Three-quark clusters at finite temperatures and densities

M. Beyer*, S. Mattiello
Fachbereich Physik, Universität Rostock, 18051 Rostock, Germany

T. Frederico
Dep. de Física, Instituto Tecnológico de Aeronáutica, Centro Técnico Aeroespacial, 12.228-900 São José dos Campos, São Paulo, Brazil.

H.J. Weber
Dept. of Physics, University of Virginia, Charlottesville, VA 22904, U.S.A.

(March 27, 2022)

We present a relativistic three-body equation to study correlations in a medium of finite temperatures and densities. This equation is derived within a systematic Dyson equation approach and includes the dominant medium effects due to Pauli blocking and self energy corrections. Relativity is implemented utilizing the light front form. The equation is solved for a zero-range force for parameters close to the confinement-deconfinement transition of QCD. We present correlations between two- and three-particle binding energies and calculate the three-body Mott transition.

PACS: 12.39.Ki, 21.65.+f, 21.45.+v

Keywords: correlations, three-body equations, light front, quark matter, Mott transition, relativistic quark models, Dyson equations

Exploring the phase structure of quantum chromodynamics (QCD) for the whole density-temperature plane is a challenging task for the standard model. Lattice Monte Carlo simulations have revealed several exciting results over the past decade [1,2]. In addition, at finite densities, modeling of QCD has added substantially to our understanding of its rich phase structure.

One particularly interesting result is the possibility of “color-superconductivity”, extensively discussed recently [3,4]. The possible appearance of color-superconductivity is related to similar ideas that led to nucleon or Cooper pairing. Another important aspect is the transition from nucleons to quarks as relevant degrees of freedom (Mott transition) that is likely to be related to the “confinement-deconfinement” phase transition. In these cases the influence of three-quark correlations has hardly been investigated. Recently, Pepin et al. addressed the question of a possible competition of three-quark clusters and two-quark condensation at finite density but zero temperature.

The problem of correlations in medium is tackled within a Green functions formalism [5] using the Dyson equation approach [6]. Three-particle correlations on the basis of a new in-medium Alt-Grassberger-Sandhas (AGS) type equation have been investigated in [7,8] and implemented into a Boltzmann-Uehling-Uhlenbeck simulation of heavy ion collisions and compared to experiments in [9,10]. Recently, the formalism has been extended to four-particle correlations and solved to study in-medium properties of the α-particle [11]. In contrast to the nucleonic phase, however, investigations of the quark phase of QCD entail several new issues such as relativity, color degrees of freedom, confinement, effective residual interactions, number of flavors, among others.

In order to derive effective few-body equations for quarks at finite temperatures and densities we have to consider in-medium effects as well as relativity. In medium the mass of the constituents may vary and in fact becomes close to zero for higher temperatures and densities, see e.g. [12]. To implement relativity we use the light front approach [13]. In the context of the quantum statistical framework it has some compelling advantages, although the concept of quasi particles in a background field that will be utilized further down and extended to include correlations introduces a special frame of reference. These are: i) Several boosts are kinematical (interaction free). As a consequence of the transitivity of kinematical boosts the Fock state decomposition is stable [14]. Since we decouple the hierarchy of Green function equations using a cluster mean field expansion [15] this decoupling of Fock spaces is therefore retained. ii) The dependence of the equations for the isolated system on the c.m. momentum of the cluster can easily be separated, see e.g. [16,17]. In a homogeneous medium this is also an important feature of the light front form, because inclusion of Pauli blocking factors that depend on the c.m. momentum of the cluster leads to a parametric dependence on the c.m. momentum only. iii) Pair creation processes are likely to be suppressed on the light front [18,19] in some frames of reference. Therefore we may presently consider particles only and no antiparticles which leads to technical simplifications that, however, can be relaxed as the investigation goes on. iv) Another advantage is that the light front form can be formulated in a Hamiltonian language and therefore all the results of the previously tested quantum statistical framework.
used here can be recovered using a proper interpretation.

