Color Superconductivity in Asymmetric Matter

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Abstract

The influence of different chemical potential for different flavors on color superconductivity is analyzed. It is found that there is a first order transition as the asymmetry grows. This transition proceeds through the formation of bubbles of low density, flavor asymmetric normal phase inside a high density, superconducting phase with a gap larger than the one found in the symmetric case. For small fixed asymmetries the system is normal at low densities and superconducting only above some critical density. For larger asymmetries the two massless quarks system stays in the mixed state for arbitrarily high densities.
A lot of progress has been made recently in mapping the phase diagram of QCD in the low temperature, high baryon density region. The main idea advanced by this research is the almost unavoidable formation of quark-quark condensates (pairing) and the spontaneous breaking (or restoration) of color-flavor symmetries of QCD \[1,2\]. These ideas may have an impact on the phenomenology of neutron stars and of heavy ion collisions at SPS energies where compressed strong interacting matter is believed to exist.

Since lattice calculations at finite density are unfeasible at this date two alternative approaches are used to study this problem. One, valid at asymptotically high densities, is the use of perturbation theory. The gaps due to perturbative QCD were believed to be small \[2\], exponentially suppressed as in weak coupling BCS theory. It was recently realized \([3–9]\), however, that the main pairing interaction in this regime is due to the exchange of magnetic gluons. These interactions make the normal phase even more unstable than short range forces do and result on a parametrically much larger gap. Extrapolating this down to more accessible densities one could expect gaps up to 100 MeV at chemical potentials $\mu \sim 1000$ MeV. It remains to be seen at which density the sub-leading correction (in powers of the coupling) become large and invalidate the perturbative expansion. Another approach to the problem is the use of phenomenological models such as NJL-like models and the instanton liquid model \([10–17]\). In particular, a large class of four-fermion interaction models was considered. Fortunately, for semi-quantitative purposes, they tend to give very similar results. This is due to the fact that the formation of the condensate depends only on the most attractive channel between quarks at the Fermi surface. Different four-fermion interactions that are both attractive in the scalar-isoscalar color $\bar{3}$ channel and of roughly the same strength produce similar BCS instabilities. Renormalization group analysis support this conclusion and help in classifying all four-fermion operators according to their importance to the formation of the gap \([18,19]\).

With the help of these phenomenological models the most likely scenarios of symmetry breaking were identified \([1,2,10,13]\). In the case of 2 massless quarks chiral symmetry is likely restored at some critical density and color $SU_c(3)$ is partially broken down to $SU_c(2)$ by the formation of the condensate (2SC)

$$\langle q_{\alpha a}^{\alpha} q_{\beta b}^{\beta} \rangle_2 = \langle q_{\beta b}^{\alpha} q_{\alpha a}^{\beta} \rangle_2 \sim \epsilon^{\alpha \beta \lambda} \epsilon^{\alpha \beta \lambda}$$

where the $\alpha, \beta, \ldots$ are color indices, $a, b, \ldots$ are flavor indices, $i, j, \ldots$ are spinor indices and $q_{L,R}$ are left (right) handed spinors. The condensate in eq. (1) singles out the third color (green from now on), only blue and red quarks pair up.

In the three massless quarks case the symmetry breaking pattern is more subtle \([20,21]\). The preferred condensate is the so-called color-flavor locked (CFL) state

$$\langle q_{\alpha a}^{\alpha} q_{\beta b}^{\beta} \rangle_2 = -\langle q_{\alpha a}^{\alpha} q_{\beta b}^{\beta} \rangle_2 \sim \epsilon_{ij} (\kappa_1 \delta_{\alpha a} \delta_{\beta b} + \kappa_2 \delta_{\alpha b} \delta_{\beta a})$$

The condensate eq. (2) breaks $SU_c(3) \times SU_L(3) \times SU_R(3)$ down to $SU_{c+L+R}(3)$, leaving the ground state invariant under a combined color-flavor rotation. The remaining symmetry $SU_{c+L+R}(3)$ is the same as the one of the vacuum after spontaneous breaking of chiral symmetry ($SU_V(3)$). This suggests that there is smooth connection between the hadronic phase to the CFL phase. The spectrum of excitations is identical in both cases: the gluon octet is mapped onto the vector meson octet, the quarks onto the baryon octet and the 8 Goldstone bosons onto the pseudo-scalar mesons.
The real world lies between the 2 and 3 massless quark case since the mass \( m_s \) of the strange quark is neither much larger nor much smaller than the QCD scale. At arbitrarily high densities the presence of the strange quark mass is irrelevant but not at lower densities. It is difficult then to determine whether the CFL phase extends all the way to lower densities and connects to the hadronic phase or whether there is a window in density with the 2SC phase. Another difference between the discussion sketched above and the potential real world application is the difference in densities among the different quark flavors. In heavy ion collisions there is an excess of \( \approx 15\% \) in the number of down quarks (minus down antiquarks) relative to the up quarks, and zero net strangeness at all. There are also significant asymmetries in neutron star cores.

