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Properties of strangelets in a new quark mass confinement model with one-gluon-exchange interaction

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The properties of strangelets at zero temperature with a new quark model that includes both the confinement and one-gluon-exchange interactions is studied in a fully self-consistent method. The charge and parameter dependence of the stability of strangelets are discussed. It is found that the one-gluon-exchange interaction lowers the energy of a strangelet, and consequently allows the strangelet to be absolutely stable. The stable strangelet radius in the present model is smaller in comparison with the absence of one-gluon-exchange interaction, and can thus be much less than that of a normal nucleus with the same baryon number, according to the strength of the confinement and one-gluon-exchange interactions.

strangelets, quark confinement model, one-gluon-exchange effect

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1 Introduction

Researchers have studied the question of whether the ordinary nuclear matter is actually metastable and decay slowly into quark matter of lower energy. Bodmer [1] first suggested a low energy nuclear state called "collapsed nuclei". Subsequently, Witten [2] studied the stability of strange quark matter (SQM) consisting of approximately equal numbers of u, d and s quarks, suggesting that SQM could indeed be stable even at zero external pressure. Since then, SQM has been intensely studied in nuclear physics, astrophysics, and other related fields.

Small lumps of SQM are normally called strangelets [3–5], or for short, slets [6]. Because of their low charge to baryon number ratio, strangelets have been proposed to be a complex ingredient of ultrahigh energy cosmic rays [7,8]. Also, they could be found in the cooling products of energetic heavy ion collisions as a unique signature for the formation of quark gluon plasma [9]. Because of the obvious difficulties in working out QCD directly and strictly, phenomenological models are needed in exploring the properties of strangelets. The most well known phenomenological model is the MIT bag model, in which the fundamental assumption is that the vacuum has a constant energy density, that is, the bag constant B which provides a negative pressure to confine quarks. Thus the confinement is assumed to be achieved by the special property of the vacuum.

Employing the bag model, many researchers have studied the properties of strangelets and presented interesting results. For example, Farhi and Jaffe [3] found that strangelets in perfect weak equilibrium are slightly positively charged. It is found that the properties of strangelets have strong parameter dependence [10,11]. Greiner et al. [12,13] studied the possible formation of strangelets in heavy-ion collisions, and they

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have suggested that strangelets are most likely highly negatively charged [9]. Other researchers have noted the important effect of the curvature contribution [14–16]. Taking into account the electrostatic effects and Debye screening, and an arbitrary surface tension of the interface between quark matter and vacuum, Alford et al. [17] determined that there is a critical surface tension below which large strangelets are unstable to fragmentation.

Another phenomenological quark model incorporates confinement through the density dependence of quark masses (CDDM) [18,19]. In this model, the key points are to use a proper quark mass scaling and to adopt a self-consistent thermodynamic treatment.

The original authors suggested an inversely linear scaling for up and down quarks [20] which was later extended to include strange quarks [21]. With the inversely linear scaling [22,23], Zhang and Su [24] studied the stability of strangelets at both zero and finite temperatures. Because the inversely linear scaling was derived according to the basic assumption of the bag model, the results from it are, in general, similar to those from the bag model [25–27].

Based on the in-medium chiral condensates and linear confinement, an inversely cubic scaling was derived [28,29]. With this scaling, Wen et al. [30,31] investigated the properties of strangelets versus the electric charge and strangeness at both zero and finite temperatures.

On application of the quark mass scaling derivation procedure and the thermodynamic formulas suggested by Wen et al. [31] (WZP model), Modarres and Gholizade [32] calculated the thermodynamic properties of strange quark matter. They introduced the one-gluon-exchange interaction obtained from the Fermi liquid picture. However, the confinement was not included, or simply they did the same as in the bag model by adding a constant vacuum energy density.

In a recent work [33], we have studied the effect of one-gluon-exchange interactions on the properties of strange quark matter and strange stars where the finite-size effect is unimportant. For a strangelet, however, the finite-size does matter. Therefore, in the present paper, we consider the important finite-size effect on strangelets, when both the confinement and the one-gluon-exchange interactions are included in a fully self-consistent way. It is found that the one-gluon-exchange interaction lowers the energy of a strangelet, and thus making the strangelet to be more stable. The stable radius of a strangelet is thus smaller compared to the case without one-gluon-exchange interactions.

