+

ОБЪЕДИНЕННЫЙ
ИНСТИТУТ
ЯДЕРНЫХ
ИССЛЕДОВАНИЙ
ДУБНА

E2-80-554

V.M.Dubovik, I.T.Obukhovsky

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

II. Blokhintsev Fluctuon in the Deuteron within the Cluster Approximation

Submitted to "Zeitschrift für Physik, C".

1980

1. INTRODUCTION

In recent year owing to the development of the dynamical quark models constructed on the basis of QCD, the Blokhintsev fluctuons [1] (the fluctuations of the nuclear density) have been realized as multiquark states. Among the latter the 6q-states are most simple and important. The study of NN-states will obviously provide the new understanding of the origin of the so-called "repulsive nucleon core", i.e., of the effect arising in the region of small internucleon distances.

The behaviour of the NN-wave function in this region is essential for describing many intensively explored phenomena: electromagnetic form factors of the proton, deuteron, tritium [2], NN- and N-nucleus scattering with high p_T [3], the cumulative nuclear effect [4], etc.

In the nonrelativistic nuclear reactions the region of the small R_{NN} is very essential also for the consideration of parity violation effects [5] as the weak interactions have a contact character.

Thus, the understanding of the dynamical nature of the core is the daily necessity. May the core be explained by the meson exchanges between nucleons only, i.e., on the particle level, or is it a specific quark phenomenon, arising due to the existence of the multiquark states? This question has been raised in the papers [6-8].

Indeed, the consideration of the NN-systems into the confinement region $R \lesssim 1$ fm requires the quark approach, as all 6q may "fall into one bag". In the colour model the antisymmetrization of the whole wave function of 6q-system can be concentrated in the space CST (Colour, Spin, Isospin). Then, the orbital part of the wave function is symmetric $[f]_X = [6]$ and all 6q may be in s-state.

Nevertheless, s^6 configuration has to be hindered in nuclei, as the large probability of the adhesion of nucleons is inadmissible in view of the nuclear matter stability. This follows also from the data on the low-energy NN-scattering, of the deuteron parameters, which demand the introduction of the repulsive core in NN-potential. The applications of the mass of $\Delta\Delta$-systems taking into account the QCD-forces give evidence for the potential barrier with the height of a few hundreds MeV [9], preventing the mutual penetration of nucleons.

However, besides s^6 there is a low-energy state of the configuration s^4p^2 due to the colour magnetic forces of QCD [8] (§2). Its wave function on the relative coordinate of 3q-clusters $\vec{R}=\frac{1}{3}(\vec{r}_1+\vec{r}_2+\vec{r}_3)-\frac{1}{3}(\vec{r}_4+\vec{r}_5+\vec{r}_6)$ has a node in the region $R \simeq (0.4 \div 0.5)$ fm, which imitates the core, i.e., such a wave function describes the data on the NN-scattering and the electromagnetic deuteron properties [6,7] in the same way, as the wave function with the core. At the same time, it gives a higher probability for producing the 6q-states (§3). In this case the effect of repulsion in the NN-systems arises as a consequence of the general Pauli principle. The CST-parts of the wave functions in the s^6 and s^4p^2 configurations differ in form, that is connected with their different permutation symmetry in the orbital space. For example, in the s^4p^2 -configuration the orbital Young scheme may be $[42]^X$; then in CST-space we have the conjugated Young scheme $[2^21^2]^{CST}$ (in the s^6 configuration there are $[6]^X$ and $[1^6]^{CST}$, respectively). This difference will be manifested in the processes depending on spins and isospins of quarks, e.g., in the weak interaction effects and in electromagnetic form factors. As a result for the calculation of such effects, it is necessary to know both the total probability of the admixture of the 6q -component in a deuteron and the relative weights of different configurations and their phase factors. Such information about quark wave functions can be provided by a full dynamical calculation using the pair qq-interaction, that is, impracticable today.

