Asymmetry of the strange sea in nucleons

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(Received 8 June 2001; published 28 November 2001)

Based on finite-temperature field theory, we evaluate medium effects in the nucleon, which can induce an asymmetry between quarks and antiquarks of the strange sea. The short-distance effects determined by the weak interaction can give rise to $\delta m \equiv \Delta m_s - \Delta m_{\bar{s}}$ where $\Delta m_{s(\bar{s})}$ is the medium-induced mass of the strange quark of a few keV at most, but the long-distance effects caused by strong interaction are sizable. Our numerical results show that there exists an obvious mass difference between strange and antistrange quarks, as large as 10–100 MeV.

DOI: 10.1103/PhysRevD.65.014003 PACS number(s): 12.40.−y, 11.10.Wx, 11.30.Na, 14.20.Dh

I. INTRODUCTION AND MOTIVATION

The existence of the Dirac sea is always an interesting topic that theoretists and experimentalists of high energy physics are intensively pursuing, and the strange content of the nucleon sea is of particular interest. Ji and Tang [1] suggested that, if a small locality of the strange sea in the nucleon is confirmed, some phenomenological consequences can result. The CCFR data [2] indicate that $s(x)/\bar{s}(x) \sim (1-x)^{-0.46\pm 0.87}$. Assuming an asymmetry between s and $\bar{s}$, Ji and Tang analyzed the CCFR data and concluded that $m_s = 260\pm 70$ MeV and $m_{\bar{s}} = 220\pm 70$ MeV [1]. So if we consider only the central values, $\delta m \equiv m_s - m_{\bar{s}} \sim 40$ MeV, which implies quark-antiquark asymmetry. However, one may alternatively conclude that the data are consistent with no asymmetry within error bars [1,2].

In the framework of the standard model $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$, we would like to look for some possible mechanisms that can induce asymmetry between quarks and antiquarks.

The self-energy of the strange quark and antiquark $\Sigma_{s(\bar{s})} = \Delta m_{s(\bar{s})}$ occurs via loops where various interactions contribute to $\Sigma_{s(\bar{s})}$ through the effective vertices. Obviously, the QCD interaction cannot distinguish between s and $\bar{s}$; in fact, neither can the weak interaction alone. Practical calculations of the self-energy also show that $\Delta m_s = \Delta m_{\bar{s}}$. In fact, because of the CPT theorem, s and $\bar{s}$ must be of exactly the same mass. So we ask ourselves what can cause asymmetry between s and $\bar{s}$, which are assumed to be the sea quark and antiquark in nucleons. We find immediately that in nucleons there are asymmetric quarks u and d, namely, the composition of u-$\bar{u}$ and d-$\bar{d}$ quark-antiquarks is asymmetric. Here u-$\bar{u}$, and d-$\bar{d}$ include the valence and sea portions of the corresponding flavors. In protons, there are two valence u quarks, but one valence d quark, while in neutrons, there is one u, but two d's. This asymmetry, as we show below, can represent a medium effect that results in an asymmetry $\Delta m_s \neq \Delta m_{\bar{s}}$ for strange sea quarks.

As discussed above, if we evaluate the self-energies Δm_s and $\Delta m_{\bar{s}}$ in vacuum, the CPT theorem demands $\Delta m_s \equiv \Delta m_{\bar{s}}$. However, when we evaluate them in an asymmetric environment of nucleons, an asymmetry $\Delta^M m_s \neq \Delta^M m_{\bar{s}}$ (where the superscript M denotes the medium effects) can be expected. In other words, we suggest that the asymmetry of the u- and d-quark composition in nucleons leads to an asymmetry of the strange sea.

There exist both short-distance and long-distance medium effects. The short-distance effects occur at quark–gauge boson level, namely, a self-energy loop including a quark-fermion line and a W-boson line or a tadpole loop (see below for details). The contributions of u and d types of quark-antiquarks to the asymmetry are realized through Kobayashi-Maskawa-Cabibbo mixing. Because of the small parton mass, the Higgs contributions can be neglected. In contrast the long-distance effects are caused by loops that include a quark-fermion line and a meson (kaon in our case) line.

