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Medium effects on kaon and antikaon spectra in heavy-ion collisions

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In the linear chiral perturbation theory, both kaon and antikaon masses decrease in dense matter. There is also a repulsive vector potential for the kaon and an attractive one for the antikaon. With these effects included in the relativistic transport model, it is found that the slope parameter of the kaon kinetic energy distribution is larger than that of the antikaon. This is consistent with the experimental data from heavy-ion collisions in the Alternating Gradient Synchrotron experiments at Brookhaven.

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In relativistic heavy-ion collisions, nuclear matter can be compressed to densities which are many times that of normal nuclei. This has recently generated great interest in theoretical studies of hadron properties under extreme conditions [1]. It was suggested [2, 3] that a novel form of dense matter, charged kaon condensation, would occur in dense nuclear matter. Because of the explicit chiral symmetry breaking, nucleons act on the kaon as an effective scalar field

$$S \sim \frac{\bar{N}N}{f_K^2}, \tag{1}$$

where $f_K \sim 93$ MeV is the kaon decay constant and N denotes the nucleon field. The effective scalar field couples to the kaon with a "coupling constant" equal to half the KN sigma term Σ^{KN}, i.e.,

$$\delta L = \tfrac{1}{2}\Sigma^{KN} S\bar{K}K. \tag{2}$$

When combined with the kaon mass term $m_K^2\bar{K}K/2$ in the free field Lagrangian it gives an effective kaon mass m_K^*

$$m_K^{*2} = m_K^2 - \Sigma^{KN}\langle S\rangle\ . \tag{3}$$

This can be rewritten as

$$m_K^* = m_K\left(1-\frac{\rho_B}{\rho_c}\right)^{1/2}, \tag{4}$$

in terms of the critical density ρ_c for kaon condensation which is given by

$$\rho_c = f_K^2 m_K^2/\Sigma^{KN}. \tag{5}$$

The KN sigma term Σ^{KN} is defined by

$$\Sigma^{KN} = \tfrac{1}{2}(m_u+m_s)\langle N|\bar{u}u+\bar{s}s|N\rangle \tag{6}$$

with m_u and m_s being, respectively, the up and strange quark masses. There is considerable uncertainty in the strangeness content of the nucleon [4], but the lower and upper limits on ρ_c can be obtained by taking $\langle N|\bar{s}s|N\rangle = \langle N|\bar{u}u|N\rangle$ and by taking $\langle N|\bar{s}s|N\rangle = 0$. With $m_s = 25\frac{m_u+m_d}{2}$ and the pion-nucleon sigma term $\Sigma_{\pi N} = \frac{m_u+m_d}{2}\langle N|\bar{u}u+\bar{d}d|N\rangle \simeq 46\,\mathrm{MeV}$, we find $300\,\mathrm{MeV} \le \Sigma_{\mathrm{KN}} \le 600\,\mathrm{MeV}$, and thus

$$2.5\rho_0 < \rho_c < 5\rho_0. \tag{7}$$

The decrease of the kaon effective mass in the nuclear matter may be important in understanding the enhancement of kaon production in heavy-ion experiments [5, 6]. We note, however, that the critical density estimated above will be modified by the energy dependence of the KN scattering amplitude. A similar effect has been shown to lead to the suppression of the reduced pion effective mass in dense matter [7]. But we believe that the kaon mass eventually will reduce to the mass of the strange quark (plus that of the up or down quark) at high densities when the chiral symmetry is restored. This is indeed demonstrated in studies with effective Lagrangians which include broken scale and chiral invariance [8].

There is also a vector interaction in the chiral Lagrangian. For an $N=Z$ system, it can be reduced to

$$L_V = -\frac{3}{8f_K^2} i\bar{N}\gamma^0 N\bar{K}\overleftrightarrow{\partial}_t K. \tag{8}$$

This leads to a repulsive vector potential for a kaon in nuclear matter [9]

$$V_{KN} = \frac{3}{8f_K^2}\rho_B = \frac{1}{3}\frac{g_\omega^2}{m_\omega^2}\rho_B. \tag{9}$$

The second equality is obtained by using the KFSR relation $m_\rho = 2\sqrt{2}f_K g_\rho$ and the SU(3) relation $g_\omega = 3g_\rho$. We see that the kaon-nucleon vector-exchange mean-field potential is just 1/3 of the nucleon-nucleon mean-field potential $V_\omega = (g_\omega^2/m_\omega^2)\rho_B$, commonly used in the Walecka theory of nucleon-nucleon interactions [10]. This potential becomes attractive for the antikaon.

