COLOR SUPERCONDUCTING NEUTRAL MATTER

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We describe the effects of the strange quark mass and of the color and electric neutrality on the superconducting phases of QCD.

1 Introduction

It is now a well established fact that at zero temperature and sufficiently high densities quark matter is a color superconductor. The study starting from first principles was done in Ref. [1][2]. At baryon chemical potentials much higher than the masses of the quarks u, d and s, the favored state is the so-called Color-Flavor-Locking (CFL) state, whereas at lower values, when the strange quark decouples, the relevant phase is called two-flavor color superconducting (2SC).

An interesting possibility is that in the interior of compact stellar objects (CSO) some color superconducting phase may exist. In fact the central densities for these stars could be up to $10^{15}$ g/cm$^3$, whereas the temperature is of the order of tens of keV. However the usual assumptions leading to prove that for three flavors the favored state is CFL should now be reviewed. Matter inside a CSO should be electrically neutral and should not carry any color. Also conditions for $\beta$-equilibrium should be fulfilled. As far as color is concerned, it is possible to impose a simpler condition, that is color neutrality, since in Ref. [5] it has been shown that there is no free energy cost in projecting color singlet states out of color neutral states. Furthermore one has to take into account that at the interesting density the mass of the strange quark is a relevant parameter. All these effects, the mass of the strange quark, $\beta$-equilibrium and color and electric neutrality, imply that the radii of the Fermi spheres of the quarks that would pair are not the same creating a problem with the usual BCS pairing. Let us start from the first point. Suppose to have two fermions of masses $m_1 = M$ and $m_2 = 0$ at the same chemical potential $\mu$. The corresponding Fermi momenta are $p_{F_1} = \sqrt{\mu^2 - M^2}$ and $p_{F_2} = \mu$. Therefore the radius of the Fermi sphere of the massive fermion is smaller than the one of the massless particle. If we assume $M \ll \mu$ the massive particle has an effective chemical potential $\mu_{\text{eff}} = \sqrt{\mu^2 - M^2} \approx \mu - M^2/2\mu$ and the mismatch between the two Fermi spheres is given by

$$\delta \mu \approx \frac{M^2}{2\mu} \tag{1}$$

This shows that the quantity $M^2/(2\mu)$ behaves as a chemical potential. Therefore for $M \ll \mu$ the mass effects can be taken into account through the introduction of the mismatch between the chemical potentials of the two fermions given by eq. (1). This is the way that we will follow in our study.

Now let us discuss $\beta$-equilibrium. If electrons are present (as generally required by electrical neutrality) chemical potentials of quarks of different electric charge are differ-
ent. In fact, when at the equilibrium for \( d \to u e \bar{e} \), we have
\[
\mu_d - \mu_u = \mu_e \quad (2)
\]
From this condition it follows that for a quark of charge \( Q_i \) the chemical potential \( \mu_i \) is given by
\[
\mu_i = \mu + Q_i \mu_Q \quad (3)
\]
where \( \mu_Q \) is the chemical potential associated to the electric charge. Therefore
\[
\mu_e = -\mu_Q \quad (4)
\]
Notice also that \( \mu_e \) is not a free parameter since it is determined by the neutrality condition
\[
Q = -\frac{\partial \Omega}{\partial \mu_e} = 0 \quad (5)
\]
At the same time the chemical potentials associated to the color generators \( T_3 \) and \( T_8 \) are determined by the color neutrality conditions
\[
\frac{\partial \Omega}{\partial \mu_3} = \frac{\partial \Omega}{\partial \mu_8} = 0 \quad (6)
\]
We see that in general there is a mismatch between the quarks that should pair according to the BCS mechanism for \( \delta \mu = 0 \). Therefore the system might go to a normal phase since the mismatch, as we shall see, tends to destroy the BCS pairing, or a different phase might be formed. In the next Sections we will explore some of these possible phases.

