Quark-exchange Effects in a Deuteron Breakup at Intermediate Energy

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Abstract

Microscopical approach to a deuteron breakup at high and intermediate energies is proposed. We show that the quark exchange effects, resulting from the full asymmetry of the $6q$-deuteron wave function with respect to the pair permutations of quark variables, strongly affect the proton momentum distribution in the deuteron, as well as the polarization observables of inclusive deuteron breakup, when the “internal momentum” in the deuteron is of order of a few hundreds MeV/c.
1 Introduction

During the last decade the breakup of relativistic nonpolarized and polarized deuterons on hydrogen/nuclear target have been a subject of experimental and theoretical studies. They, as well as studies of ed-scattering at high $Q^2$, bring light on many aspects of the nucleon-nucleon interaction and the deuteron structure at short distances, including relativistic effects, effects of non-nucleon degrees of freedom, etc. The deuteron breakup experiments, we are discussing in this paper, include measurements of differential cross-section, $E_p \frac{d^3 \sigma}{dp}$, and polarization observables (the tensor analyzing power, $T_{20}$, and the polarization transfer coefficient, $\kappa_0$) of the $(d,p)$ reaction with the final proton registered at zero angle and the momentum $p \geq \frac{1}{2} \times$ (deuteron lab. momentum) at deuteron beam energy of a few GeV.

We think that there are at least three kinematical regions where different physical pictures determine the reaction. At the region of the “internal momentum” in the deuteron, $k \leq 0.1 \text{ GeV}/c$, the relevant degrees of freedom in the deuteron are nucleons and mesons. Here the nonrelativistic impulse approximation (IA) improved by the rescattering effects and the final state interactions, both calculated with the standard two-nucleon deuteron wave functions, seems to be valid. At the region of very high $k$ (according to estimations of [1] $k \geq 1 \text{ GeV}/c$) one has to provide a smooth transition to the perturbative QCD.

In the present paper we discuss the deuteron breakup at the intermediate region ($k$ between 0.1 and 1 GeV/c) where the nucleons in the deuteron are assumed to loose gradually their individuality, but the pure perturbative QCD does not yet work. At this region the deuteron structure can be described in terms of the constituent quarks and the chiral mesons within the dynamical scheme similar to that used in the baryon spectroscopy [2]. Then one should note that the Pauli principle at the constituent quark level gives rise to a number of the short-range baryon-baryon components ($NN^*$, $N^*N$, $N^*N^*$) of the deuteron wave function [3, 4].

However, instead of the usual Fock probabilities for the baryon-baryon components, which are applicable only for the structureless (or well separated baryons), one should use a language of ”effective numbers” [5] developed in nuclear cluster physics [5].

In this work we shall restrict ourselves to a simplified model, in which the quark antisymmetrization effects are taken into account for the deuteron $S$-wave only. Due to the centrifugal effect the quark-exchange contribution for the $D$-wave is estimated to be three orders of magnitude less than the one for the $S$-wave [4]. Besides, we do not consider here a delicate question of the relativistic internal momentum $k$ in the deuteron: we use it in the same way as for the equal-mass $NN$ configuration in the framework of the light-cone dynamics (“minimal relativistic prescription”). The $3q \times 3q$ decomposition of the deuteron six-quark wave function contains, apart from the equal-mass $NN$ configuration, a set of configurations, $NN^*$, whose “constituents” ($N$ and $N^*$, respectively) have different masses. Meanwhile, the main result of this work is that the quark exchanges between $3q$ clusters in the deuteron give a visible contribution in the differential cross-section and polarization observables of the $(d,p)$ breakup in the intermediate region, and have to be taken into consideration.
The outline of the paper is organized as follows. In Sec. 2 we consider a decomposition of the deuteron six-quark wave function into $3q \times 3q$ clusters and how it modifies a matrix element of the deuteron breakup. In Sec. 3 expressions for the proton momentum distribution in the deuteron and polarization observables are obtained. The results of numerical calculations and conclusions are given in Sec. 4.

