Medium effects on the surface tension of strangelets in the extended quasiparticle model

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We propose a modification of the finite size effects due to the effective bag function in the extended quark quasiparticle model with a running coupling constant. The bag function is associated with the quark chemical potential and the radius of strangelets. Considering the medium effects, the surface tension should be redefined with an additional term described by the surface term of the bag function. With the increasing baryon number of stable strangelets, it is found that the coupling strength becomes stronger while the surface tension decreases in the vicinity of 35 MeV fm^{-2} for strangelets of the baryon number greater than 10^3. The comparison with the bag model is shown and the distinction for smaller strangelets is very clear.

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I. INTRODUCTION

Strange quark matter (SQM) is postulated to be the true ground state of quantum chromodynamics (QCD) with a new strangeness freedom. It is argued by Witten that the strange quark phase could have survived the early Universe [1]. Subsequently, the new matter form was searched for in QCD phase transitions of relativistic heavy ion collision experiments and also probed in compact objects, that is, strange stars or hybrid stars. At extremely high densities, the three flavors of *u*, *d,* and *s* quarks can be treated on an equal footing and consequently form the so-called color-flavor-locked phase [2,3]. In contrast, there was a proposal that the strange matter could not have survived the early Universe because it could boil and form bubbles of hadronic gas [4].

Up to now, lattice QCD is incapable of getting significant results for the case of finite chemical potential. Thus we have to search for effective models for QCD. There are many phenomenal models used in investigating SQM or strangelets, lumps of SQM. One is the famous MIT bag model. Calculations based on the bag model have shown many aspects of SQM and strangelets [5]. The bag constant is often introduced phenomenologically with the expectation that it simulates nonperturbative corrections. Many different values have been taken for it. It has been usually assumed to be temperature or density dependent. For example, a Gaussian parametrization for the density dependence is adopted in Ref. [6].

It is also possible that the quark mass, pion mass, and so on, change with temperature and density. These quantities are essentially due to medium effects. The effective mass produced has been extensively discussed, for example, within the Nambu-Jona-Lasinio (NJL) model [7] and within a quasiparticle model [8]. A chiral phase transition and dynamical symmetry breaking are demonstrated in the NJL model. Quark matter deconfinement properties have been successfully investigated using the density-dependent bag constant in the quasiparticle model. In the literature, some people have constructed the quasiparticle model in terms of the temperature- and density-dependent bag constant and made progress in studying the nonperturbative QCD model at zero or small density. Most importantly in this paper, the quasiparticle model with running coupling constant is used to investigate the medium effect of SQM at moderate densities. Calculation programs with phenomenological models should be well guided by the spirit of QCD. Recently, we followed the original idea [8] and demonstrated the derivation of a chemical- potential-dependent bag function and extended it to the study of finite-size strangelets [9]. Furthermore, the thermodynamic treatment was improved based on the chemical-potential-dependent running coupling constant used to study the finite-size effects.

It is well known that the surface tension is a fundamental parameter for the finite-size effects and the nucleation phase transition [10]. Berger discussed the surface tension based on a modification to the density of states [11] and concluded that the maximum value of the surface tension (105 MeV)^3 is insufficient to save the SQM. Huang *et al.* reported a calculation of the surface tension in quenched QCD on lattices [10]. Recently, Alford *et al.* gave a critical value of surface tension below which larger strangelets are unstable and the quark star surface will fragment into a crystalline crust of charged strangelets [12]. In our new version of the quark quasiparticle model [9], the confinement mechanism is described by a bag function of the strangelets' radius and chemical potential. It is different from the conventional bag constant, so the surface tension should be modified accordingly. In this paper, we shed new light on the finite-size effect by modification of the radius- and chemical-potential-dependent bag function. We develop a new formula to calculate the surface tension.

This paper is organized as follow. In the next section, we briefly give the thermodynamic treatment in the extended quasiparticle model at zero temperature based on the running coupling constant. The effective masses of quarks and the energy of strangelets are calculated. Then we present the finite-size effects of strangelets and calculate the coupling constant, the surface energy per baryon, and the surface tension versus the baryon number. The last section is a short summary.

