Neutral Larkin–Ovchinnikov–Fulde–Ferrell state and chromomagnetic instability in two-flavor dense QCD

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(Dated: March 26, 2022)

In two-flavor dense quark matter, we describe the dynamics in the single plane wave Larkin–Ovchinnikov–Fulde–Ferrell (LOFF) state satisfying the color and electric neutrality conditions. We find that because the neutral LOFF state itself suffers from a chromomagnetic instability in the whole region where it coexists with the (gapped/gapless) two-flavor superconducting (2SC/g2SC) phases, it cannot cure this instability in those phases. This is unlike the recently revealed gluonic phase which seems to be able to resolve this problem.

PACS numbers: 12.38.-t, 11.15.Ex, 11.30.Qc

It is natural to expect that cold quark matter may exist in the interior of compact stars. This fact motivated intensive studies of this system over the past few years (for a review, see Ref. [1]). Matter in the bulk of stars should be electrically neutral. It should be also a color singlet. The electric and color neutrality conditions play an important role in the dynamics of quark pairing [2, 3]. Several states have been studied in two and three-flavor dense QCD as possible ground states satisfying these conditions [2, 3, 4, 5, 6, 7].

Recently, it has been revealed that the (gapped/gapless) two-flavor superconducting (2SC/g2SC) phase suffers from a chromomagnetic instability connected with the presence of imaginary (tachyonic) Meissner masses of gluons [8]. Later a chromomagnetic instability has been also found in the color-flavor locked phase of three-flavor QCD [9]. It is clear that the resolution of this problem is intimately connected with the determination of the genuine ground state. Therefore it is one of the central issues in the field.

In this Letter, we will study this problem in neutral two-flavor dense QCD. At present, there are two types of dynamics suggested for its resolving. In Refs. [10, 11, 12], the single plane Larkin–Ovchinnikov–Fulde–Ferrell (LOFF)-like dynamics [13, 14] were considered. [Although the dynamics considered in Refs. [11, 12] look different from the conventional LOFF one, by using an appropriate gauge transformation, one can show that they are in fact equivalent to the latter.] The second type is the dynamics of the gluonic phase introduced and studied recently in Ref. [15]. In that phase, gluonic degrees of freedom play a crucial role and the dynamics is manifestly non-abelian. It is clear that it would be important to discriminate these two types of dynamics in order to establish the genuine ground state. It is the primary goal of this Letter.

We recall that in the two-flavor case, the manifestations of the chromomagnetic instability are quite different in the regimes with \( \delta \mu < \Delta < \sqrt{2} \delta \mu \) and \( \Delta < \delta \mu \) [8], where \( \delta \mu \) yields a mismatch between chemical potentials for the up and down quarks and \( \Delta \) is a diquark gap. The (strong coupling) regime with \( \delta \mu < \Delta \) corresponds to the 2SC solution, and the (intermediate coupling) regime with \( \Delta < \delta \mu \) corresponds to the gapless g2SC one [8]. While in the g2SC solution both the 4-7-th gluons and the 8-th one have tachyonic masses, in the 2SC solution, with \( \delta \mu < \Delta < \sqrt{2} \delta \mu \), only the 4-7-th gluons are tachyonic.

The main result of this Letter is that the neutral LOFF state itself suffers from a chromomagnetic instability in the whole region where it coexists with the 2SC/g2SC states. Because of that, it cannot cure this instability in those states. This is unlike the gluonic phase which seems to be able to resolve this problem.

Studying the neutral LOFF state in two-flavor quark matter is a nontrivial problem. Up to now, such studies have been done only in the weak coupling dynamical regime in which neither the neutral 2SC nor neutral g2SC solutions exist [16]. Here we will consider the neutral LOFF phase both in the strong coupling and intermediate coupling regimes that allows to clarify the role of the LOFF dynamics for curing the chromomagnetic instability.

