Quark deconfinement phase transition for improved quark mass density-dependent model

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Abstract

By using the finite temperature quantum field theory, we calculate the finite temperature effective potential and extend the improved quark mass density-dependent model to finite temperature. It is shown that this model can not only describe the saturation properties of nuclear matter, but also explain the quark deconfinement phase transition successfully. The critical temperature is given and the effect of ω- meson is addressed.

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I. INTRODUCTION

Quark-meson coupling (QMC) model suggested by Guichon [1] is a famous hybrid quark meson model, which can describe the saturation properties of nuclear matter and many other properties of nuclei successfully. In this model, the nuclear system was suggested as a collection of MIT bag, \( \omega \)-meson and \( \sigma \)-meson and the interactions between quarks and mesons are limited within the MIT bag regions because the quark cannot escape from the MIT bag. This model had been extended by many authors, for example, including \( \rho \) meson to discuss the neutron matter [2], adding hyperons to study the strange hadronic matter [3], suggesting the density-dependent vacuum energy to investigate the effect of environment [4] and etc. [5].

Although the QMC model is successful in describing many physical properties of nuclear systems, but as was pointed in our previous paper [6], it has two major shortcomings: (1) It is a permanent quark confinement model because the MIT bag boundary condition cannot be destroyed by temperature and density. It cannot describe the quark deconfinement phase transition. (2) It is difficult to do nuclear many-body calculation beyond mean field approximation (MFA) by means of QMC model, because we cannot find the free propagators of quarks and mesons easily. The reason is that the interactions between quarks and mesons are limited within the bag regions, the multireflection of quarks and mesons by MIT bag boundary must be taken into account for getting the free propagators. These two shortcomings actually inherited from the MIT bag.

To overcome these two shortcomings, in our previous paper [6], we suggested an improved quark mass density-dependent (QMDD) model. The QMDD model was suggested by Fowler, Raha and Weiner [7] many years ago. According to the QMDD model, the masses of \( u \), \( d \), \( s \) quarks and corresponding antiquarks are given by

\[
m_q = \frac{B}{3n_B}(i = u, d, \bar{u}, \bar{d})
\]  

\[
m_{s, \bar{s}} = m_{s0} + \frac{B}{3n_B}
\]

where \( m_{s0} \) is the current mass of the strange quark, \( B \) is the bag constant, \( n_B \) is the baryon number density

\[
n_B = \frac{1}{3}(n_u + n_d + n_s),
\]
$n_u, n_d, n_s$ represent the density of $u$ quark, $d$ quark, and $s$ quark, respectively. As was explained and proved in Ref. [8, 9], the basic hypothesis (1), (2) correspond to a confinement mechanism of quarks because when volume $V \to \infty$, $n_B \to 0$, and then $m_q \to \infty$, quark will be confined. Using QMDD model, many authors investigated the dynamical and thermodynamical properties of strange quark matter and strange star and found that the results given by QMDD model are nearly the same as those obtained in MIT bag model [8-16]. But a nice advantage for QMDD model is that the MIT bag boundary condition has been given up. The ansatz Eq. (1) is employed to replace the MIT bag boundary.

But QMDD model still has two shortcomings: (1) It is still an ideal quark gas model. No interactions between quarks exist except a confinement ansatz Eqs. (1-2). (2) It still cannot explain the quark deconfinement phase transition and give us a correct phase diagram as that given by lattice QCD because when $n_B \to 0, m_q \to \infty, T \to \infty$ [8]. To overcome the difficulty(2), we introduced a new ansatz that $m_q$ is not only a function of density, but also depends on temperature $T$ [8,12,15], we suggested

\[ B = B_0[1 - (T/T_C)^2], \quad 0 \leq T \leq T_C, \quad (4) \]
\[ B = 0, \quad T > T_C \quad (5) \]

and extended the QMDD model to a quark mass density- and temperature- dependent model(QMDTD) model. The ansatz(4) guarantees that $m_q \to 0$ when $T \to T_C$. It changes the permanent confinement mechanism(MIT bag) to a nonpermanent confinement mechanism (Friedberg-Lee(FL) soliton bag) in the QMDTD model. Since the vacuum density $B$ equals to the different value between the local false vacuum minimum and the absolute real vacuum minimum of nonlinear scalar field in FL model, we introduced a nonlinear scalar field to improve the QMDD model in Refs. [17, 22] and changed the ad-hoc ansatz(4) to an output $B(T)$ curve which can be calculated from the improved QMDD model.

