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Why Color-Flavor Locking is Just like Chiral Symmetry Breaking

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Abstract

We review how a classification into representations of color and flavor can be used to understand the possible patterns of symmetry breaking for color superconductivity in dense quark matter. In particular, we show how for three flavors, color-flavor locking is *precisely* analogous to the usual pattern of chiral symmetry breaking in the QCD vacuum.

While it has been known from the work of Bailin and Love that color superconductivity can occur in cold, dense quark matter [1], only recently has there been sustained interest, following work of Alford, Rajagopal, and Wilczek, and of Rapp, Schäfer, Shuryak, and Velkovsky [2–11]. Theories used to estimate the gap for color superconductivity include Nambu–Jona-Lasinio (NJL) models and the renormalization group. While Bailin and Love found tiny gaps, ~ 1 MeV, current estimates of the gap are much larger, ~ 100 MeV [2–11].

Color superconductivity is not like "ordinary" superconductivity, such as in the theory of Bardeen, Cooper, and Schrieffer (BCS). In the latter, the gap is an exponential in $1/g^2$. This result was also obtained by Bailin and Love [1], who assumed that, at zero temperature, static magnetic interactions are screened by a (non-perturbative) magnetic mass. In QCD, however, Son [7] and ourselves [6] argued independently that nearly static magnetic interactions are not screened; if g is the QCD coupling constant, collinear emission of soft magnetic gluons turns the gap from an exponential in $1/g^2$ into an exponential in $1/g$. For sufficiently small g, $\exp(-1/g)$ is much larger than $\exp(-1/g^2)$. This is one of the reasons why the gap is now found to be two orders of magnitude larger than the previous estimate by Bailin and Love.

The other reason is the following. When one explicitly computes the magnitude of the gap for a spin-zero condensate, ϕ_0, one finds

$$\frac{\phi_0}{\mu} =: b_0 \, \exp\left(-\frac{c_0}{g}\right) [1 + O(g)] \, ,$$

$$b_0 = 256 \, \pi^4 \left(\frac{2}{g^2 N_f}\right)^{5/2} b_0' \quad , \quad c_0 = \sqrt{\frac{6 N_c \pi^4}{N_c + 1}} \, . \tag{1}$$

This result is valid for N_f flavors of massless quarks at a chemical potential μ, coupled to a $SU(N_c)$ gauge theory. The value of c_0 and the prefactor of $1/g^5$ was computed first by Son [7]; the constant $256\,\pi^4$ is due to Schäfer and Wilczek [10] and ourselves [11]. Actually, this number is only an *estimate* for the prefactor; these are the factors from logarithms in the exponent which turn into a prefactor of $1/g^5$. The uncertainty from other effects which also contribute to the prefactor is represented by the number b_0'.

For a BCS gap, $\phi \sim \exp(-1/g^2)$, as a function of g the gap is usually infintesimal until it turns on suddenly. In contrast, because of the prefactor of $1/g^5$ in (1), as g increases from zero, ϕ_0/μ increases, reaches a maximum, and then *decreases* at strong g.

(Incidentally, this is just like a semiclassical tunneling probability, in which there are five zero modes. Why the color superconducting gap looks like a tunneling problem, with *five* zero modes, *independent* of the number of colors or flavors, is a complete mystery to us.)

Setting the undetermined constant $b_0' = 1$, what is so exciting is that because of the factor of $256\,\pi^4$, this maximum value of the gap is really quite large; for $N_f = 2$, the maximum is $\phi_0/\mu \simeq 0.13$ and occurs at a coupling of $g \simeq 4$, where $\alpha_s = g^2/4\pi \sim 1$. This large value of the gap supports previous estimates found using NJL models [2–5].

(Notice also that the gap *decreases* as N_f *increases*, like $1/N_f^{5/2}$, assuming (!) that b_0' does not depend significantly on N_f. The implications of this for quark matter, where there are up, down, and strange quarks, is intriguing.)

