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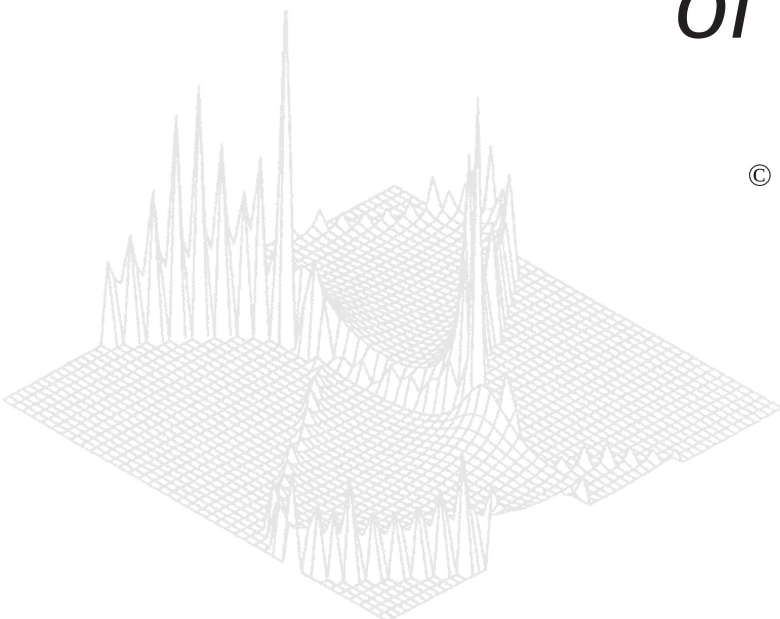

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Kaon Condensation and Dilepton Production from Strange Hadronic Matter*

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Abstract

We consider modification of kaons and the implications for dilepton production in the early stage of high-energy heavy-ion collisions. Constructing the equation of state of hadronic matter, including kaons as well as hyperons Λ with recourse to the relativistic mean-field theory, we study the production rate of dileptons. The possibility of K^+ condensation is also revisited in this framework.

1. Introduction

In the course of relativistic heavy-ion collisions with several GeV/u high density and/or high temperature, hadronic matter will be formed, where strangeness degrees of freedom as well as pions or other hadrons can be excited. Some years ago the possibility of strange hadronic matter was indicated within the relativistic mean-field (RMF) theory, where the abundance of lambda or other hyperons is very large and even overwhelms that of nucleons [1, 11].

On the other hand, kaon (K^+) condensation has been postulated from a simple consideration [2]; in heavy-ion collisions strangeness particles are created through the strangeness-conserving strong interactions, and especially kaons or lambda hyperons are rather easily created due to their small mass. Since the number of K^+ or K^- mesons becomes different due to their different creation mechanisms, there is a nonzero strangeness-chemical potential (μ_S) equal to that of the K^+ meson. When the lowest-excitation energy of kaons reaches the chemical potential μ_S, kaons begins to condense at that density. In this context, degrees of freedom of strange mesons besides hyperons become important.

The behaviour of kaons in matter is currently an interesting and important subject, triggered by K^- condensation in cold dense matter [3]. Many works have been carried out on the onset mechanism, the equation of state or its implications for neutron star physics, based on chiral symmetry [4, 5, 6]. Unfortunately a lack of precise information on the in-medium properties of kaons prevents us from a definite conclusion on the possibility of kaon condensation. In 1988 Nelson and Kaplan also suggested the possibility of K^+K^- pair condensation in heavy-ion

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collisions in a similar way [7]; the frequencies of K^+ and K^- excitations in nuclear matter (without any hyperons) *both* decrease as the density increases and K^+K^- pair condensation occurs around four times the nuclear density, where the strangeness chemical potential crosses the K^+ effective mass at the same time. However, their result seems unlikely now, to be true since subsequent calculations have given qualitatively different results, especially on the K^+ excitation energy: it should receive a repulsive effect instead of attractive [4, 5, 6].