In fact, Brodsky and Ma proposed a meson-baryon resonance mechanism and suggested that the sea quark-antiquark asymmetries are generated by a light-cone model of energetically favored meson-baryon fluctuations [3]. It was first observed by Signal and Thomas [4] that the meson-cloud model of nucleons can introduce a mechanism for strange-antistrange asymmetry in the nucleon sea, although their formalism is not consistent in treating the antiquark distributions in a strict sense [5]. An s-$\bar{s}$ asymmetry was also predicted by Burkardt and Warr [6] from the chiral Gross-Neveu model at large N_c in the LC formalism. Our physics picture is similar and the method is parallel to theirs, while all the calculations are done based on the finite-temperature field theory.

II. FORMULATION

We are going to employ the familiar formulation of quantum field theory at finite temperature and density. As is well

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known, the thermal propagator of quarks can be written as [7]

$$iS_q(k)=\frac{i(\not{k}+m_q)}{k^2-m_q^2}-2\pi(\not{k}+m_q)\,\delta(k^2-m_q^2)f_F(k\cdot u), \tag{1}$$

where u_μ is the four-vector for the medium, f_F denotes the Fermi-Dirac distribution function

$$f_F(x)=\frac{\theta(x)}{e^{\beta(x-\mu)}+1}+\frac{\theta(-x)}{e^{-\beta(x-\mu)}+1}, \tag{2}$$

$\beta=1/kT$, and μ is the chemical potential. We notice that the first term of Eq. (1) is just the quark propagator in the vacuum. Its contribution to Σ_1 is of no importance to us because this is related to the wave-function renormalization of the quark in the vacuum. We need to focus on the medium effect, which comes from the second term of Eq. (1).

It is experimentally confirmed that the total light quark number in the nucleon is 3. If we omit the small mass differences of the light quarks (u and d types, explicitly), the quark density in the nucleon is

$$\begin{aligned} n_q-n_{\bar{q}}&=\int\frac{d^3k}{(2\pi)^3}[f_F(\omega_k)-f_F(-\omega_k)]\\ &=\frac{3}{V_{eff}}\qquad (q=u,d). \end{aligned} \tag{3}$$

In this expression $\omega_k=\sqrt{\mathbf{k}^2+m_q^2}$ is the energy of the light quark and V_{eff} is the effective nucleon volume where $q(\bar{q})$ resides. Here we have ignored the possible sea quark asymmetry for light quarks of u and d flavors [8]. For up and down flavors, we have

$$n_u-n_{\bar{u}}=\frac{2}{V_{eff}},\qquad n_d-n_{\bar{d}}=\frac{1}{V_{eff}} \tag{4}$$

in protons and

$$n_u-n_{\bar{u}}=\frac{1}{V_{eff}},\qquad n_d-n_{\bar{d}}=\frac{2}{V_{eff}} \tag{5}$$

in neutrons.

A. Short-distance contribution

The corresponding Feynman diagrams are shown in Figs. 1(a) and 1(b).

The two contributions to the self-energy of the s quark $(\bar{s})$ (a) and (b) are due to the charged current $(W^{\pm})$ and neutral current (including the Z and γ), respectively; the latter is usually called the tadpole diagram [9]. The total contribution is analogous to that given in [9]; the only difference is that for neutrinos only the weak neutral current plays a role, while for quarks the electromagnetic interaction also needs to be involved.

The contribution due to the charged current is

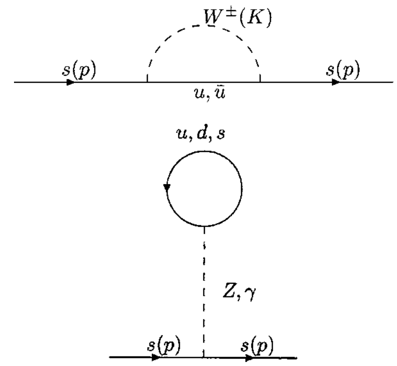

FIG. 1. (a) The self-energy Σ_1^s for a strange quark where the exchanged bosons are either W bosons or kaons corresponding to short- and long-distance effects, respectively. (b) The tadpole diagram that contributes to the self-energy Σ_2^s of the strange quark.

$$\Sigma_1^s=\sqrt{2}G_F\gamma^0 L\sin^2\theta_C(n_u-n_{\bar{u}}), \tag{6}$$

where G_F is the Fermi coupling constant and θ_C is the Cabibbo angle. The contribution due to the weak neutral current is

$$\begin{aligned}\Sigma_2^s=&3\sqrt{2}G_F\left(-1+\frac{4}{3}Q^{(s)}\sin^2\theta_W\right)\cdot\sum_f\,(T_3^{(f)}-2Q^{(f)}\sin^2\theta_W)\\ &\times(n_f-n_{\bar{f}}),\end{aligned} \tag{7}$$

where $Q^{(f)}$ refers to the charge of the corresponding quark (u,d,s). Pal and Pham pointed out that the axial part of the neutral current does not contribute [9].