Here we will analyze the two massless quarks phase in a model that conserves quark flavors. We choose our interaction to be a contact four fermion coupling with spin, flavor and color structure abstracted from the one gluons exchange projected on the color \( \bar{3} \) channel (the 6 is repulsive). Our starting point is then the Lagrangian

\[
L - \mu_a N_a = \psi_{\alpha a}^\dagger (i \sigma^\mu \partial_\mu + \mu) \delta_{ab} + \delta \tau^3_{ab} \delta_{\alpha \beta} \psi_{\beta b} - g^2 \psi_{\alpha a}^T \sigma_2 \psi_{\beta b}^\dagger \sigma_2 \psi_{\beta b}^* (\delta_{\alpha \gamma} \delta_{\beta \delta} - \delta_{\alpha \delta} \delta_{\beta \gamma}), \tag{3}
\]

where \( \mu = (\mu_u + \mu_d)/2 \) is the chemical potential averaged over flavors (equal to one third of the baryon chemical potential), \( \delta = (\mu_u - \mu_d)/2 \) is their difference, \( N_a \) is the operator that counts the total amount of flavor \( a \) (quarks minus antiquarks) and the spinor indices were omitted. The interaction in eq. (3) does not mix left and right handed quarks and a sum over handedness is implicit throughout the paper. The matrices \( \sigma^\mu \) are defined as \( \sigma^\mu = (1, -\vec{\sigma}) \) for the left handed quarks and \( \sigma^\mu = (1, \vec{\sigma}) \) for the right handed ones. The mean field approximation, although not arising as a systematic expansion in a small parameter is usually reliable to determine the phase structure of the theory. It is exact in the limit of large number of flavors (keeping \( g^2 N_F \) constant) and on the weak coupling limit. To perform the mean field approximation we rewrite the model defined by eq. (3) by introducing a dummy field \( \Delta_{\alpha a \beta b} \)

\[
L + \mu_a N_a = \psi_{\alpha a}^\dagger (i \sigma^\mu \partial_\mu + \mu) \delta_{ab} + \delta \tau^3_{ab} \delta_{\alpha \beta} \psi_{\beta b} + \frac{1}{g^2} \Delta_{\alpha a \beta b} \Delta^*_{\gamma \delta b a} \epsilon_{\alpha \beta \gamma \delta}, \tag{4}
\]

where \( \epsilon_{\alpha \beta \gamma \delta} = \delta_{\alpha \gamma} \delta_{\beta \delta} - \delta_{\alpha \delta} \delta_{\beta \gamma} \). It is convenient to use the Nambu-Gorkov formalism and write eq. (4) as

\[
L + \mu_a N_a = \left( \psi^\dagger \psi^T \sigma_2 \right) \mathcal{M} \left( \begin{array}{c} \psi \\ \sigma_2 \psi^* \end{array} \right) + \frac{1}{g^2} \Delta_{\alpha a \beta b} \Delta^*_{\gamma \delta b a} \epsilon_{\alpha \beta \gamma \delta}, \tag{5}
\]

with

\[
\mathcal{M} = \begin{pmatrix}
\mu \sigma^\mu + \mu + \delta \tau^3 \\
\Phi^* \\
\Phi \\
\Phi^* 
\end{pmatrix} \left( \begin{array}{c}
p_\mu \sigma^\mu + \mu + \delta \tau^3 \\
p_\mu \sigma^\mu + \mu + \delta \tau^3 \\
p_\mu \sigma^\mu + \mu + \delta \tau^3 \\
p_\mu \sigma^\mu + \mu + \delta \tau^3 
\end{array} \right)^{-1}, \tag{6}
\]

where \( \bar{\sigma}^\mu = (1, \bar{\sigma}) \) for the left handed quarks and \( \bar{\sigma}^\mu = (1, -\bar{\sigma}) \) for the right handed ones. The matrix \( \tau^3 \) acts on flavor space (and is diagonal in color and spin) and \( \Phi \) is the operator...
\[ \Phi_{\alpha \beta \gamma} = \Delta_{\gamma \alpha \beta} \epsilon_{\gamma \delta \alpha \beta}, \quad (7) \]

acting on color-flavor space.

