Color Superconductivity in a Dense Quark Matter

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Abstract

The color superconductivity of a dense quark matter is reviewed with emphasis on the long range nature of the pairing force and the multiplicity of the order parameter. The former gives rise to a non BCS behavior of the superconducting energy scale and the latter modifies the critical value of the Ginzburg-Landau parameter that separates the superconductivity of type I and that of type II.

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I. INTRODUCTION

In this lecture, I would like to review our works for the past years on the color superconductivity in a dense quark matter. While it is an old subject dated back to 1970’s [1] [2] [3], the revived interests on the color superconductivity since 1998 [4] [5] [6] are largely promoted by the experimental efforts of exploring the phase diagram of the strong interaction through RHIC or inside neutron stars. While lattice simulations have been successful at nonzero temperature, it becomes prohibitively difficult when the chemical potential becomes nonzero. Analytical technique remains the only means especially at large chemical potential, where an ideal gas of quarks and gluons presents the leading order approximation. The important progresses that have been made recently include the discovery of the color-flavor locked structure in the super phase [6] and the first-principle determination of the energy gap and the transition temperature [7] [8] [9] [10] [11] [12] [13]. The entire subject has been reviewed by several authors [14].

Among many novel characters of the color superconductivity, I would like to take this opportunity to highlight two of them which differ remarkably from the non-relativistic superconductivity in a metal. The first is the long range nature of the pairing force, which is responsible to the non BCS exponent of the weak coupling formula of the gap energy or the transition temperature and the suppression of the pre-exponential factor. The second is the multiplicity of the order parameters, which leads to a rich variety of inhomogeneous condensates and a different value of the critical Ginzburg-Landau parameter that separates the superconductivity of type I from that of type II. Many technical details have been suppressed and this lecture will serve a tour guide to our papers on the subject and the related works by others.
II. THE LONG RANGE PAIRING FORCE.

Let us consider a quark matter of ultra high baryon density such that the corresponding chemical potential is well above $\Lambda_{\text{QCD}}$ to warrant a perturbative treatment because of the asymptotic freedom. While this density may not be attainable in a realistic quark matter, say the core of a neutron star, the approximation there is systematically under control. The Lagrangian density of QCD at a nonzero chemical potential $\mu$ in the chiral limit reads:

$$\mathcal{L} = -\frac{1}{2} \text{tr} F_{\mu \nu}^l F_{\mu \nu}^l - \bar{\psi} \left( \frac{\partial}{\partial x^\mu} - ig A_\mu \right) \psi + \mu \bar{\psi} \gamma_4 \psi + \text{renormalization counter terms},$$

where $A_\mu = A_\mu^l T^l$ and $F_{\mu \nu}^l = \partial_x A_{\nu}^l - \partial_x A_{\mu}^l - ig [A_\mu, A_\nu]$ with $T^l$ the $SU(N_c)$ generator in its fundamental representation and $\psi$ is a Dirac spinor with both color and flavor indices. The renormalized coupling constant is defined at the chemical potential via

$$g = \frac{24\pi^2}{(11N_c - 2N_f) \ln \frac{\mu}{\Lambda}}$$

with $\mu >> \Lambda$ and $\Lambda = \Lambda_{\text{QCD}}$ for $N_c = N_f = 3$.

Perturbatively, the di-quark interaction is dominated by the process of one-gluon exchange, as is shown in Fig. 1. The amplitude is simply that of the one-photon exchange in QED multiplying the group theoretic factors $T_l^{c_1 c_2} T_l^{c_2 c_1}$, which can be decomposed into a color anti-symmetric channel (anti-triplet for $SU(3)$) and a color symmetric one (sextet for $SU(3)$), i.e.

$$T_l^{c_1 c_2} T_l^{c_2 c_1} = -\frac{N_c + 1}{4N_c} (\delta_{c_1 c_1} \delta_{c_2 c_2} - \delta_{c_1 c_2} \delta_{c_2 c_1}) + \frac{N_c - 1}{4N_c} (\delta_{c_1 c_2} \delta_{c_2 c_2} + \delta_{c_1 c_2} \delta_{c_2 c_1})$$