For the electromagnetic current the situation is different. The exchanged photon connecting the s-quark line (or $\bar{s}$ quark) and the closed loop (tadpole) possesses zero energy-momentum and its propagator

$$\frac{1}{q^2+i\epsilon}$$

results in an infrared divergence. In regular field theory, this does not cause problem because the integration over the inner momentum of the loop is exactly zero. In the case of the gluon infrared divergence, because of the non-Abelian Yang-Mills properties, there can be a tachyonic gluon mass [10] which can serve as an infrared cutoff. In the case of the photon as the gauge boson of the $\mathrm{U}_{em}(1)$ group there is no effective mass. This also means that there is no connection between the closed fermion loop and the s-quark line; thus the two parts are actually disconnected and such an electromagnetic tadpole should not be included in the phenomenological calculations even though it has a superficial infrared divergence. Thus we drop the electromagnetic tadpole from our later calculations.

As for the strange antiquark, one can obtain the corresponding self-energy by changing the direction of propagation. Obviously, if we ignore the small mass difference of u and d quarks, we have

$$\Sigma_1^{\bar{s}}=-\Sigma_1^s\qquad\text{and}\qquad\Sigma_2^{\bar{s}}=-\Sigma_2^s,$$

and the mass difference between the strange and antistrange quarks is

$$\delta m \equiv 2(\Sigma_1^s + \Sigma_2^s) \tag{8}$$

due to the short-distance interactions.

B. Long-distance effects

Up to now the dynamical picture of the sea quark interaction in QCD is not definitively understood. In our case of interest, the main contribution of the strong interaction will be attributed to the low-energy effective coupling between the internal Goldstone bosons and quarks. According to the common knowledge of the low-energy effective strong interaction, the strange quark is generated from dissociation of nucleons into hyperons plus kaons. In this picture the $s(\bar{s})$ quark in the sea interacts with the light quark and the kaon meson essentially. This is analogous to the interaction considered in Refs. [3,4,6] and by Kogan and Shifman who introduced such effects for weak radiative decay [11], but we estimate these effects with a quark-meson interaction instead of a baryon-meson interaction [12].

In the calculations, we need an effective vertex for $\bar{s}qM$ where q can be either u or d quark and M is a pseudoscalar or vector meson. Here we retain only the lowest-lying meson states such as π, K, ρ, etc. The effective chiral Lagrangian for the interaction between quarks and mesons has been derived by many authors [13,14]. For completeness, we present the well-established form of the chiral Lagrangian as [14]

$$L_x = i\bar{q}(\partial\!\!\!/ + \Gamma\!\!\!/ + g_A \Delta\!\!\!/ \gamma_5 - iV\!\!\!/)q - m\bar{q}q + \text{meson part}. \tag{9}$$

It is noted that, in general, the chiral Lagrangian applies only to the interactions between constituent quarks and mesons, on the other hand, the sea quark picture is valid for current quarks (or partons) [15]. Here we borrow the chiral Lagrangian picture just because the fundamental forms of interactions of either partons or constituent quarks with mesons are universal. The key point is the coupling constants; they may be different in the two pictures. However, we can assume that they do not deviate by orders of magnitude from each other. Therefore we can use the coupling constants for constituent quarks in our estimation of the order of magnitude. The results obtained and new data might offer us an opportunity to determine the effective coupling between parton quarks and mesons. Anyway, we point out that the numerical results obtained in this work may have a relatively large error of about 10 MeV, as we will see in the section on numerical calculations.