The mean field approximation amounts to estimating the path integral by its saddle value point. Ignoring then the fluctuations on the \( \Delta_{\alpha \beta \gamma} \) field we are left with a Gaussian integral over the fermion fields that produces the familiar \( \det(M) = \text{Tr} \log(M) \) term. The potential \( \Omega(\Delta, \mu, \delta) \) defined by

\[ e^{-i\Omega(\Delta, \mu, \delta)} = \int D\psi D\psi^\dagger e^{i\int L + \mu a N_a} \]

is the expectation value of the Hamiltonian \( H - \mu a N_a \) in the state that minimizes \( H - \mu a N_a \) in the subspace of states with a given value of \( \Delta \). In our case

\[ \Omega(\Delta, \mu, \delta) = \frac{i}{2} \int \frac{d^4p}{(2\pi)^4} \text{tr} \log M + \frac{2}{g^2} \Delta_{\alpha \beta \gamma} \Delta_{\gamma \delta \alpha \beta}^* \epsilon_{\alpha \beta \gamma \delta}, \quad (9) \]

where the trace is over spin, color and flavor spaces. The value of \( \Delta_{\alpha \beta \gamma} \) in the ground is obtained by minimizing \( \Omega(\Delta, \mu, \delta) \) at fixed \( \mu \). We make an ansatz, corresponding to the 2SC state, for the form of the order parameter

\[ \Delta_{\alpha \beta \gamma} = \frac{\Delta}{2} \epsilon_{\alpha \beta} \epsilon_{\gamma 3}. \quad (10) \]

As mentioned before, this choice breaks color \( S_c(3) \) down to \( S_c(2) \). By using a basis in spin space that diagonalizes \( \vec{p} \vec{\sigma} \) and by doing the similarity transformation

\[ M \rightarrow U M U^{-1}, \quad U = \left( \begin{array}{cc} 1 & 0 \\ 0 & \tau_2 \end{array} \right) \]

we are left with the task of computing the determinant of

\[ M = \left( \begin{array}{cc} p_0 - \epsilon^\pm(p) + \delta \tau^3 & \Phi \\ \Phi^* & p_0 + \epsilon^\pm(p) - \delta \tau^3 \end{array} \right), \quad (12) \]

where

\[ \Phi_{\alpha \beta} = -i \Delta \epsilon_{\alpha \beta 3} \quad (13) \]

acts only in color space and \( \epsilon^\pm(p) = \pm p - \mu \). Using the relations

\[ \text{tr} \log \left( \begin{array}{cc} A & B \\ C & D \end{array} \right) = \text{Tr} \log(-BC + BDB^{-1}A), \quad (14) \]

\[ \Phi^2 = -P, \quad (15) \]

where \( P \) is the projector on the first two colors subspace, and

\[ \text{Tr}_{\text{flavor}} \log(a + b \tau^3) = \log(a + b) + \log(a - b) \quad (16) \]

\[ \text{Tr}_{\text{color}} \log(a + bP) = (N_c - 1) \log(a + b) + \log(a), \quad (17) \]
we arrive at

\[
\Omega(\Delta, \mu, \delta) = \frac{\Delta^2}{g^2} + \sum_{\pm} \frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} [2 \log((p_0 + \delta)^2 - E_{\pm}^2(p)) + \log((p_0 + \delta)^2 - \epsilon_{\pm}^2(p)) + 2 \log((p_0 - \delta)^2 - E_{\pm}^2(p)) + \log((p_0 - \delta)^2 - \epsilon_{\pm}^2(p))] + C,
\]

(18)

where \(E_{\pm}(p) = \sqrt{\Delta^2 + \epsilon_{\pm}^2(p)}\) is the energy of the quasi-particle excitations. Performing the \(p_0\) integration and fixing the constant \(C\) by demanding \(\Omega(\Delta, \mu, \delta)\) has the correct (free theory) value at \(\Delta = 0\) we arrive at

\[
\Omega(\Delta, \mu, \delta) = \frac{\Delta^2}{g^2} - \frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \sum_{\pm} \left[ 2(E_{\pm}(p) - p - \theta(\delta - E_{\pm}(p))(E_{\pm}(p) - \delta)) + |\epsilon_{\pm}(p)| - p - \theta(\delta - |\epsilon_{\pm}(p)|)|\epsilon_{\pm}(p)| - \delta \right] + 3p + \theta(\delta + \epsilon^{-}(p))(\delta + \epsilon^{-}(p)).
\]

