Pauli blocking effects for nucleon-nucleus scattering

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(Received 19 January 1993)

Theoretical predictions for nucleon-nucleus scattering are customarily based on multiple-scattering formalisms in which the projectile-nucleon–struck-target-nucleon two-body subsystem is antisymmetrized and in which the target wave function is antisymmetric with respect to the nucleon constituents. Although formalisms exist for nucleon-nucleus scattering which account for full $(A+1)$-body antisymmetrization, these theories have not been adequately implemented in numerical applications. Pauli blocking of the struck target nucleon in intermediate scattering states is implicitly included in standard optical potential calculations which consistently include terms through second-order in the projectile-target nucleon scattering t matrix. In this work the multiple scattering expansion of the fully antisymmetrized nucleon-nucleus optical potential is organized so as to make explicit corrections to the standard optical potential due to Pauli blocking of the projectile nucleon in intermediate scattering states. Numerical calculations are presented for the resulting density dependent, projectile-target nucleon effective interaction and comparison with a previous density dependent model is given. It is shown that density dependent effective interaction t matrices for nucleon-nucleus scattering calculations should include Pauli blocking of just the projectile nucleon and binding potential corrections for just the target nucleon.

PACS number(s): 24.10.−i, 24.10.Cn, 25.40.Cm, 25.40.Ve

I. INTRODUCTION

In the nonrelativistic multiple-scattering formalisms of Watson [1] and Kerman, McManus, and Thaler (KMT) [2] it is assumed that the projectile particle is distinguishable from the nucleon constituents of the target nucleus. The identity of the A target nucleons is incorporated in the theory by projecting antisymmetric target nucleus states in all intermediate scattering states in the formal definition of the projectile-nucleus scattering t matrix [2,3]. Expansions of the optical potential for elastic scattering in terms of quasi-two-body t matrices result in a linear term which is summed over the A target nucleons and a second-order term which is proportional to two-body correlations in the target nucleus wave function [2,3] plus higher-order terms. The quasi-two-body t matrices are in reality $(A+1)$-body operators, but for the leading Born term only depend on the coordinates of the projectile and *one* active target nucleon. Intermediate scattering states of the active target nucleon used in constructing these quasi-two-body t matrices are not restricted with respect to the states occupied by the other $(A-1)$ target nucleons, i.e., the quasi-two-body t matrices are *not* Pauli blocked. Even so, this formalism accounts for Pauli blocking of the target nucleons in intermediate scattering states by construction. Explicit Pauli blocking corrections do not appear in the final expression for the optical potential due to certain cancellations which occur among the second-order terms [3].

Application of the Watson or KMT formalisms to the case of nucleon-nucleus scattering is usually carried out by way of the Takeda and Watson (TW) [4] prescription. According to this procedure the unsymmetrized, nucleon-nucleon (NN) quasi-two-body t matrices in the Watson or KMT formalisms are everywhere replaced with operators which are antisymmetric with respect to the projectile nucleon and struck target nucleon labels. Takeda and Watson argue that the leading correction to this prescription involves a three-body exchange potential which should be small, and indeed calculations [5] show that this term is negligible at medium energies.

Intuitively, however, one would not expect the simple TW prescription to adequately account for all of the effects arising in nucleon-nucleus scattering due to the imposition of full $(A+1)$-body antisymmetrization. In particular, Pauli blocking of the projectile nucleon in intermediate scattering states is not accounted for by simply invoking the TW procedure. Experience with Pauli blocking effects in nucleon-nucleon g-matrix calculations [6] for infinite nuclear matter and for density dependent, effective interaction t matrices [7,8] shows that such effects can be quite important.

In this work the formalism of Picklesimer and Thaler [9], which incorporates full $(A+1)$-body antisymmetrization, is used to study Pauli blocking corrections to the standard nucleon-nucleus optical potential. The additional Pauli blocking corrections explicitly obtained here are a consequence of the identity of the projectile and the nucleon constituents of the target.

