Strange quark matter with dynamically generated quark masses

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Abstract

Bulk properties of strange quark matter (SQM) are investigated within the SU(3) Nambu–Jona-Lasinio model. In the chiral limit the model behaves very similarly to the MIT bag model which is often used to describe SQM. However, when we introduce realistic current quark masses, the strange quark becomes strongly disfavored, because of its large dynamical mass. We conclude that SQM is not absolutely stable.

PACS: 12.39.Ba; 12.39.Fe
Keywords: Strange quark matter; dynamical quark masses
The properties of strange quark matter (SQM) and in particular the conjecture that SQM could be the absolute ground state of strongly interacting matter \cite{1,2} have attracted much attention in nuclear physics and astrophysics. Whereas the empirical stability of ordinary nuclei excludes the existence of absolutely stable non-strange quark matter (NSQM) this argument does not hold for SQM because the decay of nuclei into SQM would involve higher-order weak interaction processes, associated with a very long lifetime. In 1984 Farhi and Jaffe \cite{3} investigated the question of absolutely stable SQM within the MIT bag model \cite{4}. Treating the bag constant and the strange quark mass as free parameters they found a reasonable window for which SQM is stable and NSQM is unstable compared with an $^{56}$Fe-nucleus. Almost 15 years later bag model calculations still play a key role in this field (for reviews see \cite{5}-[8]).

In ref. \cite{9} we have investigated non-strange quark matter within the Nambu–Jona-Lasinio (NJL) model \cite{10} in mean field approximation. For vanishing (current) quark masses and with the appropriate choice of parameters we found that quark matter in the NJL mean field behaves very similarly to quark matter in an MIT bag. For the description of SQM we have to introduce finite quark masses, at least in the strange quark sector. In this case the correspondence between the two models becomes less accurate, with the NJL mean field behaving in a much more complex way. This is related to the fact that the bag constant in the NJL model is not a phenomenological input parameter, like in a bag model, but a dynamical property, which has its origin in the spontaneous breaking of chiral symmetry. This motivated us to extend the model of ref. \cite{9} to investigate bulk properties of strange quark matter within the NJL model.

The three-flavor NJL model has been discussed by many authors, e.g. \cite{11}-[21]. In this article we adopt the Lagrangian of ref. \cite{21}:

$$\mathcal{L} = \bar{q}(i\partial - \hat{m}_0)q + G \sum_{k=0}^{8} \left[ (\bar{q}\lambda_k q)^2 + (\bar{q}i\gamma_5\lambda_k q)^2 \right] - K \left[ \text{det}_f(\bar{q}(1 + \gamma_5)q) + \text{det}_f(\bar{q}(1 - \gamma_5)q) \right].$$

(1)

Here $q$ denotes a quark field with three flavors, $u$, $d$ and $s$, and three colors. $\hat{m}_0 = \text{diag}(m_0^u, m_0^d, m_0^s)$ is a $3 \times 3$ matrix in flavor space. We restrict ourselves to the isospin-symmetric case, $m_0^u = m_0^d$. The Lagrangian contains a $U(3)_L \times U(3)_R$-symmetric four-point interaction term and a six-point interaction, which is a determinant in flavor space and which breaks the $U_A(1)$-symmetry. $G$ and $K$ are constants with dimension energy$^{-2}$ and energy$^{-5}$, respectively.

The model is not renormalizeable and we have to specify a regularization scheme for divergent integrals. For simplicity we use a sharp cut-off $\Lambda$ in 3-momentum space. Besides $\Lambda$ we have to fix four parameters, namely the coupling constants $G$ and $K$ and the current quark masses $m_0^u$ and $m_0^s$. In most of our calculations we will adopt the parameters of ref. \cite{21}: $\Lambda = 602.3$ MeV, $G\Lambda^2 = 1.835$, $K\Lambda^5 = 12.36$, $m_0^u = 5.5$ MeV and $m_0^s = 140.7$ MeV. These parameters have been determined by fitting $f_\pi$, $m_\pi$, $m_K$ and $m_{\eta'}$ to their empirical values, while the mass of the $\eta$-meson is underestimated by about 6%.