In this expression we omit the irrelevant part which contains only mesons. Here $\bar{q} = (\bar{q}_u, \bar{q}_d, \bar{q}_s)$ and

$$V_\mu(x) = \lambda \cdot \mathbf{V}_\mu = \sqrt{2}\begin{pmatrix} \frac{\rho_\mu^0}{\sqrt{2}} + \frac{\omega_\mu}{\sqrt{2}} & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\frac{\rho_\mu^0}{\sqrt{2}} + \frac{\omega_\mu}{\sqrt{2}} & K_\mu^{*0} \\ K_\mu^{*-} & \bar{K}_\mu^{*0} & \phi_\mu \end{pmatrix}. \tag{10}$$

Δ_μ and Γ_μ are defined as

$$\Delta_\mu = \frac{1}{2}[\xi^\dagger(\partial_\mu - ir_\mu)\xi - \xi(\partial_\mu - il_\mu)\xi^\dagger],$$

$$\Gamma_\mu = \frac{1}{2}[\xi^\dagger(\partial_\mu - ir_\mu)\xi + \xi(\partial_\mu - il_\mu)\xi^\dagger], \tag{11}$$

with

$$\xi = \exp[i\lambda^a \Phi^a(x)/2f],$$

where the Goldstone bosons Φ^a are the pseudoscalar mesons in the SU(3) octet and f is the decay constant.

From the chiral effective Lagrangian, the basic effective vertex is a pure-derivative axial vector (chiral symmetric) coupling as $f_{kqs}\bar{\psi}\gamma_5\gamma_\mu\partial_\mu K\psi$. In the last part of this work, we will give more discussion on this issue. The quantum correction at one-loop level provides a self-energy Σ_3^s for the strange quark shown in Fig. 1(a) (where we only need to replace the W-boson line by a kaon line). Here we neglect the higher-loop contribution which may be induced by the higher-order graphs in the chiral Lagrangian. In general, for such an estimation, the effective coupling by itself may contain some higher-loop contributions; thus their effects do not influence the qualitative conclusions although the quantitative results may change [14]. So we can write the amplitude due to the long-distance effective interaction as

$$-i\Sigma_3^s = if_{kqs}^2 \int \frac{d^4k}{(2\pi)^4}\gamma_5\gamma_\mu iS_q(k)\gamma_5\gamma_\mu \frac{(p-k)^2}{(p-k)^2 - M_K^2}, \tag{12}$$

where M_K is the mass of kaons.

In the rest frame of the medium $[u_\mu = (1,\mathbf{0})]$ we use

$$\delta(k^2 - m_q^2) = \frac{1}{2\omega_k}[\delta(k_0 - \omega_k) + \delta(k_0 + \omega_k)], \tag{13}$$

where $\omega_k = \sqrt{\mathbf{k}^2 + \mathbf{m}_q^2}$ is the energy of the light quark. So the long-distance medium correction to the mass of the strange quark can be evaluated, and we obtain

$$\Sigma_3^s = \gamma_0 \frac{f_{kqs}^2}{2}\int \frac{d^3\mathbf{k}}{(2\pi)^3}\left[\frac{m_s^2 - 2m_s\omega_k}{m_s^2 - 2m_s\omega_k - M_K^2} f_F(\omega_k) - \frac{m_s^2 + 2m_s\omega_k}{m_s^2 + 2m_s\omega_k - M_K^2} f_F(-\omega_k)\right]. \tag{14}$$

After simple manipulations, this becomes

$$\Sigma_3^s = \gamma_0 \frac{f_{kqs}^2}{2}\left[(n_q - n_{\bar{q}}) + \int \frac{d^3\mathbf{k}}{(2\pi)^3}\frac{M_K^2}{m_s^2 - 2m_s\omega_k - M_K^2} f_F(\omega_k) - \int \frac{d^3\mathbf{k}}{(2\pi)^3}\frac{M_K^2}{m_s^2 + 2m_s\omega_k - M_K^2} f_F(-\omega_k)\right]. \tag{15}$$

In order to avoid the pole in the second term of Eq. (14), we use the familiar Breit-Wigner formulation. Thus we give Δm_s contributed by the long-distance effects as

$$\begin{aligned}\Delta m_s &= \frac{f_{kqs}^2}{2}\int \frac{d^3\mathbf{k}}{(2\pi)^3}\left[\frac{(m_s^2 - 2m_s\omega_k - M_K^2)(m_s^2 - 2m_s\omega_k)}{(m_s^2 - 2m_s\omega_k - M_K^2)^2 + (M_K\Gamma_K)^2} f_F(\omega_k) - \frac{(m_s^2 + 2m_s\omega_k - M_K^2)(m_s^2 + 2m_s\omega_k)}{(m_s^2 + 2m_s\omega_k - M_K^2)^2 + (M_K\Gamma_k)^2} f_F(-\omega_k)\right] \\ &= \frac{f_{kqs}^2}{2}\int \frac{k^2 dk}{2\pi^2}\left[\frac{(m_s^2 - 2m_s\omega_k - M_K^2)(m_s^2 - 2m_s\omega_k)}{(m_s^2 - 2m_s\omega_k - M_K^2)^2 + (M_K\Gamma_k)^2} f_F(\omega_k) - \frac{(m_s^2 + 2m_s\omega_k - M_K^2)(m_s^2 + 2m_s\omega_k)}{(m_s^2 + 2m_s\omega_k - M_K^2)^2 + (M_K\Gamma_k)^2} f_F(-\omega_k)\right],\end{aligned} \tag{16}$$

