Exclusive process ^{2}H$(e,e'p)N^*$ as a tool for investigation of the quark structure of the deuteron

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The exclusive ^{2}H$(e,e'p)N^*$ process at CEBAF energies, with the missing mass coincidence technique, offers the opportunity to reveal the baryon-baryon (BB) composition of the deuteron and to reconstruct the deuteron six-quark wave function. As an example, we project out two different kinds of currently discussed six-quark functions into various BB channels including such baryons as n, p, Δ, $N(1440)$, $N(1520)$, and $N(1535)$. Spectroscopic factors and momentum distributions are calculated using the nonrelativistic quark model. The results are that the $\Delta\Delta$ component in the deuteron is almost entirely connected with the six-quark configuration s^6 while the channels with $N(1440)$, $N(1520)$, and $N(1535)$ display the configuration s^4p^2.

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I. INTRODUCTION

One of the dominating views of the nucleon structure is that nonperturbative QCD effects produce, first of all, the constituent mass $m_q \sim \frac{1}{3}m_N$ of an individual quark as a result of its coupling to the virtual collective vacuum excitations. Here, the nucleon quark radius equals approximately 0.5–0.6 fm.

On face value it seems that the quark degrees of freedom may be seen in both elastic and inelastic form factors of electron scattering on the deuteron providing the momentum transfer is large enough. However, the quark microscopic calculations [1,2] show that the nucleon-nucleon overlap region $r \le 0.5$ fm is characterized by a destructive interference of quark configurations s^6 and s^4p^2 which results in the shape of form factors, typical for traditional repulsive core NN potentials. So, there is only evidence that quark concepts are compatible with electron form factors, low-energy NN-scattering phase shifts, etc. [1–3].

An illuminative demonstration of the essential role of the excited quark configurations is the possibility to describe the NN-scattering data (differential cross section and vector polarizations) within a rather broad energy range $0 < E_{\rm lab} \le 5$ GeV based on deep attractive potentials with forbidden states [4]. These potentials are connected to the quark configuration s^4p^2 producing the nodal S-wave function of relative motion of two nucleons.

Of course, it is important to search for more and more convincing manifestations of quark concepts in the NN system. One of the possible ways was outlined in our preliminary communications [5,6] and is presented in this paper.

Low-energy nuclear physics deals with cluster properties of nuclei having a shell-model structure. The theory of clustering, which is presented in Refs. [7,8], leads to rather successful experiments on the quasielastic knock-out of clusters from the light nuclei by fast protons or electrons with the energies 100–300 MeV [9]. First, the spectroscopic factors $S_x^A(i,f)$ ("partial probabilities," their sum is not normalized to unity) to find cluster x in the nucleus A were measured, where i and f mean initial state of the nucleus A and final state of the nucleus $A-x$, respectively. Second, cluster momentum distributions were also measured. At higher proton bombarding energies (up to 1 GeV) and correspondingly higher energies of knocked-out clusters, where the knock-out process results from the nuclear interior instead of the nuclear surface, the amplitudes of deexcitation of virtually excited clusters during the quasielastic knock-out process (e.g., $\alpha^* \to \alpha$) can be measured and result in very unusual properties of quasielastic form factors [7,8]. One of the experimental results [9] gives the simplest evidence of such deexcitation. The essential property of high-energy quasielastic knock-out is that final-state exchange terms are of no importance due to the large values of momenta of fast particles x or $A-x$ in the final state. In the present paper we use this idea to try to make quark effects measurable.

The six-quark wave function of the deuteron, due to the essential contribution of the s^4p^2 configuration, shows very rich cluster properties. So, together with the NN, $\Delta\Delta$, and CC components as the well-known characteristics of the s^6 configuration, there are many different cluster components NN^*, NN^{**}, N^*N^*, $\Delta\Delta^*$, $\Delta\Delta^{**}$, $\Delta^*\Delta^*$, CC^*,... where one or two oscillator quanta of excitation are concentrated on the internal Jacobi coordinates of the $3q$ cluster. Here, one-quantum excitations N^* have negative parities [$N_{3/2^-}(1520)$ and $N_{1/2^-}(1535)$] while the lowest of the two-quantum states N^{**} is Roper resonance $N_{1/2^+}(1440)$, etc. It is important that the $\Delta\Delta$ component of the s^4p^2 configuration is one order of magnitude small-

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er than that of s^6.

The appearance of orbitally excited baryons in baryon-baryon (BB) composition of the deuteron was noted first in Ref. [10] where some qualitative estimates of $S^d_{\rm BB}$ factors were also done within the six-quark cluster model. Our Ref. [5] has dealt with "realistic" six-quark configurations but the simplest baryonic wave functions were used to calculate spectroscopic factors $S^d_{\rm BB}$ for BB components with $N(1440)$, $N(1520)$, or $N(1535)$ as one of baryons. However, it appears [6] that the mixing of different three-quark configurations in baryonic wave functions [11] can have a very remarkable influence on $S^d_{\rm BB}$ factors, the extended and improved series of which was calculated in the following paper [6].

Collecting the knowledge of our work, in the present paper we emphasize that there is an interesting opportunity to observe the baryon-baryon composition of deuteron, to measure the spectroscopic factors $S^d_{\rm BB}$ and momentum distributions in various BB channels. In fact, it means that there is an experimental opportunity of displaying the composition of quark configurations in deuteron. Namely, we propose the exclusive process of the quasielastic knock-out of protons from deuterons by electrons $^2{\rm H}(e,e'p)R$, where various states of the resonance spectator R can be identified by coincidence measurement of fast final particles. The final proton must have an energy above 1 GeV to avoid the admixture with the process where the slow resonance R is excited off nucleon by incident photon in the deuteron (this process is connected with the usual NN component in the deuteron) and at the same time the resonance-spectator R must be quite slow to avoid the influence of final-state meson exchange triangle diagrams with the resulting redistribution of excited baryons.

We shall see that the specific features of quark structure of virtual baryons sometimes result in the interesting situations when the factorization approximation is no longer valid, the knock-out amplitudes with and without deexcitation are interfering, and the experimentally observable recoil momentum distribution appears to have some unexpected shape.

As an important ingredient of our theoretical presentation we outline the formalism of spectroscopic factors $S^d_{\rm BB}$ (which can be easily generalized, for example, to $12q$ bag, etc.) and deliver the necessary numerical values.

II. MICROSCOPIC FORMALISM FOR ^{2}H$(e,e'p)R$ PROCESS

The diagram describing the quasielastic process under discussion is presented in Fig. 1 where the spectator resonance R can be identified with the baryonic states $\Delta(1232)$, $N(1440)$, $N(1520)$, $N(1535)$, etc. The dashed circle represents the matrix element of the nuclear (quark) electromagnetic current. If the four-momentum transfer q is reasonably small (in quark scale), we can neglect the contribution of sea quarks and write down this matrix element as

$$j^{\nu}_{fi}=\sum_{k=1}^{6}\int\prod_{i=1}^{6}d\mathbf{r}_i\exp(i\mathbf{q}\cdot\mathbf{r}_k)[\psi^*_f(\mathbf{r}_1,\ldots,\mathbf{r}_6)\hat{j}^{\nu}(k)\psi_d(\mathbf{r}_1,\ldots,\mathbf{r}_6)]\ , \tag{1}$$

where f is the symbol of final state and $\hat{j}^{\nu}(k)$ is the operator of the electromagnetic current of quark k. In the nonrelativistic approximation its components look like

$$\hat{\rho}=\hat{e}\ , \tag{2}$$

$$\hat{\mathbf{j}}=\hat{e}\frac{\hat{\mathbf{p}}_i+\hat{\mathbf{p}}_f}{2m}+\hat{\mu}\frac{i}{2m}[\sigma,\mathbf{q}]\ , \tag{2'}$$

where m is the constituent quark mass, and $\hat{e}$ and $\hat{\mu}$ operators of the electric charge and of the magnetic moment, $\hat{e}=\hat{\mu}=\frac{1}{6}e+\frac{1}{2}e\tau_3$. The primary relativistic expression (Dirac particle without anomalous magnetic moment) is

$$j^{\mu}_{fi}=\hat{e}\bar{u}(p_f)\left[\frac{p^{\mu}_f+p^{\mu}_i}{2m}-\frac{\sigma^{\mu\nu}q_{\nu}}{2m}\right]u(p_i)\ . \tag{3}$$

The cross section averaged over initial spin states and summed up on final ones can be expressed as

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{k}_p}=\delta(E_f-E_i-\omega)\frac{e^2}{|\mathbf{k}|\varepsilon' q^4}\times[(k_{\mu}+k'_{\mu})(k_{\nu}+k'_{\nu})+g_{\mu\nu}q^2]\times\frac{1}{(2\pi)^3}\frac{1}{2J_d+1}\sum_{M_dM_pM_R}j^{\mu}_{fi}j^{\nu *}_{fi}\ . \tag{4}$$

