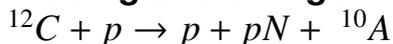


Probing short-range NN-correlations in the reaction



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Abstract. The exclusive reaction $^{12}\text{C}(p, ppN)^{10}\text{A}$ of quasi-elastic knockout of the nucleon from the short-range correlated nucleon pair $\langle NN \rangle$ in the ^{12}C nucleus by proton with an energy of few GeV is considered in the plane-wave approximation. The two-nucleon spectroscopic factors and momentum distributions of the center of mass motion of NN -pairs are considered by analogy with theory of quasi-elastic knockout of fast deuteron clusters from nuclei (p, Nd) using the translationally-invariant shell model.

1 Introduction

Quasi-elastic knock-out of fast deuterons from ^{12}C nucleus by protons at 670 MeV discovered in Dubna in 1957 [1] had demonstrated surprisingly high cross section. The transfer of high kinetic energy to such a weakly bound system as a deuteron lead D. Blokhintsev to a key idea about fluctuations of nuclear matter density in nuclei [2] later on called as fluctons [3]. Now fluctons are considered either like multi-quark configurations or as few-nucleon correlations. Two nucleons being at short relative distances $r_{NN} \sim 0.5$ fm due to the uncertainties principle have a high relative momentum $q_{rel} \sim 0.4$ GeV/c and the repulsive core in the NN-interactions which nature is still not quite clear, enhances this effect. Two-nucleon short-range correlations (SRC) in nuclei are actively studied during few last decades [4]. Under the SRC pair is assumed a pair of nucleons with small momentum of their center mass motion $k_{c.m.}$ and large as compared the Fermi momentum in heavy nuclei $p_F = 250 - 300$ MeV/c and oppositely directed individual momenta of the nucleons $\mathbf{p}_1 \approx -\mathbf{p}_2$. Experiments with electron and proton beams show that such correlations do exist in nuclei and the probability to find in nucleus the SRC pn pair is about 20 times higher than the pp - and nn -pair [5]. The dominance of the pn SRC pair is connected with the tensor forces acting in spin-triplet NN pair and vanishing in the spin-singlet 1S_0 states of the pp - and nn pairs. Furthermore, factorization over the $k_{c.m.}$ and q_{rel} momenta takes the place at enough large q_{rel} and low $k_{c.m.}$ in momentum distribution of the SRC: $n(\mathbf{p}_1, \mathbf{p}_2) \approx C_A n_{c.m.}(\mathbf{k}_{c.m.}) n_{rel}(\mathbf{q}_{rel})$, where C_A is a smooth function of number of nucleons A . For broad class of nuclei from ^4He to ^{208}Pb the distribution over internal momentum \mathbf{q}_{rel} at high $|\mathbf{q}_{rel}|$ is an universal function close to the deuteron wave function squared. The measured distribution $n_{c.m.}(\mathbf{k}_{c.m.})$ is compatible with the tree-dimensional Gaussian with the parameter $\sigma = 140 - 160$ MeV/c (see Ref. [6] and references therein).

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A new experiment for study of the SRC in the ^{12}C nucleus was performed at the BM@N in JINR [7] and analysis of the collected data is in progress now. In this experiment an inverse kinematics is used, where the ^{12}C beam with momentum 4 GeV/c per nucleon interacts with the liquid hydrogen target providing the transition $^{12}\text{C} + p \rightarrow pp + N + ^{10}\text{A}$. All free final nucleons and residual nucleus ^{10}A are registered at conditions when the missing momentum is large, $p_{miss} > 250$ MeV/c, where \mathbf{p}_{miss} is the momentum of the nucleon which is knocked out by the target proton from the SRC pair. The scattering angle in the center mass frame of the subprocess $pN \rightarrow pN$ is chosen to be close to 90° , $\theta_{cm} = 90^\circ \pm 30^\circ$, in order to minimize the off-mass-shell effects in the pp -scattering. From theory side, this reaction was analyzed in Ref. [8] within the plane-wave approximation corresponding to the pole diagrams depicted in Fig.1 and the corresponding formalism was presented. In the developed formalism the two-nucleon spectroscopic factors in the ^{12}C nucleus are calculated using the translationally-invariant shell model with mixing configurations. The short-range correlations are introduced by replacement of the internal wave function of the two-nucleon cluster by the deuteron wave function $\psi_d(q_{rel})$ for the spin-triplet states. For the spin-singlet NN -pairs the t-matrix of NN scattering in the 1S_0 can be used for realistic NN -interaction potentials. Since the internal momentum q_{rel} in the SRC pair is large, the relativistic effects are necessary to take into account. These effects are considered in [8] within the light-front dynamics.

