

Color Superconductivity: CFL and 2SC phases

- Introduction
- Hierarchies of effective lagrangians
- Effective theory at the Fermi surface (HDET)
- Symmetries of the superconductive phases

Introduction

- Ideas about CS back in 1975 (Collins & Perry-1975, Barrois-1977, Frautschi-1978).
- Only in 1998 (Alford, Rajagopal & Wilczek; Rapp, Schafer, Schuryak & Velkovsky) a real progress.
- The phase structure of QCD at high-density depends on the number of flavors with mass $m < \mu$.
- Two most interesting cases: $N_f = 2, 3$.
- Due to asymptotic freedom quarks are almost free at high density and we expect difermion condensation in the color channel 3^* .

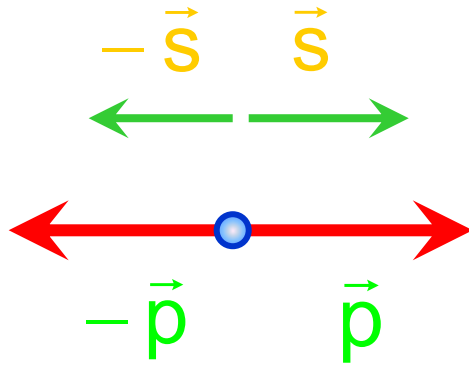
Consider the possible pairings at very high density

$$\langle 0 | \psi_{ia}^\alpha \psi_{jb}^\beta | 0 \rangle \quad \alpha, \beta \text{ color}; \quad i, j \text{ flavor}; \quad a, b \text{ spin}$$

- Antisymmetry in spin (a,b) for better use of the Fermi surface
- Antisymmetry in color (α, β) for attraction
- Antisymmetry in flavor (i,j) for Pauli principle

Only possible pairings

LL and RR



For $\mu \gg m_u, m_d, m_s$

Favorite state for $N_f = 3$, **CFL** (color-flavor locking) (Alford, Rajagopal & Wilczek 1999)

$$\langle 0 | \Psi_{iL}^\alpha \Psi_{jL}^\beta | 0 \rangle = -\langle 0 | \Psi_{iR}^\alpha \Psi_{jR}^\beta | 0 \rangle \propto \Delta \varepsilon^{\alpha\beta C} \varepsilon_{ijC}$$

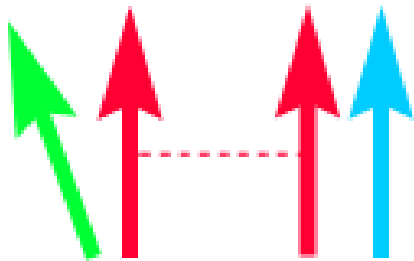
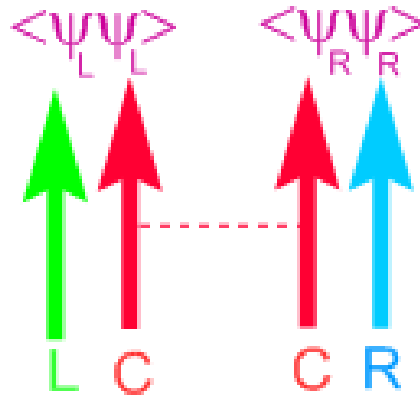
Symmetry breaking pattern

$$SU(3)_c \otimes SU(3)_L \otimes SU(3)_R \Rightarrow SU(3)_{c+L+R}$$

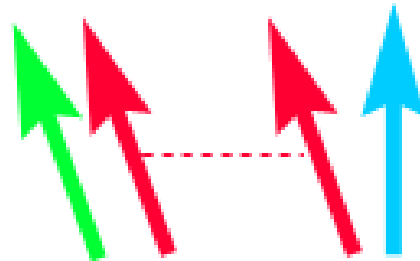
Why CFL?

$$\langle 0 | \Psi_{iL}^\alpha \Psi_{jL}^\beta | 0 \rangle \propto \Delta \varepsilon^{\alpha\beta C} \varepsilon_{ijC}$$

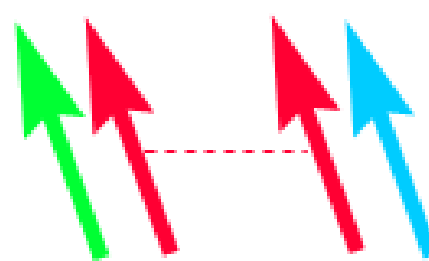
$$\langle 0 | \Psi_{iR}^\alpha \Psi_{jR}^\beta | 0 \rangle \propto -\Delta \varepsilon^{\alpha\beta C} \varepsilon_{ijC}$$



rotate left



rotate color



rotate right

What happens going down with μ ? If $\mu \ll m_s$, we get

3 colors and 2 flavors (2SC)

$$\langle 0 | \Psi_{iL}^\alpha \Psi_{jL}^\beta | 0 \rangle = \Delta \varepsilon^{\alpha\beta 3} \varepsilon_{ij}$$

$$SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \Rightarrow SU(2)_c \otimes SU(2)_L \otimes SU(2)_R$$

However, if μ is in the intermediate region we face a situation with fermions having different Fermi surfaces (see later). Then other phases could be important (LOFF, etc.)

Difficulties with lattice calculations

- Define euclidean variables: $x_0 \rightarrow -ix_E^4$, $x^i \rightarrow x_E^i$
 $\gamma_0 \rightarrow \gamma_E^4$, $\gamma^i \rightarrow -i\gamma_E^i$

- Dirac operator with chemical potential

$$D(\mu) = \gamma_E^\mu D_E^\mu + \mu \gamma_E^4$$

- At $\mu = 0$ $D(0)^\dagger = -D(0)$, $\gamma_5 D(0) \gamma_5 = -D(0)$

- Eigenvalues of $D(0)$ pure imaginary

- If $|\lambda\rangle$ eigenvector of $D(0)$, $\gamma_5 |\lambda\rangle$ eigenvector with eigenvalue $-\lambda$