The fully antisymmetrized optical potential is expanded to second order in free, antisymmetrized two-body t matrices. Cancellations occur among the second-order terms analogous to that in Appendix A of Ref. [3] such that target nucleon Pauli blocking corrections are again implicit in the final optical potential. The leading corrections to the lowest order "$t\rho$" optical potential are then shown to be due to Pauli blocking of the projectile nucleon, the action of the nuclear binding potential on the struck target nucleon in intermediate states (included in standard optical potential formalisms [2,3]), three-body exchange effects [5], and two-body correlations in the target nucleus wave function.

In Sec. II the full $(A+1)$-body antisymmetrized optical potential formalism of Ref. [9] is considered. Practical implementation of this formalism is discussed in Sec. III, followed by a brief presentation of numerical results in Sec. IV. Conclusions are given in Sec. V.

II. ANTISYMMETRIZED OPTICAL POTENTIAL FORMALISM

Reference [9] presents a nonrelativistic $(A+1)$-body antisymmetrized optical potential formalism for describing nucleon-nucleus elastic scattering. The optical potential and the resulting nucleon-nucleus t matrix are assumed to operate on *unsymmetrized* $(A+1)$-body states where the wave functions describing the $(A+1)$-body system are not antisymmetrized with respect to the labels of the projectile nucleon and target nucleon constituents [10]. It is important to note that the unsymmetrized $(A+1)$-body states referred to here and in Ref. [9] include fully antisymmetrized A-body states of the target nucleus. Dynamical effects due to projectile and target nucleon identity are therefore included in the construction of the optical potential.

From Ref. [9] the $(A+1)$-body antisymmetrized transition operator appropriate for elastic scattering is given by [11]

$$\hat{T}=\sum_{i=1}^{A} v_{0i}\overline{\mathcal{A}}(1+G_\alpha T)\ , \tag{1}$$

where v_{0i} represents the two-body interaction between the projectile (0) and the ith target nucleon. Here, the unsymmetrized channel corresponds to the projectile nucleon label being (0) and the target constituent labels running from $1,2,\ldots,A$. In Eq. (1) $\overline{\mathcal{A}}$ is the antisymmetrizer for the projectile nucleon and A target nucleon labels. It is given by

$$\overline{\mathcal{A}}=1-\sum_{i=1}^{A} E_{0i}\ , \tag{2}$$

where E_{0i} is the exchange operator for projectile (0) and identical target nucleon (i). The unsymmetrized $(A+1)$-body propagator G_α is given by

$$G_\alpha=(E-H_0-H_A+i\epsilon)^{-1}\ , \tag{3}$$

where H_0 is the projectile kinetic energy operator and H_A is the nuclear Hamiltonian. In Eq. (1) T represents the usual unsymmetrized scattering operator given by

$$T=\sum_{i=1}^{A} v_{0i}(1+G_\alpha T)\ . \tag{4}$$

The $(A+1)$-body antisymmetrized optical potential which generates $\hat{T}$ is defined by

$$\hat{T}'=\hat{U}'+\hat{U}'G_\alpha\mathcal{P}\hat{T}'\ , \tag{5}$$

where

$$\hat{T}'=\frac{A-1}{A}\hat{T}\ , \tag{6}$$

$(A-1)/A$ is the usual KMT scaling factor [2], and $\mathcal{P}$ projects the unsymmetrized elastic channel [2]. Using Eqs. (1) and (4)–(6) and eliminating the transition operators, the following expression for $\hat{U}'$ is obtained:

$$\hat{U}'=\frac{A-1}{A}\sum_i v_{0i}\overline{\mathcal{A}}+\hat{U}'G_\alpha\sum_i v_{0i} -\frac{A-1}{A}\hat{U}'G_\alpha\mathcal{P}\sum_i v_{0i}\overline{\mathcal{A}}\ . \tag{7}$$

Using the definition

$$\hat{U}'\equiv\sum_{i=1}^{A}\hat{U}_i' \tag{8}$$

and introducing an auxiliary, quasi-two-body operator $\hat{t}_{0i}$ given by

$$\hat{t}_{0i}=v_{0i}\overline{\mathcal{A}}+v_{0i}\overline{\mathcal{A}}G_\alpha t_{0i}\ , \tag{9}$$

where

$$t_{0i}=v_{0i}+v_{0i}G_\alpha t_{0i}\ , \tag{10}$$

the two-body interactions in Eq. (7) can be eliminated resulting in

$$\hat{U}_i'=\frac{A-1}{A}\hat{t}_{0i}+\hat{U}'G_\alpha t_{0i}-\frac{A-1}{A}\hat{U}'G_\alpha\mathcal{P}\hat{t}_{0i}-\hat{U}_i'G_\alpha t_{0i}\ . \tag{11}$$