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II. THE EXTENDED QUASIPARTICLE MODEL

The main purpose of this paper is to study the properties of quark masses in the deconfined SQM. For the three flavors of quarks, nonzero current quark masses are applied and the exact chiral symmetry is broken. In addition to the current mass, the effective quasiparticle mass should be introduced to include the interaction effect in the quasiparticle approach. For the medium dependence of the quark quasiparticle model, the effective quark mass m_i^* is derived at the zero-momentum limit of the dispersion relation following from the effective quark propagator by resumming one-loop self-energy diagrams in the hard dense loop approximation [8]. The in-medium effective mass of quarks can thus be expressed as [8,13,14]

$$m_i^* = \frac{m_i}{2} + \sqrt{\frac{m_i^2}{4} + \frac{g^2\mu_i^2}{6\pi^2}} \ (i = u, d, s), \tag{1}$$

where m_i is the current mass of corresponding quarks. In the present paper we take $m_{u0} = 5$ MeV, $m_{d0} = 10$ MeV, and $m_{s0} = 120$ MeV for up, down, and strange quarks, respectively. In fact, the quasiparticle idea can be followed back to the work by Fowler *et al.* [15] showing that the particle mass may change with the environment parameters. Following the original ideas, the quark mass density-dependent model was studied by Chakrabarty *et al.* [16]. Accordingly, as a phenomenological method, our quasiparticle model uses a similar treatment. They are apparently different in approach but equally satisfactory in result. The in-medium screening mass in Eq. (1) is merely a model assumption on the quasiparticle mass in the present treatment and cannot be justified field-theoretically.

In our present model, the quantity g is related to the strong coupling α_s by the equation $g = \sqrt{4\pi\alpha_s}$, which was treated as a constant in the previous work [8]. By replacing the constant coupling by an effective running coupling $G(\mu)$, the nonperturbative effects can be accommodated. The running of the coupling should be determined by the QCD renormalization equation group [17]. One can use an analytical expression phenomenologically, such as $g(Q/\Lambda) = \frac{4\pi}{3}\sqrt{\frac{1}{\ln(Q^2/\Lambda^2)} - \frac{1}{1-Q^2/\Lambda^2}}$ in Ref. [18]. One can also apply the μ and T dependence of the approximate running coupling constant in lattice QCD simulations [14,19]. In this paper, we adopt the following approximate expression for the running quantity $g(\mu)$ [20,21]:

$$g^2(T = 0, \mu) = \frac{48\pi^2}{29}\left[\ln\left(\frac{0.8\mu^2}{\Lambda^2}\right)\right]^{-1}, \tag{2}$$

where Λ is the QCD scale-fixing parameter. In the present calculation, the Λ values are taken to be 180 and 200 MeV, respectively. With the running coupling constant rather than a fixed one, the in-medium mass will result in a new thermodynamic treatment. In particular, a density-dependent bag function has to be derived. In the following sections, we discuss the cases of both bulk SQM and strangelets.

A. Strange quark matter with running coupling constant

In order to derive the thermodynamic formulas, let us start from the quasiparticle contribution to the chemical thermodynamic potential density for bulk matter, that is,

$$\Omega_i = -\frac{d_i T}{2\pi^2}\int_0^\infty \ln\left(1 + e^{-(\sqrt{p^2+m_i^{*2}}-\mu_i)/T}\right)p^2 dp. \tag{3}$$