![Deuteron inclusive breakup in spectator approximation.](image)

**Fig. 1.** Deuteron inclusive breakup in spectator approximation.

## 2 NN and NN* components of the deuteron six-quark wave function

The deuteron wave function, considered at small internucleon distances as a six-quark ($6q$) object, was shown [6, 7, 8] to be qualitatively equivalent to the Resonating Group Method (RGM) wave function

$$\psi^d(1, 2, \ldots, 6) = \hat{A} \{\varphi_N(1, 2, 3)\varphi_N(4, 5, 6)\chi(r)\},$$

(2.1)

where $\hat{A} = \frac{1}{\sqrt{10}}(1 - 9P_{36})$ is a quark antisymmetrizer and $\varphi_N(1, 2, 3)$ and $\varphi_N(4, 5, 6)$ are the wave functions of the nucleon three quark ($3q$) clusters; $\chi(r)$ is the RGM distribution function.

The presence of the quark antisymmetrizer $\hat{A}$ in expression (2.1) is the main difference between the microscopical quark and the meson-nucleon points of view on the deuteron structure at short distances. In the absence of the quark antisymmetrizer the RGM distribution function would coincide, up to the renormalization effects at short range, with the conventional deuteron wave function.

Due to the antisymmetrizer the deuteron wave function (2.1), being decomposed into $3q \times 3q$ clusters, includes, apart from the standard $pn$ component, the nontrivial $NN^*$, $N^*N$ and $N^*N^*$ components which respond to all possible nucleon resonance states, $N^*$ (see, e.g., [4]). So, in the spectator approximation (Fig.1) the reaction matrix element is

$$M_{dp \rightarrow p'X} = \sqrt{2} \sum_{B, M_B} \psi^d_{Bp'}(k_{p'})\Gamma_{Bp \rightarrow X},$$

(2.2)

where the summation over $B$ includes $pn$ and all $pN^*$ configurations of the deuteron; $M_B$ is $z$-projection of $B$-baryon spin. The other notations in (2.2) are as follows:
\[ \Gamma_{Bp \rightarrow X} \text{ is the amplitude of the reaction } B + p \rightarrow X : \]

\[
\Gamma_{pB \rightarrow X} = \int \prod_{i=1}^{3} \prod_{j=7}^{9} d^3r_i \ d^3r_j \ e^{-i(k_X + k_{p'} - k_p)r_X} \\
\times \langle \psi_X(1, 2, 3, 7, 8, 9) | \hat{\Gamma}(1, 2, 3, 7, 8, 9) | \varphi_B(1, 2, 3) \varphi_p(7, 8, 9) \rangle, \tag{2.3}
\]

\[ r_X = \frac{1}{6} \left( \sum_{i=1}^{3} r_i + \sum_{i=7}^{9} r_i \right), \tag{2.4} \]

and \( \varphi_B(1, 2, 3) \) is a wave function of the \( 3q \) cluster of the baryon \( B \); \( \varphi_N(7, 8, 9) \) is the same for the target proton; \( r \) stands for the \( i \)-th quark coordinate and \( \psi_{Bp'}(\mathbf{r}') \) is a Fourier transformation of the overlap, \( \psi_{Bp'}(\mathbf{r}) \), between the \( 6q \) wave function \( \psi_{Bp}(\mathbf{r}) \) of the deuteron and the \( pB \) wave function; \( \mathbf{r} \) stands for the relative coordinate between the target and the baryon \( B \):

\[
\psi_{Bp'}(\mathbf{r}) = \left( \frac{6!}{3!3!2} \right)^{1/2} \langle \varphi_B(1, 2, 3) \varphi_p'(4, 5, 6) | \psi_d(1, 2, 3, 4, 5, 6) \rangle, \tag{2.5}
\]

\[
r = \sum_{i=1}^{3} r_i - \sum_{i=4}^{6} r_i. \tag{2.6}
\]