Here we are focused mainly on the distribution on the momentum of the center mass motion of the SCR pair in the initial nucleus. A necessary part of the mathematical formalism [8] is given in the next section together with the obtained numerical results. Concluding remarks are given in the last section.

2 The model

Considering the contribution of the SRC pairs to the reaction in question we use as a basis the results of analysis of the reactions of quasi-elastic knock-out of fast nucleon clusters from nuclei by protons. For our purposes the most important are data on the reactions of quasielastic knock-out of the deuterons $A(p, pd)B$ from light nuclei $A = ^6\text{Li}$, [9], $A = ^7\text{Li}$, ^{12}C [10] and on the reaction $^{6,7}\text{Li}(p, nd)B$ obtained in JINR with the proton beam of 670 MeV. The reaction (p, nd) presents a special interest since it corresponds to the kinematics of the backward quasi-elastic scattering on the di-neutron $\langle nn \rangle$, $p + \langle nn \rangle \rightarrow n + d$, where the spin-isospin state of the $\langle nn \rangle$ system differs from that for the deuteron. Theory of these reactions was developed in the framework of the shell-model model for nuclear structure using realistic mechanisms

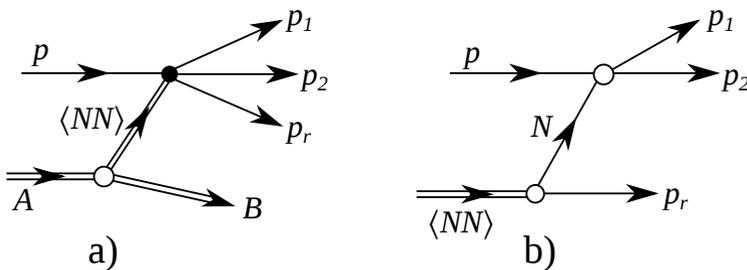


Figure 1. The pole mechanism of the reaction $A(p, 3N)^{10}B$ (a) and subprocess $p\langle NN \rangle \rightarrow pNN$ (b).

for the $p\langle NN \rangle$ -backward elastic or quasi-elastic scattering [9–12] and the inverse reaction $pd \rightarrow \langle pp \rangle + n$ [13, 14]. The spectroscopic factors obtained on this way were found to be

compatible with shell-model predictions. For this reason in our study of the SRC we use the translationally-invariant shell model for calculation of the two-nucleon spectroscopic factors. In order to account for the short-range NN-correlations we replace the shell model wave function of the NN-pair at large relative momenta between nucleons by the deuteron (or dineutron/diproton) wave function.

2.1 Elements of formalism

The transition matrix element corresponding to the Feynman diagram in Fig. 1a has the following form

$$M_{fi} = M(A \rightarrow B + \langle NN \rangle) \frac{1}{p_{\langle NN \rangle}^2 - m_{\langle NN \rangle}^2 + i\epsilon} M(p \langle NN \rangle \rightarrow pNN) \quad (1)$$

where $M(A \rightarrow B + \langle NN \rangle)$ is the amplitude of the virtual decay of the nucleus A into $\langle NN \rangle$ pair and the residual nucleus B in certain internal states and definite state of their relative motion, $p_{\langle NN \rangle}^2 - m_{\langle NN \rangle}^2 + i\epsilon^{-1}$ is the propagator of the $\langle NN \rangle$ -pair with four-momentum $p_{\langle NN \rangle}$ and the mass $m_{\langle NN \rangle}$, $M(p \langle NN \rangle \rightarrow pNN)$ is the matrix element of the subprocess of quasielastic knock-out of the nucleon N from the $\langle NN \rangle$ pair. All tree factors in Eq. (1) are relativistic invariant and therefore can be calculated in any suitable reference frame. The amplitude $M(A \rightarrow B + \langle NN \rangle)$ can be calculated in the rest frame of the nucleus A and in this case it can be expressed as a product of the spectroscopic factor S_A^x of the $\langle NN \rangle$ cluster, which we denote below as x , and the wave function of the relative motion of the center mass of the cluster and the center mass of the residual nucleus $\Psi_{\nu\Lambda M_\Lambda}(\mathbf{p}_B)$ in the following form

