Deuteron distribution in nuclear matter

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(November 4, 2018)

Abstract

We analyze the properties of deuteron-like structures in infinite, correlated nuclear matter, described by a realistic hamiltonian containing the Urbana $v_{14}$ two-nucleon and the Urbana TNI many-body potentials. The distribution of neutron-proton pairs, carrying the deuteron quantum numbers, is obtained as a function of the total momentum by computing the overlap between the nuclear matter in its ground state and the deuteron wave functions in correlated basis functions theory. We study the differences between the S- and D-wave components of the deuteron and those of the deuteron-like pair in the nuclear medium. The total number of deuteron type pairs is computed and compared with the predictions of Levinger’s quasideuteron model. The resulting Levinger’s factor in nuclear matter at equilibrium density is 11.63. We use the local density approximation to estimate the Levinger’s factor for heavy nuclei, obtaining results which are consistent with the available experimental data from photoreactions.

PACS number(s):
I. INTRODUCTION

The suggestion that the nuclear response may be interpreted as the response of a collection of neutron-proton (np) pairs carrying the quantum numbers of the deuteron was first put forward in the fifties by Levinger \[1\] and Gottfried \[2\], to explain nuclear photoabsorption data.

The basic idea underlying the Levinger’s *quasideuteron* (QD) model is that the nuclear photoabsorption cross section \(\sigma_A(E_\gamma)\), above the giant dipole resonance and below the pion threshold, is proportional to that corresponding to the break-up of a deuteron embedded in hadronic matter, and denoted hereafter as \(\sigma_{QD}(E_\gamma)\)

\[
\sigma_A(E_\gamma) = P_D \sigma_{QD}(E_\gamma) . \tag{1}
\]

The proportionality constant \(P_D\) has to be interpreted as the fraction of the \(A(A-1)/2\) nucleon–nucleon pairs, which are of QD type, and it is given by

\[
P_D = L \left[ \frac{Z(A-Z)}{A} \right] , \tag{2}
\]

where \(A\) and \(Z\) denote the nuclear mass and charge and \(L\) is the so called Levinger’s factor. \(P_D\) can be directly calculated from the ground state wave function of the nucleus with mass \(A\). Since the deuteron is a bound state, \(P_D\) scales with the number of particles \(A\).

From \(P_D\), the probability of finding a deuteron-like nucleon pair in a complex nucleus can be easily extracted. Such probability can be obtained by normalizing the number of QD pairs \(P_D\) to the total number of pairs, and, therefore, it is inversely proportional to the number of particles. The probability is zero in infinite nuclear matter, unless the nuclear matter wave function contains a long range order, providing a condensation of QD pairs.

According to the Levinger’s model \[3\] \(\sigma_{QD}(E_\gamma)\) is taken as the deuteron cross section times a damping function, of exponential form, accounting for the Pauli blocking of the final states available to the nucleon ejected from the QD:

\[
\sigma_{QD}(E_\gamma) = \sigma_d(E_\gamma) e^{-(D/E_\gamma)} . \tag{3}
\]
Subsequently, Laget [4] proposed to associate $\sigma_{QD}(E_\gamma)$ with the transition amplitudes of virtual $(\pi + \rho)$-meson exchanges between the two nucleon of the QD pair, leading to a cross section denoted $\sigma_{d}^{exch}(E_\gamma)$.

Both models fit reasonably well the existing photoreaction data in heavy nuclei, but the resulting factors, $L_{Lev}(A)$ and $L_{Laget}(A)$, have different phenomenological values, with $L_{Laget}(A)$ being about 20% larger than $L_{Lev}(A)$.

A generalization of the QD model was proposed by Frankfurt and Strikman [5], to explain the production of fast backward protons in semi-inclusive processes off nuclear targets. According to the model of Refs. [5], generally referred to as few nucleon correlation model, the structure of the nuclear wave function at short internucleon distances is dominated by strongly correlated multinucleon clusters. A quantitative understanding of the above reaction processes requires a microscopic calculation of the quasideuteron distribution $P_D(k_D)$ in the nucleus, as a function of its momentum $k_D$. Moreover, the integral of $P_D(k_D)$ over $k_D$, being proportional to $P_D$, provides an unbiased calculation of the Levinger’s factor $L$.

More recently, the occurrence and spatial structure of deuteron-like configurations in light nuclei has been studied using the Green’s Function Monte Carlo (GFMC) method [6]. It is interesting to extend such analysis to heavier nuclei and to nuclear matter.