We also showed that the ratio of the transition temperature to the gap is *twice* the same ratio in BCS, with $T_c/\phi_0 \sim 1.13$ [11]. This establishes conclusively that color superconductivity is *not* like BCS.

In this talk we ignore the magnitude of the gap, to concentrate on the possible patterns of symmetry breaking. This has been analyzed using three colors, and either two or three flavors, using NJL models [4,5]. In [6] we developed a general classification in terms of representations in color and flavor. In this Proceeding we show that there is an exact analogy between color-flavor locking for three flavors, and chiral symmetry breaking in the QCD vacuum. This was noted in our Letter [6], in refs. 19, 21, and 23, although perhaps the present relaxed discussion might be more transparent than those cryptic comments.

Our analysis is motivated by the following puzzle. Scattering between two quarks occurs in one of two possible ways, either anti-symmetric or symmetric in the two color indices for quarks in the fundamental representation. (The gluon interaction is flavor blind, but flavor will enter later through Fermi statistics.) In color $SU(3)_c$, for example, the anti-symmetric representation is a color anti-triplet, $\overline{\mathbf{3}}$, and the symmetric is a color sextet, $\mathbf{6}$. Then in all models studied, either NJL or with single-gluon exchange, the color anti-triplet channel is attractive, and the color sextet channel is repulsive.

(We comment that this need *not* be true for all values of the coupling. Since QCD is asymptotically free, the color anti-triplet channel must be the attractive channel for all densities above some value. But in strong coupling, which corresponds to intermediate densities, it is certainly possible that the attractive channel is the color sextet. The phenomenology of such an intermediate phase would be interesting to explore. We ignore this possibility for now, although it could be easily treated using our approach.)

For total spin $J = 0$, the quark-quark condensate has the form

$$\phi^{ij}_{ab}(\Gamma) \; = \; q^i_a(-\vec{p})^T \, C \, \Gamma \, q^j_b(\vec{p}) \; , \tag{2}$$

q is the quark field, q^T the Dirac transpose, C the charge conjugation matrix, and Γ a Dirac matrix. The $SU(3)_c$ color indices for a quark in the fundamental representation are $i, j = 1, 2, 3$; the flavor indices are $a, b = 1, \ldots, N_f$, for N_f flavors of massless quarks. The Dirac matrix Γ classifies the chiral and helicity structure of the condensate. We don't worry about the four kinds of condensates [6,8] which Γ represents, since all that does is to add another index to the condensate field.

For two flavors, the condensate found is of the form

$$\langle \phi^{ij}_{ab} \rangle = \epsilon^{ijk} \, \epsilon_{ab} \, \langle \phi^k \rangle \; . \tag{3}$$

This is anti-symmetric in the quark color indices, i and j, and so only involves the color anti-triplet channel, as expected. We are sloppy about indices: ϕ^k is an anti-triplet, not a triplet. It is also a flavor singlet.

For three flavors, one would also expect that the condensate is again a color anti-triplet,

$$\langle \phi^{ij}_{ab} \rangle = \epsilon^{ijk} \, \epsilon_{abc} \, \langle \phi^k_c \rangle \; , \tag{4}$$

and a flavor anti-triplet.

The puzzle is that this is *not* what is found. Instead, the "color-flavor locked" condensate is of the form [4,5]

$$\langle \phi^{ij}_{ab} \rangle = \delta^i_a \, \delta^j_b \, \kappa_1 + \delta^i_b \, \delta^j_a \, \kappa_2 \; . \tag{5}$$

This contains a piece which is anti-symmetric in the color indices, and so is a color anti-triplet, but it also contains a piece which is symmetric, and so is a color sextet. The color anti-triplet piece is proportional to $\kappa_1 - \kappa_2$, the color sextet to $\kappa_1 + \kappa_2$. The color-flavor locked condensates [4,5] are all predominantly color triplet, but there is always some color sextet piece, $0 \neq |\kappa_1 + \kappa_2| \ll |\kappa_1 - \kappa_2|$. The quandry is then, why for three flavors is there some condensation in a *repulsive* channel? And why doesn't this happen for two flavors?

