Nucleon momentum distributions in asymmetric nuclear matter

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Abstract

Nucleon momentum distributions at various densities and isospin-asymmetries for nuclear matter are investigated systematically within the extended Bruecker-Hartree-Fock approach. The shapes of the normalized momentum distributions varying with $k/k_F$ are practically identical, while the density and isospin dependent magnitude of the distribution is directly related to the depletion of the Fermi sea. Based on these properties, a parameterized formula is proposed with the parameters calibrated to the calculated result.

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I. INTRODUCTION

To determine reliably the structure and properties of nuclear matter is one of the central issues in nuclear physics and nuclear astrophysics [1–4]. One of the most important properties of nuclear matter is the neutron and proton momentum distributions which can shed light on the correlations between nucleons [5–8]. In an ideal infinite noninteracting Fermi systems at zero temperature, the momentum distribution is the step function, i.e., \( n(k) = \theta(k_F - k) \), and the Fermi sea is fully occupied. Once the interactions are turned on, the correlations induced by the interactions among fermions lead to the occupation of states with momenta \( k > k_F \) (the high-momentum distribution) and the depletion of the Fermi sea [9–11]. In addition, the depletion can be straightly obtained from the momentum distribution. As for the nuclear matter, due to the hard core and the tensor component of the \( NN \) interaction, the depletion of the Fermi sea is quite significant [12, 13]. It measures the dynamical \( NN \) correlation strength induced by the \( NN \) interaction [7], and is believed to be an indicator for testing the validity of physical picture of independent particle motion in the mean field approach or standard shell model [14, 15] in a nuclear many-body system. The knowledge of the momentum distribution in nuclear matter may provide useful information to acquaint the depletion of the deeply bound state inside finite nuclei and then to understand the structure beyond mean field theory of finite nuclei. It may as well help ones to study the effects of short-range correlations (SRCs) on the observables in the heavy-ion reactions [16, 17].

Experimentally, the high-momentum distribution and the \( NN \) correlations were unambiguously identified in series of experiments, such as \((e, e' p)\) [18], \((e, e' NN)\) [19, 20] and so on. Especially, the two-nucleon knockout experiment shows that nucleons can form short-range correlated pairs with large relative momenta and small center-of-mass momenta [21]. The number of neutron-proton (np) correlated pairs was found to be about 18 times that of proton-proton (pp) correlated pairs [22, 24] which suggests that the tensor correlations due to the strong tensor components of the \( NN \) interaction, in addition to SRCs, play also an important role in the high-momentum distributions [25]. In Ref. [8], the authors attempted to distinguish the dominant regions of tensor correlations and SRCs via comparing the momentum distributions of nuclear matter with the deuteron. They found that SRCs tend to dominate the high-momentum distributions above \( 3 \text{ fm}^{-1} \), while the tensor correlations is
of interest in the region of $k \sim 2 - 2.8 \, \text{fm}^{-1}$. However, both in the theoretical calculations and experiments, the effects of these two correlations on the momentum distribution are hard to distinguished strictly.

In theoretical calculations, the $NN$ correlations in nuclear matter have often been studied in combination with the nucleon momentum distribution. Various theoretical methods have been employed to study these distributions, such as the correlated basis functions [26, 27], quantum Monte Carlo method [28], the self-consistent Green’s function (SCGF) [8, 29, 32], the in-medium T-matrix method [33, 34] and the extended Brueckner-Hartree-Fock (EBHF) method [3, 35, 36]. In Ref. [32], the temperature, density, and isospin dependence of the depletion of the Fermi sea is clarified in the framework of SCGF. The momentum distribution at large momentum has been discussed as well which shows an exponential damping tendency. In Ref. [38], the authors have calculated the nucleon momentum distribution and quasiparticle strength in symmetric nuclear matter within EBHF approach. Parameterized three-section expression of the momentum distribution fit to the microscopic calculation has also been provided. Unfortunately, the parametrization is density independent and merely valid for symmetric case. In the present paper, we shall extend the parameterized expression of the momentum distribution to asymmetric nuclear matter and simplify the form of the expression of the momentum distribution. Moreover, the density and isospin dependence of the depletion of the Fermi sea is discussed as well within the EBHF approach. In order to obtain a more realistic momentum distribution expression, the calculated momentum distribution within the EBHF approach includes the three-body force (TBF) effects.