(19)

The first line corresponds to the contribution of the red and blue quarks, the second to the unpaired green quarks and the last one to anti-particles that can be created when \(\delta < \mu\). The mean particle number density of each flavor is obtained by

\[
n_u + n_d = -\frac{\partial}{\partial \mu} \Omega(\Delta, \mu, \delta)
\]

\[
= \int \frac{d^3 p}{(2\pi)^3} \sum_{\pm} \left[ -2 \frac{\epsilon_{\pm}(p)}{E_{\pm}(p)} + 2\theta(\delta - E_{\pm}(p))\frac{\epsilon_{\pm}(p)}{E_{\pm}(p)} - \text{sgn}(\epsilon_{\pm}(p)) + \theta(\delta - |\epsilon_{\pm}(p)|)\text{sgn}(\epsilon_{\pm}(p)) \right] + \theta(\delta + \epsilon^{-}(p))
\]

\[
n_u - n_d = -\frac{\partial}{\partial \delta} \Omega(\Delta, \mu, \delta)
\]

\[
= \int \frac{d^3 p}{(2\pi)^3} \sum_{\pm} \left[ 2\theta(\delta - E_{\pm}(p)) + \theta(\delta - |\epsilon_{\pm}(p)|) - \theta(\delta + \epsilon^{-}(p)) \right].
\]

(20)

All three dimensional integrals are cutoff at momenta \(\Lambda\). For illustration, we choose the values \(\Lambda = 900\) MeV, \(g^2/\Lambda^2 = 4.4\). They produce a gap of around \(\Delta \sim 100\) MeV at \(\mu = 300\) MeV and \(\delta = 0\). In fig. (1) we have \(\Omega(\Delta, \mu, \delta)\) as a function of \(\Delta\) for different values of \(\mu\) and \(\delta\). They all correspond to states with the same value of \(n_u + n_d\), but increasing values of the asymmetry \(n_u - n_d\). The minimum of the first curve, top to bottom, corresponds to the flavor symmetric, superconducting state. As mentioned above, the superconducting state consists of paired blue and red quarks of both flavors, with green quarks left unpaired.

There are two ways of looking at flavor asymmetric matter: one may fix different densities or different chemical potentials. Let us first consider the first possibility and imagine increasing the asymmetry \(n_u - n_d\) from its the initial zero value while keeping the total density \(n_u + n_d\) fixed. States with small \(n_u - n_d\) values can be created by adding to this symmetric state up-quasi-particles and down-quasi-holes. In the grand canonical formalism used here this is accomplished by increasing the value of \(\delta\) from zero to some small value. Since there is a gap in the spectrum of blue and red quasi-particle, no quasi-particle with these colors are created for \(\delta < \Delta\). The flavor asymmetry is made up only with green, up-quasi-particles and green, down quasi-holes. If all colors were paired we would see no change in \(\Omega\) at all in the superconducting phase until the split between chemical potentials \(\delta\) equaled the gap
The potential can be more easily lowered in the unstable state with $\Delta = 0$ since there quarks of all colors can be created, even with small $\delta$. As a result $\Omega(\Delta = 0, \mu, \delta)$ decreases, with increasing $\delta$, faster than $\Omega(\Delta \neq 0, \mu, \delta)$, as can be seen in the second curve in fig. (1). At some value of $\delta$ the superconducting and the normal phase are both stable. The value of $n_u - n_d$ is much higher on the normal phase than in the superconducting phase, by the reasons explained above. One could imagine that, as $n_u - n_d$ increases, the system proceeds towards a mixed phase of $\Delta = 0$ and $\Delta \neq 0$ phases, with the up-rich $\Delta = 0$ phase occupying more and more of the space until the superconducting phase completely disappears. The presence of two conserved charges makes things a little more involved. The reason is that the total density $n_u + n_d$ in the superconducting phase is much higher than the total density in the normal phase. Thus, even though by choosing the right amounts of the two phases one can arrange so that the overall asymmetry $n_u - n_d$ be any number between the asymmetries in the two phases, the overall total density is always smaller than the initial total density. In order to achieve higher values of $n_u - n_d$ and keep $n_u + n_d$ constant we need to increase also the value of the average chemical potential $\mu$ (the result in shown on fig. (1), fourth curve to meet the vertical axis, from top to bottom). The value of $\mu$ is chosen in such a way as to satisfy