2 Pairing Fermions with Different Fermi Momenta

We start now our discussion considering a simple model with two pairing quarks, \( u \) and \( d \), with chemical potentials
\[
\mu_u = \mu + \delta \mu, \quad \mu_d = \mu - \delta \mu \quad (7)
\]
and no further constraints. The gap equation for the LOFF phase at \( T = 0 \) is given by (see for example ref. [7])
\[
1 = \frac{g}{2} \int \frac{d^3 p}{(2\pi)^3} \frac{1}{\epsilon(\vec{p}, \Delta)} (1 - \theta(-\epsilon - \delta \mu) - \theta(-\epsilon + \delta \mu)) \quad (8)
\]
where
\[
\epsilon(\vec{p}, \Delta) = \sqrt{\xi^2 + \Delta^2}, \quad \xi = \vec{v}_F \cdot (\vec{p} - \vec{p}_F) \quad (9)
\]
with \( \vec{v}_F \) and \( \vec{p}_F \) the Fermi velocity and Fermi momentum. The meaning of the two step functions is that at zero temperature there is no pairing when \( \epsilon(\vec{p}, \Delta) < |\delta \mu| \). In other words the pairing may happen only for excitations with positive energy. However, the presence of negative energy states, as in this case, implies that there must be gapless modes. When this happens, there are blocking regions in the phase space, that is regions where the pairing cannot occur. The effect is to inhibit part of the Fermi surface to the
pairing giving rise a to a smaller condensate with respect to the BCS case where all the surface is used. In the actual case the gap equation at $T = 0$ has two different solutions (see for instance ref. [7]) corresponding to: a) $\Delta = \Delta_0$, and b) $\Delta^2 = 2\delta \mu \Delta_0 - \Delta_0^2$ where $\Delta_0$ is the BCS solution of the gap equation for $\delta \mu = 0$. The two solutions are illustrated in Fig. 1.

We see that the solution a) is always favored with respect to the solution b) (called the Sarma phase[8]). Furthermore the BCS phase goes to the normal phase at

$$\delta \mu_1 = \frac{\Delta_0}{\sqrt{2}}$$

(10)

This point is called the Chandrasekhar-Clogston (CC) point[9] (denoted by CC in Fig. 1). Ignoring for the moment that in this case, after the CC point the system goes to the normal phase, we notice that the gaps of the two solutions coincide at $\delta \mu = \Delta_0$. This is a special point, since in presence of a mismatch the spectrum of the quasi-particles is modified from

$$E_{\delta \mu = 0} = \sqrt{(p - \mu)^2 + \Delta^2} \rightarrow E_{\delta \mu} = \left|\delta \mu \pm \sqrt{(p - \mu)^2 + \Delta^2}\right|$$

(11)

Therefore for $|\delta \mu| < \Delta$ we have gapped quasi-particles with gaps $\Delta \pm \delta \mu$. However, for $|\delta \mu| = \Delta$ a gapless mode appears and from this point on there are regions of the phase space which do not contribute to the gap equation (blocking region).

The gapless modes are characterized by

$$E(p) = 0 \Rightarrow p = \mu \pm \sqrt{\delta \mu^2 - \Delta^2}$$

(12)

Since the energy cost for pairing two fermions belonging to Fermi spheres with mismatch $\delta \mu$ is $2\delta \mu$ and the energy gained in pairing is $2\Delta$, we see that the fermions begin to unpair for $2\delta \mu \geq 2\Delta$. These considerations will be relevant for the study of the gapless phases when neutrality is required.
3 The g2SC Phase

The g2SC phase\[10\] has the same condensate as the 2SC

\[\langle 0 | \psi^\alpha_{aL} \psi^\beta_{bL} | 0 \rangle = \Delta \epsilon^{\alpha \beta 3}_{ab3}, \quad \alpha, \beta \in SU_c(3), \quad a, b \in SU(2)_L\]  

and, technically, it is distinguished by the 2SC one by the presence of gapless modes starting at \(\delta \mu = \Delta\). In this case only two massless flavors are present (quarks \(u\) and \(d\)) and there are 2 quarks ungapped \(q_{udr}, q_{db}\) and 4 gapped \(q_{u}, q_{d}, q_{g}, q_{g}\), where the color indices 1, 2, 3 have been identified with \(r, g, b\) (red, green and blue). The difference with the usual 2SC phase is that color and electrical neutrality are required:

\[\frac{\partial \Omega}{\partial \mu_e} = \frac{\partial \Omega}{\partial \mu_3} = \frac{\partial \Omega}{\partial \mu_8} = 0\]  

