We review what is different and what is similar in a color superconductor as compared to an ordinary BCS superconductor. The parametric dependence of the zero-temperature gap, $\phi_0$, on the coupling constant differs in QCD from that in BCS theory. On the other hand, the transition temperature to the superconducting phase, $T_c$, is related to the zero-temperature gap in the same way in QCD as in BCS theory, $T_c/\phi_0 \simeq 0.567$.

1 Cold, dense fermionic matter

Consider a degenerate, non-interacting fermionic system, where all momentum states up to the Fermi surface are occupied. The fermion occupation number is a step function, $n_0(\epsilon_0) = \theta(-\epsilon_0)$, where $\epsilon_0 \equiv \omega - \mu$ is the energy of the fermions with respect to the Fermi energy, $\mu$, and $\omega$ is their kinetic energy. For non-relativistic particles with 3-momentum $k$ and mass $m$, $\omega = k^2/(2m)$, while for massless (ultrarelativistic) particles, $\omega = k \equiv |k|$. (The units are $\hbar = k_B = c = 1$.) The excitation spectrum of the fermions consists of a particle branch, $\epsilon_0^p \equiv \epsilon_0$, and a hole branch, $\epsilon_0^h \equiv -\epsilon_0$. The formation of particle-hole excitations at the Fermi surface costs no energy, $\epsilon_0^p + \epsilon_0^h = 0$.

Now switch on an interaction, for the sake of simplicity a point-like four-fermion interaction with interaction strength $G^2$. Let the sign of $G^2$ be defined such that $G^2 > 0$ in the case of an attractive interaction, and $G^2 < 0$, if the interaction is repulsive. Due to Pauli’s exclusion principle and energy conservation, scattering of fermions can occur exclusively at the Fermi surface. In other words, physical scattering processes are possible only for $\epsilon_0 \to 0$. The amplitude for fermion-fermion scattering is

$$\Gamma(\epsilon_0) \sim \frac{G^2}{1 - G^2 \ln(\mu/\epsilon_0)} , \quad (1)$$
where we suppress all constant factors that are irrelevant for the qualitative arguments presented here. As $\epsilon_0 \to 0$, in the case of an attractive interaction the scattering amplitude develops a singularity at an energy scale

$$\epsilon_0^* \sim \mu \exp(-1/G^2) \ .$$

(2)

For repulsive interactions, no such singularity occurs. Obviously, the singularity (2) occurs even when the attractive interaction is arbitrarily weak, $G^2 \to 0^+$, all that changes is the energy scale $\epsilon_0^*$ of the singularity.

This singularity is the famous Cooper instability. It is cured by the formation of Cooper pairs which, as bosons, condense in the true, energetically favored ground state of the system. Macroscopically, this leads to the phenomenon of superfluidity, or, if the fermions carry charge, superconductivity.

The excitation spectrum of the system also changes. A gap $\phi_0$ forms at the Fermi surface, separating the branch for quasiparticle excitations, $\epsilon^p = -\epsilon$, $\epsilon \equiv \sqrt{\epsilon_0^2 + \phi_0^2}$, from the one for quasihole excitations, $\epsilon^h = +\epsilon$. Now exciting a quasiparticle–quasihole pair costs at least an energy $\epsilon^p + \epsilon^h \geq 2\phi_0$. The quasiparticle occupation number, $n(\epsilon_0)$, is “smeared” around the (original) Fermi surface, $n(\epsilon_0) = (\epsilon - \epsilon_0)/(2\epsilon)$.

To compute the gap, one has to solve a gap equation. In the case of a point-like four-fermion interaction, this equation takes the form

$$\phi_0 = G^2 \int_{-\mu}^{0} \frac{d\epsilon_0}{\epsilon} \phi_0 .$$

(3)

(Again, irrelevant constant factors are omitted.) Besides the trivial (energetically disfavored) solution $\phi_0 = 0$, this equation has always a non-trivial (energetically favored) solution, $\phi_0 \neq 0$. Since the gap is a constant for point-like four-fermion interactions, we can divide both sides of Eq. (3) by $\phi_0$. The remaining integral can be solved exactly, with the result

$$\phi_0 \simeq 2\mu \exp(-1/G^2) .$$

(4)

Apparently, $\phi_0 > \epsilon_0^*$, cf. Eq. (2). In other words, since there are no quasiparticle states with energy $\epsilon^p \geq -\phi_0$ (or quasihole states with $\epsilon^h \leq \phi_0$), the gap has just the right order of magnitude to prevent the scattering amplitude from developing a singularity.