Hereinafter, we describe the new quark mass scaling, the consistent thermodynamic treatment, and properties of slets.

2 Quark mass scaling with one-gluon-exchange interaction

The original idea in the quark mass density-dependent model is that the quarks would become, with decreasing density, so massive that the vacuum is unable to support them, and thereby must be confined in a finite volume. Because the energy density of the system for a large particle mass m_q is $3m_q n_\mathrm{b}$ at lower baryon number density n_b while the bag model assumes a constant vacuum energy density B, Fowler et al. [20] naturally made them equal and gave the mass scaling of u and d quarks as $m_{u/d} = B/(3n_\mathrm{b})$. Subsequently, Chakrabarty et al. [21–23] extended this to include strange quarks as $m_s = m_{s0} + B/(3n_\mathrm{b})$, where m_{s0} is the current mass of strange quarks.

Based on the in-medium chiral condensates and linear confinement, an inverse cubic-root scaling was derived [28]. At zero temperature, it is

$$m_q = m_{q0} + \frac{D}{n_\mathrm{b}^{1/3}} \quad (q = u, d, s), \tag{1}$$

where m_{q0} is the current mass of quark flavor q, and D is a confinement parameter determined by stability arguments.

Eq. (1) considers merely the linear confinement interaction. Because the confinement interaction dominates at lower density, the scaling in eq. (1) describes quark matter accurately. With increasing density, however, perturbative interaction becomes more critical. We should thus include perturbative interaction at higher density.

As described elsewhere [33], a new quark mass scaling was obtained using the quark interaction

$$v(r) = \sigma r - \frac{4}{3}\frac{\alpha_\mathrm{s}}{r}, \tag{2}$$

where σ is the string tension and α_s is the running coupling constant of QCD. The first term on the right hand side is the linear confinement while the second term represents the one-gluon-exchange interaction. The new quark mass scaling can be expressed as:

$$m_q = m_{q0} + \frac{D}{n_\mathrm{b}^{1/3}} - C n_\mathrm{b}^{1/3}, \tag{3}$$

where $n_\mathrm{b} = \frac{1}{3}\sum_i n_i$ is the baryon number density, D is still the confinement parameter as in eq. (1), and C represents the strength of the one-gluon-exchange effect.

The last two terms on the right hand side of eq. (3) are, respectively, the contributions of the confinement and one-gluon-exchange interactions. They are the same for all flavors, representing the interaction between quarks. We thus specially denote them as:

$$m_\mathrm{I} = \frac{D}{n_\mathrm{b}^{1/3}} - C n_\mathrm{b}^{1/3}. \tag{4}$$

Because of quark confinement, the effective quark mass would be infinitely large that the quark can only be confined in a proper region. Therefore, the density-dependent form in eq. (4) can be understood as a Laurent series of the Fermi momentum which is proportional to the cubic root of the density, truncated to leading order in both directions. For the two coefficients D and C, they can be linked to known quantities such as the pion-nucleon sigma term, the chiral condensates, quark current masses and the coupling constant. One method to derive the linkage is provided elsewhere [33] where C is given to be proportional to the coupling constant of strong interactions. As known, the strong coupling constant is running. Therefore, C should also be density-dependent. However, the running speed is rather slow because the strong coupling is a logarithmically decreasing function [34]. On the other hand, the derivation must have some assumptions and approximations. We thus treat C as a phenomenological parameter with different values.

Previously [31], the squared root of the confinement parameter, $D^{1/2}$, is estimated in the range of $(156, 270)$ MeV [18]. Here, we employ the values $D^{1/2} = 160, 180, 200, 220$ MeV. For the strength parameter C of the one-gluon-exchange effect, we take $C = 0.1$, 0.6 as examples. Formally, if one takes $C = 0$, then the previous quark mass scaling in eq. (1) can be obtained. To compare with the previous results, we will also present the results by taking $C = 0$ at $D^{1/2} = 160$ MeV.

3 Self-consistent thermodynamic formulae

The special problem in studying strangelets is to include the finite-size effect, while the size of a strangelet is of crucial importance in analyzing the propagation and detection. Therefore, the proper density of state which includes the finite-size effect needs obtained initially. This has been done by applying the multi-expansion method originally comprised by Balian and Bloch [35], and later developed by others [3,14,24,36–38].