Keeping terms through second order, the antisymmetrized optical potential can be written as

$$\hat{U}'=\frac{A-1}{A}\sum_i\hat{t}_{0i} +\frac{A-1}{A}\left[\sum_{i\neq j}\hat{t}_{0i}G_\alpha t_{0j}-\frac{A-1}{A}\sum_{ij}\hat{t}_{0i}G_\alpha\mathcal{P}\hat{t}_{0j}\right] +\cdots\ , \tag{12}$$

which is analogous to the expression for the unsymmetrized optical potential in Ref. [3].

Introducing the free NN two-body t matrix t_{0i}^f where

$$t_{0i}^f=v_{0i}+v_{0i}gt_{0i}^f \tag{13}$$

and g is the free NN propagator, $\hat{t}_{0i}$ can be reexpressed as

$$\hat{t}_{0i}=t_{0i}^f\overline{\mathcal{A}}+t_{0i}^f\overline{\mathcal{A}}G_\alpha t_{0i}-t_{0i}^f g\hat{t}_{0i}\ . \tag{14}$$

Defining $\mathcal{A}_{0i}\equiv1-E_{0i}$ and $\varepsilon_{0i}\equiv-\sum_{j\neq i}E_{0j}$, where $\overline{\mathcal{A}}=\mathcal{A}_{0i}+\varepsilon_{0i}$, and expanding to second order in the free NN t matrix $\hat{t}_{0i}$ becomes

$$\hat{\tau}_{0i} = t^f_{0i}\mathcal{A}_{0i} + t^f_{0i}\varepsilon_{0i} + t^f_{0i}\overline{\mathcal{A}}G_\alpha t^f_{0i} - t^f_{0i} g t^f_{0i}(\mathcal{A}_{0i}+\varepsilon_{0i}) + \cdots \, . \quad (15)$$

If v_{0i} is chosen to vanish for nonantisymmetric NN states then the antisymmetrizer $\mathcal{A}_{0i}$ becomes redundant and $t^f_{0i}\mathcal{A}_{0i}$ can be written simply as t^f_{0i}. Thus $\hat{\tau}_{0i}$ can be expanded as

$$\hat{\tau}_{0i} = t^f_{0i} + t^f_{0i}\varepsilon_{0i} - t^f_{0i} g t^f_{0i}\varepsilon_{0i} + t^f_{0i}(\overline{\mathcal{A}}G_\alpha - g)t^f_{0i} + \cdots \, . \quad (16)$$

Therefore, to second order in t^f_{0i} the $(A+1)$-body antisymmetrized optical potential is written as

$$\begin{aligned}\hat{U}' = {} & \frac{A-1}{A}\sum_{i=1}^{A}[t^f_{0i} + t^f_{0i}\varepsilon_{0i} - t^f_{0i} g t^f_{0i}\varepsilon_{0i} + t^f_{0i}(\overline{\mathcal{A}}G_\alpha - g)t^f_{0i}] \\ & + (A-1)^2\left[\frac{1}{A(A-1)}\sum_{i\neq j} t^f_{0i}G_\alpha t^f_{0j} - \frac{1}{A^2}\sum_{ij} t^f_{0i}G_\alpha \mathcal{P} t^f_{0j}\right] \\ & + (A-1)^2\left[\frac{1}{A(A-1)}\sum_{i\neq j} t^f_{0i}\varepsilon_{0i}G_\alpha t^f_{0j} - \frac{1}{A^2}\sum_{ij} t^f_{0i}G_\alpha \mathcal{P} t^f_{0j}\varepsilon_{0j} - \frac{1}{A^2}\sum_{ij} t^f_{0i}\varepsilon_{0i}G_\alpha \mathcal{P} t^f_{0j} \right. \\ & \left. - \frac{1}{A^2}\sum_{ij} t^f_{0i}\varepsilon_{0i}G_\alpha \mathcal{P} t^f_{0j}\varepsilon_{0j}\right] + \cdots \, . \end{aligned} \quad (17)$$

In the limit $\overline{\mathcal{A}} \to 1$, $\hat{U}'$ in Eqs. (12) and (17) reduces to the usual unsymmetrized optical potential. In the limit $\varepsilon_{0i} \to 0$ for the above choice of v_{0i} the optical potential in Eqs. (12) and (17) reduces to the unsymmetrized optical potential with the TW substitution for t_{0i} and t^f_{0i}, respectively.

