Abstract

Properties of dense hadronic matter including strange particles are studied within the relativistic mean-field theory (RMFT). The possibility of kaon condensation is reexamined, and a simple condition is found for the parameters included in RMFT.

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In relativistic heavy-ion collisions, it is expected that the hot and dense zone includes many species of hadrons, where a lot of strange particles should be produced. Recent studies have shown that kaon dispersion relation would be much changed in nuclear medium, which gives rise to kaon condensation in neutron star matter [1] on one hand and modification of kaon production in heavy-ion collisions [2], change of the dilepton production rate from RBUU [3] or the fire ball [4] on the other hand. For the former subject, the hyperon degrees of freedom has been extensively studied by way of the relativistic mean-field theory (RMFT) [5, 6]. In this letter we consider strange hadronic matter by extending RMFT to incorporate nucleon, Λ hyperon and K meson, which are essential degrees of freedom in strange hadronic matter. First preliminary report on this subject was given in ref. [26]. Here we examine the relevance of the parameters included in RMFT in more detail and present a relation between them to give a condition for kaon condensation in high-density matter.

Some years ago a possibility of strange hadronic matter produced in relativistic heavy-ion collisions was indicated in RMFT, where abundance of lambda or other hyperons overwhelms that of nucleons [7]. They, however, considered only baryons. On the other hand kaons, the lightest strange mesons, may also carry strangeness and they are much modified in the medium. So it needs to take kaons into account properly as well as hyperons to explore strange hadronic matter.

Nelson and Kaplan [8] have suggested \( K^+K^- (K^0\bar{K}^0) \)-pair condensation in relativistic heavy-ion collisions [1]. As increasing density the lowest energy of \( K^+(K^0) \) is reduced by the \( KN \) s-wave interaction (mainly by the KN-sigma term \( \Sigma_{KN} \) ), and eventually becomes equal to the strangeness chemical potential \( \mu_s (= \mu_K) \), where the \( K^+(K^0) \) effective mass reaches \(-\mu_K\) simultaneously and \( K^+K^- (K^0\bar{K}^0) \)-pair condensation occurs. Their idea is very interesting, but they considered only nucleon matter and did not take into account any hyperon degrees of freedom. Moreover, their result seems unlikely in light of the subsequent studies about kaon in medium; the \( K^+ \) excitation energy receives a repulsive effect instead of an attractive one [1].

\[\text{The possibility of kaon condensation in relativistic heavy-ion collisions has been also suggested for a different reason [1].}\]
In relativistic heavy-ion collisions hyperons and kaons are produced in pairs; this possibility is more favorable than that of the kaon-antikaon pair production due to the different threshold energy. In previous papers [10, 11], furthermore, it has been shown that the relativistic effects moderate kaon condensation by the *self-suppression* mechanism in neutron-star matter. This self-suppression effects should also appear in symmetric nuclear matter; the main effect caused by the $KN$-sigma term $\Sigma_{KN}$ is weakened by the small nucleon effective mass.

Here we study strange hadronic matter at high-density but rather low temperature ($T < m_\pi$) by dealing hyperon and kaon equally within RMFT. In particular we pay attention to the abundance of $K^+(K^0)$ which is measured by the strangeness chemical potential $\mu_s (=\mu_K)$, and discuss the possibility of $K^+(K^0)$ condensation. We, hereafter, restrict ourselves to hadron matter in thermal equilibrium, and treat it at zero temperature for simplicity because the nuclear properties such as the nucleon and meson effective masses are not largely changed below $T \sim m_\pi$ [12], and thermal-pion effects hardly modify our conclusion. We shall give some comments on this matter at the end of this letter, where we also briefly discuss the observability of kaon condensation in the realistic situation of heavy-ion collisions.