This paper is organized as follows. In the next section, we give a brief review of the adopted theoretical approaches including the EBHF theory and spectral function. The formula of the momentum distribution is derived in Sec. III. In Sec. IV, we employ the obtained formula to study the SRC effects on the heavy-ion reactions. And finally, a summary is given in Sec. V.

II. THEORETICAL APPROACHES

The present calculations for asymmetric nuclear matter are based on the EBHF approach, for which one can refer to Ref. [39] for details. The extension of the BHF scheme to include microscopic TBF can be found in Refs. [40, 41]. After several self-consistent iteration,
the effective interaction matrix $G$ in the Brueckner-Bethe-Goldstone (BBG) theory can be obtained. This G-matrix which include all the ladder diagrams of the $NN$ interaction embodies the tensor correlations and the SRCs. Using the G-matrix, the mass operator $M(k, \omega)$ can be calculated.

A. The mass operator within the extended Brueckner-Hatree-Fock approach

Generally, the nucleon momentum distribution needs the exact knowledge of the mass operator. In practice, it is impossible to calculate the mass operator exactly. In an actual calculation, one can only evaluate some approximations to the mass operator. Within the framework of the BBG theory, the mass operator can be expanded in a perturbation series according to the number of hole lines. To the lowest-order approximation, i.e., the BHF approximation, the mass operator is written as

$$M_1(k, \omega) = \sum_{k'} \theta(k_F - k') \langle kk' | G[\omega + \epsilon(k')] | kk' \rangle_A,$$

where $\omega$ is the starting energy and $\epsilon(k)$ represents the s.p. spectrum in the BHF approximation. The step function $\theta(k_F - k)$ is the Fermi distribution at zero temperature. The subscript $A$ denotes antisymmetrization of the matrix elements.

The quantity $M_1(k, \omega)$ only has a right-hand cut which is mainly responsible for the depletion under Fermi surface \cite{11, 12}. Therefore, the calculation of the momentum distribution requires at least the first two order approximation of the mass operator. The second order in the hole-line expansion of the mass operator, which might be answerable for the high-momentum distributions above Fermi surface \cite{11}, is given by \cite{39}

$$M_2(k, \omega) = \frac{1}{2} \sum_{k'k_1k_2} \theta(k' - k_F) \theta(k_F - k_1) \theta(k_F - k_2)$$

$$\times \frac{|\langle kk' | G[\epsilon(k_1) + \epsilon(k_2)] | k_1k_2 \rangle_A|^2}{\omega + \epsilon(k') - \epsilon(k_1) - \epsilon(k_2) - i\epsilon},$$

where the step function $\theta(k' - k_F)$ guarantees the integral over $k'$ above the Fermi surface. In the present paper, we calculate the mass operator to the second order approximation, i.e., $M(k, \omega) \cong M_1(k, \omega) + M_2(k, \omega)$.  


B. The spectral function and the momentum distribution

The knowledge of $M(k, \omega)$ allows us to write down the Green’s function in the energy-momentum representation,

$$G(k, \omega) = \frac{1}{\omega - \frac{k^2}{2m} - M(k, \omega)}. \quad (3)$$

Except at the Fermi energy $\epsilon_F$, the mass operator $M(k, \omega)$ is complex and can be written as

$$M(k, \omega + i\eta) = V(k, \omega) + iW(k, \omega) \quad (4)$$

with the property $[M(k, \omega + i\eta)]^* = M(k, \omega - i\eta)$, where $\eta = +0$ to ensure the integral-path. The spectral function $S(k, \omega)$, which describes the probability density of removing a particle with momentum $k$ from a target nuclear system and leaving the final system with the excitation energy $\omega$, is thus given by

$$S(k, \omega) = \frac{i}{2\pi} [G(k, \omega) - G(k, \omega)^*] = -\frac{1}{\pi} \frac{W(k, \omega)}{[\omega - k^2/2m - V(k, \omega)]^2 + W(k, \omega)^2}. \quad (5)$$

And it should fulfill the sum rule

$$\int_{-\infty}^{\infty} S(k, \omega) d\omega = 1. \quad (6)$$

In Ref. [6], the authors show that an elaborately dealing with the integral over the energy can satisfy the sum rule quite well by adopting the mass operator up to second order in the framework of EBHF approach.