This creates a mismatch between the two Fermi spheres given by \(\delta \mu = \mu_e/2\). Furthermore one has to satisfy the gap equation. One finds two branches of solutions of the gap equation corresponding to the BCS phase and to the Sarma phase. It turns out that the solution to the present problem belongs to the Sarma branch. In\[10\] it is also shown that the solution is a minimum of the free energy following the neutrality line. On the other hand this point is a maximum following the appropriate line \(\mu_e = \text{const.}\). We see that the neutrality conditions promote the unstable phase (Sarma) to a stable one. However this phase has an instability connected to the Meissner mass of the gluons\[11\]. In this phase the color group \(SU_c(3)\) is spontaneously broken to \(SU_c(2)\) with 5 of the 8 gluons acquiring a mass; precisely the gluons 4, 5, 6, 7, 8. At the point \(\delta \mu = \Delta\) where the 2SC phase goes into the g2SC one, all the massive gluons have imaginary mass. Furthermore the gluons 4, 5, 6, 7 have imaginary mass already starting at \(\delta \mu = \Delta/\sqrt{2}\), that is at the Chandrasekhar-Clogston point. This shows that both the g2SC and the 2SC phases are unstable. The instability of the g2SC phase seems to be a general feature of the phases with gapless modes\[12\].

4 The gCFL phase

The gCFL phase is a generalization of the CFL phase which has been studied both at \(T = 0\)\[13,14\] and \(T \neq 0\)\[15\]. The condensate has now the form

\[\langle 0 | \psi^\alpha_{aL} \psi^\beta_{bL} | 0 \rangle = \Delta_1 \epsilon^{\alpha \beta 1}_{ab1} + \Delta_2 \epsilon^{\alpha \beta 2}_{ab2} + \Delta_3 \epsilon^{\alpha \beta 3}_{ab3}\]  

The CFL phase corresponds to all the three gaps \(\Delta_i\) being equal. Varying the gaps one gets many different phases. In particular we will be interested to the CFL, to the g2SC characterized by \(\Delta_3 \neq 0\) and \(\Delta_1 = \Delta_2 = 0\) and to the gCFL phase with \(\Delta_3 > \Delta_2 > \Delta_1\). Notice that here in the g2SC phase the strange quark is present but unpaired.

In flavor space the gaps \(\Delta_i\) correspond to the following pairings

\[\Delta_1 \Rightarrow ds, \quad \Delta_2 \Rightarrow us, \quad \Delta_3 \Rightarrow ud\]  

(16)
The mass of the strange quark is taken into account by shifting all the chemical potentials involving the strange quark as follows: \( \mu_{as} \rightarrow \mu_{as} - \frac{M_s^2}{2\mu} \). It has also been shown in ref. 16 that color and electric neutrality in CFL require

\[
\mu_8 = -\frac{M_s^2}{2\mu}, \quad \mu_e = \mu_3 = 0
\]  

(17)

At the same time the various mismatches are given by

\[
\delta\mu_{bd-gs} = \frac{M_s^2}{2\mu}, \quad \delta\mu_{rd-gu} = \mu_e = 0, \quad \delta\mu_{rs-bu} = \mu_e - \frac{M_s^2}{2\mu}
\]  

(18)

It turns out that in the gCFL the electron density is different from zero and, as a consequence, the mismatch between the quarks \( d \) and \( s \) is the first one to give rise to the unpairing of the corresponding quarks. This unpairing is expected to occur for

\[
2 \frac{M_s^2}{2\mu} > 2\Delta \Rightarrow \frac{M_s^2}{\mu} > 2\Delta
\]  

(19)

This has been substantiated by the calculations in a NJL model modeled on one gluon-exchange in 14. The transition from the CFL phase, where all gaps are equal, to the gapless phase occurs roughly at \( \frac{M_s^2}{\mu} = 2\Delta \). In Fig. 2 we show the free energy of

![Free energy graph](image)

Figure 2: We give here the free energy of the various phases with reference to the normal phase, named unpaired in the figure.

