\%$ and $\beta^2 \simeq 2\%$. These rather moderate values of β^2 for the deuteron give $P_\gamma = -(1+4) \cdot 10^{-7}$, i.e., the results tend to the experimental one $P_\gamma^{exp} = (-13 \pm 4.5) \cdot 10^{-7}$ [3] and have the right sign.

It should be emphasized, after [5], that although P_γ^{exp} is anomalously high, it does not contradict P.V. data on other processes. Nevertheless the experimental errors in [3] are great, and it would be desirable to make them more precise P_γ^{exp}. It is also desirable to have the significant value of the angular asymmetry α in $np \to d\gamma$ and in the scattering $\vec{p}p \to pp$ (it is known that: $\alpha_\gamma = (0.6 \pm 2.1) \cdot 10^{-7}$ [20], α_{pp} (15 MeV) $= (-1.7 \pm 0.8) \cdot 10^{-7}$ [5], α_{pp} (45 MeV) $= (-3.2 \pm 1.1) \cdot 10^{-7}$ [8]). Perhaps, we shall succeed in observing the contributions of 9q-fluctuons by precise measurements of P.V. parameters in $\vec{p}(\vec{n})d \to p(n)d$, , $\vec{n}d \to T\gamma$ and so on (it is known that: $\alpha_{pd} = (0.35 \pm 0.85) \cdot 10^{-7}$ [21], $\alpha_\gamma = (5.0 \pm 2.5) \cdot 10^{-6}$ [22]). Such a complex of experiments will help to decide the problem of weak NN-interactions and will provide serious restrictions to the probabilities of 6q- and 9q-fluctuons in nuclei; this is very important also for understanding the fluctuon mechanism of the scattering with high transverse transfers p_T.

The estimates of the weight of the 6q-component in deuteron from processes with strong interaction give much higher probabilities β^2. So, a number of extrapolations of formulas for the electromagnetic form factors, obtained by the quark counting rules, to the region of small q^2, predict the 6q-admixture in the deuteron up to $\sim 6\%$ /16/; the estimations on the tunneling of deuteron from the state $|np\rangle$ to 6q-state in the MIT-bag give the values of β^2 from 5% to 13% /23/. Close values arise in explaining the cumulative effect /24/ within the fluctuon approach /25/ (note, other approaches are possible here, e.g., a large contribution can be given by the space-time "gathering" mechanism /26/).

In conclusion let us emphasize once more that the P.V. effects on few nucleon reactions are very sensitive to very important details of NN- and qq-interactions, e.g., to the configuration content of fluctuons /27/. Therefore, the weak interactions are the natural analyzers of strong interaction at the oversmall distances in nuclei /2,27/.

ASKNOWLEDGEMENTS

The authors are grateful to I.V. Chuvilo, S.B.Gerasimov, A.B.Govorkov, A.P.Kobushkin, V.G.Neudatchin, Yu.F.Smirnov, J.J. de Swart, R.Tamagaki, Yu.M.Tchuvilsky, V.S.Zamiralov for the discussions of various problems.

One of the authors (V.D.) would like to reminisce about the conversations with D.I.Blokhintsev in the early stage of the development of the quark-nuclear approach to P.V. effects.

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