Temperature dependence of single-particle properties in nuclear matter

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Abstract

The single-nucleon potential in hot nuclear matter is investigated in the framework of the Brueckner theory by adopting the realistic Argonne $V_{18}$ or Nijmegen 93 two-body nucleon-nucleon interaction supplemented by a microscopic three-body force. The rearrangement contribution to the single-particle potential induced by the ground state correlations is calculated in terms of the hole-line expansion of the mass operator and provides a significant repulsive contribution in the low-momentum region around and below the Fermi surface. Increasing temperature leads to a reduction of the effect, while increasing density makes it become stronger. The three-body force suppresses somewhat the ground state correlations due to its strong short-range repulsion, increasing with density. Inclusion of the three-body force contribution results in a quite different temperature dependence of the single-particle potential at high enough densities as compared to that adopting the pure two-body force. The effects of three-body force and ground state correlations on the nucleon effective mass are also discussed.

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I. INTRODUCTION

The determination of the equation of state (EOS) of nuclear matter based on microscopic many-body approaches is of great interest in nuclear physics and nuclear astrophysics [1, 2, 3, 4]. During the dynamical evolution of heavy ion collisions (HIC) at intermediate and high energies, a transient state of hot and dense nuclear matter can be produced and therefore the experiments with HIC are powerful tools for constraining the nuclear EOS [4]. Since the EOS can not be measured directly in the experiments, one has to compare the experimental observables and the theoretical simulations by using transport models [5]. The single-particle (s.p.) potential felt by a nucleon in the nuclear medium is one of the basic ingredients of transport models for HIC and controls together with the nucleon-nucleon cross sections the collision dynamics.

Microscopically the s.p. potential in cold nuclear matter has been investigated by many authors based on microscopic theoretical approaches [6, 7, 8, 9, 10]. It has been pointed out [11] that the s.p. potential calculated at the lowest level of the Brueckner-Hartree-Fock (BHF) approximation cannot describe with sufficient accuracy the mean field, but it is necessary to include higher-order ground state correlations. In Refs. [6, 12], the effect of the ground-state correlations on the s.p. potential has been investigated within the Brueckner-Bethe-Goldstone (BGG) theory. Therein it is shown that inclusion of the rearrangement contribution in the s.p. potential is also crucial for restoring the Hugenholtz-Van Hove theorem, which is strongly violated at the lowest BHF level of approximation. The importance of the rearrangement term in the hole-line expansion of the mass operator has also been verified in Refs. [13, 14, 15] in connection with the optical model potential, the s.p. properties in the nuclear medium such as the nucleon effective mass and the nucleon mean free path, and the superfluidity properties in neutron matter and nuclear matter. Recently, the calculation of the high-order correlations in the mass operator has been extended to the case of isospin asymmetric nuclear matter at zero temperature [9] within the extended BHF approach.

Nowadays it is widely recognized that dressing in a non-relativistic framework the interaction with short-range correlations, either G-matrix [1] or in-medium T-matrix [16, 17, 18, 19], is not enough to reproduce the empirical saturation properties of cold nuclear matter. The saturation mechanism demands for the high-density repulsive contribution of the three-body force (TBF) [20, 21]. Microscopic TBF [21] have been used in Brueckner calculations [22, 23]. Phenomenological versions of TBF have also been used mainly in variational approaches [24]. In Ref. [25] we have extended the BHF approach with microscopic TBF to finite temperature and investigated the EOS of hot nuclear matter. It was shown that the TBF affect considerably the properties of hot nuclear matter especially at high densities and temperatures. The aim of the present paper is to extend our previous work by including the effect of the ground state correlations in the calculation of the s.p. potential.

The paper is arranged as follows. In the next section II we shall give a brief review of our theoretical model including the finite-temperature BHF approach with a microscopic TBF and the hole-line expansion for the mass operator at finite temperature [13, 26]. Our numerical results are presented and discussed in Sect. III, including the rearrangement term and its temperature dependence, the TBF effect on the ground state correlations and the s.p. potential. Finally, a summary will be given in Sect. IV along with a short comparison with other approaches.
II. THEORETICAL MODEL