We shall explore here strange hadronic matter at high density and/or temperature phase by explicitly treating kaons as well as hyperons. We use an extended RMF to include strange particles and take into account the matter effects on the kaon through the KN interactions based on chiral symmetry. We shall see that kaons are easily excited and an important constituent in matter in thermal equilibrium. Their chemical potential grows rapidly and seems to indicate K^+ condensation at relatively low density.

To extract information on such matter, especially kaon modification in dense and/or hot matter, it is most desirable to study dilepton production at an early stage of relativistic heavy-ion collisions [8].* Since the pion contribution is still dominant for the low invariant-mass region about 1 GeV, we could see the kaon contribution above this regime. In particular, the ϕ-meson peak is interesting because the threshold from the free K^+K^- is located just below the ϕ-meson mass by 30 MeV. Hence the peak should be sensitive to the modification of kaons. We also consider the broadening effect of the ϕ-meson peak due to the same modification of kaons [10]. If the K^+ meson forms a condensate, it would bring about new observable effects. We propose here a new peak for a signal of the condensation in the dilepton production, and discuss how the new peak grows up.

2. Strange Hadronic Matter

First we briefly describe our model Lagrangian within an extended RMF theory [1, 13]:

$$\begin{aligned}\mathcal{L} &= \overline{N}(i\partial\!\!\!/ - M_N)N + g_\sigma \bar{N}N\sigma + g_\omega \overline{N}\gamma_\mu N\omega^\mu \\ &+ \overline{\Lambda}(i\partial\!\!\!/ - M_\Lambda)\Lambda + g_\sigma^\Lambda \overline{\Lambda}\Lambda\sigma + g_\omega^\Lambda \overline{\Lambda}\gamma_\mu\Lambda\omega^\mu - \widetilde{U}[\sigma] + \tfrac{1}{2}m_\omega^2\omega_\mu\omega^\mu \\ &+ (D_\mu K)^\dagger D^\mu K - m_K^2 K^\dagger K + \frac{\Sigma_{KN}}{f^2}\overline{N}NK^\dagger K\,, \end{aligned} \tag{1}$$

where we have used a covariant derivative, $D^\mu = \partial^\mu + (i/2f^2)j_V^\mu$ with $j_V^\mu = 3/4\overline{N}\gamma^\mu N$, and $\widetilde{U}[\sigma]$ is the self-energy potential of the sigma-field given as [12]

$$\widetilde{U}[\sigma] = \frac{\frac{1}{2}m_s^2\sigma^2 + \frac{1}{3}B_\sigma\sigma^3 + \frac{1}{4}C_\sigma\sigma^4}{1 + \frac{1}{2}A_\sigma\sigma^2}\,. \tag{2}$$

* Gale and Kapusta suggested a remarkable change of the production rate from pion-pair annihilation if the dispersion relation of pions is much modified by the strong $\pi - N$ p-wave interaction [9].

This Lagrangian should be regarded as a minimal model including strangeness degrees of freedom, but still it retains essential features of the baryon-meson system. Indeed we have performed a full calculation with $SU(3)$ octet baryons and mesons and found that kaons and the Λ are surely the main constituent strange particles in matter. Since pions are, of course, the dominant constituent in hot and/or dense matter, we cannot discuss bulk properties of hadronic matter. Instead, we concentrate our attention here on some specific quantities which are irrelevant to pion contamination: the particle fraction of strange particles, strangeness chemical potential and excitation energy of kaons in a medium. We shall see later that these quantities provide the production rate of dileptons from matter.

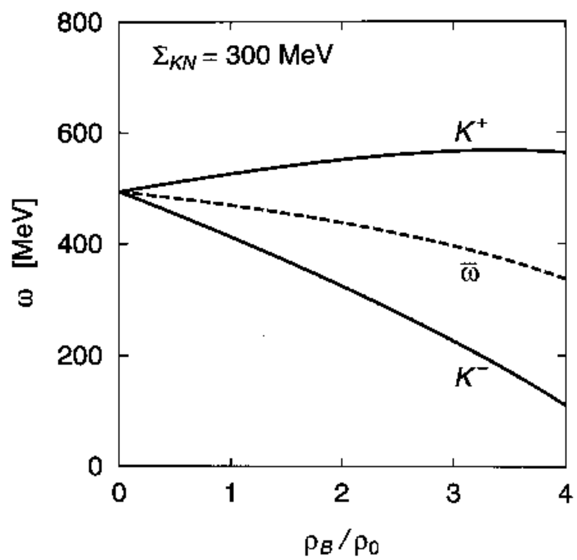

Fig. 1. Lowest excitation energies of kaons and their average $\bar{\omega}$.