Due to the interactions the quarks acquire dynamical masses $m_i$ which are in general different from their current masses $m_i^0$ and depend on temperature and density. In the following we consider quark matter at $T = 0$ and zero or non-zero baryon number.
density \( \rho_B = \frac{1}{3} n_B = \frac{1}{3}(n_u + n_d + n_s) \), with \( n_i = \langle \bar{q}_i q_i \rangle \). Then, in mean field (Hartree) approximation, the dynamical quark masses are solutions of the gap equation

\[
m_i = m_0^i - 4G \langle \bar{q}_i q_i \rangle + 2K \langle \bar{q}_j q_j \rangle \langle \bar{q}_k q_k \rangle ,
\]

with \( (q_i, q_j, q_k) \) being any permutation of \((u, d, s)\). The quark condensates \( \langle \bar{q}_i q_i \rangle \) are given by

\[
\langle \bar{q}_i q_i \rangle = -\frac{3}{\pi^2} \int_{p_{F_i}}^{\Lambda} p^2 dp \frac{m_i}{\sqrt{m_i^2 + p^2}} .
\]

They depend on the Fermi momenta \( p_{F_i} = (\pi^2 n_i)^{1/3} \) and again on the dynamical masses \( m_i \). Thus, for given quark number densities \( n_i \), eqs. (2) and (3) have to be solved self-consistently. After that we can calculate the energy density \( \varepsilon \) and the total pressure \( p \) of the system:

\[
\varepsilon = \sum_{i=u,d,s} \frac{3}{\pi^2} \int_0^{p_{F_i}} p^2 dp \sqrt{m_i^2 + p^2} - (B - B_0) , \quad p = -\varepsilon + \sum_{i=u,d,s} n_i \sqrt{m_i^2 + p_{F_i}^2} .
\]

Here \( B \) denotes the bag pressure,

\[
B = \sum_{i=u,d,s} \left( \frac{3}{\pi^2} \int_0^{\Lambda} p^2 dp \left( \sqrt{m_i^2 + p^2} - \sqrt{m_0^2 + p^2} \right) - 2G \langle \bar{q}_i q_i \rangle^2 \right) + 4K \langle \bar{u}u\rangle \langle \bar{d}d\rangle \langle \bar{s}s\rangle ,
\]

and \( B_0 = B|_{n_u=n_d=n_s=0} \). The latter was introduced in eq. (4) in order to ensure \( \varepsilon = p = 0 \) in vacuum.

The existence of a bag pressure term in eq. (3) suggests that there is a certain connection between the present model and standard bag model descriptions of quark matter. Let us have a closer look at this point. It is useful to begin with the chiral limit, i.e. \( m_0^i = 0 \) for all flavors. With \( \Lambda, G \) and \( K \) as specified above, chiral symmetry is spontaneously broken in vacuum with equal dynamical masses \( m_u = m_d = m_s = 310.6 \) MeV and a bag pressure \( B_0 = 57.3 \) MeV/fm\(^3\). (There are other solutions of the gap equation but they are energetically less favored.) Now we consider \( SU(3) \)-symmetric quark matter, \( n_u = n_d = n_s = \rho_B \). At low densities the dynamical quark masses decrease with density but they remain finite. However, when \( \rho_B \) exceeds a critical value of \( \sim 0.9 \rho_0 \) chiral symmetry gets restored and the quark masses and condensates vanish. Consequently the bag pressure is also zero in this regime and we get from eq. (4)

\[
\varepsilon = \frac{9}{4\pi^2} p_F^4 + B_0 , \quad p = \frac{3}{4\pi^2} p_F^4 - B_0 .
\]

This means the system behaves exactly like a gas of massless non-interacting quarks with three flavors and three colors inside a large MIT bag with a bag constant \( B_0 \). It should be kept in mind, however, that eq. (5) only holds in the regime of vanishing dynamical quark masses. For the present parameters it turns out that the minimal energy per baryon number \( E/A = \varepsilon/\rho_B \) is indeed found inside this region. Thus eq. (5) is appropriate to describe the system in the vicinity of the point of stability.