where in the Breit-Wigner expression we take the usual approximation that Γ_K is the measured value of the lifetime of $K^{\pm}$.

The statistical integral Eq. (16), as a function of temperature T and chemical potential μ, can be expressed in terms of the quark density and therefore the nucleon size [cf. Eq. (3)] at a given temperature. In the practical calculation, we use a numerical integral method to relate the mass correction Δm_s to the effective nucleon radius R.

III. NUMERICAL RESULTS

A. Short-distance effects

Our numerical results show that

$$\delta m = 92 \ \text{eV} \sim 0.8 \ \text{keV} \quad \text{for the proton},$$

$$\delta m = 0.38 \ \text{keV} \sim 3.0 \ \text{keV} \quad \text{for the neutron},$$

in the range of the effective nucleon radius $R \approx 0.5$–1.0 fm. So we see that the short-distance interaction cannot result in a large asymmetry in the strange sea.

B. Long-distance effects

According to the picture of chiral field theory [16–19], the effective pseudovector coupling implies $f_{kqs} = g_A/\sqrt{2}f$, where the axial-vector coupling $g_A = 0.75$. The pion decay constant $f_\pi = 93$ MeV and the kaon decay constant $f_K = 130$ MeV; for our estimation, an approximate SU(3) symmetry might be valid, so that f can be taken as an average of f_π and f_K. One can trust that the order of magnitude of the effective coupling at the vertices does not deviate too much from this value.

In the chiral quark model from the framework of the standard chiral field theory [13], one usually takes the quarks as the constituent quarks [16,20]. However, there is also a suggestion [21] to consider the quarks in the chiral dynamics as current quarks from the successful description of the proton spin data [17,18], consistent with our decision to use the chiral Lagrangian picture for an effective description of the current quark and meson interaction. Therefore we present our calculated results of Δm_s for two cases: (a) with the quark masses being the current quark masses $m_s = 150$ MeV and $m_u \approx m_d = 6$ MeV; (b) with the quark masses being the constituent quark masses $m_s = 500$ MeV and $m_u = m_d = 350$ MeV.

We present the numerical results for Δm_s as a function of the effective nucleon radius in Fig. 2. We find that the result depends sensitively on the value of the nucleon volume for small R values, which correspond to high density, but the dependence becomes mild as R becomes larger.

In Fig. 2, one can see that when R is about 0.5 fm Δm_s can be as large as 50 MeV, which results in $\delta m \equiv 2\Delta m_s \simeq 100$ MeV for the current quark mass case and even larger for the constituent quark mass case. But for the situation corresponding to a normal nucleon case, δm might be only of the order of around a few tens of MeV, which is within the uncertainties as estimated in Ref. [1].

We note that at high temperature with ordinary density of nuclear matter, or at high density with ordinary temperature, the baryons and mesons may undergo a phase transition to the quark-gluon plasma with quarks and gluons as the basic degrees of freedom, and spontaneous chiral symmetry may be restored. Thus the effective s-kaon-u coupling f_{kqs} and the pion or kaon decay constant f may decrease with increasing temperature and density, and the quark-antiquark asymmetry may also decrease with increasing temperature and density. However, here we consider the nucleon case, with T

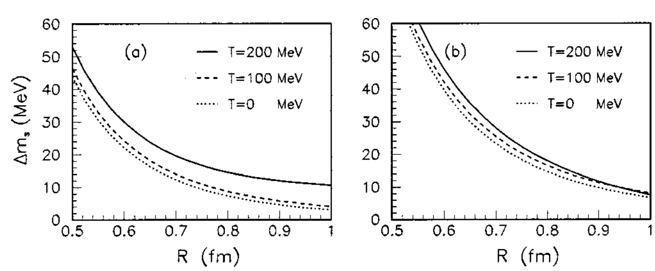

FIG. 2. The medium correction Δm_s for the strange quark mass vs the effective nucleon radius R, in two cases with different quark masses: (a) current quark masses $m_s = 150$ MeV and $m_u \approx m_d = 6$ MeV; (b) constituent quark masses $m_s = 500$ MeV and $m_u = m_d = 350$ MeV.