The assumption of a point-like four-fermion interaction can be relaxed. Assume that the interaction is mediated by a scalar boson of mass $M$. For the sake of definiteness, assume that the boson mass is generated by many-body interactions at nonzero density, $M \sim g\mu$. The boson-fermion coupling is
denoted by \( g \), the boson propagator is \( \Delta(P) = 1/(M^2 - P^2) \), \( P \equiv P^\mu = (p_0, \mathbf{p}) \), \( P^2 = p_0^2 - \mathbf{p}^2 \). In this case, the gap equation (3) becomes

\[
\phi_0(k) = g^2 \int_{-\mu}^{0} d\epsilon_0/\epsilon \, \phi_0(q) \frac{q}{k} \ln \left[ \frac{M^2 + (k + q)^2}{M^2 + (k - q)^2} \right].
\]

(5)

Here, \( q \equiv \epsilon_0 + \mu \). The frequency dependence of the boson propagator and, consequently, that of the gap function has been neglected, based on the argument that in weak coupling, \( g \ll 1 \), \( p_0 \sim \epsilon_\sim \phi_0^\sim \mu \exp(-1/g^2) \ll M \sim g\mu \). The logarithm arises from the integration over the angle between the boson 3-momentum \( \mathbf{p} \equiv k - \mathbf{q} \) and the 3-momentum of the fermion in the condensate, \( \mathbf{q} \), for details see Ref. [2]. Note that this factor enhances contributions from the region of momenta \( \mathbf{q} \approx k \) (collinear enhancement).

In the case of a boson-mediated interaction, the gap is no longer constant, but a function of momentum. Consequently, the gap equation (3) is no longer a simple fix-point equation for \( \phi_0 \), but an integral equation which has to be solved numerically. However, to estimate the order of magnitude of the gap at the Fermi surface, \( k \equiv \mu \) (for massless fermions), one may make the following approximation. First note that, due to the factor \( 1/\epsilon \), the integrand peaks at the Fermi surface. It is then sufficient to approximate the slowly varying logarithm and the gap function with their values for \( q = k = \mu \). This leads to the estimate

\[
\phi_0 \simeq 2 \mu \exp \left[ -\frac{1}{g^2 \ln(1 + 2\mu^2/M^2)} \right]
\]

(6)

for the value of the gap function at the Fermi surface, \( \phi_0 \equiv \phi_0(\mu) \), which should be compared with Eq. (4). All that changed is that the coupling constant is effectively increased by the logarithm originating from collinear enhancement. For \( M \sim g\mu \), the logarithm becomes \( \sim \ln(1/g) \) in weak coupling.

2 Cold, dense quark matter

The density in cold quark matter increases \( \sim \mu^3 \). Asymptotic freedom then implies that single-gluon exchange becomes the dominant interaction between quarks. Single-gluon exchange is attractive in the color-antitriplet channel, and therefore leads to color superconductivity in cold, dense quark matter [3, 4, 5]. Recently, considerable activity was generated [6–39] by the work of Refs. [6], which suggested that the zero-temperature color-superconducting gap \( \phi_0 \) could be as large as 100 MeV. This order of magnitude is quite surprising, because earlier work by Bailin and Love [3] estimated the gap to be \( \phi_0 \sim 1 \) MeV. While gaps of order 100 MeV could also be relevant for the physics of nuclear collisions (see below), gaps of about 1 MeV allow at most for the
possibility that neutron star cores, if consisting of quark matter, could be color superconductors.

The authors of Refs. 6, 9 based their arguments on a simple model where quarks interact via a point-like four-fermion interaction, giving rise to a gap equation of the type (3). The coupling strength \( G^2 \) was adjusted such that the model reproduced the order of magnitude of the chiral transition at nonzero temperature and zero quark chemical potential, \( T_\chi \sim 150 \text{ MeV} \). The earlier work of Bailin and Love 5 was already more sophisticated in the sense that they used one-gluon exchange, employing gluon propagators with electric and magnetic screening masses. This gives rise to a gap equation of the form (5), with \( g \) being the strong coupling constant.

Unfortunately, both approaches fail to capture an essential property of single-gluon exchange: at zero temperature, due to the absence of magnetic screening, magnetic interactions are truly long-range. Surprisingly, this fact was already known to Barrois 3, but apparently never made it into the published literature.

