Color-flavor locked strangelets in a quark mass density-dependent model

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The color-flavor locked (CFL) phase of strangelets is investigated in a quark mass density-dependent model. Parameters are determined by stability arguments. It is concluded that three solutions to the system equations can be found, corresponding, respectively, to positively charged, negatively charged, and nearly neutral CFL strangelets. The charge to baryon number of the positively charged strangelets is smaller than the previous result, while the charge of the negatively charged strangelets is nearly proportional in magnitude to the cubic-root of the baryon number. However, the positively charged strangelets are more stable compared to the other two solutions.

Keywords: Quark matter; Color-flavor locking; Strangelets.

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1. Introduction

At high densities, quark matter, or its finite lumps, the so-called strangelets, may form the color superconductor by chiral-symmetry violating condensate. Most probably strange quark matter (SQM) is more stable than the hadronic matters which may have far-reaching consequences astronomically and cosmologically. It has potential applications to the astrophysics of compact stars and to high energy heavy ion collisions. Recently, the properties of CFL phase have been investigated by many authors with various models. Rajagopal and Wilczek demonstrated that the CFL quark matter is automatically neutralized without any requirements on electrons. Madsen studied the positively charged CFL strangelets, and found that the CFL strangelets may be more stable than the ordinary strangelets without color-flavor locking. Alford investigated the new gapless CFL phase in neutral cold quark matter. Lugones and Horvath derived an approximate equation of state (EOS) to lowest order in $\Delta$ and $m$ and found the effects of pairing on EOS and the window of stability for CFL strange matter.

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In the standard MIT bag model, quark masses are constant. As is well known, however, the quark mass changes with environment, or, in other words, it depends on density. Such masses are usually called effective masses. Effective masses and effective bag constants for quark matter have been extensively discussed within the Nambu–Jona-Lasinio model. In recent years, the quark mass density-dependent model (QMDD) has been successful in describing the unpaired quark matter. Very recently, we also studied CFL strangelets within the framework of bag model. Now in this paper, we study how the density dependence of quark masses influences the properties of CFL strangelets. It is found that for a given baryon number, there are also three kinds of CFL strangelets which are, respectively, charge-positive (slet-1), negative (slet-2), and nearly neutral (slet-3). These multiple solutions are determined by the multiple values of quarks Fermi momenta satisfying the self-consistent condition. The finite-size effects of strangelets with fixed Fermi momenta influence the fractions of quarks as well as the charge of strangelets. For slet-1, the ratio of the charge to squared cubic-root baryon number is only half of that in the pure bag model, while the charge of slet-2 is proportional, in magnitude, to the cubic-root of baryon number. For the same parameters, the slet-1 is more stable.

This paper is organized as follow. In the subsequent Sec. 2, we briefly give the thermodynamic formulas used for the study of CFL strangelets in the mass density-dependent model. Then we present the numerical results and discussions in Sec. 3. The Sec. 4 is a short summary.

2. Thermodynamic formulas and mass density dependence of quark masses

In this paper, the CFL phase consists of up, down, and strange quarks in three-color QCD. The symmetrical color-flavor locked phase needs an attractive interaction between two quarks near the Fermi surface with equal and opposite momentum. The Fermi surface is vividly fixed in the momentum space. The condensate of Cooper pairs, associated with the group of ud, us and ds, spontaneously breaks the color gauge symmetry. At high densities, quarks of three colors and three flavors are allowed to pair and have the same Fermi momentum.

In the pairing ansatz, we consider the case $\Delta_3 \approx \Delta_2 = \Delta_1 = \Delta$ for common CFL phase. Here $\psi^\alpha_a$ is a quark field with color $\alpha = (r,g,b)$ and flavor $a = (u,d,s)$. The thermodynamic potential of the color-flavor locked strangelets can be written as

$$\Omega = \sum_i \Omega_{t,i} + \Omega_{\text{pair},V} + B,$$

where the bag constant $B$ is appended as normally done. Eq. (2) is derived from microscopic models under the condition that $\Delta/\mu$ is small. For pairing contri-
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bution, we include the volume term \( \Omega_{\text{pair,V}} \approx -3 \Delta^2 \bar{\mu}^2 / \pi^2 \), which will appear in the number density, energy and pressure etc. The finite-size contribution to \( \Omega_{\text{pair}} \) is assumed small and reasonably neglected. The quantity \( \bar{\mu} \equiv (\mu_u + \mu_d + \mu_s)/3. \) \( \Omega_{\ell,i} \) denotes the normal quark contribution from flavor type \( i \) (\( i = u, d, s \)), i.e.,

\[
\Omega_{\ell,i} = - T \int_{p_i}^{p_F} \ln \left\{ 1 \pm \exp \left[ - \frac{\varepsilon_i(p_i) - \mu_i}{T} \right] \right\} n_i(p, m_i, R)dp.
\]