Finally, the momentum distribution $n(k)$ is related to the spectral function by

$$n(k) = \int_{-\infty}^{\epsilon_F} S(k, \omega) d\omega \quad (7)$$

or equivalently,

$$n(k) = 1 - \int_{\epsilon_F}^{\infty} S(k, \omega) d\omega. \quad (8)$$
FIG. 1: (Color online). Neutron and proton momentum distributions in asymmetric nuclear matter at various isospin-asymmetries calculated within the EBHF approach. Fig. (a)~(f) represent the different densities $0.50\rho_0$, $0.75\rho_0$, $1.0\rho_0$, $1.50\rho_0$, $2.0\rho_0$, $2.5\rho_0$.

The Fermi energy $\epsilon_F$ follows the on shell condition $\epsilon_F = k_F^2/2m + V(k, \epsilon_F)$. Using the momentum distribution $n(k)$, one can then define the depletion parameter

$$\chi = \frac{\sum_k n(k > k_F)}{\rho} = \frac{1}{\pi^2} \int_{k_F}^{\infty} n(k)k^2dk/\rho,$$

i.e., the proportion of the particle number above the Fermi momentum. Which is related to several physical quantities such as the correlation strength or the defect function, and is believed to be an indicator for the convergence of the so-called BBG hole-line expansion.

### III. THE FORMULA OF THE MOMENTUM DISTRIBUTION

In this section, we first exhibit the numerical calculation of momentum distributions within the EBHF approach, then roughly analyze the behavior of these distributions, and finally provide a formula of calculating the distribution. The realistic Argonne $V18$ two-body interaction supplemented with a microscopic $3BF$ [40, 41] is taken as the $NN$ interaction.

In the present paper, the calculation is under zero temperature.

We systematically report the calculated neutron and proton momentum distributions at various isospin-asymmetries $\beta = 0.0, 0.1, 0.2, 0.3, 0.4, 0.5$ with different total densities $0.50\rho_0$, $0.75\rho_0$, $1.0\rho_0$, $1.50\rho_0$, $2.0\rho_0$, $2.5\rho_0$ in Fig. 1. Hereafter, the isospin-asymmetry $\beta$ is defined as $\beta = (\rho_n - \rho_p)/(\rho_n + \rho_p)$ and $\rho_0 = 0.17$ fm$^{-3}$ is the empirical saturation density of...
FIG. 2: (Color online). The normalized momentum distribution $n(k)/n(0)$ vs $k/k_F$ below the Fermi momentum at various isospin-asymmetries (upper panel) and at various densities (lower panel).

nuclear matter. The distributions present a discontinuity at their respective Fermi momenta $k_F^n$ (hereafter $\tau = n, p$). For positive asymmetries, the neutron Fermi momentum $k_F^n$ is larger than the proton Fermi momentum $k_F^p$, therefore the proton and neutron momentum distributions are located at the left and right sides of the symmetric case, respectively. One should note that the neutron momentum distribution differs only slightly from proton momentum distribution in symmetric case due to the charge-dependent interaction Argonne V18. The discrepancy is too tiny to be recognized in the Figs.