$$\frac{M(A \rightarrow B + \langle NN \rangle)}{p_{\langle NN \rangle}^2 - m_{\langle NN \rangle}^2 + i\epsilon} = S_A^x \Psi_{\nu\Lambda M_\Lambda}(\mathbf{p}_B) \sqrt{2m_A 2m_B / 2m_{\langle NN \rangle}}, \quad (2)$$

where \mathbf{p}_B is the 3-momentum of the residual nucleus B in the rest frame of the nucleus A , ν is the number of oscillator quanta, Λ is the orbital angular momentum and M_Λ is its z-projection, m_i is the mass of the nucleus i ($i = A, B$). The spectroscopic factor S_A^x is determined by the overlap integral between the wave function of the nucleus A , ψ_A , and the product of the wave functions ψ_B , ψ_x and $\psi_{\nu\Lambda}$ as (for detail see, for example, [15])

$$S_A^x = \left(\begin{matrix} A \\ x \end{matrix} \right)^{1/2} \langle \psi_A | \psi_B \psi_{\nu\Lambda}(\mathbf{R}_{A-x} - \mathbf{R}_x) \psi_x \rangle, \quad (3)$$

where the combinatoric factor $\left(\begin{matrix} A \\ x \end{matrix} \right)^{1/2}$ takes into account identity of the nucleons. The wave function of the nucleus A in the translationally-invariant shell model (TISM) [16] is

$$\psi_A = |AN[f](\lambda\mu)\alpha LSTJM_J M_T \rangle, \quad (4)$$

where N is the total number of oscillator quanta, $[f]$ is the Young scheme determining the permutational symmetry of the orbital part of the wave function, $(\lambda\mu)$ is the Elliott symbol defining the $SU(3)$ symmetry, L, S, J and T are the orbital, spin, isospin and total angular and isospin momenta, respectively; M_J and M_T are the z-projections of the total angular momentum J and of the isospin T ; the α is required for a complete labeling of the nuclear states. The TISM wave function with mixing configurations with different $[f]$, L , and S but the same J and T is the following superposition of the wave functions given in Eq.(4):

$$\psi_A^{JT} = \sum_{[f]LS} \alpha_{[f]LS}^{A,JT} |AN[f](\lambda\mu)\alpha LSTJM_J M_T \rangle. \quad (5)$$

In the TISM the parentage expansion of the antisymmetric wave function $|AN_A\alpha\rangle$ over the product of the antisymmetric wave functions of the $A-x$ ($\equiv B$) nucleus, $|BN_B\beta\rangle$, and the antisymmetric wave function of the cluster x , $|xN_x\gamma\rangle$, and the wave function of their relative motion $\psi_{\nu\Lambda M_\Lambda}$ has the following form

$$|AN_A\alpha\rangle = \hat{A} \left\{ |A-xN_B\beta\rangle |xN_x\gamma\rangle | \nu\Lambda M_\Lambda \rangle \right\} = \sum_{N_B\beta N_x\gamma \nu\Lambda M_\Lambda} \langle AN_A\alpha | A-xN_B\beta, \nu\Lambda M_\Lambda, xN_x\gamma \rangle \times |A-xN_B\beta\rangle |xN_x\gamma\rangle | \nu\Lambda M_\Lambda \rangle, \quad (6)$$

where \hat{A} is the properly normalized antisymmetrization operator in respect of permutation of the nucleons between the systems $A-x$ and x , $\langle AN_A\alpha | A-xN_B\beta, \nu\Lambda M_\Lambda, xN_x\gamma \rangle$ is the parentage coefficient. The following relation takes the place in Eq.(6) for numbers of the oscillator quanta:

$$N_A = N_B + N_x + \nu. \quad (7)$$

Using the TISM model for the nuclear wave functions of nuclei A , B ($\equiv A-x$) and the cluster x the spectroscopic factor (3) can be calculated analytically and after that the transition matrix element (1) can be written as