The integral upper limit \( p_F \) means Fermi momentum, which is equal for \( u, d, s \) quarks in CFL quark matter. When \( T = 0 \), this expression can be simplified, giving

\[
\Omega_{\ell,i} = \int_{0}^{p_F} (\varepsilon_i - \mu_i)n_i(p, m_i, R)dp,
\]

where the density of states for strangelets from the multi-reflection theory is

\[
n_i(p, m_i, R) = g_i \left[ \frac{1}{2\pi^2} p^2 + \frac{S}{V} f_i^S \left( \frac{p}{m_i} \right) p + \frac{C}{V} f_i^C \left( \frac{p}{m_i} \right) \right].
\]

Here \( V = \frac{4}{3} \pi R^3 \), \( S = 4\pi R^2 \), and \( C = 8\pi R \). The function \( f_i^S \) and \( f_i^C \) are respectively referred as surface term and curvature term

\[
f_i^S \left( \frac{p}{m_i} \right) = - \frac{1}{4\pi^2} \cot^{-1} \left( \frac{p}{m_i} \right),
\]

\[
f_i^C \left( \frac{p}{m_i} \right) = \frac{1}{12\pi^2} \left[ 1 - \frac{3p}{2m_i} \cot^{-1} \left( \frac{p}{m_i} \right) \right].
\]

For convenience, we define a parameter \( \phi_i \equiv \arctan(p_F/m_i) \). If the thermodynamic potential density in Eq. \([1] \) is divided into three parts i.e., \( \Omega_{\ell,i} = \Omega_i^V + \frac{\Delta}{R} \Omega_i^S + \frac{\Delta}{2R} \Omega_i^C \), then the volume term \( \Omega_i^V \), the surface term \( \Omega_i^S \), and the curve term \( \Omega_i^C \) can have analytical forms as,

\[
\Omega_i^V = \frac{g_i m_i^4}{16\pi^2} \left\{ \frac{8}{3} \frac{\mu_i}{m_i} \tan^3 \phi_i + \ln \left[ \tan \phi_i + \sec \phi_i \right] - \tan \phi_i \sec \phi_i^2 - \tan^3 \phi_i \sec \phi_i \right\}, \tag{8}
\]

\[
\Omega_i^S = \frac{g_i m_i^3}{24\pi^2} \left\{ \ln \left[ \tan \phi_i + \sec \phi_i \right] - \frac{3\mu_i}{m_i} \left[ \frac{\pi}{2} - \phi_i \right] \tan^2 \phi_i + \tan \phi_i - \phi_i \right\} + \sec \phi_i \left[ \tan \phi_i + \sec^2 \phi_i(\pi - 2\phi_i) \right] - \pi \right\}, \tag{9}
\]

\[
\Omega_i^C = \frac{g_i m_i^2}{48\pi^2} \left\{ \ln \left[ \tan \phi_i + \sec \phi_i \right] - \frac{3\mu_i}{m_i} \left[ \frac{1}{3} \tan \phi_i - \tan^2 \phi_i(\frac{\pi}{2} - \phi_i) + \phi_i \right] + \tan \phi_i \sec \phi_i - 2 \sec^3 \phi_i(\frac{\pi}{2} - \phi_i) + \pi \right\}. \tag{10}
\]
The common fermi momentum $p_F$ is a fictional intermediate parameter, which can be determined by minimizing the thermodynamic potential $\Omega$, i.e.,

$$\frac{\partial \Omega}{\partial p_F} = 0,$$

or, explicitly,

$$\sum_{i=u,d,s} n'_i(p_F, m_i, R) \left[ \sqrt{p_F^2 + m_i^2} - \mu_i \right] = 0.$$

By differentiating $\Omega_{t,i}$ with respect to the chemical potential $\mu_i$ we obtain

$$n_{t,i} = g_i p_F^3 / (6\pi^2) + \frac{3g_i m_i^2}{8\pi^2 R} \left[ \phi_i - \tan \phi_i - (\frac{\pi}{2} - \phi_i) \tan^2 \phi_i \right] + \frac{3g_i m_i}{8\pi^2 R^2} \left[ \phi_i + \frac{1}{3} \tan \phi_i - (\frac{\pi}{2} - \phi_i) \tan^2 \phi_i \right].$$