For comparison we consider non-strange isospin-symmetric quark matter, \( n_s = 0 \), \( n_u = n_d = \frac{3}{2} \rho_B \), still in the chiral limit. In this case strange and non-strange quarks
behave differently with density, and chiral symmetry is only partially restored at high densities: When $\rho_B$ exceeds $\sim 1.0\rho_0$ the non-strange quarks become again massless but the strange quark remains massive with a mass $m_s^*$ which is 154.2 MeV in our example. Hence the bag pressure remains also finite, $B = B^* = 3.7$MeV/fm$^3$, and the system now behaves like a bag model for a gas of massless quarks with two flavors and three colors but with an effective bag constant $B_{eff} = B_0 - B^*$, which is smaller than the bag constant $B_0$ in the SU(3)-symmetric case.

Thus, in contrast to a naive bag model, the effective bag constant which describes the behavior of NJL quark matter depends on the flavor composition of the matter. This is also shown in fig. 1 where various properties of isospin-symmetric quark matter are plotted as functions of the fraction of strange quarks, $r_s = n_s/n_B$. With the present parameters stable quark matter exists in the region $0 < r_s < 0.48$. As discussed in ref. [3] these solutions correspond to the high-density phase of a first-order chiral phase transition in equilibrium with the non-trivial vacuum. Throughout this paper we will denote the dynamical quark masses of these stable matter solutions by $m_i^*$ and the corresponding bag pressure by $B^*$. In fig. 1(a) the solid line indicates the energy per baryon number, $E/A$ of stable quark matter. $m_i^*$ and the effective bag constant $B_{eff} = B_0 - B^*$ are shown in fig. 1(b) and (c), respectively. For $0.12 < r_s < 0.48$ the system behaves like a bag model with massless quarks and a bag constant $B_{eff} = B_0$. For $r_s < 0.12$ the effective bag constant is somewhat smaller and the strange quarks become massive. In panel (a) we also show the bag model result for $E/A$ with a bag constant $B_{BM} = B_0$ and massless quarks (dashed line). For $0.12 < r_s < 0.48$ it agrees of course exactly with the NJL result, but also for $r_s < 0.12$ the deviations are relatively small. In fig. 1 all NJL curves stop at $r_s \simeq 0.48$. For $r_s > 0.48$ the NJL quark matter is unstable against evaporation of free massive s-quarks. In fact, this kind of phase separation is very typical for asymmetric matter and has been discussed for nuclear matter [22, 23] as well as for models of the deconfinement phase transition [24, 25]. However, the evaporation of free quarks is of course unphysical and reflects the missing confinement of the NJL model.

With finite current quark masses $m_i^*$ things change considerably. Fig. 2 shows the same quantities as fig. 1, but now with $m_i^u = m_i^d = 5.5$ MeV and $m_i^s = 140.7$ MeV [21]. We now find stable quark matter for $0 < r_s < 0.79$. In vacuum the dynamical masses of the up- and down-quarks are increased by $\sim 60$ MeV to 367.6 MeV and that of the strange quark by $\sim 240$ MeV to 549.5 MeV. Perhaps more important for us is the well-known fact that the quark condensates do not completely vanish even at higher densities. As a consequence, the dynamical quark masses in stable quark matter, $m_i^*$, are still well above the current masses $m_i^0$ (see fig. 2(b)). For non-strange isospin-symmetric quark matter ($r_s = 0$) we find $m_i^u = 52.6$ MeV and $m_i^d = 464.4$ MeV. Since on the other hand the Fermi energy in this system is only 361.1 MeV the conversion of a non-strange quark into a strange quark is obviously energetically not favored. Here the notion of dynamical quark masses is crucial: In a naive bag model one would assign the strange quark its current mass which in our example is more than 200 MeV lower than the Fermi energy.