$$M_{fi}(pA \rightarrow 3NB) = \left(\frac{A}{2} \right)^{1/2} \sum_{\alpha_i, \alpha_f, \nu, \Lambda, \mathcal{L}} \sum_{M_{J_x}, \bar{J}, \bar{M}, M_\Lambda} \alpha_i^{A J_i T_i} \alpha_f^{A-2J_f T_f} \langle AN_A\alpha | A-xN_B\beta, \nu\Lambda M_\Lambda, xN_x\gamma \rangle U(\Lambda L_x \bar{J} S_x; \mathcal{L} J_x) \begin{Bmatrix} L_f & S_f & J_f \\ \mathcal{L} & S_x & \bar{J} \\ L_i & S_i & J_i \end{Bmatrix} (T_f M_{T_f} T_x M_{T_x} | T_i M_{T_i}) (\Lambda M_\Lambda J_x M_x | \bar{J} \bar{M}) (J_f M_f \bar{J} \bar{M} | J_i M_i) \{(2L_i+1)(2S_i+1)(2J_f+1)(2\bar{J}+1)\}^{1/2} \Phi_{\nu\Lambda M_\Lambda}(\mathbf{k}_{c.m.}) M(p < NN \rangle \rightarrow p_1 p_2 p_r) \quad (8)$$

Here we used the standard notations for the Clebsch-Gordan and Racah coefficients and 9j-symbols of the rotation group.

The wave function of the relative motion of the c.m. of the cluster x and the center mass of the residual nucleus $B = A-x$ contains the space part $R_{\nu\Lambda}(k)$ that is the harmonic oscillator wave function, and angular wave function $Y_{\Lambda M_\Lambda}(\theta_k, \phi_k)$ that is the spherical harmonic: $\Psi_{\nu\Lambda M_\Lambda}(\mathbf{k}) = R_{\nu\Lambda}(k) Y_{\Lambda M_\Lambda}(\theta_k, \phi_k)$. If one chose the quantization axis OZ along the vector $\mathbf{k}_{c.m.}$, then the only one spin projection M_Λ is allowed in Eq. (8), $M_\Lambda = 0$.

In order to consider the SRC NN pair, one should left in the sum over N_x , L_x in Eq.(8) only the ground state of the cluster x , $N_x = 0$, $L_x = 0$. Therefore from Eq.(7) one has $N_B = N_A - \nu$, where $\nu\Lambda = 20$ and 22 for two nucleons from p -shell, $\nu\Lambda = 00$ for two nucleons from s -shell, and $\nu\Lambda = 11$ for sp -pair of nucleons in the x cluster. Furthermore, for $N_x = 0$ one has $\mathcal{L} = \Lambda$, $\bar{J} = J_f$, $\bar{M} = M_i$ in Eq. (8). Using these simplifications for the case of the ground state of the ^{12}C nucleus with $J_i = 0$, $T_i = 0$, the spin averaged matrix element (1) can be written as

$$\overline{|M_{fi}(p^{12}\text{C} \rightarrow pNN^{10}\text{A})|^2} = \sum_{\sigma_p, \sigma_1, \sigma_2, \sigma_r, M_{J_f}} \left(\frac{A}{2} \right) \left| \sum_{\alpha_f L_f J_f \Lambda} \alpha_i^{A J_i T_i} \alpha_f^{B J_f T_f} U(\Lambda 0 J_f S_x; \Lambda J_x) (\Lambda 0 J_x M_f | J_f M_f) U(S_f J_f L_i \Lambda; L_f J_x) \left(\frac{1}{2} M_{J_f} - \sigma_r \frac{1}{2} \sigma_r | J_x M_{J_f} \right) \{(2S_i+1)(2L_f+1)(2J_f+1)(2J_x+1)\}^{1/2} R_{\nu\Lambda}(p_B) \sqrt{\frac{2\Lambda+1}{4\pi}} \right|^2 ((2J_f+1)(2T_f+1))^{-1} |M_{pN}(s, t)|^2 |\psi_x(q_{rel})|^2. \quad (9)$$

When deriving Eq. (9) we neglect the spin dependence of the pN elastic scattering amplitude in the upper vertex of the diagram in Fig. 1 b.