As both the electric and magnetic parts of the one-photon exchange are repulsive between two electrons flying in opposite directions, the one-gluon exchange interaction between two quarks flying in opposite directions is attractive in the color anti-symmetric channel because of the negative sign of the first term on r. h. s. of (3).
At high baryon density, the Fermi sea of quarks tends to screen the one-gluon exchange interaction through the HDL (hard dense loop) resummed gluon propagator \cite{15}, which takes the form

\[
D^{ab}_{ij}(k, i\omega) = \frac{-i\delta^{ab}}{k^2 + \omega^2 + \sigma_M(k, \omega)} \left( \delta_{ij} - \frac{k_i k_j}{k^2} \right)
\]  

(4)

and

\[
D^{ab}_{44}(k, i\omega) = \frac{-i\delta^{ab}}{k^2 + \sigma_E(k, \omega)}
\]  

(5)

and \(D_{4j}(k, i\omega) = 0\) in Coulomb gauge with \(i\omega\) the Matsubara energy of the gluon. In the region of energy and momentum for pairing, i.e. \(k << \mu, \omega << \mu\) and \(\omega << k\), the magnetic self-energy is given by

\[
\sigma_M(k, \omega) \simeq \frac{\pi}{4} m_D^2 \frac{|\omega|}{k}
\]  

(6)

and the electric self-energy by

\[
\sigma_E(k, \omega) \simeq m_D^2
\]  

(7)

with \(m_D^2 = \frac{N_f g^2 \mu^2}{2\pi}\) the Debye mass. The absence of the screening of the long range magnetic gluon propagator in the static limit, \(\omega = 0\), leads to the forward singularity of the di-quark scattering, which dominates the pairing force. The long range magnetic gluons introduces also the non Fermi liquid behavior, that renders the quark self energy function logarithmically enhanced toward the Fermi surface, i.e.
\[ \Sigma(i\nu, p) \big|_{p=\mu} \approx -\gamma_4 \frac{N_c^2 - 1}{N_c} \frac{g^2}{24\pi^2} i\nu \ln \frac{K_c^3}{\pi m_D^2 |\nu|} \]  

(8)

with \( i\nu \) the Matsubara energy of the quark and \( K_c \) a cutoff momentum. \( k_B T << K_c << \mu \) a cutoff momentum. The energy scale of the color superconductivity is measured by the energy gap of the quark spectrum at \( T = 0 \) or the superconducting transition temperature \( T_c \). The determination of the latter quantity will be reviewed in this section.

The proper vertex function \( \Gamma(i\nu', p'|i\nu, p) \) for di-quark scattering with incoming energy-momenta \((p, \pm i\nu)\) and outgoing energy-momenta \((p', \pm i\nu')\) satisfies a Dyson-Schwinger equation, which, when projected into the color anti-symmetric channel with an angular momentum \( J \), takes the form

\[ \Gamma_J(i\nu', p'|i\nu, p) = \tilde{\Gamma}_J(i\nu', p'|i\nu, p) + k_B T \sum_{\nu''} \int_0^\infty dq K_J(i\nu', p'|i\nu'', q) \Gamma_J(i\nu'', q|i\nu, p), \]  

(9)

where \( \tilde{\Gamma}_J(i\nu', p'|i\nu, p) \) stands for the two particle irreducible part of \( \Gamma_J(i\nu', p'|i\nu, p) \), the kernel

\[ K_J(i\nu', p'|i\nu'', q) = \frac{p^2 \tilde{\Gamma}_J(i\nu', p'|i\nu, p)}{2\pi^2(2J+1)} S(i\nu, p)S(-i\nu, p) \]  