We have incorporated the kaon–nucleon interaction through chiral theory: mainly the KN sigma term Σ_{KN} and the Tomozawa–Weinberg-type vector interaction [6, 12, 13]. In Fig. 1 we show, for example, the lowest excitation energy of kaons ($\omega_{k=0}$), where we can see the net effects of these terms [see Eq. (8) below].

The parameters in the mean-field Lagrangian are chosen so as to reproduce the bulk properties of nuclear matter and lambda hypernuclei. It is known that the parameters, except for the Λ-mean-field couplings, are rather well-fixed, while there is large ambiguity in the latter ones. The $SU(3)$ symmetry suggests that values of these coupling constants are two-thirds of the nucleon ones: $g_\sigma^\Lambda = 2g_\sigma/3$ and $g_\omega^\Lambda = 2g_\omega/3$. They should, however, be different from the bare ones, since they are effective ones including many-body effects and higher-order correlations. Recently Rufa *et al.* showed that the coupling constants cannot be fixed unambiguously from the single-particle spectra of lambda hypernuclei [1]. We shall mainly use the two parameter sets in Table 1, keeping in mind that there is the error ellipsis in the two-dimensional $g_\sigma^\Lambda - g_\omega^\Lambda$ plane.

Table 1. Parameters where $(M^*/M)_0$ denotes the ratio of the effective nucleon mass to the free one at the normal nuclear density ρ_0

PRM	$(M^*/M)_0$	g_σ	g_ω	B_σ (fm^{-1})	C_σ	A_σ (fm^2)	$g_\sigma^\Lambda/g_\sigma$	$g_\omega^\Lambda/g_\omega$
PM1	0·70	9·94	9·99	23·5	0·0	5·65	0·67	0·67
PM2	0·70	9·94	9·99	23·5	0·0	5·65	0·21	0·0

Before showing our results, we would like to make a few comments. We assume thermal and chemical equilibrium matter, which should be achieved by the strong interaction. Then, there are some conserving quantities during the processes: the baryon number B, isospin charge I and strangeness S. Correspondingly we assign the chemical potentials μ_i $(i = B, I, S)$. We, hereafter, consider the isospin symmetric matter, $\mu_I = 0$. Then the chemical potential for each particle can be written in terms of them; e.g. $\mu_K = \mu_S$ and $\mu_\pi = \mu_I$. The number density of each quasiparticle can be written in the same form of the noninteracting case within the RMF theory: for kaons, e.g. it generally reads

$$\rho_K = \zeta^K + \rho_K^*, \quad \rho_K^* = \int d^3k/(2\pi)^3(f_+ - f_-)\,, \tag{3}$$

where $f_\pm$ is the distribution function of $K^\pm$, $f_\pm = 1/\{\exp[\beta(\omega_\pm \mp \mu_S)] - 1\}$, and ζ^K denotes the condensate above the critical density.

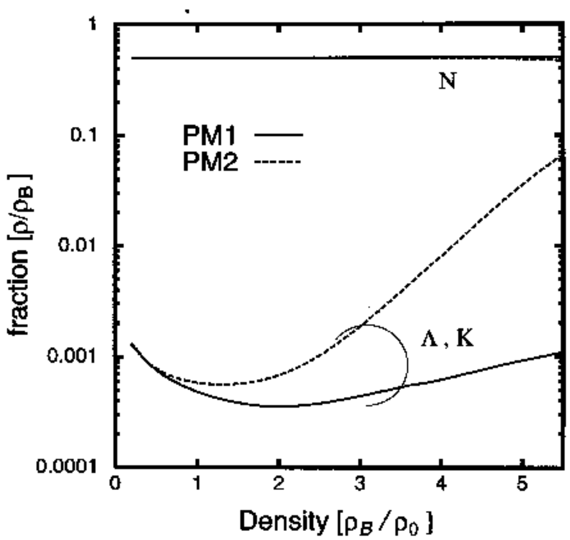

Fig. 2. Particle fractions as functions of density at $T = 50$ MeV.