The difference becomes obvious in fig. 2(a) where the energy per baryon number of stable quark matter is plotted as a function of $r_s$. The solid line indicates the NJL result while the dotted line corresponds to a bag model calculation with $B_{BM} = 105.2$ MeV/fm$^3$ (the effective bag constant at $r_s = 0$) and quarks with $m_u = m_d = 5.5$ MeV.
and $m_s = 140.7$ MeV, our current quark masses. Whereas the latter has a minimum at a finite fraction of strange quarks, $r_s = 0.27$, the former is a strictly rising function of $r_s$: additional strangeness is always disfavored in the NJL calculation. In fig. 2(a) we also show the result of a bag model calculation with $m_u = m_d = 52.6$ MeV and $m_s = 464.4$ MeV, our dynamical quark masses at $r_s = 0$ (dashed-dotted line). This curve rises even steeper with $r_s$ than the NJL result. The reason is that in the NJL calculation the dynamical strange quark mass $m_s^* = m_s$ drops with $r_s$ (see fig. 2(b)). Therefore the penalty for adding strangeness becomes smaller than it would be with a constant mass $m_s^*(r_s = 0)$, even though this effect is partially compensated by the fact that the effective bag constant $B_{eff} = B_0 - B^*$ rises with $r_s$ for not too large $r_s$.

$B_{eff}$ is shown in fig. 2(c). It varies much stronger with $r_s$ than in the chiral limit and is much larger in magnitude. Already at its minimum, at $r_s = 0$, it is about twice as large. To major extent this is the reason why the energy per baryon number at this point is almost 200 MeV larger than in the chiral limit. Both, the magnitude of the effective bag constant and its strong variation with $r_s$, have their origin in the fact that the bag pressure is very sensitive to the dynamical quark masses. As a consequence the vacuum bag pressure $B_0$ is now 291.7 MeV/fm$^3$, five times the chiral limit value, and, since the strange quark mass in stable matter varies, the corresponding bag pressure $B^*$ varies strongly with $r_s$.

It should be also noted that, even for fixed $r_s$, quark masses and the bag pressure are density dependent and the values for $m_i^*$ and $B_{eff}$ shown in fig. 2 only correspond to the densities of stable quark matter. As soon as one moves away from the minimum of $E/A$, $m_i$ and $B_0 - B$ change. Thus, whereas in the chiral limit NJL quark matter could be well described in terms of an MIT-like bag model with vanishing quark masses and a single bag constant (at least within a larger range of densities and strangeness fractions), it behaves in a much more complex way when we include finite current quark masses.

So far, focusing on the bag model aspect, we made some simplifying assumptions which we have to abandon for a more realistic study of SQM. In particular we have to take into account weak decays. This implies that we have to include electrons and (in principle) neutrinos. To large extent we will adopt the model of Farhi and Jaffe [3], with the MIT bag replaced by the NJL mean field. With the usual assumption that the neutrinos can freely leave the system, the matter is characterized by four (three quark- and one electron-) chemical potentials. They are related to the corresponding densities in the standard way:

$$ n_i = \frac{1}{\pi^2} \left( \mu_i^2 - m_i^2 \right)^{3/2} \theta(\mu_i^2 - m_i^2) \quad \text{for} \quad i = u, d, s \quad \text{and} \quad n_e = \frac{\mu_e^3}{3\pi^2} . \quad (7) $$

Here we neglected the electron mass. For the quarks one has to keep in mind that the masses entering the r.h.s. have to be calculated from the gap equation and are therefore density dependent themselves. Treating the electrons as a free gas, energy density and pressure of SQM are given by

$$ \varepsilon_{SQM} = \varepsilon + \frac{\mu_e^4}{4\pi^2} , \quad p_{SQM} = p + \frac{\mu_e^4}{12\pi^2} . \quad (8) $$

with $\varepsilon$ and $p$ as defined in eq. (4).
In chemical equilibrium maintained by weak interactions only two of the four chemical potentials are independent:

$$\mu_d = \mu_s = \mu_u + \mu_e.$$  \hspace{1cm} (9)

Furthermore we demand charge neutrality,

$$\frac{2}{3} n_u - \frac{1}{3} (n_d + n_s) - n_e = 0.$$ \hspace{1cm} (10)

Thus the system can be characterized by one independent quantity, e.g. the baryon number density $\rho_B$.