First we briefly explain our basic formulation. The lagrangian density with nucleon, lambda and kaon is given as follows:

$$\mathcal{L} = \bar{\psi}_N (i \not\partial - M_N) \psi_N + g_\sigma \bar{\psi}_N \psi_N \sigma + g_\omega \bar{\psi}_N \gamma_\mu \psi_N \not\omega^\mu$$
$$+ \bar{\psi}_\Lambda (i \not\partial - M_\Lambda) \psi_\Lambda + g_\sigma \bar{\psi}_\Lambda \psi_\Lambda \sigma + g_\omega \bar{\psi}_\Lambda \gamma_\mu \psi_\Lambda \not\omega^\mu$$
$$- \tilde{U}[\sigma] + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu$$
$$+ \{ (\partial_\mu - \frac{3i}{8f^2} \bar{\psi}_N \gamma_\mu \psi_N) K^\dagger \} \{ (\partial^\mu + \frac{3i}{8f^2} \bar{\psi}_N \gamma_\mu \psi_N) K \} - m_K^2 K^\dagger K$$
$$+ \frac{\Sigma_{KN}}{f^2} \bar{\psi}_N \psi_N K^\dagger K,$$  \hspace{1cm} (1)

where $\psi_N$, $\psi_\Lambda$ and $K$ are the nucleon, lambda and kaon fields, respectively, and $\tilde{U}[\sigma]$ is the self-energy potential of the scalar mean-field as

$$\tilde{U}[\sigma] = \frac{\frac{1}{2} m_s^2 \sigma^2 + \frac{1}{3} B_\sigma \sigma^3 + \frac{1}{4} C_\sigma \sigma^4}{1 + \frac{1}{2} A_\sigma \sigma^2}.$$  \hspace{1cm} (2)

3
In the above expression the interaction between nucleon and kaon is given by the \( KN \) sigma term \( \Sigma_{KN} \) and the Tomozawa-Weinberg type vector interaction. The latter interaction is introduced by the minimal coupling as in the usual chiral perturbation theory. Furthermore we omit the interaction between kaon and \( \Lambda \) because we treat only the system where the number of \( \Lambda \) hyperons is very small. In addition, the isospin-dependent terms are omitted because only the isospin saturation system is treated here.

From the Euler-Lagrange equations the nucleon and lambda effective masses are given as

\[
M^*_N = M_N - U_s(N) \\
M^*_\Lambda = M_\Lambda - U_s(\Lambda)
\]

\[
m^*_K = \sqrt{m^2_K - \frac{\Sigma_{KN}}{f^2} \rho_s(N)}, \tag{3}
\]

with the scalar potentials,

\[
U_s(N) = g_\sigma \langle \sigma \rangle + \frac{\Sigma_{KN}}{f^2 m^*_K} \rho_s(K), \\
U_s(\Lambda) = g_\Lambda \langle \sigma \rangle, \tag{4}
\]

where the kaon scalar density \( \rho_s(K) \) is equal to the kaon number density \( \rho_K \) at zero temperature, and other scalar densities \( \rho_s(\alpha) (\alpha = N, \Lambda) \) are defined by

\[
\rho_s(\alpha) = \frac{\gamma}{(2\pi)^3} \int d^3p f_\alpha(p) \frac{M^*_\alpha}{\sqrt{p^2 + M^2_\alpha}}, \tag{5}
\]

with the spin-isospin degeneracy factor, \( \gamma = 4 \) for nucleon and \( \gamma = 2 \) for lambda, and the Fermi momentum-distribution \( f_\alpha(p) \). The scalar mean-field \( \langle \sigma \rangle \) is given by

\[
\frac{\partial}{\partial \langle \sigma \rangle} \tilde{U}([\sigma]) = g_\sigma \rho_s(N) + g_\Lambda \rho_s(\Lambda). \tag{6}
\]

Then the baryon single-particle energy is written as

\[
\varepsilon_\alpha(p) = \sqrt{p^2 + M^2_\alpha} + U_0(\alpha), \tag{7}
\]
where time components of the vector potentials $U_0$ are given as

$$
U_0(N) = \frac{g_\omega^2}{m_\omega^2} \rho_N + \frac{g_\omega g_\Lambda^2}{m_\omega^2} \rho_\Lambda + \frac{9}{32 f^4 m_K^3} \rho_s(K) \rho_N
$$

$$
U_0(\Lambda) = \frac{g_\omega g_\Lambda^2}{m_\omega^2} \rho_N + \frac{g_\omega^2}{m_\omega^2} \rho_\Lambda.
$$