In Fig. 2 we first show the particle fractions as functions of density at $T = 50$ MeV. We can see that the nucleon fraction is dominant over a range of densities and the fractions of K and Λ are less than 0·1. This also means that the strangeness fraction in the baryon sector is at most 10%. Here it would be interesting to compare our result with the previous one. The Frankfurt group emphasized that hyperons, especially the lambda, can be easily mixed in hadronic matter; in the extreme case the strangeness fraction is almost one [1].

On the contrary, our result shows that when the strangeness mesons are properly included its fraction is not so large. On the other hand, we can see the steep growth of the strangeness chemical potential μ_S in Fig. 3; for the parameter set PM2, it will reach the lowest excitation energy of kaons around $6\rho_0$ (the nuclear density ρ_0 is $0\cdot 17\,\mathrm{fm}^{-3}$). In other words it signals kaon condensation. It is interesting to note that a 10% fraction is sufficient for condensation.

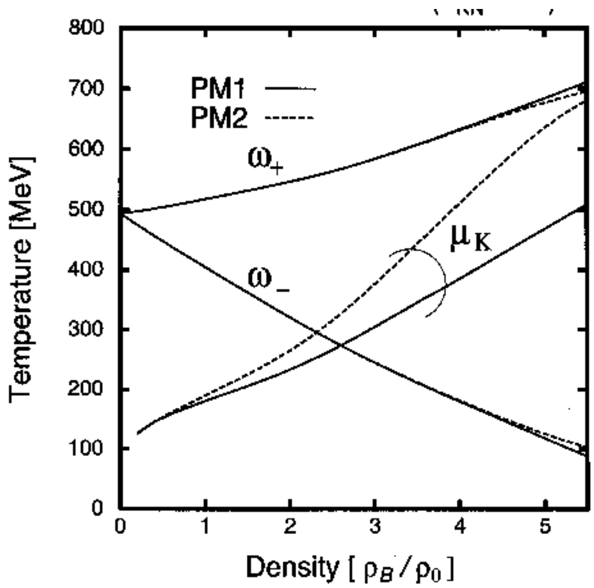

Fig. 3. Frequencies of kaonic modes and the strangeness chemical potential at $T = 50$ MeV.

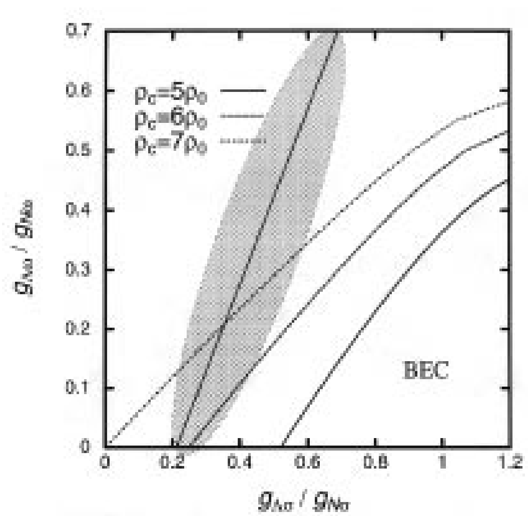

Fig. 4. Kaon condensation (BEC) in the parameter plane at $T = 50$ MeV. The error ellipsis (shaded area) is estimated following Ref. [1].