On the other hand the vacuum structure in the light front dynamics involves technical difficulties in the presence of condensates, viz. zero modes in general, see e.g. [28]. These difficulties include a proper description of zero modes and Goldstone modes which are being investigated by several groups, see e.g. [29]. Further, rotational invariance on the light front implies interaction dependent generators. As a consequence angular decomposition is technically more involved. Presently, we make use of angular averaging. This introduces uncertainties, but for homogeneous infinite matter angular averaging has proven quite useful in the past. Finally, the description of spin degrees of freedom on the light-front is an issue that involves real complications for the three-body system. For an extensive discussion of the problem see Ref. [30]. Presently, we treat the spin in a simplified manner; namely by averaging over elementary spin projections. This approximation leads to an effective one-channel boson-type three-body equation. This simplifies the calculation significantly for the sake of a transparent treatment of in-medium effects of the relativistic equations. A boson-type relativistic three-body equation with a zero range interaction has previously proven useful for studies of isolated nucleon form factors [31,32]. The technically involved full inclusion of the spin degrees of freedom along the lines of Ref. [30] will be postponed to a later stage of the calculations. As a consequence we are presently not able to distinguish between so-called scalar, pseudoscalar, vector, and axial vector quark pairs. As many aspects can be improved, we emphasize that we present and solve for the first time an in-medium relativistic three-body equation for both finite temperatures and finite densities.

The light front has been utilized earlier for infinite nuclear matter calculations at zero temperatures and normal nuclear matter densities in [23]. That approach is based on the Brückner G matrix and two-body correlations are treated with an in-medium Blankenbecler-Sugar equation including a medium-modified interaction. For a general discussion of differences between the Green function approach (for finite temperatures and densities) used here and the Brückner approach (at zero temperatures) see e.g. Ref. [3].

The equations derived here are based on a systematic quantum statistical framework using a cluster expansion for the Green functions. We start with the chronological Green function [3]

\[ i \frac{\partial}{\partial \tau} G^{\tau-\tau'}_{\alpha\beta} = \delta(\tau - \tau')N_{\alpha\beta} + \sum_\gamma \int d\tau_1 M^{\tau-\tau_1}_{\alpha\gamma} G^{\tau_1-\tau'}_{\gamma\beta} \]  

where

\[ N_{\alpha\beta} = \langle [A_\alpha(\tau), A^\dagger_{\beta}(\tau)] \rangle \]  

Here, we neglect retardation in the mass operator \( M \), that would introduce more intermediate Fock space components, viz.

\[ M^{\tau-\tau'}_{\alpha\beta} \rightarrow \delta(\tau - \tau')M^{\tau}_{\alpha\beta} \]

\[ M^{0}_{0,\alpha\beta} = \sum_\gamma \langle [A_\alpha, H](\tau), A^\dagger_\gamma(\tau) \rangle \cdot N^{-1}_{\gamma\beta} \]

So far the formalism presented is rather general and for a typical many-body Hamiltonian (with generic two-body interactions \( V_2 \)) has been proven useful for various domains of many-body physics (see [4]) including the calculation of correlations in nuclear systems at finite densities and temperatures [1,19].

At chemical potentials \( \mu \) and temperatures \( T = 1/(kB\beta) \) averaging is due to a grand canonical ensemble in equilibrium, \( \langle O \rangle = \text{Tr}(\rho_0O) \), where \( \rho_0 \) denotes the corresponding statistical operator expressed in terms of light front operators. In averaging we consider the few-body cluster embedded in a homogeneous mean field of uncorrelated particles which is a reasonable approximation. Hence, together with \( \langle \ \rangle \) the hierarchy of Green functions decouples. This leads, upon introducing particle \( b(k) \) and antiparticle \( \bar{d}(k) \) Fock operators, to the standard Fermi-Dirac single particle distribution functions. For particles this reads

\[ f(k^+, \vec{k}_T^2) = \langle b^\dagger(k)b \rangle = \left( \exp \left[ \beta \left( \frac{1}{2}(k^+ + k_\perp^2) - \mu \right) \right] + 1 \right)^{-1} \]

expressed in terms of light front form momenta given by \( \vec{k}_\perp = (k_x, k_y) \) and \( k^+ = k_0 \pm \vec{k}_\perp \). The on-shell light-front energy is given by \( k_\perp^2 = (k_\perp^2 + m(\mu, T)^2)/k^+ \) with a medium dependent mass \( m(\mu, T) \). The components of the four vector \( k_\perp \) are \( k_\perp = (k_\perp^0, \vec{k}_\perp) \). The thermodynamic Green function for the one-particle case, by use of eqs. (1), (3), and (5) is then given by
\[ G(k) = \frac{\theta(k^+)}{2k^+} (\gamma k_0 + m(\mu, T)) \]
\[ \times \left( \frac{1 - f}{\frac{1}{2}k^+ - \frac{1}{2}k_0^+ + i\varepsilon} + \frac{f}{\frac{1}{2}k^+ - \frac{1}{2}k_0^+ - i\varepsilon} \right). \] (7)