In the present model, the quark acts like a free particle with the equivalent mass m_i and the dispersion relation $\varepsilon_i = \sqrt{p^2 + m_i^2}$. Therefore, the energy density of a strangelet with radius R can be given in the multi-expansion approach by

$$E = \sum_i E_i + E_{\mathrm{C}}, \tag{5}$$

where m_i is the mass of flavor type i, R is the strangelet radius, and ν_i is the Fermi momentum corresponding to flavor i. The summation index i goes over all flavors considered. The thermodynamic contribution from flavor i is

$$E_i = E_i(\nu_i, m_i, R) = \int_0^{\nu_i} \sqrt{p^2 + m_i^2}\, n_i'(p)\mathrm{d}p, \tag{6}$$

where the density of state $n_i'(p)$ is given in the multi-expansion method by

$$\begin{aligned} n_i'(p) = &\frac{g_i}{2\pi^2}\left\{p^2 - \frac{3p}{2R}\arctan\left(\frac{m_i}{p}\right)\right. \\ &\left.+\frac{1}{R^2}\left[1 - \frac{3p}{2m_i}\arctan\left(\frac{m_i}{p}\right)\right]\right\}, \end{aligned} \tag{7}$$

with g_i being the degeneracy factor of particle type i, i.e. $g_u = g_d = g_s = 6$, $g_e = 2$. The three terms on the right hand side of eq. (7) are the volume term, surface term [4], and curvature term [14–16], respectively. The particle number density is

$$n_i = \int_0^{\nu_i} n_i'(p)\mathrm{d}p. \tag{8}$$

In principal, the Coulomb energy contribution should be included for a charged system. The corresponding contribution, E_{C} in eq. (5), is

$$E_{\mathrm{C}} = \frac{2}{15}\pi R^2 \alpha_{\mathrm{C}}(Q_{\mathrm{v}}^2 + 5Q^2), \tag{9}$$

where $\alpha_{\mathrm{C}} \approx 1/137$ is the fine structure constant and Q_{v} is the volume term of the total electric charge density Q, i.e., $Q = \sum_i q_i n_i$ and $Q_{\mathrm{v}} = \sum_i q_i n_{i,\mathrm{v}}$ with $q_u = \frac{2}{3}$, $q_d = -\frac{1}{3}$, $q_s = -\frac{1}{3}$ and $q_e = -1$.

From eqs. (5) and (8), and the fundamental thermodynamic laws, we can show that the pressure is given as:

$$P = \sum_i \left[-\Omega_i - \frac{R}{3}\frac{\partial \Omega_i}{\partial R} + n_{\mathrm{b}}\frac{\partial \Omega_i}{\partial m_i}\frac{\mathrm{d}m_i}{\mathrm{d}n_{\mathrm{b}}}\right] + P_{\mathrm{C}}, \tag{10}$$

where P_{C} represents the contribution from the Coulomb interaction. It should be pointed out that any model needs to satisfy the thermodynamic consistent conditions. For this to be fulfilled in the present model, the pressure is derived according to the fundamental thermodynamic differential equation, which gives an additional term because of the density dependence of the quark masses. The quantity Ω_i in eq. (10) is the free-particle contribution of the particle type i with a density dependent mass and an effective chemical potential. The concrete expression can be given by

$$\Omega_i = \Omega_i(\nu_i, m_i, R) = \int_0^{\nu_i}\left(\sqrt{p^2 - m_i^2} - \mu_i^*\right)n_i'(p)\mathrm{d}p. \tag{11}$$

Here we have introduced an effective chemical potential $\mu_i^* \equiv \sqrt{\nu_i^2 + m_i^2}$. In the quark mass density-dependent case, the actual chemical potential is

$$\begin{aligned} \mu_i = &\mu_i^* + \frac{4}{15}\pi R^2 \alpha_{\mathrm{C}} q_i\left(5Q + \frac{g_i \nu_i^2 Q_{\mathrm{v}}}{2\pi^2 n_i'(\nu_i)}\right) + \frac{1}{3}\sum_j \frac{\partial \Omega_j}{\partial m_j}\frac{\mathrm{d}m_j}{\mathrm{d}n_{\mathrm{b}}} \\ &-\frac{2\alpha_{\mathrm{C}} R^2 Q_{\mathrm{v}}}{45\pi}\sum_j \frac{q_j g_j \nu_j^2}{n_j'(\nu_i)}\frac{\partial n_j}{\partial m_j}\frac{\mathrm{d}m_j}{\mathrm{d}n_{\mathrm{b}}}. \end{aligned} \tag{12}$$

At finite temperature, the relevant results can be similarly given. For detailed information, one refers to the section V of ref. [31]. For a fully self-consistent thermodynamic derivation and explanation, one can refer to the appendix of ref. [18].