the temperature of normal nuclear matter in the chiral dynamics, in a range of about 100 MeV to a few hundreds of MeV, which is lower than chiral symmetry restoration scale of about 1 GeV, and also with the density not high for a nucleon in ordinary nuclear matter. Therefore we neglect the dependence of the coupling and decay constants on temperature and density in the present study. But it should be clear that the trend shown in Fig. 2, where the strange quark-antiquark mass difference δm_s increases with increasing temperature and density, will be changed if chiral symmetry restoration is considered. Thus the quark-antiquark asymmetry caused by the long-distance effects should not be much larger than what we estimated above, if such asymmetry can be large.

IV. DISCUSSION AND SUMMARY

The strange asymmetry discussed above in fact may be caused by an asymmetry of the light quark content in nucleons, but not by the interaction itself. We show that $u(\bar{u})$ and $d(\bar{d})$ quarks in nucleons can play the role of a medium, which results in an asymmetry of the self-energy of s and $\bar{s}$. The resultant values for s and $\bar{s}$ have opposite signs, which induce a net mass difference between s and $\bar{s}$. The nucleon structure determines $\delta m = m_s - m_{\bar{s}} > 0$.

There exist both short-distance and long-distance effects which are mediated by $W(Z)$ bosons and K mesons, respectively. Because the gauge bosons W and Z are very heavy, the net effects are much suppressed and their contribution to δm can only be a few keV. By contrast, the kaon is much lighter and, moreover, the effective interaction $\bar{q}qP$ where P denotes a meson with appropriate quantum number is due to nonperturbative QCD, which is a strong interaction; thus the effective coupling is much larger than that for the weak interaction. This obvious enhancement gives a value of 10–100 MeV to δm, which is consistent with the estimate [1] required to fit the experimental data within error bars.

In our approach, we keep only the leading order in the effective chiral Lagrangian and leave the coupling constant as a parameter with SU(3) symmetry. If we completely apply the chiral effective Lagrangian, the coupling constant is fully determined. On the other hand the original formulation of the effective chiral Lagrangian is for the constituent quark, whereas here we take the parton picture instead; thus the coupling constant might deviate somewhat from that in the effective chiral Lagrangian. In this work we take the value of the coupling constant according to the Lagrangian but with a factor g_A, $f_{kqs} = g_A/\sqrt{2}f$. This factor g_A partly includes the nucleon structure effects and compensates the errors in applying the chiral effective Lagrangian to the parton picture. We expect that this choice does not much deviate from reality while using the Lagrangian for the parton picture. Certainly, such an approximation may cause errors, so that we cannot give precise predictions for the mass difference of $\bar{s}$ and s, but one can be convinced that the qualitative conclusion can be made and the order of magnitude is close to reality, because here we apply all the established theories and models except the numerical value of the coupling constant to make this evaluation.

In summary, we investigate the influence of the environment on the asymmetry of the strange sea based on the well-established finite-temperature field theory. In our opinion, a natural explanation for this asymmetry is the light quark sea in nucleons, which represents a medium with certain asymmetries. There exist always three "net" light quarks, which means an excess of q over $\bar{q}$ in nucleons. The important feature is the nonzero chemical potential for the light quark (antiquark) in this theoretical framework. As a consequence, we obtain an obvious asymmetry of the strange sea ($\delta m \sim 10-100$ MeV), although the magnitude of the asymmetry should also be constrained by chiral symmetry restoration. Considering that the net strange quark number is zero, this leads to $\mu_s \neq \mu_{\bar{s}} \neq 0$. We expect that this point will be helpful for understanding the non-negligible strangeness content of the nucleon. On the other hand, we suggest that the low-energy effective interaction provides the main dynamical origin of the strange sea asymmetry. In fact, many authors have emphasized that the light quark sea asymmetry can arise from effective interactions between the internal Goldstone bosons and quarks. So the essence of our approach is consistent with the treatments of [16–19] in the dynamical sense.

ACKNOWLEDGMENTS

This work was partially supported by the National Natural Science Foundation of China.

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