Presently we consider only particle degrees of freedom on the light front and have therefore dropped the term related to \( \theta(-k^+) \). The energy variable \( k^- \) plays the role of the Matsubara frequency \[ \Im \frac{1}{k^+ + k^-} = k^0 \rightarrow \pi \lambda / (-i\beta) + \mu. \] Eventually, averaging over the elementary spin projections leads to a Bose type Green function.

For particles that are part of a larger \( n \)-body cluster it is convenient to introduce fractions \( x = k^+ / P_n^- \), where \( P_n^+ \) is the plus component of the cluster’s c.m. momentum. For a cluster at rest \( P_n^+ = M_n \), where \( M_n \) is the mass of the cluster.

![Fig. 1. Feynman-Galitzkii equation for the two-body \( t \)-matrix with zero range interaction. The crosses refer to the Pauli-blocking factor \( N_2 \). Lines represent quasi-particles.](image1)

Evaluation of Eq. (3) for two or three particles with the mass operator of eq. (3) in a homogeneous medium of independent particles leads to equations of the Green function where the single time \( \tau - \tau' \) is Fourier transformed to a Matsubara frequency. These equations include mass corrections and Pauli blocking factors in a systematic way. They can be rearranged into equations for \( t \)-matrices. The resulting equation for the two-body \( t \)-matrix \( T_2 \) has the same formal structure as the Feynman-Galitzkii equation shown in Fig. 1.

\[ T_2 = V_2 + V_2 R_0 N_2 T_2 \] (8)

where \( V_2 \) represents the two-body interaction, and \( R_0 \) the interaction independent two-body resolvent. The Pauli blocking factor \( N_2 \), represented by the crosses in Fig. 1, is given by

\[ N_2 = \bar{f}_1 f_2 - f_1 \bar{f}_2, \quad \bar{f} = 1 - f, \] (9)

where the indices of the Fermi-Dirac function \( f \) reflect particle quantum numbers.

![Fig. 2. Loop diagram corresponding to the kernel of the integral equation (4). The crosses refer to the Pauli-blocking factor \( N_2 \).](image2)

For a numerical analysis and a first calculation of in-medium effects we utilize a zero-range model studied earlier in a different context [31,32]. The equation represented by Fig. 1 can be summed and leads to a solution for the two-particle propagator \( \tau(M_2) \), i.e.

\[ \tau(M_2) = (i\lambda^{-1} - B(M_2))^{-1}. \] (10)

The expression for \( B(M_2) \) is represented by the loop diagram of Fig. 2 and, in the rest system of the two-body system \( P^0 = (M_2, 0, 0, 0) \), given by

\[ B(M_2) = -\frac{i}{(2\pi)^3} \int \frac{x}{x(1-x)} \frac{1 - F(x, k_2^2)}{M_2^2 - M_2^{20}}, \] (11)

where

\[ M_2^{20} = (k_2^2 + m^2) / x(1-x), \]

\[ F(x, k_2^2) = f(x, k_2^2) + f(1-x, \bar{k}_2^2) \] (12)

and \( f \) given in eq. (3) with \( x = k^+ / P_n^+ \). The integral can obviously be separated into a blocking independent term \( B_0(M_2) \) and the change \( \Delta B(M_2) \) depending on the Fermi functions \( F(x, k_2^2) \). The integral involving \( B_0(M_2) \) has a logarithmic divergence that can be absorbed in a redefinition of \( \lambda \). The physical information introduced in the renormalization of the amplitude is the mass of the two bound particles, \( M_{2B}(\mu, T) \). If we assume that the two particle amplitude \( \tau \) has a pole for \( M_2(\mu, T) = M_{2B}(\mu, T) \) we may write

\[ i\lambda^{-1} = B(M_{2B}) \equiv B_0(M_{2B}) + \Delta B(M_{2B}), \] (14)

where the dependence on \( T \) and \( \mu \) is suppressed in the notation. Although, we use eq. (14) for finite temperatures and densities, it is easy to assume such a bound state for the isolated case only, and calculate the corresponding in-medium bound state \( M_{2B}(\mu, T) \) from the above equations. This would require a definite regularization procedure which we presently want to omit. The assumption of a bound state implies no restriction on the conclusions of our investigation concerning the relative importance of two- and three-body correlations at a given temperature and density. This will be obvious from the results presented and discussed at the end of this Letter.