Here we present the results at zero temperature. For convenience, we define $x_i \equiv \nu_i/m_i$ and $y_i \equiv \mu_i^*/m_i$, and carry out the relevant integrations to give concrete expressions.

The particle number density in eq. (8) becomes thus

$$\begin{aligned} n_i = &\frac{g_i \nu_i^3}{6\pi^2} + \frac{3g_i m_i^2}{8\pi^2 R}\left[y_i^2 \arctan(x_i) - x_i\left(\frac{\pi x_i}{2} + 1\right)\right] \\ &+\frac{3g_i m_i}{8\pi^2 R^2}\left[y_i^2 \arctan(x_i) - x_i\left(\frac{\pi x_i}{2} - \frac{1}{3}\right)\right]. \end{aligned} \tag{13}$$

The energy density in eq. (6) and the pressure in eq. (10) respectively give

$$
\begin{aligned}
E_i = & \frac{g_i m_i^4}{8\pi^2}\bigg\{\frac{1}{2}\left[x_i(2x_i^2+1)y_i - \mathrm{arcsh}(x_i)\right] \\
& +\frac{1}{m_i R}[\pi - x_i y_i - 2y_i^3 \mathrm{arccot}(x_i) - \mathrm{arcsh}(x_i)] \\
& +\frac{1}{m_i^2 R^2}[\pi + x_i y_i - 2y_i^3 \mathrm{arccot}(x_i) + \mathrm{arcsh}(x_i)]\bigg\}, \quad (14)
\end{aligned}
$$

and

$$
\begin{aligned}
P = & \sum_i \frac{g_i m_i^4}{48\pi^2}\bigg\{(2x_i^2-3)x_i y_i + 3\mathrm{arcsh}(x_i) \\
& +\frac{2}{m_i R}\left[3\pi y_i - 4x_i y_i - 2\pi + 2\mathrm{arcsh}(x_i) - 2y_i^3 \mathrm{arccot}(x_i)\right] \\
& +\frac{1}{m_i^2 R^2}\left[\pi(3y_i - 2) - 2\mathrm{arcsh}(x_i) - 2y_i^3 \mathrm{arccot}(x_i)\right]\bigg\} \\
& -\sum_i \frac{g_i m_i^3 n_\mathrm{b}}{16\pi^2}\frac{\mathrm{d}m_i}{\mathrm{d}n_\mathrm{b}}\bigg\{4\left[\mathrm{arcsh}(x_i) - x_i y_i\right] \\
& -\frac{6}{m_i R}\left[\pi + x_i y_i - \pi y_i - \mathrm{arcsh}(x_i)\right] \\
& +\frac{1}{m_i^2 R^2}[3\pi y_i - 4\pi - 2x_i y_i - 4\mathrm{arcsh}(x_i) \\
& +2y_i^3 \mathrm{arccot}(x_i)]\bigg\} + P_\mathrm{C}, \quad (15)
\end{aligned}
$$

where the contribution from the Coulomb interaction P_C can be expressed as such

$$
\begin{aligned}
P_\mathrm{C} = & \frac{2}{9}\pi R^2 \alpha_\mathrm{C}\bigg[\frac{3Q_\mathrm{v}}{5\pi^2}\sum_i \frac{q_i g_i \nu_i^2}{n_i'(\nu_i)}\left(n_i + \frac{R}{3}\frac{\partial n_i}{\partial R} - \frac{\partial n_i}{\partial m_i}\frac{\mathrm{d}m_i}{\mathrm{d}n_\mathrm{b}} n_\mathrm{b}\right) \\
& +Q^2 - Q_\mathrm{v}^2\bigg], \quad (16)
\end{aligned}
$$

and the quark mass derivative with respect to the density is

$$\frac{\mathrm{d}m_i}{\mathrm{d}n_\mathrm{b}} = \frac{\mathrm{d}m_\mathrm{I}}{\mathrm{d}n_\mathrm{b}} = -\frac{D}{3n_\mathrm{b}^{4/3}} - \frac{C}{3n_\mathrm{b}^{2/3}}. \quad (17)$$