Systematic quantitative investigations of nucleon-nucleon (NN) correlations in nuclear matter have been carried out within microscopic many-body theories (for a recent review see Ref. [7]). In particular, Correlated Basis Function (CBF) theory has been applied to obtain the nuclear matter momentum distribution [8] and spectral functions [9–12] from realistic hamiltonians. In this paper we use the same many body framework to carry out an ab initio calculation of the momentum distribution $P_D(k_D)$ of QD pairs in infinite nuclear matter, as well as of the associated total number of QD pairs per particle $P_D/A$.

The definition of the QD total momentum distribution in terms of the overlap between the nuclear matter and the deuteron ground state wave functions is given in Sec. II, where the many-body formalism employed in the calculations is also briefly outlined. In Sec. III the results of numerical calculations, including both the QD momentum distribution and
\( \mathcal{P}_D \) in nuclear matter at the empirical saturation density, \( \rho = 0.16 \text{ fm}^{-3} \), are discussed and compared to the empirical estimates of the Levinger’s factor. Finally, the summary and conclusions are given in Sec. IV.

**II. FORMALISM**

The distribution of QD pairs with total momentum \( \mathbf{k}_D \) in nuclear matter is defined as

\[
P_D(\mathbf{k}_D) = \frac{1}{2J_D + 1} \sum_{i<j} \sum_n |M_{ij}^{n\alpha}(\mathbf{k}_D)|^2 ,
\]

where \( J_D = 1 \) is the spin of the deuteron, and

\[
M_{ij}^{n\alpha}(\mathbf{k}_D) = \int d\tilde{R} d^3 r_1 d^3 r_2 \Psi_{NM}^*(\mathbf{r}_1, \mathbf{r}_2, \tilde{R}) \Psi_D^\alpha(\mathbf{r}_1, \mathbf{r}_2) \Phi_n(\tilde{R}) ,
\]

with \( \tilde{R} \equiv (\mathbf{r}_3, \ldots, \mathbf{r}_A) \). In the above equation, \( \Psi_{NM} \) and \( \Phi_n \) denote the normalized nuclear matter ground state wave function and the wave function of the \( (A-2) \)-nucleon system in the state \( n \), respectively. The configuration space deuteron wave function (DWF) can be cast in the form

\[
\Psi_D^\alpha(\mathbf{r}_{ij}, \mathbf{R}_{ij}) = \frac{e^{i\mathbf{k}_D \cdot \mathbf{R}_{ij}}}{\sqrt{\Omega}} \psi_D^\alpha(ij) |00\rangle ,
\]

where \( \Omega \) is the normalization volume, \( \mathbf{R}_{ij} = (\mathbf{r}_i + \mathbf{r}_j)/2 \), \( \mathbf{r}_{ij} = \mathbf{r}_i - \mathbf{r}_j \), \( |00\rangle \) is the spin-isospin singlet two-nucleon state and the relative motion of the pair is described by

\[
\psi_D^\alpha(ij) = \left[ u_D(r_{ij}) \sigma_i^\alpha - \frac{w_D(r_{ij})}{\sqrt{2}} T^{\alpha\beta}(\hat{r}_{ij}) \sigma_i^\beta \right] .
\]

In Eq.(7), \( u_D(r) \) and \( w_D(r) \) are the \( \ell = 0 \) and \( \ell = 2 \) components of the deuteron wave function, normalized according to

\[
\int_0^\infty r^2 dr \left[ u^2(r) + w^2(r) \right] = 1 ,
\]

\( \sigma_i^\alpha (\alpha = 1, 2, 3) \) denote the Pauli matrices and the tensor operator is given by
\[ T^{\alpha\beta}(\tilde{r}_{ij}) = 3\tilde{r}_{ij}^\alpha\tilde{r}_{ij}^\beta - \delta^{\alpha\beta}. \]  \hfill (9)