Here, the integration over recoil momentum $\mathbf{k}_R$ (which is not fixed directly but is fixed by four-momentum conservation) is carried out, and the matrix element of current (1) no longer contains the factor $(2\pi)^3\delta(\mathbf{k}_p+\mathbf{k}_R-\mathbf{q})$. Ex-

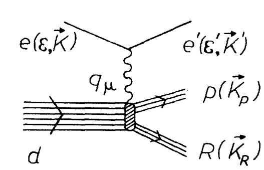

FIG. 1. General diagram corresponding to the ^{2}H$(e,e'p)R$ process.

pression (1) contains everything that we are interested in the quark structure of deuteron. Let us go into some details of it. The final-state wave function of $p+R$ system is

$$\psi_f(\mathbf{r}_1,\ldots,\mathbf{r}_6)=\exp(i(\mathbf{k}_p+\mathbf{k}_R)\cdot\mathbf{R}_C)\times\hat{A}\{\varphi_p(1,2,3)\varphi_R(4,5,6)\chi^{(-)}(\mathbf{r})\}\ , \quad (5)$$

where $\varphi_p(1,2,3)$ and $\varphi_R(4,5,6)$ are internal three-quark wave functions of the baryons and the coordinate system is demonstrated in Fig. 2 and quark antisymmetrizer is

$$\hat{A}=\begin{bmatrix}6\\3\end{bmatrix}^{-1/2}\sum_{P}(-1)^P\hat{P}\ . \quad (6)$$

Its normalization

$$\begin{bmatrix}6\\3\end{bmatrix}^{-1/2}\equiv\left[\frac{6!}{3!3!}\right]^{-1/2}=\frac{1}{\sqrt{20}}$$

assumes that $\chi^{(-)}(\mathbf{r})$ is not antisymmetric with respect to baryon permutations, and preserves the asymptotics of the mutual motion wave function $\chi^{(-)}(\mathbf{r})$ outside the overlap region of the composite particles p and R. The operator $\hat{P}$ comprises all possible quark permutations. In the case of antisymmetric function $\chi^{(-)}$ we should omit out of (6) two- and three-quark exchange terms and replace the normalization factor by

$$\left[\frac{6!}{3!3!2}\right]^{-1/2}=\frac{1}{\sqrt{10}}\ .$$

As long as the initial state is antisymmetric over permutations of quarks, and the interaction operator is symmetric we can replace [7] the antisymmetrizer $\hat{A}$ in Eq. (5) by a simple identity factor

$$\begin{bmatrix}6\\3\end{bmatrix}^{1/2}=\left[\frac{6!}{3!3!}\right]^{1/2}.$$

The matrix element (1) contains terms of two classes. The first of them corresponds to absorption of a virtual photon on one of quarks of the emitted proton while resonance R is a spectator. The second class of terms is evidently just the opposite: a photon with high momentum is absorbed by a nucleon ($\gamma^*+N\to R$) or some other baryon ($\gamma^*+N^*\to R$), while a very fast outgoing proton is a spectator. In this case the cross section is determined by the large momentum of the outgoing proton k_p, and is proportional to $|\psi_{NN}(\mathbf{k}_p)|^2$, where $\psi_{NN}(\mathbf{k})$ is a deuteron NN wave function in the momentum representation. We are interested in the quasielastic kinematical region where the first class terms dominate over the second class ones. Here, the energy of the knocked-out proton must be above 1 GeV, with recoil momentum value being small enough (typically, no more than a few hundred MeV/c). As result, the exchange terms, i.e., the terms of the second class, are suppressed due to a big value of momentum of mutual motion p-R. So, considering the quasielastic region we can write

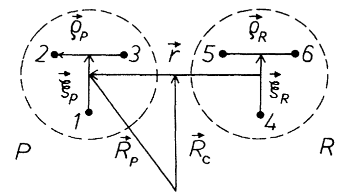

FIG. 2. Jacobi coordinates of the six-quark system.

$$j_{fi}^{\nu}=\int\prod_{i=1}^{6}d\mathbf{r}_i\exp(i\mathbf{q}\cdot\mathbf{r}_1)\{\exp[-i(\mathbf{k}_p+\mathbf{k}_R)\cdot\mathbf{R}_C]\times\left[\frac{6!}{3!3!}\right]^{1/2}\varphi_p^*(1,2,3)\varphi_R^*(4,5,6)\times\chi^{(-)*}(\mathbf{r})3\hat{j}^{\nu}(1)\psi_d(1,2,\ldots,6)\}\ , \quad (7)$$

where the factor 3 is due to identity of quarks.

Due to a large momentum transfer q_μ, in the already discussed kinematical region and, as a consequence, large relative momentum p-R in the final state, we can ignore the effect of distortions and use plane waves to describe outgoing baryons. In this respect, the situation is analogous to that of the understood process ^{2}H$(e,e'p)n$ in the quasielastic region [12], where the effect of distortion is small at smaller outgoing proton energies.

The plane-wave approximation evidently means

$$\chi^{(-)}(\mathbf{r})\to\exp(i\mathbf{Q}\cdot\mathbf{r})\ ,\qquad \mathbf{Q}=\frac{m_R\mathbf{k}_p-m_p\mathbf{k}_R}{m_p+m_R} \quad (8)$$

and necessary transformation from single-particle coordinates $\mathbf{r}_i$ to cluster Jacobi coordinates has to be done in accordance with Fig. 2.

For the sake of convenience let us insert into the matrix element (7), just after the current operator, the closure representation of unit:

$$\sum_{B,M_B}\int d\mathbf{K}_B|\varphi_B(1,2,3),\mathbf{K}_B\rangle\ \langle\varphi_B(1,2,3),\mathbf{K}_B|\equiv 1\ ,$$

where the summation is carried out over all excited states of 3q-cluster B, its spins and spin projections, etc., and $\mathbf{K}_B$ means the cluster center-of-mass momentum.

As a result,

$$j_{fi}^{\nu}=(2\pi)^{3/2}\sum_{B,M_B}\int d\mathbf{R}_p d\xi_p d\rho_p\exp(i\mathbf{q}\cdot\mathbf{r}_1)[\exp(-i\mathbf{k}_p\cdot\mathbf{R}_p)\varphi_p^*(1,2,3)3\hat{j}^{\nu}(1)\varphi_B(1,2,3)\exp(-i\mathbf{k}_R\cdot\mathbf{R}_p)]\frac{1}{(2\pi)^{3/2}}\times\int d\mathbf{r}\exp(i\mathbf{k}_R\cdot\mathbf{r})\left[\frac{6!}{3!3!}\right]^{1/2}\langle\varphi_B(1,2,3)\varphi_R(4,5,6)|\psi_d(1,2,\ldots,6)\rangle\ ; \quad (9)$$

i.e., the matrix element is now divided into two blocks. The first of them,

$$j_{pB}^{\nu}=\int d\mathbf{R}_p d\xi_p d\rho_p \exp(i\mathbf{q}\cdot\mathbf{r}_1)[\exp(-i\mathbf{k}_p\cdot\mathbf{R}_p)\varphi_p^*(1,2,3)3\hat{j}^{\nu}(1)\varphi_B(1,2,3)\exp(-i\mathbf{k}_R\cdot\mathbf{R}_p)] , \tag{10}$$

is the matrix element of the hadron current for the process $e+B\to e'+p$, where $-\mathbf{k}_R$ is the momentum of initial virtual off-shell cluster B and $\mathbf{k}_p$ is the final momentum of proton. In the second block

$$\Phi_{BR}^d(\mathbf{K}_R)=\frac{1}{(2\pi)^{3/2}}\int d\mathbf{r}\exp(i\mathbf{k}_R\cdot\mathbf{r})\Phi_{BR}^d(\mathbf{r}) ,$$

$$\Phi_{BR}^d(\mathbf{r})=\left[\frac{6!}{3!3!}\right]^{1/2}\langle\varphi_B(1,2,3)\varphi_R(4,5,6)| \times\psi_d(1,2,\ldots,6)\rangle \tag{11}$$

we see the function with the specific identity factor appearing due to the many ways one can build up cluster B of quarks with different numbers: 1, 2, 3 or 1, 2, 4, etc.

This function is observable in the $^2\mathrm{H}(e,e'p)R$ process (as well as in other knock-out processes) and is normalized to the spectroscopic factor

$$S_{BR}^d=\frac{1}{2J_d+1}\sum_{M_dM_BM_R}\int|\Phi_{BR}^d(\mathbf{r})|^2d\mathbf{r} =\frac{1}{2J_d+1}\sum_{M_dM_BM_R}\int|\Phi_{BR}^d(\mathbf{k}_R)|^2d\mathbf{k}_R . \tag{12}$$

The matrix element (9) can be presented by the generalized pole diagram of Fig. 3, where inelastic knock-out amplitudes $e+B\to e'+p$ ($B\neq p$) are incorporated, which means that we are going beyond the simplest impulse approximation. We shall see that these inelasticities can sometimes play the dominating role. The baryon-baryon composition of deuteron is just defined by Eq. (11), in particular, the "probabilities of population" of various B-R channels are defined by spectroscopic factors. This rich panorama is originated in microscopical quark treatment of both deuteron and baryons.