To study dense two-flavor quark matter in \( \beta \)-equilibrium, we consider a phenomenological Nambu–
Jona-Lasinio (NJL) model whose Lagrangian is
\[ \mathcal{L} = \bar{\psi}(i\gamma_\mu \mu)\psi + G_\Delta \left[ (\bar{\psi}C i\varepsilon e^{a}\gamma_5 \psi) (\bar{\psi} e^{a}\gamma_5 \psi C) \right], \]
where \( \varepsilon^{ij} \) and \( e^{a}_{\beta\gamma} \) are the antisymmetric tensors in the flavor and color spaces, respectively (as usual, we neglect the current quark masses). In \( \beta \)-equilibrium, the elements of the diagonal chemical potential matrix \( \mu \) for up \((u)\) and down \((d)\) quarks are \( \mu_u = \mu_d = \mu - \delta \mu, \mu_{ub} = \bar{\mu} - \delta \mu - \mu_8 \) and \( \mu_{dr} = \mu_{dg} = \mu + \delta \mu, \mu_{db} = \bar{\mu} + \delta \mu - \mu_8 \), where \( \bar{\mu} \equiv \mu - \delta \mu / 3 + \mu_8 / 3 \) and \( \delta \mu \equiv \mu / e / 2 \). Here the subscripts \( r, g, \) and \( b \) correspond to red, green and blue quark colors, \( \mu \) is the quark chemical potential (the baryon chemical potential \( \mu_B \) is given by \( \mu_B \equiv 3 \mu \)), \( \mu_c \) is the chemical potential for the electric charge, and \( \mu_8 \) is the color chemical potential. The latter is connected with the vacuum expectation value (VEV) of the time component of the 8-th gluon.\[17\]

By using the auxiliary field \( \Phi^a \sim i \bar{\psi} C \varepsilon e^{a}\gamma_5 \psi \), the Lagrangian density \[14\] can be rewritten as
\[ \mathcal{L} = \bar{\psi}(i\gamma_\mu \mu)\psi - \frac{\Phi^a \Phi^a}{4G_\Delta} - \frac{1}{2} \mu^a [\bar{\psi} e^{a}\gamma_5 \psi C] - \frac{1}{2} [\bar{\psi} C e^{a}\gamma_5 \psi] \Phi^a \Phi^a. \]
(2)
The 2SC and g2SC states have a nonzero diquark condensate \( \langle \Phi^a \rangle |_{\alpha = \beta} = \Delta \) chosen along the blue color direction. As a result, the color group \( SU(3)_c \) is spontaneously broken to \( SU(2)_c \) and blue quarks are gapless.

The order parameter of a single plane wave LOFF state has the form \( \langle \Phi^a(\vec{x}) \rangle = \Delta e^{2i\vec{q} \cdot \vec{x}} \)\[14\], with a constant modulus \( \Delta \) and a constant vector \( \vec{q} \) in the exponent. Then it is convenient to perform a gauge transformation \( \psi \rightarrow \psi' = e^{x p(-2\sqrt{3} \bar{\vec{q}} \cdot \vec{x} T^8) \psi} \), where \( T^8 = \frac{1}{2\sqrt{3}} \text{diag}(1, 1, -2) \) is the eighth generator of \( SU(3)_c \) in the fundamental representation. This gauge transformation trades the \( x \)-dependent phase of the LOFF order parameter for the \( x \)-independent term \( -2\sqrt{3} \bar{\vec{q}} \cdot \vec{x} T^8 \psi \) in the kinetic term for quarks. Obviously, the equality
\[ \bar{q} = \frac{1}{2\sqrt{3}} g \langle \bar{\Phi}^8 \rangle \]
(3)
takes place, i.e., the vector \( \vec{q} \) is proportional to the VEV of a spatial component of the eighth gluon.