A. Finite-temperature Brueckner-Hartree-Fock approach with a three-body force

The general formalism of the Brueckner-Bethe-Goldstone (BBG) theory for cold nuclear matter can be found in Refs. [8, 9] and its extension to the finite temperature case is given in Ref. [27]. Here we give a brief review for completeness. The starting point of the BBG scheme is the Brueckner reaction G-matrix, which satisfies the following Bethe-Goldstone (BG) equation,

\[ G[\rho, T, \omega] = V + V \sum_{k_1k_2} \frac{|k_1k_2\rangle Q(k_1, k_2, \rho, T)\langle k_1k_2|}{\omega - \varepsilon(k_1) - \varepsilon(k_2) + i0} G[\rho, T, \omega], \] (1)

where \( \omega \) is the starting energy and \( V \) denotes the realistic nucleon-nucleon (NN) interaction. The finite-temperature Pauli operator \( Q \) can be expressed as

\[ Q(k_1, k_2, \rho, T) = \left[ 1 - f(k_1) \right] \left[ 1 - f(k_2) \right], \] (2)

where \( f(k) \) is the Fermi distribution at finite temperature,

\[ f(k) = \left[ 1 + \exp \left( \frac{\varepsilon(k) - \mu}{T} \right) \right]^{-1}. \] (3)

In terms of the normalization condition,

\[ \rho = \sum_k f(k), \] (4)

one can determine the chemical potential \( \mu \) self-consistently by iteration for any given density \( \rho \) and temperature \( T \). The s.p. energy \( \varepsilon(k) \) in Eqs. (1) and (3) is defined as

\[ \varepsilon(k) \equiv \varepsilon(k, \rho, T) = \frac{k^2}{2m} + U(k, \rho, T), \] (5)

where the s.p. potential \( U(k, \rho, T) \) may be calculated from the real part of the on-shell anti-symmetrized G-matrix,

\[ U(k, \rho, T) = \sum_{k'} f(k') \text{Re} \langle kk' | G[\rho, T, \varepsilon(k) + \varepsilon(k')] | kk' \rangle_A. \] (6)

In the present calculations, the continuous choice [13] is adopted for the s.p. potential. On the one hand, in the zero temperature limit it provides a much faster convergence of the hole-line expansion than the gap choice [28]; on the other hand, it appears a natural choice for \( T \neq 0 \), since in the finite-temperature case any distinction between particles and holes becomes meaningless. Within the continuous choice the s.p. potential describes at the BHF level the mean field felt by a nucleon during its propagation between two successive scatterings in the nuclear medium [29].

The NN interaction \( V \) in the present calculation contains two parts: the Argonne \( V_{18} \) [30] or the Nijmegen 93 [31] two-body interaction plus the contribution of a microscopic
TBF. Two kinds of TBF have been adopted in the BHF formalism: One is the semi-
phenomenological Urbana TBF [32], which has two or few adjustable parameters determined
by fitting the empirical saturation density and energy of cold symmetric nuclear matter in
the BHF calculations [23, 33, 34]. The other one (used here) is a microscopic TBF based
on meson exchange coupled to intermediate virtual excitations of nucleon-antinucleon pairs
and nucleon resonances, which was originally proposed in [21]. It contains the contribution
of the two-meson exchange part of the NN interaction medium-modified by the intermediate
virtual excitation of nucleon resonances, the term associated to the non-linear meson-nucleon
coupling required by chiral symmetry, the simplest contribution arising from meson-meson
interactions, and finally the two-meson exchange diagram with the virtual excitations of
nucleon-antinucleon pairs. In this TBF model, the four important mesons $\pi, \rho, \sigma,$ and $\omega$
are considered [35]. The parameters of the TBF, i.e., the coupling constants and the form
factors, have been redetermined recently in Ref. [22] from the one-boson-exchange potential
(OBEP) model to meet the self-consistent requirement with the adopted two-body force.
A more detailed description of the TBF model and the approximations can be found in
Ref. [21].

In the zero-temperature case, the TBF contribution $W_3$ has been included in the BHF
calculations by constructing an effective two-body interaction $V_3$ via a suitable average with
respect to the third-nucleon degrees of freedom [20, 21]. By extending this scheme to finite
temperature, one can reduce the TBF to a temperature-dependent effective two-body force
$V_3(T)$, which reads in $r$-space [25]

$$
\langle r'_1 r'_2 | V_3(T) | r_1 r_2 \rangle = \frac{1}{4} \text{Tr} \sum_{k_n} f(k_n) \int d r_3 d r'_3 \phi_n^*(r'_3) [1 - \eta(r'_{13}, T)] [1 - \eta(r'_{23}, T)]
\times W_3(r'_1 r'_2 r'_3 | r_1 r_2 r_3) \phi_n(r_3)[1 - \eta(r_{13}, T)][1 - \eta(r_{23}, T)] ,
$$