In CBF theory \(|\Psi_{NM}\rangle\) is usually written, in coordinate space, in the form \((R \equiv (r_1, \ldots, r_A)\) specifies the nucleon positions) \[ \Psi_{NM}(R) = S\left[ \prod_{i<j} F(ij) \right] \Phi_0(R), \] where \(S\) is the symmetrization operator and \(\Phi_0\) is the Slater determinant describing a non-interacting Fermi gas of nucleons carrying momenta \(k\) with \(|k| \leq k_F = (6\pi^2 \rho/\nu)^{1/3}\), \(\nu\) being the degeneracy of the momentum states (in symmetric nuclear matter \(\nu = 4\)). The operator \(F(ij)\), accounting for the correlation structure induced by the nucleon nucleon (NN) interaction, has been chosen of the form \[ F(ij) = f_c(r_{ij}) + f_\sigma(r_{ij})(\sigma_i \cdot \sigma_j) + f_\tau(r_{ij})(\tau_i \cdot \tau_j) + f_{\sigma\tau}(r_{ij})(\sigma_i \cdot \sigma_j)(\tau_i \cdot \tau_j) \]
\[ + f_t(r_{ij})T^{\alpha\beta}(\tilde{r}_{ij})\sigma_i^\alpha\sigma_j^\beta + f_{\tau\tau}(r_{ij})T^{\alpha\beta}(\tilde{r}_{ij})\sigma_i^\alpha\sigma_j^\beta(\tau_i \cdot \tau_j), \] \hfill (11)

where \(f_c(r), f_\sigma(r), f_\tau(r), f_{\sigma\tau}(r), f_t(r)\) and \(f_{\tau\tau}(r)\) are correlation functions whose radial shapes are determined minimizing the expectation value of the hamiltonian in the ground state described by Eq.(10) \[ (13) \]. As \(r \to \infty, f_c(r) \to 1\), while all other correlation functions go to zero.

Summation over \(k_D\) of \(P_D(k_D)\) yields \(P_D/(2J_D + 1)\) in nuclear matter. This number, which corresponds to an extensive quantity and therefore is proportional to \(A\), leads to a direct evaluation of the Levinger’s factor \(L\), to be compared with the value resulting from the phenomenological analyses \[ (14,15) \] of the available experimental data on photoreactions \[ (16,17) \].

The quantity defined by Eqs.(4) and (5) is related to the fully linked part of the two–nucleon density matrix, \(\rho^{(2)}(r_1, r_2, r_1', r_2')\). This part is the only one providing extra information on the N–N correlations with respect to that carried by the one–body density matrix, or, equivalently, by the nucleon momentum distribution \[ (12) \]. Using standard cluster expansion techniques \[ (18) \], \(P_D(k_D)\) can be written as a series of terms involving an increasing number...
of particles. We have calculated the cluster contributions associated with the diagrammatic
structure shown in fig. 1 and its exchange counterpart, where the deuteron wave function
\( \Psi_D(1,2) \) is multiplied by the correlation operator \( F(1,2) \). This corresponds to a dressed
leading order approximation, whose validity has been checked in previous CBF calculations
of the response function and of the spectral function of nuclear matter, and whose expression
reads

\[
P_D(k_D) = \frac{1}{2} \frac{\rho^2}{4\pi} \int d^3r_{11'}d^3r_{12}d^3r_{1'2'} e^{i k_D \cdot (r_{11'}+r_{1'2'})/2} n(r_{11'}) n(r_{22'}) \Sigma(r_{12}, r_{1'2'}) ,
\]

with

\[
\Sigma(r_{12}, r_{1'2'}) = \frac{1}{3} Tr \left[ F^\dagger(1'2') \psi^\dagger_D(1'2') \Pi_{00} \psi_D(12) F(12) (1 - P_\sigma P_\tau) \right] .
\]

In the above equation, \( \Pi_{00} \) is the operator projecting onto the \( S = 0, T = 0 \) two-nucleon state:

\[
\Pi_{00} = \frac{1 - (\sigma_1 \cdot \sigma_2)}{4} \frac{1 - (\tau_1 \cdot \tau_2)}{4} ,
\]

while the spin- and isospin-exchange operators, \( P_\sigma \) and \( P_\tau \), are given by

\[
P_\sigma = \frac{1 + (\sigma_1 \cdot \sigma_2)}{2} , \quad P_\tau = \frac{1 + (\tau_1 \cdot \tau_2)}{2} .
\]

The function \( n(r) \) is the correlated one-body density matrix \( [8] \), normalized as \( n(r = 0) = 1 \), and trivially related to the nucleon momentum distribution, \( n(k) \), through

\[
n(r) = \frac{\nu}{(2\pi)^3 \rho} \int d^3k e^{ikr} n(k) .
\]

