Quark condensates and strange quark matter

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Based upon recent studies on quark condensates, we investigate strange quark matter at zero temperature and find that the mass parametrization popularly used for u/d quarks with the mass-density-dependent model is just a first-order approximation to a more general formula, whereas the corresponding formula for s quarks has to be modified, which leads to the result that the strangeness fraction in strange quark matter can exceed u/d fraction. This strangeness excess may have a negative influence on the search for strangelets in ultrarelativistic heavy ion collisions. [S1063-651X(97)04707-7]

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I. INTRODUCTION

Since the conjecture by Witten [1] that objects, consisting of roughly equal numbers of u, d, s quarks, might be absolutely stable and so may serve as the true QCD ground state, a great amount of investigations have been carried out into the stability and properties of strange quark matter (SQM) [2]. Most of them are based on the MIT bag model [3] in which SQM is absolutely stable around the normal nuclear density for a wide range of parameters [4,5]. Other QCD motivated phenomenological models are also applied. By using the quark mass-density-dependent model [6], Chakrabarty *et al.* [7] obtain a very different result: only at very high densities does strange quark matter have the possibility of absolute stability. However, Benvenuto and Lugones [8] point out that this is the consequence of an incorrect thermodynamical treatment of the problem. They add to the expression of the energy density an extra term, arising from the baryon density dependence of the quark masses, and get similar results to those in the bag model.

A common point in literature is that the s fraction in SQM is almost equal to, but always less than the u/d fraction. This is due to the corresponding assumptions about quark masses. In the bag model, the masses of u, d, and s quarks are all constant or independent of density. The quark confinement is mimicked by the vacuum pressure B. In the previous version of the quark mass-density-dependent model, all the masses of u, d, and s quarks decrease with density, namely [7,8]

$$m_{u,d}=\frac{B}{3n_B}, \tag{1}$$

$$m_s=m_{s0}+\frac{B}{3n_B}, \tag{2}$$

where m_{s0} is the current mass of s quarks and n_B is the baryon number density.

Because these expressions are pure parametrizations, their applicable range of densities is completely unknown. We believe there exist links to quark condensates in SQM, which to some extent resembles the situation in normal nuclear matter where the in-medium hadron or meson masses vary with density according to the mass scaling relations. Based upon recent studies on quark condensates, we investigate SQM at zero temperature and find Eq. (1) is the first-order approximation of a more fundamental formula, while Eq. (2) has to be modified, which leads to the result that the s fraction in SQM can exceed the u or d fraction.

In the subsequent section, we will derive the relation between the quark mass and density, and then present our results in the study of SQM. The conclusion and discussion are given in Sec. III.

II. FORMULAS AND RESULTS

In the mass-density-dependent model, quark confinement is achieved by requiring [6]

$$\lim_{n_B\to 0} m_q=\infty, \tag{3}$$

where m_q is the quark mass. It is also popularly believed that the quark condensate varies with density. Therefore, there must exist a relation between the quark mass and quark condensate. Inspired by Eq. (3) and the following obvious equality:

$$\lim_{n_B\to 0}\frac{\langle\bar{q}q\rangle_{n_B}}{\langle\bar{q}q\rangle_0}=1, \tag{4}$$

where $\langle\bar{q}q\rangle_0$ and $\langle\bar{q}q\rangle_{n_B}$ are the quark condensates, respectively, in vacuum and in strange quark matter with baryon number density n_B, we propose the following concise expression:

$$\frac{m_q}{m_{q0}}=\frac{1}{1-\langle\bar{q}q\rangle_{n_B}/\langle\bar{q}q\rangle_0}, \tag{5}$$

where m_{q0} is a parameter. It may be regarded as the quark mass at the chiral restoration density. We will refer to it again a little latter.

