Phase diagram of two-flavor quark matter: Gluonic phase at nonzero temperature

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The phase structure of neutral two-flavor quark matter at nonzero temperature is studied. Our analysis is performed within the framework of a gauged Nambu–Jona-Lasinio model and the mean-field approximation. We compute the free energy of the gluonic phase (gluonic cylindrical phase II) in a self-consistent manner and investigate the phase transition from the gluonic phase to the $2\text{SC}/g2\text{SC}/NQ$ phases. We briefly consider the phase diagram in the plane of coupling strength versus temperature and discuss the mixed phase consisting of the normal quark and $2\text{SC}$ phases.

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I. INTRODUCTION

The properties of cold and dense quark matter are of great interest in astrophysics and cosmology. In particular, at moderate densities of relevance for the interior of compact stars, quark matter is a color superconductor and has a rich phase structure with important implications for compact star physics [1, 2, 3, 4, 5, 6, 7, 8, 9].

Bulk matter in the interior of compact stars should be color and electrically neutral and be in $\beta$-equilibrium. In the two-flavor case, these conditions separate the Fermi momenta of up and down quarks and, as a consequence, the ordinary BCS state ($2\text{SC}$) is not always energetically favored over other unconventional states. The possibilities include crystalline color superconductivity and gapless color superconductivity ($g2\text{SC}$) [10, 11, 12, 13]. However, the $2\text{SC}/g2\text{SC}$ phases suffer from a chromomagnetic instability, indicated by imaginary Meissner masses of some gluons [14, 15]. The instability related to gluons of color 4–7 occurs when the ratio of the gap over the chemical potential mismatch, $\Delta/\delta\mu$, decreases below a value $\sqrt{2}$. Resolving the chromomagnetic instability and clarifying

![Phase diagram of two-flavor quark matter](image)

FIG. 1: The phase diagram of electrically neutral two-flavor quark matter in the plane of $\Delta_0$ and $T$. At $T = 0$, the $g2\text{SC}$ phase exists in the window $92 \text{ MeV} < \Delta_0 < 134 \text{ MeV}$ and the $2\text{SC}$ window is given by $\Delta_0 > 134 \text{ MeV}$. The unstable region for gluons 4–7 is depicted by the region enclosed by the thin solid line. The $g2\text{SC}$ phase and a part of the $2\text{SC}$ phase ($92 \text{ MeV} < \Delta_0 < 162 \text{ MeV}$) suffer from the chromomagnetic instability at $T = 0$. The quark chemical potential is taken to be $\mu = 400 \text{ MeV}$.

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the nature of true ground state of dense quark matter are central issues in the study of color superconductivity. In models without dynamic gauge fields, one has to introduce color and electric chemical potentials ($\mu$) and phases. A self-consistent analysis of the gluonic phases at $T < 65$ MeV and they found that the gluonic phase (strictly speaking, the gluonic cylindrical phase II) exists in the window $\mu_{\text{unpaired}} = 400$ MeV, which is a value typical for the cores of compact stars. The region enclosed by the thick solid line is unstable (because gluons of color 4–7 have tachyonic Meissner masses there) and, therefore, should be replaced by other chromomagnetically stable phases, for instance, gluonic phases [24, 25]. For more detailed discussions of the gluonic phases, see Refs. [45, 46]. Note, however, that we did not consider the global structure of a free energy in extracting the unstable region, but only the tendency toward the vector condensation.

In this paper, we study the gluonic cylindrical phase II at nonzero temperature and revisit the phase diagram shown in Fig. 1, computing the free energy of the gluonic phase in a self-consistent manner. The result would be useful for the phase diagram of QCD, and the compact star phenomenology as well.