If we want to compare a B_1-B_2 "relative motion wave function" in the deuteron with a quantity derived from a microscopic six-quark deuteron wave function $\psi_d(1,\ldots,6)$, we should use

$$\psi_{B_1B_2}(\mathbf{r})=\left[\frac{6!}{3!3!2}\right]^{1/2}\langle\varphi_{B_1}(1,2,3)\varphi_{B_2}(4,5,6)| \times\psi_d(1,\ldots,6)\rangle . \tag{13}$$

The additional factor of $2^{-1/2}$ is motivated by the fact that $\psi_d(1,\ldots,6)$ is antisymmetric under the interchange of baryons B_1 and B_2 (even if they are not identical), because it is antisymmetric under the interchange $(1,2,3)\leftrightarrow(4,5,6)$. But if we describe the process $^2\mathrm{H}(e,e'B_1)B_2$ on the baryon level instead of the quark level and use an antisymmetric B_1B_2 wave function in the initial state, we must multiply the reaction matrix element by $2^{1/2}$ [see the discussion found after formula (6)]. Hence, the full "structure part" of matrix element (1) will again be (11).

But we should underline the fact that (13), in some sense, differs from the usual B_1-B_2 Fock component. In the latter baryons are regarded as elementary structureless particles and the square of the norm of each line in the deuteron Fock column is the probability of the given baryon-baryon component. The total probability for all possible components is unity. In our case, only the wave functions on the quark level are probability amplitudes $\psi_d(1,\ldots,6)$, $\varphi_{B_1}(1,2,3)$, and $\varphi_{B_2}(4,5,6)$, such that each is antisymmetric and normalized to unity. $\psi_{B_1B_2}(\mathbf{r})$, strictly speaking, has no such interpretation. Only in the limit where the size of each cluster B_i approaches zero, in both the bra part of (13) and in the six-quark deuteron wave function, the function $\psi_{B_1B_2}(\mathbf{r})$ has a direct probability interpretation and corresponds literally to the given B_1B_2 component in the deuteron Fock column.

The specific interest from the viewpoint of BB spectroscopy can be seen in such channels of $^2\mathrm{H}(e,e'p)R$ reaction when the sum over B in Eq. (9) is reduced to one term only. Here, the expression (4) describing cross section can be written down in the usual form as product of cross section of basic process $e+B\to e'+p$ and B-R momentum distribution in the deuteron. So, this momentum distribution can be extracted from experiment by the usual straightforward procedure.

Indeed, if we write down all the spin projections, we can present the block of currents in Eq. (4) as

$$\frac{1}{(2\pi)^3}\frac{1}{2J_d+1}\sum_{M_dM_pM_R}j_{fi}^{\mu}j_{fi}^{\nu*}=\frac{1}{2J_d+1}\sum_{\substack{M_dM_pM_R\\M_BM_B'}}j_{pB}^{\mu}(\mathbf{K}_p,M_p;-\mathbf{K}_R,M_B)j_{pB}^{\nu*}(\mathbf{K}_p,M_p;-\mathbf{K}_R,M_B')\Phi_{BR}^d(\mathbf{K}_R;M_B,M_R,M_d) \times\Phi_{BR}^{d*}(\mathbf{K}_R;M_B',M_R,M_d) , \tag{14}$$

with some evident simplification due to orthonormality of Clebsch-Gordan coefficients:

$$\frac{1}{2J_d+1}\sum_{M_dM_R}\Phi^d_{BR}(\mathbf{K}_R;M_B,M_R,M_d)\Phi^{d*}_{BR}(\mathbf{K}_R;M'_B,M_R,M_d)=\delta_{M_BM'_B}\frac{1}{2J_d+1}\sum_{M_dM_R}|\Phi^d_{BR}(\mathbf{K}_R;M_B,M_R,M_d)|^2 .$$

The remaining coupling in Eq. (14) between blocks $j^{\mu}_{pB}j^{\nu*}_{pB}$ and $|\Phi^d_{BR}|^2$ by means of summation over projections M_B can be originated only in a small spin-orbital contribution to the transitional charge density j^0_{pB} and is neglected (see Ref. [12] with discussion of process $A(e,e'p)A-1$ for details). As a result, we can present Eq. (4) as

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{k}_p}=\frac{e^2}{|\mathbf{k}|\varepsilon'q^4}[(k_\mu+k'_\mu)(k_\nu+k'_\nu)+g_{\mu\nu}q^2]\frac{1}{2J_B+1}\sum_{M_pM_B}j^{\mu}_{pB}j^{\nu*}_{pB}\frac{1}{2J_d+1}\sum_{M_BM_RM_d}|\Phi^d_{BR}(\mathbf{K}_R)|^2\delta(E_f-E_i-\omega) . \tag{15}$$

It is natural here to pick out the cross section of a free (inelastic) scattering of ultrarelativistic electron on baryon B with momentum $-\mathbf{k}_R$:

$$\sigma_{eB\to e'p}=\frac{|\mathbf{k}'|e^2}{|\mathbf{k}|q^4}[(k_\mu+k'_\mu)(k_\nu+k'_\nu)+g_{\mu\nu}q^2]\frac{1}{2J_B+1}\sum_{M_pM_B}j^{\mu}_{pB}j^{\nu*}_{pB} . \tag{16}$$

This expression when multiplied to the energy-conserving delta function for the $e+B\to e'+p$ process gives just the laboratory system cross section

$$\frac{d^2\sigma_{eB\to e'p}}{d\Omega'd\varepsilon'}$$

if $\mathbf{k}_R=0$.

So, Eq. (15) can be rewritten as

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{k}_p}=\frac{\sigma_{eB\to e'p}(\theta_e,\mathbf{q},\mathbf{K}_p)}{\mathbf{k}'^2}\frac{1}{2J_d+1}\sum_{M_BM_RM_d}|\Phi^d_{BR}(\mathbf{K}_R)|^2\delta(E_f-E_i-\omega) , \tag{17}$$

where the momentum distribution

$$\rho(\mathbf{K}_R)=\frac{1}{2J_d+1}\sum_{M_BM_RM_d}|\Phi^d_{BR}(\mathbf{K}_R)|^2 \tag{18}$$

is a characteristic of B-R component in deuteron and its normalization

$$S^d_{BR}=\int d\mathbf{K}_R\rho(\mathbf{K}_R) \tag{19}$$

is the corresponding spectroscopic factor, $S^d_{BR}=2N^d_{BR}$, where N^d_{BR} is the "effective number" of a given B-R pair in deuteron.

III. BARYON-BARYON CONTENT OF THE SIX-QUARK WAVE FUNCTION OF DEUTERON

As noted in the preceding section, the quark-cluster content of six-quark composite system has to be explored by means of Eq. (11).

As far as the binding energy of the baryon-baryon components under consideration in the deuteron exceed 0.5 GeV, they are localized within the spatial region of very essential nucleon-nucleon overlap $r\le 1-1.3$ fm. Here, the six-quark wave function can be expanded over the simplest shell-model quark configurations s^6 and s^4p^2 with different symmetries, i.e.,

$$\begin{gathered}|s^6[6]_X[2^3]_{CS}\rangle,\quad |s^4p^2[6]_X[2^3]_{CS}\rangle ,\\ |s^4p^2[42]_X[42]_{CS}\rangle,\quad |s^4p^2[42]_X[321]_{CS}\rangle ,\\ |s^4p^2[42]_X[2^3]_{CS}\rangle,\quad |s^4p^2[42]_X[31^3]_{CS}\rangle ,\\ |s^4p^2[42]_X[21^4]_{CS}\rangle,\quad |s^52s[6]_X[2^3]_{CS}\rangle .\end{gathered} \tag{20}$$

As usual Young schemes $[f]_X$ and $[f]_{CS}$ characterize the spatial and color-spin parts of wave function [13,14], all the remaining symmetries are defined by the quantum numbers of deuteron.