(8)

with the nucleon density $\rho_N$ and the lambda density $\rho_\Lambda$. Finally, the excitation energy of kaons $\omega_K(p)$ can be written as

$$
\omega_K(p) = \sqrt{p^2 + m_K^2} + U_0(K)
$$

(9)

with the Tomozawa-Weinberg vector potential, $U_0(K) = 3 \rho_N/8 f^2$ for $K^+(K^0)$. It is to be noted that this form is different from the one given in Refs. [1, 3, 4, 10, 11]. The difference stems from how to treat the vector interaction; this coupling scheme respects the gauge invariance of vector interactions [1].

If the nucleon Fermi energy $\varepsilon_F(N)$ becomes larger than $M_\Lambda^2 + U_0(\Lambda) + m_K^2 + U_0(K)$, the $\Lambda$ and kaon appear in matter, namely $K^+(K^0)$ condensation occurs. Since the situation strongly depends on the value of the $\Lambda$ coupling constants $g_\sigma^\Lambda$ and $g_\omega^\Lambda$, we first need to examine it in detail.

Naive $SU(3)$ symmetry relation suggests that the values of $\Lambda$ coupling constants are two-third of nucleon ones: $g_\sigma^\Lambda = 2 g_\sigma/3$ and $g_\omega^\Lambda = 2 g_\omega/3$. However, the coupling constants in RMFT should be different from the bare ones because they represent the effective strength of the mean-fields including many-body effects and higher-order correlations. The force between two baryons is strong but short-range, so that the mean-field cannot be perturbatively given, but its non-locality is rarely seen in low-energy phenomena: the $\sigma$ and $\omega$ fields are brought about by not only the Hartree contribution but also Fock and higher-order contributions. For the short-range interaction such as nuclear force the non-locality of the Fock part is very small, and the Hartree and Fock contributions cannot be distinguished in the spin-isospin saturated system at the low-energy limit. Furthermore we have not had enough information about the interaction between nucleon and lambda. At present, then, we cannot determine the effective coupling constant $g_\sigma^\Lambda$ and $g_\omega^\Lambda$ sufficiently.

Studies of hypernuclei have indicated that depth of the $\Lambda$ potential is about $28 - 30$ MeV, and that its $LS$-splitting is very small. Boguta and Bohrman [13]
have given $g_\Lambda/\sigma = g_\omega/\omega = 0.33$ by fitting the $\Lambda$ single-particle level under the restriction, $g_\Lambda^2/\sigma = g_\omega^2/\omega$. Rufa et al. \cite{4}, however, have shown that the parameters allow more ambiguities for the case without the restriction. \footnote{For other discussions about hypernuclei and RMFT, see Ref. \cite{14} and references cited therein.}

On the other hand Glendenning \cite{15} have pointed out that such small coupling constants too much soften EOS of neutron matter and cannot make sufficiently heavy neutron star with 1.5 times the solar mass. This point must be very important to study strange hadronic matter, but does not constrain the parameters so severely. As pointed out before, the $\sigma$ and $\omega$ fields do include not only the Hartree contribution but also the Fock and higher-order contributions. In the hyperon-rich system, the Fock field between hyperons should work. Within RMFT we should be able to stiffen EOS of strange hadronic matter by introducing the repulsive mean-field affecting only hyperons such as the $\phi$-field in place of the Hartree-Fock calculation. In this letter we do not introduce such fields because we treat only the hyperon-poor system.