II. MODEL

In order to study the gluonic phase, we use the gauged NJL model with massless up and down quarks:

\[
\mathcal{L} = \bar{\psi}(i \slashed{D} + \hat{\mu}\gamma^0)\psi + G_D \left( \bar{\psi}i\gamma_5\varepsilon\sigma^b C\psi^T \right) \left( \psi C\gamma_5\varepsilon\sigma^b \psi \right) - \frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu},
\]

where the quark field $\psi$ carries flavor ($i, j = 1, \ldots, N_f$ with $N_f = 2$) and color ($\alpha, \beta = 1, \ldots, N_c$ with $N_c = 3$) indices, $C$ is the charge conjugation matrix; $(\varepsilon)^{ik} = \varepsilon^{ik}$ and $(\lambda)^{\alpha\beta} = \delta^{\alpha\beta}$ are the antisymmetric tensors in flavor and color spaces, respectively. The diquark coupling strength in the scalar ($J^P = 0^+$) color-antitriplet channel is denoted by $G_D$. The covariant derivative and the field strength tensor are defined as

\[
D_\mu = \partial_\mu - igA_\mu^a T^a,
\]

\[
F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + gf^{abc} A_\mu^b A_\nu^c.
\]

To evaluate loop diagrams we use a three-momentum cutoff $\Lambda = 653.3$ MeV throughout this paper. In NJL-type models without dynamic gauge fields, one has to introduce color and electric chemical potentials ($\mu_8$ and $\mu_e$) by hand to ensure color- and electric-charge neutrality. In $\beta$-equilibrated neutral two-flavor quark matter, the elements of the diagonal matrix of quark chemical potentials $\bar{\mu}$ are given by

\[
\mu_{ar} = \mu_{ug} = \bar{\mu} - \delta\mu, \\
\mu_{dr} = \mu_{dg} = \bar{\mu} + \delta\mu, \\
\mu_{ub} = \bar{\mu} - \delta\mu - \mu_8, \\
\mu_{db} = \bar{\mu} + \delta\mu - \mu_8,
\]

with

\[
\bar{\mu} = \mu - \frac{\delta\mu}{3} + \frac{\mu_8}{3}, \quad \delta\mu = \frac{\mu_e}{2}.
\]

In Nambu-Gor’kov space, the inverse full quark propagator $S^{-1}(p)$ is written as

\[
S^{-1}(p) = \begin{pmatrix}
(S_0^+)^{-1} & \Phi^- \\
\Phi^+ & (S_0^-)^{-1}
\end{pmatrix},
\]

with

\[
(S_0^+)^{-1} = \gamma^\mu p_\mu + (\bar{\mu} - \delta\mu_T^3)\gamma^0 + g\gamma^\mu A_\mu^a T^a,
\]

\[
(S_0^-)^{-1} = \gamma^\mu p_\mu - (\bar{\mu} - \delta\mu_T^3)\gamma^0 - g\gamma^\mu A_\mu^a T^a T, \tag{6a}
\]

\[
(S_0^-)^{-1} = \gamma^\mu p_\mu - (\bar{\mu} - \delta\mu_T^3)\gamma^0 - g\gamma^\mu A_\mu^a T^a T, \tag{6b}
\]
and

\[ \Phi^- = -i \varepsilon^b \gamma_5 \Delta, \quad \Phi^+ = -i \varepsilon^b \gamma_5 \Delta. \]  

(7)

Here \( \tau^3 = \text{diag}(1, -1) \) is a matrix in flavor space. Following the usual convention, we have chosen the diquark condensate to point in the third direction in color space.

For the gluonic cylindrical phase \( II = (gA_5^0) \) is the most relevant condensate, because the chromomagnetic instability related to gluons \( 4-7 \) corresponds to the tachyonic mode in the direction of \( B \). \( 24, 25, 30 \) Besides \( B \), we have to introduce a color chemical potential \( \mu_3 = (gA_5^0) \) to ensure color neutrality at \( B \neq 0 \). Taking into account these condensates, the free energy of the gluonic phase in the one-loop approximation is given by

\[
\Omega(\Delta, \mu_c, \mu_8, B, \mu_3; \mu, T) = -\frac{\mu_3^2 B^2}{2g^2} + \frac{\mu_3 \mu_8 B^2}{2g^2} - \frac{\mu_3^2 B^2}{8g^2} - \frac{1}{12\pi^2} \left( \mu^4_c + 2\pi^2 T^2 \mu^2_c + \frac{7\pi^4 T^4}{15} \right) + \frac{\Delta^2}{4G_D} - \frac{1}{2} \sum \frac{d^3 p}{(2\pi)^3} \left[ |\epsilon_a| + 2T \ln(1 + e^{-|\epsilon_a|}) \right],
\]