It should be noted that the Coulomb contribution P_C in eq. (16) differs from the corresponding equation given elsewhere [15], such that

$$P_\mathrm{C} = \frac{2}{45}\pi R^2 \alpha_\mathrm{C}(Q_V^2 + 5Q^2). \quad (18)$$

Only when the numbers of particles in the surface are not changed with the size at a fixed total baryon number, and quark masses do not depend on density, eq. (16) is reduced to eq. (18).

4 Properties of strangelets in the new quark mass scaling

In the early 1990s it was realized that electric field is screened over a distance inside strange matter so that the electron number density, is not zero inside a strangelet [39]. The β equilibrium requires $\mu_u + \mu_e = \mu_d = \mu_s$.

A strangelet, in principal, has electric charge at its surface due to the important finite-size effect [40]. If the baryon number is less than about $A \ll 10^7$, electrons can not exist within a strangelet because the Compton wavelength of the electron exceeds the size of the strangelet [41,42]. Therefore, in the present calculations, the electron is ignored and the chemical potential is treated as zero to achieve lower energy, that is, $\mu_e \backsimeq 0$. Thus, the chemical equilibrium condition becomes $\mu_u = \mu_d = \mu_s$. It should also be noted that the inclusion of electrons barely affects our final results.

If we use A, S and Z to represent the baryon number, strangeness and electric charge, and N_u, N_d and N_s to denote the number of u, d and s quarks respectively, then we have three equations as follows:

$$A = \frac{1}{3}(N_u + N_d + N_s), \quad (19)$$

$$Z = \frac{2}{3}N_u - \frac{1}{3}N_d - \frac{1}{3}N_s, \quad (20)$$

$$S = N_s. \quad (21)$$

For a mechanically stable strangelet, the pressure inside the strangelet must be zero, namely

$$P = 0, \quad (22)$$

from which the stable radius of the strangelet can be obtained. With a view to the pressure expression in eq. (15), the volume term, surface term, and curvature term are included.

To allow for a clearer depiction, we define the electric charge to baryon number ratio and strangeness fraction as $f_\mathrm{Z} = Z/A$ and $f_\mathrm{S} = S/A$, respectively. If the strangelet radius R is given, we can obtain the strangelet volume as $V = (4/3)\pi R^3$, then the baryon number density is $n_\mathrm{b} = A/V$, and from eq. (3) the quark mass m_i $(i = u, d, s)$ can be gained. From the expressions in eqs. (19)–(21) we have $N_u = A + Z$, $N_d = 2A - Z - S$ and $N_s = S$, then we obtain the number density n_i $(i = u, d, s)$. The fermi momentum ν_i $(i = u, d, s)$ can be derived by solving $n_i = 3N_i/(4\pi R^3)$ with the n_i expression in eq. (13).

For all the numerical calculations in this paper, we used $m_{u0} = 5$ MeV, $m_{d0} = 10$ MeV, and $m_{s0} = 100$ MeV. Other parameters are specifically presented.

We calculated the charge dependence of the mechanically stable radius and energy per baryon E/n_b of the strangelet first. Figures 1 and 2 show the stable radii and energy per baryon of the strangelet as functions of their charge fraction with $A = 50$ and $f_S = 0.12$ as was done elsewhere [30]. To compare with the previous results, we also plot the curve in the case of $C = 0$ at $D^{1/2} = 160$ MeV. We can see that the stable radius and energy per baryon are not monotonic functions of the ratio of charge to baryon number. There is a minimum point on each curve, which depends on the values of the parameters C and D. We note, from Figures 1 and 2, that the stable radius of a strangelet in the present model is smaller

