Microscopic optical potentials including breakup effects for d-nucleus elastic scattering

Shoya Ogawa,1,* Ryo Horinouchi,1,‡ Masakazu Toyokawa,1,‡ and Takuma Matsumoto1,§

1Department of Physics, Kyushu University, Fukuoka 819-0395, Japan

(Dated: April 9, 2018)

We construct a microscopic optical potential including breakup effects for d scattering within the Glauber model, in which a nucleon-nucleus potential is derived by the g-matrix folding model. The derived microscopic optical potential is referred to as the eikonal potential. For d scattering, the calculated elastic cross section with the eikonal potential well reproduces the result with an exact calculation estimated by the continuum-discretized coupled channels method. We found that the eikonal potential takes into account breakup effects of d within the adiabatic approximation, which is valid for scattering at high incident energies. The eikonal potential model is expected to be the Glauber model excluded the eikonal approximation.

I. INTRODUCTION

Microscopic understanding of nucleon-nucleus (NA) and nucleus-nucleus (AA) optical potentials is one of the most important issues in nuclear reaction theory. The optical potentials are necessary to describe not only elastic scattering but also reactions including higher-order processes. For example, the distorted-wave Born approximation and the continuum discretized coupled-channels method (CDCC) [1–3] require the optical potentials between constructs of a projectile (P) and a target (T) to describe inelastic scattering, breakup and transfer reactions.

The g-matrix folding model has been widely used as a reliable method to describe the optical potential [4–9]. The g matrix is an effective nucleon-nucleon interaction [10–20] in nuclear matter, and depends on the density ρm of the nuclear matter. In the g-matrix folding model, the optical potentials are derived by folding the g-matrix with the target density ρT for NA scattering, and with ρT and the projectile one ρP for AA scattering. The procedures are referred to as the single-folding (SF) model for NA scattering and double-folding (DF) model for AA scattering, respectively.

In the SF model, ρT is referred as ρm in the g matrix with the local density approximation, and the calculated NA optical potential becomes nonlocal by taking into account knock-on exchange processes. The nonlocal potential is not practical in many applications, but can be localized by the Brieva-Rook approximation [12] in good accuracy. The localized potential is quite successful in reproducing experimental data systematically, when reliable g matrices such as Melbourne [17], CEG [15, 19], and χEFT [20] g matrices are adopted.

On the other hand, it is more difficult to derive the optical potential with the DF model for AA scattering than that with the SF model for NA scattering. The main problem of the DF model is how to treat ρm in the g matrix for AA scattering. In general, the frozen-density approximation (FDA) is applied, in which the sum of ρP and ρT is taken as ρm in the g matrix. The optical potential derived by the DF model with the FDA often needs a normalization factor for the real and imaginary parts to reproduce experimental data. Thus the choice of ρm in the g matrix for the DF model is a longstanding open problem.

In the previous works for 3,4He scattering on various targets [21, 22], we proposed the target-density approximation (TDA) in the DF model, where ρm is estimated with only ρT. The DF model with the TDA (the DF-TDA model) well reproduces experimental data for 3,4He scattering with no adjustable parameter, and its theoretical validity can be confirmed by using multiple scattering theory [23–25]. Furthermore, as a practical model of the DF-TDA model, we also proposed the double-single folding (DSF) model for 3,4He scattering, in which the optical potential between 3,4He and the target is derived by folding the localized NA optical potential of the SF model with the 3,4He densities.

The DF-TDA model and the DSF model don’t take into account projectile-excitation effects. In fact, the effects are quite small for 3He and 4He scattering at high incident energies. For 3He scattering at low incident energies, breakup effects of 3He are slightly important, and CDCC with the SF model well reproduces the experimental data as mentioned in Ref. [22]. Thus we have concluded that CDCC with the SF model is a microscopic approach to describe for scattering of 0s-shell nuclei. However it is hard to apply CDCC with the DSF model to scattering of heavier projectiles such as 12C and 16O, because of a high computational cost.