Crucial to the analysis is that for a spin-zero condensate of massless particles, Fermi statistics requires that the condensate field is symmetric in the simultaneous interchange of both color and flavor indices [1,6,8]:

$$\phi^{ij}_{ab}(\Gamma) \;=\; +\phi^{ji}_{ba}(\Gamma) \; . \tag{6}$$

We assume that because of instantons, we can ignore the difference between left- and right-handed chiral symmetries, and speak only of a (global) $SU(N_f)_f$ symmetry. Then we only need some trivial group theory: in $SU(2)$, $\mathbf{2} \times \mathbf{2} = \mathbf{1}_a + \mathbf{3}_s$ in $SU(2)$, while for $SU(3)$, $\mathbf{3} \times \mathbf{3} = \overline{\mathbf{3}}_a + \mathbf{6}_s$. Here "$a$" and "$s$" denote anti-symmetric and symmetric representations, respectively. By (6), then, we can have either representations in color and flavor which are both anti-symmetric, or both symmetric.

For two flavors, under $(SU(3)_c, SU(2)_f)$ we can have a color anti-triplet, flavor singlet, $(\overline{\mathbf{3}}, \mathbf{1})$, and a color sextet, flavor triplet $(\mathbf{6}, \mathbf{3})$. The $(\overline{\mathbf{3}}, \mathbf{1})$ field is the ϕ^k of (3).

For three flavors, under $(SU(3)_c, SU(3)_f)$ we can have a color anti-triplet, flavor anti-triplet field $(\overline{\mathbf{3}}, \overline{\mathbf{3}})$, as the ϕ^i_a field in (4), and a color sextet, flavor sextet field $(\mathbf{6}, \mathbf{6})$. The sextet field is denoted as χ^{ij}_{ab}; the latter is symmetric in each pair of indices, $\chi^{ij}_{ab} = +\chi^{ji}_{ab} = +\chi^{ij}_{ba}$.

Effective lagrangians are constructed in the usual manner. The kinetic terms include

$$\mathcal{L}_0 = \mathcal{L}_{\text{gauge}} + |D_\mu \phi|^2 + |D_\mu \chi|^2 \; . \tag{7}$$

Here the gauge action includes *both* the usual gauge action for a nonabelian gauge theory, *and* the contribution of "Hard Dense Loops" [6,7,10,11]. The kinetic terms for the condensate fields involve the appropriate covariant derivative for that representation. Thus if the covariant derivative for a triplet field is $D_\mu = \partial_\mu - igA_\mu$, for the anti-triplet field $D_\mu \phi = (\partial_\mu + igA_\mu)\phi$; the covariant derivative for the sextet field is more involved.

Mass terms for the condensate fields are written to drive condensation in the anti-triplet channel:

$$\mathcal{L}_2 = -\mu_3^2 \, |\phi|^2 + \mu_6^2 \, |\chi|^2 \, . \tag{8}$$

If mixing between ϕ and χ could be neglected, because of its negative mass squared term, only ϕ condenses, and χ not.

(Condensation in the sextet channel is modeled simply by taking a negative mass squared term for χ, and a positive one for ϕ.)

There are many quartic terms which couple ϕ and χ together:

$$\mathcal{L}_4 = \lambda_1 \, [\mathrm{tr}\,(\phi^\dagger \phi)]^2 + \lambda_2 \, \mathrm{tr}\,(\phi^\dagger \phi)^2 + \lambda_3 \, [\mathrm{tr}\,(\chi^\dagger \chi)]^2 + \lambda_4 \, \mathrm{tr}\,(\chi^\dagger \chi)^2$$
$$+ \lambda_5 \, \mathrm{tr}\,(\phi^\dagger \phi) \, \mathrm{tr}\,(\chi^\dagger \chi) + \ldots \tag{9}$$