(7)

where the trace is taken with respect to the spin and isospin of the third nucleon. The defect
function $\eta(r, T)$ [21, 21] is defined as $\eta(r, T) = \phi(r) - \psi(r, T)$, where $\psi(r, T)$ is the correlated
wave function for the relative motion of two nucleons in the nuclear medium and $\phi(r)$ is
the corresponding unperturbed one. A detailed description and justification of the above
scheme can be found in Ref. [21]. As has been pointed out in Refs. [25, 36], the TBF $W_3$
itself is the same as the one adopted in our previous calculations for the zero-temperature
case [22] and is independent of temperature. However, in the finite-temperature case, the
effective two-body force $V_3(T)$ depends on temperature due to the medium effects caused by
the Fermi distribution $f(k)$ and the defect function $\eta(r, T)$, which is strongly temperature
dependent.

In the BHF approximation with the TBF, Eqs. (1) and (5-7) are solved self-consistently.
First, the temperature-dependent G-matrix is calculated along with the auxiliary potential
by solving the BG equation, then one evaluates the defect function with only the two-body
force, further constructs the effective two-body force, and finally adds the effective two-body
force to the bare two-body force. This procedure is repeated until convergence is reached.
Obviously, the effect of microscopic TBF is automatically included into the G-matrix by
iteration.
B. Hole-line expansion of the mass operator

Within the BBG theory the mass operator can be expanded into a perturbation series in terms of the number of hole lines \([13, 37, 38]\):

\[
M(k, \omega) = M_1(k, \omega) + M_2(k, \omega) + M_3(k, \omega) + \ldots ,
\]

which is schematically illustrated in Fig. 1. In this expansion the two-hole line contribution \(M_2(k, \omega)\) is the so-called rearrangement term, representing the Pauli rearrangement, i.e., the medium dependence of the effective G-matrix interaction via the ground state Pauli blocking effect. The third-order term \(M_3(k, \omega)\) accounts for the fact that the hole state is partially empty in the correlated ground state of nuclear matter \([7, 13]\).

The mass operator \(M(k, \omega) = V(k, \omega) + iW(k, \omega)\) is a complex quantity. When evaluated on the energy shell, its real part \(V(k) \equiv V(k, \varepsilon(k))\) describes the s.p. potential felt by a nucleon in the nuclear medium and can be compared with the empirical potential depth extracted from the optical potential model \([33]\), whereas its imaginary part \(W(k) \equiv W(k, \varepsilon(k))\) is related to the nucleon mean free path \([9, 13, 14]\). The on-shell condition is given by the following energy-momentum relation \([9]\),

\[
\frac{k^2}{2m} + \Re M(k, \omega) = \omega .
\]

To the lowest order of approximation, the on-shell condition is simplified as \(k^2/2m + \Re M(k, \varepsilon(k)) = \varepsilon(k)\), where \(\varepsilon(k)\) is the BHF s.p. energy. In the present work, we shall consider only the on-shell properties of the mass operator. Hereafter we denote the real and imaginary parts of the on-shell mass operator as \(V(k)\) and \(W(k)\), respectively. Their hole-line expansions can be written as

\[
V(k) = V_1(k) + V_2(k) + V_3(k) + \ldots ,
\]

\[
W(k) = W_1(k) + W_2(k) + W_3(k) + \ldots .
\]

Since the renormalization contribution \(M_3(k)\) is quite small compared to the lowest-order BHF term \(M_1(k)\) and the rearrangement term \(M_2(k)\) \([9, 38]\), in the present paper we concentrate on the investigation of the rearrangement contribution \(M_2(k)\), which in the finite-temperature case can be expressed as follows,

\[
M_2(k_1, \omega) = \frac{1}{2} \sum_{k_2k_1k_2} \left[ 1 - f(k_2) \right] f(k_1') f(k_2') \frac{|\langle k_1k_2|G|k_1'k_2'\rangle|_A|^2}{\omega + \varepsilon(k_1) - \varepsilon(k_1') - \varepsilon(k_2') + i0} .
\]