Hence the number density for flavor $i$ is

$$n_i = n_{t,i} + \frac{2}{\pi^2} \Delta^2 \bar{\mu}.$$  

The second term is from the paring effect. We should accordingly note that the density depends directly on the chemical potential $\mu$ and the paring parameter $\Delta$, not merely the Fermi momentum “$p_F$”. The corresponding energy density can then be expressed as

$$E = \Omega + \sum_i \mu_i n_i$$

$$= \sum_i \left[ \Omega_{t,i} + \mu_i (n_{t,i} + n_{\text{pair},i}) \right] + \Omega_{\text{pair},i} + B.$$  

In the conventional bag model, quark masses are constant. As is well known, however, quark masses vary with environment. In fact, not only quark masses will change but also the coupling constant will run in the medium. Effective masses and effective bag constants for quark matter have been extensively discussed, e.g., within the Nambu-Jona-Lasinio model and within a quasiparticle model. Therefore, in recent years, the quark mass density-dependent model has been shown to be successful in the study of quark matter. The question now is how to parameterize the density dependence of quark masses. In principle, it should be treated dynamically and be consistent with the overall chiral symmetries of QCD. In studying the unpaired phase of quark matter, the equivalent quark mass is

$$m_q = m_{q0} + \frac{D}{n_b^{1/3}}.$$  

The baryon number $A$ is connected to the quark number $N_u, N_d$ and $N_s$ by

$$A = \frac{1}{3} (N_u + N_d + N_s).$$
and the chemical equilibrium requires

$$\mu_d = \mu_s = \mu_u + \mu_e .$$  \hspace{1cm} (18)$$

When the baryon number $A \ll 10^7$, electrons can’t exist in a strangelet because the electron Compton wavelength exceeds the sphere radius $R$. So the electron “effective” chemical potential is zero in strangelets, which is consistent with the viewpoints of Ref. 54 and 55 to some extent. Figure 11.1 of Ref. 54 and Figure 3 of Ref. 55 point that when the mass of quark matter is less than $10^9$ GeV, the electron could be outside the quark bag. So when $\mu_e = 0$, Eq. (18) becomes $\mu_u = \mu_d = \mu_s$. With this relation, the chemical potential can be obtained from Eq. (12) for a fixed $A$ and a given radius $R$.

The radius of a strangelet with a given baryon number is determined by minimizing the thermodynamic potential with respect to the radius,

$$\frac{\partial \Omega}{\partial R} = 0 .$$ \hspace{1cm} (19)$$

Or, equivalently, by setting the pressure

$$P = -\frac{1}{3} \frac{\partial \Omega}{\partial R} + n_b \frac{\partial m_i}{\partial n_b} \frac{\partial \Omega}{\partial m_i} + B$$ \hspace{1cm} (20)$$
to zero, i.e., $P = 0$. The first two terms on the right hand side of Eq. (20) is normal, while the third term is special when quark masses are density-dependent. As has been shown in literature, this extra term is necessary to satisfy the basic thermodynamic requirement, i.e., the energy minimum must be exactly located at the zero pressure [39,40].

According to Eqs. (8), (9), and (10), the derivative of thermodynamic potential density with respect to mass can be divided into three parts as,

$$\frac{\partial \Omega_V}{\partial m_i} = \frac{g_i m_i^2}{4 \pi^2} \left\{ \tan \phi_i \sec \phi_i - \ln(\tan \phi_i + \sec \phi_i) \right\},$$ \hspace{1cm} (21)$$

$$\frac{\partial \Omega_S}{\partial m_i} = - \frac{g_i m_i^2}{8 \pi^2} \left\{ \tan \phi_i \sec \phi_i - 2 \phi_i \sec \phi_i + \ln(\tan \phi_i + \sec \phi_i) - (1 - \sec \phi_i) \pi + 2 \frac{\mu_i}{m_i} [\phi_i - \tan \phi_i] \right\},$$ \hspace{1cm} (22)$$

$$\frac{\partial \Omega_C}{\partial m_i} = \frac{g_i m_i}{96 \pi^2} \left\{ 4 \ln(\tan \phi_i + \sec \phi_i) + 2 \tan \phi_i \left( \frac{3 \mu_i}{m_i} - 2 \sec \phi_i \right) \right.$$

$$+ \left. \pi \left( 2 \tan^2 \phi_i \sec \phi_i - 4 \sec \phi_i - \frac{3 \mu_i \tan^2 \phi_i}{m_i} + 4 \right) - 2 \phi_i \left( 2 \tan^2 \phi_i \sec \phi_i + \frac{3 \mu_i}{m_i} - 4 \sec \phi_i - \frac{3 \mu_i \tan^2 \phi_i}{m_i} \right) \right\}. \hspace{1cm} (23)$$