To derive the free energy (effective potential) we need to calculate the fermion determinant. For this purpose, it is convenient to introduce the Nambu–Gor’kov spinor \( \Psi^T = (\psi_{ir}, \psi_{ir}^C, \psi_{ig}, \psi_{ig}^C, \psi_{ib}) \). Then, noting that the contribution of the gapless blue quarks factorizes and applying the standard formulas for 2 by 2 matrices in the red and green quarks sector, we find
\[ \det S^{-1}(p_0, \vec{p}) = \prod \left[ \frac{(p^0 - \bar{\mu} + \delta \mu \tau^3 - \mu_8)^2 - (\vec{p} - 2\vec{q})^2}{(p^0 - \bar{\mu} + \delta \mu \tau^3 + \mu_8)^2 - (\vec{p} + 2\vec{q})^2} \right] \prod \left[ \left( (p^0 - \bar{\mu} \tau^3 - Q)^2 - (E_{\Delta-q}^+)\right)^2 \left( (p^0 - \bar{\mu} \tau^3 + Q)^2 - (E_{\Delta-q}^-)^2 \right) \right] + 4\Delta^2(Q^2 - q^2)^2, \]
(4)
where we defined
\[ Q = \frac{1}{2} \left( |\vec{p} + \vec{q}| - |\vec{p} - \vec{q}| \right), q \equiv |\vec{q}|, \]
(5)
\[ E_{\Delta-q}^\pm = \sqrt{\frac{1}{4} \left( E_{p+q}^\pm + E_{p-q}^\pm \right)^2 + \Delta^2}; \]
(6)
\[ E_{p+q}^\pm = |\vec{p} + \vec{q}| \mp \bar{\mu}. \]
(7)
The free energy is expressed through an integral of \( \ln \det S^{-1} \). Because the vector \( \vec{q} \) in the LOFF phase is just a redundant variable for gapless blue quarks, one can put \( \vec{q} = 0 \) in the part of the free energy connected with these quarks. Then it is obvious that their contribution to the free energy coincides with that of free massless fermions. In order to perform the loop integral for the red and green quark parts, we make a usual assumption of the dominance of the region around the Fermi surface, with \( p \equiv |\vec{p}| \sim \bar{\mu} \). Then, neglecting contributions suppressed by \( 1 / \bar{\mu}^2 \), we can omit the last term \( 4\Delta^2(Q^2 - q^2)^2 \) in Eq.\[14\] and use the approximate relations
\[ Q = q \cos \theta + O \left( \frac{q^3}{p^2} \right), \quad \cos \theta = \frac{\vec{p} \cdot \vec{q}}{pq}, \]
(8)
\[ E_{\Delta-q}^\pm = E_{\Delta}^\pm + \frac{q^2}{2p} \frac{p + \bar{\mu}}{E_{\Delta}^\pm} (1 - \cos^2 \theta) + O \left( \frac{q^4}{p^2 E_{\Delta}^\pm} \right), \]
(9)
where \(E_{\lambda}^2 \equiv (p + \mu)^2 + \Delta^2\). Evaluating now the loop integral in the hard dense loop approximation, we obtain the following expression for the free energy of the two-flavor LOFF state:

\[
\Omega(\bar{\mu}, \delta \mu, \mu_s, \Delta, q) = \Omega_{2SC} + \frac{2\mu^2 q^2}{3\pi^2} + \frac{\mu^2}{\pi^2} \left[ \frac{(q + \delta \mu)^3}{q} \left( \frac{1}{2} (1 - x_1^2) \ln \frac{1 + x_1}{1 - x_1} - x_1 + \frac{2}{3} x_1^3 \right) + (q \to -q) \right],
\]

with the 2SC part \(\Omega_{2SC}\) being

\[
\Omega_{2SC} = -\frac{\mu_s^4}{12\pi^2} - \frac{\mu_b^4}{12\pi^2} - \frac{\mu_s^4}{12\pi^2} + \frac{\bar{\mu}^4}{3\pi^2} + \frac{\Delta^2}{4G_\Delta} - \frac{\bar{\mu}^2 \Delta^2}{\pi^2} \ln \frac{4(\Lambda^2 - \bar{\mu}^2)}{\Delta^2} - \frac{\Delta^2}{\pi^2} (\Lambda^2 - 2\mu^2),
\]

where \(\Lambda\) is a ultraviolet cutoff in the NJL model, \(x_1 \equiv \theta(1 - \frac{\Delta^2}{(\mu^2 + m^2)}) \sqrt{1 - \frac{\Delta^2}{(\mu^2 + m^2)}}\) with the step function \(\theta(x)\), and the contribution of electrons was also included in \(\Omega_{2SC}\).

The following remarks are in order. Because we use the sharp cutoff \(\Lambda\) in evaluating the loop integral, it is necessary to subtract the corresponding vacuum contribution with \(\bar{\mu} = \mu_s = \delta \mu = \Delta = 0\) and the same value of the vector \(\bar{q}\). This leads to removing the spurious term \(\Lambda^2 \bar{q}^2/\pi^2\) in the free energy. One can easily check that for \(q = 0\), expression (10) yields the correct 2SC/g2SC free energy, and for \(\Delta = 0, q = 0\), it reproduces the correct free energy for the normal phase \(18\). Last but not least, as will be discussed below, the one-loop (i.e., mean field) approximation we use is reasonable for realistic values of the chemical potential \(\mu\).