The terms in the first set of brackets in Eq. (17) include the usual "$t\rho$" term plus antisymmetrization and binding potential corrections. The "$t\rho$" potential corresponds to folding the free NN t matrix with the one-body ground-state nuclear density matrix [2,3]. The second term in the first set of brackets in Eq. (17) represents a three-body antisymmetrization correction linear in t^f_{0i}, which was calculated in Ref. [5] and shown to be negligible for medium energies. The third term inside the first set of brackets in Eq. (17) represents a higher-order multiple-scattering correction to the three-body term of Ref. [5]. The fourth term in the first set of brackets represents the leading correction due to Pauli blocking of the projectile nucleon (0) in intermediate scattering states (as a result of the $\overline{\mathcal{A}}$ antisymmetrizer) and binding potential effects on the struck target nucleon (resulting from H_A in G_α). This fourth term is the one of primary interest here.

The terms in the first set of parentheses in Eq. (17) give rise to two-body correlation contributions [2,3] while the terms in the second set of parentheses represent antisymmetrization corrections to the second-order optical potential of Ref. [3]. The latter should be small and are not evaluated here.

The antisymmetrized optical potential defined by

$$\hat{T} = \hat{U} + \hat{U}G_\alpha \mathcal{P}\hat{T} \quad (18)$$

is related to $\hat{U}'$ according to

$$\hat{U} = \frac{A}{A-1}\hat{U}' - \frac{1}{A-1}\hat{U}'G_\alpha \mathcal{P}\hat{U} \, . \quad (19)$$

An expansion for $\hat{U}$ in terms of t^f_{0i} can be readily obtained using Eqs. (17) and (19). Calculations [12] show, however, that the second term on the right-hand side of Eq. (19) is negligible for characteristic optical potentials at medium energies for the moderate to heavy weight target nuclei considered here and in most analyses (e.g., those in Ref. [13]). This term will not be considered further here.

III. APPLICATION OF THE ANTISYMMETRIZED OPTICAL POTENTIAL

The NN-isobar coupled channels model of Refs [8,14] was used to calculate an approximation to the fourth term in the first set of brackets in Eq. (17). In order to do so the intractable $(A+1)$-body operators, $\overline{\mathcal{A}}$ and G_α, in this term were replaced by a projectile nucleon Pauli blocking factor ($Q_{\rm proj}$) for infinite nuclear matter and a modified two-nucleon energy denominator ($1/e$) where the energy of the target nucleon includes kinetic energy and binding potential contributions appropriate for infinite nuclear matter [8]. To the order which Eq. (17) is expanded this resulting correction term, $t^f_{0i}(Q_{\rm proj}/e - g)t^f_{0i}$, can be calculated using the method for solving the Bethe-Goldstone equation developed in Ref. [8] where the projectile Pauli blocking factor and the target nucleon binding potential are assumed to act in intermediate NN scattering states. Therefore the equation to be solved is [subscripts $(0i)$ are suppressed]

$$\psi = \phi + \frac{Q_{\rm proj}}{e} v\psi \, , \quad (20)$$

where ϕ is a NN plane wave and v represents the NN-isobar coupled-channel interaction of Ref. [14]. The resulting t matrix obtained from

$$t\phi = v\psi \tag{21}$$

is related to the desired correction term by way of the following expansion:

$$t = t^f + t^f\left[\frac{Q_{\rm proj}}{e} - g\right]t$$
$$= t^f + t^f\left[\frac{Q_{\rm proj}}{e} - g\right]t^f + \cdots \ . \tag{22}$$

Thus to second order the difference, $t - t^f$, provides an approximation for the fourth term in Eq. (17). Technical details pertaining to the partial-wave solutions of Eqs. (20) and (21) are given in Ref. [8].