To improve the shortcoming(1) of QMDD model, we introduced the $\omega-$ meson and $\sigma-$ meson in QMDD model to mimic the repulsive and attractive interactions between quarks in Refs. [6, 23]. The interaction between quarks and nonlinear $\sigma$ field forms a FL soliton bag [17, 22]. The $qq\omega$ and $qq\sigma$ interaction guarantees that we can get valid saturation properties, the equation of state and the compressibility of nuclear matter [6]. But the deconfinement phase transition has not yet been studied in Ref. [6]. This motivates us to study the deconfinement properties of the improved QMDD(IQMDD) model in this paper. We would
like to emphasize that this is the basic important advantage for IQMDD model, because
the saturation properties can be explained not only by IQMDD model but also by QMC
model. The reason for the explanation of quark deconfinement by IQMDD model is that
MIT boundary constraint has been dropped and interactions between quark and mesons
have been extended to the whole space. Since instead of the MIT bag in QMC model, a
FL soliton bag is introduced in IQMDD model, we can use our model to discuss the quark
deconfinement phase transition. The spontaneous breaking symmetry of nonlinear \( \sigma \) field
will be restored and the soliton bag will disappear at critical temperature [24-27].

The organization of this paper is as follows. In the next section, we give the main formulae
of the IQMDD model and effective potential at finite temperature. The soliton solutions
of IQMDD model at different temperature and other numerical results are presented in the
third section. The last section contains a summary and discussions.

II. THE IQMDD MODEL AT ZERO AND FINITE TEMPERATURE

The effective Lagrangian density of the IQMDD model is given by

\[
\mathcal{L} = \bar{\psi} \left[ i \gamma^\mu \partial_\mu - m_q - f \sigma - g \gamma^\mu \omega_\mu \right] \psi
+ \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - U(\sigma) - \frac{1}{4} F_{\mu \nu} F^{\mu \nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu,
\]  

(6)

where \( F_{\mu \nu} = \partial_\mu \omega_\nu - \partial_\nu \omega_\mu \), \( \psi \) represents the quark field, \( m_q = \frac{B}{3m_B} \) is the mass of \( u(d) \) quark,
the \( \sigma \) and \( \omega \) field are not dependent on time, \( f \) is the coupling constant between the quark
field \( \psi \) and the scalar meson field \( \sigma \), \( g \) is the coupling constant between the quark field \( \psi \)
and the vector meson field \( \omega_\mu \), \( U(\sigma) \) is the self interaction potential for \( \sigma \) field. We omit
the contribution of the \( s \) quark and consider the nuclear system only in this paper. The
potential field \( U(\sigma) \) is chosen as [18]

\[
U(\sigma) = \frac{a}{2!} \sigma^2 + \frac{b}{3!} \sigma^3 + \frac{c}{4!} \sigma^4 + B,
\]  

(7)

\[
b^2 > 3ac
\]  

(8)

The condition \( b^2 > 3ac \) ensures that the absolute minimum of \( U(\sigma) \) is at \( \sigma = \sigma_v \neq 0 \). The
potential \( U(\sigma) \) has two minima: one is the absolute minimum \( \sigma_v \)

\[
\sigma_v = \frac{3|b|}{2c} \left[ 1 + \left( 1 - \frac{8ac}{3b^2} \right)^{\frac{1}{2}} \right],
\]  