the various phases with reference to the normal phase. The CFL phase is the stable one up to \( \frac{M_s^2}{\mu} \approx 2\Delta \). Then the gCFL phase takes over up to about 130 MeV where the system goes to the normal phase. Notice that except in a very tiny region around this point, the CFL and gCFL phases win over the corresponding 2SC and g2SC ones. The thin short-dashed line represents the free energy of the CFL phase up to the point where it becomes equal to the free-energy of the normal phase. This happens for \( \frac{M_s^2}{\mu} \approx 4\Delta \). This point is the analogue of the Chandrasekhar-Clogston point of the two-flavor case.

Also the gCFL phase has gapless excitations and as a consequence the chromomagnetic instability discussed in the case of the g2SC phase shows up here too. This has been shown in 17,18.
The existence of the chromomagnetic instability is a serious problem for the gapless phases (g2SC and gCFL) but also for the 2SC phase, as we have discussed previously. A way out of this problem would be to have gluon condensation. For instance, if one assumes artificially \( \langle A_3^\mu \rangle \) and \( \langle A_8^\mu \rangle \) not zero and with a value of about 10 MeV it can be shown that the instability disappears\(^{17}\). Also, very recently in\(^{19}\), it has been shown the possibility of eliminating the chromomagnetic instability in the 2SC phase through a gluonic phase. However it is not clear if the same method can be extended to the gapless phases.

Another interesting possibility has been considered in three papers by Giannakis and Ren, who have considered the LOFF phase, that is a nonhomogeneous phase first studied in a condensed matter context\(^{20,21}\) and then in QCD in\(^{22,23}\) (for recent reviews of the LOFF phase, see\(^{7,24}\)). The results obtained by Giannakis and Ren in the two-flavor case are the following:

- The presence of the chromomagnetic instability in g2SC is exactly what one needs in order the LOFF phase to be energetically favored\(^{25}\).
- The LOFF phase in the two-flavor case has no chromomagnetic instabilities (though it has gapless modes) at least in the weak coupling limit\(^{26,27}\).

Of course these results make the LOFF phase a natural candidate for the stable phase of QCD at moderate densities. In the next Sections we will describe the LOFF phase in its simplest version and then a simplified approach to the problem with three flavors will be presented.

### 5 The LOFF Phase

According to the authors of refs.\(^{20,21}\) when fermions belong to different Fermi spheres, they might prefer to pair staying as much as possible close to their own Fermi surface. The total momentum of the pair is not zero, \( \vec{p}_1 + \vec{p}_2 = 2 \vec{q} \) and, as we shall show, \(|\vec{q}|\) is fixed variationally whereas the direction of \( \vec{q} \) is chosen spontaneously. Since the total momentum of the pair is not zero the condensate breaks rotational and translational invariance. The simplest form of the condensate compatible with this breaking is just a simple plane wave (more complicated possibilities will be discussed later)

\[
\langle \psi(x)\psi(x) \rangle \approx \Delta e^{2i\vec{q}\cdot\vec{x}} \tag{20}
\]

It should also be noticed that the pairs use much less of the Fermi surface than they do in the BCS case. For instance, if both fermions are sitting at their own Fermi surface, they can pair only if they belong to circles fixed by \( \vec{q} \). More generally there is a quite large region in momentum space (the so called blocking region) which is excluded from pairing. This leads to a condensate generally smaller than the BCS one.

Let us now consider in more detail the LOFF phase. For two fermions at different densities we have an extra term in the Hamiltonian which can be written as

\[
H_I = -\delta \mu \sigma_3 \tag{21}
\]