In this paper as a rough estimation of the probabilities of the 6q-components we suggest using the phenomenological NN-wave functions for describing the relative motion of the centre masses of two 3q-clusters (nucleons). In that case using the usual methods of cluster nuclear physics [10] we can calculate the probabilities and the phases of any 6q-configurations. This approach has first been used to the deuteron [11] for the calculation of the probability of the $\Delta\Delta$-component determined by the configuration s^6. Like in [11] we use here the oscillator basis of the translational invariant shell model (TISM) [10] for the 3q- and 6q-states. Thus, the quark dynamics is taken into account only effectively through the oscillator parameter of the basis $\Omega = 7.56$ fm^{-2}, which has been fitted in [12] by the data on the nucleon form factors. On the other hand, the phenomenological NN-wave functions permit the correct description of the NN-scattering. We use two alternative phenomenological models of NN-interactions: a model with the repulsive core [13] and a model with the node wave function [6]. In conclusion we discuss the correspondence of our fluctuon probabilities with the earlier results [11,18-20] and with experimental restrictions.

2. REPRESENTATION OF THE NN-STATES IN TERMS OF THE 6q -QUARK SHELL BASIS

In the region of overlapping of the nucleon wave functions the identity of quarks is very important, as the quarks are the fermions evidently. The identity of fermion quarks from different nucleons implies a full antisymmetrization of $\Psi_{NN}(6q)$ with respect to any quark permutation.

Let $A=\frac{1}{10}[I-\sum_{i=1}^{3}\sum_{j=4}^{6}P_{ij}]\,(A^2=A)$ be the antisymmetrizer of quarks from different nucleons. In the rough approximation, we can write:

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

where $\Phi^{\ell}_{NN}(R)$ is the orbital part of the usual nuclear wave function (of the deuteron or the NN-scattering state) and

$$\Psi_{NN}(pn)=\Phi^{\ell}_{NN}(R)\,\{\chi^{S'T'}_{N'}\,\chi^{S''T''}_{N''}\}_{ST}\ , \qquad (2)$$

where χ^{ST}_{N} is the spin-isospin parts. Here we substitute the quark wave functions of the nucleons constructed in the quark TISM [10]:

$$\Psi_{N'}(3q)=|0(00)0[3]^{X'},\ [1^3]^{C'}S'=\tfrac{1}{2}[21]^{CS'}T'=\tfrac{1}{2}[1^3]^{CST'}>_{TISM}\ \equiv \qquad (3)$$

$$\equiv\Phi_{00}(x_1,\tfrac{\Omega}{2})\Phi_{00}(y_1,\tfrac{2}{3}\Omega)|[1^3]^{C'}S'=\tfrac{1}{2}[21]^{CS'}T'=\tfrac{1}{2}[1^3]^{CST'}>\ ,$$

where $\Phi_{n\ell}$ are the oscillator wave functions, $\Omega^{1/2}$ being the oscillator radius of the lowest states

$$\Phi_{00}(x,\Omega)=(\frac{\Omega}{\pi})^{3/4}\,e^{-\frac{1}{2}\Omega x^2}\ , \qquad (4)$$

$$\Phi_{20}(x,\Omega)=\sqrt{\frac{2}{3}}(\frac{\Omega}{\pi})^{3/4}\,(\frac{3}{2}-\Omega x^2)\,e^{-\frac{1}{2}\Omega x^2}\ ,\ \ldots\ \text{etc},$$

x_i, y_i, R, X are the Jacobi coordinates for the two $3q$ -clusters:

$$\vec{x}_{1(2)} = \vec{r}_{1(4)} - \vec{r}_{2(5)}, \quad \vec{y}_{1(2)} = \frac{1}{2}(\vec{r}_{1(4)} + \vec{r}_{2(5)}) - \vec{r}_{3(6)}, \qquad \vec{X} = \frac{1}{6}(\vec{r}_1 + \vec{r}_2 + \dots + \vec{r}_6), \tag{5}$$

having the properties:

$$d^3r_1 d^3r_2 \dots d^3r_6 = d^3x_1 d^3y_1 d^3x_2 d^3y_2 d^3R\, d^3X, \qquad r_1^2 + r_2^2 + \dots + r_6^2 = \frac{1}{2}x_1^2 + \frac{2}{3}y_1^2 + \frac{1}{2}x_2^2 + \frac{2}{3}y_2^2 + \frac{3}{2}R^2 + 6X^2. \tag{6}$$