The difference in the sign of the vector potential for the kaon and the antikaon can also be understood in terms of their quark content. The quark content of a kaon is $(q\bar{s})$ and of an antikaon is $(\bar{q}s)$, where q and $\bar{q}$ denote the light

u and d quarks and antiquarks, respectively. Like the nucleon, the kaon would feel a repulsive vector potential via the q quark. But the antikaon feels an attractive one because of the $\bar{q}$ antiquark, as in the case of the antinucleon.

It has been pointed out in Ref. [11] that the difference between the kaon and antikaon vector potentials might lead to a difference in their apparent temperatures in heavy-ion collisions. It has also been shown there that the similar difference in the proton and antiproton vector potentials may be responsible for the observed difference between the apparent temperatures of the proton and the antiproton in the Alternating Gradient Synchrotron (AGS) heavy-ion experiments at Brookhaven [12].

It is, of course, dangerous to identify the vector coupling g_ω with the one of the Walecka model, which has been adjusted in order to reproduce the saturation properties of nuclear matter. Typically, in the Walecka model one uses vector couplings which are smaller than the ones extracted from nucleon–nucleon scattering data. It is, however, interesting to note that, according to a detailed analysis of kaon-nucleon scattering data by the Jülich group [13, 14], only about 40% of the short-range repulsion is due to G-parity odd vector exchange and thus leads to attraction in the $\bar{K}N$ channel. The other 60% does not have a clear G-parity and essentially does not show up in the $\bar{K}N$ scattering. In this sense, it is not a bad choice to take the coupling constant from the Walecka model. The difference between slopes of K^+ and K^- spectra in the following study is thus a lower limit. A larger difference is expected if the realistic K^+ repulsive vector potential is used.

Is there any experimental evidence for the medium effect on the kaon and antikaon dynamics in dense matter? In the recent heavy-ion experiments at the AGS [6], both the K^+ and K^- transverse kinetic energy spectra have been measured. Although the experimental uncertainty in the K^- slope parameter is large, the data indicate that there is about 28 MeV difference in the slope parameters of K^+ and K^-.

The above discussed medium effects can be incorporated in the relativistic transport model [15, 16] based on the relativistic mean-field theory of Walecka [10]. It describes the propagation of nucleons in the self-consistent nuclear mean-field potential generated by the scalar and vector mesons. The nucleon mass in the nuclear medium is thus modified by the scalar field. Following the scaling law of Ref. [17] for hadron masses, the in-medium masses of other nonstrange hadrons are taken to have a similar density dependence as the nucleon. For the hyperons (Λ and Σ), their in-medium masses are determined by taking the effect of the scalar field to be 2/3 of that for the nucleon [18, 19].

The relativistic transport model allows also for nucleon-nucleon elastic and inelastic collisions with the latter via the excitation of nucleons to deltas. Since kaons can be produced not only from nucleon-nucleon interactions, such as $NN \to NYK$ [20], but also from pion-nucleon interactions [21] $\pi N \to YK$ and pion-pion interactions [22] $\pi\pi \to K\bar{K}$, we extend the transport model developed in Ref. [15] to include the pion via the reaction $\Delta \leftrightarrow N\pi$. In this way, the pion is produced from the delta decay and can also be absorbed by the nucleon to form the delta again. For the cross section $\sigma_{\pi N\to\Delta}$, we take it from Ref. [23]. The propagation of the pion between collisions is treated as that for a free particle as in most transport models [24–26], except in Ref. [27] where the medium effect on the pion is included via the delta-hole model [28, 29]. Our transport model includes, thus, the following particles: N, Δ, Y, π, K, $\bar{K}$ and ρ. The relevant cross sections will be discussed in the following.

Both kaon and antikaon undergo collisions with other particles. The interaction of the kaon with the nucleon is relatively weak and the cross section [30] is $\sigma_{KN} \sim 10$ mb. Due to the resonance nature of the interaction, the antikaon-nucleon cross section is much larger and can be approximately taken to be [30] $\sigma_{\bar{K}N} \sim 40$ mb. Both the kaon and antikaon can also interact with the pion via the K^* resonances. The isospin averaged cross section is given by [31]

$$\sigma_{K\pi} = \frac{\sigma_0}{1 + 4(s^{1/2} - m_{K^*})^2/\Gamma_{K^*}^2}, \tag{10}$$

where $\sigma_0 \sim 60$ mb. The mass and width of K^* are given by $m_{K^*} \sim 895$ MeV and $\Gamma_{K^*} \sim 50$ MeV, respectively.