In \cite{4} the expressions for the overlaps \( \psi_{Bp'}(\mathbf{k}_{p'}) \) were obtained for all \( pB \) configurations produced by the antisymmetrization of the \( S \)-wave component of the deuteron wave function. As the first step to an appropriate choice of the RGM distribution function, \( \chi(\mathbf{r}) \), we use the conventional \( NN \) deuteron wave function, \( \chi_{NN}(\mathbf{r}) \), consisting of the \( S \)- and \( D \)-components, and modify it according to the standard RGM renormalization condition \cite{3, 10}:

\[
\chi(\mathbf{r}) = \int \hat{\mathcal{N}}^{-1/2}(\mathbf{r}, \mathbf{r'}) \chi_{NN}(\mathbf{r'}) d^3r', \tag{2.7}
\]

where \( \hat{\mathcal{N}}(\mathbf{r}, \mathbf{r'}) \) is a norm operator

\[
\hat{\mathcal{N}}(\mathbf{r}, \mathbf{r'}) = \delta(\mathbf{r} - \mathbf{r'}) - 9 \{ \varphi_N(1, 2, 3) \varphi_N(4, 5, 6) \}_{ST=10} \delta(\mathbf{r} - \mathbf{r}'') \\
\times \hat{\mathcal{P}}_{36}( \{ \varphi_N(1, 2, 3) \varphi_N(4, 5, 6) \}_{ST=10} \delta(\mathbf{r}' - \mathbf{r}'') ). \tag{2.8}
\]

Fig.2 displays the ratios \( \chi^{L=0}(r)/\chi_{NN}^{L=0}(r) \) and \( \chi^{L=2}(r)/\chi_{NN}^{L=2}(r) \). In the further calculations we approximate the radial parts of the \( S \)- and \( D \)-components of the RGM distribution function by gaussians:

\[
\frac{\chi^{L=0}(r)}{r} = \sum_k A_k e^{-\alpha_k r^2}, \quad \frac{\chi^{L=2}(r)}{r} = r^2 \sum_k B_k e^{-\beta_k r^2}. \tag{2.9}
\]

As the nucleon and excited states wave functions we have used the wave functions of the Translationally Invariant Shell Model \cite{11} for three particles, i.e. the harmonic oscillator wave functions with oscillator parameter \( b \), with the center-of-mass oscillations
being removed. All details can be found in \([4]\). The parameters of the approximation (2.10) for the RGM distribution function based on the Paris potential are given in Table 1. Making use the gaussian approximation all necessary calculations could be performed analytically \([4]\) and the Fourier transformation (2.5) is given by

\[
\psi^{d}_{Bp'}(k_{p'}) = \chi(k_{p'})\delta_{B,N} + 3 \sum_{m_B,\mu_B} \langle L_B S_B m_B \mu_B \mid J_B M_B \rangle \frac{1}{\sqrt{\text{dim}[f_B]}} \times \gamma^X_B \sum_{k} A_{k} I_{N_B L_B,00}^{L_B}(k_{p'}; \alpha_k) Y_{L_B m_B}^{*}(\hat{k}_{p'}) (-1)^{\frac{1}{2} + S_B + 2T_B} \times \sqrt{(2T_B + 1)(2S_B + 1)} \langle S_B \frac{1}{2} \mu_B \mu_{p'} \mid 1M_d \rangle \times \sum_{S_{12}=T_{12}=0.1, S_{45}=T_{45}=0.1} \langle [f_B] S_B T_B \mid [2] S_{12} T_{12} \mid \frac{1}{2} \frac{1}{2} \rangle \times (-1)^{S_{12} + S_{45}} \left\{ \begin{array}{c} \frac{1}{2} \frac{1}{2} \\ T_{12} \\ T_{45} \end{array} \right\} \left\{ \begin{array}{c} \frac{1}{2} \frac{1}{2} \\ S_{45} \frac{1}{2} \frac{1}{2} \\ S_{12} \frac{1}{2} \frac{1}{2} \end{array} \right\}. \quad (2.11)
\]

For the explanation of notations in (2.11), see \([4]\). The functions \(I_{N_B L_B,00}^{L_B}\) are summarized in Appendix A. The fractional parentage coefficients (CFP), \(\gamma^X_B\), for \(N_B = 3\) could be found in Appendix B (the CFP for \(N_B = 0, 1, 2\) are given in \([4]\). The physical meaning of the parameter \(b\) appearing in these functions is the m.s. radius of a nucleon quark core, \(i.e.\) a characteristic distance where quark exchanges becomes noticeable.