Evaluation of the trace appearing in Eq.(13) leads to the simple result

\[
\Sigma(r, r') = \frac{1}{16} \left[ U(r)U(r') + W(r)W(r')Q(\tilde{r}, \tilde{r}') \right] ,
\]

where

\[
Q(\tilde{r}, \tilde{r}') = \frac{1}{2} \left[ 3 (\tilde{r} \cdot \tilde{r}')^2 - 1 \right] ,
\]
and the functions $U(r)$ and $W(r)$ are defined as
\begin{equation}
U(r) = u_D(r) - \Delta u(r) ,
\end{equation}
\begin{equation}
W(r) = w_D(r) - \Delta w(r) .
\end{equation}

The explicit expression of the functions $\Delta u(r)$ and $\Delta w(r)$, yielding the deviation of $U(r)$ and $W(r)$ from the bare components of the DWF, are
\begin{equation}
\Delta u(r) = u_D(r) [h_c(r) - f_\sigma(r) + 3f_\tau(r) + 3f_{\sigma\tau}(r)]
- \sqrt{8}w_D(r) [f_t(r) - 3f_{t\tau}(r)] ,
\end{equation}
and
\begin{equation}
\Delta w(r) = w_D(r) [h_c(r) - f_\sigma(r) + 3f_\tau(r) + 3f_{\sigma\tau}(r)]
- \sqrt{8} \left(u_D(r) - \frac{w_D(r)}{\sqrt{2}}\right) [f_t(r) - 3f_{t\tau}(r)] ,
\end{equation}
with $h_c(r) = 1 - f_c(r)$. Note that in absence of correlations, i.e. setting $f_c(r) \equiv 1$ and all other correlation functions identically equal to zero, $U(r)$ and $W(r)$ reduce to $u_D(r)$ and $w_D(r)$, respectively.

Using the functions defined in Eqs.(19) and (20), the wave function describing the motion of the QD pair in nuclear matter can be written in the same form as the DWF (see Eqs.(4), (5) and (7)):
\begin{equation}
\Psi_{QD}^\alpha(r_{ij}, R_{ij}) = \frac{e^{i k_0 \cdot R_{ij}}}{\sqrt{\Omega}} \left[U(r_{ij}) \sigma_i^\alpha - \frac{W(r_{ij})}{\sqrt{2}} T^{\alpha\beta}(\hat{r}) \sigma_i^\beta\right] |00\rangle .
\end{equation}

Using the Fourier tranforms of $U(r)$ and $W(r)$, defined as
\begin{equation}
U(k) = \sqrt{\frac{2}{\pi}} \int_0^\infty r^2 dr j_0(kr) U(r) ,
\end{equation}
\begin{equation}
W(k) = \sqrt{\frac{2}{\pi}} \int_0^\infty r^2 dr j_2(kr) W(r) ,
\end{equation}
\( j_0(kr) \) and \( j_2(kr) \) being spherical Bessel functions, and the nucleon momentum distribution in nuclear matter, \( n(k) \), Eq.(12) can be rewritten in the form

\[
P_D(k_D) = \int d^3k \ P(k_D, k),
\]

(26)

where

\[
P(k_D, k) = n \left( \left| \frac{k_D}{2} - k \right| \right) n \left( \left| \frac{k_D}{2} + k \right| \right) |\Psi_{QD}(k)|^2,
\]

(27)

and

\[
|\Psi_{QD}(k)|^2 = \frac{1}{4\pi} \left[ U^2(k) + W^2(k) \right].
\]

(28)

The above equations have been used to carry out the numerical calculations.

It has to be noticed that the contributions arising from the non commuting structure of the correlations reaching the four external vertices, 1, 2, 1’ and 2’, of the diagrammatical structure of fig. [4], are not exactly accounted for, but only according to the dressed leading order approximation.

**III. RESULTS**

Fig. 2 shows the behavior of \( U(r) \) and \( W(r) \) evaluated using a many body hamiltonian including the Urbana \( v_{14} \) NN potential and supplemented by the TNI model of many-body forces [19]. For comparison, we also show the components of the Urbana \( v_{14} \) DWF and the functions \( \Delta u \) and \( \Delta w \) defined in Eqs.(21) and (22), respectively. It appears that the main differences between deuteron and QD occur at \( r < 2 \) fm. At small relative distance \( (r < 1 \) fm), the effect of the nuclear medium leads to an appreciable suppression of \( U(r) \) with respect to \( u_D(r) \), whereas \( W_D(r) \) turns out to be substantially enhanced, compared to \( w_D(r) \).