In the right-hand side of eq. (1) the factor N_A^{-1} takes into account the change of the normalization of the wave function under the action of the projector A:

$$N_A = \frac{\langle\{\Phi_{NN}\Psi_{N'}\Psi_{N''}\}|A|\{\Phi_{NN}\Psi_{N'}\Psi_{N''}\}\rangle}{\langle\{\Phi_{NN}\Psi_{N'}\Psi_{N''}\}|\{\Phi_{NN}\Psi_{N'}\Psi_{N''}\}\rangle}. \tag{7}$$

The curly brackets $\{\dots\}_{CST\ell}$ in the right-hand side of eq. (2) mean the vector summation of all the moments of the $3q$ -clusters into the total moments of the $6q$ -system:

$$C = C' + C'', \quad S = S' + S'', \quad T = T' + T'', \quad L = L' + L'' + \ell;$$

in the given case we have

$$C' = C'' = 0, \quad S' = S'' = \frac{1}{2}, \quad T' = T'' = \frac{1}{2}, \quad L' = L'' = 0.$$

In the first row (3) the symbols, standing before the comma, refer to the orbital part of the wave function and are standard for the s,p -shell of the TISM [10]: $|N_p'(\lambda'\mu')L'[f']^X, \dots\rangle$. Here N_p' is the number of p-quanta of excitation of the oscillator; $(\lambda'\mu')$ are the symbols of SU_3-scheme [10] in the space of oscillator p-excitations (if $[f_p']$ is the corresponding Young scheme, characterizing the permutation symmetry of p -quanta, then $\lambda' = f'_{p_1} - f'_{p_2}$, $\mu' = f'_{p_2} - f'_{p_3}$, where f'_{p_i} are the lengths of the Young scheme rows); L' is the orbital moment; $[f']^X$ is the permutation symmetry in the coordinate space.

An approximative representation of a many-particle wave function in the form (1) is used in analogous problems of the nuclear physics for a long time, in particular, in the calculation by the resonating group method (RGM) [17], suggested by

Wheeler as early as 1937 [14]. Recently, the first RGM calculations have appeared in the quark approach [15]. We should like to emphasize that our method is essentially simpler than the RGM. The RGM is the variational approach for determining the unknown function $\Phi_{NN}(R)$, while we use in eq. (5) already known functions $\Phi_{NN}(R)$ from different phenomenological models of the NN-interactions [13,6] (Figure).

Thus, constructing the full orthonormalized basis of the TISM for completely antisymmetrized oscillator 6-quark states [21] $\Psi_n(6q)$, $n=0,1,2,\ldots$ and expansing over this basis the wave function (1) of the NN-system we have

$$\Psi_{NN}(6q) = \sum_n |\Psi_n(6q)\rangle\langle\Psi_n(6q)|\frac{1}{N_A}A\{\Phi_{NN}\Psi_{N'}(3q)\Psi_{N''}(3q)\}\rangle . \tag{8}$$

In a calculation of effects sensible only to the small internucleon distances, we may take into account only the low-lying states of the TISM in the expansion (8). We use the states of the s-shell (configuration s^6) and the first excited states of the same (positive) parity (configuration s^4p^2). We consider the states with the deuteron quantum numbers only: $(S,T)=(1,0)$. In the configuration s^6 the Pauli principle fixes uniquely the only allowed state in the TISM basis:

$$\Psi_0(6q) = |0(00)0[6]^X, [2^3]^C C=0 S=1 [2^3]^{CS} T=0 [1^6]^{CST}\rangle_{TISM} . \tag{9}$$

In the case of the orbital symmetry $s^4p^2[42]^X$ the Pauli principle requires the Young scheme $[2^21^2]^{CST}$ in the CST-space, which is compatible with any scheme $[f]^{CS}$ in the expansion

$$[2^3]^C \circ [42]^S = [42]^{CS} + [321]^{CS} + [2^3]^{CS} + [31^3]^{CS} + [21^4]^{CS} . \tag{10}$$