In Table 2 we summarize the baryons \(B\) which are taken into account in our calculations and give the effective numbers, \(N^d_{Bp}\), for three values of the parameter \(b\); for the definition of the effective number \(N^d_{Bp}\) see, \(e.g.\) \([4]\). The baryons incorporated in the components with the effective numbers \(N^d_{Bp} < 10^{-5}\) were ignored.

It should be mentioned here that the baryons from lines 2-5 and 7-10 of Table 2, have the negative parity and thus produce effective \(P\)-waves of the relative motion in the deuteron. It was already mentioned (\([12]\) and \([13]\)) that the \(P\)-waves should strongly affect polarization phenomena of the reaction at the region of internal momentum in the deuteron of few hundreds MeV/c.

There are at least two problems connected to the relativistic description which we would like to discuss here shortly. First, RGM is based on nonrelativistic dynamical equations. To establish a correspondence between the observed momentum of the spectator and the wave function argument we use the “minimal relativization prescription”. This means that we identify the light-cone variable \(k\) with the relative momentum of the \(3q\) clusters in the deuteron, considering their dynamics nonrelativistically, \(but\) in terms of new variable \(k\) (see, \(e.g.\), \([14]\)). Second, the different components of the deuteron wave function (\(NN, NN^*\ etc.) consists of “particles” (\(3q\) clusters) with different masses. In this case every component, involved in our calculations, in the limit of infinite distances should depend on its own “internal momentum”. However such a limit is allowed only for the usual \(n-p\) component and the effective \(NN^*\) components appear only in the nucleon overlap region and they are direct consequence of the underlying quark structure of a baryon. Thus this question cannot be solved at the baryon
In the present work we assume for simplicity that the "internal momenta" in the n-p and $N - N^*$ channels are the same.

3 Momentum distribution of the protons in the deuteron and polarization observables in the spectator approximation

According to the spectator approximation (2.2) the reaction matrix element squared is

$$\sum_X |M_{dp \rightarrow p'X}|^2 = 2\sigma_0 \sum_B \left| \psi_{Bp'}^d(k_{p'}) \right|^2.$$  \hspace{1cm} (3.12)

In (3.12) it is assumed that

$$\sum_X \Gamma_{pB' \rightarrow X}^\dagger \Gamma_{pB \rightarrow X} = \sigma_0 \delta_{BB'},$$  \hspace{1cm} (3.13)

where $\sigma_0$ is of order of the total $NN$ cross section. This means that the size of $B$ is supposed to be the same as the proton’s one and inelasticities in baryon-baryon collisions are small at the energy scale of few GeV. Now the proton momentum distribution in the deuteron is given by

$$n(k_{p'}) = \bar{u}^2 + \bar{w}^2 + u_1^2 + v_1^2 + \frac{1}{2} v_2^2 + v_3^2 + \frac{1}{2} v_4^2 + 2 v_5^2 + v_6^2 + 2 v_7^2 + v_8^2,$$  \hspace{1cm} (3.14)

where

$$\bar{u} = \chi_{L=0}^{L=0}(k_{p'}) - \frac{1}{9} \sum_k A_k I_{0,0,0}^0(k_{p'}; \alpha_k),$$  \hspace{1cm} (3.15)

$$\bar{w} = \chi_{L=2}^{L=2}(k_{p'}),$$  \hspace{1cm} (3.16)

$$u_1 = \frac{4}{9} \sum_k A_k I_{2,0,0}^0(k_{p'}; \alpha_k),$$  \hspace{1cm} (3.17)

$$v_1 = v_2 = v_5 = v_6 = \frac{4}{9} \sum_k A_k I_{1,1,0}^1(k_{p'}; \alpha_k),$$  \hspace{1cm} (3.18)

$$v_3 = v_4 = v_7 = v_8 = \frac{2}{9} \sum_k A_k I_{3,1,0}^1(k_{p'}; \alpha_k).$$  \hspace{1cm} (3.19)