For \( M_2(\bar{P}_2 = 0, \mu, T) < 2\mu \) (for \( \mu < m \)) and \( 2m < M_2(\bar{P}_2 = 0, \mu, T) < 2\mu \) (for \( \mu > m \)) instabilities of the Fermi gas against formation of Bose condensation or Cooper pairing can occur. As a consequence the distribution functions need to be modified to include the gap energy. The critical temperature \( T_c \) is given by the condition \( M_2(\bar{P}_2 = 0, \mu, T_c) = 2\mu \), see e.g. [35] for a discussion of similar aspects in nuclear matter. Implementation of such effects into the three-body problem would require the notion of condensate/pairing at finite \( P_2 \) which is a difficult problem that has hardly been investigated. We presently neglect this as well as the appearance of other
two-body correlations in the Fermi functions when solving the three-body problem. However, in its turn, the Dyson approach (cluster mean field expansion) utilized here has the potential to re-evaluate the condition of critical temperature in the presence of strong three-body correlations; that is, however, not the goal of our present work.

The subtraction imposed by condition (14) in the denominator of eq. (13) makes \( \tau(M_2) \) finite. Note that \( \Delta B(M_{2B}) \to 0 \) for large momenta \( \vec{k}_\perp \), so that \( \Delta B(M_{2B}) \) is finite. The resulting expression for the two-body propagator is then given by

\[
\tau(M_2) = \left( i \left[ \kappa(M_{2B}) \arctan 2\kappa^{-1}(M_{2B}) \right] \right) / (2\pi)^2
- \kappa(M_2) \arctan 2\kappa^{-1}(M_2) \right) / (2\pi)^2
+ \Delta B(M_{2B}) - \Delta B(M_2) \right)^{-1},
\]

where

\[
\kappa(M_2) = \sqrt{\frac{m^2}{M_2^2} - \frac{1}{4}}.
\]

The subtraction imposed by condition (14) in the denominator of eq. (13) makes \( \tau(M_2) \) finite. Note that \( \Delta B(M_{2B}) \to 0 \) for large momenta \( \vec{k}_\perp \), so that \( \Delta B(M_{2B}) \) is finite. The resulting expression for the two-body propagator is then given by

\[
\tau(M_2) = \left( i \left[ \kappa(M_{2B}) \arctan 2\kappa^{-1}(M_{2B}) \right] \right) / (2\pi)^2
- \kappa(M_2) \arctan 2\kappa^{-1}(M_2) \right) / (2\pi)^2
+ \Delta B(M_{2B}) - \Delta B(M_2) \right)^{-1},
\]

where

\[
\kappa(M_2) = \sqrt{\frac{m^2}{M_2^2} - \frac{1}{4}}.
\]

The quark mass \( m(\mu, T) \) as well as the two-body masses \( M_2(\mu, T), M_{2B}(\mu, T) \), and \( M_{20}(\mu, T) \) depend on the chemical potential and the temperature of the medium. Also there is an additional dependence on the cluster’s c.m. momentum related to the momentum dependence of the blocking factors. Hence several changes of kinematical variables are necessary in the two-body amplitude before eq. (13) may be employed in a three-body cluster. They are discussed in the following. The two-body mass is

\[
M_2^2 = (P_3 - q)^2
= (M_3 - q^+) \left( M_3 - \frac{q_\perp^2 + m^2}{q^+} \right) - q_\perp^2,
\]

where \( q \) denotes the momentum of the odd particle. The three-body system is taken at rest, \( P_3^\mu = (M_3, 0, 0, 0) \). Also the blocking dependent part \( \Delta B \) needs to be revisited, since the blocking factors depend on the cluster momentum and hence \( M_2 \) is different for a moving system not only because of eq. (17), but also due to medium effects that depend on the momentum. As a consequence the arguments of the blocking factors of eq. (11) have to be properly replaced by

\[
F(x; y; \vec{k}_\perp, q_\perp) = f(x, \vec{k}_\perp^2) + f(1 - x - y, \vec{k} + q_\perp^2).
\]

In the three-body rest frame \( x = k^+/M_3 \) and \( y = q^+/M_3 \).