As another approach to describe excitation effects of P and T, the Glauber model [26] has been applied to analyses of AA scattering at intermediate energies. In the Glauber model, nucleon degrees of freedom of P and T are treated by the adiabatic approximation and the eikonal approximation. One of advantages of the
Glauber model is to construct the optical potential including multiple-scattering effects from the phase-shift function [26–28]. The derived optical potential is called the eikonal potential. Although the eikonal potential has been used to analyses of scattering of weakly-binding projectiles in the previous works [27, 28], the validity of the eikonal potential for incident energies and its property have never been clarified.

In this paper, we investigate the validity of the eikonal potential compared with results of CDCC with the SF model for \(d\) scattering at 20–200 MeV/nucleon, which is described by \(n + p + T\) three-body model. The optical potentials for \(n-T\) and \(p-T\) systems are derived by the SF model with the Melbourne \(g\) matrix. In the present analysis of \(d\) scattering, the calculated elastic cross section with the eikonal potential is in good agreement with that with CDCC at high incident energies, and the eikonal potential well simulates the adiabatic calculation. Thus we conclude that the eikonal potential is regarded as a microscopic optical potential including projectile-excitation effects within the adiabatic approximation.

The paper is organized as follows. In Sec. II, we describe the theoretical framework to derive the eikonal potential for \(d\) scattering. In Sec. III, numerical results are shown, and properties of the eikonal potential are discussed. Finally, we give a summary in Sec. IV.

II. THEORETICAL FRAMEWORK

A. Glauber model

In the present study, we describe \(d\) scattering on \(T\) by a \(p + n + T\) three-body model. The scattering is described by the three-body Schrödinger equation

\[
-\frac{\hbar^2}{2\mu} \nabla^2 \Psi + U(r, R) + h_d - E \Psi = 0,
\]

where \(R\) and \(\mu\) are the relative coordinate and the reduced mass of the \(d-T\) system, respectively. The potential \(U\) between \(d\) and \(T\) is represented by

\[
U(r, R) = U_n(r, R) + U_p(r, R) + V_C(R),
\]

where \(U_n\) (\(U_p\)) is the optical potential between \(n\) (\(p\)) and \(T\). In the present analysis, we neglect Coulomb breakup processes, and the the Coulomb potential \(V_C\) thus depends on only \(R\). As the internal Hamiltonian \(h_d\) for \(d\), we adopt a simple form with the Ohmura potential between \(p\) and \(n\) [29].

Initially, \(d\) is the ground state, and has \(\hbar K\) as the relative momentum according to \(R\) at \(Z = -\infty\), where \(Z\) is the \(z\)-axis component of \(R\), and \(K\) is set to be parallel to \(Z\). Under this initial condition, \(\Psi\) is represented by

\[
\Psi(r, R) \rightarrow e^{iKZ+...} \Phi_0(r),
\]

where the “…” represents effects of the Coulomb distortion. The ground state wave function \(\Phi_0\) of \(d\) satisfies

\[
h_d \Phi_0(r) = \epsilon_0 \Phi_0(r),
\]

and the total energy conservation is defined as

\[
\frac{\hbar^2 K^2}{2\mu} = E - \epsilon_0 = E_0.
\]

The Glauber model is based on the adiabatic approximation and the eikonal approximation. In the adiabatic approximation, \(\Psi\) is described as a product of a plane wave by a new function \(\hat{\Psi}\),

\[
\Psi(r, R) = e^{iKZ} \hat{\Psi}(r, R).
\]

Inserting Eq. (6) into Eq. (1) with the adiabatic approximation, the equation for \(\hat{\Psi}\) is obtained as

\[
\left[ -\frac{\hbar^2}{2\mu} \nabla^2_R - i\hbar v \frac{\partial}{\partial Z} + U(r, R) \right] \hat{\Psi}(r, R) = 0,
\]

where \(v = \hbar K/\mu\). As an essence of the eikonal approximation, the second derivatives of \(\hat{\Psi}\) for \(R\) is neglected by assuming that \(\hat{\Psi}\) varies smoothly with \(R\). Accordingly, Eq. (7) is rewritten as

\[
i\hbar v \frac{\partial}{\partial Z} \hat{\Psi}(r, R) = U(r, R) \hat{\Psi}(r, R).
\]

The conditions for the validity of the eikonal approximation are well known as

\[
\frac{|U|}{E_0} << 1, \quad a_U K >> 1,
\]

where \(a_U\) is the radius of \(U\).