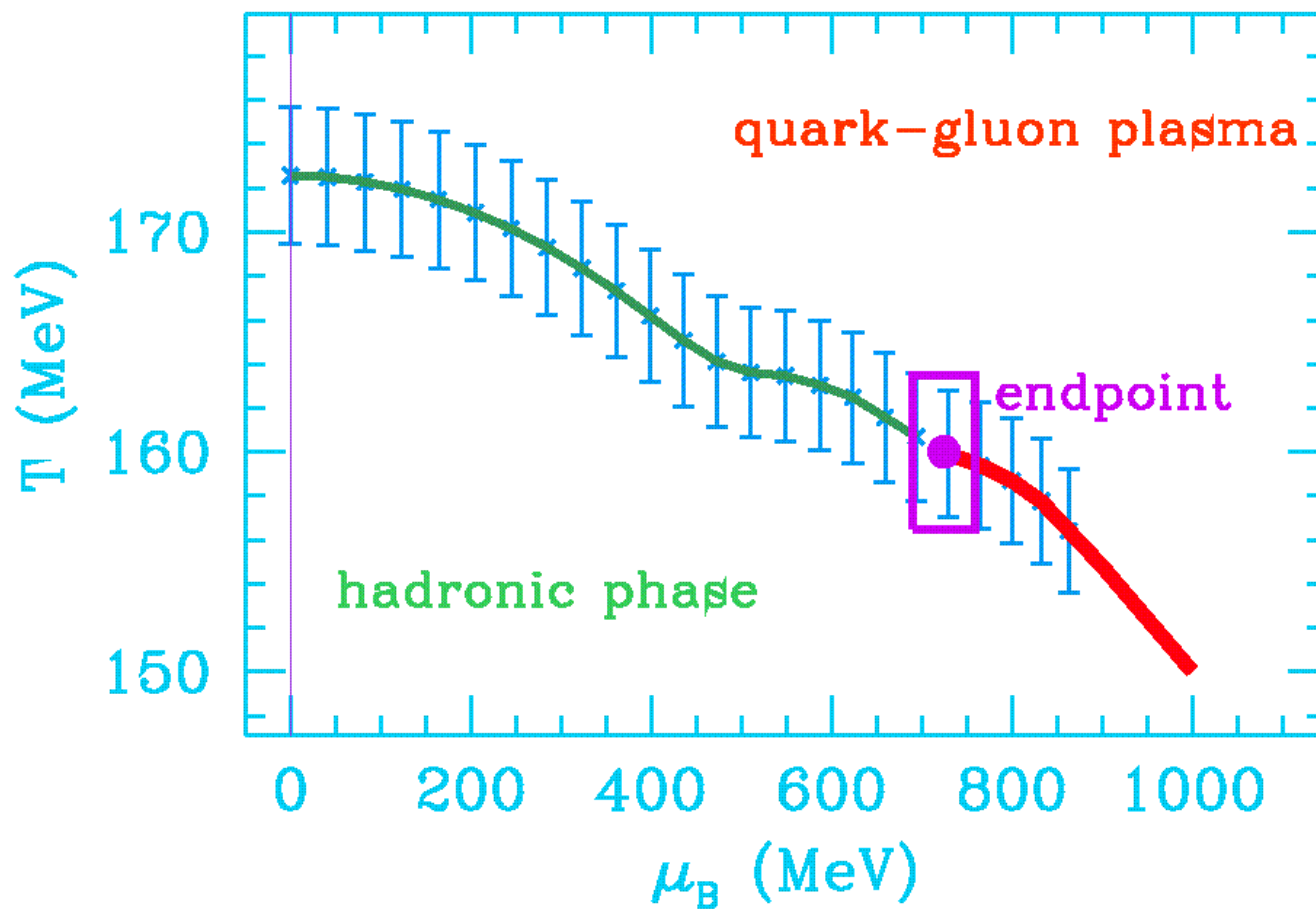




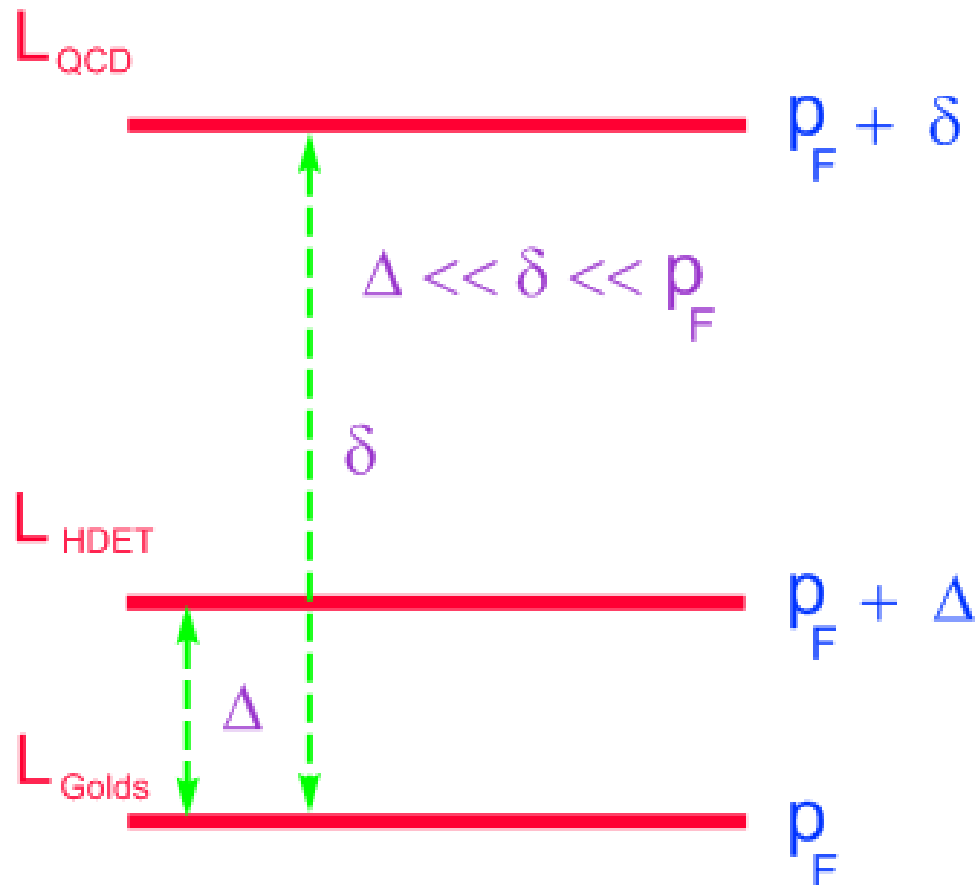
$$\det[D(0)] = \prod (\lambda)(-\lambda) > 0$$

For μ not zero the argument does not apply and one cannot use the sampling method for evaluating the determinant. However for isospin chemical potential and two degenerate flavors one can still prove the positivity.

For finite baryon density no lattice calculation available except for small μ and close to the critical line (Fodor and Katz)



Hierarchies of effective lagrangians



Integrating out heavy degrees of freedom we have two scales. The gap Δ and a cutoff, δ above which we integrate out. Therefore:

two different effective theories, L_{HDET} and L_{Golds}

- L_{HDET} is the effective theory of the fermions close to the Fermi surface. It corresponds to the Polchinski description. The condensation is taken into account by the introduction of a mean field corresponding to a Majorana mass. The d.o.f. are quasi-particles, holes and gauge fields. This holds for energies up to the cutoff.
- L_{Golds} describes the low energy modes ($E \ll \Delta$), as Goldstone bosons, ungapped fermions and holes and massless gauge fields, depending on the breaking scheme.

Effective theory at the Fermi surface (HDET)

Starting point: \mathcal{L}_{QCD}

$$\mathcal{L}_{\text{QCD}} = \bar{\psi} i \not{D} \psi - \frac{1}{4} F_{\mu\nu}^a F^{a\mu\nu} + \mu \bar{\psi} \gamma_0 \psi, \quad a = 1, \dots, 8$$

$$D_\mu = \partial_\mu + i g_s A_\mu^a T^a, \quad \not{D} = \gamma_\mu D^\mu, \quad T^a = \frac{\lambda_a}{2}$$

at asymptotic $\mu \gg \Lambda_{\text{QCD}}$,

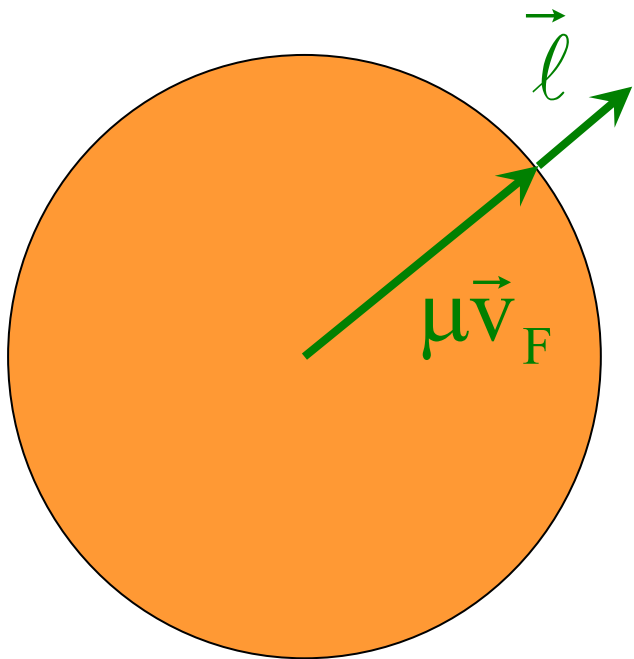
$$(\not{p} + \mu \gamma_0) \psi(p) = 0 \Rightarrow (p_0 + \mu) \psi(p) = \vec{\alpha} \cdot \vec{p} \psi(p)$$

$$(p^0 + \mu)^2 = |\vec{p}|^2 \Rightarrow p^0 = E_\pm = -\mu \pm |\vec{p}|$$

Introduce the projectors:

$$P_{\pm} = \frac{1 \pm \vec{\alpha} \cdot \vec{v}_F}{2}, \quad \vec{v}_F \equiv \vec{v} = \left. \frac{\partial E(\vec{p})}{\partial \vec{p}} \right|_{\vec{p}=\vec{p}_F} = \left. \frac{\partial |\vec{p}|}{\partial \vec{p}} \right|_{\vec{p}=\vec{p}_F} = \hat{p}$$

and decomposing: $\vec{p} = \mu \vec{v}_F + \vec{\ell}$



$$\begin{aligned} H\psi_+ &= \vec{\alpha} \cdot \vec{\ell} \psi_+ \\ H\psi_- &= (-2\mu + \vec{\alpha} \cdot \vec{\ell}) \psi_- \end{aligned}$$

- States ψ_+ close to the FS
- States ψ_- decouple for large μ

Field-theoretical version:

$$\psi(\mathbf{x}) = \int \frac{d^4 p}{(2\pi)^4} e^{-ip \cdot \mathbf{x}} \psi(p)$$

$$p^\mu = \mu v^\mu + l^\mu$$

$$v^\mu = (0, \vec{v}), |\vec{v}| = 1, \quad l^\mu = (l^0, \vec{l})$$

$$\vec{l} = \vec{v} l_{\parallel} + \vec{l}_{\perp}, \quad \vec{l}_{\perp} = \vec{l} - (\vec{l} \cdot \vec{v}) \vec{v}$$

4 d.o.f.