(9)
it corresponds to the physical vacuum, and the other is at $\sigma_0 = 0$, it represents a metastable local false vacuum. We take $U(\sigma_v) = 0$ and the bag constant $B$ can be expressed as

$$-B = \frac{a}{2!}\sigma_v^2 + \frac{b}{3!}\sigma_v^3 + \frac{c}{4!}\sigma_v^4.$$  \(10\)

From Eq. (6), we obtain the equation of motion for quark as

$$(i\gamma^\mu\partial_\mu - m_q - f\sigma - g\gamma^\mu\omega)\psi = 0,$$  \(11\)

and the equations for the scalar meson field and vector meson field as

$$\partial_\mu\partial^\mu\sigma + \frac{dU(\sigma)}{d\sigma} = -f\bar{\psi}\psi,$$  \(12\)

$$\partial_\nu F_{\nu\mu} + m_\omega\omega^2 = g\bar{\psi}\gamma^\mu\psi.$$  \(13\)

respectively. Using an approximation as that of the QMC model, we replace $\sigma(r, t) \rightarrow \sigma(r)$, $\omega_\mu(r, t) \rightarrow \delta_{\mu 0}\omega(r)$ and consider a fixed occupation number of valence quarks (3 quarks for nucleons, and quark-antiquark pair for mesons) only. In the following, we will discuss the ground state solution of the system. The Hamiltonian density is

$$\mathcal{H} = \psi^+\left[\frac{1}{i}\vec{\alpha} \cdot \vec{\nabla} + \beta(m_q + f\sigma) + g\omega\right]\psi + \frac{1}{2}\Pi^2$$

$$+ \frac{1}{2}((\nabla\sigma)^2 + U(\sigma) - \frac{1}{2}(\nabla\omega)^2)$

$$- \frac{1}{2}m_\omega^2\omega^2\omega^\mu. \quad \text{(14)}$$

where $\vec{\alpha}$ and $\beta$ are the Dirac matrices $\Pi_\sigma$ is conjugate field of the scalar meson field. One can construct a Fock space of quark states and expand the operator $\psi$ in terms of annihilation and creation operators on the space with spinor function $\varphi_n^\pm$, which satisfies the Dirac equation [17]:

$$[\vec{\alpha} \cdot \vec{p} + \beta(m_q + f\sigma) + g\omega]\varphi_n^\pm = \pm\epsilon_n\varphi_n^\pm. \quad \text{(15)}$$

The functions $\varphi_n$ satisfies the normalized condition $\int \varphi_n^+\varphi_n d^3r = 1$. From Eq. (14), the total energy of the system is given by

$$E(\sigma, \omega) = \sum_n \epsilon_n + \int \frac{1}{2}((\nabla\sigma)^2 + U(\sigma) - \frac{1}{2}(\nabla\omega)^2) - \frac{1}{2}m_\omega^2\omega^2)d^3r. \quad \text{(16)}$$

Substituting Eq. (15) into Eqs. (12, 13), and using the variational principle under the spherical symmetric condition, we find when $\sigma$ and $\omega$ satisfy the follow equations

$$-\nabla^2_\sigma + \frac{dU(\sigma)}{d\sigma} = -f\sum_n\varphi_n\varphi_n,$$
\[-\nabla_r^2 \omega + m_\omega^2 \omega = g \sum_n \varphi_n^+ \varphi_n. \quad (17)\]

respectively where \(\varphi_n = \varphi_n^+ \gamma_0\), we have the minimum of \(E(\sigma, \omega)\).