If \(U\) does not include \(V_C\), Eq. (8) can be solved analytically under the initial condition as

\[
\hat{\Psi}(r, R) = \exp \left[ -\frac{i}{\hbar v} \int_{-\infty}^{Z} U(r, b, Z')dZ' \right] \Phi_0(r),
\]

where \(b\) is the projection of \(R\) onto the \(x-y\) plane. In the eikonal approximation, the elastic scattering amplitude is given by

\[
f_{el}(q) = \frac{iK}{2\pi} \int db e^{-iqb} [1 - S(b)],
\]

where the forward-scattering approximation is adopted, and \(q\) is the transfer momentum. The eikonal \(S\) matrix, \(S\), is defined as

\[
S(b) = \langle \Phi_0 | e^{i\chi(r, b)} | \Phi_0 \rangle
\]

with the phase-shift function

\[
\chi(r, b) = -\frac{1}{hv} \int_{-\infty}^{\infty} U(r, R)dZ.
\]
This expression is valid for $U$ without $V_C$. When $U$ includes $V_C$, we need a special treatment because of the well-known logarithm divergence of $\chi$.

To describe scattering with the Coulomb interaction in the Glauber model, various approaches have been proposed so far. In this study, we apply the simplest way with the sharp cut screened Coulomb potential with the cut off radius $a_C$, and the elastic scattering amplitude is thus rewritten as

$$f_{el}^{\text{G}}(q) = [f_{\text{Ruth}}(q) + f_N(q)] e^{-2i \eta \ln(2ka_C)}, \quad (14)$$

where $f_{\text{Ruth}}$ is the the Rutherford amplitude and $\eta$ is the Zommerfeld parameter. $f_N$ represents the nuclear part of the scattering amplitude given by

$$f_N(q) = \frac{iK}{2\pi} \int d b \, e^{-i q \cdot b + 2i \eta \ln(kb)} \left[ 1 - S_N(b) \right] \quad (15)$$

with the nuclear part of the eikonal $S$ matrix

$$S_N(b) = \langle \Phi_0 | e^{i \chi_N(r, b)} | \Phi_0 \rangle \quad (16)$$

and the phase-shift function

$$\chi_N(r, b) = -\frac{1}{\hbar v} \int_{-\infty}^{\infty} [U_n(r, R) + U_p(r, -R)] dZ. \quad (17)$$

Here it should be noted that although $f_{el}$ of Eq. (14) depends on $a_C$, the differential cross section defined as $|f_{el}^{\text{G}}(q)|^2$ does not depend on $a_C$. In this paper, we refer to the calculation with $f_{el}^{\text{G}}(q)$ as the Glauber model.

### B. Eikonal potential model

According to the approach proposed by Glauber, we can derive the local optical potential $U_{\text{opt}}$ in the following. First, we introduce a new phase-shift function $\tilde{\chi}_N$ to reproduce $S_N$ obtained from Eqs. (16) and (17);

$$S_N(b) = e^{i \tilde{\chi}_N(b)}. \quad (18)$$

Here it should be noted that $S_N$ depends on only $b = |b|$ since $U_n$ and $U_p$ are taken into account only the central parts as mentioned later. Then we assume that $\tilde{\chi}_N$ is described by the local and spherical optical potential $U_E$ as

$$\tilde{\chi}_N(b) = -\frac{1}{\hbar v} \int_{-\infty}^{\infty} U_E(R) dZ. \quad (19)$$

Finally, Solving Eq. (19) for $U_E$, we obtain the following form

$$U_E(R) = \frac{\hbar v}{\pi} \frac{1}{R} \frac{d}{dR} \int_{R}^{\infty} b db \, \frac{\tilde{\chi}_N(b)}{\sqrt{b^2 - R^2}}. \quad (20)$$

The details for derivation of $U_E$ are shown in Refs. [26–28]. In this paper, the calculation with $U_E$ is called the eikonal potential model (the EP model), in which the two-body Schrödinger equation for $d$-$T$ scattering is defined as

$$\left[ -\frac{\hbar^2}{2\mu} \nabla_R^2 + U_E(R) + V_C(R) - E_0 \right] \Psi_{\text{EP}}(R) = 0. \quad (21)$$

Here it should be noted that $U_E$ includes breakup effects of $d$ through the Glauber model. Meanwhile, the optical potential in the DSF model (DSF potential) is defined as

$$U_{\text{DSF}}(R) = \langle \Phi_0 | U_n(r, R) + U_p(r, -R) | \Phi_0 \rangle. \quad (22)$$

in which breakup effects of $d$ are not taken into account.