(8)

where \( \beta = 1/T \), the \( \epsilon_a \)'s are quasi-quark energies and the sum runs over all particle and anti-particle \( \epsilon_a \)'s. Here, we added tree-level contributions from gluons (first line on the r.h.s.),

\[
\Omega_g^{(\text{tree})} = \frac{g^2}{4} \epsilon_a \epsilon_a \epsilon_a \epsilon_a A^\mu_\alpha A_\mu A_\nu A_\nu,
\]

(9)

and electrons (second line on the r.h.s.). Note also that the \( \epsilon_a \)'s depend on the vector condensates through the covariant derivatives in the quark propagator. In what follows, we neglect the color chemical potentials \( \mu_{3,8} \) and, consequently, the tree-level contributions of gluons. We have carefully checked that their effect on the free energy is negligible for realistic values of \( \alpha_s \simeq 1 \) (see also Ref. \( 36 \).

In this work, in order to remove the ultraviolet divergence in the Meissner screening masses we shall use the following subtraction

\[
\Omega_R = \Omega(\Delta, \mu_c, B; \mu, T) - \Omega(0, 0, B; 0, 0).
\]

(10)

It is known that this free energy subtraction is not adequate to remove the cutoff dependence of the free energy at \( T > 0 \). In fact, Eq. (10) leads to positive Meissner screening masses in the normal phase at \( T > 0 \). In this work we do not go into this problem because this unphysical behavior of the Meissner masses is nothing but a cutoff artifact and moreover is negligibly small at \( \mu = 400 \text{ MeV} \) and at the temperatures of interest (20 MeV at most).

In order to find the neutral gluonic phase, we first solve a set of coupled equations (the gap equation and the electrical charge neutrality condition),

\[
\frac{\partial \Omega_R}{\partial \Delta} = \frac{\partial \Omega_R}{\partial \mu_c} = 0,
\]

(11)

as a function of \( B \) and, then, compute the free energy \( \Omega_R(B) \). Finally, the minimum of \( \Omega_R(B) \) determines the neutral gluonic phase. (In the following Figs. 3 and 4 we plot the free energy evaluated along the solution of the coupled equations (11).)

III. NUMERICAL RESULTS

Figure 2 shows \( \Delta, \delta \mu \) and \( B \) in the gluonic phase at \( T = 0 \) as a function of \( \Delta_0 \). First, let us note that the results of Fig. 2 are in good agreement with those shown in Figs. 1, 2, and 3 of Ref. \( 36 \), where the color chemical potentials \( \mu_{3,8} \) were treated self-consistently. In Fig. 2 one can see that the gluonic phase exists in the window

\[
66 \text{ MeV} < \Delta_0 < 162 \text{ MeV}.
\]

(12)

The gluonic phase is energetically favored over the 2SC/g2SC/NQ phases in this whole window (see also Fig. 5 of Ref. \( 36 \)). One also sees that the phase transition between the gluonic phase and the NQ (2SC) phase at \( \Delta_0 = 66 \text{ MeV} \) (162 MeV) is strongly (weakly) of first order.
FIG. 2: The gap parameter $\Delta$ (solid line), the chemical potential mismatch $\delta \mu$ (dotted line) and the gluonic vector condensate $B$ (dashed line) versus $\Delta_0$ in the gluonic phase at $T = 0$. The quark chemical potential is taken to be $\mu = 400$ MeV.

FIG. 3: The free energy $\Omega_R(B)$ as a function of $B$ at $T = 0$ for $\Delta_0 = 75$ MeV (solid line), $\Delta_0 = 85$ MeV (dotted line), $\Delta_0 = 100$ MeV (dot-dashed line), and $\Delta_0 = 140$ MeV (dashed line). Note that the free energy is measured with respect to the 2SC/g2SC/NQ phases at $B = 0$. The results are plotted for $\mu = 400$ MeV.

Now let us take a closer look at the free energy at $T = 0$. Figure 3 shows the behavior of $\Omega_R(B)$ measured with respect to the 2SC/g2SC/NQ phases at $B = 0$. The results are plotted for $\mu = 400$ MeV at several values of $\Delta_0$.