By using the usual angular averaging procedure in order to remove the angular dependence of the energy denominator and the anti-Pauli operator \([1, 40, 41]\), \(M_2(k, \omega)\) can be readily calculated in terms of the G-matrix expanded in the partial wave representation \([9, 38]\):

\[
M_2(k, \omega) = \frac{2}{\pi^2k} \sum_{JSTLL'} (2J + 1) \int dq dP dP \left[ 1 - f(\sqrt{P^2/2 + 2q^2 - k^2}) \right] \times \int q^2 dq' R(q', P) \frac{|G_{JST}'[q, q', P, \varepsilon_12(q', P)]|^2}{\omega + \varepsilon(\sqrt{P^2/2 + 2q^2 - k^2}) - \varepsilon_12(q', P) + i0} ,
\]
where \( q' = (k_1 - k_2)/2 \) and \( P = k_1 + k_2 \) are the relative and total momenta, respectively. In the above equation, \( \varepsilon_{12}(q', P) \equiv \langle \varepsilon(k_1) + \varepsilon(k_2) \rangle \) denotes the angular average of the energy denominator and the angular-averaged anti-Pauli operator is defined as
\[
R(q', P) \equiv \frac{1}{2} \int_0^{\pi} \sin \theta d\theta f(k_1)f(k_2),
\]
where \( \theta \) is the angle between \( q' \) and \( P \).

Thus in terms of the G-matrix obtained from Eq. (1), the first- and second-order contributions in the hole-line expansion of the mass operator can be calculated. Hereafter we shall denote the on-shell mass operator containing the contribution of the rearrangement term as \( M_{12}(k) = M_1(k) + M_2(k) \), and its real and imaginary parts as \( V_{12}(k) = V_1(k) + V_2(k) \) and \( W_{12}(k) = W_1(k) + W_2(k) \), respectively.

### III. RESULTS AND DISCUSSION

The self-consistent BHF procedure extended to TBF has been applied to study symmetric nuclear matter at finite temperature. Two realistic interactions, the Argonne \( V_{18} \) and Nijmegen 93 potentials, have been used to describe the two-body force, whereas the model of TBF is the meson exchange interaction discussed in Sec. II A. The results for the EOS have been presented elsewhere [22]. Here we only report the saturation properties: The main effect of TBF is to reduce the saturation density from 0.26 fm\(^{-3} \) to 0.19 fm\(^{-3} \) due to the high-density extra repulsion. The energy per particle rises from \(-18 \) MeV to \(-15 \) MeV, whereas the compression modulus is lowered from 230 MeV to 210 MeV.

#### A. Rearrangement contribution

In Fig. 2 are plotted the real and imaginary parts of the rearrangement contribution \( M_2(k) \). It is seen that they depend sensitively on temperature in both cases with and without including the TBF contribution. The real part of \( M_2(k) \), i.e., the rearrangement term \( V_2(k) \) of the s.p. potential is repulsive and its contribution is mainly concentrated in the region below the Fermi momentum \( k_F \), where the ground state hole-hole correlations are expected to be most effective. Around \( k_F \) the magnitude of \( V_2(k) \) decreases rapidly as a function of momentum and vanishes at high enough momentum. As the nuclear matter is heated up, \( V_2(k) \) is strongly reduced due to the softening of the Pauli blocking around the Fermi surface at high temperature, which weakens the effect of the ground state hole-hole correlations.

Also the TBF contribution suppresses considerably these correlations in cold and hot nuclear matter due to its strong short-range repulsion, and leads to a sizable reduction of the rearrangement correction \( V_2(k) \). This feature is much more pronounced in the low-momentum region, where the ground state correlations are more significant. The influence of the TBF diminishes as the temperature increases, since the ground state correlations themselves become smaller at higher temperature. For instance, in the case of the \( V_{18} \) potential, at \( T = 0 \), the reduction of \( V_2(k = 0) \) due to the TBF is about 6 MeV, from 30 MeV to 24 MeV, while at \( T = 20 \) MeV, the reduction is about 3 MeV, from 13 to 10 MeV. Comparing the curves in the upper panels with the corresponding ones in the lower panels, it is readily seen that in both cases with the \( V_{18} \) potential and with the Nijmegen 93 potential, the influence of the TBF decreases as the temperature rises, implying that the temperature dependence of the TBF effect on the correlation potential \( V_2(k) \) is not much sensitive to the adopted two-body realistic NN interactions. However, the effect of the TBF
is more pronounced in the case with the Nijmegen 93 potential in the temperature range considered here.