Thus, the Pauli principle selects six TISM-states with the deuteron quantum numbers:

$$\begin{gathered}\Psi_1([42]^X[42]^{CS}),\ \Psi_2([42]^X[321]^{CS}),\ \Psi_3([42]^X,[2^3]^{CS}),\\ \Psi_4([42]^X[31^3]^{CS}),\ \Psi_5([42]^X[21^4]^{CS}),\ \Psi_6([6]^X,[2^3]^{CS}) .\end{gathered} \tag{11}$$

The degeneracy of the states vanishes due to the colour-magnetic forces H_M of QCD [16]. If the dependence of the orbital part of H_M of the quark orbital state (s or p) is emitted, we can express the mean value of H_M through the proper value of

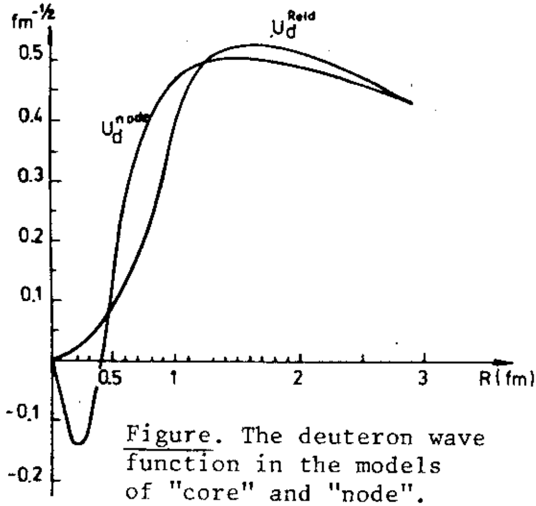

Figure. The deuteron wave function in the models of "core" and "node".

the Casimir operator $\hat{C}_2^{(6)}$ of the group $SU_{CS}(6)$ [16]

$$\Delta_n = 8N \pm \frac{4}{3} S(S+1) - \qquad (12)$$
$$-4 \langle \Psi_n | \hat{C}_2^{(6)} | \Psi_n \rangle ,$$

Δ_n depending on the Young scheme $[f]^{CS}$. As a result we find the following Table 1 of Δ_n for the vectors Ψ_n.

Thus, the state $\Psi_1(6q)$ is the most low-lying in the configuration s^4p^2. The more exact calculation, that takes into account the orbital dependence of H_M, has been made in [8] by the model of the MIT quark bag[17] and the oscillator model. They confirmed, the state Ψ_1 is peculiar in the configuration $s^4_{1/2}\ p^2_{1/2}$ and close in energy to the state $s^6_{1/2}$. Therefore, in expansion (8) we may leave the vector (9) and the vector

$$\Psi_1(6q) = |2(20)0[42]^X, [2^3]^C C=0 S=1 [42]^{CS} T=0 [2^2 1^2]^{CST} \rangle_{TISM} . \qquad (13)$$

Then we shall denote Ψ_0 (6q) as Ψ_{s^6}, and Ψ_1(6q) as $\Psi_{s^4p^2}$, and the quantities, related to Ψ_0 and Ψ_1, will be indexed s^6 and s^4p^2, respectively. For example, the overlapping coefficients of Ψ_0 and Ψ_1 with Ψ_d (6q) in (8) will be denoted by

$$C_{s^6} = \langle \Psi_{s^6} | \frac{1}{N_A} \Phi_d \{ \Psi_{N'}(3q)\, \Psi_{N''}(3q) \}_{S=1, T=0} \rangle \qquad (14a)$$

$$C_{s^4p^2} = \langle \Psi_{s^4p^2} | \frac{1}{N_A} \Phi_d \{ \Psi_{N'}(3q)\, \Psi_{N''}(3q) \}_{S=1, T=0} \rangle . \qquad (14b)$$

The antisymmetrizer A is removed from (13) because the TISM bra-vectors Ψ_n act on the ket-vectors as projection operators [21].

Table 1

The energy of the colour-magnetic interaction $E_M = \langle\Psi_n|H_M|\Psi_n\rangle$ (The data are obtained by the quark bag model [16,17,18])

n	1	2	3	4	5	6
$[f]^{CS}$	[42]	[321]	$[2^3]$	$[31^3]$	$[21^4]$	$[2^3]$
Δ_n	-88/3	-28/3	8/3	8/3	80/3	8/3
E_M(MeV)	-430	-137	39	393	393	39

The values of $|C_{s^6}|^2$ and $|C_{s^4p^2}|^2$ may be considered as the probabilities of 6q- -components with different symmetries in deuteron, the matter being in the statistical object like the Blokhintsev fluctuon[1], directly defined by the identity of nucleon constituents, rather than in a real exotic 6q-resonance.