(10)

with \( S(i\nu, p) \) the full quark propagator. The spinor indices and their contraction structure have been suppressed in (9) and (10). To the lowest order in coupling constant, \( \tilde{\Gamma} \) is given by the one-gluon exchange diagram in Fig.1 and the quark propagators in the kernel (10) are replaced by bare ones. For \( \nu < \mu, \nu' \), and \( p \sim p' \sim \mu \), we have

\[ \Gamma_J(i\nu', p'|i\nu, p) \simeq -\frac{g^2}{12\mu^2} \left( 1 + \frac{1}{N_c} \right) \left( \ln \frac{1}{|\nu' - \nu|} + 3c_J \right) \simeq -\frac{g^2}{12\mu^2} \left( 1 + \frac{1}{N_c} \right) \left( \ln \frac{1}{|\nu'|} + 3c_J \right) \]  

(11)

with \( \hat{\nu} = \frac{g^2}{256\pi^4} \left( \frac{N_f}{2} \right) \frac{\hat{\xi}}{\mu} \) with \( c_J = 1 \) for \( J = 0 \) and \( c_J = \exp \left( -2 \sum_n \frac{1}{n} \right) \) for \( J \neq 0 \). The last step of (11) follows from Son’s approximation [7] with \( \nu_> = \max(\nu, \nu') \).

The pairing temperature within each angular momentum channel, \( T_c^{(J)} \), corresponds to the highest temperature at which the Fredholm determinant of eq.(9), \( D_J = \det(1 - K_J) \) vanishes and the transition temperature is \( T_c = \max(T_c^{(J)}) \) for all \( J \). Here we have assumed the pairing between two quarks of the same herlicity, which contains the s-wave channel \((J = 0)\) and will take the maximum advantage of the pairing force.
To highlight the impact of the long range pairing force in the weak coupling formula for $T_C$, we employ the standard expansion of the Fredholm determinant,

$$D_J = 1 - k_B T \sum_\nu \int_0^\infty dp \int_0^{\infty} dq \int_0^{\infty} dp K(i\nu, p| i\nu, p) + \frac{(k_B T)^2}{2} \sum_\nu \sum_{\nu'} \int_0^\infty dp \int_0^{\infty} dq \int_0^{\infty} dp K(i\nu, q| i\nu, p) K(i\nu', q| i\nu', q) \right| + \ldots$$

(12)

For pairing mediated by phonons, which is of a short range, the $m$-th term of the expansion is of the order of $g^{2m} \ln \frac{\omega_D}{k_B T}$ with $\omega_D$ the Debye frequency. At the transition, $g^2 \ln \frac{\omega_D}{k_B T} \sim 1$ and the remaining terms are of the order of $g^{2m-2}$. This gives rise to the scaling formula

$$k_B T_c = c_{\text{BCS}} \exp\left(-\frac{\kappa_{\text{BCS}}}{g}\right)$$

with $\kappa_{\text{BCS}}$ fixed by the term of $m = 1$ and the leading order of $c_{\text{BCS}}$ fixed by the term of $m = 2$. In case of QCD, the logarithm inside the kernel (10) makes the $m$-th term of the expansion go like $g^{2m} \ln^{2m} \frac{\mu}{k_B T}$ and the Fredholm determinant is dominated by a function of $g \ln \frac{\mu}{k_B T}$. It follows then that the transition temperature scales like

$$k_B T_c = c_{\text{QCD}} \exp\left[-\frac{\kappa_{\text{QCD}}}{g}\right]$$

[7] [16] [17] leaving both $\kappa_{\text{QCD}}$ and $c_{\text{QCD}}$ non-trivial to determine.

A perturbative method was developed in [11] [12], which yields the formula of the pairing temperature

$$k_B T_c^{(J)} = cc'c''c_J \frac{\mu}{g^5} e^{-\frac{\pi}{\delta}} [1 + O(g \ln g)],$$