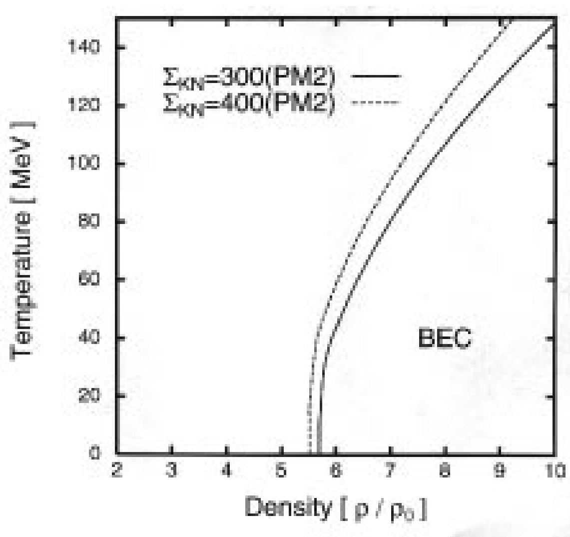

Fig. 5. Phase diagram of kaon condensation for two values of the sigma term Σ_{KN}.

Taking into account the ambiguities of parameters, we depict the regions of kaon condensation in the two-dimensional $g_\sigma^\Lambda - g_\omega^\Lambda$ plane (Fig. 4). The border lines between condensation and non-condensation are drawn for the critical densities 5, 6 and 7 times ρ_0, below which kaon condensation is possible. We also plotted the line with $U_\Lambda = 2/3U_N$, around which the error ellipsis is formed. The interesting region is then the intersection of both domains. We also examined another case with $\Sigma_{KN} = 400\,\mathrm{MeV}$ and the results are summarized as a phase diagram (Fig. 5).

3. Dilepton Production I

In this section we look at the dilepton production through the $K^+K^- \to e^+e^-$ process, which carries information about the abundance of kaons in hot and dense nuclear matter.

We assume thermally equilibrated hadronic matter at rest and the vector dominance hypothesis with ideal mixing in the electromagnetic (EM) interaction of kaons. The dilepton production rate with invariant mass M and total momentum $\mathbf{q}$ is readily computed with obvious notation as [14, 8]

$$\begin{aligned}
\frac{d^4R^{ee}}{dq^3dM} &= e^2/(q^2)^2 H_{\mu\nu}L^{\mu\nu} \\
&= \prod_{i=\pm}\left(\int \frac{d^4k_i}{(2\pi)^4}\frac{\mathrm{Im}D_K^R(k_i)}{e^{\beta(k_i^0-\mu_i)}-1}\right)\prod_{i=\pm}\left(\int\frac{d^3p_i}{(2\pi)^3 2E_i}\right) \\
&\quad\times |\mathcal{M}|^2(2\pi)^4\delta^4(k_+ + k_- - p_+ - p_-) \\
&\quad\times \delta^3(\mathbf{k}_+ + \mathbf{k}_- - \mathbf{q})\delta(M - \sqrt{(\omega_+ + \omega_-)^2 - \mathbf{q}^2})\,, \qquad (4)
\end{aligned}$$

$$|\mathcal{M}|^2 = (e^4/q^4)\Gamma^\mu\Gamma^\nu \mathrm{Tr}(\not{p}_+\gamma_\mu\ \not{p}_-\gamma_\nu)|F(q^2)|^2\,, \tag{5}$$

$$F(q^2) = \tfrac{1}{2}\,\frac{m_\rho^2}{q^2 - m_\rho^2 + i\Gamma_\rho m_\rho} + \tfrac{1}{6}\,\frac{m_\omega^2}{q^2 - m_\omega^2 + i\Gamma_\omega m_\omega} + \tfrac{1}{3}\,\frac{m_\phi^2}{q^2 - m_\phi^2 + i\Gamma_\phi m_\phi}\,. \tag{6}$$

Here $k_\pm$ ($p_\pm$) is the $K^\pm$ ($e^\pm$) momentum, $L^{\mu\nu}$ the leptonic tensor and $H^{\mu\nu}$ the hadronic tensor;

$$H^{\mu\nu} = 2\ \mathrm{Im}\ \mathrm{FT}\ i\langle\!\langle [J_\mu(x), J_\nu(0)]\rangle\!\rangle/(\exp(q_0/T) - 1)\,, \tag{7}$$

where 'FT' means taking the Fourier transform. The kaon dispersion relation is described by the in-medium propagator [6];

$$D_K^{R-1}(k_0, \mathbf{k}) \equiv (k - j_V/(2f^2))^2 - m_K^2 + \Sigma_{KN}\rho_s/f^2 + [d_{Ki1}(k\cdot v)^2 + d_{Ki2}k^2]\rho_B(i)$$