The homogenous AGS-type in-medium equations for finite temperatures and densities for the nonrelativistic bound states have been given in Refs. [14,19]. A diagrammatic representation of these equations for a zero range interaction is given in Fig. 3. Based on the Fock space representation, a derivation of relativistic three-body equations on the light front is formally identical to the nonrelativistic case, if we neglect antiparticle degrees of freedom. It appears that the modifications required to arrive at an in-medium relativistic three-body equation for (spin averaged) quasi particles on the light front are close to the modifications needed in the nonrelativistic case. The reason is the formal similarity of the light front and nonrelativistic approaches after the approximations that have been mentioned above. Hence a relativistic AGS-type equation on the light front including the effects of finite temperature and density can then be written as (compare Ref. [31,32] for the isolated case)

\[
\Gamma(y, q_\perp) = i \left( \frac{2\pi)^3}{2\pi} \right) \tau(M_2) \int_{M_2/M_3}^{1-y} dx \frac{1}{x(1-y-x)} \int_{k_{\perp}^\text{max}}^{} d^2k_\perp \frac{1-F(x, y; \vec{k}_\perp, q_\perp)}{M_3^2 - M_{20}^2} \Gamma(x, \vec{k}_\perp),
\]

where we have introduced vertex functions \( \Gamma \) and \( \tau(M_2) \) given before. Here

\[
k_{\perp}^\text{max} = \sqrt{(1-x)(xM_3^2 - m^2)},
\]

and the mass of the virtual three-particle state (in the rest system) is

\[
M_{20}^2 = \frac{k_\perp^2 + m^2}{x} + \frac{q_\perp^2 + m^2}{y} + \frac{(\vec{k} + q_\perp)^2 + m^2}{1-x-y},
\]

which is the sum of the on-shell minus-components of the three particles.

The equations suggested here, which result from a systematic Dyson equation approach, differ formally from those of Ref. [7]. For the case \( T = 0 \) at finite densities the blocking factors could be replaced by \( f(\mu, T) = n(k_F(\mu)) = \theta(k - k_F) \) where in the notation of Ref. [7]

\[
\mu(T = 0) = \sqrt{k_F^2 + m^2}.
\]

In Ref. [7] all three quark momenta are restricted according to eq. (22)

\[
(1 - n_1)(1 - n_2)(1 - n_3)
\]

instead of \( 1 - n_1 - n_2 \) (and permutations) from eqs. (18) and (19). Note that in these three-body equations an additional blocking of the spectator particle does not arise in a three-body equation driven by a two-body interaction.

FIG. 3. Diagrammatic representation of the in-medium AGS-type equation for a zero range interaction. The crosses refer to the Pauli-blocking factor \( N_2 \). The two-body input is given in Fig. 1 and Eq. 10.
The correlations between the two-body and three-body binding energies, \( B_2 \) and \( B_3 \) are shown in Fig. 4 for a temperature of \( T = 10 \) MeV. The binding energies are defined by

\[
B_3(\mu, T) = m(\mu, T) + M_{2B}(\mu, T) - M_{3B}(\mu, T) \quad (23)
\]

\[
B_2(\mu, T) = 2m(\mu, T) - M_{2B}(\mu, T). \quad (24)
\]

Binding energies are shown in units of the respective quark masses at the temperature and chemical potential indicated. The dotted line indicates \( M_2/2 = M_3/3 \) (compare \[7\]) and in a simple chemical picture of an ideal gas (law of mass action) the relative importance of the clusters. The isolated case (i.e. no medium) is shown as a solid line. As the chemical potential increases (and the quark mass becomes smaller) three-body correlation become weaker for a given \( B_2 \). In this particular case \( T = 10 \) MeV the three-body bound state disappears for \( \mu \approx 310 \) MeV, even for a model that allows for a two-quark bound state.

In the weak coupling case \( B_2 \sim 0 \) three-body correlations appear stronger, however, for a given temperature the corresponding bound states vanish above a certain value for the chemical potential. For a given chemical potential \( \mu = 100 \) MeV and in the weak coupling limit the respective values for \( B_3 \) are given in Table II.