The projectile Pauli blocking factor for infinite nuclear matter for NN intermediate states is determined by the requirements

$$Q^{NN}_{\rm proj}(\mathbf{k}'_0, \mathbf{k}'_1, k_F) = \begin{cases} 1 & \text{if } |\mathbf{k}'_0| > k_F \\ 0 & \text{if } |\mathbf{k}'_0| \le k_F \ , \end{cases} \tag{23}$$

where $\mathbf{k}'_0$ and $\mathbf{k}'_1$ are the intermediate momenta in the laboratory for the projectile and target nucleon, respectively, and k_F is the Fermi momentum of nuclear matter. The angle averaged values for $Q_{\rm proj}$, assuming relativistic kinematics, are straightforward to obtain as in Appendix A of Ref. [8]. The result is

$$Q^{NN,\rm AV}_{\rm proj} = \begin{cases} 1 & \text{if } |\gamma\sqrt{1-1/\gamma^2}\epsilon_{N_0} - \gamma k'_{NN}| > k_F \\ 0 & \text{if } |\gamma\sqrt{1-1/\gamma^2}\epsilon_{N_0} + \gamma k'_{NN}| \le k_F \\ \dfrac{1}{2}\left\{1 + \dfrac{\gamma\epsilon_{N_0}}{\sqrt{\gamma^2-1}k'_{NN}}\left[1 - \left[1 - \dfrac{k'^2_{NN} + (\gamma^2-1)\epsilon^2_{N_0} - k_F^2}{\gamma^2\epsilon^2_{N_0}}\right]^{1/2}\right]\right\}, & \text{otherwise} \ , \end{cases} \tag{24}$$

where k'_{NN} is the magnitude of the NN center-of-momentum (COM) system momentum in intermediate states,

$$\epsilon_{N_0} = (k'^2_{NN} + m^2_{N_0})^{1/2} \ ,$$

where m_{N_0} is the projectile nucleon mass, and

$$\gamma = \left[1 - \frac{P^2}{4\epsilon^2_{N_0} + P^2}\right]^{-1/2} , \tag{25}$$

where P was fixed equal to the incident nucleon laboratory momentum.

For $N\Delta$(1232 MeV) intermediate channels the blocking effect computed in Eqs. (A7)–(A9) of Ref. [8] was divided by 2 since in the present application Pauli blocking of the $N\Delta$ channel occurs only when N is the projectile nucleon which occurs half the time. Therefore, the angle averaged projectile Pauli blocking factor for $N\Delta$ intermediate states is assumed to be

$$Q^{N\Delta,\rm AV}_{\rm proj} = \tfrac{1}{2}(1 + Q^{N\Delta}_{\rm AV}) \ , \tag{26}$$

where $Q^{N\Delta}_{\rm AV}$ is given in Eq. (A8) of Ref. [8].

IV. RESULTS

Calculations of the medium modified t matrix in Eq. (21) were done for incident laboratory kinetic energies of 320 and 500 MeV. For the calculations in Ref. [8] explicit Pauli blocking of both the projectile and struck target nucleon was assumed as well as target nucleon binding potential corrections since in this previous work the in-