$$\equiv Z^{-1}\omega^2 + \omega j_V^0/f^2 - \mathbf{k}\cdot \mathbf{j}_V/f^2 - Z'^{-1}\mathbf{k}^2 - \tilde{m}_K^2 = 0\,, \tag{8}$$

where $Z^{-1} = 1 + (d_{Ki1} + d_{Ki2})\rho_B(i)$, $Z'^{-1} = 1 + d_{Ki2}\rho_B(i)$, $\tilde{m}_K^2 = m_K^2(1 - \Sigma_{KN}\rho_s/f^2m_K^2)$, and ρ_s is the scalar density. Here v is the nucleon velocity field, $v^\mu = (1, \mathbf{0})$, while d_{Ki1} and d_{Ki2} are the parameters so chosen as to reproduce the S-wave KN scattering data (summation for $i = p, n$ should be understood). We can easily see that our Lagrangian (1) gives the same dispersion equation except for the $d_{K1,2}$ terms. In accord with the in-medium kaon dispersion relation, we take the vertex function as

$$\Gamma^\mu(k, -p) = \left\{\begin{array}{c} Z^{-1}(k_0 - p_0) \\ Z'^{-1}(\mathbf{k} - \mathbf{p}) \end{array}\right\} + j_V^\mu/f^2\,, \tag{9}$$

in order to respect EM gauge invariance (the Ward–Takahashi identity).

The most transparent connection between the kaon dispersion relation and the dilepton spectrum occurs in the back-to-back kinematics [$\mathbf{q}$=0, $M = 2\bar{\omega}(\mathbf{k})$ with $2\bar{\omega} = k_+^0 + k_-^0$], where Eq. (4) is reduced to

$$\left.\frac{d^4R^{ee}}{dq^3dM}\right|_{\mathbf{q}=0} = \frac{\alpha^2}{3(2\pi)^4}\frac{k^4}{\bar{\omega}^4}\left|\frac{d\bar{\omega}}{dk}\right|^{-1} f_+f_-|F(M^2)|^2\,. \tag{10}$$

For a given value $M = 2\bar{\omega}$, the medium effect will appear in (i) the size of the allowed phase space k^3, and in (ii) the EM interaction, $|F(M^2)|^2$.

We show in Fig. 6 the dilepton production rate from kaon–antikaon annihilation at $\rho_B = 3\rho_0$ and $T = 100$ MeV, as well as the contribution from the pion-pair annihilation (dash–dot curve). First, when we just neglect the modification of the ϕ meson width due to the kaon softening, we see the larger enhancement of the rate (dashed curve) than that from the free kaons (thin solid curve), which manifests the enlargement of the corresponding kaon phase space. Kaon softening will cause the broadening of the ϕ-meson width at the same time. In our model, it will be 63 MeV at $\rho_B = 3\rho_0$ compared with the free width, 4·4 MeV. Altogether, the sharp ϕ-meson peak is reduced to a bump structure in the dilepton production rate.

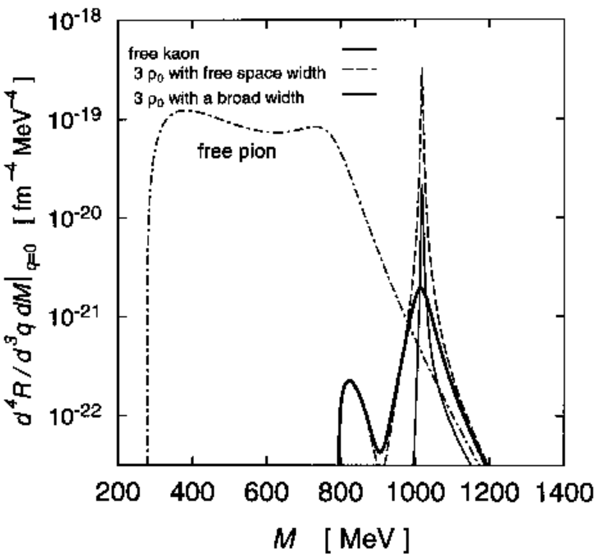

Fig. 6. Dilepton production rate from kaon-pair annihilation at $\rho_B = 3\rho_0$ and $T = 100$ MeV.