The imaginary part $W_2(k)$ of $M_2(k)$ is related directly to the lifetime or the width of a hole state [13], which vanishes above the Fermi momentum $k_F$ for cold nuclear matter due to the Pauli blocking effect [9]. However, at finite temperature, the tail of $W_2$ slightly extends to the momentum region above the Fermi surface, since the Fermi surface becomes diffusive and the Pauli blocking is weakened in this case. The temperature dependence of $W_2(k)$ is somewhat more complicated than that of the real part $V_2(k)$. With increasing temperature, $W_2(k)$ gets smaller in the momentum region well below $k_F$, while it becomes larger in the upper part of the Fermi sea. This can be explained as follows. On the one hand, the ground state hole-hole correlations decrease with temperature. On the other hand, at a higher temperature, a hole state intends to decay faster, especially close to the Fermi surface. The competition between these two effects determines the final variation of $W_2(k)$ as a function of temperature. The TBF suppresses somewhat the ground state hole-hole correlations and its contribution reduces the magnitude of the imaginary part $W_2(k)$ in agreement with the results obtained for the real part of $M_2(k)$.

In order to discuss the density dependence, we report in Fig. 3 the rearrangement contribution $M_2$ at the fixed momentum $k = 0$ as a function of density for the two cases with and without the TBF. It is seen that the real and the imaginary parts of $M_2(k)$ increase monotonically with density in both cases, indicating that the effect of the ground state hole-hole correlations is stronger in denser nuclear matter, where the number of hole state is larger. The TBF leads to a reduction of the ground state correlations in the whole density region considered. Its effect is fairly small at low densities and becomes stronger rapidly as the density increases.

**B. Mass operator**

The complex mass operator $M_{12}(k)$ including the BHF contribution and the rearrangement term is reported in Fig. 4 for several different values of density and two temperatures $T = 0, 20$ MeV. Since the real part of the rearrangement term is repulsive, its contribution reduces to a large extent the attraction of the pure BHF s.p. potential $V_1(k)$ [25]. In the zero-temperature case at normal nuclear matter density $\rho = 0.17$ fm$^{-3}$, the repulsive contribution of the ground state correlations causes the depth of the s.p. potential $V_{12}(k = 0)$ to rise to $\approx -60$ MeV from its BHF value $V_1(k = 0) \approx -85$ MeV [9]. This improves considerably the agreement of our predicted s.p. potential with the optical model potential extracted from nucleon-nucleus scattering experiments [13, 39]. As the nuclear density increases, the total potential $V_{12}(k)$ becomes more attractive in the low-momentum region, while at high enough momenta it becomes slightly less attractive. Such a density dependence is mainly attributed to the density behavior of the pure BHF s.p. potential $V_1(k)$, as in the zero-temperature case [2].

Concerning the temperature dependence, it is seen from Fig. 4 that in the case without TBF, the total potential becomes more attractive at low momenta and less attractive in the high momentum region above $k_F$ when the nuclear matter is heated up. This is quite different from the temperature behavior of the lowest-order BHF s.p. potential, where an increase in temperature results in an overall reduction of the attraction of the BHF s.p. potential in the whole momentum region [25]. At low momenta, the enhancement of the attraction of $V_{12}(k)$ with temperature is readily understood in terms of the temperature effect on the ground
state correlations: The correlation potential $V_2(k)$ applies mainly to states below the Fermi momentum and its repulsive contribution is reduced largely by increasing temperature as shown in Fig. 2. At high momenta, the contribution of $V_2(k)$ is very small and consequently the variation of the total potential $V_{12}(k)$ with temperature is essentially determined by that of the BHF potential $V_1(k)$.

At low densities (for example, $\rho = 0.085 \text{ fm}^{-3}$), the TBF effect is fairly weak. However, at high densities, the TBF modification of the s.p. potential becomes significant, especially for high temperatures. In the two cases of $\rho = 0.225$ and 0.34 fm$^{-3}$, inclusion of the TBF contribution even makes the temperature behavior of $V_{12}(k)$ quite different from that without TBF. It can be seen from the figure that the calculated $V_{12}$ with the TBF contribution gets more repulsive in the whole momentum region. This may be understood as a consequence of the competition between the following two effects. Below the Fermi momentum, the TBF suppresses the ground state hole-hole correlations due to its strong short-range repulsion and thus it reduces the repulsive contribution of the rearrangement term $V_2(k)$. On the other hand, the TBF contribution to the BHF s.p. potential $V_1(k)$ is repulsive and reduces the attraction of $V_1(k)$. At high densities, as the temperature increases, the TBF effect on the ground state correlations becomes weaker while its contribution to the pure BHF s.p. potential $V_1(k)$ gets larger. Above the Fermi momentum, the correlation term $V_2(k)$ tends to vanish and the temperature dependence of the total potential $V_{12}$ is dominated by the BHF one $V_1(k)$. Accordingly at high densities the total potential gets more repulsive in the whole momentum region as the matter is heated. Another feature related to the temperature dependence that can be observed from Fig. 4 is that the curvature around the Fermi momentum becomes more smooth with increasing temperature in both cases with and without TBF, which is in agreement with the prediction obtained for the pure BHF s.p. potential [20, 25] and is attributed to the thermal excitations around the Fermi surface at finite temperature.