2.2 Numerical results and discussion

A sensitivity of the cross section of this reaction to the high-momentum part of the $\psi_{<NN>}(q_{rel})$ was studied in Ref. [8] using the Light-Front dynamics for the matrix element of the process $p + \langle NN \rangle \rightarrow p + N + N$. The ratio of squared deuteron wave function $R = |\psi_d(q_{LFD})|^2 / |\psi_d(q_{nr})|^2$ at the relativistic q_{LFD} and non-relativistic q_{nr} relative momentum is shown in Fig. 2 versus the momentum of the spectator nucleon p_r in the rest frame of the ^{12}C at zero momentum of the residual nucleus ^{10}B for different type of the NN -interaction potential [17–19]. The scattering angles of the nucleon p_r are $\theta_r = 6^\circ$ and $\phi_r = 0$. Below we concentrated on the center mass momentum distribution of the SCR pair. In the considered model this distribution is determined by the (squared) wave function $\psi_{\nu\Lambda M_\Lambda}(\mathbf{k}_{c.m.})$, where $\mathbf{k}_{c.m.} = \mathbf{p}_B$. The radial parts of this wave functions $R_{\nu\Lambda}(p/p_0)$ corresponding to the $s^4 p^6$, $s^3 p^7$ and $s^2 p^8$ configurations of the residual nucleus ^{10}B are, respectively, the following

$$\begin{aligned} R_{22}(p) &= C \frac{4}{15} x^2 \exp\{-x^2/2\}, \quad R_{20}(p) = C \sqrt{\frac{3}{2}} (1 - \frac{2}{3} x^2) \exp\{-x^2/2\}, \\ R_{11}(p) &= C \frac{1}{\sqrt{3}} x \exp\{-x^2/2\}, \quad R_{00}(p) = C \exp\{-x^2/2\}, \end{aligned} \quad (10)$$

where $C = \sqrt{4/\sqrt{\pi} p_0^3}$, $x = p/p_0$ and p_0 is the oscillator parameter.

For calculation of the parentage coefficients we use the formalism of Ref. [15]. The coefficients of the mixing configurations $a_{[f]LS}^{A,JT}$ in Eq. (5) for the ground state of the ^{12}C nucleus and for the lowest levels of the ^{10}B nucleus corresponding to $s^4 p^6$ configuration are taken from Ref. [20]. For the states with the destroyed s -shell corresponding to the $s^2 p^8$ and $s^3 p^7$ configurations, the TISM model with mixing configurations, for our knowledge, was not performed. Therefore, in this case for estimate we use the ordinary TISM model. For the oscillator parameter p_0 in Eqs.(10) we use the values found in Ref. [10] from the analysis of the data on the reaction $^{12}\text{C}(p, pd)^{10}\text{B}$, i.e. $p_0 = 155.3 \text{ MeV}/c$ ($144.7 \text{ MeV}/c$) for transition to the $s^4 p^6$ ($s^2 p^8$) configuration of the ^{10}B . In that experiment the resolution over the excitation energy E_B^* of the residual nucleus was high enough to separate the transitions to the $s^4 p^6$ and $s^2 p^8$ states. It was found in [10] that the p_B momentum distribution for transitions to the ground state and lowest excited states of the ^{10}B with $E_B^* \leq 5 \text{ MeV}$ is well described by the $2S$ -type distribution (corresponding to $\nu\Lambda = 20$), whereas for the transitions to the excited states with the $s^2 p^8$ configuration the $0S$ -type distribution ($\nu\Lambda = 00$) is well suitable, as was expected within this model. On the other hand, it was concluded in [10] that the $1D$ ($\nu\Lambda = 22$) transitions to the $s^4 p^6$ states are strongly suppressed as well as the knockout of the sp pairs is suppressed as compared to the p^2 and s^2 pairs. The latter result can be understand qualitatively, since the size of the averaged radius of the trajectory for the s nucleon differs from that for the p nucleon and therefore the probability to find these nucleons at short distances in the quasi-deuteron cluster is lower than for the p^2 and s^2 configurations. A reason for suppression of the $1D$ transitions in the (p, pd) reaction is not quite clear.