$$l_0, l_{\parallel}, \vec{v}$$

Choosing $\vec{v} \parallel \vec{p}$



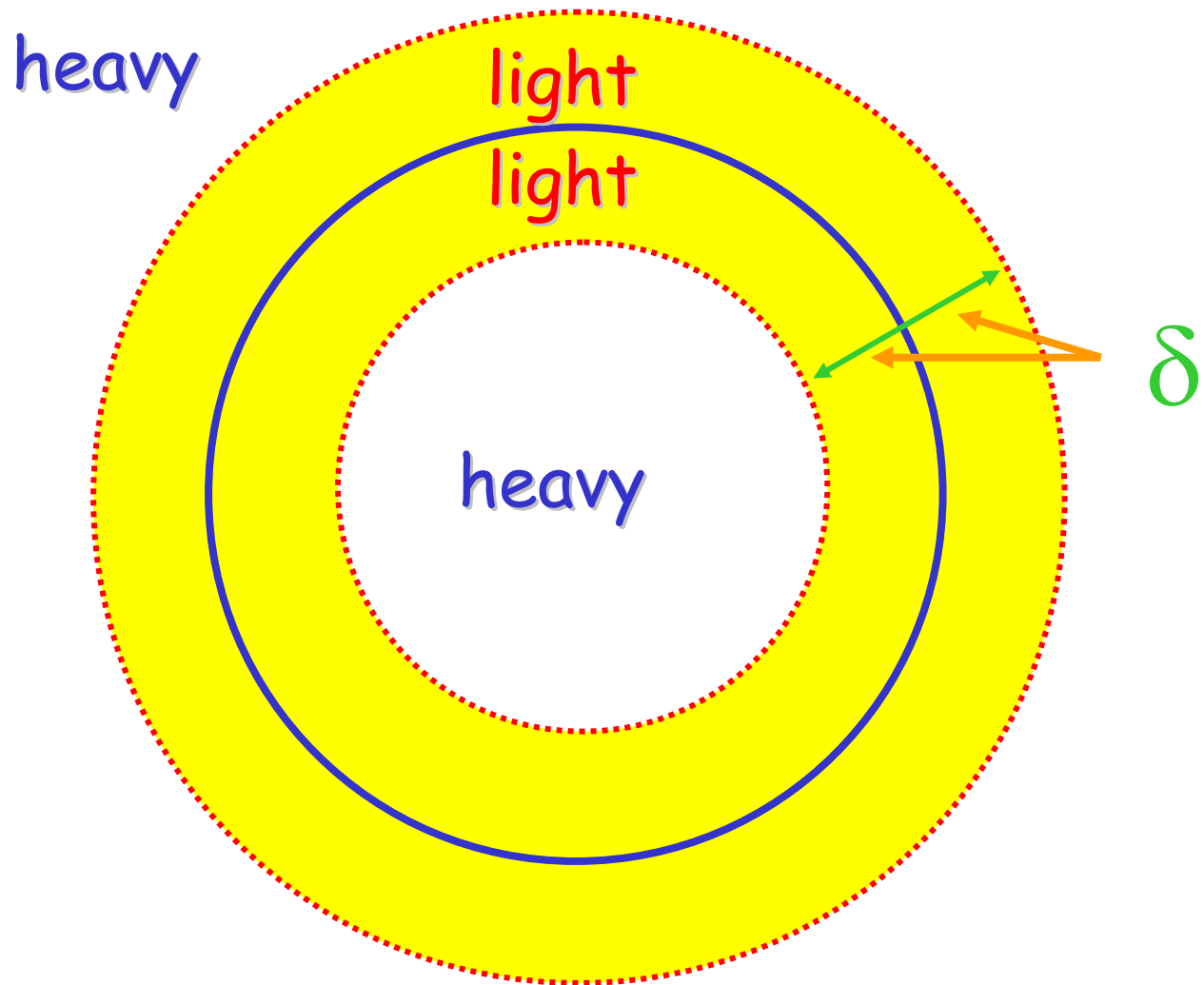
$$\vec{l}_{\perp} = 0$$

Separation of light and heavy d.o.f.

$$\text{light d.o.f.} \quad \mu - \delta \leq |\vec{p}| \leq \mu + \delta, \quad -\delta \leq l_{\parallel} \leq +\delta$$

$$\text{heavy d.o.f.} \quad |\vec{p}| \leq \mu - \delta, |\vec{p}| \geq \mu + \delta, \quad l_{\parallel} \leq -\delta, l_{\parallel} \geq +\delta$$

Separation of light and heavy d.o.f.



Momentum integration for the light fields

$$\int \frac{d^4 p}{(2\pi)^4} \Rightarrow \frac{\mu^2}{(2\pi)^4} \int d\Omega \int_{-\delta}^{+\delta} d\ell_{\parallel} \int_{-\infty}^{+\infty} d\ell_0 = \int \frac{d\vec{v}}{4\pi} \frac{\mu^2}{\pi} \int \frac{d^2 \ell}{(2\pi)^2}$$

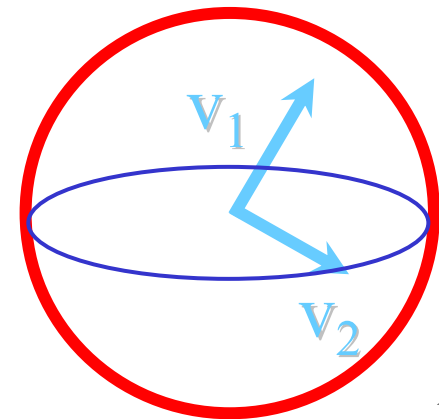
The Fourier decomposition becomes

$$\psi(\mathbf{x}) = \int \frac{d\vec{v}}{4\pi} e^{-i\mu\mathbf{v}\cdot\mathbf{x}} \psi_{\vec{v}}(\mathbf{x})$$

$$\psi_{\vec{v}}(\mathbf{x}) = \frac{\mu^2}{\pi} \int \frac{d^2 \ell}{(2\pi)^2} e^{-i\ell\cdot\mathbf{x}} \psi_{\vec{v}}(\ell)$$

$$(\psi_{\vec{v}}(\ell) \equiv \psi(p))$$

For any fixed \mathbf{v} , **2-dim theory**



In order to decouple the states corresponding to E_-

$$\psi(\mathbf{x}) = \int \frac{d\vec{v}}{4\pi} e^{-i\mu\vec{v}\cdot\mathbf{x}} [\psi_+(\mathbf{x}) + \psi_-(\mathbf{x})]$$

Momenta from the Fermi sphere

$$\psi_{\pm}(\mathbf{x}) \equiv P_{\pm} \psi_{\vec{v}}(\mathbf{x}) = P_{\pm} \frac{\mu^2}{\pi} \int \frac{d^2\ell}{(2\pi)^2} e^{-i\ell\cdot\mathbf{x}} \psi_{\vec{v}}(\ell)$$

substituting inside L_{QCD} and using

$$\int d^4x \psi^{\dagger}(\mathbf{x}) \psi(\mathbf{x}) \xrightarrow{\mu \rightarrow \infty} \frac{\mu^2}{\pi} \int \frac{d\vec{v}}{4\pi} \frac{d^2\ell}{(2\pi)^2} \psi_{\vec{v}}^{\dagger}(\mathbf{x}) \psi_{\vec{v}}(\mathbf{x})$$

$$\int d^4x \bar{\psi} (i\not{D} + \mu\gamma_0) \psi \rightarrow \int \frac{d\vec{v}}{4\pi} \left[\psi_+^\dagger iV \cdot D \psi_+ + \psi_-^\dagger (2\mu + i\tilde{V} \cdot D) \psi_- + (\bar{\psi}_+ i\not{D}_\perp \psi_- + \text{h.c.}) \right]$$

$$V^\mu = (1, \vec{v}), \quad \tilde{V}^\mu = (1, -\vec{v})$$

$$\gamma_{\parallel}^\mu = (\gamma^0, (\vec{v} \cdot \vec{\gamma}) \vec{v}), \quad \gamma_{\perp}^\mu = \gamma^\mu - \gamma_{\parallel}^\mu$$

$$\bar{\psi}_+ \gamma^\mu \psi_+ = V^\mu \psi_+^\dagger \psi_+$$

$$\bar{\psi}_- \gamma^\mu \psi_- = \tilde{V}^\mu \psi_-^\dagger \psi_-$$

$$\bar{\psi}_+ \gamma^\mu \psi_- = \bar{\psi}_+ \gamma_{\perp}^\mu \psi_-$$

$$\bar{\psi}_- \gamma^\mu \psi_+ = \bar{\psi}_- \gamma_{\perp}^\mu \psi_+$$

$$iV \cdot D \psi_+ + i\gamma^0 \not{D}_\perp \psi_- = 0$$

$$(2\mu + i\tilde{V} \cdot D) \psi_- + i\gamma^0 \not{D}_\perp \psi_+ = 0$$

Eqs. of motion:

At the leading order in μ :

$$iV \cdot D\psi_+ = 0$$

$$\psi_- = 0$$

At the same order:

$$L_D = \int \frac{d\vec{v}}{4\pi} \psi_+^\dagger iV \cdot D\psi_+$$

Propagator:

$$\frac{1}{V \cdot \ell}$$

$$\frac{1}{\not{p} + \mu\gamma_0} = \frac{(p_0 + \mu)\gamma^0 - \vec{p} \cdot \vec{\gamma}}{(p_0 + \mu)^2 - |\vec{p}|^2} \xrightarrow{\mu \rightarrow \infty} \frac{\cancel{V}}{2} \frac{1}{V \cdot \ell} \quad (\langle\langle T(\bar{\psi}\psi) \rangle\rangle)$$

$$\frac{\cancel{V}}{2} = \frac{1}{2} \gamma^0 (1 - \vec{\alpha} \cdot \vec{v}) = P_+ \gamma^0 \quad (\langle\langle T(\psi^\dagger \psi) \rangle\rangle)$$

Integrating out the heavy d.o.f.

For the heavy d.o.f. we can formally repeat the same steps leading to:

$$\int d^4x \bar{\psi}(i\not{D} + \mu\gamma_0)\psi \rightarrow \int \frac{d\vec{v}}{4\pi} \left[\psi_+^\dagger iV \cdot D \psi_+ + \psi_-^\dagger (2\mu + i\tilde{V} \cdot D) \psi_- + (\bar{\psi}_+ i\not{D}_\perp \psi_- + \text{h.c.}) \right]$$

Eliminating the E_- fields one would get the non-local lagrangian:

$$L_D = \int \frac{d\vec{v}}{4\pi} \left[\psi_+^\dagger iV \cdot D \psi_+ - P^{\mu\nu} \psi_+^\dagger \frac{1}{2\mu + i\tilde{V} \cdot D} D_\mu D_\nu \psi_+ \right]$$
$$P^{\mu\nu} = g^{\mu\nu} - \frac{1}{2} \left[V^\mu \tilde{V}^\nu + V^\nu \tilde{V}^\mu \right]$$