Interestingly, if focusing on the shapes of the momentum distributions in Fig.1, one would notice that these distributions are quite similar except the magnitudes. Inspired by this quality, we show the normalized momentum distributions $n(k)/n(0)$ as a function of the ratio $k/k_F^\tau$ below the Fermi momentum at various isospin asymmetries and densities in Fig. 2. Where $k_F^\tau$ is the respective Fermi momentum corresponding to the different isospin-asymmetry $\beta$ and density $\rho$. The shapes of the normalized distributions are practically identity except slightly small discrepancies near the Fermi momentum. In Fig. 3, the same normalized momentum distributions as Fig. 2 but above the Fermi moment exhibit the coincidence of the shapes as well. In other words, the normalized momentum distributions as a function of the ratio $k/k_F^\tau$ below (above) the Fermi momentum at various densities and isospin-asymmetries can be described by the same expression with tolerable errors.

In the domain $0 < k/k_F^\tau < 1$, the normalized momentum distribution varying with $k/k_F^\tau$.
FIG. 3: (Color online). The normalized momentum distribution \( n(k)/n(1.05k_F) \) vs \( k/k_F \) above the Fermi momentum at various isospin-asymmetries (upper panel) and various densities (lower panel).

below the Fermi momentum can be described by the following parametrization:

\[
\frac{n^\tau_<(k)}{n(0)} = 1.00329 - 0.02876x - 0.09053x^7, \tag{10}
\]

where \( n^\tau_<(k) \) corresponds to \( n(k < k_F^\tau) \) and \( x \) represents the ratio \( k/k_F^\tau \). A comparison between the calculated \( n^\tau_<(k)/n(0) \) and the parametrization is shown in the upper panel of Fig. 4. The polynomial fit is in good agreement with the calculation within the EBHF approach.

FIG. 4: (Color online). The normalized distributions as a function of \( k/k_F \) and the fittings. The upper and the lower panels correspond to the momentum below and above the momentum, respectively.
For the high-momentum distributions, i.e., \( k > k^*_{\tau} \), Ref. [42] reports that the momentum distributions appear to decrease as \( k^{-4} \), following the Tan’s relation [43, 44]. However, the Tan’s relation is simply valid for dilute system with contact interaction whereas the \( NN \) interaction is much more complicated. And the microscopic calculations including the EBHF approach [38] and the SCGF method [8, 32] indicate a nearly linear relation between \( \ln n(k) \) and \( k \) at large momentum, i.e., \( n(k) \propto \exp(-ck) \) (\( c \) is a positive constant). In addition, if one adopts the form of \( k^{-4} \) to describe the high-momentum distributions, a cutoff \( k_\Lambda \) is always supplemented owing to the slow convergence of the number density. When \( k_\Lambda \) is employed, the neglect number density is

\[
\int_{k_\Lambda}^{\infty} n(k)k^2dk \propto \frac{1}{k_\Lambda}.
\]

In calculations, the maximum value of \( k_\Lambda \) is usual about \( 5 \text{ fm}^{-1} \), which implies three to five percent missing of the number density. Most importantly, our calculations within the EBHF approach reveal the same behavior of the high-momentum distributions as Ref. [38]. On account of the above reasons, we employ the exponential form replenished by a Gauss function to describe the high-momentum behavior. The normalized momentum distribution above the Fermi momentum can be expressed as

\[
\frac{n^\tau_{\tau}(k)}{n(1.05k^*_{\tau})} = 3.548e^{-1.799x} + 52.2e^{-4.2766x^2},
\]

with \( n^\tau_{\tau}(k) \equiv n(k > k^*_{\tau}) \) and \( x \equiv k/k^*_{\tau} \). We display the expression of \( k^{-4} \) (1/\( k^4 \)), the parametrization (12) (Exp) and the calculation within EBHF approach in the lower panel of Fig. 4. Obviously, the exponential fit is more approaching to the calculation than the \( k^{-4} \) fit. But we should stress that owing to the approximations adopted in the calculations and fittings, the possibility of Tan’s relation in nucleon momentum distribution could not be ruled out.

To obtain the momentum distributions, the magnitudes of \( n(k) \) below and above the Fermi momentum remain to be identified once the shapes are provided. Actually, the mag-
FIG. 5: (Color online). The depletion parameter $\chi$ calculated within EBHF approach varying with isospin-asymmetry $\beta$ for two densities $\rho_0$ and $2.0\rho_0$.