Equation (5) is the simplest among relations satisfying the two basic requirements Eqs. (3) and (4). It is obviously different from the mass scaling relation

$$\frac{m_N^*}{m_N}=\frac{\langle\bar{q}q\rangle_\rho}{\langle\bar{q}q\rangle_0} \quad \text{or} \quad \frac{m_N^*}{m_N}=\frac{\langle\bar{q}q\rangle_\rho^{1/3}}{\langle\bar{q}q\rangle_0^{1/3}}$$

presently available for nucleons in normal nuclear matter. When the nuclear density $\rho\rightarrow 0$, one has $m_N^*/m_N\rightarrow 1$. This difference is due to the fact that quarks are confined whereas the nucleon is free in vacuum.

Since the validity of Eq. (5) is of crucial importance to the present investigation, we derive it formally from more fundamental principles below.

As is well known, the chiral symmetry of QCD is explicitly broken by the mass term $\bar{\Psi}M\Psi$ which fails to commute with flavor-SU(3) axial charges Q_5^a (a is the isospin index). The mass matrix M can be brought to diagonal form through flavor-mixing transformation. So, the quark mass contribution to the Hamiltonian may be written as

$$H_m=\int d^3x\sum_f m_f\bar{\Psi}_f\Psi_f, \tag{6}$$

where f is the flavor index with color index suppressed. For the study of SQM,

$$H_m=\int d^3x\{m_u\bar{u}u+m_d\bar{d}d+m_s\bar{s}s\}. \tag{7}$$

Regarding the SQM as a giant "nucleon," we can similarly define a "sigma term" [9]:

$$\sigma_{\pi-\mathrm{SQM}}=\frac{1}{3}\sum_{a=1}^{3}\{\langle\mathrm{SQM}|[Q_5^a,[Q_5^a,H_{\mathrm{QCD}}]]|\mathrm{SQM}\rangle -\langle 0|[Q_5^a,[Q_5^a,H_{\mathrm{QCD}}]]|0\rangle\}, \tag{8}$$

where H_{QCD} is the QCD Hamiltonian, $|\mathrm{SQM}\rangle$ is the state vector for the SQM at rest, and $|0\rangle$ is the vacuum. Substituting Eq. (7) for H_{QCD} and performing the double commutator, one has

$$\sigma_{\pi-\mathrm{SQM}}=2m_q\int d^3x\{\langle\mathrm{SQM}|\bar{q}q|\mathrm{SQM}\rangle-\langle 0|\bar{q}q|0\rangle\}, \tag{9}$$

where $m_q=\frac{1}{2}(m_u+m_d)$, $\bar{q}q=\frac{1}{2}(\bar{u}u+\bar{d}d)$, the term proportional to $(m_d-m_u)(\langle\mathrm{SQM}|\bar{d}d-\bar{u}u|\mathrm{SQM}\rangle)$, and $(m_d-m_u)\times(\langle 0|\bar{d}d-\bar{u}u|0\rangle)$ has been dropped.

Because the SQM is homogeneous, we replace the integration in Eq. (9) by multiplying the volume V of SQM

$$\sigma_{\pi-\mathrm{SQM}}=2m_qV\{\langle\mathrm{SQM}|\bar{q}q|\mathrm{SQM}\rangle-\langle 0|\bar{q}q|0\rangle\}. \tag{10}$$

Therefore,

$$m_q=\frac{-\sigma_{\pi-\mathrm{SQM}}/(2V\langle 0|\bar{q}q|0\rangle)}{1-\langle\mathrm{SQM}|\bar{q}q|\mathrm{SQM}\rangle/\langle 0|\bar{q}q|0\rangle}. \tag{11}$$

The numerator on the right-hand side (rhs) of the above equation has the same dimension with mass, and we denote it by m_{q0}. Consequently, Eq. (5) follows.