The calculations [2] do not include quark configurations $|s^4p^2[6]_X[2^3]_{CS}\rangle$ and $s^52s[6]_X[2^3]_{CS}\rangle$ which are important when projecting onto various BB channels. To use the results [2] and to take into account these configurations we can consider the identities [13]

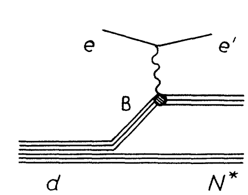

FIG. 3. Pole diagram corresponding to the $e+A\to e'+p$ elementary process and the spectator particle N^*.

$$\hat{A}[\varphi_N(1,2,3)\varphi_N(4,5,6)\Phi_{0s}(\mathbf{r})]_{S=1\ T=0}=\left[\frac{10}{9}\right]^{1/2}|s^6[6]_X[2^3]_{CS}\rangle_{\rm TISM}\ , \tag{21}$$

$$\begin{aligned}\hat{A}[\varphi_N(1,2,3)\varphi_N(4,5,6)\Phi_{2s}(\mathbf{r})]_{S=1\ T=0}\\ =\left[\frac{410}{405}\right]^{1/2}\Bigg\{\left[\frac{9}{41}\right]^{1/2}\Bigg|\left[\left[\frac{5}{6}\right]^{1/2}s^5 2s-\left[\frac{1}{6}\right]^{1/2}s^4p^2\right][6]_X[2^3]_{CS}\Bigg\rangle_{\rm TISM}\\ -\left[\frac{81}{205}\right]^{1/2}|s^4p^2[42]_X[42]_{CS}\rangle_{\rm TISM}+\left[\frac{64}{205}\right]^{1/2}|s^4p^2[42]_X[321]_{CS}\rangle_{\rm TISM}\\ +\left[\frac{1}{41}\right]^{1/2}|s^4p^2[42]_X[2^3]_{CS}\rangle_{\rm TISM}-\left[\frac{2}{41}\right]^{1/2}|s^4p^2[42]_X[31^3]_{CS}\rangle_{\rm TISM}\Bigg\}\ .\end{aligned} \tag{22}$$

Here, symbol TISM (translation-invariant shell model) means the exclusion of nonphysical c.m. oscillations,

$$\varphi_N\equiv|s^3[3]_X[21]_{CS}L=0S=\tfrac{1}{2}T=\tfrac{1}{2}\rangle_{\rm TISM}\ , \tag{23}$$

and $\Phi_{0s}(\mathbf{r})$ and $\Phi_{2s}(\mathbf{r})$ are usual oscillator wave functions

$$\Phi_{0s}(\mathbf{r})=\sqrt{4\pi}\left[\frac{2}{3}\pi b^2\right]^{-3/4}\exp\left[-\frac{3}{4}\frac{r^2}{b^2}\right]Y_{00}(\hat{r})\ , \tag{24}$$

$$\Phi_{2s}(\mathbf{r})=\left[\frac{3}{2}\right]^{1/2}\left[1-\frac{r^2}{b^2}\right]\Phi_{0s}(\mathbf{r})\ . \tag{25}$$

Considering the resonating group method (RGM) representation of six-quark deuteron wave function [3]

$$\psi_d(6q)=\hat{A}[\varphi_N(1,2,3)\varphi_N(4,5,6)\chi(\mathbf{r})]\ , \tag{26}$$

we can say that the expansion under investigation corresponds to approximate representation of function $\chi(\mathbf{r})$ at $r\le 1-1.3$ fm as the superposition of oscillator functions (24) and (25).

In the above formulas (21), (22), and (26) it is implied that antisymmetrization operator $\hat{A}$ is normalized as

$$\hat{A}=\frac{1}{\sqrt{10}}\left[1-\sum_{i=1}^{3}\sum_{j=4}^{6}\hat{P}_{ij}\right]\ . \tag{27}$$

The numerical results can be very sensitive to the specific features of microscopic quark model under consideration, and it is just one of the central points. For instance, well-known calculations [3] are based on (renormalized) one-gluon exchange between quarks and pion exchange in combination with RGM ansatz for two nucleons in the deuteron. Here, $\hat{N}^{1/2}\chi$ appears to be rather close to the traditional "realistic" wave functions (e.g., Paris potential) and the short-range part of the function χ which we are interested in can be obtained from Paris wave function by the action of the inverse operator $\hat{N}^{-1/2}$ ($\hat{N}$ is the nonlocal normalization operator of RGM which reflects antisymmetrization on quarks, i.e., the quark exchange between nucleons when they move "slowly" in deuteron). In total, the model under discussion can be characterized by the wave function of the next kind:

$$\begin{aligned}\psi_d(1,2,\ldots,6)|_{r\le 1-1.3\ \rm fm}=0.1819|s^6[6]_X[2^3]_{CS}\rangle_{\rm TISM}-0.1191|[(\tfrac{5}{6})^{1/2}s^5 2s-(\tfrac{1}{6})^{1/2}s^4p^2][6]_X[2^3]_{CS}\rangle_{\rm TISM}\\ +0.1634|s^4p^2[42]_X[42]_{CS}\rangle_{\rm TISM}-0.1273|s^4p^2[42]_X[321]_{CS}\rangle_{\rm TISM}\\ -0.0538|s^4p^2[42]_X[2^3]_{CS}\rangle_{\rm TISM}+0.0464|s^4p^2[42]_X[31^3]_{CS}\rangle_{\rm TISM}\\ +0.0057|s^4p^2[42]_X[21^4]_{CS}\rangle_{\rm TISM}\ .\end{aligned} \tag{28}$$

Here, s^6 configuration has the weight 3.3% with its variation between 2.5% and 3.3% for different sets of initial parameters.

Each quark configuration entering deuteron wave function (28) can be presented by means of the Talmi-Moshinsky-Smirnov (TMS) transformation in combination with fractional parentage technique as a superposition of various cluster components

$$|s^\nu p^{6-\nu}\alpha LST\rangle_{\rm TISM}=\sum_{B_1,B_2,nl}\Gamma^{\alpha}_{B_1,B_2,nl}(\varphi_{B_1}(1,2,3)\varphi_{B_2}(4,5,6)\Phi_{nl}(\mathbf{r}):LST)\ . \tag{29}$$

Here, α means all the necessary set of quantum numbers (including symmetries $[f]_X$, $[f]_{CS}$, etc.) to define the six-quark state unambiguously and $\Gamma^{\alpha}_{B_1,B_2,nl}$ is the corresponding spectroscopic amplitude. Its square multiplied by the identity factor just gives the "probability" to find the B_1B_2 component in the quark configuration under discussion. Further, $\Phi_{nl}(\mathbf{r})$ is the oscillator wave function (n is the number of excitation quanta) of the baryon-baryon mutual motion and summation in Eq. (29) is carried out over all baryonic internal states and quantum numbers nl of mutual motion.

The specific feature of expansion (29) is that the antisymmetric six-quark wave function is expanded into the sum of orthogonal but not fully antisymmetric terms. Each term is antisymmetric only within the cluster B_1 and B_2 (using the "fast" process we deal with the "instantaneous portrait" of virtual baryons in $6q$ system—see Introduction).

The formalism of spectroscopic amplitudes $\Gamma^{\alpha}_{B_1,B_2,nl}$ can be characterized as follows (our previous scheme [14,15] is modified here to include the internal excitations of baryons).

First, in the common notations, we use the reduction chain

$$\begin{aligned}&SU(24)_{XCST}\supset SU(6)_{CS}\times SU(4)_{XT}\supset\\&\supset SU(6)_{CS}\times SU(2)_X\times SU(2)_T\supset SU(3)_C\times SU(2)_S\times SU(2)_X\times SU(2)_T\ .\end{aligned}\tag{30}$$

Here, the physically important CS symmetry is considered and, in addition the Young schemes $[f]_{CS}$ and $[f]_{XT}$ are to be mutually conjugated in what simplifies the numerical procedures.

Second, the expression for amplitude Γ is of factorizable form

$$\Gamma^{\alpha}_{B_1,B_2,nl}=F^{XCST}_{CS\cdot XT}F^{CS}_{C\cdot S}F^{XT}_{X\cdot T}F^{X}_{s+p}\Gamma_{\rm TISM}\ ,\tag{31}$$

which corresponds to the reduction (30). Here, we mean scalar factors of SU(n) group Clebsch-Gordan coefficients [14,15], e.g.,

$$F^{CS}_{C\cdot S}=\left[\begin{array}{cc|c}[f]_C & [f]_S & [f]_{CS}\\ ([f']_C\times[f'']_C) & ([f']_S\times[f'']_S) & ([f']_{CS}\times[f'']_{CS})\end{array}\right]\ ,\tag{32}$$

and, in particular,

$$F^{XCST}_{CS\cdot XT}=\left[\frac{n_{[f']_{CS}}n_{[f'']_{CS}}}{n_{[f]_{CS}}}\right]^{1/2}\ ,$$

where it is implied that $[f]_{XCST}=[1^6]$.