FIG. 6: (Color online). The depletion parameter $\chi$ vs densities. The symbols correspond to the calculations within the EBHF approach. The lines are obtained form the Eq. (16). The upper and lower lines are related to proton and neutron, respectively.
nitudes connect with the depletion parameter $\chi$ via the relations

\begin{align*}
1 - \chi &= \frac{\frac{1}{\pi^2} \int_{0}^{k_F} n_\tau^r(k)k^2dk}{\rho_\tau} \\
&= 3 \int_{0}^{1} n_\tau^r(k)x^2dx = 0.9546n(0), \quad (13) \\
\chi &= \frac{\frac{1}{\pi^2} \int_{k_F}^{\infty} n_\tau^r(k)k^2dk}{\rho_\tau} \\
&= 3 \int_{1}^{\infty} n_\tau^r(k)x^2dx = 2.9537n(1.05k_F). \quad (14)
\end{align*}

Consequently, the magnitudes of the $n(k)$ below and above the Fermi momentum directly related to the depletion parameter, i.e.,

\begin{align*}
n(0) &= \frac{1 - \chi}{0.9546}, \quad n(1.05k_F) = \frac{\chi}{2.9537} . \quad (15)
\end{align*}

The depletion parameter $\chi$ calculated within the EBHF approach with various isospin-asymmetries $\beta = 0.0, 0.1, 0.2, 0.3, 0.4$ and 0.5 at two typical densities $\rho_0$ and $2.0\rho_0$ are exhibited in Fig. 5. Obviously, the proton (neutron) depletion of the Fermi sea increases (decreases) almost linearly with varying isospin-asymmetry. Due to Eq. (15), $n(0)$ emerges the analogous behavior which have been reported in Ref. [5, 45]. The experiments show that the np correlation is much stronger than nn or pp correlation [23, 24]. One should notice that the probability of a proton (neutron) encounters a neutron (proton) increases (decreases) linearly as a function of isospin-asymmetry. If supposing equal correlation in each correlated np pair and neglecting the nn/pp correlation, the linear isospin dependence of $\chi$ comes very naturally. In Fig. 6 we illustrate the density dependence of the depletion parameter. The different types of dots correspond to the calculated $\chi$ within the EBHF approach. According to the shapes of $\chi$ varying with $\rho$, we propose an expression with the parameters calibrated to the calculated $\chi$. The expression reads

\begin{align*}
\chi(\rho, \beta) &= 0.1669[1 + \lambda(0.1407 \frac{\rho}{\rho_0} - 0.7296)\beta] \\
&\times [1 + 2.448e^{-4.1854\frac{\rho}{\rho_0}} + 0.1382(\frac{\rho}{\rho_0})^{1.5}] . \quad (16)
\end{align*}

Where $\lambda = 1/-1$ corresponds to neutron/proton. The isospin and density dependence of $\chi$ are mainly included in the first and second square brackets on the right-hand side of Eq. (16), respectively. One would find that there is a slight discrepancy between the slopes of
curves in Fig. 5 indicating a weak density dependence of $\partial \chi / \partial \beta$. We actually account this dependence in the first square bracket of Eq. (16). In fact, a simple analysis of the calculated data on the density dependence of the slope reveals a roughly linear dependence. In Ref. [8], the authors have also mentioned a similar behavior of momentum distribution in asymmetric nuclear matter at finite temperature. The expression (16) for various densities and isospin-asymmetries are shown by lines in Fig. 6. Below the saturation density, the depletion of the Fermi sea becomes stronger with decreasing density which might mainly result from the increasing effect of the tensor correlation. While above the saturation density, the hard-core effect and the depletion get larger and large with increasing density [7].