As for s quarks, we have a similar expression:

$$m_s=\frac{-\sigma_{s-\mathrm{SQM}}/(V\langle 0|\bar{s}s|0\rangle)}{1-\langle\mathrm{SQM}|\bar{s}s|\mathrm{SQM}\rangle/\langle 0|\bar{s}s|0\rangle}, \tag{12}$$

where $\sigma_{s-\mathrm{SQM}}$ satisfies the following equality [10,11]:

$$\sigma_{s-\mathrm{SQM}}=m_s\int d^3x\{\langle\mathrm{SQM}|\bar{s}s|\mathrm{SQM}\rangle-\langle 0|\bar{s}s|0\rangle\}. \tag{13}$$

By now we can see that m_{q0} in Eq. (5) is flavor dependent. It will be taken as parameters in the present investigation.

According to Eq. (4), the Taylor series of the relative condensate at zero density has the following general form:

$$\frac{\langle\bar{q}q\rangle_{n_B}}{\langle\bar{q}q\rangle_0}=1-\frac{n_B}{\alpha_q'}+\text{ higher orders in } n_B+\cdots, \tag{14}$$

where

$$\alpha_q'=-\left(\frac{d}{dn_B}\frac{\langle\bar{q}q\rangle_{n_B}}{\langle\bar{q}q\rangle_0}\right)^{-1}_{n_B=0}. \tag{15}$$

Because the ratio $\langle\bar{q}q\rangle_{n_B}/\langle\bar{q}q\rangle_0$ is expected to decrease with increasing density, α_q' is positive. Its dimension is the same with that of the density.

If the density is not too high, we can ignore all terms in Eq. (14) with orders in n_B higher than one and obtain

$$\frac{\langle\bar{q}q\rangle_{n_B}}{\langle\bar{q}q\rangle_0}\approx 1-\frac{n_B}{\alpha_q'}. \tag{16}$$

For u/d quarks, substituting Eq. (16) into Eq. (5), we get

$$m_{u,d}=\frac{\alpha_0'm_0}{n_B}\equiv\frac{\beta}{n_B}, \tag{17}$$

where we have denoted $\alpha_{u,d}'$ by α_0' and also ignored the mass difference between u and d quarks as usually done: $m_{u0}=m_{d0}\equiv m_0$. With

$$B=3\beta=3\alpha_0'm_0, \tag{18}$$

Eq. (1) is obtained naturally.

At zero temperature, the energy density of a ud canonical system is

$$\varepsilon=\frac{3}{\pi^2}\sum_{i=u,d}\int_0^{P_{f,i}}\sqrt{p^2+m_i^2}\,p^2dp, \tag{19}$$

where $P_{f,i}=(\pi^2n_i)^{1/3}$ is the Fermi momenta of quark flavor i that relate to the flavor number density n_i. If we substitute Eq. (1) into Eq. (19) and let $n_B\rightarrow\infty$, then obviously we can get