Further, factor F^X_{s+p} describes division of the quark configuration $s^kp^n[f]_X$ into subsystems $s^{k'}p^{n'}[f']_X$ and $s^{k''}p^{n''}[f'']_X$. If the orbital Young scheme $[f]_X$ has no more than two rows, it is possible to use quasispin formalism [14,15], which enables us to reduce the above factor to the usual Clebsch-Gordan coefficient

$$\begin{aligned}\Gamma^X_{s+p}&\equiv\Gamma^{[f]_X,[f']_X\times[f'']_X}_{s^kp^n,s^{k'}p^{n'}\times s^{k''}p^{n''}}\\&=(lm\,|\,l'm',l''m'')\ ,\end{aligned}\tag{33}$$

where $l=f_1-f_2$, $m=k-n$, $l'=f_1'-f_2'$ and so on. Finally, the last factor, $\Gamma_{\rm TISM}$(TMS coefficient) [7,8] describes the transformation of the many-particle shell-model wave function to cluster Jacobi coordinates with the excluded c.m. motion

$$\begin{aligned}&\boldsymbol{\rho}_1=\mathbf{r}_1-\mathbf{r}_2,\quad \xi_1=\tfrac{1}{2}(\mathbf{r}_1+\mathbf{r}_2)-\mathbf{r}_3\ ,\\&\boldsymbol{\rho}_2=\mathbf{r}_4-\mathbf{r}_5,\quad \xi_2=\tfrac{1}{2}(\mathbf{r}_4+\mathbf{r}_5)-\mathbf{r}_6\ ,\\&\mathbf{r}=\tfrac{1}{3}(\mathbf{r}_1+\mathbf{r}_2+\mathbf{r}_3)-\tfrac{1}{3}(\mathbf{r}_4+\mathbf{r}_5+\mathbf{r}_6)\ ,\end{aligned}\tag{34}$$

and the distribution of n oscillator quanta over various Jacobi coordinates. Resulting amplitudes $\Gamma^{\alpha}_{B_1,B_2,nl}$ for colorless channels are given in Table I. Here we comment briefly on these data. It is well known [14,18] that quark configuration s^6 is a superposition of three cluster components $\varphi_N\varphi_N\Phi_{0s}(\mathbf{r})$, $\varphi_\Delta\varphi_\Delta\Phi_{0s}(\mathbf{r})$, and $\varphi_C\varphi_C\Phi_{0s}(\mathbf{r})$. If we deal with the quark configuration $s^4p^2-s^52s$, it is possible to distribute two oscillator quanta of excitation over various Jacobi coordinates—both internal coordinates of clusters B_1 and B_2 and the coordinate of their relative motion. Thus, side by side with three of the above components, expansion (29) will contain many new components

$$\begin{aligned}&\varphi_N\varphi_N\varphi_{2s}(\mathbf{r}),\quad \varphi_{N^*}\varphi_N\Phi_{1p}(\mathbf{r}),\quad \varphi_{N^*}\varphi_{N^*}\Phi_{0s}(\mathbf{r})\ ,\\&\varphi_{N^{**}}\varphi_N\Phi_{0s}(\mathbf{r}),\quad \varphi_\Delta\varphi_\Delta\varphi_{2s}(\mathbf{r}),\quad \varphi_{\Delta^{**}}\varphi_\Delta\Phi_{0s}(\mathbf{r})\ ,\\&\varphi_{\Delta^*}\varphi_\Delta\Phi_{1p}(\mathbf{r}),\quad \varphi_{\Delta^*}\varphi_{\Delta^*}\Phi_{0s}(\mathbf{r}),\quad \varphi_C\varphi_C\varphi_{2s}(\mathbf{r})\ ,\\&\varphi_{C^{**}}\varphi_C\Phi_{0s}(\mathbf{r}),\quad \varphi_{C^*}\varphi_C\Phi_{1p}(\mathbf{r}),\ldots\ ,\end{aligned}$$

TABLE I. Spectroscopic amplitudes for six-quark configurations s^6, s^4p^2, and $s^5 2s$ in deuteron with $L=0$, $S=1$, $T=0$; the orbital momenta, spins, and isospins of cluster components are to be converted to total values $L=0$, $S=1$, $T=0$.

Cluster component \ TISM configuration	$s^6[6]_X$ $[2^3]_{CS}$	$(\frac{5}{6})^{1/2}s^5 2s-$ $(\frac{1}{6})^{1/2}s^4p^2$ $[6]_X[2^3]_{CS}$	$s^4p^2[42]_X$ $[42]_{CS}$	$[321]_{CS}$	$[2^3]_{CS}$	$[31^3]_{CS}$	$[21^4]_{CS}$
$\varphi_N\varphi_N\Phi_{0s}(\mathbf{r})$	$(\frac{1}{9})^{1/2}$						
$\varphi_N\varphi_N\Phi_{2s}(\mathbf{r})$		$(\frac{1}{45})^{1/2}$	$-(\frac{1}{25})^{1/2}$	$(\frac{64}{2025})^{1/2}$	$(\frac{1}{405})^{1/2}$	$-(\frac{2}{405})^{1/2}$	
$\varphi_\Delta\varphi_\Delta\Phi_{0s}(\mathbf{r})$	$-(\frac{4}{45})^{1/2}$						
$\varphi_\Delta\varphi_\Delta\Phi_{2s}(\mathbf{r})$		$-(\frac{4}{225})^{1/2}$			$(\frac{64}{2025})^{1/2}$		$(\frac{4}{81})^{1/2}$
$\varphi_N\varphi_{N^*}\Phi_{1p}(\mathbf{r})$			$(\frac{1}{100})^{1/2}$	$(\frac{1}{2025})^{1/2}$	$-(\frac{5}{324})^{1/2}$	$-(\frac{1}{810})^{1/2}$	
$\varphi_{N^*}\varphi_N\Phi_{1p}(\mathbf{r})$			$(\frac{1}{100})^{1/2}$	$(\frac{1}{2025})^{1/2}$	$-(\frac{5}{324})^{1/2}$	$-(\frac{1}{810})^{1/2}$	
$\varphi_{N^*}\varphi_{N^*}\Phi_{0s}(\mathbf{r})$			$-(\frac{1}{50})^{1/2}$	$(\frac{32}{2025})^{1/2}$	$-(\frac{10}{324})^{1/2}$	$+(\frac{4}{405})^{1/2}$	
$\varphi_N\varphi_{N^{**}}\Phi_{0s}(\mathbf{r})$		$(\frac{2}{45})^{1/2}$	$(\frac{1}{200})^{1/2}$	$-(\frac{8}{2025})^{1/2}$	$-(\frac{1}{3240})^{1/2}$	$(\frac{1}{1620})^{1/2}$	
$\varphi_{N^{**}}\varphi_N\Phi_{0s}(\mathbf{r})$		$(\frac{2}{45})^{1/2}$	$(\frac{1}{200})^{1/2}$	$-(\frac{8}{2025})^{1/2}$	$-(\frac{1}{3240})^{1/2}$	$(\frac{1}{1620})^{1/2}$	

where the number of asterisks means the number of internal excitation quanta in a baryon. In the first approximation

$$\varphi_\Delta=|s^3[3]_X[1^3]_{CS}L=0,S=\tfrac{3}{2},T=\tfrac{3}{2}\rangle_{\rm TISM}\ , \quad (35)$$

$$\varphi_{N^*}=|s^2p\,[21]_X[21]_{CS}L=1,S=\tfrac{1}{2},T=\tfrac{1}{2}\rangle_{\rm TISM}\ , \quad (36)$$

$$\varphi_{N^{**}}=|sp^2[3]_X[21]_{CS}L=0,S=\tfrac{1}{2},T=\tfrac{1}{2}\rangle_{\rm TISM}\ . \quad (37)$$

This approximation seems, in general, to be reasonable for our purposes as far as, for example, the corrections to wave functions (35) and (36) give an effect of the order of 10–30 % to amplitudes for BB channels with Δ or N^* where N^* is either $N_{3/2^-}(1520)$ or $N_{1/2^-}(1535)$ [6]. However, the model under consideration has an interesting property that for the particular case of projecting wave function (28) by means of Eq. (29) into the NN^{**} channel where N^{**} is Roper resonance $N_{1/2^+}(1440)$ the interference of configuration

$$\left[\left[\frac{5}{6}\right]^{1/2}s^5 2s-\left[\frac{1}{6}\right]^{1/2}s^4p^2\right][6]_X[2^3]_{CS}\rangle$$

with the sum of others is negative and the two amplitudes almost cancel each other. So we refine wave function (37) taking into account the small admixture

$$|s^3[3]_X[21]_{CS}L=0,S=\tfrac{1}{2},T=\tfrac{1}{2}\rangle_{\rm TISM}$$

according to Ref. [2] (the rest of the correcting terms give a small effect).

Spectroscopic factors calculated by the method mentioned above have the following values (see also Ref. [6]):

$$\left.\begin{array}{l} S^d_{\Delta\Delta}\simeq(4-6)\times10^{-2} \\ S^d_{NN(1440)}\simeq(2-4)\times10^{-3} \\ S^d_{NN(1535)}\simeq(1.5-2)\times10^{-3} \\ S^d_{NN(1520)}\simeq 2S^d_{NN(1535)} \\ S^d_{N(1535)N(1535)}\simeq(2-3)\times10^{-4} \\ S^d_{N(1520N(1520)}\simeq 10S^d_{N(1535)N(1535)} \\ S^d_{N(1520)N(1535)}\simeq 8S^d_{N(1535)N(1535)} \end{array}\right\}\times\langle T_1t_1T_2t_2|00\rangle^2\ .$$

Other calculated spectroscopic factors $S^d_{N(1440)N(1440)}$, $S^d_{N(1440)N(1520)}$, and $S^d_{N(1440)N(1535)}$ are of the order of 10^{-5}. All the above values, to be compact, correspond to summation over all the possible charge distributions between baryons B_1 and B_2. So, if each charge is fixed, we must multiply given numbers to $(T_1t_1,T_2t_2|00)^2$, where T_1,t_1 is the isospin of baryon B_1 and its projection, etc.