After integration at zero temperature, we get the analytic result, which is called $\Omega_{V,i}$,

$$\Omega_{V,i} = -\frac{d_i}{48\pi^2}\left(\mu_i\nu_i\left(2\mu_i^2 - 5m_i^{*2}\right) + 3m_i^{*4}\ln\frac{\mu_i+\nu_i}{m_i^*}\right), \tag{4}$$

where $\nu_i \equiv \sqrt{\mu_i^2 - m_i^{*2}}$ is the Fermi momentum of the particle type i. The total thermodynamic potential density of bulk SQM is written as

$$\Omega(\{\mu_k\}) = \sum_i \Omega_{V,i}(\mu_i, m_i^*) + B_V^*, \tag{5}$$

where the sum goes over u, d, and s quarks. The effective bag function reads

$$B_V^* = \sum_i B_{V,i}(\mu_i) + B_0. \tag{6}$$

Once the thermodynamic potential is known, the pressure P and energy density E are calculated by the normal thermodynamic relations

$$P = -\Omega, \quad E = \Omega + \sum_i \mu_i n_i. \tag{7}$$

In order to satisfy the fundamental thermodynamic equation

$$E = -P + \sum_i \mu_i\frac{\partial P}{\partial\mu_i}, \tag{8}$$

we must require the relation [9] $\partial P/\partial m_i^* = -\partial\Omega_i/\partial m_i^* - \partial B_{V,i}/\partial m_i^* = 0$. So $B_{V,i}$ in Eq. (6) is determined by the integration formula

$$\begin{aligned} B_{V,i}(\mu_i) &= -\int_{m_i^*}^{\mu_i}\frac{\partial\Omega_{V,i}}{\partial m_i^*}\frac{\partial m_i^*}{\partial\mu_i}d\mu_i \\ &= -\int_{m_i^*}^{\mu_i}\frac{d_i m_i^*}{4\pi^2}\left(\mu_i\nu_i - m_i^{*2}\ln\frac{\mu_i+\nu_i}{m_i^*}\right) \\ &\quad\times\frac{\partial m_i^*(g(\mu_i),\mu_i)}{\partial\mu_i}d\mu_i, \end{aligned} \tag{9}$$

where the integration constant is included in the B_0 of Eq. (6). The adopted value of $B_0^{1/4} = 145$ MeV is equivalent to the conventional bag constant. Because of the complication in analytic derivations, we can get only numerical results. It should be noted that the quark effective mass is a function of μ_i and $g(\mu_i)$. From the quark mass scale, we calculate the relation of m_i^* versus the baryon chemical potential μ_B divided by 3 in Fig. 1 (there $\mu_i \simeq \mu_B/3$). The coupling constant indicated by the dotted line on the right axis decreases with increasing chemical potential. Accordingly, the quark mass

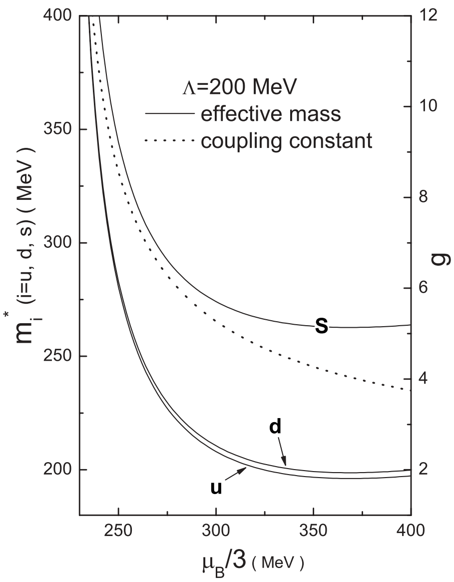

FIG. 1. The in-medium effective mass $m_i^*(i = u, d, s)$ and coupling constant as functions of the chemical potential. The solid curve on the left axis indicates the effective mass, which is in agreement with the lattice QCD data and satisfies the requirement of the asymptotic freedom of QCD. On the right axis is the coupling constant g indicated by a dotted curve.

is also a decreasing function. In any case, it is physically required by QCD that the dynamical quark mass must decrease with increasing μ_i. Considering the range of g $(0, \sqrt{6}\pi)$ in Ref. [8], we can realize it by adjusting the parameter Λ. With the chemical potential decreasing, the quark mass will become much larger so that the vacuum can support it, which indicates that our model is consistent with the quark confinement feature of QCD. In order to keep the coupling constant as a positive real number in the logarithm form of Eq. (2), there should be a lower limit value for the chemical potential.