Decomposing $\Psi_+ = \Psi_+^{\ell} + \Psi_+^{\text{h}}$ one gets

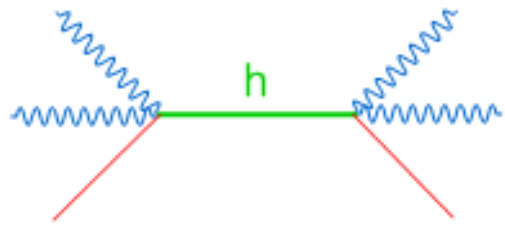
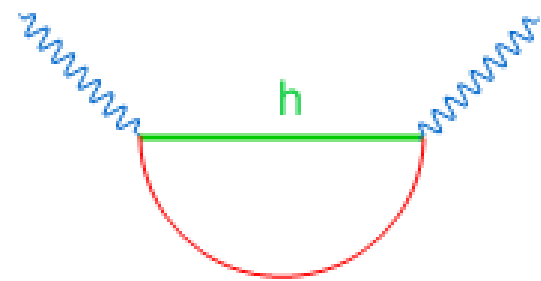
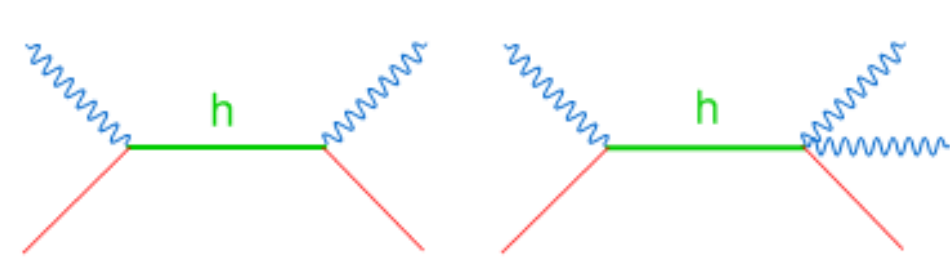
$$L_D = L_D^{\ell} + L_D^{\ell\text{h}} + L_D^{\text{h}}$$

$$L_D^{\ell} = \int \frac{d\vec{v}}{4\pi} \Psi_+^{\ell\dagger} iV \cdot D \Psi_+^{\ell}$$

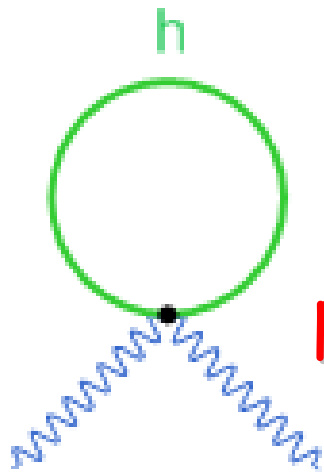
$$L_D^{\ell\text{h}} = \int \frac{d\vec{v}}{4\pi} \left(\Psi_+^{\ell\dagger} iV \cdot D \Psi_+^{\text{h}} - P^{\mu\nu} \Psi_+^{\ell\dagger} \frac{1}{2\mu + i\tilde{V} \cdot D} D_{\mu} D_{\nu} \Psi_+^{\text{h}} + \ell \leftrightarrow \text{h} \right)$$

$$L_D^{\text{h}} = \int \frac{d\vec{v}}{4\pi} \left(\Psi_+^{\text{h}\dagger} iV \cdot D \Psi_+^{\text{h}} - P^{\mu\nu} \Psi_+^{\text{h}\dagger} \frac{1}{2\mu + i\tilde{V} \cdot D} D_{\mu} D_{\nu} \Psi_+^{\text{h}} \right)$$

When integrating out the heavy fields



Contribute only if some gluons are hard, but suppressed by asymptotic freedom



From L_D^h gives the bare Meissner mass

HDET in the condensed phase

Assume $\langle \psi^A C \psi^B \rangle = \Delta_{AB}$ (A, B collective indices)

due to the attractive interaction:

$$L_I = -\frac{G}{4} \varepsilon_{ab} \varepsilon_{\dot{a}\dot{b}} V_{ABCD} \psi_a^A \psi_b^B \psi_{\dot{a}}^{C\dagger} \psi_{\dot{b}}^{D\dagger}$$

$$V_{ABCD} = V_{CDAB}^*, \quad V_{ABCD} = V_{BACD} = V_{ABDC}$$

Decompose



$$\mathbf{L}_I = \mathbf{L}_{\text{cond}} + \mathbf{L}_{\text{int}}$$

$$\mathbf{L}_{\text{cond}} = \frac{\mathbf{G}}{4} V_{\text{ABCD}} \Gamma^{\text{CD}*} \psi^{\text{AT}} \mathbf{C} \psi^{\text{B}} - \frac{\mathbf{G}}{4} V_{\text{ABCD}} \Gamma^{\text{AB}} \psi^{\text{C}\dagger} \mathbf{C} \psi^{\text{D}*}$$

$$\mathbf{L}_{\text{int}} = -\frac{\mathbf{G}}{4} V_{\text{ABCD}} \left(\psi^{\text{AT}} \mathbf{C} \psi^{\text{B}} - \Gamma^{\text{AB}} \right) \left(\psi^{\text{C}\dagger} \mathbf{C} \psi^{\text{D}*} + \Gamma^{\text{CD}*} \right)$$

We define $\Delta_{\text{AB}} = \frac{\mathbf{G}}{2} V_{\text{CDAB}} \Gamma^{\text{CD}}$, $\Delta_{\text{AB}}^* = \frac{\mathbf{G}}{2} V_{\text{CDAB}}^* \Gamma^{\text{CD}*}$

$$\mathbf{L}_{\text{cond}} = \frac{1}{2} \Delta_{\text{AB}}^* \psi^{\text{AT}} \mathbf{C} \psi^{\text{B}} - \frac{1}{2} \Delta_{\text{AB}} \psi^{\text{A}\dagger} \mathbf{C} \psi^{\text{B}*}$$

and neglect \mathbf{L}_{int} . Therefore

$$L_D = \int \frac{d\vec{v}}{4\pi} \frac{1}{2} \sum_{AB} \left[\psi_+^{A\dagger} i\mathbf{V} \cdot \mathbf{D} \psi_+^B + \psi_-^{A\dagger} i\mathbf{V} \cdot \mathbf{D} \psi_-^B + \Delta_{AB}^* \psi_-^{A\dagger} C \psi_+^B - \Delta_{AB} \psi_+^{A\dagger} C \psi_-^B \right]$$

$$\psi_{\pm}(\mathbf{x}) \equiv \psi_{\pm}(\pm\vec{v}, \mathbf{x})$$

Nambu-Gor'kov basis

$$\chi^A = \frac{1}{\sqrt{2}} \begin{pmatrix} \psi_+^A \\ C\psi_-^{A*} \end{pmatrix}$$

$$L_D = \int \frac{d\vec{v}}{4\pi} \chi^{A\dagger} \begin{bmatrix} i\mathbf{V} \cdot \mathbf{D}_{AB} & -\Delta_{AB} \\ -\Delta_{AB}^* & i\tilde{\mathbf{V}} \cdot \mathbf{D}_{AB} \end{bmatrix} \chi^B$$

$$S(\ell) = \frac{1}{(\mathbf{V} \cdot \ell)(\tilde{\mathbf{V}} \cdot \ell) - \Delta\Delta^\dagger} \begin{bmatrix} \tilde{\mathbf{V}} \cdot \ell & \Delta \\ \Delta^\dagger & \mathbf{V} \cdot \ell \end{bmatrix} \left([\Delta, \Delta^\dagger] = 0 \right)$$

From the definition: $\Delta_{AB}^* = -\frac{G}{2} V_{ABCD} \langle \Psi^{C\dagger} C \Psi^{D*} \rangle$

one derives the gap equation (e.g. via functional formalism)

$$\Delta_{AB}^* = iGV_{ABCD} \int \frac{d\vec{v}}{4\pi} \frac{\mu^2}{\pi} \int \frac{d^2\ell}{(2\pi)^2} \Delta_{CE}^* \frac{1}{D_{ED}}$$

$$\frac{1}{D_{AB}} = \left(\frac{1}{(\mathbf{V} \cdot \ell)(\tilde{\mathbf{V}} \cdot \ell) - \Delta\Delta^\dagger} \right)_{AB}$$

Four-fermi interaction one-gluon exchange inspired

$$L_I = \frac{3}{16} G \bar{\psi} \gamma_\mu \lambda^a \psi \bar{\psi} \gamma^\mu \lambda^a \psi$$

Fierz using: $\sum_{a=1}^8 (\lambda^a)_{\alpha\beta} (\lambda^a)_{\delta\gamma} = \frac{2}{3} (3\delta_{\alpha\gamma} \delta_{\beta\delta} - \delta_{\alpha\beta} \delta_{\gamma\delta})$

$$(\sigma_\mu)_{\dot{a}b} (\tilde{\sigma}^\mu)_{dc} = 2\varepsilon_{\dot{a}c} \varepsilon_{bd}$$

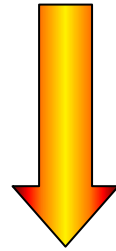
$$\sigma^\mu = (1, \vec{\sigma}), \quad \tilde{\sigma}^\mu = (1, -\vec{\sigma})$$

$$L_I = -\frac{G}{4} V_{(\alpha i)(\beta j)(\gamma k)(\delta \ell)} \psi_i^\alpha \psi_j^\beta \psi_k^{\gamma \dagger} \psi_\ell^{\delta \dagger}$$

$$V_{(\alpha i)(\beta j)(\gamma k)(\delta \ell)} = -(3\delta_{\alpha\delta} \delta_{\beta\gamma} - \delta_{\alpha\gamma} \delta_{\beta\delta}) \delta_{ik} \delta_{j\ell}$$

In the 2SC case $\Delta_{(\alpha i)(\beta j)} = \epsilon_{\alpha\beta 3} \epsilon_{ij} \Delta$

$$\Delta = 4iG \int \frac{d\vec{v}}{4\pi} \frac{\mu^2}{\pi} \int \frac{d^2\ell}{(2\pi)^2} \frac{\Delta}{\ell_0^2 - \ell_{\parallel}^2 - \Delta^2}$$



$$\Delta = \frac{G}{2} \rho \int_0^{\delta} d\xi \frac{\Delta}{\sqrt{\xi^2 + \Delta^2}}, \quad \rho = 4 \frac{\mu^2}{\pi^2}$$