The equations of motion for the free energy (10) yield four equations: Two gap equations for \(\Delta\) and \(q\) and two neutrality conditions for \(\delta \mu\) and \(\mu_s\):

\[
a) \partial \Omega / \partial \Delta = 0, \quad b) \partial \Omega / \partial q = 0, \quad c) \partial \Omega / \partial \delta \mu = 0, \quad d) \partial \Omega / \partial \mu_s = 0.
\]

We analyzed both numerically and (in some limiting cases) analytically these equations. The analysis of Eq. (12d) for \(\mu_s\) shows that it is suppressed as \(\mu_s \sim \mathcal{O} \left( \frac{\Delta^2}{\mu} \right)\) or \(\mu_s \sim \mathcal{O} \left( \frac{x^2}{\mu} \right)\) that allows to put \(\mu_s = 0\) in Eqs. (10), (12a), (12b), and (12c).

It is expected that the typical values of the quark chemical potential \(\mu\) in compact stars should be in the window 300 – 500 MeV. In the numerical analysis of equations (12), we take \(\mu = 400\) MeV and \(\Lambda = 653.3\) MeV \(18\) [as will be discussed below, the results are not sensitive to the choice of the value of \(\mu\) in the interval 300 – 500 MeV]. The NJL coupling constant \(G_\Delta\) determines physical properties of the model. In practice, it is more convenient to trade \(G_\Delta\) for \(\Delta_0\), which is the value of the 2SC gap at \(\delta \mu = 0\). Each of Eqs. (12a), (12b), and (12c) with \(\mu_s = 0\) determines a surface in three dimensional space with coordinates \(\Delta, q\), and \(\mu_c = 2\delta \mu\). The intersection of the neutrality surface, determined by Eq. (12c) and the surface determined by gap equation (12c) yields a neutrality curve. Its projections on the \((\mu_c, \Delta)\) and \((\mu_c, q)\) planes are shown in Figs. (a) and (b), respectively. For a fixed \(\Delta_0\), the solution for a neutral LOFF state is depicted by a black bold point which is an intersection of the neutrality curve and the “gap” curve determined by two gap equations (12a) and (12b). In Fig it is shown the neutrality curve of the LOFF state (bold line) and three gap curves (thin lines) for three characteristic values of \(\Delta_0\): 40 MeV (weak coupling), 110 MeV (intermediate coupling), and 190 MeV (strong coupling). The neutral LOFF solution occurs only for the intermediate coupling, corresponding to the gapless g2SC solution.

The analysis of Eqs. (12a)-(12c) shows that the neutral LOFF solution exists in the window

\[
63 \text{ MeV} < \Delta_0 < 137 \text{ MeV} \quad (\text{LOFF window}),
\]

which can be compared with the g2SC window

\[
92 \text{ MeV} < \Delta_0 < 130 \text{ MeV} \quad (\text{g2SC window}).
\]

We emphasize that for the same value \(\Delta_0\) (i.e., the same value of the coupling \(G_\Delta\)), the values of the parameters \(\Delta\) and \(\delta \mu\) are different for the neutral 2SC/g2SC and neutral LOFF phases.

In the LOFF window \(13\), we find that \(\Delta/\Delta_0 = 0.83\). The calculated free energies in different phases are shown in Fig. 2 (the free energy of the normal phase was chosen as a reference point) \(13\). One can see that the neutral LOFF phase is more stable than the neutral normal phase in the whole LOFF window \(13\) where these phases coexist. It is also more stable than the neutral g2SC/2SC phase in the whole g2SC window \(14\) plus
the narrow region $130 \text{ MeV} < \Delta_0 < 136 \text{ MeV}$ in the 2SC phase near its edge with the g2SC one. However, because the chromomagnetic instability in the 2SC phase occurs at $\delta \mu = \Delta / \sqrt{2}$, which corresponds to $\Delta_0 = 177 \text{ MeV}$, the neutral LOFF solution cannot cure this instability in a wide region, with $136 \text{ MeV} < \Delta_0 < 177 \text{ MeV}$, in that phase.