For the different values of QCD parameter, $\Lambda = 180$ MeV (solid line) and 200 MeV (dashed line), the energy per baryon of SQM is plotted in Fig. 2. The minimum values of the are curves located at the stable zero-pressure points marked by open circles; this is the thermodynamic self-consistency requirement and will be satisfied in the calculation of strangelets. For the sake of simplicity, the symbols for the zero-pressure points are omitted in the following figures.

B. The properties of strangelets

To study strangelets, the special problem is to include the finite-size effect. We do this by applying the multireflection expansion, originally suggested by Balian and Bloch [22], and later developed by Madsen [23] and Farhi, Berger, and Jaffe [5,11] among others. We express the quasiparticle contribution to the thermodynamic potential density of a strangelet at zero temperature by inserting a density of state into Eq. (3) as

$$\Omega_i = \int_0^\infty \left(\sqrt{p^2 + m_i^{*2}} - \mu_i\right) \frac{dN_i}{dp}(p, m_i^*, R)\, dp, \qquad (10)$$

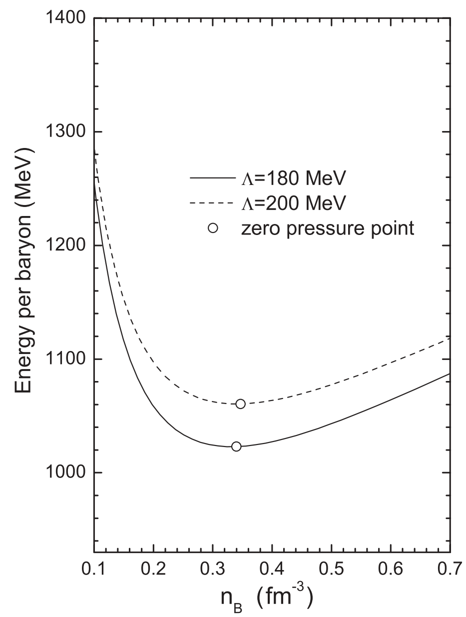

FIG. 2. The energy per baryon versus the baryon number density for different values of the QCD scale parameter Λ. The zero pressure stable points are marked by open circles “○”.

where the density of states $\frac{dN_i}{dp}$ is given in the multiexpansion approach by

$$\frac{dN_i}{dp}(p, m_i, R) = d_i\left(\frac{p^2 V}{2\pi^2} + p f_S(x_i) S + f_C(x_i) C\right). \qquad (11)$$

Here $x_i \equiv m_i^*/p$, the area $S = 4\pi R^2$, and the extrinsic curvature $C = 8\pi R$ for a sphere. The functions $f_S(x_i)$ [5,11] and $f_C(x_i)$ [23] are

$$f_S(x_i) = -\frac{1}{4\pi^2}\arctan(x_i) \qquad (12)$$

and

$$f_C(x_i) = \frac{1}{12\pi^2}\left(1 - \frac{3}{2x_i}\arctan(x_i)\right). \qquad (13)$$

After integrating Eq. (10), the thermodynamic potential density can be divided into three parts with respect to the radius dependence as $\Omega_i = \Omega_{V,i} + \Omega_{S,i} + \Omega_{C,i}$. In addition to the conventional volume part $\Omega_{V,i}$ in Eq. (4), the other two parts have the following explicit expressions:

$$\Omega_{S,i} = -\frac{d_i}{8\pi^2 R}\Bigg(m_i^{*3}\ln\frac{\mu_i + \nu_i}{m_i^*} + \mu_i^3 \arctan\frac{\nu_i}{m_i^*} - \frac{\pi}{2}(\mu_i + 2m_i^*)(\mu_i - m_i^*)^2 - 2m_i^*\mu_i\nu_i\Bigg), \qquad (14)$$

$$\Omega_{C,i} = -\frac{d_i}{8\pi^2 R^2}\Bigg[-m_i^{*2}\ln\frac{\mu_i + \nu_i}{m_i^*} + \frac{\mu_i^3}{m_i^*}\arctan\frac{\nu_i}{m_i^*} - \frac{\pi}{2}\left(\frac{\mu_i}{m_i^*} + 2\right)(\mu_i - m_i^*)^2\Bigg]. \qquad (15)$$