We discuss the ground state solution of the system now. The quark spinor in the lowest state is assumed [17-19]:

\[
\varphi = \begin{pmatrix} u(r) \\ i(\vec{\sigma} \cdot \vec{v}(r)) \end{pmatrix} \chi_m, \quad \chi_m = \begin{pmatrix} 1 \\ 0 \end{pmatrix} \text{ or } \begin{pmatrix} 0 \\ 1 \end{pmatrix}, \quad (18)
\]

where \(\vec{\sigma}\) are the Pauli matrices. Substituting Eq. (18) into Eq. (15), we get the equations of spinor components \(u\) and \(v\) as

\[
\frac{du(r)}{dr} = -[\epsilon + m_q + f\sigma(r) - g\omega]v(r),
\]

\[
\frac{dv(r)}{dr} = -\frac{2}{r}v(r) + [\epsilon - m_q - f\sigma(r) - g\omega]u(r).
\]

(19)

respectively. The normalized condition then reads as \(4\pi \int_0^\infty [u^2(r) + v^2(r)]r^2dr = 1\).

From Eq. (17) and Eq. (18), after summation of the quark states, we obtain the equation of motion of the \(\sigma, \omega\) field

\[
\frac{d^2\sigma}{dr^2} + \frac{2}{r} \frac{d\sigma}{dr} - \frac{dU(\sigma)}{d\sigma} = Nf(u^2 - v^2),
\]

(20)

\[
\frac{d^2\omega}{dr^2} + \frac{2}{r} \frac{d\omega}{dr} - m_\omega^2 \omega = -Ng(u^2 + v^2) \equiv F(r).
\]

(21)

where, the number of quarks is \(N = 3\) for baryons and \(N = 2\) for mesons. In the following discussions, we only consider the case \(N = 3\). To get a self-consistent solution of Eqs. (19, 20, 21), we select the boundary conditions for quark field and \(\sigma\) field as [17, 22]

\[
v(r = 0) = 0, u(r = \infty) = 0,
\]

\[
\sigma'(r = 0) = 0, \sigma(r = \infty) = \sigma_v.
\]

(22)

respectively. Noting that the \(r \to \infty\) asymptotic behavior of the \(\omega\) field given by Eq. (21) tends to an exponent decay wave because \(F(r \to \infty) \to 0\) for a soliton bag, we can find the corresponding Green function \(G(r, r')\) easily and obtain the \(\omega\) field by integral [20]:

\[
\omega(r) = \int_0^\infty r'^2dr' F(r')G(r, r')
\]

(23)

The numerical results of \(u(r), v(r), \sigma(r)\) and \(\omega(r)\) will be shown in next section.
In order to study the deconfinement phase transition, we turn to extend IQMDD model to finite temperature. The appropriate framework is the finite temperature quantum field theory. The finite temperature effective potential plays a central role within this framework. Under mean field approximation, the meson field operators can be replaced by their expectation values,

\[ \tilde{\omega}_\mu \rightarrow \bar{\omega} = \delta_{\mu 0}, \quad \tilde{\bar{\omega}} = \delta_{\mu 0} \]

Using the method of Dolan and Jackiw [24], up to one-loop approximation, the effective potential reads:

\[ V(\sigma; T; \mu; V_\omega) = U(\sigma) + V_B(\sigma; T) + V_F(\sigma; T; \mu; V_\omega), \quad (24) \]

where

\[ V_\omega = g \bar{\omega} = \frac{g^2}{m^2_\omega} \rho_{B0} \quad (25) \]

is the contribution of \( \omega \)-field, \( \rho_{B0} \) is saturation density of nuclear matter, \( T \) is temperature and

\[ V_B(\sigma; T) = \frac{T^4}{2\pi^2} \int_0^\infty dx x^2 \ln \left( 1 - e^{-\sqrt{(x^2 + m_\sigma^2)/T^2}} \right), \quad (26) \]

\[ V_F(\sigma; T; \mu; V_\omega) = -12 \sum_n \frac{T^4}{2\pi^2} \int_0^\infty dx x^2 \ln \left( 1 + e^{-\left(\sqrt{(x^2 + m_n^2)/T^2} - \mu_n/T + V_\omega/T\right)} \right), \quad (27) \]

where the minus sign of Eq. (27) is the consequence of Fermi-Dirac statistics. The degenerate factor 12 comes from: 2(particle and antiparticle), 2(spin), 3(color). \( m_\sigma \) and \( m_q \) are the effective masses of the scalar field \( \sigma \) and the quark field respectively:

\[ m^2_\sigma = a + b\sigma(T) + \frac{c}{2}\sigma^2(T) \quad (28) \]

\[ m_q = \frac{B(T)}{n_q} + f\sigma(T). \quad (29) \]