One final important remark is that $\Delta\Delta$ component in deuteron is practically entirely called forth by s^6 configuration (the influence of the s^4p^2 configuration is one order of magnitude less). On the contrary, all the components containing negative-parity baryon with one oscillator quantum of internal excitation are caused entirely by configuration s^4p^2.

It is reasonable to repeat that the example of wave function (28) under discussion corresponds well to deuteron M_1 and other electromagnetic form factors as its most valuable experimental foundation [1]. However, NN-scattering data taken in the broad energy range $0<E_{\rm lab}<6$ GeV are well described by deep attractive NN potentials with forbidden states [4] which correspond to the nodal wave function $\Phi_{2S}(\mathbf{r})$ of the deuteron ground state [4,16]. This fact is natural to ascribe to quark configuration

$$s^2p^2[42]_X[42]_{CS} ,$$

where the color magnetic interaction of quarks creates the powerful NN attraction which is just able to compensate the energy loss due to excitation of two quarks into p state [15].

This model does not contain the $\Delta\Delta$ component in deuteron at all, and spectroscopic factors $S_{B_1B_2}$ for BB pairs with orbitally excited baryons with negative parity should grow approximately 2–4 times as much as numerical values discussed above (the sum of squared coefficients in formula (28) now equals the probability of configuration $s^4p^2[42]_X[42]_{CS}$). But the value of $S^d_{NN(1440)}$ decreases two times as far as this value is conditioned mainly by the small admixture of the s^3 component in the $N(1440)$.

In this connection it is interesting to mention that there exists one old inclusive experiment $\gamma+d\to\Delta+X$ [24] which showed the 3% $\Delta\Delta$ probability in the deuteron which is in good agreement with our result, $N^d_{\Delta\Delta}\simeq(2-3)\times10^{-2}$. But there is a lot of skepticism in the literature [17] concerning the experiments of the inclusive character. We propose the exclusive experiment which can be helpful in ascertaining the baryon-baryon composition of deuteron including the problem of the $\Delta\Delta$ component and as a result to find out the structure of the six-quark shell-model configurations in deuteron.

Finally, note that we used a six-quark shell expansion of the deuteron wave function, which is valid only at quite short ranges ($r\leq1-1.3$ fm). Of course, the asymptotics of the $\Delta\Delta$, NN^*, . . . components, which do not possess six-quark origin, stretch far distances up to 2–2.5 fm. On the other hand, the influence of these asymptotic parts on the spectroscopic factors is small [25].

IV. CONSIDERATION OF SPECIFIC TRANSITIONS

A. Description of process ^{2}H$(e,e'p)N(1440)$

This reaction is called forth by the $NN(1440)$ component in deuteron where the proton serves as a virtual cluster B [see formulas (9), (16)–(18), and Fig. 3] with elastic electron-proton scattering $e+p\to e'+p$ being primary fast process:

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{K}_p}=\frac{\sigma_{ep\to e'p}(\theta_e,q_\mu,\mathbf{K}_p)}{\mathbf{k}'^2}\times\rho_{pN(1440)}(\mathbf{K}_R)\delta(E_f-E_i-\omega) . \quad (38)$$

Momentum distribution $\rho_{pN(1440)}(\mathbf{K}_R)$ is defined by formulas (11), (12), (18), and (19). Its normalization can be clarified by the formula

$$\rho_{pN(1440)}(\mathbf{K}_R)=S^d_{NN(1440)}|\Psi_{NN(1440)}(K_R)|^2|Y_{00}(\hat{\mathbf{K}}_R)|^2 , \quad (39)$$

where function $\Psi_{NN(1440)}(k)$ is normalized to unity. The shape of the momentum distribution presented by expression $|\Psi_{NN(1440)}(k)|^2$ is shown in Fig. 4. It is calculated by means of microscopic projection procedure (12) with the wave function (28). As it was noted above, the wave function $\Psi_{NN(1440)}(k)$ is the superposition of oscillator functions $\Phi_{0s}(k)$ and $\Phi_{2s}(k)$. In fact, it is rather close to the result of the calculation with the deep attractive $NN(1440)$ potential [19], i.e., to the result of the second model under consideration (six-quark function $s^4p^2[42]_X[42]_{CS}$). Therefore, the two models differ only by the value of the cross section (not by their form).

Here we finally comment on the cross section of electron-proton elastic scattering $\sigma_{ep\to e'p}$. We can separate in Eq. (16) the longitudinal and transversal form factors and write

$$\sigma_{eB\to e'p}=4\pi\sigma_{\text{Mott}}\left[\frac{q^4}{\mathbf{q}^4}F_L^2+\left[\frac{1}{2}\frac{q^2}{\mathbf{q}^2}+tg^2\frac{\theta}{2}\right]F_T^2\right] , \quad (40)$$

where F_T^2 is determined by the $\mathbf{q}$-transverse component of current and F_L^2 by the timelike component:

$$F_T^2=\frac{1}{4\pi(2J_B+1)}\sum_{M_BM_p}\mathbf{j}^T_{pB}\mathbf{j}^{T*}_{pB} , \quad (41)$$

$$F_L^2=\frac{1}{4\pi(2J_B+1)}\sum_{M_BM_p}\mathbf{j}^0_{pB}\mathbf{j}^{0*}_{pB} . \quad (42)$$

Here the $\mathbf{q}$-transversal component of electromagnetic current for our case of the $p\to p$ transition is

$$\mathbf{j}^T_{pp}=\frac{\hat{e}_p}{2m_p}(\mathbf{p}_i+\mathbf{p}_f)^T\delta_{M_iM_f}+\left\langle\frac{1}{2}M_f\left|i\frac{\hat{\mu}_p}{2m_p}(\sigma\times\mathbf{q})\right|\frac{1}{2}M_i\right\rangle , \quad (43)$$

where

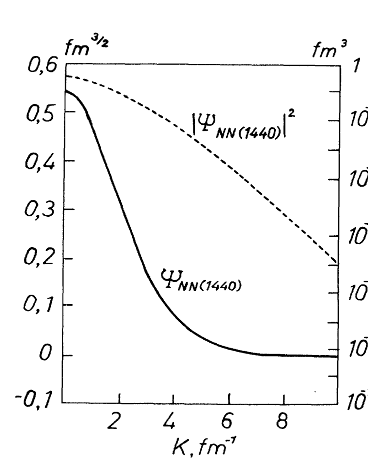

FIG. 4. $NN(1440)$ relative motion wave function in deuteron (left scale) and its square (right scale).

$$(\mathbf{p}_i+\mathbf{p}_f)^T=2[(\mathbf{K}_R\mathbf{q})(\mathbf{q}/q^2)-\mathbf{K}_R]$$

due to the conservation law $-\mathbf{K}_R+\mathbf{q}=\mathbf{K}_p$ (see Fig. 1 for notations) and $\hat{e}_p$ and $\hat{\mu}_p$ are set by the modified electric and magnetic form factors of proton, respectively:

$$\begin{aligned}
\hat{e}_p&=\bar{G}_E(q^2)=\frac{G_E(q^2)}{(1-q^2/4m_p^2)^{1/2}}\,,\\
\hat{\mu}_p&=\bar{G}_M(q^2)=\frac{G_M(q^2)}{(1-q^2/4m_p^2)^{1/2}}\,,\\
G_E(q^2)&=\left[1-\frac{q^2}{0.71\ \mathrm{GeV}^2}\right]^{-2},\\
G_M(q^2)&=\mu_p G_E(q^2)\,.
\end{aligned} \tag{44}$$

The charge component of the current is defined by the charge operator $\hat{e}_p$. As a result of the above definitions we can write

$$F_T^2=\frac{1}{4\pi}\left[\frac{\bar{G}_E(q^2)}{m_p^2}\left[\mathbf{K}_R^2-\frac{(\mathbf{K}_R\cdot\mathbf{q})^2}{\mathbf{q}^2}\right]+\frac{\mathbf{q}^2}{2m_p^2}\bar{G}_M(q^2)\right], \tag{45}$$

$$F_L^2=(1/4\pi)\bar{G}_E(q^2)\,. \tag{46}$$

After substitution of these two expressions into formula (40) when $k_R=0$ (free proton) we get the well-known Rosenbluth formula (within the recoil factor). If $k_R\neq0$ the difference from the Rosenbluth formula is caused by the contribution into Eq. (45) of the convective current.

Here we remind the reader also that the expression for hadronic current written above and, as result, expressions for F_T^2 and F_L^2 are perfectly reproduced in the quark model (10) with the use of quark currents (2) and (2′) and the s^3 function for the nucleon ($m_q=\frac{1}{3}m_p$). Here, $\mu_p=3$ and the partial Fourier transform of internal nucleon density distribution

$$F_{p\to p}(\mathbf{q})=\langle\varphi_N(\rho,\xi)|\exp(-i\tfrac{2}{3}\mathbf{q}\cdot\xi)|\varphi_N(\rho,\xi)\rangle \tag{47}$$

serves as form factor $G_E(q^2)$. It equals $\exp(-b^2\mathbf{q}^2/6)$ if the oscillator s^3 function is chosen for the nucleon, and the observed values of $G_E(q^2)$ are reproduced satisfactorily in the momentum transfer range $q<2$–3 GeV/c when $b=0.5$–0.6 fm.