When calculating the distribution over the momentum of the center of mass motion of the NN -pair $k_{cm} = p_B$ we use Eq.(9) and put the factors $|M_{pN}(s, t)|^2 |\Psi_x(q_{rel})|^2$ to be constant. For the $s^4 p^6$ configuration of the ^{10}B nucleus we put $\nu = 2$ and perform summation over the lowest states of the ^{10}B nucleus from the region of $E_B^* = 0 - 5 \text{ MeV}$ (see [8] for details) taking into account the mixing configurations [20]. As follows from Eqs. (10), for this case

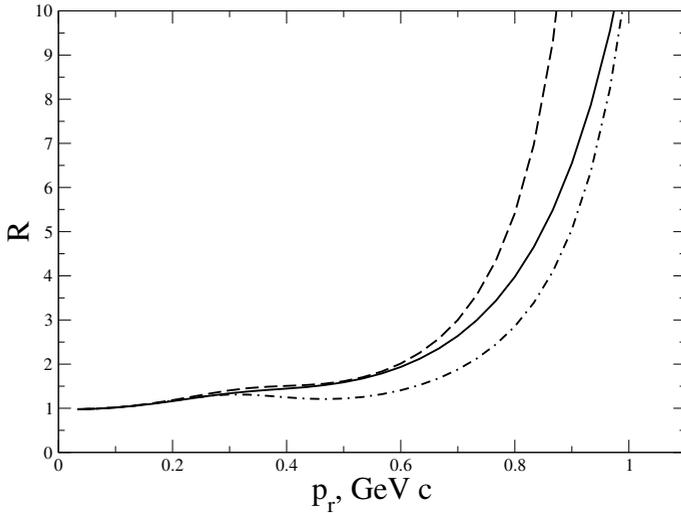


Figure 2. The ratio R (see text for details) versus the spectator momentum p_r for different type of the NN -interaction potential: Bonn [17] (dashed line), Paris [18] (dash-dotted), CD-Bonn [19] (full).

the p_B -distribution is either of the $1D$ (for $\nu\Lambda = 22$) or $2S$ (for $\nu\Lambda = 20$) type. Since both this distributions are not of the $0S$ type, it is clear that their coherent sum in Eq.(9) could not provide the $0S$ -type of the p_B - distribution, and this is indeed obtained in results our calculations shown in Fig. 3 by the full thick line. Inclusion of other states of the ^{10}B nucleus corresponding to the same s^4p^6 configuration but for higher values of E_B^* could not change the shape of the distribution qualitatively.

The transitions to the s^2p^8 states of the ^{10}B nucleus correspond to $\nu\Lambda = 00$ and therefore the p_B distribution in this case is of the $0S$ type (dot-dashed line in Fig.3). If the resolution on E_B^* is rather poor corresponding, for example, to the range $E_B^* = 0 - 30$ MeV, one should make sum over transitions to the states with all configurations s^4p^6 , s^3p^7 , s^2p^8 . Keeping in mind the above mentioned results for suppression of the sp pairs knock-out in the (p, pd) reaction [10], we neglect the transitions to the s^3p^7 configuration and make the sum of the s^4p^6 and s^2p^8 configurations. The obtained result is shown in Fig. 3 by the full thin line which is neither $0S$ nor $1D$ type. For comparison, the pure $1S$ type distribution is also shown in Fig.3 by the dotted line that corresponds to the transition to the ^{10}B state with quantum numbers $J_f = 0$, $T_f = 1$, $E_{exp} = 1.74$ MeV (with $E = 1.51$ MeV from the theoretical model [20]) corresponding to $\nu\Lambda = 20$. One can find from Fig. 3 that both $0S$ and $1S$ distributions, which well describe the p_B -distributions in the $^{12}C(p, pd)^{10}B$ reaction [10], are considerably narrower than the Gaussians $n(p_B) \sim \exp[-p_B^2/2\sigma^2]$ with $\sigma = 140 - 160$ MeV/c found from the (e, epp) reaction in Ref. [6].

If the experimental data on the p_B distribution in the reaction $^{12}C+p \rightarrow ppN + ^{10}B$ will be obtained and found to be similar to that from the (e, epp) reaction with electron beam [6], then one will have to conclude that the formalism which allows one to explain data on the reactions of quasi-elastic knock-out of fast deuterons from the light nuclei [10, 21, 22] and used here for describing the nuclear structure of the ^{12}C , has to be modified in the region of SRC. Presumably, antisymmetrization effects in realistic nuclear wave functions with cluster configurations are of importance here. This issue will be considered in a separate publication. In any case, it is important to get together with the results of the SRC experiment [7] some data at the same beam energy on the reaction of quasi-elastic knockout of the fast deuterons $^{12}C(p, pd)^{10}B$ at high transferred momenta corresponding to the $p\langle NN\rangle \rightarrow pd$ scattering at $\approx 90^\circ$ in the center mass of the pd system.