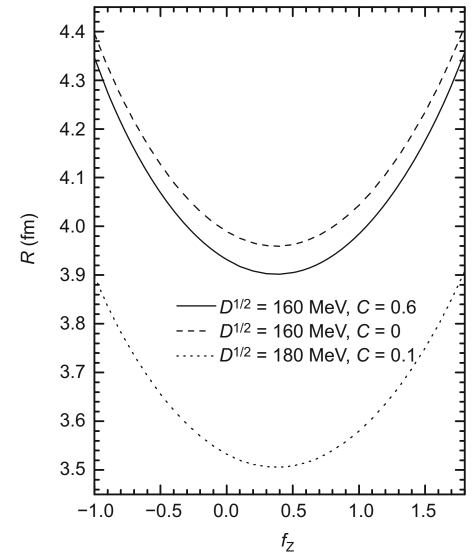

Figure 1 Mechanically stable radii of strangelets as functions of the ratio of charge to baryon number at $A = 50$ and a given strangeness fraction $f_S = 0.12$. A minimum is found on each curve.

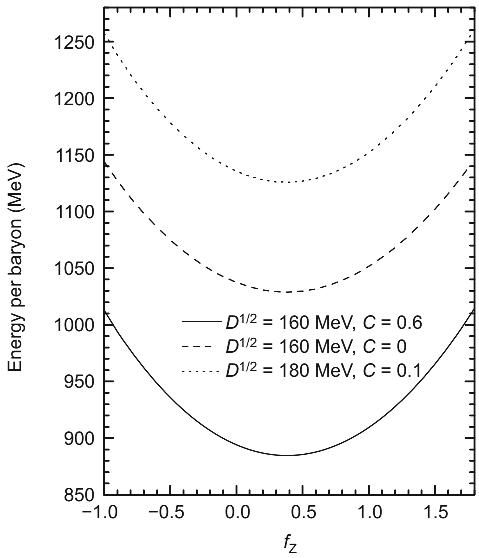

Figure 2 Energy per baryon of a strangelet with the fixed baryon number $A = 50$ and the strangeness fraction $f_S = 0.12$ versus the ratio of charge to baryon number f_Z. A minimum is found on each curve.

and the energy becomes lower in comparison with absence of one-gluon-exchange effect.

Figures 3 and 4 show how the parameters affect the radius and energy per baryon. The figures were given with $A = 10, Z = 1, S = 10$. From the two figures we can draw the conclusion that both the stable radius and the energy per baryon are monotonic functions of the parameter C for a given D value, namely, they are decreasing functions of C.

Figure 5 gives the baryon number dependence of the mechanically stable radius. To compare with the ordinary nuclei radius relation $R = 1.12A^{1/3}$ fm, we also plotted it with a dash-dotted line. Comparing with previous works where the one-gluon-exchange effect was not involved, the mechanically stable radii of a strangelet is smaller. It is readily apparent that the radius of a strangelet in the present model is also smaller than that of a nucleus with the same baryon number.

Figure 6 shows the baryon number dependence of the energy per baryon and the charge fraction of stable strangelets. Generally, they decrease with increasing baryon number A. We find that once the one-gluon-exchange interaction is assumed, then both the energy per baryon and the charge fraction are reduced.

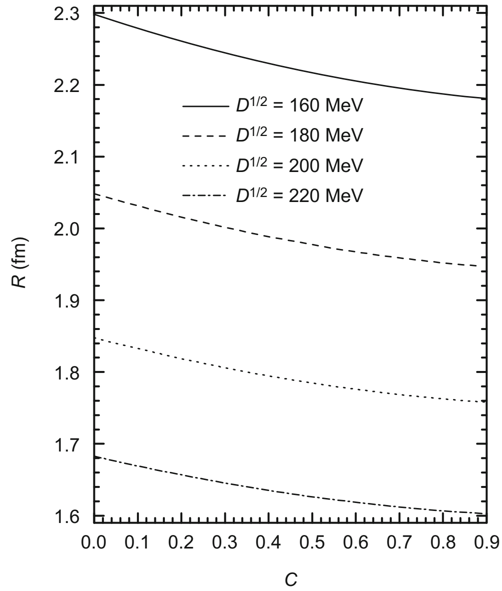

Figure 3 Mechanically stable radius of a strangelet as functions of the parameter C.