In total, the expected cross section value of process

$${}^2\mathrm{H}(e,e'p)N(1440)$$

integrated over momentum distribution equals 10^{-3} of that for the "basic" process ${}^2\mathrm{H}(e,e'p)n$ and depends on the model considered (see above).

B. Process ${}^2\mathrm{H}(e,e'p)\Delta$

This process in the kinematical region of quasielastic knock-out (see Sec. II above) is originated from the $\Delta\Delta$ component in deuteron (the contribution of $\Delta^*\Delta$ and $\Delta^{**}\Delta$ components is small). So it is of principal importance for the inspection of various quark models. The primary fast process here is the inelastic electron scatterings $e+\Delta\to e'+p$ with the quark spin-isospin flip. The expression for differential cross-section, therefore, looks like

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{K}_p}=\frac{\sigma_{e\Delta\to e'p}(\theta_e,q_\mu,\mathbf{K}_p)}{\mathbf{k}'^2}\rho_{\Delta\Delta}(\mathbf{K}_R)\delta(E_f-E_i-\omega) \tag{48}$$

and the corresponding momentum distribution (which has the sense at $\mathbf{K}_R\le0.4$–0.5 GeV/c) can be presented by the formula

$$\rho_{\Delta\Delta}(\mathbf{K}_R)=S^d_{\Delta\Delta}|\Phi_{0s}(K_R)|^2|Y_{00}(\hat{\mathbf{K}}_R)|^2\,. \tag{49}$$

The "reduced" distribution $|\Phi_{0s}(K_R)|^2$ normalized to unity is shown in Fig. 5. Its oscillatorlike shape is connected with the large value of deuteron binding energy in the $\Delta\Delta$-channel (localization of system at small distance). Something similar was noted in connection with Fig. 4—the approximation of the $NN(1440)$-momentum distribution by the superposition of the two lowest oscillator functions appears to be rather close to the "perfect" one, i.e., calculated within the $NN(1440)$ potential model [19]. The wrong character for the Gaussian asymptotics (instead of the pole one) occurs here practically only at $k_R=0.7$ GeV/c.

The primary (elementary) process $e+\Delta\to e'+p$ can be classified as an almost pure M_1 transition and, as a result, the longitudinal form factor F_L^2 in Eq. (40) can be neglected. The transversal form factor calculated by means of the quark currents (2) and (2′) can be written as

$$\begin{aligned}
F_T^2&=\frac{1}{2J_\Delta+1}\left|\left\langle\varphi_N(1,2,3)\left|\left|\sum_{i=1}^{3}\exp(i\mathbf{q}\cdot\mathbf{r}_i)\hat{\mu}(i)\frac{i}{2m_q}[\sigma(i),\mathbf{q}]\right|\right|\varphi_\Delta(1,2,3)\right\rangle\right|^2\\
&=\frac{1}{4\pi}\frac{2}{9}\frac{|\mathbf{q}|^2}{m_q^2}F^2_{\Delta\to p}(\mathbf{q})\,.
\end{aligned} \tag{50}$$

If the oscillator s^3 functions are chosen for N and Δ then the form factor of transition

$$F^2_{\Delta\to p}(\mathbf{q})=\langle\varphi_N(\rho,\xi)|\exp(-i\tfrac{2}{3}\mathbf{q}\cdot\xi)|\varphi_\Delta(\rho,\xi)\rangle \tag{51}$$

has the same q dependence $\exp(-\frac{1}{6}b^2\mathbf{q}^2)$ as the elastic form factor $p\to p$ (see above). Of course such q dependence of both form factors is not realistic if we have in mind the broad range of momentum transfer. Therefore, it is reasonable to just use the experimentally measured form factor of electroexcitation $N\to\Delta$ [20] which decreases with increasing q somewhat more steep than nu-

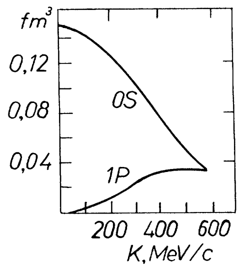

FIG. 5. Momentum distributions $\Phi^2_{0s}(k)$ and $\Phi^2_{1p}(k)$.

cleon elastic form factor (44).

The expected cross-section is approximately 10^{-2} of that for process ^{2}H$(e,e'p)n$ if the common description (28) of the six-quark deuteron is valid and two orders of magnitude smaller if the quark state of deuteron is given by configuration $s^4p^2[42]_X[42]_{CS}$.

C. Reactions ^{2}H$(e,e'p)N(1520)$ and ^{2}H$(e,e'p)N(1535)$

These reactions are interesting by their complicated character—the deexcitation amplitudes of virtually excited baryons B are of primary importance here and the rich interference of a few terms corresponding to various virtual states of these baryons in Eq. (9) appears as a result.

Indeed, the spectators $N(1520)$ and $N(1535)$ have one oscillator quantum of internal excitation. So the reactions considered are connected with two different cluster components of six-quark configuration s^4p^2, namely, $\varphi_N\varphi_{N^*}\Phi_{1p}(\mathbf{r})$ and $\varphi_{N^*}\varphi_{N^*}\Phi_{0s}(\mathbf{r})$. Evidently, the second component contributes to our reactions by mean of elementary deexcitation process $e+N^*\to e'+N$.

Having in mind process ^{2}H$(e,e'p)N(1520)$ we can write expression (9) in terms of baryon-baryon composition of the deuteron as follows

$$j^{\nu}_{fi}\sim j^{\nu}_{pp}\Phi^d_{pN(1520)}+j^{\nu}_{pN(1520)}\Phi^d_{N(1520)N(1520)}+j^{\nu}_{pN(1535)}\Phi^d_{N(1535)N(1520)}\ , \tag{52}$$

and for process ^{2}H$(e,e'p)N(1535)$

$$j^{\nu}_{fi}\sim j^{\nu}_{pp}\Phi^d_{pN(1535)}+j^{\nu}_{pN(1535)}\Phi^d_{N(1535)N(1535)}+j^{\nu}_{pN(1520)}\Phi^d_{N(1520)N(1535)}\ . \tag{53}$$

To go forth we must now compare the magnitudes and q dependence of matrix elements of hadronic electromagnetic current both for elastic scattering j^{ν}_{pp} and for deexcitation processes $j^{\nu}_{pN(1520)}$ and $j^{\nu}_{pN(1535)}$. The above currents can be calculated easily in the nonrelativistic quark model by means of wave functions (23) and (36) for nucleon and resonances, respectively:

$$\mathbf{j}^T_{pN^*}=i((11,\tfrac{1}{2}M_p|J^*M^*_J)\mathbf{e}_{-1}+(1-1,\tfrac{1}{2}M_p|J^*M^*_J)\mathbf{e}_{+1})\frac{1}{\sqrt{6}m_qb}$$
$$\times\exp\left[-\frac{\mathbf{q}^2b^2}{6}\right]+\left[\frac{(\mathbf{P}_i+\mathbf{P}_f)^T}{6m_q}\delta_{M_pM^*_J}+\frac{i}{2m_q}\left\langle\frac{1}{2}M_p\left|[\sigma,\mathbf{q}]\right|\frac{1}{2}M^*_J\right\rangle\right]\frac{i(|\mathbf{q}|b)}{\sqrt{6}}\exp\left[-\frac{\mathbf{q}^2b^2}{6}\right](10,\tfrac{1}{2}M^*_J|J^*M^*_J)\ ,$$
$$j^0_{pN^*}=\frac{i}{\sqrt{6}}(|\mathbf{q}|b)\exp\left[-\frac{\mathbf{q}^2b^2}{6}\right](10,\tfrac{1}{2}M^*_J|J^*M^*_J)\delta_{M_pM^*_J}\ . \tag{54}$$

Here, $J^*=\frac{3}{2}$ for $N(1520)$ and $J^*=\frac{1}{2}$ for $N(1535)$, M^*_J and M_p are projections of J^* and proton spin, respectively, $\mathbf{e}_{\pm1}$ are two $\mathbf{q}$-transversal cyclic unit vectors, and b represents the quark oscillator parameter.