4 pairing fermions

G determined at $T = 0$. M , constituent mass ~ 400 MeV

$$1 = 8G \int_0^\Lambda \frac{d^3\vec{p}}{(2\pi)^3} \frac{1}{\sqrt{|\vec{p}|^2 + M^2}}$$

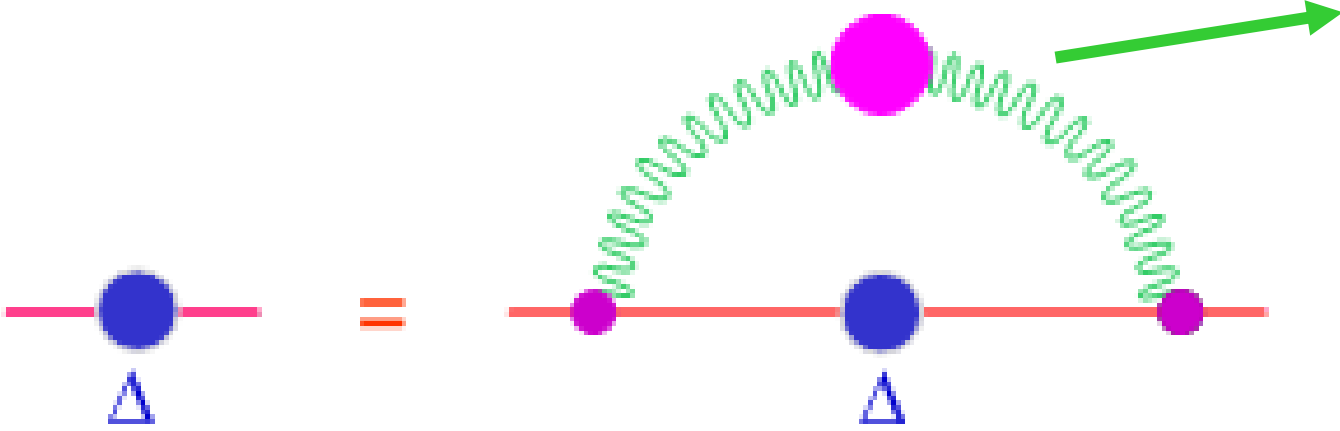
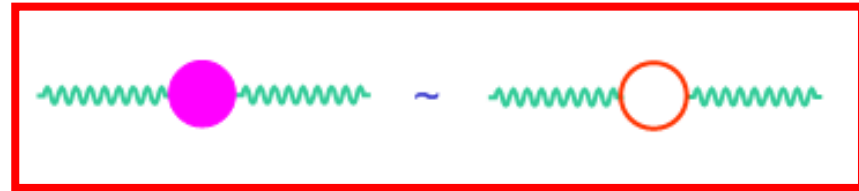
with $\Lambda = \mu + \delta$

For $\mu = 400 - 500$ MeV, $\Lambda = 800$ MeV, $M = 400$ MeV

$$\Delta_{2SC} = 33 - 88 \text{ MeV}$$

Similar values for CFL.

Gap equation in QCD



$$\Delta(p_0) = \frac{g^2}{12\pi^2} \int dq_0 \int d(\cos \theta) \left(\frac{\frac{3}{2} - \frac{1}{2} \cos \theta}{1 - \cos \theta + G/(2\mu^2)} + \frac{\frac{1}{2} + \frac{1}{2} \cos \theta}{1 - \cos \theta + F/(2\mu^2)} \right) \frac{\Delta(q_0)}{\sqrt{q_0^2 + \Delta(q_0)^2}}$$

Hard-loop approximation $q_0 \ll |\vec{q}| \rightarrow 0$

$$F = m_D^2$$

electric

$$G = \frac{\pi}{4} m_D^2 \frac{q_0}{|\vec{q}|}$$

magnetic

$$m_D^2 = N_f \frac{g^2 \mu^2}{2\pi^2}$$

For small momenta magnetic gluons are unscreened and dominate giving a further logarithmic divergence

$$b = 256\pi^4 \left(\frac{2}{N_f} \right)^{5/2} g^{-5}$$

$$\Delta(p_0) = \frac{g^2}{18\pi^2} \int dq_0 \log \left(\frac{b\mu}{|p_0 - q_0|} \right) \frac{\Delta(q_0)}{\sqrt{q_0^2 + \Delta(q_0)^2}}$$

Results:

$$\Delta(p_0) \approx \Delta_0 \sin \left(\frac{g}{3\sqrt{2}\pi} \log \left(\frac{b}{p_0} \right) \right)$$

$$\Delta_0 = 2b\mu e^{-\frac{3\pi^2}{\sqrt{2}g}}$$

$$\left(1 \approx \frac{6g^2}{c} (\log(\delta/\Delta))^2 \rightarrow \Delta \approx \delta e^{-c/g} \right)$$

from the double log

To be trusted only for $\mu > 10^5 \text{ GeV}$ but, if extrapolated at 400-500 MeV, gives values for the gap similar to the ones found using a 4-fermi interaction.

However condensation arises at asymptotic values of μ .

Symmetries of superconducting phases

Consider 3 flavors, **u,d,s**. Antisymmetry requires

$$\langle \Psi_{iL}^{\alpha} \Psi_{jL}^{\beta} \rangle = \varepsilon^{\alpha\beta\gamma} \varepsilon_{ijk} A_{\gamma}^k$$

Diagonalize A through:

$$A \rightarrow g_{\text{SU}(3)_c} A g_{\text{SU}(3)_L} = A_D$$

From the gap equations:

$$A_{(1,1,1)} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad A_{(1,1,0)} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad A_{(1,0,0)} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}$$

with gaps $\Delta_{(1,1,1)} < \Delta_{(1,1,0)} < \Delta_{(1,0,0)}$

but $F_{(1,1,1)} < F_{(1,1,0)} < F_{(1,0,0)}$

since more fermions pair in the more symmetric configurations.

Confirmed by RG analysis (Evans et al.)

From symmetry alone:

$$[(3_c, 3_{L(R)}) \otimes (3_c, 3_{L(R)})]_S = (3_c^*, 3_{L(R)}^*) \oplus (6_c, 6_{L(R)})$$

implying in general

$$\begin{aligned} \langle \Psi_{iL}^\alpha \Psi_{jL}^\beta \rangle &= \Delta \varepsilon^{\alpha\beta I} \varepsilon_{ijI} + \Delta_6 (\delta_i^\alpha \delta_j^\beta + \delta_i^\beta \delta_j^\alpha) = \\ &= \Delta (\delta_i^\alpha \delta_j^\beta - \delta_i^\beta \delta_j^\alpha) + \Delta_6 (\delta_i^\alpha \delta_j^\beta + \delta_i^\beta \delta_j^\alpha) = \\ &= (\Delta + \Delta_6) \delta_i^\alpha \delta_j^\beta + (\Delta_6 - \Delta) \delta_i^\beta \delta_j^\alpha \end{aligned}$$

In NJL case with cutoff **800 MeV**,
constituent mass **400 MeV** and $\mu = 400 \text{ MeV}$

$$\Delta = 85.3 \text{ MeV}, \quad \Delta_6 = -1.3 \text{ MeV}$$

Original symmetry:

$$G_{\text{QCD}} = \text{SU}(3)_c \otimes \text{SU}(3)_L \otimes \text{SU}(3)_R \otimes \text{U}(1)_B \otimes \underbrace{\text{U}(1)_A}_{\text{anomalous}}$$

broken to

$$G_{\text{CFL}} = \text{SU}(3)_{c+L+R} \otimes \text{Z}_2 \otimes \underbrace{\text{Z}_2}_{\text{anomalous}}$$

Goldstones $3 \times 8 + 1 + 1 - 8 = 8 + 8 + 1 + 1 = 17 + \underbrace{1}_{\text{massive}}$

8 give mass to the gluons and 8+1 are true massless Goldstone bosons

Notice:

- Breaking of $U(1)_B$ makes the CFL phase superfluid
- The CFL condensate is not gauge invariant, but consider

$$X_\gamma^k = \varepsilon_{\alpha\beta\gamma} \varepsilon^{ijk} (\Psi_{iL}^\alpha \Psi_{jL}^\beta)^\dagger, \quad Y_\gamma^k = \varepsilon_{\alpha\beta\gamma} \varepsilon^{ijk} (\Psi_{iR}^\alpha \Psi_{jR}^\beta)^\dagger$$

$$\Sigma_j^i = (Y^\dagger X)_j^i = \sum_\alpha (Y_\alpha^j)^* X_\alpha^i$$

Σ is gauge invariant and breaks the global part of G_{QCD} . Also

$$\det(X), \quad \det(Y) \quad \text{break} \quad U(1)_B \otimes U(1)_A \Rightarrow Z_2 \otimes Z_2$$

$U(1)_A$ is broken by the anomaly but induced by a 6-fermion operator irrelevant at the Fermi surface. Its contribution is parametrically small and we expect a very light NGB (massless at infinite chemical potential)

Spectrum of the CFL phase

Choose the basis: $\psi_i^\alpha = \frac{1}{\sqrt{2}} \sum_{A=1}^9 (\lambda_A)_i^\alpha \psi^A$

