Limit Temperatures for Meson and Diquark Resonances in a Strongly Interacting Quark Matter

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We investigate mesons and diquarks as resonant states above chiral critical temperature \( T_c \) in flavor SU(2) Nambu–Jona-Lasinio model. For each kind of resonance, we solve the pole equation for the resonant mass in the complex energy plane, and find an ultimate temperature where the pole starts to disappear. The phase diagram including these limit temperatures in \( T - \mu \) plane is obtained. The maximum limit temperature at \( \mu = 0 \) is approximately two times \( T_c \).

PACS numbers: 11.30.Rd, 11.15.Ex, 25.75.Nq

From the collective phenomena observed in high energy heavy ion collisions at RHIC, the formed new matter is not a weakly coupled quark-gluon plasma (QGP), but in a strongly coupled region. Theoretically, the lattice theory and the perturbative resummation techniques have told us that QCD matter at temperatures not far above \( T_c \), where \( T_c \) is the critical temperature for deconfinement and chiral restoration phase transitions, is not a perturbative QGP. Since the interaction is strong enough, there might be glueballs, mesons, diquarks and baryons in the system at temperatures \( T > T_c \). Recently, the meson mass above \( T_c \) is studied through lattice calculation in chiral limit, and the lattice baryonic susceptibility is explained with a picture including baryons at \( T > T_c \). If these composite states of quarks and gluons do exist above \( T_c \), they should be melted away when the temperature of the system is high enough. Therefore, a natural question one asks is what the ultimate temperature for the composite states is.

In the idealized case at asymptotically high baryon density, the phase transitions and the related new matters have been widely discussed from first principle QCD calculations. For the region above but close to the phase transition line, the study depends on effective models. It is well-known that the Nambu–Jona-Lasinio (NJL) model applied to quarks offers a simple but effective scheme to study chiral symmetry restoration, color superconductivity, and pion superfluidity at finite temperature and moderate number densities. In the chiral restoration phase, the quarks keep massless, and the massive mesons and diquarks can decay into quark-antiquark and quark-quark pairs. Therefore, the mesons and diquarks above \( T_c \) are not stable bound states, but rather resonant states, the pole equations for meson and diquark masses should be regarded in their complex forms in order to determine the resonant masses and widths self-consistently. This was done for mesons in the case without considering color superconductivity. In this paper, we discuss the effect of chiral restoration and color superconductivity on the meson and diquark resonances above \( T_c \), and try to extract the ultimate temperatures of these resonances from the thermodynamics of the NJL model.

The flavor SU(2) NJL model is defined through the Lagrangian density,

\[
\mathcal{L} = \bar{\psi} (i\gamma^\mu \partial_\mu + \mu \gamma_0) \psi + G_s \left( (\bar{\psi} \psi)^2 + (\bar{\psi} i \gamma_5 \tau \psi)^2 \right) + G_d \left( \bar{\psi}_1^c i \gamma^5 \epsilon^{ij} \epsilon^{\alpha \beta \gamma} \psi_{j\beta}^c \right) \left( \bar{\psi}_{i\alpha} i \gamma^5 \epsilon^{ij} \epsilon^{\alpha \beta \gamma} \psi_{j\beta} \right),
\]

where \( \mu \) is the quark chemical potential, \( G_s \) and \( G_d \) are coupling constants in color singlet channel and anti-triplet channel, \( \tau = (\tau_1, \tau_2, \tau_3) \) are Pauli matrices in the flavor space, and \( \epsilon_{ij} \) and \( \epsilon_{\alpha \beta \gamma} \) are totally antisymmetric tensors in the flavor and color spaces.

The quark-antiquark and diquark condensates which are order parameters of chiral and color superconductivity phase transitions, respectively, are defined as

\[
\sigma = -2G_s \langle \bar{\psi} \psi \rangle, \quad \Delta = -2G_d \langle \bar{\psi}_{1\alpha} i \gamma^5 \epsilon^{ij} \epsilon^{\alpha \beta \gamma} \psi_{j\beta} \rangle,
\]

where it has been regarded that only the first two colors participate in the diquark condensate, while the third one does not. In mean field approximation, the two condensates as functions of \( T \) and \( \mu \) are determined by the gap equations,

\[
\sigma (1 - 2G_s I_1) = 0, \quad \Delta (1 - 2G_d I_2) = 0,
\]