3. THE CALCULATION OF ADMIXTURES OF s^6 AND s^4p^2 COMPONENTS ON DEUTERON

In order to calculate the admixtures of s^6 and s^4p^2 we need the expansions of 6q-states Ψ_n into sums of products of two 3q-cluster states (especially we are needed a coefficient at $\Psi_n\Psi_n$ -term in this expansion), i.e.,

$$\begin{aligned}\Psi_{s^6} &= \frac{1}{3}\Phi_{00}(R,\tfrac{3}{2}\Omega)\{\Psi_{N'}(3q)\Psi_{N''}(3q)\}_{S=1,\,T=0} + \\ &+ \sqrt{\frac{4}{45}}\Phi_{00}(R,\tfrac{3}{2}\Omega)\{\Psi_{\Delta'}(3q)\Psi_{\Delta''}(3q)\}_{S=1,\,T=0} + \\ &+ \sqrt{\frac{4}{5}}\Phi_{00}(R,\tfrac{3}{2}\Omega)\{\sum_i \Psi^{C\neq 0}_{B_i'}(3q)\Psi^{C''\neq 0}_{B_i''}(3q)\}_{C=0,\,S=1,\,T=0}\end{aligned} \qquad (15)$$

(the latter is a contribution of colour pairs);

$$\begin{aligned}\Psi_{s^4p^2} &= -\frac{1}{5}\Phi_{20}(R,\tfrac{3}{2}\Omega)\{\Psi_{N'}(3q)\Psi_{N''}(3q)\}_{S=1,\,T=0} + \\ &+ \sqrt{\frac{1}{50}}\Phi_{00}(R,\tfrac{3}{2}\Omega)\{\Psi_{\vec{N}^{**}}(3q)\Psi_{N}(3q)\}_{S=1,\,T=0} +\end{aligned}$$

$$+\frac{1}{10}\Phi_{00}(R,\frac{3}{2}\Omega)\{\Psi_N(3q)\Psi_{N^{**}}(3q)\}_{S=1,\ T=0}\quad +$$

$$+\sqrt{\frac{1}{50}}\Phi_{11}(R,\frac{3}{2}\Omega)\{Y_{1m}(\hat{R})\ \Psi_{N^*}(3q)\ \Psi_N(3q)\}_{S=1,\ T=0,\ L=0}\quad +$$

$$+\frac{3}{10}\Phi_{00}(R,\frac{3}{2}\Omega)\{\Psi_{\Delta^{*\prime}}(3q)\ \Psi_{\Delta^{*\prime\prime}}(3q)\}_{S=1,\ T=0,\ L=0}\quad +$$

$$+\sqrt{\frac{4}{5}}\{\sum_i\Phi_{n_i\ell_i}(R)\ \Psi_{B_i'}^{C\neq 0}(3q)\ \Psi_{B_i''}^{C''\neq 0}(3q)\}_{C=0,\ S=1,T=0,L=0}\qquad \text{(the same, as above).}\tag{16}$$

Here we denote: $\Psi_\Delta(3q)$ is the quark wave function of Δ-isobar, $\Psi_{N^*}(3q)$ and $\Psi_{N^{**}}(3q)$ are the orbital excitations of the nucleon state with one or two quanta, respectively. In particular, we use:

$$\Psi_{\tilde{N}^{**}}(3q)\equiv|2(20)0[21]^X,\ [1^3]^C C=0\ S=\frac{1}{2}[21]^{CS}\ T=\frac{1}{2}[21]^{CST}>_{TISM}=$$

$$=\sqrt{\frac{1}{2}}\Phi_{11}(x_1,\frac{1}{2}\Omega)\Phi_{11}(y_1,\frac{2}{3}\Omega)\{Y_{1m}(\hat{x}_1)Y_{1m}(\hat{y}_1)\}_{L=0}\ \times$$