plus other quartic terms, some of which mix ϕ and χ. We are opaque with respect to indices: $\mathrm{tr}\,(\phi^\dagger \phi) = (\phi^i_a)^* \phi^i_a$, while $\mathrm{tr}\,(\phi^\dagger \phi)^2 = (\phi^i_a)^* \phi^j_a (\phi^j_b)^* \phi^i_b$. Whatever the forms of the quartic terms, however, they are all rather innocuous, since they are at least quadratic in the χ field. Consequently, even if ϕ condenses, unless a coupling such as λ_5 is large and negative, it doesn't drive the condensation of χ. This is what is typical with two coupled fields: one field may condense, but generally the other doesn't, so we can ignore the field which doesn't.

For three flavors, however, there *are* operators which mix ϕ and χ in a special way. Consider

$$\mathrm{tr}\,(\phi \chi \phi) + \mathrm{h.c.} \equiv \phi^i_a \, (\chi^{ij}_{ab} + \chi^{ij}_{ba}) \, \phi^j_b + \mathrm{h.c.} \tag{10}$$

This is a cubic operator which is invariant under both $SU(3)_c$ color and $SU(3)_f$ flavor rotations. Our sloppiness in notation obscures this point: the indices on ϕ are anti-triplets, and so can be contracted with the triplet indices on χ.

Nevertheless, such an operator does not appear in an effective lagrangian, because it is not invariant under global $U(1)$ rotations of baryon number. This is easily corrected, however, by multiplying by the operator $\det(\phi)^*$; this is also invariant under $SU(3)_c$ color and $SU(3)_f$ flavor, but precisely soaks up the requisite factors for baryon number. Thus there is no symmetry reason to prevent the following term from appearing in the effective lagrangian:

$$\mathcal{L}_6 = \lambda_6 \, \det(\phi)^* \, \mathrm{tr}\,(\phi \chi \phi) + \mathrm{h.c.} \tag{11}$$

This is a six-point operator, and so in the sense of the renormalization group, is marginal relative to $\mathcal{L}_4$, and doesn't affect the critical behavior.

It does, however, affect the *pattern* of condensation. In particular, assume that ϕ^i_a condenses in such a manner that $\det(\phi) \neq 0$. This is true for color-flavor locking, where

$$\langle \phi^i_a \rangle = \delta^i_a \, \phi_0 \, . \tag{12}$$

Then our funny operator becomes

$$\mathcal{L}_6 = \lambda_6 \, \phi_0^5 \, (\chi_{ij}^{ij} + \chi_{ji}^{ij}) + \text{h.c.} \, . \tag{13}$$

Thus when ϕ condenses, $\mathcal{L}_6$ generates a term *linear* in χ. The special thing about an operator linear in a field is that even if the mass squared term is positive, a nonzero vacuum expectation value is always generated by a linear term. Thus even when $\mu_6^2 > 0$, χ acquires a nonzero expectation value from its mixing with ϕ,

$$\langle \chi_{ab}^{ij} \rangle \sim \lambda_6 \frac{\phi_0^5}{\mu_6^2} \, (\delta_a^i \delta_b^j + \delta_b^i \delta_a^j) \, . \tag{14}$$

There are of course many other terms which couple six fields together. None of these other terms, however, produce a term linear in χ when ϕ condenses.

(One can also write down a six-point term like that in (11), except with $\det(\phi)$ replaced by $\det(\chi)$. Even with such a term, however, if for intermediate densities condensation is for the sextet instead of the anti-triplet field, by switching the sign of the mass terms in (8), χ condensation does *not* drive ϕ condensation, since the six-point term like that in (11) is still quadratic in ϕ, not linear.)