Now let us turn to the imaginary part $W_{12}$, which is also called absorptive potential [13]. It is seen from Fig. 4 that for cold nuclear matter the absorptive potential crosses zero at the Fermi momentum $k_F$. The BHF contribution $W_1(k)$ vanishes below $k_F$ and the correlation term $W_2(k)$ above $k_F$ due to the Pauli blocking. At finite temperature, $W_1(k)$ may extend to the momentum region below $k_F$ and $W_2(k)$ to above $k_F$. The total absorptive potential crosses zero at a momentum close to $k_F$, since the chemical potential does not change very much up to $T = 20$ MeV [38]. For each fixed density considered here, the absolute value of $W_{12}(k)$ increases above the Fermi momentum $k_F$ while it decreases below $k_F$. The temperature dependence of $W_{12}(k)$ turns out to be stronger in nuclear matter with a smaller density. As expected, the TBF effect on the imaginary part $W_{12}(k)$ increases with density, while it is relatively weak compared to that on the real part $V_{12}(k)$, especially at high densities.

It is instructive to make a comparison with the predictions of the in-medium T-matrix approach, where the self-energy is self-consistently calculated with the T-matrix [16, 17, 18, 19]. The main difference from the G-matrix is that the T-matrix embodies both particle-particle and hole-hole correlations. Since its convergence is not protected by any hole line expansion, one needs a very accurate determination of the self-energy. Also, zero-temperature calculations are usually not possible in this approach without removing in some way the too strong pairing instabilities. Doing so, however, nuclear matter appears to be underbound [17, 18, 19]. Concerning the s.p. properties, one observes an overall qualitative agreement between the T-matrix self-energy and the Brueckner one (see for instance Fig. 5 of Ref. 17).
or Fig. 13 of Ref. [18]), which is an indication that the rearrangement term contains most of the hole-hole correlations.

C. Effective mass

The effective mass describes the nonlocality of the s.p. potential and makes its local part less attractive. It is of great interest [42], since it is closely related with many nuclear phenomena such as the dynamics of heavy ion collisions at intermediate and high energies [43], the damping of nuclear excitations and giant resonances [44], and the adiabatic temperature of collapsing stellar matter [45]. The effective mass \( m^*(k) \) is defined as [13]

\[
m^*(k) = m \left[ \frac{dE(k)}{dk} \right]^{-1},
\]

(14)

where \( E(k) \) is the s.p. energy determined by the momentum-energy relation Eq. (9). When the mass operator is expanded up to the second order, one can readily calculate the effective mass as follows:

\[
m^*(k) = \left[ 1 + \frac{m}{k} \frac{dV_{12}(k)}{dk} \right]^{-1}.
\]

(15)

The calculated \( m^*(k)/m \) versus momentum is reported in Fig. 5, where the upper panel shows the results for nuclear matter at the empirical saturation density \( \rho = 0.17 \text{ fm}^{-3} \) and for two temperatures \( T = 0, 20 \text{ MeV} \) with and without including TBF, while the lower panel displays the results for three values of density \( \rho = 0.085, 0.17, \text{ and } 0.34 \text{ fm}^{-3} \) at zero temperature. Due to the high possibility for particle-hole excitations near the Fermi surface [13], the momentum dependence of \( m^*(k)/m \) in cold nuclear matter is characterized by a wide bump around the Fermi momentum \( k_F \). Recently such a structure has also been found [46] within the relativistic Dirac-Brueckner-Hartree-Fock approach for the momentum dependence of the non-relativistic type of effective mass introduced in terms of the Schrödinger equivalent s.p. potential [39]. Inclusion of the contribution of the rearrangement term makes the peak of the effective mass more pronounced as compared to the results obtained at the lowest BHF level of approximation [25], consistent with previous investigations adopting pure two-body NN interactions [6, 9, 13, 38]. The value \( m^*_F \equiv m^*(k_F) \) of the effective mass obtained in the present calculation is around 1.02 when the TBF contribution is included and 1.08 in the case without the TBF. Both values are larger than the BHF value \( \approx 0.8 \) [9], which is mainly attributed to the contribution of the ground state correlations, i.e., the arrangement term \( V_2(k) \). Inclusion of even higher-order terms, i.e., the third- and the fourth-order terms in the hole-line expansion of the mass operator may reduce \( m^*_F/m \) to about 0.9 as discussed in Ref. [9].