Our results are shown in fig. 3. Panel (a) shows the energy per baryon number $E/A = \varepsilon_{SQM}/\rho_B$ as a function of $\rho_B$. The solid line corresponds to the model described above, the dotted line to non-strange quark matter, where $\mu_s$ was set equal to zero by hand. Obviously the strangeness degree of freedom is only important at densities above $\sim 4 \rho_0$. This can also be seen in panel (c) where the fractions $r_i = n_i/n_b$ of the various particles are plotted. Since electrons (dotted line) play practically no role (and are hardly visible in the plot) the fraction of u-quarks (dashed-dotted) is fixed by charge neutrality (eq. (10)) to $r_u \simeq 1/3$. For $\rho < 3.85 \rho_0$ the fraction of strange quarks, $r_s$, (solid) is zero and thus $r_d$ must be the remaining $2/3$ (dashed). For $\rho > 3.85 \rho_0$ $r_s$ becomes non-zero and $r_d$ drops accordingly.

The fact that there is no strangeness at lower densities is again due to the relatively large mass of the strange quark. The dynamical quark masses are plotted in fig. 3(b) as functions of $\rho_B$. For comparison we also show the chemical potential $\mu_s = \mu_d$ (dotted line). As long as $\mu_s \leq m_s$ the density of strange quarks is zero. For $\rho_B \lesssim 2 \rho_0$, $m_s$ drops with density. This can be mainly attributed to the decrease of $\langle \bar{u}u \rangle \langle \bar{d}d \rangle$ due to the rising density of non-strange quarks. At higher densities $\langle \bar{u}u \rangle \langle \bar{d}d \rangle$ is practically zero and $m_s$ stays almost constant until $n_s$ becomes non-zero and causes $|\langle \bar{s}s \rangle|$ to drop. This behavior is rather different from standard parameterizations which have been used in the literature to study SQM and which depend on the total baryon number density $\rho_B$ only \cite{20,27}.

At those densities where strange quarks exist, they lead to a reduction of the energy, as can be seen by comparing the solid curve in fig. 3(a) with the dotted one. However, the minimum of $E/A$ does not lie in this regime, but at a much lower density, $\rho_B = 2.25 \rho_0$. Here we find $E/A = 1102$ MeV. Compared with the energy per baryon in an iron nucleus, $E/A \simeq 930$ MeV, this is still very large. In this sense our results are consistent with the empirical fact that stable NSQM does not exist. However, since the energy of strange quark matter is even higher, our calculation predicts that also SQM is not the absolute ground state of strongly interacting matter.

This result is very robust with respect to changes of the model parameters. Obviously stable quark matter with finite strangeness is only possible if the in-medium strange quark mass $m_s^*$ is much lower than the value we obtained above. The easiest way to achieve this is to choose a lower value of the current mass $m_s^0$. If we leave all other parameters unchanged, $m_s^0$ should be at most $85$ MeV if we require $n_s \neq 0$ at the minimum of $E/A$. With this value we obtain much too low masses for $K$ and $\eta$ ($m_K \simeq 390$ MeV, $m_\eta \simeq 420$ MeV), while $E/A$ is still relatively large (1075 MeV). If we want to come down to $E/A \simeq 930$ MeV we have to choose $m_s^0 = 10$ MeV or, alternatively, $m_s^0 = 25$ MeV and $m_s^0 = m_d^0 = 0$. This is of course completely out of range.
We could also try to lower $m_s^*$ by choosing a smaller coupling constant $G$. (Since $\langle \bar{u}u \rangle \langle \bar{d}d \rangle$ is already very small at the densities of interest, $m_s^*$ is almost insensitive to $K$). However, with a lower $G$ the vacuum masses of all quarks drop and correspondingly the vacuum bag pressure $B_0$. This causes the minimum of $E/A$ to move to lower densities. It turns out that the corresponding chemical potential $\mu_s^*$ drops even faster than $m_s^*$ and it is thus not possible to obtain stable SQM in that way. To avoid the strong density decrease we could increase the coupling constant $K$ while decreasing $G$, e.g. in such a way that the vacuum masses of u- and d-quarks are kept constant. In order to get $n_s \neq 0$ at the minimum of $E/A$ we have to lower $G\Lambda^2$ to 1.5 and to increase $K\Lambda^5$ to 21.27, almost twice the value of our original parameter set. For these parameters the energy per baryon number is still 1077 MeV. On the other hand we cannot further decrease the ratio $G/K$ because that would flip the sign of the effective $q\bar{q}$ coupling in the pseudoscalar-flavor singlet channel, which is dominated by the combination $2G + \frac{2}{3}K(\langle \bar{u}u \rangle + \langle \bar{d}d \rangle + \langle \bar{s}s \rangle)$. In that case there would be no solution for the $\eta'$-meson in vacuum. Hence there seems to be no realistic way to find absolutely stable SQM within our model.