We see from Eqs. (24-27) that the scalar-like interaction \( f\psi^+\sigma\psi \) gives contribution to effective masses of quark and \( \sigma \) meson and then forms a confined soliton bag, the vector-like interaction \( g\psi^+\gamma^\mu\omega_\mu\psi \) gives contribution to an effective chemical potential of quarks. In fact, this finite temperature effective potential for FL model had been calculated by many others authors including us [22, 25-27], but without \( \omega \)-field.

In the soliton bag model, the finite temperature vacuum energy density \( B(T) \) is defined as

\[ B(T; \mu; V_\omega) = V(\sigma_0; T; \mu; V_\omega) - V(\sigma_0; T; \mu; V_\omega). \quad (30) \]
It is the different from the values at the perturbative false vacuum state and the values at the physical real vacuum state of the finite temperature effective potential. At critical temperature $T_C$ of quark deconfinement phase transition, $B$ equal to zero: $B(T_C) = 0$.

III. RESULTS

Before numerical calculations, we consider the parameters of IQMDD model at first. To guarantee our model can be used not only to explain the saturation properties of nuclear matter, but also to describe the quark deconfinement phase transition, we fix our parameters as those in Ref. [6]. Hereafter we fix the parameters as: at zero temperature, the bag constant $B = 174 \text{ MeVfm}^{-3}$, the masses of $\omega$- meson and $\sigma$- meson $m_\omega = 783 \text{ MeV}$ and $m_\sigma = 509 \text{ MeV}$, the coupling constant $f = 5.45$ and $g = 3.37$ respectively and the mass of nucleon $M_N = 939 \text{ MeV}$, $b = -8400 \text{ MeV}$. As was shown in Ref. [6], this set of parameters can give us reasonable properties of nuclear matter, including the value of saturation point, the equation of state and the compressibility. We will prove in this section that we can also obtain a reasonable quark deconfinement critical temperature $T_C$ by means of this set of parameters.

The set of coupled differential equations (19), (20) with boundary conditions Eq. (22) can be solved numerically at zero temperature. Our results at zero temperature are shown in Figs. (1-4). The variation of the scalar field $\sigma$ as a function of the radius $r$ is presented in Fig. 1. We see that the value of $\sigma$ inside the hadron is very different from that of outside: inside $\sigma$ is less than zero, and outside $\sigma \to \sigma_v$. The transitional values of $\sigma$ field through the surface from inside to outside is abrupt. The curve in Fig. 1 is very similar as that given by Refs. [17, 22] where the $\omega$- meson is omitted. This is reasonable because the main contribution to form a soliton bag comes from the nonlinear $\sigma$ field. The variation of wave function of quark field is shown in Fig. 2 and Fig. 3. In Fig. 2, we plot the curves of quark wave function $u, v$ vs. $r$ respectively. In Fig. 3, we plot the curve of quark density $u^2 - v^2$ vs. $r$. The soliton bag is exhibited in this figure transparently. The curve of $\omega$- field vs. $r$ given by Eq. (21) is shown in Fig. 4. This solution is obtained from Green function method and Eq. (23). We see from this Figure that the curve of $\omega$- field decays considerably when $r$ is large than the soliton bag radius.

Now we turn to investigate the case of finite temperature. The effective potential at finite
temperature can be obtained by numerical calculations from the set of Eqs. (24-27). The curves of effective potential $U(\sigma; T; \mu; V_\omega)$ at different temperatures $T = 0$ MeV, $T = 80$ MeV, $T = 127$ MeV and $T = 204$ MeV are shown in Fig. 5 respectively. The shape of the effective potential confirms that a first order phase transition will take place [22, 25-27]. We see from Fig. 5 there are two vacua where one corresponds to the physical vacuum and the other corresponds to the false vacuum when $T < 127$ MeV. But when $T$ increases to a critical temperature $T_C = 127$ MeV, these two vacua degenerate and the values of bag constant tends to zero.