The differential cross section of the reaction is proportional to the proton momentum distribution, $n(k_{p'})$, (see, e.g., [14]) and the tensor analyzing power, $T_{20}$, and polarization transfer coefficient, $\kappa_0$, are given by

$$T_{20} = \frac{1}{\sqrt{2}} \frac{2 \sqrt{2} \bar{u} \bar{w} - \bar{w}^2 - \frac{1}{2} v_2^2 - \frac{1}{2} v_1^2 + \frac{4}{9} v_6^2 + \frac{4}{9} v_8^2}{n(k_{p'})},$$  \hspace{1cm} (3.20)

$$\kappa_0 = \frac{(\bar{u}^2 - \bar{w}^2 - \sqrt{1/2} \bar{u} \bar{w} + v_1^2 + v_2^2 - 1/2 v_4^2 + v_3^2 - 1/2 v_4^2 + 2 v_5^2 - \frac{1}{10} v_6^2 + 2 v_7^2 - \frac{1}{10} v_8^2)}{n(k_{p'})}.$$  \hspace{1cm} (3.21)
4 Numerical calculations and conclusions

The results of our calculations and comparison with the experimental data are shown in Figs. 3-9. Figs. 5 and 7 explore the dependence of the polarization observables $T_{20}$ and $\kappa_0$ on the parameter $b$ of the oscillator quark potential. In Figs. 3, 4, 6 and 8 we demonstrate insensitivity of the current approach on the choice of input $NN$ potential.

In Fig. 9 we compare the results with the experimental data on the $T_{20} - \kappa_0$ plane. It was recently shown [12], that a correlation between $T_{20}$ and $\kappa_0$ should give important information about the non-nucleon degrees of freedom in the deuteron. Particularly, in the case, when the deuteron wave function consists of the two standard ($S-$ and $D-$wave) components only, the parametric curve describing the correlation between $T_{20}$ and $\kappa_0$ must lie on the circle. Our calculations give an example of a deformation of this relation produced by additional components of the deuteron wave function.

The points for the proton momentum distribution in the deuteron $n(k_p')$ (Figs. 3 and 4) were extracted from data for the $p(d,p)$ reaction cross section [15] (see also this reference, as well as [16], [17] for $A(d,p)$ data). $T_{20}$ was measured in [17]-[22] and $\kappa_0$ in [21, 23, 24, 25].

The conclusions of the present work are as follows:
(i) When the internal momentum in the deuteron is of a few hundreds MeV/c, the effects of the quark exchange between three-quark clusters in the deuteron are as important as the ones originating from the relative motion of the nucleons (estimated by the standard two-nucleon potentials).
(ii) The $pN^*$ components produce effective $P$-waves in the deuteron which correct the momentum distribution of protons in the deuteron and polarization observables of the deuteron breakup at high/intermediate energy in the “right” direction.
(iii) The results of the current approach depend weekly on the choice of an input two-nucleon potential, but, at the same time, very sensitive to the quark core radius of the nucleon, $b$. However, it should be noted the value $b = 0.8$ fm which provide the best fit to data is somewhat larger than the m.s.r. of the quark core of a free nucleon, $b = 0.5 \pm 0.6$ fm, commonly used in theoretical calculations.