We see immediately that with the exception of the first term in the expression for transversal current $\mathbf{j}^T_{pN^*}$ the structure of others terms in the expressions for $\mathbf{j}^T_{pN^*}$ and $j^0_{pN^*}$ is the same as the structure of analogous terms in $\mathbf{j}^T_{pp}$ and j^0_{pp}. The literal identity has a place if we substitute

$$\hat{e}_p\to\left[\frac{m_p}{3m_q}\right]i\frac{(|\mathbf{q}|b)}{\sqrt{6}}\exp\left[-\frac{1}{6}b^2\mathbf{q}^2\right](10,\tfrac{1}{2}M^*_J|J^*M^*_J)\ , \tag{55}$$

$$\hat{\mu}_p\to\left[\frac{m_p}{m_q}\right]i\frac{(|\mathbf{q}|b)}{\sqrt{6}}\exp\left[-\frac{1}{6}b^2\mathbf{q}^2\right](10,\tfrac{1}{2}M^*_J|J^*M^*_J)\ . \tag{56}$$

As a result, the transversal and longitudinal form factors F^2_T and F^2_L in expression (40) for cross sections of processes $e+N(1520)\to e'+p$ and $e+N(1535)\to e'+p$ can be written as

$$F^2_T=\frac{1}{4\pi(2J^*+1)}\sum_{M^*_JM_p}\mathbf{j}^T_{pN^*}\mathbf{j}^{T^*}_{pN^*}$$
$$=\frac{1}{36\pi m_q^2}\left[\frac{1}{b^2}F^2_{1,N^*\to p}(\mathbf{q})+\left[\frac{9}{2}\mathbf{q}^2+\mathbf{K}^2_R-\frac{1}{\mathbf{q}^2}(\mathbf{K}_R\cdot\mathbf{q})^2\right]F^2_{2,N^*\to p}(\mathbf{q})\right]\ , \tag{57}$$

$$F_L^2=\frac{1}{4\pi(2J^*+1)}\sum_{M_J^* M_p} \mathbf{j}^0_{pN^*}\mathbf{j}^{0*}_{pN^*}=\frac{1}{4\pi}F^2_{2,N^*\to p}(\mathbf{q}) \ , \tag{58}$$

where the transitional form factors $F_{1,N^*\to p}(\mathbf{q})$ and $F_{2,N^*\to p}(\mathbf{q})$ are

$$F_{1,N^*\to p}(\mathbf{q})=\exp(-\tfrac{1}{6}b^2\mathbf{q}^2) \ , \tag{59}$$

$$F_{2,N^*\to p}(\mathbf{q})=\frac{(|\mathbf{q}|b)}{3\sqrt{2}}\exp\left[-\frac{1}{6}b^2\mathbf{q}^2\right] . \tag{60}$$

This kind of q-dependence is probably oversimplified, although it gave some reasonable results for the quark description of elastic $p\to p$ form factor (see above). A more adequate q dependence can be extracted in the future from the planned experiments on electroexcitation of $N(1520)$ and $N(1535)$ (the present data [21] are still rather fragmentary). Nevertheless, it is important that in the momentum transfer range $|\mathbf{q}|\sim 1.3-2$ GeV/c, which we are interested in, the values of deexcitation amplitudes $N(1520)\to p$ and $N(1535)\to p$ are comparable to these of elastic scattering $p\to p$ or even surpass them.

Relative motion functions $\Phi^d_{pN(1520)}(\mathbf{K}_R)$ and $\Phi^d_{pN(1535)}(\mathbf{K}_R)$ correspond to the $1p$ state, i.e., they are going to zero at $k_R\to 0$. The functions

$$\Phi^d_{N(1535)N(1535)}(\mathbf{K}_R) \ ,$$

$$\Phi^d_{N(1535)N(1520)}(\mathbf{K}_R) \ ,$$

and

$$\Phi^d_{N(1520)(1520)}(\mathbf{K}_R)$$

correspond, evidently, to the $0s$ state and have just maximum at $k_R=0$ (see Fig. 5 for this difference, which can display itself in the experiments under discussion only if the momentum of the knocked-out proton is high enough, i.e., a few GeV/c). Considering the above comments exposed after formula (60) and also taking into account the values of spectroscopic factors we conclude that in the quasielastic region $k_R\le 300$ MeV/c the terms $j^\nu_{pp}\Phi^d_{pN(1520)}$ and $j^\nu_{pp}\Phi^d_{pN(1535)}$ in Eqs. (52) and (53) can be neglected and the reactions ^{2}H$(e,e'p)N(1520)$ and ^{2}H$(e,e'p)N(1535)$ proceed here by means of the knockout of virtually excited baryon N^* accompanied by its deexcitation. Concerning Eq. (53) it has to be noted also that the spectroscopic factor for the $N(1520)N(1535)$ channel is one order of magnitude bigger than for the $N(1535)N(1535)$ one. Therefore, the last term in Eq. (53) predominates very much over the others, i.e., the ^{2}H$(e,e'p)N(1535)$ reaction in the kinematical region of the quasielastic peak proceeds through $N(1520)N(1535)$ component of deuteron with the elementary process $e+N(1520)\to e'+p$. As a consequence, the factorization approximation (17) is valid here; however, its realization is unusual as far as deexcitation cross section $e+N(1520)\to e'+p$ is picked out instead of the usual elastic one $e+p\to e'+p$. So we can present the cross section of the $^2N(e,e'p)N(1535)$ process as

$$\frac{d^2\sigma}{d\mathbf{k}'d\mathbf{K}_p}=\frac{\sigma_{eN(1520)\to e'p}(\theta_e,q_\mu,\mathbf{K}_p)}{\mathbf{k}'^2}\rho_{N(1520)N(1535)}(\mathbf{K}_R)\times\delta(E_f-E_i-\omega) \ , \tag{61}$$

where the momentum distribution is

$$\rho_{N(1520)N(1535)}(\mathbf{K}_R)=S^d_{N(1520)N(1535)}|\Phi_{0s}(K_R)|^2\times|Y_{00}(\hat{\mathbf{K}}_R)|^2 \ . \tag{62}$$

It has the typical shape of the S-state and is close to that presented in Fig. 4 (with no deexcitation it was of $1p$ shape [5]). The predicted cross section is expected to be approximately 10^{-3} of that for process ^{2}H$(e,e'p)n$ for both models examined.

Concerning ^{2}H$(e,e'p)N(1520)$ process it is difficult to pick out one of two last terms in Eq. (52) and their essential interference is expected.

As far as we know the model of the baryon-baryon structure of deuteron exposed here is the only one where the components with negative parity baryons appear, e.g., $NN(1520)$, $NN(1535)$, $N(1520)N(1520)$, etc. (and they appear on equal footing as the positive parity components). This circumstance, most probably, allows us to explain by a rather simple way the features of η-meson production in reactions $d+\gamma\to d+\eta$ and $n+p\to d+\eta$ [22] [η-decay of "spectator" $N(1535)$ plays the essential role here].

V. CONCLUSION

Let us summarize. We have examined the possibility of the unified microscopic six-quark description of various baryon-baryon components in the deuteron considering the importance of fast exclusive processes of nucleon quasielastic knock-out. In this way we were able to calculate the spectroscopic factors, i.e., the generalized "probabilities" to find these components and to calculate also momentum distributions of baryon-baryon relative motion here.

Of course, we may try to describe the baryon-baryon content of deuteron within the standard meson theory where various baryon-baryon components are generated only by means of meson exchanges. However, each separate component here needs its individual vertex constants and phenomenological form factors in each vertex πNN^*, ρNN^*, πN^*N^*, etc. As a result, this description has really no heuristic power [23].

In the quark microscopic approach all the parameters of qq interaction and of the quark bases are fixed, in principle, by the baryonic excitation spectrum taken in a broad energy range. It appears to be sufficient to calculate both six-quark wave functions of deuteron including its projections onto various baryon-baryon components and nucleon-nucleon phase shifts, etc. [2]. So the predictive power here seems to be rather convincing.

As a tool for investigating the above baryon-baryonic components we propose the process $^2\mathrm{H}(e,e'p)N^*$ to be studied with the knocked-out proton being fast and spectator N^* slow. This process offers the opportunity to extract from the exclusive experimental data both the spectroscopic factors and the momentum distributions. It is of importance here that the orbitally excited baryons N^* such as $N(1440)$, $N(1520)$, and $N(1535)$ are just originating from the excited quark configuration s^4p^2-$s^5 2s$ in the deuteron, while $^2\mathrm{H}(e,e'p)\Delta$ process in a very convincing way answers the question whether a $\Delta\Delta$-component (i.e., s^6 configuration) is present in the deuteron or not.

Very nonstandard shapes of momentum distributions are expected for the $N(1520)$ and $N(1535)$ spectator particles which is connected with an interesting interplay of various quark amplitudes. However, when extracting the necessary spectroscopic information from $^2\mathrm{H}(e,e'p)N^*$ quasielastic experimental data we have to take into account the "dangerous" competing diagram of Fig. 6. Here, the final-state baryon N^* does not reflect the corresponding baryon-baryon component in deuteron but reflects intermediate meson rescattering. Preliminary calculations by Yu. L. Dorodnykh and us, however, show that this triangle diagram gives a negligible contribution to the reaction amplitude if just the knocked-out proton is fast ($E_p > 1$ GeV) and the spectator N^* is slow, $k \leq 400-500$ MeV$/c$. This question is planned to be discussed in our next paper.

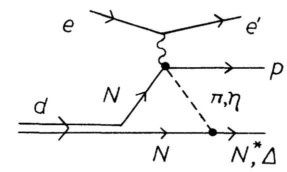

FIG. 6. Triangle diagram where N^* or Δ are born due to meson rescattering.

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