The dependence of \( B_3(T) \) for a given chemical potential is opposite from the nuclear case, compare Ref. [14]. This is in agreement with the expectation from other approaches that give a negative slope for the phase transition, see e.g. [3]. For strong couplings, i.e. \( B_2/m \gtrsim 0.65 \) and \( T = 10 \) MeV, two-body correlations dominate for all chemical potentials.

### TABLE I

| \( \mu \)/MeV | 10  | 50  | 100 |
|----------------|-----|-----|-----|
| 100            | 300 | 299 | 279 |
| 200            | 300 | 293 | 236 |
| 300            | 289 | -   | -   |
| 307            | 279 | -   | -   |

Quark masses are chosen for values of \( \mu \) and \( T \) to study the Mott transition.

### TABLE II

| \( T \)/MeV | 10  | 50  | 100 |
|-------------|-----|-----|-----|
| 10          | 23  | 19  | 5.0 |
| \( B_3 \)/MeV | 7.6 | 6.5 | 1.8 |

Three-body binding energies for weak coupling \( (B_2/m = 10^{-3}) \) and \( \mu = 100 \) MeV.
In Fig. 5 the same correlations are shown for different temperatures and chemical potentials. For higher temperatures and chemical potentials the three-body bound state disappears and also the three-body correlations become weaker as the lines are below the long-dashed line. In all examples the isolated case is limiting. An increase in $\mu$ means more particles in the medium that preferentially occupy the momentum components necessary to form bound states. A similar effect arises for increasing temperatures. For a given chemical potential $\mu = 100$ MeV Fig. 5 reflects stronger correlations at lower temperatures.

In Fig. 6 we show the dependence of the three-body binding energy on the chemical potential for different temperatures as indicated. $B_3 = 0$ corresponds to the Mott density.

In conclusion, we have derived for the first time a consistent relativistic three-body equation for particles embedded in a medium of both finite temperature and finite density. This equation systematically includes the effects of Pauli blocking and mass shift of the quasi-particles involved. The equations are solved for parameters close to the phase transition of QCD. We find that three-body clusters become less stable for a denser or hotter system than the two-body cluster. This justifies, a posteriori our restriction to the two-body bound states in this investigation. However correlations may still exist as the “pole” moves into the continuum (e.g. corresponding to anti-bound states in this approximation).

For further investigations and to achieve more quantitative results a specific model can be chosen that leads to a specific two-body $t$-matrix and a corresponding three-body bound state. Furthermore, the approximations we used in the treatment of spin of the particles have to be relaxed and a full treatment needs to be implemented along the lines discussed at length in Ref. [30]. The neglect of confinement in the vicinity of the Mott line may be justified by color screening discussed in Ref. [38]. However a more realistic treatment of confinement is highly desirable [34]. The light front approach leading to 3-dim equations is suited to include relativistic confining potentials into the formalism.

In view of the present simplifications, related to the averaging of spin degrees of freedom and the use of a zero
range force, we would like to emphasize that the main result of this Letter is to show that an in-medium three-body equation that reflects the basic requirements of relativity can be derived and solved in the physical region of interest close to the Mott transition. The light-front framework appears to be very useful in this context as it allows for a Hamiltonian formalism close to an intuitive interpretation based on nonrelativistic approaches.

Addressing the question of color-superconductivity in a next step the full three-body $t$-matrix has to be implemented in a calculation of the critical temperature of condensation/pairing. This can also be achieved in a systematic fashion within the Dyson equation approach used here.

Acknowledgment: Work supported by the Deutsche Forschungsgemeinschaft, grant BE 1092/10-1.