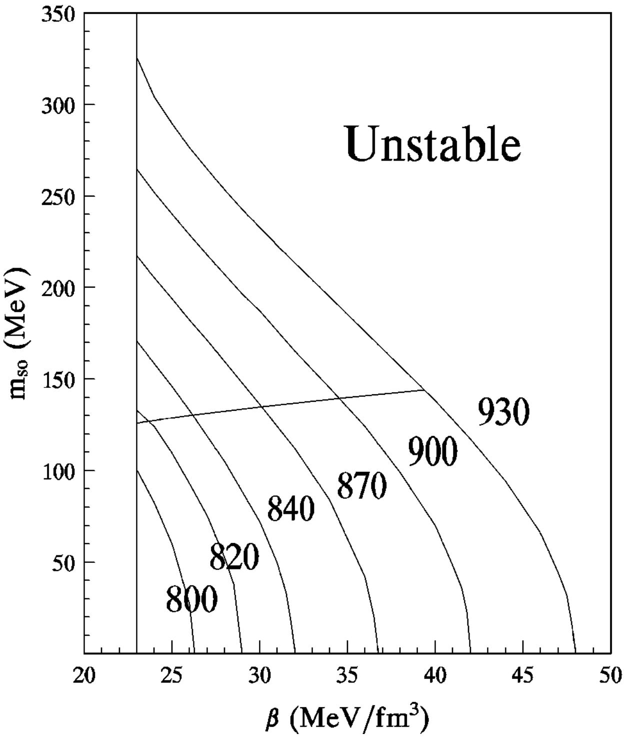

FIG. 1. Counters of fixed E/n_B in the β-m_{s0} plane. The stability region is where the energy per baryon is less than 930 MeV. The vertical line at the left is the minimum β for which two-flavor quark matter is unbound. The nearly horizontal line is the equal mass line where the numbers of three flavors are exactly the same.

$$\lim_{n_B \to 0} \varepsilon = B = \text{const} > 0. \tag{20}$$

So in this context, the parameter B can have the same meaning with the bag constant in the MIT bag model (not necessarily interpreted as the vacuum pressure, however).

For strange quarks, the lattice calculation [12] has shown that $\langle \overline{s s}\rangle_0$ is nearly one order of magnitude higher than $\langle \overline{u u}\rangle_0$ or $\langle \overline{d d}\rangle_0$ in the quark phase. So we expect $\langle \overline{s s}\rangle_{n_B}$ / $\langle \overline{s s}\rangle_0$ to be relatively small and can be neglected. From Eq. (5), we thus have

$$m_s \approx m_{s0}. \tag{21}$$

In fact, we can compare the two ratios to a leading order as such.

According to the model-independent research [11,13], we have in nuclear matter

$$\frac{\langle \overline{u u}\rangle_\rho}{\langle \overline{u u}\rangle_0} \quad \text{or} \quad \frac{\langle \overline{d d}\rangle_\rho}{\langle \overline{d d}\rangle_0} = 1 - \frac{\sigma_N}{m_\pi^2 f_\pi^2}\rho + \cdots, \tag{22}$$

where m_π is the pion mass, f_π is the pion decay constant, ρ is the medium density and σ_N is the pion-nucleon sigma term.

Similarly, one can write down [14]

$$\frac{\langle \overline{s s}\rangle_\rho}{\langle \overline{s s}\rangle_0} = 1 - \frac{\sigma_K}{m_K^2 f_K^2}\rho + \cdots, \tag{23}$$

where

$$\sigma_K = \frac{1}{2}(m_u + m_s)\langle \overline{u u} + \overline{s s}\rangle \tag{24}$$

is the kaon-nucleon sigma term and f_K is the kaon decay constant. In the chiral limit, $f_K = f_\pi$. Because the kaon mass drops rapidly with increasing density [15] while the π's remains remarkably unchanged [16], the difference between $m_K^2 f_K^2$ and $m_\pi^2 f_\pi^2$ would not be too much at high enough densities. On the other hand, we have already known [17]

$$\frac{\sigma_K}{\sigma_N} \approx 6 \gg 1. \tag{25}$$

This implies the lhs of Eq. (22) is much greater than that of Eq. (23). Therefore, we expand Eq. (5) as

$$\frac{m_q}{m_{q0}} = 1 + \frac{\langle \overline{q q}\rangle_{n_B}}{\langle \overline{q q}\rangle_0} + \left(\frac{\langle \overline{q q}\rangle_{n_B}}{\langle \overline{q q}\rangle_0}\right)^2 + \cdots \tag{26}$$

and only take its zero-order approximation for s quarks.

It is very clear from the derivation process that the condition for Eqs. (17) and (21) is: the density should be high enough but not too high. We will estimate the range as follows.

From Eq. (16), the highest density $n_{B\max}$ is less than α_0'. According to Eq. (18), we have

$$\alpha_0' = \frac{\beta}{m_0}. \tag{27}$$

The quantity β should be no less than 23 MeV fm^{-3} (the reason is to be recounted below). Taking the smallest permissible value for β and 10 MeV for m_0, we have $\alpha_0' = 2.3$ fm^{-3}. This indicates that the upper restriction of n_B is not strict.