λ_A $A = 1, \dots, 8$ Gell – Mann matrices

$$\lambda_9 = \lambda_0 = \sqrt{\frac{2}{3}} \times 1$$

$$\text{Tr}[\lambda_A \lambda_B] = 2\delta_{AB}$$

Inverting:
$$\psi^A = \frac{1}{\sqrt{2}} \sum_{\alpha i} (\lambda_A)_\alpha^i \psi_i^\alpha = \frac{1}{\sqrt{2}} \text{Tr}[\lambda_A \Psi]$$

$$\langle \psi^A \psi^B \rangle = \frac{1}{2} \sum (\lambda_A)_\alpha^i (\lambda_B)_\beta^j \Delta \varepsilon^{\alpha\beta I} \varepsilon_{ijI} = \frac{\Delta}{2} \text{Tr} \left[\sum_I (\lambda_A \varepsilon_I \lambda_B^T \varepsilon_I) \right]$$

↓

$$(\varepsilon_I)_{\alpha\beta} = \varepsilon_{\alpha\beta I}$$

$$\sum_I \varepsilon_I g^T \varepsilon_I = g - \text{Tr}[g] \quad \text{for any } 3 \times 3 \text{ matrix } g$$

We get $\langle \psi^A \psi^B \rangle = \Delta_A \delta_{AB}$ $\Delta_A = \begin{cases} A=1, \dots, 8 & \Delta_A = \Delta \\ A=9 & \Delta_9 = -2\Delta \end{cases}$

quasi-fermions

$$\varepsilon_A(\vec{p}) = \sqrt{(\vec{v} \cdot \vec{\ell})^2 + \Delta_A^2}$$

gluons

Expected

$$m_g^2 \approx g^2 F^2$$



NG coupling constant

but wave function renormalization effects important (see later)

NG bosons

Acquire mass through quark masses except for the one related to the breaking of $U(1)_B$

NG boson masses quadratic in m_q since the approximate invariance

$$(Z_2)_L \otimes (Z_2)_R : \Psi_{L(R)} \rightarrow -\Psi_{L(R)}$$

and quark mass term: $\bar{\Psi}_L M \Psi_R + \text{h.c.}$

$$M \rightarrow -M$$

Notice: anomaly breaks $(Z_2)_L \otimes (Z_2)_R$ through instantons, producing a chiral condensate (6 fermions \rightarrow 2), but of order $(\Lambda_{\text{QCD}}/\mu)^8$

In-medium electric charge

$$D_\mu \psi = \partial_\mu \psi - ig_\mu^a T_a \psi - i\psi Q A_\mu$$

The condensate breaks $U(1)_{em}$ but leaves invariant a combination of Q and T_8 .

CFL vacuum: $X_\alpha^i = Y_\alpha^i = \delta_\alpha^i$

Define: $Q_{SU(3)_c} \equiv -\frac{2}{\sqrt{3}} T_8 = \text{diag}\left(-\frac{1}{3}, -\frac{1}{3}, +\frac{2}{3}\right) = Q$

$\tilde{Q} = Q_{SU(3)_c} \otimes 1 - 1 \otimes Q = Q \otimes 1 - 1 \otimes Q$ leaves invariant the ground state

$$Q \langle X \rangle - \langle X \rangle Q \rightarrow Q_{\alpha\beta} \delta_{\beta i} - \delta_{\alpha j} Q_{ji} = 0$$

Eigs(\tilde{Q}) = 0, ± 1 Integers as in the Han-Nambu model

$$A_\mu = \tilde{A}_\mu \cos \theta - \tilde{G}_\mu \sin \theta$$

$$g_\mu^8 = \underbrace{\tilde{A}_\mu \sin \theta + \tilde{G}_\mu \cos \theta}_{\text{rotated fields}}$$

rotated fields

new interaction:

$$g_s g_\mu^8 T_8 \otimes 1 + e A_\mu 1 \otimes Q \rightarrow \tilde{e} \tilde{Q} \tilde{A}_\mu + \tilde{g}_s \tilde{G}_\mu \tilde{T}$$

$$\tan \theta = \frac{2}{\sqrt{3}} \frac{e}{g_s}, \quad \tilde{e} = e \cos \theta, \quad \tilde{g}_s = \frac{g_s}{\cos \theta}$$

$$\tilde{T} = -\frac{\sqrt{3}}{2} \left[(\cos^2 \theta) Q \otimes 1 + (\sin^2 \theta) 1 \otimes Q \right]$$

The rotated "photon" remains massless,
whereas the rotated gluons acquires a mass
through the Meissner effect

A piece of CFL material for massless quarks
would respond to an em field only through NGB:

"bosonic metal"

For quarks with equal masses, no light modes:

"transparent insulator"

For different masses one needs non zero
density of electrons or a kaon condensate
leading to massless excitations

Quark-hadron continuity

CFL main properties:

- Integer charges (confinement)
- Breaking of chiral symmetry
- Baryon number superfluidity

Except for last point as the hadronic phase.

However for 3-flavors a di-baryon condensate

$$H \sim \langle udsuds \rangle \approx \det[X]$$

is possible. For instance:

$(p, n, \Sigma^+, \Sigma^-, \Sigma^0, \Lambda)$

In this phase (**hypernuclear matter phase, HMP**) again the baryon number is broken and superfluidity arises. This phase has the same symmetries as CFL:

No phase transition between
CFL and HMP

(Schafer & Wilczek)

Why? Remember complementarity idea

(Banks & Rabinovici; Fradkin & Shenker;)

- In some case two phases are indistinguishable. From one phase to the other varying continuously the parameters.

- Consider a confined and a broken phase:

$\psi_i \in R$ of G , elementary states in the broken phase

$Q_\alpha \in \tilde{R}$ of \tilde{G} , composite states in the confined phase

$$G \sim \tilde{G}, \quad R \sim \tilde{R}$$

Suppose there are Higgs fields such that

$$Q_\alpha = \psi_i \varphi_\alpha^i \quad \text{with} \quad \langle \varphi_\alpha^i \rangle \approx \delta_\alpha^i \quad \text{in the broken phase}$$

Fermions in the two phases are the same but for a re-definition of the quantum numbers (neutrality of the Higgs in the broken phase)

Vector mesons in the broken phase from gauge fields:

$$(Z_\mu)_j^i = -\varphi_j^{\beta*} \left(\partial_\mu - (g_\mu)_\beta^\alpha \right) \varphi_\alpha^i$$

For CFL and HMP we have: $G = SU(3)_c$, $\tilde{G} = SU(3)$ with fermions in the fundamental in three copies

CFL phase	HYP phase
$\Psi_{\alpha}^i \langle D_k^{\alpha} \rangle$	$B_k^i = \Psi_{\alpha}^i D_k^{\alpha}$

$$D_i^{\alpha} = \varepsilon_{ijk} \varepsilon^{\alpha\beta\gamma} \psi_{\beta}^j \psi_{\gamma}^k, \quad \langle D_i^{\alpha} \rangle \approx \delta_i^{\alpha}$$

New definition of the electric charge in the CFL phase. Add to the old generators the charges of the diquark condensates.

Quarks with the same charge as the baryons. Vector mesons $G_k^i = D_{\alpha}^{i*} g_{\beta}^{\alpha} D_k^{\beta}$

Electric charges of the vector mesons

G_k^i	u	d	s
u	0	-1	-1
d	1	0	0
s	1	0	0

$$G_k^i = D_{\alpha}^{i*} g_{\beta}^{\alpha} D_k^{\beta}$$

- In the HYP phase a nonet of vector bosons, but if the dibaryon H exists the singlet vector becomes unstable. Notice that in the CFL phase only 8 vector bosons.
- In the CFL phase 8+1 quarks, but the gap of the singlet is twice the one of the octet. The singlet baryon has the structure

$$S = \varepsilon^{ijk} \varepsilon_{\alpha\beta\gamma} \Psi_i^\alpha \Psi_j^\beta \Psi_k^\gamma$$

and $H=SS$

Spectrum of the 2SC phase

Remember $\Delta_{(\alpha i)(\beta j)} = \varepsilon_{\alpha\beta 3} \varepsilon_{ij} \Delta$

No Goldstone bosons

ψ_i^α $\alpha = 1, 2$ gapped (equal gap)

ψ_i^3 ungapped

$SU(3)_c \rightarrow SU(2)_c$



$8 - 3 = 5$ massive gluons



Light modes:

$\psi_i^3 + 3$ gluons ($M = 0$)

2+1 flavors

It could happen $\mu \approx m_s, \mu \gg m_u, m_d$

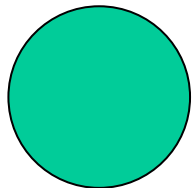
$$\mu \gg m_u, m_d, m_s$$



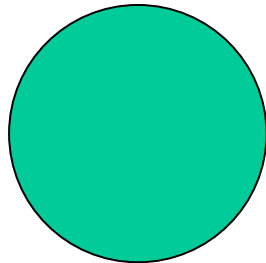
Phase transition expected

$$m_u, m_d < \mu < m_s$$

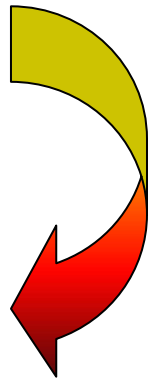
$$E_F = \mu = \sqrt{p_F^2 + M^2} \rightarrow p_F = \sqrt{\mu^2 - M^2}$$



$$M_1 > M_2$$



The radius of the Fermi sphere decreases with the mass



Simple model: $p_{F_1} = \sqrt{\mu^2 - m_s^2}, \quad p_{F_2} = \mu$

$$\Omega_{\text{unpair}} = 2 \int_0^{p_{F_1}} \frac{d^3 \vec{p}}{(2\pi)^3} \left(\sqrt{\vec{p}^2 + m_s^2} - \mu \right) + 2 \int_0^{p_{F_2}} \frac{d^3 \vec{p}}{(2\pi)^3} (|\vec{p}| - \mu)$$