Now we can give the energy density and pressure according to the basic thermodynamic relations [9],

$$E = \Omega + \sum_i \mu_i n_i = \sum_i \left(\Omega_i - \mu_i \frac{\partial \Omega_i}{\partial \mu_i} + B_i(\mu_i, R) \right) + B_0, \tag{16}$$

$$\begin{aligned} P &= -\Omega - \frac{R}{3}\frac{\partial \Omega}{\partial R} \\ &= -\sum_i \left(\Omega_i + \frac{R}{3}\frac{\partial \Omega_i}{\partial R} + B_i(\mu_i, R) + \frac{R}{3}\frac{\partial B_i(\mu_i, R)}{\partial R} \right) - B_0. \end{aligned} \tag{17}$$

For the detailed derivation of the above energy density and pressure formulas, please see Ref. [24]. In the finite-size case, $B_i(\mu_i, R)$ depends not only on the chemical potential but also on the radius R of the strangelet. In Ref. [9], we divided $B_i(\mu_i, R)$ into three parts according to the R dependence:

$$B_i(\mu_i, R) = B_{V,i} + \frac{3}{R} B_{S,i} + \frac{6}{R^2} B_{C,i}, \tag{18}$$

where the three terms, $B_{V,i}$, $B_{S,i}$, and $B_{C,i}$, correspond respectively to the volume, surface, and curvature terms. Similarly to the derivation of Eq. (9), we can obtain the surface and curvature terms as follows:

$$B_{S,i}(\mu_i) = -\int_{m_i^*}^{\mu_i} \frac{\partial}{\partial m_i^*} \left(\Omega_{S,i} + \frac{R}{3} \frac{\partial \Omega_{S,i}}{\partial R} \right) \frac{\partial m_i^*}{\partial \mu_i} d\mu_i, \tag{19}$$

$$B_{C,i}(\mu_i) = -\int_{m_i^*}^{\mu_i} \frac{\partial}{\partial m_i^*} \left(\Omega_{C,i} + \frac{R}{3} \frac{\partial \Omega_{C,i}}{\partial R} \right) \frac{\partial m_i^*}{\partial \mu_i} d\mu_i, \tag{20}$$

where the partial derivatives with respect to the radius can be obtained from Eq. (14). All the interaction constants have been collected into the parameter B_0 in Eq. (16).

In Fig. 3 for the fixed baryon number $A = 100$, the energy per baryon of the strangelet is plotted as a function of radius R. The lines for $\Lambda = 180$ MeV (solid line) and $\Lambda = 200$ MeV (dashed line) are located above the dotted line, which is calculated by the bag model. It can be concluded that the Λ values have a proper influence on the result of the calculation of the strangelet energy. In fact, because of the medium effects, the effective quark masses are larger than the current constant masses, and so the energy of strangelets will be bigger than in the bag model.

In Fig. 4 we demonstrate the coupling constant as an increasing function of the baryon number A of the strangelet. The solid and dashed curves, respectively, are for $\Lambda = 180$ and 200 MeV. According to the stable conditions of weak equilibrium and zero pressure, we can calculate the chemical potentials μ_u and μ_s; we then find that a larger value of coupling strength g is required for increasing baryon number and it gradually approaches the horizontal dotted lines, which are the corresponding values for the bulk SQM limit.