To show this behaviour more clearly, we plot the bag constant $B$ vs. $T$ curve by solid curve in Fig. 6, we see that the bag constant $B$ decreases as temperature $T$ increases. When $T$ approaches to $T_C$, $B$ approaches to zero and quark deconfined phase transition happens. When temperature increases to the regions $127$ MeV $< T < 204$ MeV, as shown in Fig. 5 the physical vacuum becomes instable and the perturbative vacuum becomes stable. The soliton solution tends to disappear [22]. When $T$ approached to 204 MeV, the effective potential becomes to have an unique minimum only. Such potential no longer ensures the existence of the soliton bag anymore because the spontaneous breaking symmetry has been restored. We have zero solution $\sigma = 0$ in the regions $T \geq 204$ MeV only.

To illustrate our soliton solutions more transparently, we plot the soliton solutions by fixed the temperature $T = 0$ MeV, $T = 80$ MeV and $T = 125$ MeV in Fig. 7. It is shown in Fig. 7 that the curves stretch slowly to infinite with increasing temperature. The radius of soliton increases and the skin of the soliton becomes thicker when the temperature increases. It is of course very reasonable.

All characters of the soliton solutions found by IQMDD model are almost the same as those given by Ref. [22] where the $\omega$ field has not been take into account. But we hope to emphasize although the $\omega$ field will not affect the main characters of the soliton solutions, but it will change their detailed behavior. To address the effect of $\omega$ field on the deconfinement transition, as an example, we plot the $B(T)$ curve without $\omega$ field by dotted line in Fig. 6. We find the critical temperature from the condition $B(T_C) = 0$ changes to $T_C = 140$ MeV which is higher than 127 MeV where the $\omega$ field existed. If we notice that the interaction of $qq\omega$ is repulsive, and then it can help to deconfine quarks, the decrease of $T_C$ due to $\omega$-field is natural.
IV. SUMMARY AND DISCUSSION

By using the finite temperature quantum field theory, we calculate the finite temperature effective potential and extend our previous discussions to finite temperature. It is shown that the IQMDD model can not only describe the saturation properties of nuclear matter, but also explain the quark deconfinement phase transition successfully. We have shown the soliton solution curves for different temperatures and found the critical temperature $T_C = 127$ MeV. The $\omega$- field in IQMDD model is important. The repulsive $qq\omega$ interaction plays the central role to describe the saturation properties, and affect the critical temperature remarkably. Comparing to naive QMC model, the advantage of IQMDD model is obvious because instead of MIT bag, a Friedberg-Lee soliton bag exists in this model.

Of course, there are still shortcomings in the IQMDD model. For example, the interactions between quarks and mesons are still isospin independent. The chiral phase transition cannot be discussed by this model because it is lack of chiral symmetry. To Overcome these shortcomings, we hope to add the $\rho$- meson and $\pi$- meson in IQMDD model in the near future.

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FIG. 1: The σ field as a functions of r at zero temperature.
FIG. 2: Quark wave functions in arbitrary unit as a functions of $r$. 
FIG. 3: The quark density $u^2(r) - v^2(r)$ in arbitrary unit as functions of $r$. 
FIG. 4: The $\omega$ field as a functions of $r$. 

\(\omega (\text{MeV})\)

\(r (\text{fm})\)
FIG. 5: The temperature-dependent effective potential.
FIG. 6: The bag constant $B(T)$ as functions of $T$, the solid line corresponds to IQMDD model with $\omega$-meson and the dotted line corresponds to the same model without $\omega$-meson.
FIG. 7: The soliton solutions for different temperature $T=0$ MeV, 80 MeV, 125 MeV