$$\times|[1^3]^C S=\frac{1}{2}[21]^{CS}T=\frac{1}{2}[21]^{CST}r^{(1)}_{CST}>-\tag{17}$$

$$-\sqrt{\frac{1}{2}}[\sqrt{\frac{1}{2}}\Phi_{20}(x_1,\frac{1}{2}\Omega)\Phi_{00}(y_1,\frac{2}{3}\Omega)-\frac{1}{\sqrt{2}}\Phi_{00}(x_1,\frac{1}{2}\Omega)\Phi_{20}(y_1,\frac{2}{3}\Omega)]\times$$

$$\times|[1^3]^C S=\frac{1}{2}[21]^{CS}T=\frac{1}{2}[21]^{CST}r^{(2)}_{CST}>,$$

where $r^{(i)}$ are the Yamanouchi symbols, which characterize the basic vectors of irreducible representations /10/ of the permutation group:

$$\Psi_{N^{**}}(3q)=|2(20)0[3]^X,[1^3]^C C=0\,S=\frac{1}{2}[21]^{CS}T=\frac{1}{2}[1^3]^{CST}>_{TISM}\quad=$$

$$= [\sqrt{\frac{1}{2}}\Phi_{20}(x_1,\frac{1}{2}\Omega)\Phi_{00}(y_1,\frac{2}{3}\Omega)+\sqrt{\frac{1}{2}}\Phi_{00}(x_1,\frac{1}{2}\Omega)\Phi_{20}(y,\frac{2}{3}\Omega)]\cdot\times$$

$$\times |[1^3]^C S=\frac{1}{2}[21]^{CS} T=\frac{1}{2}[1^3]^{CST}\rangle . \tag{18}$$

The state $\Psi_N(3q)$ is given by (3), and other wave functions are not needed.

Using these formulas we can reduce the expression (14) to:

$$C_{s^6} = \frac{1}{3}\frac{1}{N_A}\int\Phi_{00}(R,\frac{3}{2}\Omega)\Phi_d(\vec{R})\,d^3R \tag{19a}$$

and

$$C_{s^4p^2} = -\frac{1}{5}\frac{1}{N_A}\int\Phi_{20}(R,\frac{3}{2}\Omega)\,\Phi_d(\vec{R})\,d^3R . \tag{19b}$$

The final results are obtained by the numerical integration in (19) of the wave functions (4) and the phenomenological deuteron functions of two alternative models of the NN-interactions [13,6] (Table 2).

We see, the fluctuon probabilities $|C_{s^6}|^2$ and $|C_{s^4p^2}|^2$ increase with increasing radius $\Omega^{-1/2}$ of the oscillator 6q-state. However, as we have already noted, this radius is not arbitrary; it must take into account effectively a quark dynamics. In paper [12] $\Omega = 7.56\ fm^{-2}$ was fitted by the nucleon form factors. This value gives evidently the radius $r_{3q}=\Omega^{-1/2} = 0{,}364$ fm of the "quark core" of a nucleon, while the most part of the full charge radius $r_{ch} = 0.8$ fm may arise due to a ρ-meson cloud. We emphasize, in the model of NN-interaction with the forbidden state [6], the position of the node of the wave function of the NN-state coincides practically with the radius of the quark oscillator $R^{FS}_{node} = (0.36\div 0.4)$ fm. Then, at $\Omega = 7.56\ fm^{-2}$ the orthogonality condition $(\Phi^{node}_{NN}(R), \Phi_{00}(R,\Omega_{FS}))=0$ is fulfilled also for the nodeless configuration s^6. This is the reason that magnitude of $C^{FS}_{s^6}$ is not stable with respect to the deviations of the parameter Ω. Thus, it is necessary to fix $\Omega = 7.56\ fm^{-2}$ by the physical considerations, this value corresponding to all restrictions following from the phenomenological models of NN-interactions.