The momentum space behavior of \( |U(k)|, |W(k)|, |u_D(k)|, |w_D(k)|, |\Delta u(k)| \) and \( |\Delta w(k)| \) is displayed in fig. 3. The main effect of the nuclear medium appears to be a shift of the second minimum of both \( |U(k)| \) and \( |W(k)| \) towards lower values of \( k \).
Eqs. (21) and (22) show that the nuclear medium modifications to the DWF are driven by the functions
\[ H_t(r) = f_t(r) - 3 f_{t\tau}(r) \]
and
\[ \Delta H_c(r) = -f_\tau(r) + 3 f_\tau(r) + 3 f_{\sigma\tau}(r), \]
resulting from the combination of different components of the NN correlation operator. The radial dependence of \( H_t(r) \) and \( \Delta H_c(r) \), illustrated in fig. 4, shows that the effect of scalar and spin-isospin correlations, described by \( \Delta H_c(r) \), dominates at very short relative distance, whereas \( H_t(r) \), accounting for tensor correlations, has a significantly longer range.

The distribution of deuteron pairs with total momentum \( k_D, P_D(k_D) \), resulting from our approach is displayed in fig. 5 as a solid line. Within the Fermi gas model, \( P_D(k_D) \equiv 0 \) at \( |k_D| > 2k_F \), implying that the high momentum tail of \( P_D(k_D) \) is entirely due to NN correlations. The distribution of deuterons in a Fermi gas is represented by a dashed line in the figure. The comparison between the two curves clearly shows that the correlations deplete the distribution with respect to the Fermi gas at \( |k_D| < 2k_F \). The depletion is mostly due to the non-central, tensor correlations.

Similarly, one can define the relative momentum distribution of the nucleons belonging to a QD pair in nuclear matter
\[ P_{rel}^D(k) = \phi(k)|\Psi_{QD}(k)|^2, \quad (29) \]
where
\[ \phi(k) = \int d^3k_D n\left(\left|\frac{k_D}{2} - k\right|\right) n\left(\left|\frac{k_D}{2} + k\right|\right). \quad (30) \]

For example, in the Fermi gas model \( n(k) = \theta(k_F - k) \), and \( \phi(k) \) takes the simple form
\[ \phi(k) = (2\pi)^3 2\rho \left(1 - \frac{3}{2} x + \frac{1}{2} x^3\right) \theta(1 - x), \quad (31) \]
with \( x = k/k_F \).

Fig. 6 shows the relative momentum distribution of a QD pair in nuclear matter, as well as the functions \( |\Psi_{QD}(k)|^2 \) and \( \phi(k) \) defined by Eqs. (28) and (30), respectively. For comparison the relative momentum distribution of a deuteron in free space is also displayed.
The total number of pairs of the QD type in nuclear matter, $P_D$, can be obtained by momentum integration of either $P_D^{rel}(k)$ or $P_D(k_D)$ times the spin multiplicity, $2J_D + 1 = 3$, of the deuteron:

$$
\frac{P_D}{A} = \frac{3}{\rho} \int \frac{d^3k_D}{(2\pi)^3} P_D(k_D) = \frac{3}{\rho} \int \frac{d^3k}{(2\pi)^3} P_D^{rel}(k). \quad (32)
$$

The calculation carried out using the correlated model of nuclear matter and Eq.(26) yields $P_D/A = 2.895$, to be compared with the Fermi gas model result of 3.406.

In order to compare the calculated $P_D$ to the number of QD pairs extracted from the analysis of photonuclear data we have to make a connection with the Levinger’s formula given in Eqs.(1) and (2). The relation is given by

$$
P_D = L \left( \frac{Z(A-Z)}{A} \right), \quad (33)
$$

and, for symmetrical matter ($Z = A/2$), one has

$$
L(A) = 4 \frac{P_D}{A}. \quad (34)
$$

The nuclear matter value resulting from our calculation gives $L(\infty) = 11.63$. This value should be compared with that given by the phenomenological formula