### C. NA optical potential

In the present analysis, we derive $U_n$ and $U_p$ by the $g$-matrix folding model. In the $g$-matrix folding model, the optical potential is obtained by folding the $g$ matrix with the target density, and the knock-on exchange processes between the interacting two nucleons are considered as the dominant component of the antisymmetrization [30, 31]. The folding potential is generally nonlocal because of the knock-on exchange process, but the nonlocality is well localized by the Brieva-Rook method [12]. The folding potential is composed of the direct (DR) and exchange (EX) terms

$$U_j(R_j) = U_j^{\text{DR}}(R_j) - U_j^{\text{EX}}(R_j) \quad (23)$$

where $j = n$ or $p$, and $R_j$ is the relative coordinate between particle $j$ and $T$. Each term is written with the one-body and mixed densities, $\rho_T$ and $\tilde{\rho}_T$, as

$$U_j^{\text{DR}}(R_j) = \sum_{\nu} \int d r_T g^{\text{DR}}_{\nu}(s, \rho) \rho^{(\nu)}_T (r_T), \quad (24)$$

$$U_j^{\text{EX}}(R_j) = \sum_{\nu} \int d r_T g^{\text{EX}}_{\nu}(s, \rho) \tilde{\rho}^{(\nu)}_T (r_T, R_j) J_0(K(R_j) s). \quad (25)$$

Here the index $\nu$ represents the $z$-component of isospin of the nucleons in T, $r_T$ is the internal coordinate of T, and $s = r_T - R_j$. The $g$ matrix $g^{\text{DR(EX)}}$ is the direct (exchange) term of the $g$ matrix among the considering nucleons. The mixed density can be described by using the one-body density with the local Fermi gas approximation [32] as

$$\tilde{\rho}_T (r_T, R_j) = \rho_T (|r_T - s/2|) \frac{3j_1(k^F_E s)}{k^F_E s}. \quad (26)$$

In the present study, we employ the Melbourne $g$ matrix [17] constructed from Bonn-B nucleon-nucleon interaction [33]. The Melbourne $g$ matrix well reproduce nucleon elastic scattering with and without localization of the exchange term [17, 34]. For $^{208}$Pb, the matter density is calculated by the spherical Hartree-Fock method with the Gogny-D1S interaction [35] and the spurious
Differential cross sections \( d\sigma/d\Omega \) as a function of transfer momentum \( q \) for \( d \) scattering from (a) \( ^{12}\text{C} \) and (b) \( ^{208}\text{Pb} \) targets at \( E_{\text{in}}/A_{\text{P}} = 20-200 \text{ MeV} \). The solid line represents the result of the EP model, the dashed line denotes the result of the DSF model, and the dot-dashed line corresponds to the result of CDCC, respectively. Each cross section is multiplied by the factor shown in the figure.

Center-of-mass (c.m.) motion is removed with the standard procedure [9]. For \( ^{12}\text{C} \), we take the phenomenological proton-density determined from electron scattering [36]; here the finite-size effect of proton charge is unfolded with the standard procedure [37], and the neutron density is assumed to have the same geometry as the proton one.