Finally, the formula of the momentum distribution can be summarized as

$$n(k) = \begin{cases} 
\frac{1-\chi}{0.9546} [1.00329 - 0.02876 \frac{k}{k_F} - 0.09053 (\frac{k}{k_F})^7] & \text{if } k \leq k_F^\tau, \\
\frac{\chi}{2.9537} [3.548 e^{-1.799 \frac{k}{k_F}} + 52.2 e^{-4.2766 (\frac{k}{k_F})^2}] & \text{if } k \geq k_F^\tau, \end{cases}$$

(17)

with the expression of the depletion parameter Eq. (16). A comparison between the momentum distributions from formula (17) and from the EBHF approach is given in Fig. 7. It can be clearly seen that the formula is quite accurate except a slight difference near the Fermi momentum. This formula can be applied to calculate the momentum distribution in finite nuclei assisted by the local density approximation. As is well known, at low densities the

FIG. 7: (Color online). Neutron and proton momentum distributions form the formula (17) and the calculation within the EBHF approach at two isospin-asymmetries and densities.
nuclear matter system can minimize its energy by forming light cluster such as deuterons, or particularly strongly bound $\alpha$ particle \cite{46}. In theoretical calculations such as EBHF approach, the in-medium T-matrix method and SCGF method, the effective interaction including all the ladder-diagram contribution always encounters a singularity leading to unstable results at low densities \cite{47,48}. Therefore, we emphasize that the achieved formula (17) of the momentum distribution might be solely reliable for the density of $0.1\rho_0 < \rho < 3.0\rho_0$ and the isospin-asymmetry of $\beta \in (-0.5, 0.5)$ for uniform nuclear matter. Otherwise, one should be careful of the depletion parameter.

IV. APPLICATION TO THE TRANSPORT MODEL

As an example of application, the obtained density and asymmetry-dependent nucleon momentum distribution Eq. (17) and the fraction of high-momentum nucleons Eq. (16) were both involved in the isospin-dependent Boltzmann-Uehling-Uhlenbeck (IBUU) transport model \cite{49}. The free neutron to proton ratio and the $\pi^-/\pi^+$ ratio as a function of momentum in the central Au+Au reaction at 400 MeV/nucleon are demonstrated in the upper and middle windows of Fig. 8. As comparisons, nucleon momentum distribution from Cai & Li is also used \cite{17,50}. The lower window shows the integrations of the $\pi^-/\pi^+$ ratio and comparison with the FOPI data \cite{51}. From Fig. 8 with the formula (17) and the work of Cai & Li, it is seen that both the momentum distribution of the free n/p ratio and the $\pi^-/\pi^+$ ratio are quite similar, except for energetic n/p ratio. While the difference of the total $\pi^-/\pi^+$ yields ratios is evidently shown in the lower window of Fig. 8. The value of $\pi^-/\pi^+$ yields ratio with the formula (17) is higher than that with the form of Cai & Li. The reason is that the fraction of high-momentum nucleons with the formula (17) is smaller than that with the form of Cai & Li \cite{50} and larger number of neutron-proton correlation causes a lower value of the $\pi^-/\pi^+$ ratio in heavy-ion collisions.

V. SUMMARY AND OUTLOOK

Summarily, we have systematically calculated the nucleon momentum distributions and the depletions of the Fermi sea at various densities and isospin-asymmetries for nuclear matter within the EBHF approach. The identity of these shapes of the normalized mo-
FIG. 8: (Color online) Upper panel: Free neutron to proton ratio as a function of momentum in the central Au+Au reaction at 400 MeV/nucleon. Middle panel: Same as upper panel, but for $\pi^-/\pi^+$ ratio. Lower panel: Comparison of calculated total $\pi^-/\pi^+$ yields ratios and FOPI data [51].

momentum distributions below (above) the Fermi momentum varying with $k/k_F$ is detected, indicating an uniform expression of the of the momentum distribution for different densities and isospin-asymmetries. Whereas the magnitude of the momentum distribution is directly related to the depletion of the Fermi sea, which first decreases and then increases with densities resulting from the tensor and hard core effects of the $NN$ interaction [7]. Using these properties, the parameterized formula of momentum distribution is proposed with the
expression of the depletion. Moreover, a heavy-ion reaction example adopting the obtained formula is given to test its reliability.