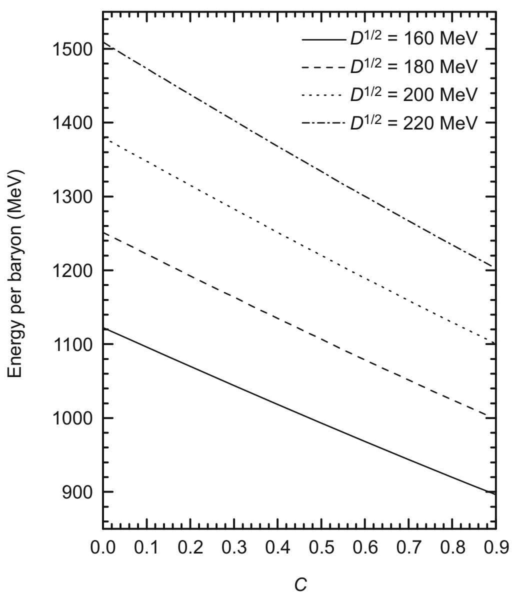

Figure 4 Energy per baryon of a strangelet as functions of the parameter C.

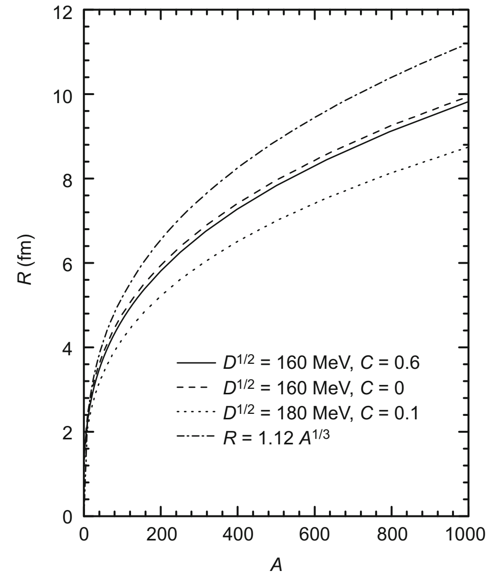

Figure 5 Baryon number dependence of the strangelet radius. For comparison purpose, the ordinary nuclei radius $R = 1.12A^{1/3}$ fm is also shown with the dash-dotted line.

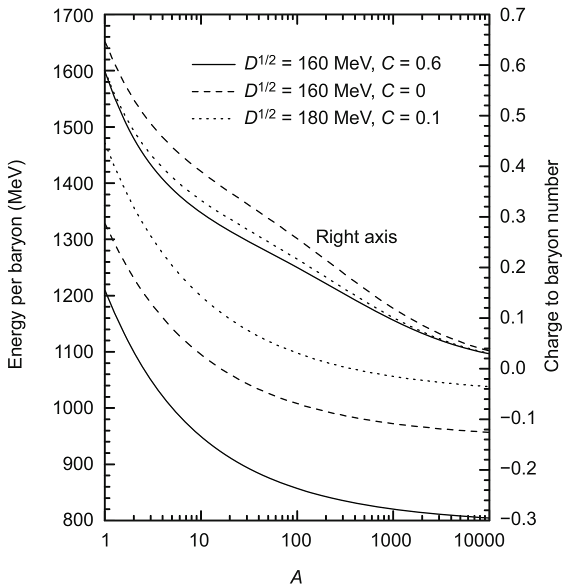

Figure 6 Baryon number dependence of the energy per baryon and the charge fraction of stable strangelets. Both the energy per baryon and the charge fraction are decreasing functions of the baryon number.

In general, gluons contribute to a quark system because the quarks interact with each other via gluon exchanges. Presently we have tried to include the one-gluon-exchange effect on the properties of zero-temperature strangelets at leading order. In further studies at finite temperature, a free gluon contribution term can be included. Concurrently, the quark masses needs to be dependent on both the density and the temperature as well.

5 Summary

On application of a new quark mass scaling which includes one-gluon-exchange effect and considering the weak chemical equilibrium, we have studied the thermodynamical properties of strangelets. The results show that the radius and energy per baryon of a strangelet is not a monotonic function of the charge fraction. The stable strangelet radius in the present model is smaller than in other cases [30] where the important one-gluon-exchange interaction was not included. It is also smaller than that of a normal nucleus with the same baryon number, which might be relevant to the propagation and detection of strangelets [43–45].

We have also investigated the baryon number dependence of the energy per baryon, and the charge fraction. It was found that the one-gluon-exchange interaction lowers the energy of a strangelet, and consequently allows the strangelet to be absolutely stable.

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