(13)

where the exponent factor $\kappa = \sqrt{\frac{6N_c}{N_c + 1}} \pi^2$ was first obtained in [7], the pre-exponential factor $c = 1024 \sqrt{2 \pi^3 N_f}^{-\frac{3}{2}}$ was found in [8] and [9], the factor $c' = 2e^\gamma$ with $\gamma$ the Euler constant was found in [9] and [12]. The factor $c'' = \exp\left[-\frac{1}{16}(\pi^2 + 4)(N_c - 1)\right]$ stems from the non Fermi liquid behavior of the quark self energy (8) and was calculated in [12] and reproduced in [13] via the gap equation (The existence of this correction was suggested in [7]). The gauge invariance of the formula (13) was discussed in [18]. The angular momentum dependent factor $c_J$ was calculated in [12] and the consequence of the universal exponent for different angular momentum channels in (13) was discussed in [19] in the context of the crystalline superconductivity. Some higher order terms have also been identified [20]. It follows from (13) that the superconducting transition temperature is $T_c = T_c^{(0)}$. 

6
A similar perturbative method for the gap energy has been developed in [21].

**III. THE MULTIPLICITY OF THE ORDER PARAMETER.**

The Ginzburg-Landau theory is a powerful tool to explore the superconducting state right below the transition temperature. Unlike the ordinary superconductors, the order parameter of the color superconductivity which describes the di-quark condensate, $\Psi_{f_1 f_2}^{c_1 c_2}$ is of multi-components, carrying the color-flavor indices and the chirality. For three color and three flavors, the symmetry group of theory in the chiral limit is $SU(3)_c \times SU(3)_{f_R} \times SU(3)_{f_L} \times U(1)_B$ of which the electromagnetic gauge group, $U(1)_{em}$ is a subgroup. Restricting within the even parity sector and neglecting the small projection in the color-sextet representation, the order parameter takes the form

$$\Psi_{f_1 f_2}^{c_1 c_2} = \epsilon^{c_1 c_2 c} \epsilon_{f_1 f_2 f} \Phi^c.$$  \hspace{1cm} (14)

and the Ginzburg-Landau free energy functional which is consistent with the symmetry group reads [22] [23]

$$\Gamma = \int d^3r \left[ \frac{1}{4} F_{ij}^l F_{ij}^l + \frac{1}{2} (\vec{\nabla} \times \vec{A})^2 + 4 \text{tr}(\vec{D} \Phi)^\dagger (\vec{D} \Phi) + 4 \text{tr} \Phi^i \Phi + b_1 (\text{tr} \Phi^\dagger \Phi)^2 + b_2 (\text{tr} \Phi^\dagger \Phi)^2 \right]$$ \hspace{1cm} (15)

where the covariant derivative $\vec{D} \Phi = \vec{\nabla} \Phi - ig \vec{A} \Phi + i \frac{2}{\sqrt{3}} e \vec{A} \Phi T^8$ with $\vec{A}$ the color vector potential and $\vec{A}$ the ordinary electromagnetic vector potential. The $3 \times 3$ matrix $\Phi = \Phi_0 + \Phi_1 T^l$ with $\Phi_0$ the singlet and $\Phi_1$ the octet under a simultaneous color-flavor rotation. The generator here, $T^l$, should be understood to be that of the anti-triplet representation of $SU(3)_c$. At weak coupling, the one-gluon exchange lends us the following expressions of the Ginzburg-Landau coefficients [22] [23]

$$a = \frac{48\pi^2}{\zeta(3)} k^2 B T_C (T - T_C),$$

$$b_1 = b_2 = \frac{576\pi^4}{\zeta(3)} \left( \frac{k B T_C}{\mu} \right)^2.$$  \hspace{1cm} (16)
which follows also from the mean field approximation of the NJL effective action at moderate coupling if \( k_B T_c \ll \mu \).