Long-range magnetic interactions have the important consequence that one can no longer neglect the frequency dependence of the boson propagator, as done in the derivation of (5). For massless boson exchange, the gap equation assumes the (approximate) form

\[
\phi_0(k_0) \simeq g^2 \int_{-\mu}^{0} \frac{d\epsilon}{\epsilon} \phi_0(\epsilon) \frac{1}{2} \ln \left( \frac{\mu^2}{\epsilon^2 - k_0^2} \right) .
\]

Again, as in (3), there is an additional logarithm, representing collinear enhancement. In QCD it arises from the exchange of ultrasoft, magnetic gluons. While the collinear enhancement in Eq. (5) is cut off by the mass of the scalar boson, \( M \), here it is cut off by the energy of the magnetic gluon, \( p_0 = \epsilon \pm k_0 \). For \( \epsilon \sim k_0 \sim \phi_0 \), \( p_0 \) is of the order of \( \phi_0 \), too, while in weak coupling, \( M \sim g\mu \gg \phi_0 \). For magnetic gluon exchange in QCD, the contribution of the collinear region, \( q \simeq k \), to the gap integral is therefore much larger than in the case of massive boson exchange.

To estimate the effect of the logarithm on the solution of the gap equation, let us neglect the energy dependence of the gap function in the integrand and consider its value at the Fermi surface, \( k = \mu, \ k_0 = \phi_0 \). One may also make the approximation

\[
\ln \left( \frac{\mu^2}{\epsilon^2 - \phi_0^2} \right) \simeq 2 \ln(\mu/\epsilon) .
\]

Then, the integral is again exactly solvable, with the (order of magnitude) result

\[
\phi_0 \simeq 2\mu \exp(-1/g) .
\]

Due to the explicit \( \epsilon \) dependence of the logarithm in (3), the power of \( g \) in the exponent is reduced as compared to the BCS result (3). The case of
massive scalar boson exchange, Eq. (6), interpolates between these two cases, as $g^2 \ll g^2 \ln(1/g) \ll g$ for $g \ll 1$. This reflects the fact mentioned earlier that, while the collinear singularity $q = k$ in (5) is cut off by the mass of the scalar boson, $M \sim g\mu$, the singularity in (7) is cut off by the gluon energy, $p_0 \sim \phi_0$, which is, in weak coupling, much smaller than $M$. In the literature, the parametric dependence on $g$ of the solution (8) to the gap equation (7) was first discussed in Refs. 14, 21, 24. As of today, various refinements of the solution (8) have been discussed 12, 15, 23, 24, 27. The value of the gap function at the Fermi surface, $k = \mu$, $k_0 = \phi_0$, is

$$\phi_0 = b \mu g^{-5} \exp \left( \frac{-c}{g} \right) \left[ 1 + O(g) \right], \quad c = \frac{3\pi^2}{\sqrt{2}}. \quad (9)$$

Furthermore, the gap function has a non-trivial (on-shell) energy dependence,

$$\phi_0(\epsilon) = \phi_0 \sin \left[ \frac{\pi}{2} \frac{g}{c} \ln \left( \frac{b \mu}{f(\epsilon)} \right) \right]. \quad (10)$$

The constant $c$ was first computed by Son 24. To obtain the correct numerical value for $c$, one has to account for the modifications of the gluon propagator in the presence of a dense medium 40. Then, what dominates the gap equation (7) and determines $c$ is the contribution from nearly static, Landau-damped magnetic gluons. Furthermore, Son showed that in computing $c$, it is essential to retain the energy dependence (10) of the gap function. To the level of accuracy considered by Son 24, $b = 1$ and $f(\epsilon) \equiv \epsilon$. The prefactor $g^{-5}$ arises from subleading contributions of static electric and non-static magnetic gluons.

The constant $b$ collects constant factors in these subleading contributions. It was first computed by Schäfer and Wilczek 12 and the present authors 15,