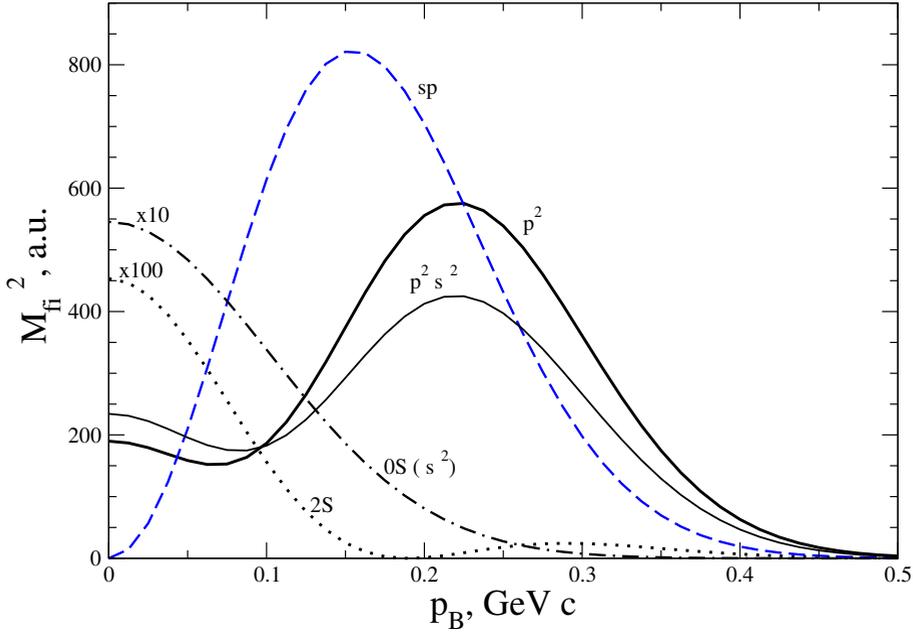


Figure 3. The squared transition matrix element $|M_{fi}(p + {}^{12}\text{C} \rightarrow ppN + {}^{10}\text{B})|^2$ from Eq. (9) divided by $|M_{pN}(s, t)|^2 |\Psi_x(q_{rel})|^2$ for transitions to different states of the ${}^{10}\text{B}$ versus the p_B -momentum: $s^4 p^6$ at $E_B^* \leq 5$ MeV (full thick line), $s^3 p^7$ (dashed), $s^2 p^8$ (dashed-dotted, multiplied by factor 10), $s^4 p^6 + s^2 p^8$ (full thin), the pure $2S$ -type contribution (dotted, multiplied by factor 100). See text for details.

3 Concluding remarks

The exclusive reaction ${}^{12}\text{C}(p, ppN){}^{10}\text{A}$ of quasi-elastic knockout of the nucleon from the short-range correlated nucleon pair $\langle NN \rangle$ in the ${}^{12}\text{C}$ nucleus by proton with an energy of few GeV is considered in the plane-wave impulse approximation by analogy with theory of reactions of quasi-elastic knockout of fast deuterons (p, pd). The nuclear structure is considered within the translationally-invariant shell model with mixing configurations. Short-range correlation effects in the wave function of the ${}^{12}\text{C}$ nucleus ψ_A are taken into account by modification of the internal wave function of the two-nucleon cluster at high relative momenta in the parentage expansion of the ψ_A . Within this assumption the “spectator” $A - 2$ system and its center of mass motion are not modified by the SRC. Relativistic effects in the subprocess $p + \langle NN \rangle \rightarrow p + N + N$ are taken into account within the light-front dynamics and found to be very important at the spectator momenta (in the rest frame of ${}^{12}\text{C}$ nucleus) $p_r \geq 0.5 - 0.6$ GeV/c. In this region of p_r a sensitivity to details of the NN -interaction potential at short NN -distances is very high. The distributions over the center of mass momentum of the $\langle NN \rangle$ -pair are calculated for transitions to different states of the residual nucleus ${}^{10}\text{B}$ and compared with similar results for the reaction (p, pd) [10, 21, 22] and with available data on distribution of the center mass of SRC pairs from experiments with electron beam [6]. The next step of our study will be an account for initial and final state interactions in the considered reaction.

Acknowledgments. I am thankful to O. Hen, M. Patsyuk and E. Piassetzky for stimulating discussion. This work is supported in part by RFBR grant number 18-02-40046.

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