In the present paper, the mass operator is just calculated up to the second order, the missing higher order perhaps enhances the depletion of the Fermi sea and eventually influences the momentum distribution. Especially, the missing higher order might as well reduce the particle strength around Fermi surface [52]. Thus, the parameterized formula cannot be considered as definite. In addition, the calculation is based on realistic Argonne V18 only. With different interactions, the depletions of the Fermi sea and the momentum distributions would differ from each other [7, 8, 32]. Furthermore, the normal state of symmetric nuclear matter becomes unstable owing to pairing tendency of np [47, 48, 52, 53] and one should account the effect of the pairing on the momentum distribution. An improvement of the calculations including these effects is under way.

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[1] B. A. Li, L.W. Chen, and C. M. Ko, Phys. Rep. 464, 113 (2008).
[2] I. Bombaci and U. Lombardo, Phys. Rev. C 44, 1892 (1991).
[3] A. Ramos, A. Polls and W. H. Dickhoff, Nucl. Phys. A 503, 1 (1989).
[4] M. Baldo and H. R. Moshfegh, Phys. Rev. C 86, 024306 (2012).
[5] P. Yin, J. Y. Li, P. Wang, and W. Zuo, Phys. Rev. C 87, 014314 (2013).
[6] P. Wang, and W. Zuo, Phys. Rev. C 89, 054319 (2014).
[7] Z. H. Li and H.-J. Schulze, Phys. Rev. C 94, 024322 (2016).
[8] A. Rios, A. Polls and W. H. Dickhoff, Phys. Rev. C 89, 044303 (2014).
[9] A. B. Migdal, Sov. Phys. -JETP 5, 333 (1957) [Zh. Eksp. Teor. Fiz. (USSR) 32 339 (1957)].
[10] J. M. Luttinger, Phys. Rev. 119, 1153 (1960).
[11] C. Mahaux and R. Sartor, Phys. Rep. 211, 53 (1992).
[12] J. P. Jeukenne, A. Lejeune, and C. Mahaux, Phys. Rep. 25, 83 (1976).
[13] B. E. Vonderfecht, W H. Dickhoff, A. Polls, and A. Ramos, Phys. Rev. C 44, R1265 (1991); Nucl. Phys. A 555, 1 (1993).
[14] V. R. Pandharipande, I. Sick, and P. K. A. deWitt Huberts, Rev. Mod. Phys. 69, 981 (1997).
[15] J. M. Cavedon, B. Frois, D. Goutte et al., Phys. Rev. Lett. 49, 978 (1982).
[16] G. C. Yong and B. A. Li, Phys. Rev. C 96, 064614 (2017).
[17] Z. X. Yang, X. H. Fan, G. C. Yong and W. Zuo, Phys. Rev. C 98, 014623 (2018).
[18] D. Rohe et al., Phys. Rev. Lett. 93, 182501 (2004).
[19] C. J. G. Onderwater et al., Phys. Rev. Lett. 81, 2213 (1998).
[20] R. Starink et al., Phys. Lett. B 474, 33 (2000).
[21] R. Shneor et al., Phys. Rev. Lett. 99, 072501 (2007).
[22] E. Piasezky, M. Sargsian, L. Frankfurt, M. Strikman, and J. W. Watson, Phys. Rev. Lett. 97, 162504 (2006).
[23] R. Subedi et al., Science 320, 1476 (2008).
[24] O. Hen et al. (The CLAS Collaboration), Science 346, 614 (2014).
[25] R. Schiavilla, R. B. Wiringa, S. C. Pieper, and J. Carlson, Phys. Rev. Lett. 98, 132501 (2007).
[26] S. Fantoni and V. Pandharipande, Nucl. Phys. A 427, 473 (1984).
[27] O. Benhar, A. Fabrocini, and S. Fantoni, Nucl. Phys. A 505, 267 (1989); Phys. Rev. C 41, R24 (1990).
[28] A. Gezerlis and J. Carlson, Phys. Rev. C 81, 025803 (2010).
[29] Y. Dewulf, D. Van Neck, and M. Waroquier, Phys. Rev. C 65, 054316 (2002); Y. Dewulf, W. H. Dickhoff, D. Van Neck, E. R. Stoddard, and M. Waroquier, Phys. Rev. Lett. 90, 152501 (2003).
[30] T. Frick, H. Muther, A. Rios, A. Polls, and A. Ramos, Phys. Rev. C 71, 014313 (2005).
[31] A. Rios, A. Polls, and I. Vidana, Phys. Rev. C 79, 025802 (2009).
[32] A. Rios, A. Polls, and W. H. Dickhoff, Phys. Rev. C 79, 064308 (2009).
[33] P. Bozek, Phys. Rev. C 59, 2619 (1999); 65, 054306 (2002).
[34] V. Soma and P. Bozek, Phys. Rev. C 78, 054003 (2008).
[35] R. Sartor and C. Mahaux, Phys. Rev. C 21, 1546 (1980); P. Grangé, J. Cugnon, and A. Lejeune, Nucl. Phys. A 473, 365 (1987); M. Jaminon and C. Mahaux, Phys. Rev. C 41, 697 (1990); M. Baldo, I. Bombaci, G. Giansiracusa, and U. Lombardo, Nucl. Phys. A 530, 135.
(1991); C. Mahaux and R. Sartor, Nucl. Phys. A 553, 515 (1993).