In the weak coupling regime, $66 \text{ MeV} < \Delta_0 < 92 \text{ MeV}$, the chromomagnetic instability does not exist in the NQ phase (see Fig. 1). We note that the curvature of $\Omega_R(B)$ at $B = 0$,

$$m_M^2 = \left. \frac{d^2 \Omega_R(B)}{dB^2} \right|_{B=0},$$

(13)

can be regarded as the Meissner mass squared $\partial^2 \Omega_R/\partial B^2|_{B=0}$ in the 2SC/g2SC/NQ phases, since the solutions of Eq. (11) satisfy $\Delta = \bar{\Delta} + O(B^2)$ and $\mu_e = \bar{\mu}_e + O(B^2)$ for small values of $B$, where $\Delta$ and $\bar{\mu}_e$ denote their values at $B = 0$ [24, 25]. We found that $m_M^2$ is indeed zero in the weak coupling regime. In addition, we observed that, for small values of $B$, the system is in the ungapped ($\Delta = 0$) phase and the free energy behaves like $\Omega_R \sim O(B^4)$. 
However, contrary to the result of Fig. 1, the free energy has a global minimum at \( B \neq 0 \) and the gluonic phase is energetically favored over the NQ phase. For \( 92 \text{ MeV} < \Delta_0 \lesssim 162 \text{ MeV} \), one finds tachyonic modes at \( B = 0 \) because the g2SC phase and a part of the 2SC phase suffer from the chromomagnetic instability and, therefore, are unstable against the formation of \( B \). Consequently, the gluonic phase is realized in this region, as expected. For strong coupling, \( \Delta_0 \gtrsim 162 \text{ MeV} \), the 2SC phase is chromomagnetically stable in this regime and the free energy has a global minimum at \( B = 0 \), though it is not plotted in Fig. 3.

We now turn to the free energy of the gluonic phase at \( T > 0 \). Figure 4(a) display the temperature dependence of the free energy for \( \Delta_0 = 75 \text{ MeV} \). As \( T \) grows, the free-energy gain gets reduced, but the change of the vacuum expectation value of \( B \) is rather small. As a result, we observe a strong first-order transition from the gluonic phase to the NQ phase at \( T \approx 14 \text{ MeV} \). Note that \( m_{3\lambda}^2 \) remains positive at any value of \( T \), which is consistent with the result shown in Fig. 1.

In Fig. 4(b), the same plot is displayed for \( \Delta_0 = 85 \text{ MeV} \). At low temperature, like in the case of \( \Delta_0 = 75 \text{ MeV} \), the gluonic phase is more favored than the chromomagnetically stable NQ phase. At \( T \approx 9 \text{ MeV} \), \( m_{3\lambda}^2 \) turns negative, meaning that the stable NQ phase undergoes a phase transition into the unstable g2SC phase (see Fig. 1). The gluonic phase is energetically favored until the temperature reaches \( T \approx 20 \text{ MeV} \). Above this temperature, the g2SC phase becomes stable and therefore is favored.

In Figs. 4(c) and 4(d), we plot the free energy for the cases of \( \Delta_0 = 100 \text{ MeV} \) and \( \Delta_0 = 140 \text{ MeV} \), respectively. In both cases, the gluonic phase is energetically favored at low temperature. In contrast, at high temperature, the global minima of the free energy are realized at \( B = 0 \), which means that, as expected from the result of Fig. 1, the
FIG. 5: Schematic phase diagram of neutral two-flavor quark matter at moderate density in $\Delta_0-T$ plane. The thick solid line denotes the line of second-order or weakly first-order transitions and strong first-order transitions are indicated by a thick dashed line. In the region enclosed by the thick solid and dashed lines, the gluonic phase is energetically more favored than the 2SC/g2SC/NQ phases.

chromomagnetically stable 2SC/g2SC phases are favored. For $\Delta_0 = 100$ MeV, the phase transition from the gluonic phase to the g2SC phase takes place at $T \simeq 21$ MeV. In the case of $\Delta_0 = 140$ MeV, the phase transition takes place at $T \simeq 18$ MeV.