$$
L_{Lev}(A) = 13.82 \frac{A}{R^3[fm^3]}, \quad (35)
$$

reported in Ref. [14], providing $L_{Lev}(\infty) = 9.26$. Notice that, for a deuteron in a Fermi gas, $L_{FG}(\infty) = 13.6$. Surface contributions to $L(A)$ can be estimated by exploiting the calculation of the enhancement factor $\mathcal{K}$ in the electric dipole sum rule for finite nuclei of Ref. [20], performed within the CBF theory and Local Density Approximation (LDA). The enhancement factor is related to experimental data on photoreactions through the equation:

$$
1 + \mathcal{K}_{exp} = \frac{1}{\sigma_0} \int_{m_{\pi}c^2}^{m_{pi}c^2} \sigma_A(E_\gamma) dE_\gamma, \quad (36)
$$

where $\sigma_0 = 60 \left[ Z(A-Z)/A \right] \text{MeV mb}$ and $m_{\pi}c^2$ is the $\pi$-meson production threshold.

Therefore, the Levinger’s factor can be related to $\mathcal{K}$ in the mass number range where the coefficient $D$ in Eq.(3) is fairly $A$-independent, namely for sufficiently large values of $A$. 

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By adding the surface contributions, as extracted from Ref. [20], to the nuclear matter bulk result, we get:

\[ L(A) = 11.63 - 9.76 \, A^{-1/3}. \]  

(37)

Fig. 7 shows our results for \( L(A) \) compared with \( L_{Lev}(A) \) and \( L_{Laget}(A) \), as extracted from the available experimental data on photoreactions. The computed Levinger’s factors are almost A–independent for heavy nuclei (\( A > 100 \)), and result to be \( \sim 25\% \) larger than \( L_{Lev}(A) \) and \( \sim 15\% \) smaller than \( L_{Laget}(A) \), and therefore they are consistent with the experimental data. This differs from what happens for the enhancement factor \( K \), where essentially the same theory as the one used in this paper leads to a value which is \( \sim 60\% \) larger than the experimental one. Therefore, this disagreement between theory and experiment has to be mainly traced back to the sizeable tail contributions to the electric dipole sum rule, absent in the definition of Eq.(36).

IV. CONCLUSIONS

Correlated Basis Function theory of the two–body density matrix has been applied to compute the distribution \( P_D(k_D) \) of neutron–proton pairs characterized by the deuteron wave function and having total momentum \( k_D \).

It has been found that this distribution in nuclear matter is mostly concentrated at \( 0 \leq |k_D| \leq 2k_F \). Besides being responsible for the appearance of the tail of \( P_D(k_D) \) at \( |k_D| > 2k_F \), NN correlations produce an appreciable effect at low momenta. The inclusion of correlations associated with the tensor component of the one pion exchange interaction leads to a \( \sim 15\% \) decrease of \( P_D(k_D) \) at \( |k_D| < k_F \). In general, inclusion of correlations reduces the prediction of the Fermi gas model in this region.

Summation of \( P_D(k_D) \) over \( k_D \) provides the total number \( P_D \) of QD pairs, and, consequently, allows for an \textit{ab initio} calculation of the Levinger’s factor, \( L(A) \). The CBF results for symmetrical nuclear matter, \( L(\infty) = 11.63 \), is about 20\% larger than \( L_{Lev}(\infty) = 9.26 \).
given in the literature. In the case of heavy nuclei $L_{Lev}(A)$ and $L_{Laget}(A)$ bracket our CBF results, which are therefore consistent with the available photoreaction data within the quasideuteron model phenomenology.

It should be noticed that the theoretical estimate of $L(A)$ in the range $150 \leq A \leq 250$ is fairly constant, its increase with $A$ being of $\sim 3\%$. The $A^{-1/3}$ surface behavior leads to a very slow increase of $L(A)$ with $A$, and at $A \sim 200$ we are still quite far away from the asymptotic region.

In addition, the analysis described in this paper shows that when a deuteron is embedded in nuclear matter at equilibrium density, its wave function gets appreciably modified by the surrounding medium. While in the case of the S-wave component the difference is mostly visible at small relative distance ($r < 1$ fm), the D-wave component of the QD appears to be significantly quenched, with respect to the deuteron $w_D(r)$, over the range $0 < r < 2$ fm. It has to be pointed out, however, that the radius of the QD configuration is very close to the deuteron radius, the difference being $\sim 2\%$. This result is in agreement with the conclusions of a recent study of deuteron-like configurations in light nuclei [6]. The authors of Ref. [6] find that the density distributions of $np$ pairs carrying the deuteron quantum numbers in $^3$He, $^4$He, $^6$Li, $^7$Li and $^{16}$O exhibit size and structure similar to those observed in the deuteron.