[36] Kh. S. A. Hassaneen and H. Muther, Phys. Rev. C 70, 054308 (2004).

[37] P. Wang, S.-X. Gan, P. Yin, and W. Zuo, Phys. Rev. C 87, 014328 (2013).

[38] M. Baldo, I. Bombaci, G. Giansiracusa, U. Lombardo, C. Mahaux, and R. Sartor, Phys. Rev. C 41, 1748 (1990); Nucl. Phys. A 545, 741 (1992).

[39] W. Zuo, I. Bombaci, and U. Lombardo, Phys. Rev. C 60, 024605 (1999).

[40] P. Grangé, A. Lejeune, M. Martzolff, and J.-F. Mathiot, Phys. Rev. C 40, 1040 (1989).

[41] W. Zuo, A. Lejeune, U. Lombardo, and J.-F. Mathiot, Nucl. Phys. A 706, 418 (2002); Eur. Phys. J. A 14, 469 (2002).

[42] O. Hen, L. B. Weinstein, E. Piasecki, G. A. Miller, M. M. Sargsian, and Y. Sagi, Phys. Rev. C 92, 045205 (2015).

[43] S. Tan, Annals of Physics 323, 2952 (2008); 323, 2971 (2008); 323, 2987 (2008).

[44] J. T. Stewart, J. P. Gaebler, T. E. Drake, and D. S. Jin, Phys. Rev. Lett. 104, 235301 (2010).

[45] H. Muther, G. Knehr, and A. Polls, Phys. Rev. C 52, 2955 (1995).

[46] S. Typel, G. Ropke, T. Klahn, D. Blaschke, and H. H. Wolter, Phys. Rev. C 81, 015803 (2010).

[47] V. J. Emery, Nucl. Phys. A 12, 69 (1959); H. F. Arellano and J. P. Delaroche, Eur. Phys. J. A 51, 7 (2015);

[48] W. H. Dickhoff, Phys. Lett. B 210, 15 (1988); B. E. Vonderfecht, C. C. Gearhart, W. H. Dickhoff, A. Polls and A. Ramos, Phys. Lett. B 253, 1 (1991).

[49] G. C. Yong, Phys. Rev. C 96, 044605 (2017).

[50] B. J. Cai, B. A. Li, and L. W. Chen, Phys. Rev. C 94, 061302 (2016).

[51] W. Reisdorf et al. (FOPI Collaboration), Nucl. Phys. A 848, 366 (2010).

[52] X. H. Fan, X. L. Shang, J. M. Dong and W. Zuo, Phys. Rev. C. 99, 065804 (2019)

[53] X. L. Shang, and W. Zuo, Phys. Rev. C. 88, 025806 (2013); X. L. Shang, P. Wang, P. Yin and W. Zuo, J. Phys. G: Nucl. Part. Phys. 42, 055105 (2015).