Let us review first the structure of a homogeneous condensate [22], i.e. \( \vec{A} = \vec{\nabla} \Phi = 0 \), for which three regions of the parameter space need to be considered respectively. 1) \( b_1 > 0 \) and \( b_1 + 3b_2 > 0 \) The minimum free energy corresponds to the color-flavor locked condensate [6] with \( \Phi = \phi_0 U \), where

\[
\phi_0 = \sqrt{-\frac{2a}{b_1 + 3b_2}},
\]

and \( U \) is a unitary matrix (The nontrivial winding of \( U \) onto the gauge group, \( SU(3)_c \times U(1)_{em.} \) gives rise to vortex filaments [23]). 2) \( b_1 < 0 \) but \( b_1 + b_2 > 0 \): In this case the color-flavor locked condensate (17) becomes a saddle point of the free energy, which is nevertheless bounded from below. The minimum corresponds to an isoscalar condensate [22], given by \( \Phi = \text{diag}(\sqrt{-\frac{2a}{b_1 + b_2}}e^{i\alpha}, 0, 0) \). 3) For \( b_1 \) and \( b_2 \) outside the region specified by 1) and 2), the free energy is no longer bounded from below. Higher powers of the order parameter have to be restored and the superconducting transition becomes the first order one. In the following we shall focus our attention to the case 1) of the parameters.

The color-flavor locked condensate breaks the original symmetry group down to a \( SU(3) \) of simultaneous color-flavor rotation, among which there is an unbroken \( U(1) \) gauge symmetry and the corresponding gauge potential \( \vec{V} \) is obtained through the rotation [24]

\[
\vec{V} = -\vec{A}^8 \sin \theta + \vec{A} \cos \theta
\]

\[
\vec{V} = \vec{A}^8 \cos \theta + \vec{A} \sin \theta,
\]

with \( \tan \theta = \frac{-2e}{\sqrt{3}g} \). In comparison with the electroweak theory, the field \( \vec{V} \) and \( \vec{V} \) are the analogs of the photon and the \( Z \)-boson, and the rest of the \( A \)'s are like the \( W \)-boson’s. The mass of the \( \vec{V} \) and that of \( \vec{A}^l \) (\( l = 1, ..., 7 \)) are

\[
m_Z^2 = 4g^2 \phi_0^2 \sec^2 \theta = \frac{1}{\delta^2}, \quad m_W^2 = m_Z^2 \cos^2 \theta = \frac{1}{\delta'^2},
\]

with \( \delta \) and \( \delta' \) the corresponding penetration depths. Other low-lying excitations of even parity consist of a Goldstone boson associated to the broken baryon number, \( U(1)_B \), the
Higgs bosons associated to $\vec{V}$ and to $\vec{A}$ ($l = 1, \ldots, 7$) with masses

$$m_H^2 = (b_1 + 3b_2)\phi_0^2 = \frac{2}{\xi^2}, \quad m_H'^2 = b_1\phi_0^2 = \frac{2}{\xi'^2},$$

and $(\xi, \xi')$ the corresponding coherence lengths. The excitations of odd parities includes the Goldstone bosons associated to the chiral symmetry breaking triggered by the color-flavor locking [6].

An interesting issue to address is the type of the superconductivity in response to an external magnetic field, which is partially screened because of its projection onto the broken $\vec{V}$ field through (18). This amounts to calculate the free energy of a domain wall separating the super phase and the normal phase with the bulk of super phase and that of the normal phase held in thermal equilibrium under an external magnetic field of the critical strength,

$$H_c = 2\sqrt{\frac{6a^2}{b_1 + 3b_2}}|\csc \theta|.$$  \hspace{1cm} (19)

The simplest ansatz of the solution to the Ginzburg-Landau equations that minimizes the domain wall free energy consists of only nonzero $\vec{V}$, parallel to the external magnetic field, and nonzero components $\Phi_0$ and $\Phi_8$, which maintains the maximum symmetry spared by the boundary condition. Upon introducing dimensionless quantities via

$$s = \delta x, \quad \Phi_0 + \frac{1}{\sqrt{3}}\Phi_8 = \phi_0 u, \quad \chi = \sqrt{2}(\Phi_0 - \frac{1}{2\sqrt{3}}\Phi_8) = \sqrt{2}\phi_0 v, \quad V = -\frac{\sqrt{-3a}}{g}A\cos \theta,$$

the corresponding Ginzburg-Landau equations become

$$-A'' + \frac{1}{3}A(2u^2 + v^2) = 0,$$

$$-\frac{1}{\pi^2}u'' + (A^2 - 1)u + \frac{1}{3}(2u^2 + v^2)u + \frac{1}{3}\rho(u^2 - v^2)u = 0,$$