where, in the original LOFF papers\(^{20,21}\), \( \delta \mu \) is proportional to the magnetic field due to the impurities, whereas in the actual case \( \delta \mu = (\mu_1 - \mu_2)/2 \) and \( \sigma_3 \) is a Pauli matrix.
acting on the two fermion space. According to refs. [20][21] this favors the formation of pairs with momenta
\[ \vec{p}_1 = \vec{k} + \vec{q}, \quad \vec{p}_2 = -\vec{k} + \vec{q} \] (22)
We will discuss in detail the case of a single plane wave (see eq. (20)). The interaction term of eq. (21) gives rise to a shift in \( \xi \) due both to the non-zero momentum of the pair and to the different chemical potentials
\[ \xi = E(p) - \mu \rightarrow E(\pm \vec{k} + \vec{q}) - \mu \mp \delta \mu \approx \xi \mp \bar{\mu} \] (23)
with
\[ \bar{\mu} = \delta \mu - \vec{v}_F \cdot \vec{q} \] (24)
Notice that the previous dispersion relations show the presence of gapless modes at momenta depending on the angle of \( \vec{v}_F \) with \( \vec{q} \). Here we have assumed \( \delta \mu \ll \mu \) (with \( \mu = (\mu_1 + \mu_2)/2 \)) allowing us to expand \( E \) at the first order in \( \vec{q}/\mu \).

The gap equation for the present case is obtained simply from eq. (8) via the substitution \( \delta \mu \rightarrow \bar{\mu} \). By studying this equation one can show that increasing \( \delta \mu \) from zero we have first the BCS phase. Then at \( \delta \mu = \delta \mu_1 \) there is a first order transition to the LOFF phase [20][22]. And at \( \delta \mu = \delta \mu_2 > \delta \mu_1 \) there is a second order phase transition to the normal phase [20][22]. We start comparing the grand potential in the BCS phase to the one in the normal phase. Their difference is given by (see for example ref. [7])
\[ \Omega_{\text{BCS}} - \Omega_{\text{normal}} = -\frac{\rho_F^2}{4\pi^2v_F} \left( \Delta_0^2 - 2\delta \mu^2 \right) \] (25)
where the first term comes from the energy necessary to the BCS condensation, whereas the last term arises from the grand potential of two free fermions with different chemical potential. We recall also that for massless fermions \( p_F = \mu \) and \( v_F = 1 \). We have again assumed \( \delta \mu \ll \mu \). This implies that there should be a first order phase transition from the BCS to the normal phase at \( \delta \mu = \Delta_0/\sqrt{2} \), since the BCS gap does not depend on \( \delta \mu \). In order to compare with the LOFF phase one can expand the gap equation around the point \( \Delta = 0 \) (Ginzburg-Landau expansion) to explore the possibility of a second order phase transition [20][22]. The result for the free energy is
\[ \Omega_{\text{LOFF}} - \Omega_{\text{normal}} \approx -0.44 \rho(\delta \mu - \delta \mu_2)^2 \] (26)
At the same time, looking at the minimum in \( q \) of the free energy one finds
\[ qv_F \approx 1.2 \delta \mu \] (27)
Since we are expanding in \( \Delta \), in order to get this result it is enough to minimize the coefficient of \( \Delta^2 \) in the free-energy (the first term in the Ginzburg-Landau expansion).

We see that in the window between the intersection of the BCS curve and the LOFF curve and \( \delta \mu_2 \), the LOFF phase is favored. Also at the intersection there is a first order transition between the LOFF and the BCS phase. Furthermore, since \( \delta \mu_2 \) is very close to \( \delta \mu_1 \) the intersection point is practically given by \( \delta \mu_1 \). The window of existence of the LOFF phase \( (\delta \mu_1, \delta \mu_2) \approx (0.707, 0.754)\Delta_0 \) is rather narrow, but there are indications that considering the realistic case of QCD [28] the window opens up. Such opening occurs also for different structures than the single plane wave [23][30].
6 The LOFF phase with three flavors

In the last Section we would like to illustrate some preliminary result about the LOFF phase with three flavors. This problem has been considered in\cite{29} under various simplifying hypothesis:

- The study has been made in the Ginzburg-Landau approximation.
- Only electrical neutrality has been required and the chemical potentials for the color charges $T_3$ and $T_8$ have been put equal to zero (see later).
- The mass of the strange quark has been introduced as it was done previously for the gCFL phase.
- The study has been restricted to plane waves, assuming the following generalization of the gCFL case:

$$\langle \psi_\alpha^a \psi_\beta^b \rangle = \sum_{I=1}^3 \Delta_I(\vec{x}) \epsilon^{\alpha\beta I} \epsilon_{abI}, \quad \Delta_I(\vec{x}) = \Delta_I e^{2i\vec{q}_I \cdot \vec{x}} \quad (28)$$