It should be also stressed that we consider here the intermediate region between the pure $NN$ and the perturbative QCD regimes and our model cannot be used for a region of very high $k$. 
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Appendix A

$$I_{00,00}^0 = \left(\frac{18 + 30\alpha k b^2}{18 + 30\alpha k b^2}\right)^{\frac{3}{2}} b^3 \exp\left\{ -\frac{b^2 k^2}{12} \left(\frac{15 + 16\alpha k b^2}{3 + 5\alpha k b^2}\right) \right\}$$

$$I_{11,00}^1 = I_{00,00}^0 \ g_1 \ b \ k$$

$$I_{20,00}^0 = I_{00,00}^0 \ \sqrt{6} \left( g - \frac{3}{2} g_1 g + \frac{g_1^2}{2} b^2 k^2 \right)$$

$$I_{22,00}^2 = I_{00,00}^0 \ \sqrt{\frac{3}{g_1^2 b^2 k^2}}$$

$$I_{31,00}^1 = I_{00,00}^0 \ \sqrt{10} g_1 \ b \ k \left( g - \frac{3}{2} g_1 g + \frac{3g_1^2}{10} b^2 k^2 \right),$$

where

$$g = \frac{4\alpha k b^2 - 3}{15 + 16\alpha k b^2}, \quad g_1 = \frac{4\alpha k b^2 - 3}{18 + 30\alpha k b^2}.$$

Appendix B.

Baryon orbital wave functions $|3(\lambda \mu)[f] L(r)>$

The choice of Jacobi coordinates is

$$r_{12} = \frac{r_1 - r_2}{\sqrt{2}}, \quad \rho = \frac{r_1 + r_2 - 2r_3}{\sqrt{6}}.$$

$$|3(30)[3]1(111)> = \frac{1}{2} \left[ \sqrt{\frac{5}{3}} \varphi_{200}(r_{12}) \varphi_{11M}(\rho) + \frac{1}{\sqrt{3}} \sum_{m_1, m_2} \langle 21 \ m_1 m_2 | 1M \rangle \times \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho) - \frac{1}{2} \varphi_{000}(r_{12}) \varphi_{31M}(\rho) \right]$$

$$|3(30)[21]1(112)> = -\frac{\sqrt{5}}{6} \varphi_{200}(r_{12}) \varphi_{11M}(\rho) - \frac{1}{3} \sum_{m_1, m_2} \langle 21 \ m_1 m_2 | 1M \rangle \times \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho) - \frac{\sqrt{3}}{2} \varphi_{000}(r_{12}) \varphi_{31M}(\rho)$$
\[
|3(30)[21]1(121)\rangle = \frac{\sqrt{3}}{2} \varphi_{31M}(r_{12}) \varphi_{000}(\rho) - \frac{\sqrt{5}}{6} \varphi_{11M}(r_{12}) \varphi_{200}(\rho) \\
- \frac{1}{3} \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 1M \rangle \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]

\[
|3(11)[21]1(121)\rangle = \frac{2}{3} \varphi_{200}(r_{12}) \varphi_{11M}(\rho) - \frac{\sqrt{5}}{3} \sum_{m_1,m_2} \langle 21 \ m_1m_2 \ 1M \rangle \\
\times \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho)
\]

\[
|3(11)[21]1(121)\rangle = \frac{2}{3} \varphi_{11M}(r_{12}) \varphi_{200}(\rho) - \frac{\sqrt{5}}{3} \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 1M \rangle \\
\times \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]

\[
|3(30)[1^3][1(123)]\rangle = \frac{1}{2} \varphi_{31M}(r_{12}) \varphi_{000}(\rho) - \frac{1}{2} \sqrt{\frac{5}{3}} \varphi_{11M}(r_{12}) \varphi_{200}(\rho) \\
- \frac{1}{\sqrt{3}} \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 1M \rangle \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]

\[
|3(11)[21]2(112)\rangle = - \sum_{m_1,m_2} \langle 21 \ m_1m_2 \ 2M \rangle \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho)
\]

\[
|3(11)[21]2(121)\rangle = \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 2M \rangle \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]

\[
|3(30)[3][3(111)]\rangle = \frac{\sqrt{3}}{2} \sum_{m_1,m_2} \langle 21 \ m_1m_2 \ 3M \rangle \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho) \\
- \frac{1}{2} \varphi_{000}(r_{12}) \varphi_{33M}(\rho)
\]