$$-\frac{1}{\pi^2}v'' + \frac{1}{4}(A^2 - 4)v + \frac{1}{6}(u^2 + 5v^2)v - \frac{1}{6}\rho(u^2 - v^2)v = 0,$$  \hspace{1cm} (20)

subject to the boundary conditions

$$u \mapsto 0, \quad v \mapsto 0, \quad A' \mapsto 1 \quad \text{at} \quad s \mapsto -\infty$$

$$u \mapsto 1, \quad v \mapsto 1, \quad A' \mapsto 0 \quad \text{at} \quad s \mapsto \infty$$  \hspace{1cm} (21)
where $\kappa = \frac{\delta}{\xi}$ is the Ginzburg-Landau parameter, $\rho = \frac{b_1 - \sqrt{b_1^2 + 3b_2}}{b_1 + 3b_2}$ with $\rho = -\frac{1}{2}$ for mean field approximation, and the prime denotes the derivative with respect to $s$. In case of an ordinary superconductors with one component of the order parameter, Ginzburg-Landau found analytically [26] that the domain wall energy vanishes at

$$\kappa = \kappa_c = \frac{1}{\sqrt{2}} \simeq 0.707,$$

which was later clarified by Arikosov [27] as the demarcation between the type I superconductivity ($\kappa < \kappa_c$) and the type II one ($\kappa > \kappa_c$). The equations (20) are more complicated and depend on two dimensionless parameters, $\kappa$ and $\rho$. Nevertheless a set of inequalities have been established following the variational arguments [25]. We have shown that the domain wall energy is a decreasing function of $\kappa$ and an increasing function of $\rho$. Consequently,

$$\kappa_c(\rho) \leq \frac{1}{\sqrt{2}}, \quad \frac{d\kappa_c}{d\rho} \geq 0.$$

The numerical solution of the equations (21) shows that

$$\kappa_c = 0.589$$

for $\rho = -\frac{1}{2}$. While the color-flavor locked component $\Phi_0$ dominates the bulk of the condensate, the unlocked component, $\Phi_8$, shows up near the domain wall.

Various fluctuation effects of the Ginzburg-Landau theory have also been considered in the literature, see e.g. [25] [28].

**IV. CONCLUDING REMARKS**

In this lecture, I have reviewed some novel properties of the color superconductivity. A systematical approach to determine various physical quantities at weak coupling has been sketched. While the accuracy of the perturbative formula for the transition temperature (13) and that for the Ginzburg-Landau coefficients (16) is uncertain when extrapolated to a realistic quark matter, it is instructive to substitute in the observationally attainable quark density for an order of magnitude estimation of the quantities of interests.
Taking the generic chemical potential of in the core of a neutron star, $\mu = 400\text{MeV}$, as a benchmark, we have $\alpha_S \equiv \frac{\alpha^2}{4\pi} \simeq 1$ and the mixing angle $\theta = -5.6^\circ$ for $N_c = N_f = 3$ and $\Lambda = \Lambda_{\text{QCD}} = 200\text{MeV}$. Using eq. (16) and (19), we find the thermodynamical critical field $H_c = 1.47 \times 10^{20} \left(\frac{k_B T}{\mu}\right) \left(1 - \frac{T}{T_c}\right) \text{Gauss}$. Furthermore, the criterion (23) implies that the color superconductivity is type I if $k_B T_C < 14\text{MeV}$ and type II if $k_B T_C > 14\text{MeV}$. It follows from (13) that $k_B T_c \simeq 3.5\text{MeV}$ with one-gluon exchange approximation and the color superconductivity is of type I. Beyond the one-gluon exchange approximation, $k_B T_c$ could be high enough to cross over to the type II region, but it is unlikely to be a strong type II system. The critical magnetic field, which is by orders of magnitude higher than the typical magnetic field in a neutron star, is too strong to lead to observational distinctions between the two types of the color-superconductivity.

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