It is now easy to understand why we don't have to worry about this kind of mixing for two flavors. When $N_f = 2$, the $(\overline{\mathbf{3}}, \mathbf{1})$ is ϕ^i, and the $(\mathbf{6}, \mathbf{3})$ field χ_{ab}^{ij}. There is simply no cubic coupling possible between two ϕ^i's and one χ_{ab}^{ij}; there is also no $\det(\phi)$ term. Thus when the color anti-triplet ϕ condenses, the color sextet χ does *not*. This is what detailed calculations in NJL models find [2,3].

The physically realistic case is $2+1$ flavors, with the up and down quarks essentially degenerate in mass, but the strange quark with a large mass. One can then show that $2+1$ flavors is analogous to three flavors, although the proliferation of fields becomes a little dizzying. For the condensates of up quarks with down quarks, denote the color anti-triplet field as ϕ^i, and the color sextet field as χ_{ab}^{ij}. For the condensates of up or down quarks with strange quarks, we write the color anti-triplet field as $\widetilde{\phi}_a^i$, and the color sextet field as $\widetilde{\chi}_a^{ij}$. Here the flavor index is only for up and down, $a, b = 1, 2$. Then there is an operator analogous to $\det(\phi)$; it is

$$\det(\phi) \rightarrow \epsilon^{ijk} \epsilon^{ab} \phi^i \widetilde{\phi}_a^j \widetilde{\phi}_b^k \, . \tag{15}$$

The existence of this operator is interesting in its own right. For three flavors, Schäfer and Wilczek [5] pointed out that $\det(\phi)$ can be used as a global order parameter to distinguish different phases. The operator of (15) is the analogy for $2+1$ flavors.

For $2+1$ flavors, the operator similar to (11) is

$$\widetilde{\mathcal{L}}_6 = \epsilon^{ijk}\epsilon^{ab}\phi^i\widetilde{\phi}^j_a\widetilde{\phi}^k_b\ \left(\lambda_6'\ \widetilde{\phi}^l_c\ \chi^{lm}_{cd}\ \widetilde{\phi}^m_d + \lambda_6''\ \phi^l\ \widetilde{\chi}^{lm}_c\ \widetilde{\phi}^m_c\right) + \ldots \qquad (16)$$

Because of this operator, when *both* ϕ^i and $\widetilde{\phi}^i_a$ condense, so will the two color sextet fields, χ^{ij}_{ab} and $\widetilde{\chi}^{ij}_a$.

This conclusion agrees with all previous work done using NJL models [5]. It is not apparent from their analysis, which did not use our color and flavor decomposition, but it does explain why, in terms of (5), $\kappa_1 + \kappa_2 \neq 0$. It also agrees with Alford, Berges, and Rajagopal [5], who find that for $2+1$ flavors, the color sextet field condenses *only* along the direction in which the color anti-triplet field condenses. This is what one expects from the detailed form of the cubic coupling in (16). One doesn't see the color sextet field in weak coupling [7,9–11], since it is of fifth order in the color anti-triplet condensate.

Now we are ready to understand the analogy to chiral symmetry breaking. In our analysis, the admixture of a color symmetric piece in the condensate of (5) is really a red herring; what matters is the condensation of ϕ^i_a in (12). Once we recognize that we can basically forget about the color sextet field, and concentrate on the color anti-triplet, we see that the pattern of symmetry breaking is most familiar. Consider not $SU(3)_c$ color and $SU(3)_f$ flavor, but chiral symmetry breaking, where a global flavor symmetry of $SU(3)_\ell \times SU(3)_r$ breaks to $SU(3)_f$. The order parameter is again a 3×3 matrix, a $(\overline{\mathbf{3}}, \mathbf{3})$ under $(SU(3)_\ell, SU(3)_r)$ flavor:

$$U^{ab} = \overline{q}^i_{\ell,a}\ q^i_{r,b} \qquad (17)$$

so a and b are triplet indices for the flavor symmetries of $SU(3)_\ell$ and $SU(3)_r$, respectively. The flavor matrix U^{ab} is of course a color singlet.