with the functions \( I_1 \) and \( I_2 \) defined as

\[
I_1(\sigma, \Delta) = 4 \sum_{p, \alpha} \frac{1}{E_p} \left[ \frac{E^\sigma_p}{E^\Delta_p} \tanh \frac{E^\sigma_p}{2T} + \frac{1}{2} \frac{E^\sigma_p}{2T} \right],
\]

\[
I_2(\sigma, \Delta) = 4 \sum_{p, \alpha} \frac{1}{E^\Delta_p} \tanh \frac{E^\Delta_p}{2T},
\]

where \( E^\pm_p = \sqrt{(E^\sigma_p)^2 + \Delta^2} \) are the energies of the effective quarks which participate in the diquark condensate, \( E^\pm_p = E_p \pm \mu \) with \( E_p = \sqrt{p^2 + M_q^2} \) and quark mass \( M_q = \sigma \) are the energies of the other quarks which are not involved in the diquark condensate, and \( \sum_{p, \alpha} = \int \frac{d^3p}{(2\pi)^3} \sum_{\alpha = \pm} \) includes momentum integration and energy summation. The two gap equations determine self-consistently the chiral and color superconductivity phase transition lines in the temperature and...
chemical potential plane, which are, respectively, shown as dashed and dot-dashed lines in Fig. In the following we focus on the meson and diquark properties above the chiral critical temperature $T_c$ where $\sigma$ keeps zero, while $\Delta$ is zero in normal phase but finite in color superconductivity phase.

In the NJL model, the meson and diquark modes are regarded as the poles of their effective propagators constructed by the mean field quarks. When the quark propagator in Nambu-Gorkov space is diagonal in normal phase, the summation of bubbles in RPA selects its specific channel by choosing at each stage the same proper polarization function, any of the four mesons and six diquarks is related to its own polarization function only. However, when the quark propagator is with off-diagonal elements in the phase of color superconductivity, we must consider carefully all possible channels in the bubble summation in RPA.

The mesons $\sigma$, $\pi_+$, $\pi_-$, $\pi_0$ are the eigen models of the system in any phase, the degenerated mass $M_m$ is determined through the pole equation

$$1 - 2G_s \Pi_{mm}(k_0) = 0,$$

with the meson polarization function

$$\Pi_{mm}(k_0) = I_1 + 4k_0^2 J(k_0),$$

where $I_1$ is related to the gap equations and the function $J(k_0)$ is given by

$$J(k_0) = -\sum_{p,\alpha,\beta} \left[ \frac{1}{2} \frac{\tanh E_0^\alpha}{2E_0^\alpha} + \frac{E_0^+ E_0^- + E_0^+ E_0^- + \Delta^2}{E_0^-(E_0^- + E_0^+)} k_0^2 - (E_0^- + E_0^+)^2 \right],$$

$$\tanh \frac{E_0^\alpha}{2} + \frac{E_0^+ E_0^- + E_0^+ E_0^- + \Delta^2}{E_0^-(E_0^- + E_0^+)} k_0^2 - (E_0^- + E_0^+)^2 \right].$$

Since the diquarks condense only in the color 3 direction, see the definition of the diquark condensate, the color symmetry is spontaneously broken from SU(3) to SU(2) in color superconductivity phase. While the mass $M_d$ of the four diquarks $d_1$, $d_2$, $\bar{d}_1$, and $\bar{d}_2$, constructed by the quarks with colors 2 and 3, 1 and 3, $\bar{1}$ and $\bar{3}$, and $\bar{2}$ and $\bar{3}$, is characterized only by the diquark polarization $\Pi_{dd}(k_0)$ and $\Pi_{\bar{d}\bar{d}}(k_0)$ in any phase,

$$(1 - 2G_d \Pi_{dd}(k_0)) (1 - 2G_d \Pi_{\bar{d}\bar{d}}(k_0)) = 0,$$

the diquarks $d_3$ and $\bar{d}_3$, constructed by the quarks with colors 1 and 2, and $\bar{1}$ and $\bar{2}$, are no longer the eigen modes of the system in color superconductivity phase, the new eigen modes are the linear combinations of $d_3$ and $\bar{d}_3$, and their mass is controlled by the diagonal and off-diagonal elements,

$$\det \begin{pmatrix} 1 - 2G_d \Pi_{dd}(k_0) & -2G_d \Pi_{d\bar{d}d}(k_0) \\ -2G_d \Pi_{\bar{d}\bar{d}d}(k_0) & 1 - 2G_d \Pi_{\bar{d}\bar{d}}(k_0) \end{pmatrix} = 0,$$