Here we would like to suggest one more important hint for the $\Lambda$ mean-field from the optical potential analysis in proton-nucleus elastic scatterings \cite{16,17}. In order to reproduce the proton-nucleus optical potential observed experimentally, the strength of the vector mean-field must be inversely proportional to the incident energy \cite{16,17} around 1 GeV of proton incident energy. In general only the Hartree (local) parts contribute to the nucleon field at the high-energy limit because the Pauli blocking does not influence nucleon with momentum far from Fermi sea. From the optical potential analysis, thus, we have noticed that the Hartree contribution to the vector mean-field is very small \cite{16}. Since lambda is not affected by the Pauli effects, we can suppose that the mean-field for $\Lambda$ does not receive so strong many-body effects, and consequently the Hartree contribution should be dominant. Hence the $\Lambda$ potential $U_0(\Lambda)$ is estimated to be very small.

According to the above considerations we prepare two sets of parameters for the coupling of $\Lambda$ to the nucleon mean-fields. One set (L1) is that $g_\sigma^\Lambda = 2g_\sigma/3, g_\omega^\Lambda = 2g_\omega/3$ with the $SU(3)$ symmetry relation, and the other set (L2) is that $g_\omega^\Lambda = 0$ and

\begin{align*}
\frac{g_\Lambda}{g_\sigma} &= \frac{2g_\sigma}{3}, \quad \frac{g_\Lambda}{g_\omega} = \frac{2g_\omega}{3} \\
\end{align*}
obtained by the following condition.

\[ U_s(\Lambda) - U_0(\Lambda) = \frac{2}{3}(U_s(N) - U_0(N)). \]  

(10)

In addition we use two kinds of parameter-sets for nucleon, PM1 \((K = 200 \text{MeV},\ M_N^*/M_N = 0.7\ \text{at}\ \rho = \rho_0)\) and PM4 \((K = 200 \text{MeV},\ M_N^*/M_N = 0.55\ \text{at}\ \rho = \rho_0)\), and \(\Sigma_{KN} = 300 \text{MeV}\) for kaons. The depth of the \(\Lambda\) potential given by Eq. (10) is a little larger than the experimental value 30 MeV. However we have not known the value in the infinite-matter limit without the surface contribution, and the small difference does not affect a rather qualitative study in this letter.

Now we calculate the density-dependence of

\[ \Delta \mu = \varepsilon_F(N) - M_N^* - U_0(\Lambda) - m_K^* - U_0(K) \]  

(11)

in the no-strangeness system: \(\rho_\Lambda = \rho_K = 0\). At the critical density where \(\Delta \mu = 0\), a pair of \(\Lambda\) and \(K^+ (K^0)\) appears in the ground state, namely \(K^+ (K^0)\)-condensation occurs. Above the critical density chemical equilibrium among nucleon, lambda and kaon ever holds,

\[ \varepsilon_F(N) = \varepsilon_F(\Lambda) + \mu_K \]  

(12)

with the nucleon and lambda Fermi energies, \(\varepsilon_F(\alpha),\ \alpha = N, \Lambda,\) and the kaon chemical potential \(\mu_K = m_K^* + U_0(K)\), which is identical to the strangeness chemical potential \(\mu_s\). Fig. 1 shows the results with PM1-L1, PM1-L2, PM4-L1 and PM4-L2. We add that of free lambda and kaon with PM1 (PM1-free). In all cases kaon condensation occurs, while the critical density is different. In the parameter-sets with L1 the critical density is very high: \(\rho_c = 16 \rho_0\) for PM1-L1 and \(\rho_c = 7.6 \rho_0\) for PM4-L1. The former critical density is too high to be attained in relativistic heavy-ion collisions. In the parameter-sets with L2, however, the critical density is not very high: \(\rho_c = 6.2 \rho_0\) for PM1-L1 and \(\rho_c = 3.9 \rho_0\) for PM4-L1. We should furthermore note the difference between results of PM1-L2 and PM1-free; the critical density in PM1-L2 is larger than that in PM1-free though lambda and kaon feel an attractive field in PM1-L2. This difference is generated only by the Tomozawa-Weinberg term. In fact the KN-sigma term cannot much contribute to the final results because of the self-suppression mechanism [10, 11]
In Fig. 2 we show EOS of the normal phase and kaon condensed phase (a), the effective masses of nucleon and kaon as the ratio to their bare masses (b), and the \( \Lambda \) fraction \((\rho_\Lambda/\rho_B, \rho_B = \rho_N + \rho_\Lambda)\), (c) for PM1-L2 and PM4-L2. We can see that kaon condensation reduces the total energy per nucleon above the critical density. The effective mass decreases for nucleon, while increases for kaon; the latter feature stems from the reduction of the nucleon fraction. Although the reduction of EOS and the change of the effective masses are not so large, the \( \Lambda \) density increases rapidly, especially for PM4-L2. Since our model does not involve any field which works between hyperons and the kaon-lambda interaction, properties of the system above the threshold are not quantitatively realistic except near the critical density. So we should not take these results so seriously.