### III. RESULTS AND DISCUSSIONS

We analyze \( d \) scattering from \( ^{12}\text{C} \) and \( ^{208}\text{Pb} \) targets at \( E_{\text{in}}/A_{\text{P}} = 20-200 \text{ MeV} \) by using the DSF model, the EP model, and CDCC, in which the calculation of CDCC is regarded as the exact calculation. Figure 1 shows the \( q \) dependence of \( d\sigma/d\Omega \) for \( d \) scattering from (a) \( ^{12}\text{C} \) target and (b) \( ^{208}\text{Pb} \) target at the incident energies \( E_{\text{in}}/A_{\text{P}} = 20-200 \text{ MeV} \). The dashed and dotted lines stand for the results of the DSF model and CDCC, respectively. The result of the EP model is shown by the solid line. The difference between the dashed and dotted lines represents breakup effects of \( d \) on the elastic scattering that become more important at forward angles as the incident energy decreases. One sees that the EP model well reproduces the results of CDCC for both \( ^{12}\text{C} \) and \( ^{208}\text{Pb} \) targets as the incident energy increases. For \( E_{\text{in}}/A_{\text{P}} \leq 40 \text{ MeV} \), the solid line is slightly different from the dot-dashed line for the oscillation behavior. We will discuss the difference later.

As shown in Fig. 1, the EP model well simulates the results of CDCC. One of advantages of the EP model is available to derive an optical potential including breakup effects of \( d \) on the elastic scattering as the eikonal potential. In Fig. 2, we show the optical potentials derived by the EP model and the DSF model for \( d \) scattering from (a) \( ^{12}\text{C} \) and (b) \( ^{208}\text{Pb} \) targets at \( E_{\text{in}}/A_{\text{P}} = 40 \text{ MeV} \). The solid line represents the eikonal potential, and the dashed line corresponds to the DSF potential that doesn’t include breakup effects of \( d \). The dot-dashed line stands for the dynamical polarization (DP) potential \( U_{\text{DP}} \) defined as

\[
U_{\text{DP}}(R) = U_{E}(R) - U_{\text{DSF}}(R),
\]

which represents breakup effects of \( d \) on the optical potential between \( d \) and \( T \). One sees that the DP potential is repulsive in the peripheral region of nuclei and absorptive in the whole region. The behavior of the DP potential is the same as that in the previous works [27, 28].
FIG. 2: The $R$ dependence of the potentials for $d$ scattering from (a) $^{12}$C and (b) $^{208}$Pb targets at $E_{\text{in}}/A_{P} = 40$ MeV. The solid and dashed lines represent the eikonal potential and the DSF potential. The dot-dashed line stands for the dynamical polarization potential.

important point of the present work is that the eikonal potential is derived by the microscopic approach based on the folding model with the Melbourne $g$ matrix.

In order to investigate properties of the EP model, we compare the result of the EP model with that of the Glauber model. From the definition of the EP model, it is clear that if we apply the eikonal approximation to the EP model, the obtained eikonal $S$ matrix becomes the same as $S_{\text{N}}$ of the Glauber model. Thus the EP model might be regarded as the Glauber model without the eikonal approximation, i.e. the calculation with only the adiabatic approximation. To clear this point, we also perform CDCC with the adiabatic approximation (adiabatic-CDCC), where energies for all excited states are replaced by the ground-state energy.

Figure 3 shows the results of the EP model, the Glauber model, and adiabatic-CDCC for $d$ scattering from the $^{12}$C target (a) and the $^{208}$Pb target (b). The solid, dashed, and dot-dashed lines represent the results of the EP model, the Glauber model, and adiabatic-CDCC, respectively.

First, we discuss the results of the EP model and adiabatic-CDCC. For $E_{\text{in}}/A_{P} \geq 80$ MeV, the results of adiabatic-CDCC are in good agreement with those of CDCC, and the adiabatic approximation well works. For $E_{\text{in}}/A_{P} \leq 40$ MeV, the EP model well simulates the result of the adiabatic-CDCC, but the discrepancy between the EP model and CDCC is not negligible for the oscillation behavior as mentioned before. Therefore the discrepancy comes from ambiguities of the adiabatic approximation. Then we conclude that the eikonal potential corresponds to the microscopic optical potential including breakup effects evaluated by the adiabatic approximation.

Next, we focus on the difference between the EP model and the Glauber model, which represents the validity of the eikonal approximation. The eikonal approximation is valid for the scattering at high incident energies as mentioned in Eq. (9). In fact, one can see that the discrepancy of the two models at forward angles becomes small as the incident energy increases for both $^{12}$C and $^{208}$Pb targets. For $d$ scattering from $^{208}$Pb target at $E_{\text{in}}/A_{P} \leq 40$ MeV, the oscillation behavior of the Glauber model is much different from that of the EP model. The large difference comes from the strong Coulomb interaction in addition to the scattering at the low energy because the eikonal approximation cannot take into account the Coulomb interaction accurately as mentioned in Sec. II. On the other hand, the EP model simulates the results of adiabatic-CDCC at forward angles for scattering at low incident energies, and then the EP model is regarded as the Glauber model without ambiguities of the eikonal approximation.