Here, we would like to make a comment regarding the order of the phase transitions. As mentioned above, we observed the strong first-order transition (gluonic phase ↔ NQ phase) at $\Delta_0 = 75$ MeV. On the other hand, for the cases of $\Delta_0 = 85, 100, 140$ MeV, the phase transition (gluonic phase ↔ 2SC/g2SC phases) is likely to be of second order. However, evaluating the free energy self-consistently near the critical temperatures is not easy and hence we do not exclude the possibility of weak first-order transitions. Furthermore, it should be also mentioned that, because of the cutoff artifact in Eq. (10), it might be impossible to distinguish a weak first-order transition from a second-order one.

IV. SUMMARY, CONCLUSIONS, AND OUTLOOK

A. Phase diagram

We studied the gluonic cylindrical phase II at nonzero temperature. Using the gauged NJL model and the one-loop approximation, we computed the free energy of the gluonic phase self-consistently and investigated the phase structure of the gluonic phase. Although we neglected the color chemical potentials, we have checked that, for $\alpha_s \simeq 1$, their effect on the free energy is negligible.

In the weak coupling regime, we found that the gluonic phase undergoes a strong first-order transition into the NQ phase as it is heated. This is a new aspect of the gluonic phase at $T > 0$, which is not shown in Fig. 1. On the other hand, since the phase transitions from the gluonic phase to the chromomagnetically stable 2SC/g2SC phases are of second-order or weakly first-order, we expect that the corresponding critical line shown in Fig. 1(i.e., the right branch of the thick solid line) is not drastically altered by the self-consistent analysis. (In other words, the Meissner masses squared can be a rough criterion for choosing the energetically favored phase in this regime.) We thus are able to make a sketch of a schematic phase diagram of two-flavor quark matter, which is free from the chromomagnetic instability related to gluons 4–7 (see Fig. 5).  

1 It is interesting to note that the neutral single plane-wave Larkin-Ovchinnikov-Fulde-Ferrell state [48, 49] has a similar phase structure as the gluonic phase [26, 31, 32, 54].
enlarges its region. We argue therefore that the gluonic phase which could resolve the chromomagnetic instability related to gluons 4–7 is a strong candidate for the ground state of a neutral two-flavor color superconductor in the intermediate coupling regime. Alternatives include other types of the gluonic phases [24, 25, 36, 37], the crystalline phases [10, 11, 43] and the mixed phase [19]. It should be mentioned that, at $T = 0$, the gluonic color-spin locked phase is more stable than the gluonic cylindrical phase II in some region of $\Delta_0$ and moreover is free from the chromomagnetic instability at moderate densities [54, 55].

Although we concentrated on the phase diagram in $T$-$\Delta_0$ plane in this work, it is obviously worthwhile to revisit the phase diagram in $T$-$\mu$ plane. A preliminary study [50] indicates that currently known phase diagrams [51, 52, 53] must be significantly altered. In addition, the critical temperature for the gluonic phase could reach a few tens of MeV and, therefore, it is interesting to study astrophysical implications of the gluonic phase, e.g., the quark matter equation of state, neutrino emission from compact star cores, and so on.

### B. Gluonic phase versus mixed phase

Finally we briefly look at a mixed phase consisting of the NQ and the 2SC phases [19, 20, 21]. For the mixed phase to exist, it must satisfy the Gibbs conditions, which are equivalent to chemical and mechanical equilibrium conditions between the NQ and the 2SC phases. These conditions end up as follows

$$P^{(\text{NQ})}(\mu, \mu_e) = P^{(\text{2SC})}(\mu, \mu_e).$$

(14)

Beside Eq. (14) two components must have opposite electrical charge densities. Otherwise a globally neutral mixed phase could not exist. We solved Eq. (14) and found that the globally neutral mixed phase exists in the window

$$67 \text{ MeV} < \Delta_0 < 201 \text{ MeV},$$

(15)

where the quark chemical potential was taken to be $\mu = 400$ MeV.