4. DISCUSSION OF THE RESULTS. RELATION TO THE CORE PROBLEM

It is interesting to compare the obtained results with estimations of the admixture of 6q-components in the deute-

Table 2

The admixtures of the s^6 and s^4p^2 components in the deuteron for the Reid model with the soft core (SC) [13] and the model with the forbidden state (FS) [6]

$\Omega(\mathrm{fm}^{-2})$		2	4	6	7.56	9
$r_0 = \Omega^{-1/2}$	(fm)	0.707	0.500	0.408	0.364	0.323
C_{s^6}	SC	0.273	0.135	0.0842	0.063	0.052
	FS	0.315	0.134	0.052	0	-0.008
$C_{s^4p^2}$	SC	0.264	0.149	0.0983	0.075	0.060
	FS	0.280	0.123	0.122	0.141	0.116
$\lvert C_{s^6}\rvert^2$	SC	0.074	0.018	0.0071	0.0040	0.0027
	FS	0.099	0.018	0.0027	0	0.00006
$\lvert C_{s^4p^2}\rvert^2$	SC	0.070	0.022	0.0097	0.0056	0.0036
	FS	0.078	0.015	0.0148	0.0199	0.0135

ron [11,18-20]. We can write conventionally:

$$|d\rangle = \alpha|NN\rangle + \beta|6q\rangle , \qquad (20)$$

where

$$|6q\rangle = \frac{1}{\beta}\left[C_{s^6}\Psi_{s^6}(6q) + C_{s^4p^2}\Psi_{s^4p^2}(6q) + \dots \right],$$
$$\alpha \cong 1, \quad \beta^2 = |C_{s^6}|^2 + |C_{s^4p^2}|^2 + \dots \; . \qquad (21)$$

Let us note, the states $|NN\rangle$ and $|6q\rangle$ are not orthogonal as the NN -component contributes in the fractional parentage expansion (15), (16) of the shell-states s^6 and s^4p^2 with the coefficients $\frac{1}{3}\Phi_{00}(R)$ and $-\frac{1}{5}\Phi_{20}(R)$, respectively. However, the orbital overlapping coefficients $\gamma_{n\ell} = \int \Phi_{n\ell}(R)\,\Phi_d(R)\,d^3R$ are small, and with an accuracy up to the factors $\frac{1}{3}\gamma_0$ and $-\frac{1}{5}\gamma_{20}$ the admixture of the 6q -components can be determined as a magnitude of β^2 in the expressions (20) and (21). From Table 2 we find $\beta^2_{sc} \simeq 1\%$ and $\beta^2_{FS} \simeq 2\%$. A number of extrapolations of the "quark counting rule" formulas to the region of small q^2 predicts the 6q-admixture in the deuteron up to 6% [18]. The close values [20] arise from the estimations of the data on the cumulative processes [20]. (But in this case the fluctuon approach is not uniquely possible; see, e.g., [22]).

The question what configuration (s^6 or s^4p^2) dominates at small distances has a direct relation to the core problem. In [9] the balance of quark forces has been studied in the adiabatic approximation as a function of the parameter R. In this approach the point $R=0$ corresponds to the shell configuration $s^6[6]^X$ and the position of the level E_0 in the s^6-configuration ($E_0 = \langle\psi_0|H|\psi_0\rangle$) gives some information on the NN-forces at small distances. In the s^6-configuration the repulsive colour-magnetic forces prevail, and this result agrees well with the phenomenological NN-potentials [13], which include the repulsive core at $R_C \simeq 0.5$ fm.

However, in the overlapping region of the 3q-clusters the real situation may differ badly from the adiabatic picture. Particularly, the RGM computations urge us [15] that the wave function depending on R has the node form. The node function may correspond to the case, when at small distances in NN-system the s^4p^2 configuration dominates, as it has first been noted in [7,23]. Really in the expansion (15) of s^4p^2 state the NN-term has the nodal character (the oscillator function $\Phi_{20}(R, \frac{3}{2}\Omega)$ has a node in point $R = \Omega^{-1/2}$). There is a possibility also, that in the NN-forces the repulsive core is in fact only a reflection of the node of the wave function $\Psi_{NN}(6q)$ like in the cluster system α-α, etc. [6,24]. At small distances $\Psi_{NN}(6q)$ is rather to a superposition of the TISM basis vectors in the s^6 and s^4p^2 configurations; this is reflected in expansion (8). In the subsequent paper of our cycle we shall show that the parity violation effects restrict the fluctuon probability in deuteron and perhaps are crucial to its configuration content [23].