$$\Omega_{\text{pair}} = 2 \int_0^{p_{F_{\text{comm}}}} \frac{d^3 \vec{p}}{(2\pi)^3} \left(\sqrt{\vec{p}^2 + m_s^2} - \mu \right) + 2 \int_0^{p_{F_{\text{comm}}}} \frac{d^3 \vec{p}}{(2\pi)^3} (|\vec{p}| - \mu) - \underbrace{\frac{\mu^2 \Delta^2}{4\pi^2}}_{\text{condensation energy}}$$

$$\frac{\partial \Omega_{\text{pair}}}{\partial p_{F_{\text{comm}}}} = 0 \quad \rightarrow \quad p_{F_{\text{comm}}} = \mu - \frac{m_s^2}{4\mu}$$

condensation energy

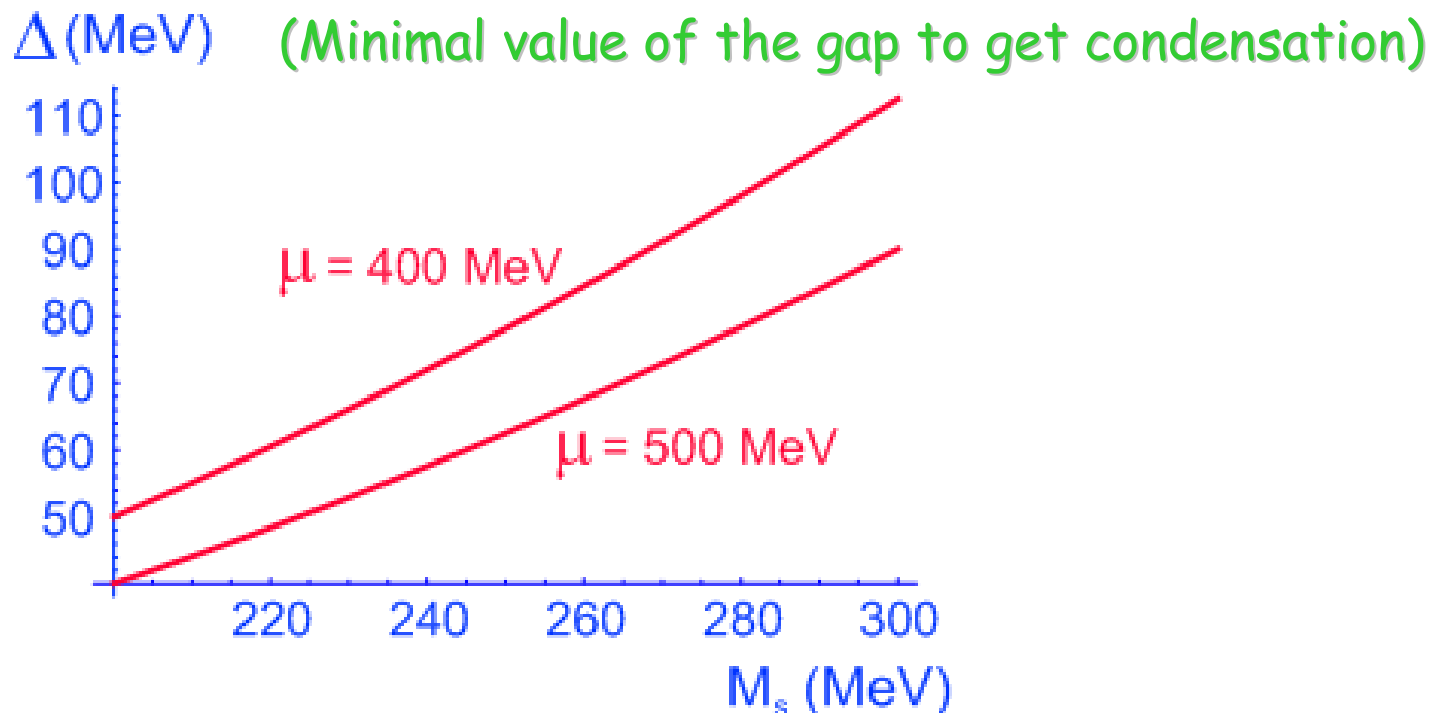
$$\Omega_{\text{pair}} - \Omega_{\text{unpair}} \approx \frac{1}{16\pi^2} \left(m_s^4 - 4\Delta^2 \mu^2 \right)$$

Condensation if:

$$\mu > \frac{m_s^2}{2\Delta}$$

Notice that the **transition must be first order** because for being in the pairing phase

$$\Delta > \frac{m_s^2}{2\mu} \neq 0$$



Effective lagrangians

- Effective lagrangian for the CFL phase
- Effective lagrangian for the 2SC phase

Effective lagrangian for the CFL phase

NG fields as the phases of the condensates in the $(\bar{3}, \bar{3})$ representation

$$X_\gamma^k = \varepsilon_{\alpha\beta\gamma} \varepsilon^{ijk} \left\langle \Psi_{iL}^\alpha \Psi_{jL}^\beta \right\rangle^*, \quad Y_\gamma^k = \varepsilon_{\alpha\beta\gamma} \varepsilon^{ijk} \left\langle \Psi_{iR}^\alpha \Psi_{jR}^\beta \right\rangle^*$$

Quarks and X, Y transform as

$$\Psi_L \rightarrow e^{i(\alpha+\beta)} g_c \Psi_L g_L^T, \quad \Psi_R \rightarrow e^{i(\alpha-\beta)} g_c \Psi_R g_R^T$$
$$g_c \in SU(3)_c, \quad g_{L,R} \in SU(3)_{L,R}, \quad e^{i\alpha} \in U(1)_B, \quad e^{i\beta} \in U(1)_A$$

$$X \rightarrow g_c X g_L^T e^{-2i(\alpha+\beta)}, \quad Y \rightarrow g_c Y g_R^T e^{-2i(\alpha-\beta)}$$

Since $X, Y \in U(3)$

The number of NG fields is

$$\#X + \#Y = (1 + 8) + (1 + 8) = 18$$

8 of these fields give mass to the gluons. There are only **10 physical NG bosons** corresponding to the breaking of the global symmetry (we consider also the NGB associated to $U(1)_A$)

$$SU(3)_L \otimes SU(3)_R \otimes U(1)_V \otimes U(1)_A$$



$$SU(3)_{L+R} \otimes Z_2 \otimes Z_2$$

Better use fields belonging to $SU(3)$. Define

$$X = \hat{X} e^{2i(\varphi+\theta)} \quad Y = \hat{Y} e^{2i(\varphi-\theta)}$$

and

$$d_X = \det(X) = e^{6i(\varphi+\theta)} \quad d_Y = \det(Y) = e^{6i(\varphi-\theta)}$$

transforming as

$$\begin{aligned} \hat{X} &\rightarrow g_c \hat{X} g_L^T & \hat{Y} &\rightarrow g_c \hat{Y} g_R^T \\ \varphi &\rightarrow \varphi - \alpha & \theta &\rightarrow \theta - \beta \end{aligned}$$

The breaking of the global symmetry can be described by the gauge invariant order parameters

$$\Sigma_j^i = (\hat{Y}_\alpha^j)^* \hat{X}_\alpha^i \rightarrow \Sigma = \hat{Y}^\dagger \hat{X}, \quad d_X, \quad d_Y$$

Σ, d_X, d_Y are 10 fields describing the physical NG bosons. Also

$$\Sigma \rightarrow g_R^* \Sigma g_L^T$$

shows that Σ^T transforms as the usual chiral field.

Consider the currents:

$$J_X^\mu = \hat{X} D^\mu \hat{X}^\dagger = \hat{X} (\partial^\mu \hat{X}^\dagger + \hat{X}^\dagger g^\mu) = \hat{X} \partial^\mu \hat{X}^\dagger + g^\mu$$

$$J_Y^\mu = \hat{Y} D^\mu \hat{Y}^\dagger = \hat{Y} (\partial^\mu \hat{Y}^\dagger + \hat{Y}^\dagger g^\mu) = \hat{Y} \partial^\mu \hat{Y}^\dagger + g^\mu$$

$$g^\mu = i g_s g_a^\mu T^a$$

$$J_{X,Y}^\mu \rightarrow g_c J_{X,Y}^\mu g_c^\dagger$$

Most general lagrangian up to two derivatives
invariant under G , the space rotation group $O(3)$
and Parity (R.C. & Gatto 1999)

$$P: (X \leftrightarrow Y, \varphi \leftrightarrow \varphi, \theta \leftrightarrow -\theta)$$

$$L = -\frac{F_T^2}{4} \text{Tr} \left[\left(J_X^0 - J_Y^0 \right)^2 \right] - \alpha_T \frac{F_T^2}{4} \text{Tr} \left[\left(J_X^0 + J_Y^0 \right)^2 \right] + \frac{1}{2} (\partial_0 \varphi)^2 + \frac{1}{2} (\partial_0 \theta)^2 +$$

$$+ \frac{F_S^2}{4} \text{Tr} \left[\left(\vec{J}_X - \vec{J}_Y \right)^2 \right] + \alpha_S \frac{F_S^2}{4} \text{Tr} \left[\left(J_X^0 + J_Y^0 \right)^2 \right] - \frac{v_\varphi^2}{2} |\vec{\nabla} \varphi|^2 - \frac{v_\theta^2}{2} |\vec{\nabla} \theta|^2$$

$$L = -\frac{F_T^2}{4} \text{Tr} \left[\left(\hat{X} \partial_0 \hat{X}^\dagger - \hat{Y} \partial_0 \hat{Y}^\dagger \right)^2 \right] - \alpha_T \frac{F_T^2}{4} \text{Tr} \left[\left(\hat{X} \partial_0 \hat{X}^\dagger + \hat{Y} \partial_0 \hat{Y}^\dagger + 2g_0 \right)^2 \right] +$$

$$+ \frac{F_S^2}{4} \text{Tr} \left[\left(\hat{X} \vec{\nabla} \hat{X}^\dagger - \hat{Y} \vec{\nabla} \hat{Y}^\dagger \right)^2 \right] + \alpha_S \frac{F_S^2}{4} \text{Tr} \left[\left(\hat{X} \vec{\nabla} \hat{X}^\dagger + \hat{Y} \vec{\nabla} \hat{Y}^\dagger + 2\vec{g} \right)^2 \right] +$$

$$+ \frac{1}{2} (\partial_0 \varphi)^2 + \frac{1}{2} (\partial_0 \theta)^2 - \frac{v_\varphi^2}{2} |\vec{\nabla} \varphi|^2 - \frac{v_\theta^2}{2} |\vec{\nabla} \theta|^2$$