Because no electrons are included in the system, finite-size CFL strangelets are charged. The net electric charge is $Z = \frac{2}{3} n_u - \frac{1}{3} n_d - \frac{1}{3} n_s$ in unit of the electric
charge of an electron. In principle, the Coulomb effect should be included, though it is small compared with the strong interaction. We include it in the numerical calculations.

For the charged CFL strangelets, the Coulomb energy is,

\[ E_{\text{Coul}} = \frac{1}{10} \frac{\alpha Z^2}{R} + \frac{1}{2} \frac{\alpha Z^2}{R}, \]

(24)

where the fine structure constant is \( \alpha \approx \frac{1}{137} \) and \( Z \) is the volume term of the total electric charge \( Z \) of the CFL strangelets.

3. Numerical results and discussions

We now discuss the parameters adopted in our calculations. The current quark masses are taken to be \( m_{u0} = 5 \text{ MeV}, m_{d0} = 10 \text{ MeV} \) except for specific indications somewhere. The superconducting gap varies from several tens to several hundreds of MeV in previous papers. In this paper we take the value \( \Delta = 100 \text{ MeV} \), as in Ref. 15.

It is known that the lowest energy per baryon in ordinary nuclei is 930 MeV for iron. In order not to contradict with standard nuclear physics, the energy per baryon at zero temperature should be greater than 930 MeV for two flavor quark matter, and less than 930 MeV for three flavor quark matter so that SQM can have a chance to be absolutely stable. We show the stability window in Fig. 1. The horizontal axis, namely \( D = 0 \) and constant quark mass \( m_q = m_{q0} \) (\( q = u, d, s \)),
shows a range of (144.297, 157.3634) MeV for the constant $B^{1/4}$ in bag model. If we require the energy per baryon less than the mass of nucleons $E/n_b = 939$ MeV, we can derive the same (meta)stable range of $B$ as the result in Ref. 3. The vertical axis shows the previous range of (154.8278, 156.1655) MeV for $\sqrt{D}$ in QMDD model 15. The parameter pairs $(D^{1/2}, B^{1/4})$ under the solid line or above the dashed line will not be adopted for stability purpose. If $B^{1/4} = 150$ MeV, a range for $D$ can be found. For baryon number $A=200$, the three lines of energy per baryon are shown respectively for $D^{1/2}=80, 40$ and 0 in Fig. 2. The full dots denote the zero pressure points, which are exactly the minima of the energy per baryon. The corresponding Fermi momenta, indicated on the right axis, are also affected by the mass parameter $D$.

We choose the parameters $B^{1/4} = 144$ MeV, $D^{1/2} = 120$ MeV in calculating the data for Figs. 3 and 4. The energy per baryon and electric charge per $A^{2/3}$ as a function of baryon number $A$ for CFL strangelets is given in Fig. 3. The strange current mass $m_{\text{si}}$ is taken to be 80, 120 MeV respectively. We find the net electric charge $Z \approx 0.15A^{2/3}$, which is half of the result in the pure bag model 15. In the same figure, the energy of the neutral ordinary strangelets is also plotted by a dash-dotted line. When we solve Eq. (12) we find that there are three different solutions for the common Fermi momentum. In Fig. 4 the charge of the other two solutions are given. The charge of CFL slet-2 (dashed lines) is proportional to cubic-root of baryon number, while CFL slet-3 (dotted lines) is nearly neutral. The corresponding Fermi momenta are also indicated on the right axis. Specially, given $A = 200$ we can get the smaller radius $R = 5.55$ with $p_F \approx 267$ MeV, which is marked by slet-1 in...
Fig. 5. The energy per baryon of CFL strangelets as functions of the radius. The mechanically stable radius, marked by full dots and given numerically by $R_{\text{stable}}$, is located at the minimum, where the pressure is exactly zero. For a fixed baryon number ($A = 200$ for the figure), there are three solutions: Figure (a), (b), (c) respectively for slet-1, slet-2, and slet-3.