$$n_u + n_d = x(n_u + n_d)_{\text{super}} + (1 - x)(n_u + n_d)_{\text{normal}}$$
$$n_u - n_d = x(n_u - n_d)_{\text{super}} + (1 - x)(n_u - n_d)_{\text{normal}},$$  \hspace{1cm} (21)

for some $x$, $0 \leq x \leq 1$. $\delta$ is determined by the condition that $\Omega(\Delta, \mu, \delta)$ should have two degenerate minima as a function of $\Delta$ (so phases can coexist), so we are left with two equations eq. (21) and two independent parameters ($\mu$ and $x$) so equations eq. (21) can, generically, be (uniquely) satisfied. The transition, then, proceeds through a series of

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**FIG. 1.** Potential $\Omega$ in units of $10^{10}$ MeV$^4$ as a function of the gap $\Delta$ in MeV. By the order the curves meet the vertical axis, from top to bottom: i) $\mu = 300$ MeV, $\delta = 0$, ii) $\mu = 300$ MeV, $\delta = 40$ MeV, iii) $\mu = 300$ MeV, $\delta = 77.5$ MeV, $(n_u + n_d)_{\text{normal}} = 3.1 \times 10^6$ MeV$^3$, $(n_u + n_d)_{\text{super}} = 3.91 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{normal}} = 2.34 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{super}} = 0.90 \times 10^6$ MeV$^3$, iv) $\mu = 310$ MeV, $\delta = 83$ MeV, $(n_u + n_d)_{\text{normal}} = 4.31 \times 10^6$ MeV$^3$, $(n_u + n_d)_{\text{super}} = 3.45 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{normal}} = 1.04 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{super}} = 2.69 \times 10^6$ MeV$^3$, v) $\mu = 322.5$ MeV, $\delta = 89$ MeV, $(n_u + n_d)_{\text{normal}} = 4.92 \times 10^6$ MeV$^3$, $(n_u + n_d)_{\text{super}} = 3.91 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{normal}} = 1.22 \times 10^6$ MeV$^3$, $(n_u - n_d)_{\text{super}} = 3.14 \times 10^6$ MeV$^3$, vi) $\mu = 310$ MeV, $\delta = 120$ MeV.
mixed states, each one with different values of $\mu$ and $\delta$ (arranged in such a way as to keep $n_u + n_d$ constant while increasing $n_u - n_d$). Each one of these mixed states is formed by a slightly asymmetric, high density superconducting with a larger gap than the one before the transition started and a highly asymmetric, low density normal phase. The mixed phase exists until its normal component occupies the whole space and the total density of the normal phase equals the initial value of $n_u + n_d$ (fifth curve to meet the vertical axis in fig. (1), top to bottom). Increasing $n_u - n_d$ even further is accomplished by increasing $\delta$ again, at roughly fixed $\mu$. At this point the superconducting state is not even metastable anymore.

The different phases as a function of $n_u + n_d$ and $(n_u - n_d)/(n_u + n_d)$ are sketched on fig. (2). The horizontal axis, corresponding to the symmetric case, is in the superconducting phase for any density, since the BCS instability exists for any density. The vertical axis correspond to the completely asymmetric case, with $\mu_u = -\mu_d$. There are only up-quarks and down-anti-quarks present, no pairing is possible so those states are in the normal phase. In between these two phases there is a large mixed phase region that occupy most of the $n_u, n_d \geq 0$ region. The precise location of the phase boundaries are somewhat sensitive to the parameters of the model. The example shown in fig. (2) was obtained with the parameters discussed above. The distance between the superconductor-mixed phase is on the large side among those obtained with other reasonable parameters. As fig. (2) shows, for asymmetries larger than about 20% the system remains in the mixed phase for arbitrarily high densities. The phase boundaries in the low density region, where the model considered here has no resemblance to real QCD, continue all the way to the origin. Unfortunately it is very hard to determine the properties and precise location of the mixed phase with a reasonable degree of certainty. The bubbles of the normal phase are charged and, if formed in heavy ion collisions, may lead to charge separation that could be observable.

The color-flavor locked state, expected in the 3 massless flavor case, is expected to behave under asymmetries in a similar manner. Actually, since color is completely broken and all
quark colors have a gap, one could expect the mechanisms discussed here to be even more effective than in the 2SC state. In the case of massive strange quarks though, it remains to be seen whether existence of “gapless superconductors” pointed out in [22] may hinder this conclusion. A discussion of this question, as well as a more complete treatment of the two and three flavors asymmetric matter will be presented elsewhere [23].

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