Using $SU(3)_c$ gauge invariance we can choose:

$$\hat{X} = \hat{Y}^\dagger = e^{i\Pi^a T^a / F_T}$$

Expanding at the second order in the fields

$$L \approx \frac{1}{2}(\partial_0 \Pi^a)^2 + \frac{1}{2}(\partial_0 \varphi)^2 + \frac{1}{2}(\partial_0 \theta)^2 - \frac{v^2}{2} |\vec{\nabla} \Pi^a|^2 - \frac{v_\varphi^2}{2} |\vec{\nabla} \varphi|^2 - \frac{v_\theta^2}{2} |\vec{\nabla} \theta|^2$$

$$v^2 = \frac{F_S}{F_T}$$

Gluons acquire Debye and Meissner masses (not the rest masses, see later)

$$m_D^2 = \alpha_T g_s^2 F_T^2, \quad m_S^2 = \alpha_S g_s^2 F_S^2 = \alpha_S v^2 g_s^2 F_T^2$$

Low energy theory supposed to be valid at energies \ll gap. Since we will see that gluons have masses of order Δ they can be decoupled

$$\begin{aligned}
 L = & -\frac{F_T^2}{4} \text{Tr} \left[\left(\hat{X} \partial_0 \hat{X}^\dagger - \hat{Y} \partial_0 \hat{Y}^\dagger \right)^2 \right] - \alpha_T \frac{F_T^2}{4} \text{Tr} \left[\left(\hat{X} \partial_0 \hat{X}^\dagger + \hat{Y} \partial_0 \hat{Y}^\dagger + 2g_0 \right)^2 \right] + \\
 & + \frac{F_S^2}{4} \text{Tr} \left[\left(\hat{X} \vec{\nabla} \hat{X}^\dagger - \hat{Y} \vec{\nabla} \hat{Y}^\dagger \right)^2 \right] + \alpha_T \frac{F_S^2}{4} \text{Tr} \left[\left(\hat{X} \vec{\nabla} \hat{X}^\dagger + \hat{Y} \vec{\nabla} \hat{Y}^\dagger + 2\vec{g} \right)^2 \right] + \\
 & + \frac{1}{2} (\partial_0 \phi)^2 + \frac{1}{2} (\partial_0 \theta)^2 - \frac{v_\phi^2}{2} |\vec{\nabla} \phi|^2 - \frac{v_\theta^2}{2} |\vec{\nabla} \theta|^2
 \end{aligned}$$

$$g_\mu = -\frac{1}{2} \left(\hat{X} \partial_\mu \hat{X}^\dagger + \hat{Y} \partial_\mu \hat{Y}^\dagger \right)$$

we get the gauge invariant result:

$\sim \chi$ -lagrangian

$$\mathbf{L}_{\text{NGB}} = \frac{F_{\text{T}}^2}{4} \left(\text{Tr} \left[\partial_t \Sigma \partial_t \Sigma^+ \right] - v^2 \text{Tr} \left[\vec{\nabla} \Sigma \vec{\nabla} \Sigma^+ \right] \right) -$$
$$-\frac{1}{2} \left((\partial_t \varphi)^2 - v_\varphi^2 |\vec{\nabla} \varphi|^2 \right) - \frac{1}{2} \left((\partial_t \theta)^2 - v_\theta^2 |\vec{\nabla} \theta|^2 \right)$$

$$\Sigma = \hat{Y}^+ \hat{X}$$

Effective lagrangian for the 2SC phase

Light modes: 2 ungapped fermions and 3 massless gluons of $SU(2)_c$. Consider gluons:

➤ Gauge invariance: $E_i^a = F_{oi}^a$, $B_i^a = \frac{1}{2} \epsilon_{ijk} F_{jk}^a$

➤ $O(3) + \text{Parity}$
$$L = \frac{1}{g^2} \sum_{a=1}^3 \left(\frac{\epsilon}{2} \vec{E}^a \cdot \vec{E}^a - \frac{1}{2\lambda} \vec{B}^a \cdot \vec{B}^a \right)$$

➤ Gluon velocity
$$v = \frac{1}{\sqrt{\epsilon\lambda}}$$

Effective gauge coupling ($\lambda = 1$, see later):

$$V_{\text{Coul}} = \frac{g^2}{\epsilon r} = \frac{g_{\text{eff}}^2}{r} \quad \Rightarrow \quad g_{\text{eff}}^2 = \frac{g^2}{\epsilon}$$

$$\alpha_s = \frac{g^2}{4\pi c} \quad \Rightarrow \quad \alpha'_s = \frac{g_{\text{eff}}^2}{4\pi v} = \frac{g^2}{4\pi c\sqrt{\epsilon}}$$

Since $\epsilon \gg 1$ (see later), gluons move slowly and

$$\frac{\alpha'_s}{\alpha_s} = \frac{1}{\sqrt{\epsilon}} \ll 1$$

Equivalently do the re-scaling

$$x^{0'} = \frac{x^0}{\sqrt{\varepsilon}}, \quad A_0^{\alpha'} = \sqrt{\varepsilon} A_0^\alpha, \quad g' = \frac{g}{\varepsilon^{1/4}}$$

⇓

$$L = -\frac{1}{g'^2} F_{\mu\nu}^{a'} F^{\mu\nu a'}$$

Since $\varepsilon = 1 + \frac{g^2 \mu^2}{18\pi^2 \Delta^2} \approx \frac{g^2 \mu^2}{18\pi^2 \Delta^2}$ from $\alpha'_s = \frac{g^2}{4\pi c \sqrt{\varepsilon}}$

$$\alpha'_s = \frac{3}{2\sqrt{2}} \frac{g\Delta}{\mu}$$

This defines the coupling at the matching scale Δ

α'_s grows going to lower energies and becomes ~ 1 at Λ'_{QCD} . The coupling is small at the matching scale and we expect Λ'_{QCD} to be small. From one-loop beta-function

$$\Lambda'_{\text{QCD}} \approx \Delta e^{-2\pi/\beta_0\alpha'_s} \approx \Delta \exp\left(-\frac{2\sqrt{2}\pi}{11} \frac{\mu}{g\Delta}\right), \quad \beta_0 = \frac{22}{3}$$

Numerical estimate very difficult, it depends on

$$\Delta = c \mu g^{-5} e^{-3\pi^2/\sqrt{2}g}$$

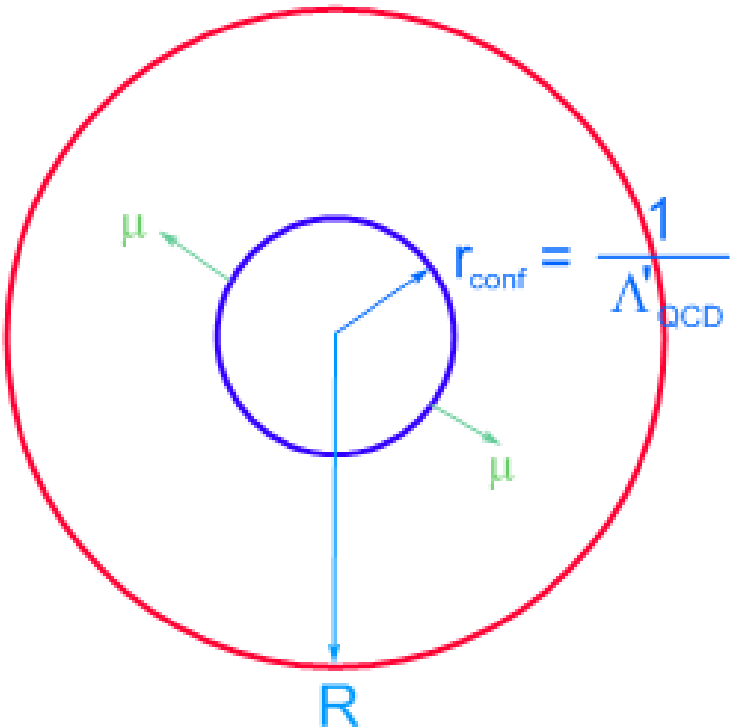
$$c = 512\pi^4$$

$$c = 512\pi^4 \exp\left(-\frac{4 + \pi^2}{8}\right)$$

$$0.3 \text{ KeV} \leq \Lambda'_{\text{QCD}} \leq 10 \text{ MeV}$$

For $\Lambda_{\text{QCD}} = 200 \text{ MeV}$, $\mu = 600 \text{ MeV}$

For increasing μ exponential decreasing
Confinement radius $\frac{1}{\Lambda'_{\text{QCD}}}$ grows
exponentially



Looking at fixed distance
and increasing the density
there is a crossover at
which the color degrees of
freedom are unconfined.