The lowest density $n_{B\min}$ can be estimated like this. From [14,15]

$$\Delta m_K^2(\rho) = -\frac{\sigma_K}{f_K^2}\rho \approx 2 m_{K0} \Delta m_K, \tag{28}$$

we have

$$n_{B\min} \sim -2\frac{\Delta m_K}{m_{K0}}\frac{m_{K0}^2 f_K^2}{\sigma_K} \tag{29a}$$

$$\simeq 2\frac{m_{K0} - m_\pi}{m_{K0}}\frac{m_\pi^2 f_\pi^2}{\sigma_K} \tag{29b}$$

$$\simeq 2\left(1 - \frac{m_\pi}{m_{K0}}\right)\frac{\sigma_N}{\sigma_K}\frac{m_\pi^2 f_\pi^2}{\sigma_N}, \tag{29c}$$

where $m_{K0} \approx 495$ MeV is the kaon mass in free space. An extrapolation from low-energy pion-nucleon scattering data gives [18]: $\sigma_N = (45 \pm 8)$ MeV. Taking the center value 45 MeV for σ_N, and 6 for the ratio σ_K/σ_N, together with $m_\pi \approx 140$ MeV, $f_\pi \approx 93$ MeV, we get: $n_{B\min} \approx 0.12$ fm^{-3}.

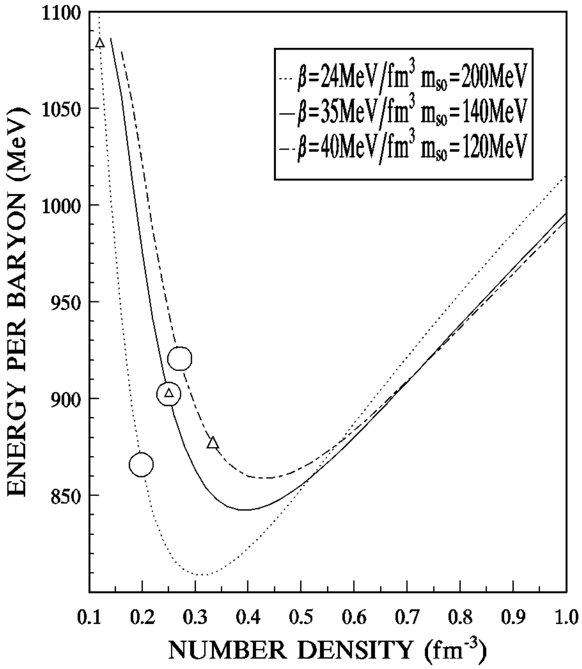

FIG. 2. The energy per baryon vs baryon number density. The zero pressure occurs at the points marked with ○. The triplicate points (see text) are marked with △.

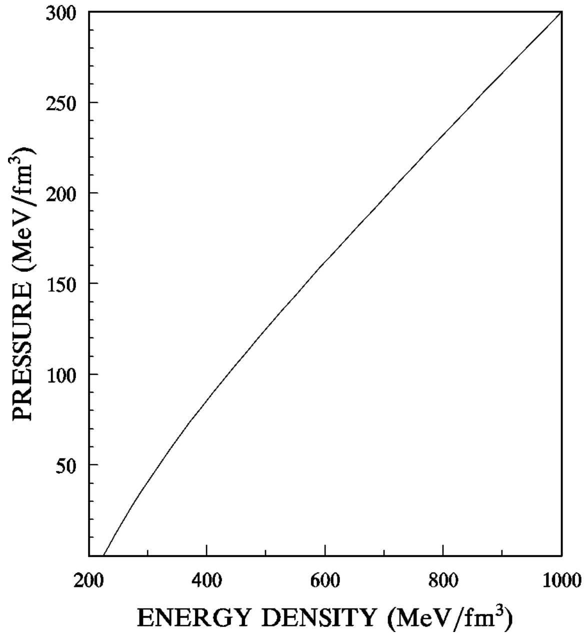

FIG. 3. The equation of state for parameters $\beta = 34.9$ MeV fm^{-3}, $m_{s0} = 140$ MeV. It asymptotically shows a similar behavior to the ultrarelativistic case as should be expected.