In order to calculate the free energy of the mixed phase we take account of finite-size effects, i.e., the surface and Coulomb energies associated with phase separation. The surface and Coulomb energy densities are given by

$$\epsilon_s = \frac{dx\sigma}{r_0}, \quad \epsilon_C = 2\pi\alpha_{\text{em}} f_d(x) x (\Delta n_e)^2 r_0^2,$$

(16)

where $\sigma$ is the surface tension, $x$ is the volume fraction of the rarer phase, $\Delta n_e$ is the difference of the electric charge density between NQ and 2SC phases, and $\alpha_{\text{em}} = 1/137$. These energy densities also depend on the dimension $d$ ($d = 1, 2,$ and 3 correspond to slabs, rods, and droplets configurations, respectively) and $r_0$, which denotes the radius of the rarer phase. The geometrical factor $f_d(x)$ is given by

$$f_d(x) = \frac{1}{d+2} \left( \frac{2-d x^{1-2/d}}{d-2} + x \right).$$

(17)
Minimizing the sum of $\epsilon_S$ and $\epsilon_C$ with respect to $r_0$, we obtain

$$\epsilon_S + \epsilon_C = \frac{3}{2} \left(4\pi\alpha_{em}d^2 f_d(x)x^2\right)^{1/3} \left(\Delta n_e\right)^{2/3} \sigma^{2/3}. \quad (18)$$

The actual value of the surface tension in quark matter is poorly known, in this work we assume $d = 3$ (droplets configuration) and try relatively small surface tension.

Figure 6 displays the free energy of the 2SC/g2SC phase, the gluonic phase, and the mixed phase. For a very small surface tension $\sigma = 3$ MeV/fm$^2$, the mixed phase is the most favored in a wide range of $\Delta_0$, $103$ MeV $< \Delta_0 < 166$ MeV. The gluonic phase is energetically more favored than the mixed phase only in the weak coupling regime. Note that the value of the surface tension, $\sigma = 3$ MeV/fm$^2$ at $\mu = 400$ MeV, is close to that calculated by Reddy and Rupak [19]. For a surface tension $\sigma = 10$ MeV/fm$^2$, there still is a wide window where the mixed phase is more stable than the g2SC phase, but the mixed phase is less favored than the gluonic phase.

It should be mentioned here that, however, we did not take into account the thickness of the boundary layer, which has been estimated to be comparable to the value of the Debye screening length in each of the two phases, and therefore the results shown in Fig. 6 is not a final conclusion [56, 57]. The effect of charge screening would increase the surface energy substantially [58, 59].