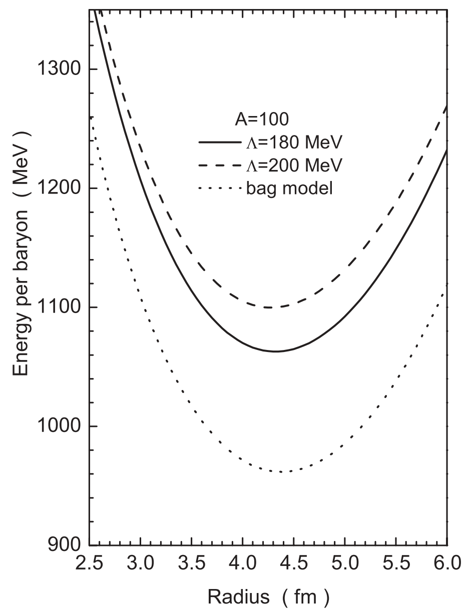

FIG. 3. The energy per baryon as a function of the radius of a strangelet with baryon number $A = 100$ for $\Lambda = 180$ MeV (the solid curve) and $\Lambda = 200$ MeV (the dashed curve). The bag model result (dotted line) is used for comparison.

III. MODIFICATIONS OF THE SURFACE TENSION WITH MEDIUM EFFECTS

For strangelets with finite baryon number, especially small strangelets, the finite-size effects will play an important role. The finite-size effects include surface and curvature terms. In the present version of the quasiparticle model, the bag function is related to the finite size and describes the nonperturbative interaction between quarks. Our previous work [9] demonstrates that the surface term is more important than the curvature term when talking about their contribution to the energy. The surface modification is the main purpose of this paper.

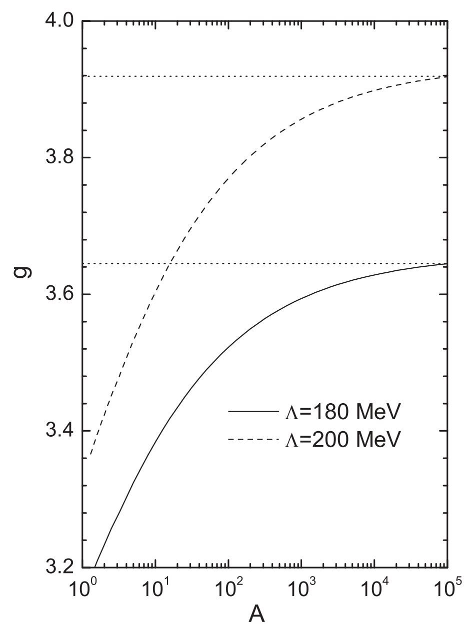

FIG. 4. The coupling constant g as a function of the baryon number for stable strangelets. The two horizontal dotted lines are the corresponding values for bulk strange quark matter.

In the literature, the intrinsic surface tension is characteristic of the phase boundary between the true vacuum and the perturbative phase and it can be neglected in comparison to the confinement bag constant. So only the calculation of the dynamical surface tension is carried out in our work. The early work is typically traced back to the paper [11] by Berger *et al.* They suggested that the surface tension parameter σ, which arises in finite strangelets, can be calculated from the surface modification of the fermion density of states. Strangelets with larger values of the surface tension could survive the early Universe [11]. In contrast, there is a suggestion that the surface tension for the interface separating the quark and the hadron phase should be smaller to make the mixed phase occur [6,25]. Recently, a critical value (70 MeV/fm^2 in Ref. [26]) has been suggested, above which the structure of the mixed phase will become unstable. Lattice gauge simulations suggest a range of $\sigma \approx$ 10–100 MeV/fm^2 at finite temperature [10,27]. However, the value of the surface tension is poorly known and typically values used in literature range within 10–50 MeV fm^{-2} [28–30].