As the nuclear matter is heated, the peak of \( m^*(k) \) becomes flatter and the peak value lower, which is related directly to the temperature effect on the s.p. potential \( V_{12} \) around the Fermi momentum and is similar to the result obtained at the BHF level of approximation [25]. It is seen that in the case of \( T = 20 \text{ MeV} \) the peak structure almost disappears. The TBF effect on the effective mass is significant only at low temperatures and in the low-momentum region below \( k_F \), because the TBF effect on the BHF s.p. potential \( V_1(k) \) is an overall enhancement in the whole momentum region [25] and thus the momentum-dependence of the total potential \( V_{12}(k) \) is mainly affected by the correlation potential \( V_2(k) \), which is important below \( k_F \). As the temperature increases, the correlation potential
$V_2(k)$ itself and the influence of the TBF on it become weaker as has been shown in Fig. 2, and as a result also the TBF effect on the effective mass is smaller at higher temperature. From the lower part of Fig. 5 it is clear that with increasing density the peak structure of the effective mass around $k_F$ becomes less pronounced and the value of the effective mass at $k_F$ decreases, which is comparable with the previous calculations adopting pure two-body nucleon-nucleon forces [6, 14].

IV. SUMMARY

In the present work, we have reported the investigation of the s.p. properties for cold as well as hot nuclear matter within the framework of the Brueckner theory by adopting the Argonne $V_{18}$ or the Nijmegen 93 two-body nucleon-nucleon interaction plus a microscopic three-body force based on the meson-exchange model. The mass operator has been calculated up to the second order of the hole-line expansion. Special attention has been paid to the effect of ground state correlations on the s.p. potential in hot nuclear matter and the influence of the TBF on the ground state correlations. Our result shows that these correlations give a repulsive contribution to the s.p. potential mainly within the Fermi sphere in agreement with previous investigations [9, 13]. As the temperature rises, the real and the imaginary part of the rearrangement contribution become less repulsive due to the weakening of the correlations.

The TBF contribution turns out to reduce the two-hole line correlation term in magnitude due to its strong short-range repulsion. When the nuclear matter is heated, the TBF effect on the ground state correlations becomes weaker. At $T = 0$, the imaginary part of the second-order mass operator vanishes below the Fermi surface due to the Pauli blocking, while in the finite-temperature case its tail may extend slightly above the Fermi momentum, since the Fermi surface becomes diffusive and the Pauli blocking is weakened. Both the real and the imaginary parts of the rearrangement term are shown to be increasing functions of density, indicating that the effect of the ground state correlations is stronger at higher density.

The consideration of the two-hole line diagram of the mass operator reduces remarkably the attraction of the lowest-order BHF s.p. potential and improves significantly the agreement of our microscopic s.p. potential with the empirical optical model potential [13]. For low density up to about $\rho_0 = 0.17$ fm$^{-3}$, the total s.p. potential $V_2(k)$ is shown to become more attractive in the low-momentum region when the nuclear matter is heated up, which can be attributed to the reduction of the repulsive contribution from the ground state correlations in the finite-temperature case. The role played by the TBF turns out to become stronger as the density increases. The TBF affects the total s.p. potential in two ways. On the one hand it provides a repulsion to the BHF part and makes the BHF s.p. potential more repulsive. On the other hand it suppresses somewhat the ground state correlations and reduces the repulsive contribution of the rearrangement term. As a combined result, at high densities the s.p. potential becomes less attractive in the whole momentum region when the TBF is taken into account. At high enough densities, the TBF contribution even changes the temperature dependence of the s.p. potential, i.e., it makes the s.p. potential more repulsive at high temperature than at low temperature.