We could ask whether this result can change if we introduce additional interaction terms to the Lagrangian eq. (1). In particular vector interactions are known to be important in dense matter. However, since vector mean fields are repulsive, the energy per baryon number will be even larger than before and SQM remains strongly disfavored compared with ordinary nuclei. The same is true for the medium effects described in ref. [28] which also tend to increase the energy. In fact, these effects are complementary to ours because they are most important (and most reliable) at high densities, whereas our dynamical quark masses are most important at low densities. In ref. [19] the chiral phase transition was studied within a “scaled” NJL model. There is no six-point interaction in this model, but the different flavors are coupled through a dilaton field. The coupling was found to be weak except for rather low choices of the gluon condensate [20]. This makes it very unlikely to find absolutely stable SQM within that model, although a quantitative study of this question has not yet been done.

In summary, we investigated strange and non-strange quark matter within the NJL model. In the chiral limit the model behaves very similarly to the MIT bag model which has been mostly used in the literature to describe SQM. For realistic model parameters including finite current quark masses the NJL model shows a more complex behavior. This is due to the fact that bag pressure and effective quark masses are dynamical properties of the model and are therefore in general density dependent. We find that the dynamical mass of the strange quark stays very large at the relevant densities. As a consequence SQM is not favored compared with NSQM in the model and there seems to be no chance to find SQM with energies per baryon number lower than in ordinary nuclei. This almost rules out the original idea of absolutely stable SQM as a simple consequence of the Pauli exclusion principle. Of course we cannot completely exclude that more sophisticated mechanisms, like color superconductivity [29, 30], in particular the recently discovered mechanism of color-flavor locking [31], change our results. This is, however, beyond the scope of this letter.
Acknowledgement:

We would like to thank J. Schaffner-Bielich for having drawn our attention to the topic of strange quark matter.

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Figure 1: Properties of isospin-symmetric ($n_u = n_d$) stable quark matter as a function of $r_s = n_s/n_B$ in the chiral limit (parameters of ref. [23], but with $m_0^i = 0$). (a) Energy per baryon number. The result of the present model (solid line) is compared with a bag model calculation with bag constant $B_{BM} = 57.3 \text{MeV/fm}^3$ (dotted). (b) Dynamical masses of the strange (solid) and non-strange quarks (dotted). (c) Effective bag constant.
Figure 2: The same as fig. 1 but with finite current masses (parameters of ref. [21]). In panel (a) the result of the present model (solid line) is compared with bag model calculations with bag constant $B_{BM} = 105.2 \text{MeV/fm}^3$ and $m_u = m_d = 5.5 \text{MeV}$, $m_s = 140.7 \text{MeV}$ (dotted) or $m_u = m_d = 52.6 \text{MeV}$, $m_s = 464.4 \text{MeV}$ (dashed-dotted).
Figure 3: Properties of charge-neutral quark matter with electrons as a function of baryon number density $\rho_B$ (parameters of ref. [21]). (a) Energy per baryon number in chemical equilibrium ($\mu_s = \mu_d = \mu_u + \mu_e$, solid line) and for non-strange matter ($\mu_s = 0, \mu_d = \mu_u + \mu_e$, dotted). (b) Dynamical quark masses $m_u$ (dashed-dotted), $m_d$ (dashed) and $m_s$ (solid). The dotted line denotes the chemical potential $\mu_s = \mu_d$. (c) $r_i = n_i/n_B$ for the different particle species: up quarks (dashed-dotted), down quarks (dashed), strange quarks (solid), and electrons (dotted).