All cross sections are expected to be modified in the medium [5, 32–35]. Since the density dependence of these cross sections are not well known, we will use the free cross sections in the present study as in other transport models except the following cross sections, $\sigma_{\pi N\to\Delta}$, $\sigma_{\pi\pi\to\rho}$, $\sigma_{K\pi\to K^*}$, $\sigma_{\bar{K}N\to\Lambda\pi}$, and $\sigma_{MM\to K\bar{K}}$, where M denotes either a pion or a rho meson. The medium dependences of the first three cross sections are determined from the Breit-Wigner form with the in-medium masses and widths. The in-medium widths of the resonances are evaluated using the in-medium masses of the particles. Both $\sigma_{\bar{K}N\to\Lambda\pi}$ and $\sigma_{MM\to K\bar{K}}$ in the nuclear medium have been calculated in Ref. [5] and will be used in the relativistic transport model.

In Ref. [5], kaon production from heavy-ion collisions at the AGS energies has been studied in the hydrochemical model with in-medium hadron masses. The development there is, however, not thermodynamically consistent because only the masses are changed in the noninteracting particle model without introducing corrections to the energy and pressure of the system from the particle interactions. With the relativistic transport model, the medium dependence of hadron effective masses and the thermodynamics of the system are treated in a consistent framework.

As in Ref. [5], we assume that initially a highly excited and compressed fireball is formed in heavy-ion collisions. Its expansion is then described by the relativistic transport equation. Initially, the fireball is assumed to have a temperature of 190 MeV and a density of $4\rho_0$, where ρ_0 is the normal nuclear matter density. The initial energy density, including both the particle kinetic energy and the collective scalar and vector field energy, is then about 1.6 GeV/nucleon. This initial condition corresponds approximately to that one expects for heavy-ion collisions at the AGS energies. In the initial fireball, we include nucleons, deltas, pions, and rho mesons by assuming that they are in thermal and chemical equilibrium. Since the strange particles are not expected to be in equilibrium, we fix the initial number of kaons and antikaons to be that of proton nucleus collisions, i.e., about ten kaons and four antikaons. From the strangeness conservation, the hyperon number is thus 6. All particles are assumed to be uniformly distributed inside the fireball of radius ~ 3.5 fm, which roughly corresponds to the radius of the silicon target and leads to a baryon number of the fireball of about 100.

The fireball then expands, and kaons and antikaons are produced from the reaction $MM \to K\bar{K}$. We have ignored kaon production from the processes $NN \to NYK$ and $\pi N \to YK$, as they are unimportant in the expanding fireball [5]. Both kaons and antikaons with the in-medium masses propagate through the hadronic matter under the influence of the mean-field potential and undergo collisions with both nucleons and pions. Furthermore, antikaons can be destroyed via the reaction $\bar{K}N \to Y\pi$. In our calculations, we take the critical density for kaon condensation to be $\rho_c = 5\rho_0$. The scalar and vector meson coupling constants $g_\sigma^2/(4\pi) \sim 3.82$ and $g_\omega^2/(4\pi) \sim 4.54$ are taken from the parameter set in Ref. [15], which gives a nucleon effective mass $m_N^* = 0.83\,m_N$ and a nuclear compressibility $K = 380$ MeV of the normal nuclear matter. With this value of the vector meson coupling constant and the vector meson mass $m_\omega = 783$ MeV, the strength of the vector mean-field potential is about 38 MeV for the kaon and -38 MeV for the antikaon at the normal nuclear density.

Unlike the hydrochemical model, there is no need in the transport model to introduce a freeze out density. We simply let the system expand to a very low density, $\sim 0.3\rho_0$ in the present calculation, when it is practically free of both mean-field potential and particle collisions. The fireball reaches such a low density after about 10 fm/c. In Fig. 1 we show the time evolution of the particle abundance. As in Ref. [5], most kaons are created within the first 2 fm/c of the fireball expansion. The number of hadron resonances such as delta and rho meson are less in the present calculations than in Ref. [5], as different in-medium hadron masses are employed in the two studies.

To take into account possible acceptance limitations and the nuclear transparency effect, an overall normal-

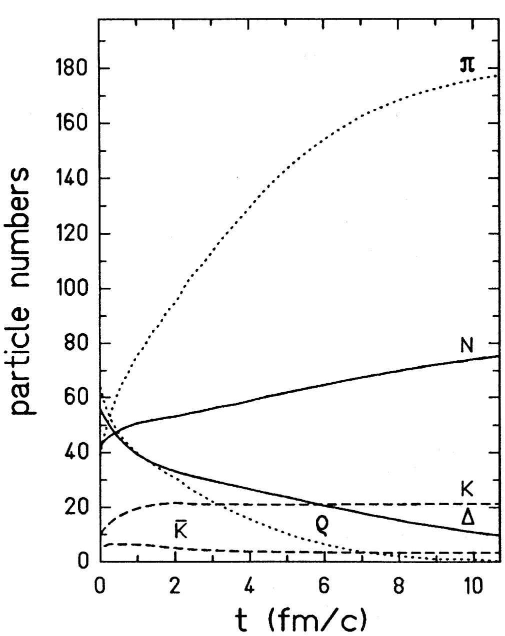

FIG. 1. The time evolution of the particle abundance.