[1] F. Karsch, Nucl. Phys. Proc. Suppl. 83, 14 (2000).
[2] M. Alford, A. Kapustin and F. Wilczek, Phys. Rev. D 59, 054502 (1999).
[3] M. Alford, hep-ph/0102047.
[4] M. Alford, K. Rajagopal and F. Wilczek, Phys. Lett. B 422, 247 (1998).
[5] R. Rapp, T. Schafer, E. V. Shuryak and M. Velkovsky, Phys. Rev. Lett. 81, 53 (1998).
[6] Understanding Deconfinement In QCD, proceedings, edited by D. Blaschke, F. Karsch, and C.D. Roberts (Singapore, World Scientific, 2000)
[7] S. Pepin, M. C. Birse, J. A. McGovern and N. R. Walet, Phys. Rev. C 61, 055209 (2000).
[8] L.P. Kadanoff and G. Baym, Quantum Theory of Many-Particle Systems (Mc Graw-Hill, New York, 1962); A.L. Fetter and J.D. Walecka, Quantum Theory of Many-Particle Systems, (McGraw-Hill, New York, 1971).
[9] J. Dukelsky, G. Röpke, and P. Schuck, Nucl. Phys. A628, 17 (1998), and refs. therein.
[10] E. O. Alt, P. Grassberger and W. Sandhas, Nucl. Phys. B 2, 167 (1967).
[11] M. Beyer, G. Röpke and A. Sedrakian, Phys. Lett. B 376, 7 (1996).
[12] M. Beyer and G. Röpke, Phys. Rev. C 56, 2636 (1997).
[13] M. Beyer, Few Body Syst. Suppl. 10, 179 (1999).
[14] M. Beyer, W. Schadow, C. Kuh�ts and G. Röpke, Phys. Rev. C 60, 034004 (1999).
[15] C. Kuh�ts, M. Beyer and G. Röpke, Nucl. Phys. A 668, 137 (2000).
[16] M. Beyer, Nucl. Phys. A684, 556c (2001).
[17] M. Beyer, C. Kuh�ts, G. Röpke and P. D. Danielewicz, in “Progress in Nonequilibrium Green’s Functions”, ed. M. Bonitz (World Scientific, Singapore, 2000) p. 161.
[18] C. Kuh�ts, M. Beyer, P. Danielewicz and G. Röpke, Phys. Rev. C 63, 034605 (2001).
[19] M. Beyer, S. A. Sofianos, C. Kuh�ts, G. Röpke and P. Schuck, Phys. Lett. B 488, 247 (2000).
[20] S. P. Klevansky, Rev. Mod. Phys. 64 (1992) 649.
[21] P. A. Dirac, Rev. Mod. Phys. 21, 392 (1949).
[22] R. J. Perry, A. Harindranath and K. G. Wilson, Phys. Rev. Lett. 65, 2959 (1990).
[23] M. V. Terentev, Sov. J. Nucl. Phys. 24, 106 (1976).
[24] B. L. Bakker, L. A. Kondratyuk and M. V. Terentev, Nucl. Phys. B 158, 497 (1979).
[25] L. A. Kondratyuk and M. V. Terentev, Sov. J. Nucl. Phys. 31, 561 (1980).
[26] L. L. Frankfurt and M. I. Strikman, Nucl. Phys. B 148, 107 (1979).
[27] L. L. Frankfurt and M. I. Strikman, Phys. Rept. 76 (1981) 215.
[28] F. Lenz, Nucl. Phys. (Proc. Suppl.) B 90 (2000), 46.
[29] H.C. Pauli and L.C.L. Hollenberg (eds.) Non-perturbative QCD and hadron phenomenology, Nucl. Phys. (Proc. Suppl.) B 90 (2000).
[30] M. Beyer, C. Kuh�ts and H. J. Weber, Annals Phys. 269, 129 (1998).
[31] T. Frederico, Phys. Lett. B 282, 409 (1992).
[32] W. R. de Araujo, J. P. de Melo and T. Frederico, Phys. Rev. C 52, 2733 (1995).
[33] G.A. Miller and R. Machleidt, Phys. Rev. C60 (1999) 035202; G.A. Miller, Prog. Part. Nucl. Phys. 45 (2000) 83.
[34] J. H. Sales, T. Frederico, B. V. Carlson and P. U. Sauer, Phys. Rev. C 63 (2001) 064003.
[35] J. P. de Melo, T. Frederico, H. W. Naus and P. U. Sauer, Nucl. Phys. A 660 (1999) 219.
[36] T. Alm, G. Röpke and M. Schmidt, Phys. Rev. C 50, 31 (1994).
[37] S. M. Schmidt, diploma thesis, Rostock 1993 (unpublished) and private communication.
[38] F. Karsch, M. T. Mehr and H. Satz, Z. Phys. C 37, 617 (1988).
[39] V. N. Gribov, Eur. Phys. J. C 10, 91 (1999).