Since the scalar density \( \rho_s \) approaches to a finite value at the infinite density limit in RMFT, the vector mean-field dominantly affects the nuclear EOS in high-density regime. Hence the difference between the results of PM1 and PM4 comes from the different strength of the vector mean-fields. Namely kaon condensation is brought about by the balance among the different vector mean-fields of nucleon, lambda and kaon. Therefore the condition for \( K^+(K^0) \)-condensation is given as

\[
\frac{g_\omega^2}{m_\omega^2} - \frac{g_\omega^\Lambda g_\omega}{m_\omega^2} - \frac{3}{8}f^2 \leq 0.
\]  

(13)

As for the lambda potential the present parameterizations L1 and L2 should be the extreme ones; the realistic value must be laid between two cases. In order to clarify the ambiguity coming from the parameterization, in Fig. 3, we plot the relation between the critical density \( \rho_c \) and the ratio \( x_\omega = g_\omega^\Lambda / g_\omega \) under the condition (14). The critical density dose not strongly depends on the coupling \( g_\omega^\Lambda \) as far as \( g_\omega^\Lambda \leq 0.5g_\omega \), while it steeply varies around \( g_\omega^\Lambda = 0.6g_\omega \). In our view the value of \( g_\omega^\Lambda \) must be much smaller than L1 because of the small LS-splitting in lambda hypernuclei and very small Hartree contribution to the nucleon vector mean-field. Hence the realistic critical density should not be so far from that in L2.

In this letter we have not discussed the effects of temperature, while the compressed system in relativistic heavy-ion collisions must be hot.

\( ^3 \) We consider here the case \( T < m_\pi \), since we are interested in moderately high-energy heavy-ion
that the vector mean-field is only proportional to density in RMFT and independent of temperature. Furthermore it has been shown in Ref. [12] that the temperature dependence of the scalar mean-field is also small below $T \sim 200\text{MeV}$. Secondly, in a realistic situation, we should take into account another temperature effect as well, where many pions may be excited. One may wonder whether such thermal pions significantly modify our conclusion through the interaction with baryons or kaons. However, this is not the case at temperature below the pion mass ($T < m_\pi$). Actually the number density of pions ($\propto T^3$ in the chiral limit) is less than $0.2\text{fm}^{-3}$ there, which is much less than the baryon density we are interested in ($\rho \geq 4\rho_0$). Besides that, there is another reason why $\pi - N$ and $\pi - K$ interaction are not so important in our discussion. Experimental data shows that the $s$-wave interaction hardly contributes, at low momentum, in isospin-symmetric matter, which can be also explained by the chiral-symmetry argument. Hence the leading contribution comes from the $p$-wave interaction. Its contribution to the effective energies of baryons or kaons in the medium, which is essentially proportional to the square of pion momentum or that of kaon momentum, then enters through the modification of the dispersion relations (Eqs. (7),(9)). Its magnitude for the effective masses of kaons and baryons is proportional to $T^4$ in the chiral limit, and then much suppressed for $T < m_\pi$; e.g. $\delta M_N < \text{several tens MeV}$ by way of the scattering-volume approximation, which should be compared with the one from RMFT, $\delta M_N \sim \text{several hundreds MeV}$ even at $\rho_0$. From these discussions we may safely say that thermal effects play little role in our discussion and the effects of RMFT are dominant at temperature below $m_\pi$, that is, the density effects overwhelm the temperature effects in short.