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[1] K. Rajagopal and F. Wilczek, in *At the Frontier of Particle Physics/Handbook of QCD*, edited by M. Shifman (World Scientific, Singapore, 2001).
[2] M. G. Alford, Annu. Rev. Nucl. Part. Sci. 51, 131 (2001).
[3] D. K. Hong, Acta. Phys. Pol. B 32, 1253 (2001).
[4] S. Reddy, Acta. Phys. Pol. B 33, 4101 (2002).
[5] D. H. Rischke, Prog. Part. Nucl. Phys. 52, 197 (2004).
[6] R. Casalbuoni and G. Nardulli, Rev. Mod. Phys. 76, 263 (2004).
[7] M. Buballa, Phys. Rept. 407, 205 (2005).
[8] M. Huang, Int. J. Mod. Phys. E 14, 675 (2005).
[9] I. A. Shovkovy, Found. Phys. 35, 1309 (2005).
[10] M. Alford, J. A. Bowers, and K. Rajagopal, Phys. Rev. D 63, 074016 (2001).
[11] J. A. Bowers and K. Rajagopal, Phys. Rev. D 66, 065002 (2002).
[12] I. Shovkovy and M. Huang, Phys. Lett. B 564, 205 (2003).
[13] M. Huang and I. Shovkovy, Nucl. Phys. A729, 835 (2003).
[14] M. Huang and I. A. Shovkovy, Phys. Rev. D 70, 051501(R) (2004).
[15] M. Huang and I. A. Shovkovy, Phys. Rev. D 70, 094030 (2004).
[16] I. Giannakis and H. C. Ren, Phys. Lett. B 611, 137 (2005).
[17] I. Giannakis and H. C. Ren, Nucl. Phys. B723, 255 (2005).
[18] I. Giannakis, D. f. Hou, and H. C. Ren, Phys. Lett. B 631, 16 (2005).
[19] S. Reddy and G. Rupak, Phys. Rev. C 71, 025201 (2005).
[20] F. Neumann, M. Buballa, M. Oertel, Nucl. Phys. A714, 481 (2003).
[21] I. Shovkovy, M. Hananske, M. Huang, Phys. Rev. D 67, 103004 (2003).
[22] M. Huang, Phys. Rev. D 73, 045007 (2006).
[23] D. K. Hong, hep-ph/0506097.
[24] E. V. Gorbar, M. Hashimoto, and V. A. Miransky, Phys. Lett. B 632, 305 (2006).
[25] E. V. Gorbar, M. Hashimoto, and V. A. Miransky, Phys. Rev. D 75, 085012 (2007).
[26] E. V. Gorbar, M. Hashimoto, and V. A. Miransky, Phys. Rev. Lett. 96, 022005 (2006).
[27] K. Fukushima, Phys. Rev. D 73, 094016 (2006).
[28] E. V. Gorbar, M. Hashimoto and V. A. Miransky, and I. A. Shovkovy, Phys. Rev. D 73, 111502(R) (2006).
[29] M. Hashimoto, Phys. Lett. B 642, 93 (2006).
[30] O. Kiriyama, D. H. Rischke, and I. A. Shovkovy, Phys. Lett. B 643, 331 (2006).
[31] O. Kiriyama, Phys. Rev. D 74, 074019 (2006).
[32] O. Kiriyama, Phys. Rev. D 74, 114011 (2006).
[33] K. Iida and K. Fukushima, Phys. Rev. D 74, 074020 (2006).
[34] L. He, M. Jin, and P. Zhuang, Phys. Rev. D 75, 036003 (2007).
[35] R. Gatto and M. Ruggieri, Phys. Rev. D 75, 114004 (2007).
[36] M. Hashimoto and V. A. Miransky, Prog. Theor. Phys. 118, 303 (2007).
[37] E. J. Ferrer and V. de la Incera, Phys. Rev. D 76, 045011 (2007).
[38] R. Casalbuoni, R. Gatto, M. Mannarelli, G. Nardulli, and M. Ruggieri, Phys. Lett. B 605, 362 (2005); 615, 297(E) (2005).
[39] K. Fukushima, Phys. Rev. D 72, 074002 (2005).
[40] A. Kryjevski, hep-ph/0508180.
[41] T. Schäfer, Phys. Rev. Lett. 96, 012305 (2006).
[42] M. Ciminale, G. Nardulli, M. Ruggieri, and R. Gatto, Phys. Lett. B 636, 317 (2006).
[43] K. Rajagopal and R. Sharma, Phys. Rev. D 74, 094019 (2006).
[44] X. B. Zhang and J. I. Kapusta, Phys. Rev. D 75, 054012 (2007).
[45] E. V. Gorbar, J. Jia, and V. A. Miransky, Phys. Rev. D 73, 045001 (2006).
[46] A. Buchel, J. Jia, and V. A. Miransky, Nucl. Phys. B772, 323 (2007).
[47] M. Buballa and I. A. Shovkovy, Phys. Rev. D 72, 097501 (2005).
[48] A. I. Larkin and Y. N. Ovchinnikov, Zh. Eksp. Teor. Fiz. 47, 1136 (1964) [Sov. Phys. JETP 20, 762 (1965)].
[49] P. Fulde and R. A. Ferrell, Phys. Rev. 135, A550 (1964).
[50] O. Kiriyama, arXiv:0805.3301 [hep-ph].
[51] S. B. Rüster, V. Werth, M. Buballa, I. A. Shovkovy, and D. H. Rischke, Phys. Rev. D 72, 034004 (2005).
[52] D. Blaschke, S. Fredriksson, H. Grigorian, A. M. Öztas, and F. Sandin, Phys. Rev. D 72, 065020 (2005).
[53] H. Abuki and T. Kunihiro, Nucl. Phys. A768, 118 (2006).
[54] M. Hashimoto and J. Jia, Phys. Rev. D 76, 114019 (2007).
[55] M. Hashimoto, Phys. Rev. D 78, 031501 (2008).
[56] H. Heiselberg, C. J. Pethick, E. F. Staubo, Phys. Rev. Lett. 70, 1355 (1993).
[57] N. K. Glendenning, S. Pei, Phys. Rev. C 52, 2250 (1995).
[58] T. Norsen and S. Reddy, Phys. Rev. C 63, 065804 (2001).
[59] D. N. Vorkresensky, M. Yasuhira, and T. Tatsumi, Nucl. Phys. A723, 291 (2003).