with the diquark polarization functions

$$\Pi_{dd}(k_0) = \Pi_{\bar{d}\bar{d}}(-k_0) = I_2 - 8k_0 K_1(k_0) + 4k_0^2 K_2(k_0),$$

$$\Pi_{d\bar{d}d}(k_0) = \Pi_{\bar{d}\bar{d}d}(k_0) = I_2 + 8k_0 K_3(k_0) + 4 \left(k_0^2 - 8\Delta^2\right) K_4(k_0),$$

where $I_2$ is related to the gap equations and the functions $K_i(k_0)$ are defined as

$$K_1(k_0) = -\frac{1}{2} \sum_{p,\alpha,\beta} \alpha \left[ \frac{E_0^\alpha}{E_0^\alpha + 1} \frac{\tanh \frac{E_0^\beta}{2}}{k_0^2 - \left(E_0^\alpha + E_0^\beta\right)^2} + \frac{\tanh \frac{E_0^\alpha}{2}}{k_0^2 - \left(E_0^\alpha - E_0^\beta\right)^2} \right],$$

$$K_2(k_0) = -\frac{1}{2} \sum_{p,\alpha,\beta} \frac{1}{2E_0^\alpha} \left[ \frac{\tanh \frac{E_0^\beta}{2}}{k_0^2 - \left(E_0^\alpha + E_0^\beta\right)^2} + \frac{\tanh \frac{E_0^\alpha}{2}}{k_0^2 - \left(E_0^\alpha - E_0^\beta\right)^2} \right],$$

$$K_3(k_0) = -\sum_{p,\alpha,\beta} \frac{E_0^\alpha}{E_0^\alpha k_0^2 - 4 \left(E_0^\alpha\right)^2},$$

$$K_4(k_0) = -\sum_{p,\alpha,\beta} \frac{1}{E_0^\alpha k_0^2 - 4 \left(E_0^\alpha\right)^2}.$$

It is easy to check that, the pole equation is reduced to $\Delta = 0$ in normal phase.

With the help of the gap equation in color superconductivity phase, $1 - 2G_d I_2(0, \Delta) = 0$, the mass equation
for the diquarks \(d_1, d_2, \bar{d}_1 \) and \(\bar{d}_2\) is simplified to

\[
k_0^2 (k_0^2 K_2^2(k_0) - 4 K_1^2(k_0)) = 0. \tag{12}
\]

Obviously, one solution is \(k_0^2 = 0\), and the nonzero root is determined by

\[
k_0^2 K_2^2(k_0) - 4 K_1^2(k_0) = 0. \tag{13}
\]

Similarly, the mass equation \(9\) for the two new eigen modes in color superconductivity phase is reduced to

\[
k_0^2 ((k_0^2 - 4 \Delta^2) K_2^2(k_0) - 4 K_3^2(k_0)) = 0. \tag{14}
\]

Again, one of its solution is \(k_0^2 = 0\), and the nonzero root is calculated through

\[
(k_0^2 - 4 \Delta^2) K_2^2(k_0) - 4 K_3^2(k_0) = 0. \tag{15}
\]

Therefore, there are three massless diquarks and three massive diquarks in color superconductivity phase. Two of the massive diquarks determined by \(13\) are degenerated and very light, and the other massive diquark controlled by \(15\) is heavy. Taking into account the Higgs mechanism, the three massless and the double degenerated light diquarks will be eaten up by gauge fields, only the heavy diquark is left in color superconductivity phase. In the following, we consider the four mesons in normal and color superconductivity phases, the six diquarks in normal phase, and the heavy diquark in color superconductivity phase.

Above \(T_c\) any of the mass equations should be regarded in its complex form at

\[k_0 = M - i \Gamma/2 \tag{16}\]

in order to determine the resonant mass \(M\) and width \(\Gamma\) self-consistently. From the above pole equations, one sees that this corresponds to solving the conditions

\[
M_m^2 - (ImM_1(M_m)/2M_m)^2 - ReM_1(M_m) = 0,
\]

\[
\Gamma_m = -ImM_1(M_m)/M_m
\]

for the meson \(M_m\) and width \(\Gamma_m\),

\[
M_d^2 - \left(Im\left(M_2(M_d) + \sqrt{M_3(M_d)}\right)/2M_d\right)^2 - ReM_2(M_d) = 0,
\]

\[
\Gamma_d = -Im\left(M_2(M_d) + \sqrt{M_3(M_d)}\right)/M_d
\]