$$b = 512 \pi^4 \left( \frac{2}{N_f} \right)^{5/2} b', \quad (11)$$

with an undetermined constant $b'$. Schäfer and Wilczek 12 obtained $b' = 1/2$ and, as before, $f(\epsilon) \equiv \epsilon$. We showed 15 that actually $b' = 1$ and $f(\epsilon) \equiv \epsilon - \epsilon_0$. The additional factor 2 is the same that occurs in Eqs. (4), (6), and (8). It arises from the measure of integration in the gap equation, $d\epsilon_0/\epsilon \equiv -d \ln(\epsilon - \epsilon_0)$, and $-\ln(\epsilon - \epsilon_0)|_{-\mu}^{\phi_0} \simeq \ln(2\mu/\phi_0)$. This also explains the modification of $f(\epsilon)$. In addition, we pointed out 15 that it is important to compute the gap function on the correct quasiparticle mass shell, and we considered for the first time the case of non-zero temperature. As will be discussed in more detail below, the
Brown, Liu, and Ren computed $T_c$ from the quark-quark scattering amplitude. Using the result $T_c/\phi_0 = e^\gamma/\pi \simeq 0.567$, where $\gamma \simeq 0.577$ is the Euler–Mascheroni constant, they obtain

$$b' = \exp \left( -\frac{\pi^2 + 4}{8} \right) \simeq 0.176 ,$$

where the correction to the previous estimate $b' = 1$ arises from a finite, $\mu$ dependent contribution to the wavefunction renormalization for quarks in a dense medium. The authors of Ref. 27 also assert that there are no further corrections to $b'$ at this order in $g$. If correct, this is a remarkable result, because from previous calculations it appeared that computing $b'$ exactly to leading order in $g$ would be a formidable task.

Unlike calculations of the free energy or the Debye mass, where the perturbative expansion in powers of $g^2$ appears to be well-behaved even when extrapolated down to moderate values of $\mu$, this result for the wavefunction
renormalization (which is equivalent to non-Fermi liquid behavior) indicates that perturbation theory is not such a good approximation, at least for quarks near the Fermi surface. The confirmation of the results of Ref. 27 is clearly an outstanding problem for the field.

Figure 1 shows the value of the zero-temperature gap $\phi_0$ at the Fermi surface according to Eq. (9) as a function of $\mu$ for $N_f = 2$ and 3 massless quark flavors with $b' = 1$ (full and dotted lines), and for $N_f = 3$ flavors with $b' = \exp[-(\pi^2 + 4)/8]$ (dashed line). The running of the coupling $g(\mu)$ with the chemical potential $\mu$ was computed from the 3-loop QCD $\beta$ function, however, not for 6 but only for 3 flavors of massless quarks. Therefore, the QCD scale $\Lambda = 364$ MeV is chosen somewhat larger than the standard value, to give the value $\alpha_s(\mu = 2$ GeV) $\simeq 0.309$. Although an extrapolation of the weak-coupling result (9) to large $g$ (small $\mu$) appears audacious, it is interesting to note that the maximum value of the gap for $b' = 1$ is of the order of 100 MeV, quite in agreement with the earlier estimates of Refs. 6, 9. However, taking into account the quark wavefunction renormalization reduces the gap to values of a few MeV. These values are of the order of typical superfluid gaps in ordinary hadronic matter. This lends credibility to the conjecture that quark and hadronic matter are continuously connected, although symmetry arguments suggest that, at zero temperature, there is a first order phase transition between these two phases of nuclear matter.

3 Not so cold, dense quark matter

To understand how the color-superconducting gap changes with temperature, it is instructive to first consider the simpler BCS case. At nonzero temperature $T$, the gap equation (9) becomes

$$\phi = G^2 \int_{-\mu}^{0} \frac{d\epsilon}{\epsilon(\phi)} \phi \tanh \left(\frac{\epsilon(\phi)}{2T}\right).$$

(13)

Here, $\phi$ is the value of the gap at temperature $T$, $\phi \equiv \phi(T)$, and as before, $\phi_0 \equiv \phi(0)$ denotes the zero-temperature gap in the following. The $\phi$ dependence of the quasiparticle excitation energy $\epsilon(\phi) \equiv \sqrt{\epsilon_0^2 + \phi^2}$ has been made explicit, to distinguish $\epsilon(\phi)$ from $\epsilon \equiv \epsilon(\phi_0)$ used previously.

Again, $\phi$ is constant and can be divided out on both sides of Eq. (13), with the result

$$1 = G^2 \int_{-\mu}^{0} \frac{d\epsilon}{\epsilon(\phi)} \tanh \left(\frac{\epsilon(\phi)}{2T}\right).$$

(14)

The effect of the tanh is to reduce the value of the integrand, such that $\phi$ in the factor $1/\epsilon(\phi)$ has to decrease in order to balance the 1 on the left-hand
side. At some critical temperature $T_c$, this balance can no longer be achieved and $\phi = 0$ is the only solution of Eq. (13). $T_c$ is the temperature where the superconducting condensate melts. Physically, the random thermal energy of the fermions exceeds their binding energy in a Cooper pair. Thus, $T_c$ must be of the same order as $\phi_0$.