\[
|3(30)[21]3(112)\rangle = - \frac{1}{2} \sum_{m_1,m_2} \langle 21 \ m_1m_2 \ 3M \rangle \varphi_{22m_1}(r_{12}) \varphi_{11m_2}(\rho) \\
- \sqrt{\frac{3}{4}} \varphi_{000}(r_{12}) \varphi_{33M}(\rho)
\]

\[
|3(30)[21]3(121)\rangle = - \sqrt{\frac{3}{4}} \varphi_{33M}(r_{12}) \varphi_{000}(\rho) \\
- \frac{1}{2} \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 3M \rangle \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]

\[
|3(30)[1^3][3(123)]\rangle = \frac{1}{2} \varphi_{33M}(r_{12}) \varphi_{000}(\rho) \\
- \frac{\sqrt{3}}{2} \sum_{m_1,m_2} \langle 12 \ m_1m_2 \ 3M \rangle \varphi_{11m_1}(r_{12}) \varphi_{22m_2}(\rho)
\]
The parameters of the approximation (2.10) for the RGM distribution function based on the Paris potential. The oscillator parameter $b = 0.8 \text{ fm}$.

| $n$ | $A_n \text{ (fm}^{-3/2}\text{)}$ | $\alpha_n \text{ (fm}^{-2}\text{)}$ | $B_n \text{ (fm}^{-7/2}\text{)}$ | $\beta_n \text{ (fm}^{-2}\text{)}$ |
|-----|-------------------------------|----------------|----------------|----------------|
| 1   | $1.3863 \cdot 10^{-2}$        | $8.845 \cdot 10^{-3}$ | $3.1946 \cdot 10^{-4}$ | $4.1231 \cdot 10^{-2}$ |
| 2   | $7.1097 \cdot 10^{-2}$        | $3.1668 \cdot 10^{-2}$ | $3.6241 \cdot 10^{-3}$ | $1.1980 \cdot 10^{-4}$ |
| 3   | $1.7493 \cdot 10^{-1}$        | $1.0619 \cdot 10^{-1}$ | $1.7823 \cdot 10^{-2}$ | $2.6855 \cdot 10^{-4}$ |
| 4   | $2.9944 \cdot 10^{-1}$        | $3.8459 \cdot 10^{-1}$ | $7.5863 \cdot 10^{-2}$ | $5.5486 \cdot 10^{-4}$ |
| 5   | $-6.8075 \cdot 10^{-1}$       | $2.4318$           | $3.3025 \cdot 10^{-1}$ | $1.3350$           |
| 6   | $1.8345 \cdot 10^{-1}$        | $3.5360$           | $-2.6708 \cdot 10^{-1}$ | $5.4722$           |

The effective numbers $N_{Bp}^d$ of different baryon-proton configurations in the deuteron. The calculations were done for the Paris potential.