Equations (17) and (21) indicate that u and d quark masses decrease with density whereas s quark mass remains constant for the range of density considered. This conclusion gets support from a recent work with the chiral color dielectric model [19].

Equation (17) is identical to the corresponding Eq. (1) directly used by the previous authors [7,8] while Eq. (21) is completely not. It is this difference that might make the s fraction exceed the u/d fraction.

We assume the SQM to be a Fermi-gas mixture of u, d, s quarks and electrons with chemical equilibrium maintained by the weak interactions:

$$d,s \leftrightarrow u+e+\overline{\nu_e}, \quad s+u \leftrightarrow u+d.$$

Neutrinos enter and leave the system freely. For a given n_B, the chemical potentials $\mu_i (i=u,d,s,e)$ are determined by the following equations [4]:

$$\mu_d=\mu_s\equiv\mu, \tag{30}$$

$$\mu_u+\mu_e=\mu, \tag{31}$$

$$n_B=\frac{1}{3}(n_u+n_d+n_s), \tag{32}$$

$$\frac{2}{3}n_u-\frac{1}{3}n_d-\frac{1}{3}n_s-n_e=0, \tag{33}$$

where

$$n_i=\frac{g_i}{6\pi^2}(\mu_i^2-m_i^2)^{3/2}, \tag{34}$$

which is derived from the relation

$$n_i=-\frac{\partial\Omega_i}{\partial\mu_i} \tag{35}$$

with

$$\Omega_i=-\frac{g_i}{48\pi^2}\left[\mu_i(\mu_i^2-m_i^2)^{1/2}(2\mu_i^2-5m_i^2) + 3m_i^4\ln\frac{\mu_i+\sqrt{\mu_i^2-m_i^2}}{m_i}\right], \tag{36}$$

where g_i is the degeneracy factor with values 6 and 2, respectively, for quarks and for electrons, $m_{u,d}$ and m_s to be replaced by Eqs. (17) and (21).

As in Ref. [8], the pressure and energy density are given, respectively, by

$$P=\sum_i\left(-\Omega_i+n_B\frac{\partial\Omega_i}{\partial n_B}\right), \tag{37}$$

$$E=-P+\sum_i \mu_i n_i. \tag{38}$$

Since the baryonic matter is known to exist in the hadronic phase, we require β to be such that the ud system is unbound. This constrains β to be bigger than 23 MeV fm^{-3}, i.e., at $P=0$, $E/n_B>930$ in order not to contradict standard nuclear physics. On the other hand, we are interested in the possibility that SQM may be bound, i.e., at $P=0$, $E/n_B<930$, which gives an upper bound 325 MeV to m_{s0}. Also, when $m_{s0}\rightarrow 0$, β approaches its maximum value 48 MeV fm^{-3}.

In Fig. 1, we give the counters of fixed E/n_B in the β-m_{s0} plane. The vertical line at the left is the minimum β for which nonstrange quark matter is unbound. The ''stability window'' is also trianglelike but wider than that in Ref. [8]. A noticeable feature not seen before is that there appears