(While it does not happen in the vacuum, in a phase in which color superconductivity occurs, if chiral symmetry is broken in the color singlet channel, it may well condense in a color octet channel as well. To be precise, one should classify the chiral symmetry with respect to the unbroken color/flavor group.)

Then in QCD, the observed pattern of chiral symmetry breaking is *identical* to that of (12). For chiral symmetry breaking, the condensate like that of (12) is

$$\langle U^{ab} \rangle = \delta^{ab}\, U_0\ , \qquad (18)$$

and corresponds to *all* quark flavors developing the *same* constituent quark mass. It is crucial to drop the color sextet piece from the condensate in order to understand this analogy; it is not at all obvious from (5).

This is not the only way in which the symmetry may break; instead, only *one* flavor of quark might develop a constituent quark mass, $\langle U^{ab} \rangle \sim \delta^a_1 \delta^b_1$. (These are the only two possibilities [6].) For chiral symmetry breaking in QCD, nature obviously favors a common constituent quark mass for all flavors over just one flavor developing a quark mass. This was proven by Coleman and Witten to be true in the limit of a large number of colors, $N_c \to \infty$ [12].

For color superconductivity, the pattern of symmetry breaking analogous to one flavor of quark developing a constituent quark mass is

$$\langle \phi^i_a \rangle = \delta^{i1}\, \delta_{a1}\, \phi'_0 \,. \qquad (19)$$

It is easy to argue why the symmetry breaking in (12) is favored over that in (19): (12) ensures that all colors develop a gap, while (19) leaves most colors ungapped. Gapping all colors is energetically favorable.

We conclude by making a simple point about the effects of terms in the effective lagrangian for nonzero current quark masses. For chiral symmetry, these terms are linear in the U field,

$$\mathcal{L}_{\text{quark mass}} \sim \text{tr}\left[m_q(U + U^\dagger)\right] . \qquad (20)$$

Consequently, the pseudo-Goldstone bosons, such as pions, *etc.*, have a mass squared which is *linear* in the quark mass, $m_\pi^2 \sim m_q$.

A term linear in $\text{tr}\,(U)$ is possible because we only need to preserve the flavor symmetry of $SU(N_f)_f$. In contrast, mass terms are very different for color superconductivity, since we still have to preserve the invariance under color and flavor. For example, if we were to add a term for the strange quark mass to a three flavor lagrangian, it would be *quadratic* in the strange quark mass:

$$\mathcal{L}_{\text{quark mass}} \sim m_s^2\, (\phi^i_s)^* \phi^i_s \;, \qquad (21)$$

where m_s is the strange quark mass, and "s" denotes the flavor direction for the strange quark. This is rather heavy: this pseudo-Goldstone boson has a mass like the strange quark mass, ~ 100 MeV.

A more interesting example, which we do not have space to explain, arises when one considers the "instanton-free" region which arises at very high densities. In that case we do effectively restore the full flavor symmetry, and neglecting instanton and mass effects, parity is spontaneously broken [6,13]. Then we must introduce left- and right-handed condensate fields, and like (21), the mass term is

$$\mathcal{L}_{\text{quark mass}} \sim (m_u + m_d)^2 |\phi^i_{\ell,a} + \phi^i_{r,a}|^2 \;, \qquad (22)$$

where m_u and m_d are the up and down quark masses. As in (21), the pseudo-Goldstone boson for the breaking of parity has a mass which is *linear* in the quark mass. But now we are dealing with the up and down quark masses, which are *very* light, ~ 5 or 10 MeV; thus the mass for the pseudo-Goldstone boson is also of this order.

In heavy-ion collisions, this implies that if one reaches a phase in which cool, dense quark matter is produced in an instanton-free regime, then the time scale for a parity violating condensate to relax is not ~ 1 fm, as it is with (say) disoriented chiral condensates, but a factor of ten or twenty larger, $\sim 10-20$ fm/c. This might produce observable signals, if one is careful to trigger on collisions in which the produced quark matter is not hot, but cool.

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