4. Dilepton Production II

Next we consider dilepton production in the kaon condensed phase. As we have already seen in Section 2, there is a possibility of kaon condensation. If the kaon condensed state is created over an extended space–time region in the relativistic heavy-ion collisions, it can generate observable effects in the production process of the dilepton. We know that such a state can be described as a chiral rotated one in the $SU(3) \times SU(3)$ space [6, 15];

$$|K\rangle = U_K|0\rangle, \quad U_K = \exp(i\mu_S t\hat{Y})\exp(i\theta F_4^5)\,, \tag{11}$$

where $F_\alpha(F_\alpha^5)$ are the vector (axial-vector) charges satisfying current algebra, and $\hat{Y}$ the hypercharge. Here θ denotes the chiral angle and it is an order parameter of the condensed phase.

First, consider the matrix element of the EM current $J_\mu = V_{3\mu} + 1/\sqrt{3}V_{8\mu}$ of hadrons in the condensed phase $|K\rangle$;

$$\langle K; f|J_\mu|i; K\rangle = \langle f|\tilde{J}_\mu|i\rangle\,, \tag{12}$$

for a transition $i \to f$. Then the chiral-rotated current $\tilde{J}_\mu$ can be calculated by way of current algebra, as

$$\tilde{J}^\mu = U_K^\dagger J_\mu U_K = aV_3^\mu + bV_8^\mu + cA_5^\mu\,, \tag{13}$$

where the coefficients are $a = \frac{1}{2}(1+\cos\theta), b = \frac{1}{2}\sqrt{3}(-1+3\cos\theta)$ and $c = -\sin\theta$. For the current–current commutator, which directly comes in the hadronic tensor $H_{\mu\nu}$ (7), we then find

$$[\tilde{J}_\mu(x), \tilde{J}_\nu(0)] \simeq (1-\theta^2/2)[V^3_\mu(x), V^3_\nu(0)] + 1/3(1--3/2\theta^2)[V^8_\mu(x), V^8_\nu(0)]$$

$$+ \theta^2[A^5_\mu(x), A^5_\nu(0)] - \theta\{[A^5_\mu(x), J_\nu(0)] + [J_\mu(x), A^5_\nu(0)]\} \qquad (14)$$

for the small condensate. From these equations we can clearly see *vector–axial—vector mixing* in the condensed phase.* The first two terms are ordinary ones with some reduction, while the last two terms are inherent in the condensed phase. The vector dominance hypothesis tells us that hadrons can couple with not only vector mesons but also axial-vector strange mesons (K_1 for the lowest mass). Considering the concrete process $K^+K^- \to e^+e^-$, we depict each contribution in Fig. 7. Substituting Eq. (14) in Eqs (4) and (7) we finally get the production rate of dileptons. In Fig. 8 we show one result to demonstrate the effect of the condensate. It would be interesting to compare this with Fig. 6.

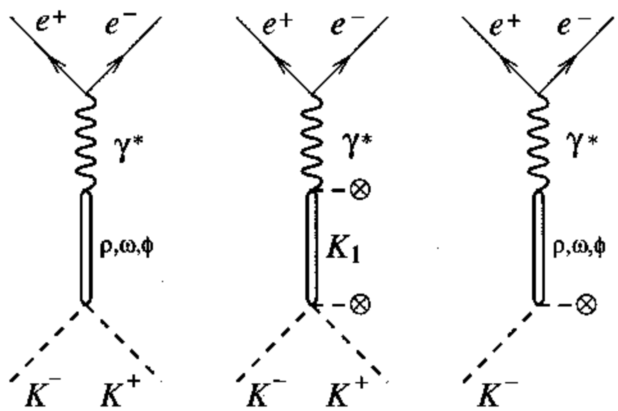

Fig. 7. Three contributions read from Eq. (14) where the cross-circle represents the condensate.