The relative momentum distribution of a QD pair, $P_{rel}^{D}(k)$, extends into the region $|k| > k_F$, where it appears to be strongly suppressed with respect to the corresponding deuteron momentum distribution $|\Psi_D(k)|^2$, although $|\Psi_{QD}(k)|^2$ is larger than $|\Psi_D(k)|^2$ at high $k$. It has to be pointed out that the behavior of $P_{rel}^{D}(k)$ at $k > k_F$ is entirely dictated by the high momentum tail of the nuclear matter momentum distribution, produced by strong short range NN correlations. Within the Fermi gas model $n(k > k_F) \equiv 0$, and $P_{rel}^{D}(k > k_F)$ vanishes identically.

Higher order cluster terms, neglected in this paper and arising from the inclusion of additional bonds in the diagrammatical structure of fig. [4] are not expected to change the main conclusions of the present paper, neither regarding the behavior of the deuteron distribution
in nuclear matter, nor as far as the discussion on the Levinger’s factor is concerned.

In view of the relevance that $P_D(k_D)$ and $|\Psi_{QD}(k)|^2$ may assume in the study of those lepton–nucleus reactions where the ejected hadron is in kinematical regions forbidden to lepton–nucleon processes, the calculations presented in this paper need to be extended $i)$ by introducing higher order cluster terms in the expansion of the two–body density matrix, and $ii)$ by explicitly considering finite nuclei wave functions. Work in these directions is in progress.

ACKNOWLEDGMENTS

Two of us (O.B. and A.F.) thank the Institute for Nuclear Theory at the University of Washington for its hospitality and the Department of Energy for partial support during the completion of this work. The support from the International Centre for Theoretical Physics and from RFFI grants N98-02-17463 and N99-02-17727 is gratefully acknowledged by A.Yu.I. and G.I.L.
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FIGURES

FIG. 1. Diagram showing the cluster contribution to $P_D(k_D)$ of eqs.(4) and (5) considered in this paper. The oriented solid lines represent the correlated one-body density matrix, whereas the wiggly lines correspond to the dressed deuteron-like np pairs.

FIG. 2. Upper panel: the solid line shows the radial dependence of $U(r)$, defined by Eq.(19), while the dashed and dot-dash lines correspond to $u_D(r)$ and $\Delta u(r)$, respectively. Lower panel: same as the upper panel, but for the $\ell=2$ components of the QD and deuteron wave functions. All wave functions are given in units of (GeV/c)$^{3/2}$.

FIG. 3. Same as in fig. 2 in momentum space. All wave functions are given in units of (GeV/c)$^{-3/2}$.

FIG. 4. Radial dependence of the functions $\Delta H_c(r)$ and $H_t(r)$, entering the definitions of $\Delta u(r)$ and $\Delta w(r)$ (see Eqs.(21) and (22)).

FIG. 5. Momentum distribution of QD pairs in nuclear matter at equilibrium density as a function of the total momentum $|k_D|$ (see. Eqs.(4) and (5)). Solid line: correlated model; dashed line: deuterons in a Fermi gas model. The insert shows a blow up of the region $|k_D|/2k_F < 1$, plotted in linear scale.

FIG. 6. The solid line shows the relative momentum distribution of a QD pair in nuclear matter at equilibrium density (Eq.(29)). The dashed and dot-dash lines correspond to $\phi(k)$ (in units of (GeV/c)$^3$) and $|\Psi_{QD}(k)|^2$ (in units of (GeV/c)$^{-3}$), defined by Eqs.(30) and (28). The diamonds show the squared momentum space wave function of a free deuteron.

FIG. 7. The CBF Levinger’s factor $L(A)$ of heavy nuclei (solid line) and nuclear matter (indicated by the arrow). The LDA approximation of ref. [20] has been used for heavy nuclei. The phenomenological values of $L_{Lev}(A)$ corresponding to photoreaction data of Lepretre et al. [16] (squares) and Ahrens et al. [17] (crosses and diamonds) are taken from ref. [14]. The empirical values of $L_{Lev}(A)$ represented by circles are from ref. [21].
\[ \Delta H_c = f_c(r) + 3f_\tau(r) + 3f_{\sigma\tau}(r) \]

\[ H_t = f_t(r) - 3f_{t\tau}(r) \]