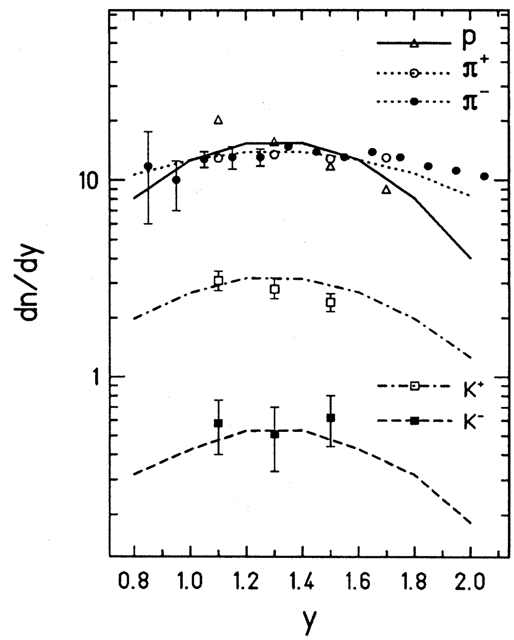

FIG. 2. The rapidity distribution of particles. Curves are from the theoretical calculations while data are from Ref. [6].

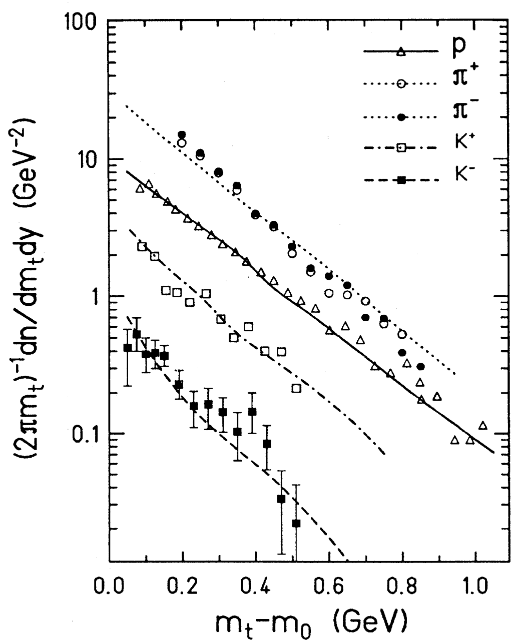

FIG. 3. Same as Fig. 2 for the transverse kinetic-energy distributions of particles.

ization factor of 0.6 is introduced and the center-of-mass rapidity is taken to be 1.4 as in Ref. [5]. The final rapidity distributions of the particles are shown in Fig. 2. The result is very similar to that of Ref. [5]. We see that the pion distribution agrees reasonably with the data. The failure of the calculated proton distribution at smaller rapidities is due to the neglect of protons from the target spectator. Both kaon and antikaon distributions agree also fairly well with the data. The lack of kaons in the small rapidity compared with the data can be understood if we allow some of the initial kaons to be produced from the interaction of pions with the target spectator nucleons. The transverse kinetic energy distributions of the particles are shown in Fig. 3. We see that all particles have essentially exponential distributions and they agree again with the data as in the hydrochemical model. The pion spectrum has a slope parameter that is slightly higher than the measured one. Including the collective mean-field potential of the pion, which has been neglected in the present study, is expected to soften the pion spectrum as shown recently in Ref. [27] for the pion transverse energy spectrum from the Bevalac experiments. The slope parameters for kaons and antikaons indeed show the difference expected from the mean-field effects, i.e., the effective temperature of antikaons is seen to be lower than that of kaons.

In summary, we have carried out a transport model calculation including the medium effects on the kaon and antikaon dynamics in dense matter. The theoretical results in both rapidity and transverse kinetic energy distributions of the particles agree with the data from recent AGS experiments. While the rapidity distribution and the relative yields may be affected by the aforementioned uncertainties in the effective mass of the kaon, the difference in the slopes should remain unchanged, and, as already mentioned, should be considered as a lower limit. The initial states of the systems such as the temperature and density, have been treated as parameters. It will be more satisfactory if it can also be determined consistently from the transport model. Encouraged by the recent success of the relativistic hadronic cascade model of Pang *et al.* [36] for heavy-ion collisions at the AGS energies, we are extending our relativistic transport model to treat the initial stage of the collisions as well.

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