for the masses \(M_d\) and widths \(\Gamma_d\) of the two triple degenerated diquarks in normal phase, and

\[
M_d^2 - (ImM_4(M_d)/2M_d)^2 - ReM_4(M_d) = 0,
\]

\[
\Gamma_d = -ImM_4(M_d)/M_d
\]

for the mass \(M_d\) and width \(\Gamma_d\) of the heavy diquark in color superconductivity phase, where the complex functions \(M_i\) are defined as

\[
M_1(k_0) = \frac{1 - 2G_s I_1(0, 0)}{8G_s J(k_0 - i\epsilon)},
\]

\[
M_2(k_0) = \frac{1 - 2G_d I_2(0, 0)}{8G_d K_2(k_0 - i\epsilon)} + \frac{K_1^2(k_0 - i\epsilon)}{K_2^2(k_0 - i\epsilon)},
\]

\[
M_3(k_0) = \frac{4 K_1^2(k_0 - i\epsilon)}{K_2^2(k_0 - i\epsilon)} \left( \frac{K_1^2(k_0 - i\epsilon)}{K_2^2(k_0 - i\epsilon)} + \frac{1 - 2G_d I_2(0, 0)}{8G_d K_2(k_0 - i\epsilon)} \right),
\]

\[
M_4(k_0) = \frac{4 K_1^2(k_0 - i\epsilon)}{K_2^2(k_0 - i\epsilon)} + \Delta^2.
\]

The real parts \(ReJ\) and \(ReK_i\) are just the functions \(J\) and \(K_i\) themselves, and the imaginary parts \(ImJ\) and \(ImK_i\) can be written as

\[
ImJ(k_0 - i\epsilon) = \pi \sum_{p,\alpha} \left[ \frac{\tanh \frac{\alpha}{2T}}{E_p^\alpha} \delta_1 + \frac{E_p^\alpha E_{\Delta} + \alpha E_p^\alpha E_{\Delta}}{2E_p^\alpha E_{\Delta} \left( E_{\Delta} + \alpha E_p^\alpha \right)^2} \left( \frac{\tanh \frac{E_p^\alpha}{2T} + \alpha \tanh \frac{E_{\Delta}}{2T}}{\tanh \frac{E_p^\alpha}{2T} + \tanh \frac{E_{\Delta}}{2T}} \right) \right],
\]

\[
ImK_1(k_0 - i\epsilon) = \pi \sum_{p,\alpha} \frac{1}{E_{\Delta}^\alpha} \left[ \alpha \left( \frac{\tanh \frac{E_p^\alpha}{2T} + \tanh \frac{E_{\Delta}^\alpha}{2T}}{\tanh \frac{E_p^\alpha}{2T} - \tanh \frac{E_{\Delta}^\alpha}{2T}} \right) \delta_3 - \left( \frac{\tanh \frac{E_p^\alpha}{2T} - \tanh \frac{E_{\Delta}^\alpha}{2T}}{\tanh \frac{E_p^\alpha}{2T} - \tanh \frac{E_{\Delta}^\alpha}{2T}} \right) \delta_3 \right],
\]

\[
ImK_2(k_0 - i\epsilon) = \pi \sum_{p,\alpha} \frac{1}{E_{\Delta}^\alpha} \left[ \frac{\tanh \frac{E_p^\alpha}{2T} + \tanh \frac{E_{\Delta}^\alpha}{2T}}{E_p^\alpha + E_{\Delta}^\alpha} \delta_3 - \frac{\tanh \frac{E_p^\alpha}{2T} - \tanh \frac{E_{\Delta}^\alpha}{2T}}{E_p^\alpha - E_{\Delta}^\alpha} \delta_3 \right],
\]

\[
ImK_3(k_0 - i\epsilon) = \pi \sum_{p,\alpha} \frac{E_p^\alpha}{(E_{\Delta}^\alpha)^2} \tanh \frac{E_p^\alpha}{2T} \delta_2,
\]

\[
ImK_4(k_0 - i\epsilon) = \pi \sum_{p,\alpha} \frac{E_p^\alpha}{(E_{\Delta}^\alpha)^2} \delta_2,
\]

with the \(\delta\) functions \(\delta_1 = \delta(k_0 - 2E_p), \delta_2 = \delta(k_0 - 2E_{\Delta}), \delta_3 = \delta(k_0 - (E_p^\alpha \pm E_{\Delta}^\alpha)), \) and \(\delta_\alpha = \delta(k_0 \pm (E_{\Delta}^\alpha \pm \alpha E_p^\alpha)),\) corresponding to energy conservations for different decay channels. In deriving \(17, 18\) and \(19\), we have assumed that the imaginary part \(\Gamma\) in \(M_i\) may be
neglected, so that the equations for $M$ and $\Gamma$ are decoupled.