In both cases with and without the TBF, the momentum dependence of the effective mass displays a broad peak around the Fermi momentum in the zero-temperature case, which is similar to the result obtained in the lowest BHF approximation. The effects of the TBF and the ground state correlations on the effective mass are found to be important mainly
below the Fermi surface. The ground state correlations make the peak more pronounced as compared to the BHF one, whereas the TBF reduces the peak via its effect on the rearrangement term. Increasing temperature smoothes the s.p. potential around the Fermi surface and consequently leads to a flattening of the peak, making it almost vanish at high enough temperature. The value of the effective mass at the Fermi surface decreases with density, as in the BHF approximation.

In summary, it is shown by our investigation that both the ground state correlations and the TBF affect considerably the s.p. potential in the nuclear medium and its temperature dependence. It is therefore important to take into account their effects in the application to transport-model simulations of HIC.

This investigation reinforces the role played by TBF in nuclear matter. The G-matrix, as two-body effective interaction, which embodies the short-range particle-particle correlations, while preventing hard-core divergences, is not able to reproduce the empirical saturation density at the convergence of the hole line expansion [28]. Other non-relativistic approaches, the variational method [24] as well as in-medium T-matrix theory seem to confirm the latter conclusion. On the other hand, it has been proven that the Dirac-Brueckner theory includes TBF (Z-diagrams) as the effect of coupling to the negative energy states.

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Figures

FIG. 1: Hole-line expansion of the mass operator.

FIG. 2: The real part $V_2(k)$ and the imaginary part $W_2(k)$ of the rearrangement term $M_2$ versus nucleon momentum $k$ at normal nuclear density $\rho_0 = 0.17$ fm$^{-3}$. The results in the upper and lower panels are obtained by adopting the $V_{18}$ and Nijmegen 93 two-body nucleon-nucleon interactions, respectively. The solid (dashed) curves represent the results with (without) the TBF contribution for temperatures $T = 0, 10, 20$ MeV from top to bottom.

FIG. 3: The density dependence of the real and the imaginary parts of the rearrangement term at $k = 0$, $M_2(k = 0)$, with (solid lines) and without (dashed lines) the TBF contribution for three different temperatures $T = 0, 10, 20$ MeV from the top to bottom.

FIG. 4: The real and imaginary parts of the total mass operator up to the second-order term for four different nuclear densities as a function of momentum $k$. In each panel the rising curves correspond to the real part $V_{12}(k)$ and the dropping curves to the imaginary part $W_{12}(k)$. For both quantities, the solid and dashed (dotted and dot-dashed) lines stand for the results at $T = 0, 20$ MeV with (without) the TBF contribution, as indicated in the first panel.

FIG. 5: The effective mass $m^*/m$ including the BHF and the second-order correlation contributions as a function of momentum $k$. The upper panel shows the results at normal nuclear matter density for two temperatures $T = 0, 20$ MeV with and without the TBF. The lower panel depicts the effective mass including the TBF contribution at $T = 0$ for three densities $\rho = 0.085, 0.17, \text{ and } 0.34$ fm$^{-3}$. 
\[ M(k) = \quad + \quad + \quad + \cdots \]
AV18

$\rho = 0.085 \text{ fm}^{-3}$

$\rho = 0.17 \text{ fm}^{-3}$

$\rho = 0.255 \text{ fm}^{-3}$

$\rho = 0.34 \text{ fm}^{-3}$

$T=0$ with TBF

$T=0$ without TBF

$T=20\text{MeV}$ with TBF

$T=20\text{MeV}$ without TBF

$M_{12} (\text{MeV})$

$k (\text{fm}^{-1})$
\[ \rho = 0.085 \text{ fm}^{-3} \]

\[ \rho = 0.17 \text{ fm}^{-3} \]

\[ \rho = 0.34 \text{ fm}^{-3} \]

\[ m*/m \]

\[ k (\text{fm}^{-1}) \]

- \[ T = 0 \]
- \[ T = 20 \text{ MeV} \]

- \( \rho = 0.085 \text{ fm}^{-3} \) without TBF
- \( \rho = 0.17 \text{ fm}^{-3} \) without TBF
- \( \rho = 0.34 \text{ fm}^{-3} \) without TBF

- \( \rho = 0.085 \text{ fm}^{-3} \) with TBF
- \( \rho = 0.17 \text{ fm}^{-3} \) with TBF
- \( \rho = 0.34 \text{ fm}^{-3} \) with TBF