Two remarks are in order: (i) since the hadronic matter $|0\rangle$ is isospin-symmetric, $U(\alpha)|0\rangle = |0\rangle, U(\alpha) = \exp(i\alpha \cdot \mathbf{Q}) \in SU(2)_{isospin}$, and we may generate various condensed phases depending the parameter α. Our choice here corresponds to $\alpha = 0$, which in turn means there is only the K^+ condensate. For other choices of the parameter we see that there are K^+ and K^0 condensates at the same time and Eq. (13) is slightly changed. The degree of freedom for this choice is, however, redundant and they should give the same result for the physical quantities as it does. (ii) There is another interesting effect coming from the condensate. It is a generation of photon mass m_γ; since the order parameter is complex, the photon can interact with surrounding condensate and get a finite mass [15]. We can easily estimate its value as $m_\gamma^2 = 2e^2|\langle K\rangle|^2 \simeq e^2 f^2\theta^2$, and find $m_\gamma \sim$ several MeV. Although it is very small, it would modify the photon propagator and thereby a new dilepton peak appears at very small invariant mass. We think, however, it would be very hard to detect this peak, since in this region there is a huge background due to Bremsstrahlung.

* We can refer to similar phenomena caused by thermal mesons [16] or a disoriented chiral condensate [17].

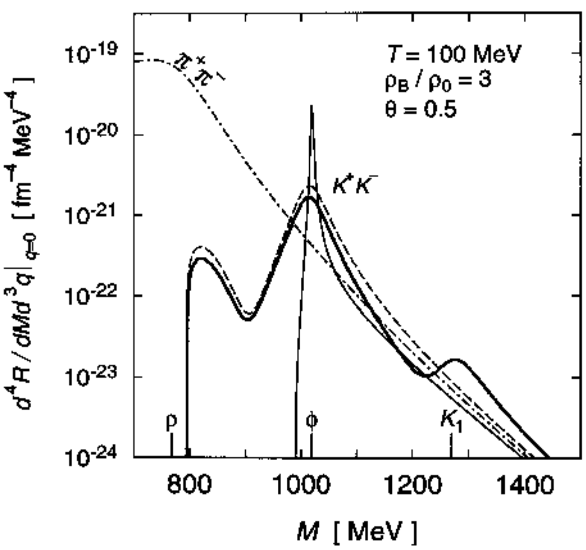

Fig. 8. Dilepton production from the kaon condensed phase. The dashed curve corresponds to the case $\theta = 0$. The case without any interaction is also depicted for reference (thin curve). Besides the ϕ-meson peak we can see a new peak coming from the K_1 meson.

5. Summary and Concluding Remarks

We have studied the bulk properties of strange hadronic matter in the relativistic mean-field theory, where the single-particle energy levels of Λ hypernuclei are reproduced. Contrary to previous studies of strange hadronic matter, we have taken into account the kaonic degrees of freedom besides hyperons. We have found the steep rise of the strangeness chemical potential, while the strangeness fraction is not so large. In some parameter regions we have seen that K^+ condensation would occur at several times the nuclear density. In the light of ambiguities in the values of the coupling constants we cannot hope to decide conclusively on the basis of relativistic mean-field theory whether or not K^+ condensation occurs. We can say, however, that the critical density is much reduced from the non-interacting case, the values of the parameters which lead to K^+ condensation are not unusual and condensation does not necessarily demand a large abundance of strangeness. If stopping power is enough to create high-density 'matter' or a shock wave rapidly compresses nuclear matter to very high density in heavy-ion collisions [18], we can expect this possibility.

Next we examined how the medium effect on kaons shows up in the production rate of dileptons. We have seen a large enhancement of the rate around the ϕ-meson peak, if the modification of its width is discarded. However, once the broadening effect of the ϕ-meson width is taken into account, kaon softening results in a reduction of the sharp ϕ-meson peak to a bump in the dilepton production rate. If the K^+ condensate is created over an extended space–time region following heavy-ion collisions, the chiral rotation can generate observable effects in the process of dilepton production: the vector–axial-vector mixing induced by the condensate gives rise to a new peak or bump around the K_1 meson.

Acknowledgments

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