| $J^P\ell$ | State | $N_{Bp}^d$ | $b = 0.5 \text{ fm}$ | $b = 0.7 \text{ fm}$ | $b = 0.8 \text{ fm}$ |
|-----------|-------|-------------|----------------|----------------|----------------|
| 1 $1^+\frac{1}{2}\frac{1}{2}$ | $|0(00)[3][0\frac{1}{2}\frac{1}{2}]\rangle$ | 0.982 | 0.996 | 1.005 |
| 2 $3^-\frac{1}{2}\frac{1}{2}$ | $|1(10)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $1.52 \times 10^{-3}$ | $3.60 \times 10^{-3}$ | $4.46 \times 10^{-3}$ |
| 3 $1^-\frac{1}{2}\frac{1}{2}$ | $|1(10)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $7.59 \times 10^{-4}$ | $1.80 \times 10^{-3}$ | $2.23 \times 10^{-3}$ |
| 4 $1^-\frac{1}{2}\frac{1}{2}$ | $|1(10)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $3.79 \times 10^{-4}$ | $9.01 \times 10^{-4}$ | $1.16 \times 10^{-3}$ |
| 5 $3^-\frac{1}{2}\frac{1}{2}$ | $|1(10)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $7.59 \times 10^{-4}$ | $1.80 \times 10^{-3}$ | $2.23 \times 10^{-3}$ |
| 6 $1^+\frac{1}{2}\frac{1}{2}$ | $|2(20)[21][0\frac{1}{2}\frac{1}{2}]\rangle$ | $2.32 \times 10^{-3}$ | $5.47 \times 10^{-3}$ | $6.75 \times 10^{-3}$ |
| 7 $1^-\frac{1}{2}\frac{1}{2}$ | $|3(30)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $2.52 \times 10^{-4}$ | $3.63 \times 10^{-4}$ | $3.67 \times 10^{-4}$ |
| 8 $3^-\frac{1}{2}\frac{1}{2}$ | $|3(30)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $5.04 \times 10^{-4}$ | $7.25 \times 10^{-4}$ | $7.33 \times 10^{-4}$ |
| 9 $1^-\frac{1}{2}\frac{1}{2}$ | $|3(30)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $1.26 \times 10^{-4}$ | $1.81 \times 10^{-4}$ | $1.83 \times 10^{-4}$ |
| 10 $3^-\frac{1}{2}\frac{1}{2}$ | $|3(30)[21][1\frac{1}{2}\frac{1}{2}]\rangle$ | $2.52 \times 10^{-4}$ | $3.63 \times 10^{-4}$ | $3.67 \times 10^{-4}$ |
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Fig. 2. Ratios $\chi^{L=0}(r)/\chi^{L=0}_{NN}(r)$ and $\chi^{L=2}(r)/\chi^{L=2}_{NN}(r)$. Calculations are performed with the Paris potential and the parameter $b = 0.8 \text{ fm}$.

Fig. 3. Momentum distribution of the proton in the deuteron. Solid line stands for calculations within the current approach with $b = 0.8 \text{ fm}$, dashed and short-dashed lines – to $b = 0.5$ and 0.7 fm, respectively. Calculations are performed with the Paris potential. The points are extracted from $^1H(d,p)$ cross section data [13].

Fig. 4. Momentum distribution of the proton in the deuteron. Solid and dashed lines stand for calculations within the current approach with $b = 0.8 \text{ fm}$, with Paris and Reid-soft-core potentials, respectively. By short-dashed line we denote the curve obtained in the IA with Paris potential.

Fig. 5. $T_{20}$ in the current approach. Solid line corresponds to $b = 0.8 \text{ fm}$, dashed and short-dashed lines – to $b = 0.7$ and 0.5 fm, respectively. Calculations are performed with the Paris potential. The experimental data are from [17]-[19].

Fig. 6. $T_{20}$ in the current approach ($b = 0.8 \text{ fm}$), with the Paris potential (solid line) and the Reid-soft-core potential (short-dashed line). Dashed line stands for the IA calculations with the Paris potential. The experimental data are from [17]-[19].

Fig. 7. The same as in Fig. 5, for $\kappa_0$. The experimental data are from [21, 23, 24].

Fig. 8. The same as in Fig. 6, for $\kappa_0$.

Fig. 9. The $\kappa_0$ vs $T_{20}$ plot in the current approach (dashed line). The solid line represents the results of IA [12]. Calculations are performed with the Paris potential. Data are from [23].
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S-wave

D-wave

$r$ (fm)
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Fig. 3

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Fig. 7

The graph depicts the relationship between $\kappa_0$ and $k$ (GeV/c) for various datasets.

- Cheung et al. (diamonds)
- Nomofilov et al. (crosses)
- Dzikowski et al. (triangles)
- Kuehn et al. (squares)

The data points are accompanied by error bars, indicating the uncertainty in the measurements.

The axes are labeled as $\kappa_0$ on the vertical axis and $k$ (GeV/c) on the horizontal axis.
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Fig. 9