In chiral limit there are three parameters in the NJL model, the momentum cutoff $\Lambda$ and the two coupling constants $G_s$ and $G_d$. $\Lambda$ and $G_s$ can be fixed through fitting the constituent quark mass and pion decay constant in the vacuum, which leads to $\Lambda = 0.65$ GeV and $G_s = 5.01$ (GeV)$^{-2}$ \cite{21}. While one can not fix $G_d$, we can determine its low and high limits by taking into account the physical constraints on the diquark mass in the vacuum \cite{21}. $G_d^{\min} < G_D < G_d^{\max}$ with $G_d^{\min} \sim 0.8 G_s$ and $G_d^{\max} = 1.5 G_s$.

It is now ready to calculate the resonant masses and widths above the chiral phase transition line. At the critical temperature $T_c$ where one has $1 - 2 G_s J_1(0, \Delta) = 0$, the pole equation for the meson mass \cite{15} is reduced to $k_0^2 J(k_0) = 0$, and the solution is $k_0^2 = 0$. Therefore, the degenerated meson mass evolves with starting value $M_m(T_c) = 0$, reflecting correctly the chiral property at the critical point. From our numerical calculation, the meson mass goes up monotonously with increasing temperature, and finally ends at a maximum temperature $T_r$ where there is no more root for the complex mass equation and the pole of the meson propagator disappears. This limit temperature as a function of $\mu$ is plotted as a solid line in Fig. 1(a) in the case without considering the diquark channel. The mesons are in bound states at low temperature $T < T_c$, but in resonant states in the region $T_c < T < T_r$. The maximum limit temperature $T_r(0)$ is about two times the chiral critical value $T_c(0)$. Above the limit temperature $T_r$, the meson resonances disappear and there are only quarks in the system.

Including the diquark channel, we calculated the meson and diquark masses and widths as functions of $T$ and $\mu$ for $G_d = G_d^{\min}$ and $G_d = G_s$. For each kind of resonance, there exists an ultimate temperature where the pole of the corresponding propagator vanishes, and the resonance disappears from the system. These limit temperatures as functions of $\mu$ are shown in Fig. 1(b) and (c) as solid lines. The dashed and dot-dashed lines are the phase transition lines for the chiral restoration and color superconductivity calculated through the gap equations \cite{4}. The symbols $M, M_D, M_{DH}$ and $M_{DL}D_H$ in normal phase indicate the resonant regions with only mesons, mesons and light diquarks, mesons and heavy diquarks, and mesons and light and heavy diquarks, respectively, and $M_{DH}$ means the meson and heavy diquark resonances in color superconductivity phase.

In normal phase, the limit temperature for mesons is independent of the coupling $G_d$, while the regions with diquark resonances depend strongly on the strength of $G_d$. The light triple degenerated diquarks can be generated at any possible $G_d$, but the heavy triple degenerated diquarks are created only in the case with strong coupling. In color superconductivity phase, the heavy diquark left can be produced everywhere at any coupling $G_d$. Compared with the case without considering diquark channel shown in Fig. 1(h), the resonant states in Fig. 1(b) and (c) at high baryon density can exist in a wider region, due to the contribution from the diquark condensate. When the temperature of the system is higher than the maximum of the limit temperatures, all the resonances disappear, and there are only quarks in the system.

In summary, we have investigated the meson and di-
quark resonances in a strongly interacting quark matter above but close to the chiral critical temperature $T_c$ in the NJL model. Above $T_c$, the massive mesons and diquarks can decay into two massless quarks and become resonant states. In mean field approximation to quarks and random phase approximation to mesons and diquarks, we derived the pole equations for the resonances in complex energy plane. For each kind of massive resonances, there exists a limit temperature $T_r$ where the pole disappears and the resonances are melted in the hot and dense quark matter. The maximum limit temperature at $\mu = 0$ is approximately two times $T_c$, which agrees well with the estimated critical temperatures for mesons, diquarks and baryons in lattice and phenomenological calculations\[1, 2, 5\]. In color superconductivity phase, the meson and heavy diquark resonances are survived everywhere.

**Acknowledgments:** The work is supported by the grants NSFC10428510, 10435080 10575058 and SRFDP20040003103.

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