In summary we have discussed the possibility of $K^+(K^0)$-condensation in high-density baryon matter and suggested that kaon condensation is caused by the difference of the vector mean-fields for nucleon, lambda and kaon; this difference becomes linearly larger with density. Thus the possibility of this condensation depends on the strength of their vector mean-fields. We have found a simple criterion for kaon condensation within RMFT (Eq.(13)). If parameters satisfy this condition, we can collisions around several tens GeV/u, e.g., the AGS energy region [2, 21].
expect kaon condensation at some density. The value of the critical density depends on the value of the nucleon effective mass $M_N^*$ at the saturation density. We have found $4 - 16 \rho_0$ for the critical density within our parameter sets, PM1 and PM4. The typical value of effective mass is empirically known as $M_N^*/M_N = 0.55 - 0.7 \rho_0$ [7, 13, 17, 18, 19], and our parameter-sets are consistent with them. Furthermore the small LS-splitting of lambda hypernuclei and the very tiny Hartree contribution to the nucleon vector mean-field suggests the small vector mean-field for lambda.

There remains a problem of the possibility of $K^+(K^0)$-condensation in high-energy heavy-ion collisions. There is controversy at present about how a system reaches thermal equilibrium and to what extent density and/or temperature are raised. Recent numerical simulations have shown that baryon density $\rho/\rho_0 = 7 - 10$ can be achieved by the high-energy heavy-ion collisions with several tens GeV/u like in the AGS energy region [20, 21]. In this case the system is expected to be quasi-equilibrium for the duration of 4 - 8 fm/c with the temperature $T \approx 120$MeV [20], which is still below the pion mass. Hence $K^+(K^0)$-condensation is very plausible to be generated in the high-density regime. In the heavy-ion collisions around hundreds GeV/u like in the SPS energy region it depends on the model whether the system is stopped and equilibrium is realized [22]; e.g., temperature becomes $T \approx 140$MeV in the RQMD simulation [23] which may be a marginal temperature for our discussion to be applied. $K^+(K^0)$-condensation, once occurs in course of relativistic heavy-ion collisions around the AGS energy region or higher, would give rise to a new phenomenon for dilepton production [26]. Of course we must wait for numerical simulations in order to conclude whether this phenomenon can be observed in the actual heavy-ion collisions. One promising way for this purpose is the RBUU approach [3, 24, 25], but we need to introduce the momentum-dependence for the mean-fields in this high-energy region [25]. It may be a very hard task in the light of the present computer power.

In this letter we have taken into account only the essential degrees of freedom for kaon condensation in symmetric nuclear matter. They would be sufficient to discuss the onset of the condensation [27], while we must consider other strange
particles besides Λ and kaon to get a realistic description of the condensed phase. Full argument including $SU(3)$ octet baryons and mesons will be given in a separate paper [27].

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Figure Captions

Fig. 1 Density-dependence of $\Delta \mu$. The dotted, solid, chain-dotted, dashed and thick dotted lines indicate the results in PM1-L1, PM1-L2, PM4-L1, PM4-L2 and PM1-free, respectively.

Fig. 2 Density-dependence of the total energy per nucleon (a), the ratio of the effective masses to the bare masses for nucleon and kaon (b), and the kaon fraction (c). The dashed and thick dotted lines indicate the results for PM1-L2 and PM4-L2, respectively. In the second column (b) the upper two lines show the effective mass of kaon and the lower ones show that of nucleon. The thick and thin lines in (a) denote EOS under kaon condensation and normal phases, respectively.

Fig. 3 Parameter-dependence of the critical density. The $x$-axis corresponds to the ratio $x_\omega = g_\omega^A / g_\omega$. The solid and dashed lines indicate the results for PM1 and PM4, respectively.
EQUATION OF STATE

(a) 

(b) 

(c) 

\[ \frac{E_T}{A} \text{ (MeV)} \]

\[ \frac{M^*}{M} \]

\[ \frac{\rho_K}{\rho_B} \]

\[ \frac{\rho}{\rho_0} \]