- The condensate depends on three momenta, meaning three lengths of the momenta $q_i$ and three angles. In\cite{29} only four particular geometries have been considered: 1) all the momenta parallel, 2) $\vec{q}_1$ antiparallel to $\vec{q}_2$ and $\vec{q}_3$, 3) $\vec{q}_2$ antiparallel to $\vec{q}_1$ and $\vec{q}_3$, 4) $\vec{q}_3$ antiparallel to $\vec{q}_1$ and $\vec{q}_2$.

The minimization of the free energy with respect to the $|q_I|$’s leads to the same result as in eq. (27), $|q_I| = 1.2 \delta \mu_I$. Then, one has to minimize with respect to the gaps and $\mu_e$ in order to require electrical neutrality. It turns out that the configurations 3 and 4 have an extremely small gap. Furthermore for $M_s^2/\mu$ greater than about $80 \, MeV$ the solution gives $\Delta_1 = 0$ and $\Delta_2 = \Delta_3$. In this case the configurations 1 and 2 have the same free energy. The results for the free energy and for the gap of this solution are given in Fig. 4. In this study, the following choice of the parameters has been made: the BCS gap, $\Delta_0 = 25 \, MeV$, and the chemical potential $\mu = 500 \, MeV$. The values are the same discussed previously for gCFL in order to allow for a comparison of the results. We are now in the position to compare these results with the ones obtained in\cite{14} for the gCFL phase. The comparison is made in Fig. 4. Ignoring the chromomagnetic instabilities of the gapless phases and of 2SC we see that LOFF takes over with respect to gCFL at about $M_s^2/\mu = 128 \, MeV$ and goes over to the normal phase for $M_s^2/\mu \approx 150 \, MeV$. However, since the instability exists, it should be cured in some way. The results for the LOFF phase, assuming that also for three flavors the chromomagnetic instability does not show up, say that the LOFF phase could take over the CFL phase before the transition to gCFL. For this it is necessary that the window for the LOFF phase gets enlarged. However, in\cite{30} it has been show that for structures more general than the plane wave the windows may indeed becomes larger. If we define the window for the single plane wave as $(\delta \mu_2 - \delta \mu_1)/\delta \mu_2$ (see the previous Section) we would get 0.06. The analogous ratio in going from one to three plane waves goes to about $(150 - 115)/150 = .23$, with a gain of almost a factor 4. On the other hand, in\cite{22} it has been shown that considering some of the crystalline structures already taken in exam in\cite{22}, as the face centered cube or the cube
the window becomes $(1.32 - 0.707)/1.32 = 0.46$ with a gain of about 7.7 with respect to the single plane wave. If these gains would be maintained in going from two to three flavors with the face centered cube structure, one could expect a gain from 4 to 7.7 with an enlargement of the window between 88 and 170 $MeV$, which would be enough to cover the region of gCFL (which is about 70 $MeV$). At last we want to comment about the approximation in neglecting the color neutrality condition and assuming $\mu_3 = \mu_8 = 0$. The results of ref.\cite{ref29} show that $\mu_e \approx M_2^2/(4\mu)$ as for the case of 3 color and 3 flavor unpaired quarks\cite{ref16}. Furthermore the unpaired quarks have also $\mu_3 = \mu_8 = 0$. Also, from Fig.\ref{fig:fig3} we see that in our approximations the transition from the LOFF to the normal phase is very close to be continuous. Since we expect the chemical potentials to be continuous at the transition point, we expect $\mu_3 = \mu_8 = 0$ also on the LOFF side, at least when close to the critical point. This means the color neutrality condition should be $\mu_3 = \mu_8 = 0$ in the neighborhood of the transition. Therefore we expect the determination of the point $M_2^2/\mu = 150$ $MeV$ to be safe. On the other hand, the requirement of color neutrality
could change the intersection point with gCFL. Nevertheless, since the critical point for LOFF is higher than the one of gCFL, for increasing $M_s$ the system must to go into the LOFF phase.

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