In this paper, a chemical-potential- and radius-dependent bag function and running coupling constant are applied. So the surface term will provide a modification to the surface tension. We can use the above formulas in the preceding section to investigate the finite-size effects. In 1987, Berger and Jaffe made an original definition that the surface tension equals the free energy per unit surface area [11]. At a later time, the surface term of the density of states was developed in Ref. [31]. Mardor and Svetitsky corrected an error of a factor of 2 in the coefficient of the area term and gave a suitable formula in the MIT bag model [32]. Their density of states is defined as

$$\frac{dN_i}{dp} = \frac{p^2 V}{2\pi^2} - \frac{pS}{8\pi}\left[1 - \frac{2}{\pi}\arctan\left(\frac{p}{m_i}\right)\right]. \tag{21}$$

The corresponding surface tension within the bag model is

$$\sigma_i^{\text{bag}} = d_i T\int_0^\infty \frac{p}{4\pi^2}\arctan\left(\frac{m_i}{p}\right)\ln\left(1+e^{-(\sqrt{p^2+m_i^2}-\mu_i)/T}\right)dp. \tag{22}$$

In fact, the surface part in Eq. (21) is consistent with Eq.(12). In this paper, we adopt the density of states derived by the multiexpansion approach in Sec. II. The total energy density per baryon is the sum of three parts,

$$E = \sum_i (E_{V,i} + E_{S,i} + E_{C,i}). \tag{23}$$

In particular, in the quasiparticle model, the effective bag constant at finite size will have a modification needed to calculate the volume energy density $E_{V,i}$, which is associated with the particle number n_i and chemical potential as

$$\sum_i E_{V,i} = \sum_i (\Omega_{V,i} + B_{V,i} + \mu_i n_i) + B_0. \tag{24}$$

The surface energy density $E_{S,i}$, and curvature energy density $E_{C,i}$ due to the i-type quark are

$$E_{S,i} = \Omega_{S,i} + B_{S,i}, \tag{25}$$

$$E_{C,i} = \Omega_{C,i} + B_{C,i}. \tag{26}$$

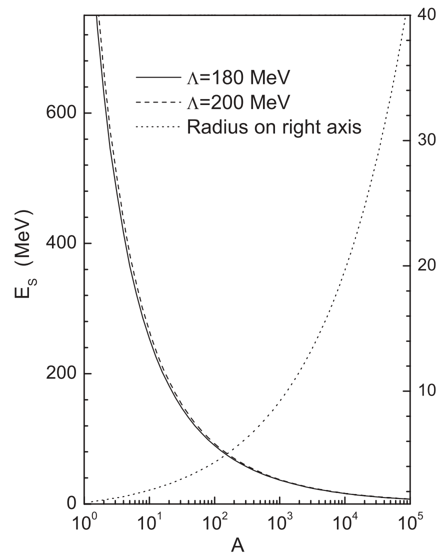

FIG. 5. The surface energies of strangelets are showed on the left axis by solid and dashed lines respectively for Λ =180, and 200 MeV. The corresponding radius (dotted line) are on the right axis.

Consequently, the surface tension σ_i, the free energy per unit surface area, should be redefined as [33]

$$\sigma_i \equiv \frac{E_{S,i}V}{4\pi R^2} = \frac{R}{3}(\Omega_{S,i} + B_{S,i}). \tag{27}$$

The first term on the right side of Eq. (27) is the kinetic contribution to the surface tension at zero temperature. It has the same expression as the previous definition σ_i^{bag} of Eq. (22) in the conventional bag model [11]. The important difference is that the constant quark mass in the bag model is replaced by the effective mass m_i^* in Eq. (1), $\sigma_i^{\text{bag}}(m_i \to m_i^*) = \frac{R}{3}\Omega_{S,i}$. The second additional term is dominated by $B_{S,i}$ in Eq. (19). Generally, it will give a positive modification to σ_i^{bag} depending on the value of the running coupling constant $g(\mu_i)$. Combining the conventional term $\sigma_i^{\text{bag}}(m_i^*)$ with the additional term $B_{S,i}$ multiplied by the factor $R/3$, the surface tension is only associated with the chemical potentials and independent of R. The total surface tension σ_i is a sum over the three flavors, $\sigma = \sum_i \sigma_i$.