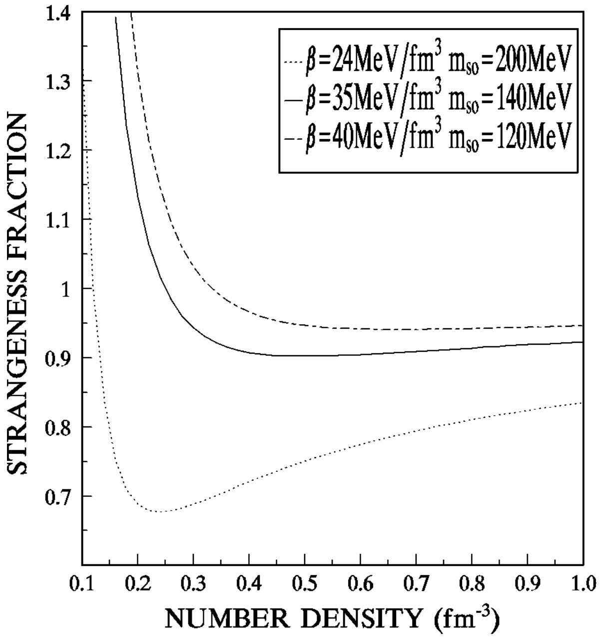

FIG. 4. The strangeness fraction (n_s/n_B) vs baryon number density (n_B). When the density is less than β/m_{s0}, the strangeness fraction is greater than 1.

another line where electrons are not present and the numbers of u, d, and s quarks are exactly the same. This is due to their equal masses on the line. When E/n_B is not too low (>813 MeV), the corresponding counters intersect the equal mass line.

The energy per baryon vs baryon number density is given in Fig. 2 for three pairs of parameters. On each line, there are two special points marked, respectively, by ○ and △. The former is the zero pressure point. The latter is the triplicate point where the s content is just a triplet. For case 2, the two points coincide with each other. However, they generally do not correspond to the minimum energy per baryon.

The resulting equation of state is plotted in Fig. 3. Because it is insensitive to parameters, we have only chosen one parameter pair: $\beta=34.9$ MeV fm^{-3} and $m_{s0}=140$ MeV.

In Fig. 4, we give the dependence of the s fraction on densities. It asymptotically tends to the result obtained with a previous version of the mass-density-dependent model but is not so at low densities. When the density is smaller than β/m_{s0}, the s fraction exceeds the u/d fraction.

III. CONCLUSION AND DISCUSSION

We have obtained the important Eq. (5) which reflects the relation between the quark mass and quark condensate. When the relative condensate is taken to first-order terms in density, the popularly used mass parametrization formula for u/d quarks is naturally obtained whereas the corresponding formula for s quarks has to be modified.

With the quark condensate results in normal nuclear matter, we estimate that the applicable range of Eqs. (17) and (21) is (0.12–2.3) fm^{-3}. When applied to study SQM, the results are similar to those obtained before, except that the s content in SQM can exceed u or d content.

This quite unusual result may give an explanation for the failure of searches for strangelets in ultrarelativistic heavy ion collisions. Unlike the situation in the interior of superdense stars [20] where the time scale is enough to establish flavor equilibrium by the weak interaction as mentioned above, the strangeness content in QGP is believed via the strangeness enrichment due to the early black-body radiation of more kaons ($q\bar{s}$) than antikaons ($\bar{q}s$) off the fireball [21]. The probability to produce $\bar{q}q$ pairs can be calculated by [22]

$$|M|^2=\exp\left(-\frac{\pi m^2}{\kappa}\right), \tag{39}$$

where κ is the string tension in the color forcefield. The string tension of 1 GeV/fm leads to a suppression of the strange quarks as compared to u/d quarks. This may suppress the strangeness enrichment process to such an extent that the strangeness fraction does not possibly exceed the u/d fraction.

However, this should not discourage experimental searches for strangelets or cold QGP formation. After all, the quark condensate in SQM is a new domain which needs more investigations in its own right. In our present study, we have resorted to the corresponding results in normal nuclear matter and only considered the lowest-order approximation. The higher-order contribution is not available up to now. Also, other factors, for example, a strong magnetic field [23], may change the configuration of the constituents in SQM.

It should be mentioned that the strangeness excess was also observed by Farhi and Jaffe in their early work [see Fig. 2(a) in Ref. [4]]. Because of the large coupling constant and small strange quark mass, the phenomenon was not taken seriously.

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