The net positive charge $Z$ satisfies $Z/A^{2/3} \approx 0.15$ as mentioned above. The corresponding high $p_F$ is the usual case previously obtained by other authors. With the Fermi momentum decreasing we can find the other negatively charged solutions marked by slet-2 and almost neutralized solution marked by slet-3, respectively, in Fig. 5(b) and (c). They have larger radius $7.8$ fm. The corresponding Fermi momenta (42 MeV and 12 MeV or so) are smaller than that for slet-1 which will be interpreted in the next paragraph. It can been claimed that slet-1 is more stable at the same parameters. In Fig. 5(c) for slet-3, the charge to baryon number ration is $Z/A \approx -0.0006$, close to neutral state because $n_u \approx n_d \approx n_s$ is nearly satisfied.

There exist in literature two treatments concerning the deconfined quark phase. One is the bulk quark matter which is infinite in volume and without regard to boundary of the sphere. The other is finite-size lump quark matter, i.e. the so-called strangelets. The ordinary quark matter can support electrons within it and is charge neutral in $\beta$ equilibrium. CFL quark matter is automatically electrically neutral without participation of electrons because BCS-pairing requires the Fermi momenta of different flavor quarks to be equal. For ordinary strangelets or CFL strangelets, we should fall back on how the density of states depend on the quark numbers. The surface term has a negative contribution to the total number of massive quarks no matter how large the pairing Fermi momentum is. The curvature term, however, has a positive or negative contribution, depending on the values of quark mass and Fermi momentum. So ordinary strangelets is positively charged because the number of heavier, negatively charged strange quarks is suppressed compared to the number of lighter up and down quarks. Similarly, the CFL strangelets (slet-1) is positively charged as in Madsen’s work. The Fermi momenta of three cases...
satisfy (slet-3)<(slet-2)<(slet-1). In the case of smaller $p_F$, the surface and curvature terms make the number of strange quarks greater than, or nearly equal to, that of $u/d$ quarks. As a consequence, the slet-2 is negatively charged. For slet-3, the very small Fermi momentum make the surface and curvature effect so weak that slet-3 is nearly neutral. Although there are multiple solutions for a fixed baryon number, it is necessary to declare that the positively charged slet-1 is more stable than and the other two if the density dependence of the pairing parameters is not considered.

It should be emphasized that the common Fermi momentum “$p_F$” is only a fictional intermediate parameter in CFL matter. Different from the usual phase, the density of CFL strangelets depends not only indirectly on the chemical potential through $p_F$, but also directly on the chemical potential. As a consequence the density can be large even if $p_F$ is small as long as the chemical potential is big. Therefore, the density is still higher than, or at most, near the normal nuclear saturation density, because the chemical potential is big, though $p_F$ is small. In fact, the corresponding chemical potentials $\bar{\mu}$ of slet-2 and slet-3 are larger than that of slet-1. The results satisfy the validity condition that the ratio $\Delta/\mu$ is small. So the densities of new CFL strangelets “slet-2” and “slet-3” are still in reasonable range. We can still distinguish them through the difference of the electric charge $Z$ per $A^{1/3}$ and chemical potentials $\bar{\mu}$ in Fig. 6. The smaller the baryon number, the more visible the difference is. The Coulomb energy of the positively charged slet-1 has been shown in Fig. 7. With increasing baryon number $A$, the trend becomes placid. The effect on the free energy per baryon will not exceed 0.02 MeV. For slet-2 and slet-3 it will be much smaller than 0.001 MeV.

It should be noted that the new solutions, slet-2 and 3, have an unusual common
momentum. When two paired quarks have a very small momentum, their global behavior may looks like a boson. Therefore, the new solution may indicate that boson condensation or diquark condensate appears to some extent, and the formation mechanism needs further investigations in the future.

4. Summary

We have studied CFL strangelets within the framework of the quark mass density-dependent model. We have add the Coulomb interaction to the charged strangelets. It is found that the positive net charge is $Z/A^{2/3} \approx 0.15$, nearly half of the previous result in the pure bag model. Importantly, with decreasing Fermi momentum, we can find the other two solutions which have larger radius than the ordinary solution of CFL phase. The new solutions are slightly negatively charged or nearly neutral. Although with small Fermi momentum, the strangelets have large chemical potential. Due to the pairing effect, they have a comparable density with dense matter. The charge to baryon number of the positively charged strangelets is smaller than previously found, while that of the negatively charged strangelets is nearly proportional in magnitude to the cubic-root of the baryon number. However, present results depend on the parameter choice, and so, further studies are needed.

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