There is an easy way to compute the change of $\phi$ with $T$, which was to our knowledge first suggested in Ref. 15. Note that \[ \tanh \left( \frac{\epsilon(\phi) + \epsilon(\phi)}{2T} \right) \simeq 1 \] far from the Fermi surface, $\epsilon(\phi) \gtrsim |\epsilon_0| \gg \phi_0 \sim T$. A nonzero temperature influences the integrand only in the region close to the Fermi surface, $|\epsilon_0| \leq \phi_0 \sim T$. Let us therefore divide the range of integration into two parts, $0 \geq \epsilon_0 \geq -\kappa \phi_0$, and $-\kappa \phi_0 \geq \epsilon_0 \geq -\mu$, where $\kappa \gg 1$. Then, the tanh need only be kept in the first region, and the integral over the second region can be performed similarly as at zero temperature,

\[ 1 \simeq G^2 \int_{-\kappa \phi_0}^{0} \frac{d\epsilon_0}{\epsilon(\phi)} \tanh \left( \frac{\epsilon(\phi)}{2T} \right) + G^2 \ln \left( \frac{\mu}{\kappa \phi_0} \right) . \]  

Using the solution (4) for the zero-temperature gap $\phi_0$, the second term becomes $1 - G^2 \ln(2\kappa)$. Apparently, the 1 on the left-hand side is almost completely saturated by this second term. Cancelling the 1 and writing \[ \ln(2\kappa) \equiv \int_{-\kappa \phi_0}^{0} \frac{d\epsilon_0}{\epsilon(\phi)} \left( \frac{1}{\epsilon(\phi)} \tanh \left( \frac{\epsilon(\phi)}{2T} \right) - \frac{1}{\epsilon(\phi_0)} \right) = 0 . \]  

The dependence on $\kappa$ is spurious: one might as well send $\kappa \rightarrow \infty$ because the integrand vanishes when $\kappa \gg 1$. Equation (16) determines $\phi(T)/\phi_0$ as a function of $T$. In particular, at $T_c$, where $\phi = 0$, one derives the well-known result

\[ \frac{T_c}{\phi_0} = \frac{e}{\pi} \simeq 0.567 \]  

mentioned above.

In QCD, it turns out that the effect of temperature on the gap equation is essentially identical to that in BCS theory: the integrand in (9) is multiplied with $\tanh[\epsilon(\phi)/2T]$. For the same reasons as in the BCS case, this factor is negligible far away from the Fermi surface. One may again divide the range of integration into two parts, and neglect the effects of temperature in the one far from the Fermi surface. The integral over this region can be computed as for $T = 0$. Quite similarly to the treatment in the BCS case, it is found to saturate the left-hand side of the gap equation up to corrections of order $O(g)$.
Therefore, the integral over the region close to the Fermi surface must also be of order $O(g)$, in order to cancel these corrections. To see this, it is permissible to compute this integral to leading order in $g$. One then derives the same condition as in the BCS case, except that $G^2$ is replaced by $g$, see Ref.\[ for details. Consequently, to leading order in $g$, the $T$ dependence of the gap at the Fermi surface, normalized to the zero-temperature gap, $\phi(T)/\phi_0$, is the same as in BCS theory. In particular, the ratio $T_c/\phi_0$ is again given by Eq. (17). In retrospect, this is not surprising: the prefactor $b$ of the zero-temperature gap, Eq. (11), was seen to be determined by subleading terms in the gap equation [terms of order $O(g)$ relative to the leading terms due to Landau-damped magnetic gluons]. As explained above, temperature affects the gap equation at the same subleading order.

An immediate consequence is that when multiplying the ordinate of Fig. 1 with 0.567, one obtains the location of the phase transition to the color-superconducting phase in the $T-\mu$ phase diagram of nuclear matter. In the 2-flavor case without wavefunction corrections, the transition temperature is of order $\sim 100$ MeV. The color-superconducting phase could then be accessible in heavy-ion collisions at BNL–AGS or GSI–SIS energies, which explore the range of moderate temperatures and high (net) baryon density in the nuclear matter phase diagram. However, in the 3-flavor case, including the effects of the quark wavefunction renormalization, the transition temperature is at most $\sim 6$ MeV. For such small temperatures, color superconductivity occurs at best in neutron star cores, if they consist of quark matter.

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