IV. SUMMARY

We construct a microscopic optical potential including breakup effects of $d$ scattering with the Glauber model, in which a NA potential is derived by the folding model with
FIG. 3: Differential cross sections $d\sigma/d\Omega$ as a function of transfer momentum $q$ for $d$ scattering from (a) $^{12}$C and (b) $^{208}$Pb target at $E_{in}/A_{p} = 20$–200 MeV. The solid line represents the result of the EP model, the dashed line denotes the result of the Glauber model, and the dot-dashed line corresponds to the result of adiabatic-CDCC, respectively. Each cross section is multiplied by the factor shown in the figure.

the Melbourne $g$ matrix. The microscopic optical potential is referred to as the eikonal potential. In order to confirm the validity of the eikonal potential, we compared the EP model with CDCC for the elastic cross sections for $d$ scattering from $^{12}$C and $^{208}$Pb at $E_{in}/A_{p} = 20$–200 MeV. As the result, the EP model well reproduces the results of CDCC as the incident energy increases. Furthermore we found that the difference between the results of the EP model and CDCC at low incident energies comes from mainly the adiabatic approximation. In fact, the EP model simulates the result of adiabatic-CDCC for $E_{in}/A_{p} \leq 40$ MeV. Thus, the eikonal potential constructed by the $g$-matrix folding model seems to be a microscopic optical potential including projectile-excitation effects within the adiabatic approximation, and the EP model is regarded as the Glauber model without ambiguities of the eikonal approximation.

In the present analysis, we presented the validity of the EP model for $d$ scattering, which is one of the simplest system for scattering of a two-body projectile. One of the advantage of the EP model is to be applicable to analyses for scattering of many-body projectiles as well as the Glauber model. In the forthcoming paper, we will report analyses for scattering of two-neutron halo nuclei such as $^6$He and $^{11}$Li with the EP model. Furthermore, we will discuss Coulomb breakup effects, which are omitted in the present calculation. The treatment of Coulomb breakup processes within the eikonal approximation has been discussed in some papers [38, 39], and the problem is an important subject for the eikonal calculation.

Acknowledgements

The authors are grateful to Prof. M. Yahiro for fruitful discussions. This work is supported in part by by Grant-in-Aid for Scientific Research (Nos. 25400266, and 26400278) from Japan Society for the Promotion of Science (JSPS).