For a given baryon number, the strangelets can be calculated by solving the chemical potentials μ_u and μ_s. In Fig. 5, the surface energy per baryon is shown on the left axis and the radius is plotted by a dotted line on the right axis. The choice of the Λ value has little influence on the numerical results. According to Eq. (27), the total surface tension of strangelets is shown for Λ = 180 MeV (solid line) and 200 MeV (dashed line) in Fig. 6. When the baryon number of strangelets increases to $A > 10^3$, the surface tension has a slight decline and comes gradually down to the vicinity of 30 MeV fm^{-2}. The comparison with the bag model result is marked by the dotted line in the figure, and the difference is obvious especially for smaller strangelets. It is physically reasonable that the surface tension decreases with increasing baryon number. With a larger Λ, the strangelets can have a larger surface tension. Therefore, we can predict to some extent that a larger

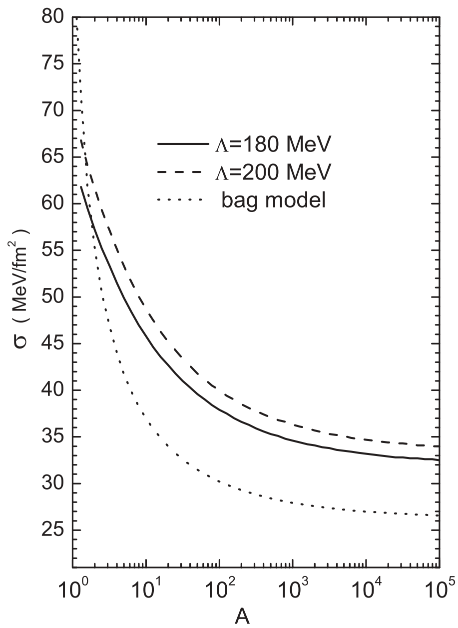

FIG. 6. Surface tension σ decreases with the increasing baryon number A.

Λ together with a stronger coupling strength can save the SQM in the early Universe. These properties can also be applied when studying the propagation of strangelets as cosmic rays [34].

IV. SUMMARY

In this paper we have studied the finite-size effects on the properties of strangelets in the framework of the extended quasiparticle model at zero temperature. In this model, the effective quark mass, characteristic of the variational mass system, is associated with the environment, that is, the chemical potential. In the calculations, we take into account the density-dependent running coupling constant $g(\mu)$ derived from the effective quark propagator by resumming one-loop self-energy diagrams in the hard dense loop approximation. So we can give a new calculation formula for the running coupling instead of assuming a constant value as in previous work. It is found by numerical calculations that the coupling constant and the quark effective masses are sensitively dependent on the effects of the medium. They decrease with increasing chemical potential in a proper range. This result is consistent with the spirit of QCD. From the thermodynamic consistency relation, a chemical-potential- and radius-dependent bag function can be derived and can be divided into three parts, the volume term $B_{V,i}$, the surface term $B_{S,i}$, and the curvature term $B_{C,i}$. They play an important role in finite-size strange quark matter. The surface term will have a positive influence on the energy of the strangelets depending on the interaction constant $g(\mu_i)$. Subsequently, the contribution of the surface term $B_{S,i}$ to the redefined surface tension is described by an additional term of $B_{S,i}R/3$. Ultimately, the total surface tension is dominated by the interaction and independent of the radius.

With the new definition, we find the surface tension equals 35 MeV fm^{-2} for strangelets of $A > 10^3$. The relations of the coupling constant, the radius, the surface energy per baryon, and the surface tension with the variational baryon number are investigated. For stable strangelets, we find that the surface tension will decrease with increasing baryon number but become larger for a big Λ value. Comparisons with the MIT bag model indicate an obvious difference for smaller strangelets. One can get a larger surface tension with a larger Λ value. Therefore, it can be predicted that the larger Λ together with a stronger coupling strength can save the SQM in the early Universe.

The work is carried out for the sake of simplicity. It is desirable to perform a more realistic calculation including electrons and quark charge distributions at the surface [35]. The influences on the quark-hadron mixed phase structure are expected to be calculated in future work.

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