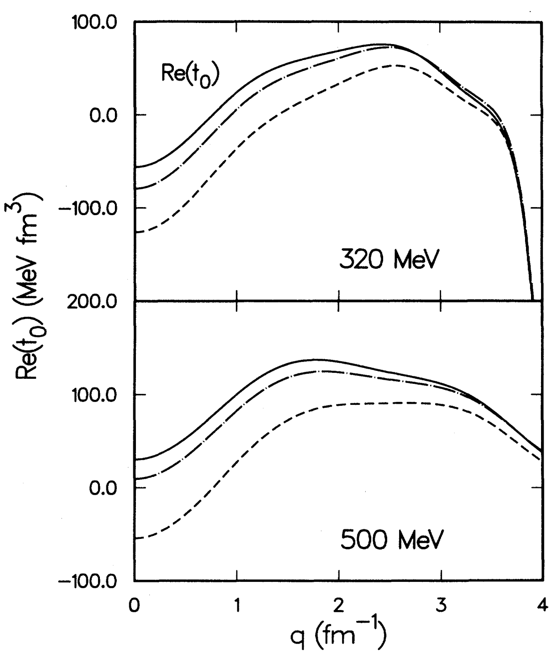

FIG. 1. Nucleon-nucleon–isobar coupled-channel model calculations of the density dependent $\mathrm{Re}(t_0)$ at 320 MeV (upper half) and 500 MeV (lower half). On-momentum-shell t-matrix elements are shown as a function of momentum transfer (q). The dashed, dash-dotted, and solid curves represent the results of calculations made with $k_F=0.0$ fm^{-1}, $k_F=1.4$ fm^{-1} with Pauli blocking of the projectile and binding potentials acting on the target nucleon as described in the text, and $k_F=1.4$ fm^{-1} with, in addition, explicit Pauli blocking of the target nucleon (from Ref. [8]), respectively.

tent was to provide an approximation for the Watson-type [1] *NN* effective interaction t matrix. The calculations presented here are the same as that done in Ref. [8] except for the use of the *projectile only* Pauli blocking factors in Eqs. (24) and (26) rather than the *projectile and target nucleon* Pauli blocking factors in Appendix A of Ref. [8].

The results for the real and imaginary parts of the isoscalar spin-independent and spin-orbit *NN* t matrices [dominant terms for nucleon—(spin zero) nucleus elastic scattering] in the *NN* COM system are shown in Figs. 1–4. The spin-independent and spin-orbit t matrices are denoted in the figures by t_0 and t_0^{LS}, respectively. The t matrices are evaluated "on-momentum-shell" where the magnitudes of the initial and final *NN* COM momenta correspond to the value of the incident beam laboratory momentum in the *NN* COM system (target nucleon assumed at rest in the laboratory) [8]. The values at zero density ($k_F=0$) are shown by the dashed lines. These correspond to the free *NN* scattering values predicted by the *NN*-isobar model [14]. The values corresponding to $k_F=1.4$ fm^{-1} from Ref. [8] are shown by the solid lines while the new results computed here for $k_F=1.4$ fm^{-1} are indicated by the dash-dotted lines. Except for Re(t_0^{LS}) the density dependence is weakened by not Pauli blocking the target nucleon as would be expected. For Re(t_0^{LS}) removing Pauli blocking on the struck target nucleon actually results in somewhat stronger density dependence. At energies above 500 MeV the binding potential corrections dominate so that elimination of Pauli blocking for the struck target nucleon will have minor effect on the density dependence computed in Ref. [8].

The density dependent interaction developed here was used to predict 500 MeV proton elastic scattering from ^{40}Ca as explained in Ref. [8]. The resulting predictions for the differential cross section, analyzing power, and spin rotation parameter [13] were negligibly different from similar calculations shown in Refs. [8,13] in which explicit Pauli blocking of both the projectile and the target nucleon and target nucleon binding potential corrections were assumed. The antisymmetrization effects considered here should be more important at lower energies.

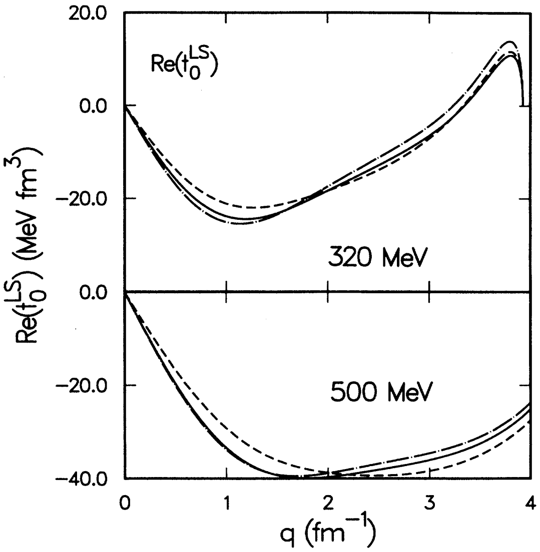

FIG. 3. Same as Fig. 1, except Re(t_0^{LS}).