It is however not the end of the story. Combining Eq. (12) with Eq. (68) for Meissner masses in the second paper in Ref. [10], we found that the Meissner masses for the 4–7-th gluons in neutral LOFF state are imaginary for all values of the parameter $\Delta_0 > \Delta_0^{cr} = 81 \text{ MeV}$, see Fig. 3 (the Meissner mass of the 8-th gluon is zero [10]). In other words, for these values of $\Delta_0$, the neutral LOFF state itself suffers from a chromomagnetic instability. Because the values of $\Delta_0$ both in the 2SC phase and in the g2SC window [15] are larger than this critical value $\Delta_0^{cr} = 81 \text{ MeV}$ (see Fig. 2), we conclude that in the whole region where the neutral 2SC/g2SC and LOFF phases coexist, all these three phases are unstable. Therefore, the neutral LOFF state with a single plane wave cannot cure the chromomagnetic instability.

We checked that this conclusion is not sensitive to the choice of the value of the quark chemical potential in the interval $300 \text{ MeV} < \mu < 500 \text{ MeV}$. In particular, we found that for $\mu = 300 \text{ MeV}$, the LOFF window is $49 \text{ MeV} < \Delta_0 < 121 \text{ MeV}$ and the g2SC window is $72 \text{ MeV} < \Delta_0 < 115 \text{ MeV}$. The neutral LOFF state has the chromomagnetic instability for $\Delta_0 > \Delta_0^{cr} = 64 \text{ MeV}$ in this case. As to $\mu = 500 \text{ MeV}$, the LOFF and g2SC windows are $74 \text{ MeV} < \Delta_0 < 141 \text{ MeV}$ and $109 \text{ MeV} < \Delta_0 < 134 \text{ MeV}$, respectively. The neutral LOFF state has the chromomagnetic instability for $\Delta_0 > \Delta_0^{cr} = 94 \text{ MeV}$. Thus, the conclusions of the present analysis seem to be robust.

As to the reliability of the mean field (one-loop) approximation we use in this problem, Fig. 1 (a) shows that the values of the ratio $\Delta / \mu$ in the neutral LOFF state are typically less than one fourth. This suggests that the mean field approximation should be at least qualitatively reliable. Still, it would be certainly important to clarify more explicitly the role of fluctuations in this problem.

Because of the relation between the LOFF vector $\vec{q}$ and the VEV of $\vec{A}^{(8)}$ in Eq. (8), the chromomagnetic instability in the neutral LOFF state implies that the vector condensate $\langle \vec{A}^{(8)} \rangle$ alone is not enough for washing out the chromomagnetic instability. We consider this as a strong indication on the relevance of the gluonic phase [15] for curing this problem. Indeed, in the gluonic phase there are additional vector condensates directly connected with the 4–7-th gluons and, as was shown in Ref. [15], there is
FIG. 3: Meissner mass square $M^2$ for the 4-7-th gluons in the neutral LOFF state. The QCD coupling constant $\alpha_s = g^2/(4\pi)$ is chosen to be equal to 1.

no chromomagnetic instability in that phase in the strong coupling regime close to the value $\Delta_0 = 177$ MeV corresponding to the end point $\delta\mu = \Delta/\sqrt{2}$ of this instability in the 2SC solution.

This makes the study of the gluonic phase both far away from the edge $\Delta_0 = 177$ MeV of the chromomagnetic instability in the 2SC solution and in the g2SC window to be quite desirable. Also, it is clear that it would be interesting to study the neutral LOFF state in the three-flavor QCD \cite{21} and to examine the relevance of the LOFF state with more than one plane waves \cite{21} for curing the chromomagnetic instability.

The work was supported by the Natural Sciences and Engineering Research Council of Canada.

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\[18\] While the $\vec{q}$ dependent part of the free energy \cite{16} coincides with that in the second paper in Ref. \cite{16}, the 2SC parts of these two expressions are essentially different. As a result, the free energy in Ref. \cite{16} does not reproduce correctly the limits with $q = 0$ and with $q = 0$, $\Delta = 0$. Because of that, it is not appropriate for studying a neutral LOFF state.

\[19\] In the part of LOFF window \cite{18} with $130$ MeV $< \Delta_0 < 137$ MeV, there exists a second branch of the neutral LOFF solution. By using expression \cite{18} one can show that its free energy is larger than the free energy of the first branch. This energetically unfavorable branch is not shown in Fig.\textit{2}.

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