[1] M. Kamimura, M. Yahiro, Y. Iseri, Y. Sakuragi, H. Kameyama, and M. Kawai, Prog. Theor. Phys. Suppl. 89, 1 (1986).
[2] N. Austern, Y. Iseri, Y. Sakuragi, M. Kawai, G. Rawitscher, and M. Yahiro, Phys. Rep. 154, 125 (1987).
[3] M. Yahiro, K. Ogata, T. Matsumoto, and K. Minomo, Prog. Theor. Exp. Phys. (2012) 01A206.
[4] H. F. Arellano, F. A. Brieva, and W. G. Love, Phys. Rev. C 52, 301 (1995).
[5] K. Minomo, K. Ogata, M. Kohno, Y. R. Shimizu, and M. Yahiro, J. Phys. G 37, 085011 (2010)
[6] B. Sinha, Phys. Rep. 20, 1 (1975).
[7] B. Sinha and S. A. Moszkowski, Phys. Lett. B 81, 289 (1979).
[8] D. T. Khoa, W. von Oertzen, H. G. Bohlen, and S. Ohkubo, J. Phys. G 34, R111 (2007).
[9] T. Furumoto, Y. Sakuragi, and Y. Yamamoto, Phys. Rev. C 82, 044612 (2010).
[10] G. Bertsch, J. Borysowicz, H. McManus, and W. G. Love, Nucl. Phys. A 284, 299 (1977).
[11] J. -P. Jeukenne, A. Lejeune, and C. Mahaux, Phys. Rev. C 16, 80 (1977); J. -P. Jeukenne, A. Lejeune, and C. Mahaux, Phys. Rep. 25, 83 (1976).
[12] F. A. Brieva and J. R. Rook, Nucl. Phys. A 291, 299 (1977); ibid. 291, 317 (1977); ibid. 297, 206 (1978).
[13] G. R. Satchler and W. G. Love, Phys. Rep. 55, 183 (1979).
[14] G. R. Satchler, “Direct Nuclear Reactions”, Oxford University Press, (1983).
[15] N. Yamaguchi, S. Nagata, and T. Matsuda, Prog. Theor. Phys. 70, 459 (1983); N. Yamaguchi, S. Nagata, and J. Michiyama, Prog. Theor. Phys. 76, 1289 (1986).
[16] L. Rikus, K. Nakano, and H. V. Von Geramb, Nucl. Phys. A 414, 413 (1984); L. Rikus, and H. V. Von Geramb, Nucl. Phys. A 426, 496 (1984).
[17] K. Amos, P. J. Dortmans, H. V. Von Geramb, S. Karataglidis, and J. Raynal, in Advances in Nuclear Physics, edited by J. W. Negele and E. Vogt(Plenum, New York, 2000) Vol. 25, p. 275.
[18] S. M. Saleem and W. Haider, J. Phys. G 28, 1313 (2002).
[19] T. Furumoto, Y. Sakuragi, and Y. Yamamoto, Phys. Rev. C 78, 044610 (2008); ibid., 79, 011601(R) (2009); ibid., 80, 044614 (2009).
[10] M. Toyokawa, M. Yahiro, T. Matsumoto, K. Minomo, K. Ogata and M. Kohno, Phys. Rev. C 92, 024618 (2015).
[11] K. Egashira, K. Minomo, M. Toyokawa, T. Matsumoto and M. Yahiro, Phys. Rev. C 89, 064611 (2014).
[12] M. Toyokawa, T. Matsumoto, K. Minomo and M. Yahiro, Phys. Rev. C 91, 064610 (2015).
[23] K. M. Watson, Phys. Rev. 89, 575 (1953).
[24] A. K. Kerman, H. McManus, and R. M. Thaler, Ann. Phys. 8, 551 (1959).
[25] M. Yahiro, K. Minomo, K. Ogata, and M. Kawai, Prog. Theor. Phys. 120, 767 (2008).
[26] R. J. Glauber, in Lectures in Theoretical Physics, edited by W. E. Brittin (Interscience, New York, 1959), Vol. 1, pp 315–414.
[27] K. Yabana, Y. Ogawa, and Y. Suzuki, Phys. Rev. C 45, 2909 (1992).
[28] B. Abu-Ibrahim and Y. Suzuki, Phys. Rev. C 62, 034608 (2000).
[29] T. Ohmura, B. Imanishi, M. Ichimura, and M. Kawai, Prog. Theor. Phys. 43, 347 (1970).
[30] Y. C. Tang, M. LeMere, and D. R. Thompson, Phys. Rep. 47, 167 (1978).
[31] K. Aoki and H. Horiuchi, Prog. Theor. Phys. 69, 857 (1983), and references therein.
[32] J. W. Negele and D. Vautherin, Phys. Rev. C 5, 1472 (1972).
[33] R. Machleidt, K. Holinde, and Ch. Elster, Phys. Rep. 149, 1 (1987).
[34] M. Toyokawa, K. Minomo, and M. Yahiro, Phys. Rev. C 88, 054602 (2013).
[35] J. F. Berger, M. Giroud, and D. Gogny, Comput. Phys. Commun. 63, 365 (1991).
[36] H. de Vries, C. W. de Jager, and C. de Vries, At. Data Nucl. Data Tables 36, 495 (1987).
[37] R. P. Singh, M. W. S. Macauley, and P. K. A. De Witt Hubertus, Nucl. Instrum. and Method 148, 113 (1978).
[38] K. Ogata, M. Yahiro, Y. Iseri, T. Matsumoto, and M. Kamimura, Phys. Rev. C 68, 064609 (2003).
[39] P. Capel, D. Baye, and Y. Suzuki, Phys. Rev. C 78, 054602 (2008).