V. CONCLUSIONS

The $(A+1)$-body antisymmetrized nucleon-nucleus optical potential formalism of Ref. [9] was used as a basis

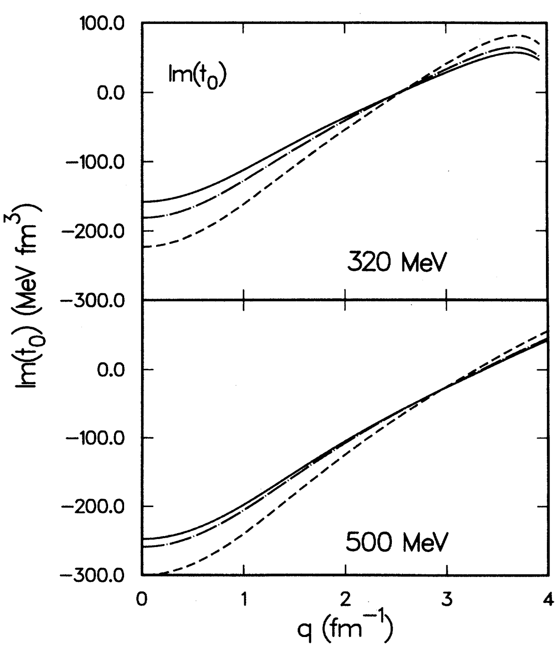

FIG. 2. Same as Fig. 1, except Im(t_0).

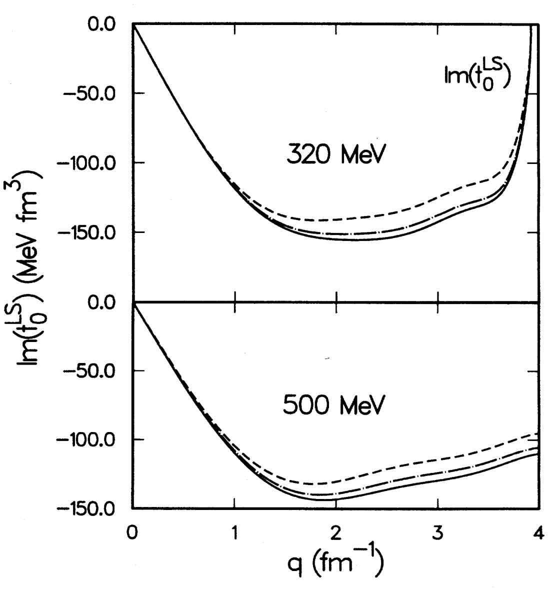

FIG. 4. Same as Fig. 1, except Im(t_0^{LS}).

for obtaining a practical estimate of the effects of projectile particle and target constituent identity on the standard, unsymmetrized optical potential and in predicted nucleon-nucleus elastic scattering observables. The antisymmetrized optical potential was expanded to second order in terms of the free, antisymmetrized NN t matrix. Struck-target-nucleon Pauli blocking corrections (implicit), target-nucleon binding potential corrections, and two-body target-nucleon correlation contributions were obtained as in the usual unsymmetrized theory [2,3]. In addition, the leading $(A+1)$-body antisymmetrization corrections were identified. One of these corresponds to a three-body exchange correction which was calculated previously [5] and shown to have negligible effects at medium energies. Another results in additional medium modifications to the standard, unsymmetrized optical potential. This correction can be combined with the first-order term, as in Eq. (22), to produce a new density dependent NN effective interaction. This latter many-body correction was estimated by computing NN scattering with a projectile nucleon Pauli blocking factor appropriate for infinite nuclear matter [8].

Density dependent effects were computed at 320 and 500 MeV and were shown to be similar to, but somewhat weaker than the density dependent effects calculated in Ref. [8] where both nucleons were explicitly Pauli blocked. It is recommended that density dependent effective interactions, which are to be used for nucleon-nucleus scattering calculations in the context of a nonrelativistic multiple-scattering optical potential formalism, be computed as described here where only the projectile nucleon is explicitly Pauli blocked in intermediate states and only the struck target nucleon is acted on by the nuclear binding potential in intermediate scattering states.

ACKNOWLEDGMENTS

Computing resources for this work were provided in part by The University of Texas System Center for High Performance Computing. This research was supported in part by the U.S. Department of Energy.

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