We are indebted to V.G.Neudatchin, Yu.F.Smirnov and Yu.M.Tchuvilsky for interesting remarks.

REFERENCES

1. Blokhintsev D.I. Sov.Phys.JETP, 1958, 6, p.995.
2. Matveev V.A., Muradyan R.M., Tavkhelidze R.M. Lett.Nuovo Cim., 1973, 7, p.719; Brodsky S.J., Farrar G.R. Phys. Rev.Lett., 1973, 31, p.1153.

4. Baldin A.M. In: Proceedings of the XIX International Conference on High Energy Physics. Tokyo, 1978, Edited by S.Homma et al., p.455. Phys.Soc. of Japan, Tokyo, 1978.
5. Dubovik V.M., Kobushkin A.P. Report ITP-78-85, Kiev, 1978; Dubovik V.M. Contribution to "Int.Symp. of Few Particles Problems in Nucl.Phys.", JINR, D4-80-271, Dubna, 1980; Dubovik V.M., Obukhovsky I.T. Extend.Contr., ibid.
6. Neudatchin V.G. et al. Phys.Rev., 1975, C11, p.128.
7. Neudatchin V.G., Smirnov Yu.F., Tamagaki R. Progr.Theor. Phys., 1977, 58, p.1-72; Smirnov Yu.F., Obukhovsky I.T., Neudatchin V.G. Yad.Fiz., 1978, 27, p.860.
8. Obukhovsky I.T. et al. Phys.Lett., 1979, B88, p.231; Yad.Fiz., 1980, 31, p.516.
9. Liberman D.A. Phys.Rev., 1977, D16, p.1542; De Tar C.B. Phys.Rev., 1978, D17, p.301.
10. Neudatchin V.G., Smirnov Yu.F., Golovanova N.F. Adv. in Nucl.Phys.,11, ed. by S.W.Negele and F.Vogt. Plenum Press, N.Y., 1979.
11. Smirnov Yu.F., Tchuvilsky Yu.M. Preprint Moscow State University, 1978.
12. Light A.L., Pagnamenta A. Phys.Rev., 1970, D2, p.1150.
13. Reid R.V., Jr. Ann.Phys., 1968, 50, p.411.
14. Wheeler J.A. Phys.Rev., 1937, 52, pp.1083, 1107.
15. Ribeiro J.E.F.T. Comm. CFMC E6/78, 1978; Toki H. Zeitschr. für Phys., 1980, A294, p.173; Oka M., Yazaki K. Phys. Lett., 1980, 90B, p.41.
16. Jaffe R.L. Phys.Rev.Lett., 1977, 38, p.195; Er.p.617.
17. Aerts A.Th.M., Mulders P.J.G., de Swart J.J. Phys.Rev., 1978, D17, p.260.
18. Matveev V.A., Sorba P. Nuovo Cim.Lett., 1977, 20, p.145; Nuovo Cim., 1978, 45A, p.357; Kizukuri Y., Namiki M., Okano K. Progr.Theor.Phys., 1979, 61, p.559.
19. Kobushkin A.P. Kiev report ITP-77-113E, 1977.
20. Burov V.V., Lukyanov V.K., Titov A.I. Phys.Lett., 1977, 67B, p.46; Lukyanov V.K., Titov A.I. Elem.Part. and Atomic Nuclei, 1979, 10, p.815.
21. Dubovik V.M., Obukhovsky I.T. JINR, P2-80-501, Dubna,1980.
22. Kalinkin B.N., Cherbu A.V., Shmonin V.L. Acta Phys.Pol., 1979, B9, pp.375,385,393.
23. Neudatchin V.G. In: Clustering Aspects of Nucl.Struct. and Nucl.Matter.Proc. of III Int.Conf., Winnipeg, 1978, eds. I.M.Van Oers et al. AIP Conf.Proc., 1978, No.47, AIP, N.Y., p.469.
24. Neudatchin V.G., Kukulin V.I., Korotkikh V.L. Phys.Lett., 1971, 34B, p.581.

Received by Publishing Department
on August 7 1980.