Medium Modifications of the Rho Meson at CERN Super Proton Synchrotron Energies (200 GeV/nucleon)

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Rho meson propagation in hot hadronic matter is studied in a model with coupling to $\pi\pi$ states. Medium modifications are induced by a change of the pion dispersion relation through collisions with nucleons and Δ's in the fireball. Maintaining gauge invariance dilepton production is calculated and compared to the recent data of the CERES Collaboration in central S + Au collisions at 200 GeV/u. The observed enhancement of the rate below the rho meson mass can be largely accounted for.

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In central heavy-ion reactions at ultrarelativistic energies hot and dense hadronic matter is formed at the early stages of the collision. From such experiments one hopes to infer properties of mesons and baryons in the vicinity of the chiral phase transition. Because of its electromagnetic origin, lepton pair emission is an ideal probe. Once produced, the dilepton will decouple from the strongly interacting particles and thus carry undistorted information of the dynamical properties of the hot nuclear system. For this reason the recent low-mass dilepton spectra measured by the CERES Collaboration in p + Be, p + Au, and S + Au [1] collisions at 450 and 200 GeV/u, respectively, have produced considerable interest. While the proton data are well described by known decays from a hadronic fireball an excess of dielectrons with invariant masses of 0.2 to 1 GeV/c^2 has been observed in the S + Au case and has been attributed to $\pi^+\pi^-$ annihilation from the pionic component of the hadron gas [1]. When this process is included in the description of the data the excess of strength between the $2m_\pi$-vacuum threshold and the rho peak is still unexplained. This led the authors of Ref. [2] to the conclusion that one is seeing a signal of partial restoration of chiral symmetry through a decrease of vector meson masses (here especially the rho meson) near the phase boundary. Since the rho meson strongly couples to the $\pi^+\pi^-$ channel this effect may also be explained more conventionally, however, namely by modifications of the pion propagation in the presence of nucleons and isobars [3–9].

To study the dilepton emission from $\pi^+\pi^-$ annihilation we invoke the vector-dominance model (VDM) in which the electromagnetic current is related to the third isospin component of the rho meson field by the field current identity $J^\mu = (m_\rho^2/g)\rho_3^\mu$ where m_ρ and g denote the rho meson mass and the universal VDM coupling constant, respectively. In terms of the retarded (thermal) rho-meson propagator

$$D_\rho^{\mu\nu}(q;T) = -i\int d^4x e^{-iqx}\sum_i \frac{e^{-\beta E_i}}{Z} \times \langle i|[\rho_3^\mu(x),\rho_3^\nu(0)]|i\rangle\Theta(x_0) \quad (1)$$

($\beta = 1/T$) the dilepton emission rate per unit volume, $R = dN_{l^+l^-}/d^4x$, is given by

$$\frac{dN_{l^+l^-}}{d^4xd^4q} = \frac{\alpha^2}{3\pi^2M^2}H(q;T), \quad (2)$$

where $q = p_+ + p_- = (q_0,\vec{q})$ denotes the total momentum for (massless) dileptons, $M^2 \equiv q_\mu^2 > 0$ is the invariant mass, and the hadronic piece $H(q;T)$ is given by

$$H(q;T) = f^\rho(q_0,T)[(m_\rho^{(0)})^4/\pi g^2]g_{\mu\nu}\,\mathrm{Im}D_\rho^{\mu\nu}(q;T) \quad (3)$$

$[f^\rho(q_0,T) = (e^{q_0/T}-1)^{-1}]$. To reduce the numerical complexity in evaluating $D_\rho^{\mu\nu}$ we shall restrict ourselves to back-to-back kinematics, i.e., $\vec{q} = 0$. In that case, from gauge invariance, only the space components contribute. Denoting $D_\rho^{ij}(q_0,\vec{q}=0;T) = \delta^{ij}D_\rho(q_0;T)$ Eq. (3) then reduces to

$$H(q_0,\vec{q}=0;T) = -f^\rho(q_0,T)\frac{3(m_\rho^{(0)})^4}{\pi g^2}\,\mathrm{Im}D_\rho(q_0;T), \quad (4)$$

where

$$D_\rho(q_0;T) = [q_0^2 - (m_\rho^{(0)})^2 - \Sigma_\rho(q_0;T)]^{-1}, \quad (5)$$

involving the scalar part of the rho self-energy Σ_ρ.

To study medium modifications of the ρ propagator we first need a realistic model for the ρ meson in free space. As is well known [10] an appropriate description is provided by a bare pole graph renormalized by $\pi\pi$ rescattering. This amounts to solving a Lippmann-Schwinger type equation

$$M_{\pi\pi}(E,q_1,q_2) = V_{\pi\pi}(E,q_1,q_2) + \int_0^\infty dk\,k^2V_{\pi\pi} \times (E,q_1,k)G_{\pi\pi}^0(E,k)M_{\pi\pi}(E,k,q_2), \quad (6)$$

where E is the starting energy of the pion pair, q_1 and q_2 denote the moduli of the pion three-momenta in the center

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of mass system frame, and $G^0_{\pi\pi}(E,k)$ is the free two-pion propagator:

$$G^0_{\pi\pi}(E,k) = \frac{1}{\omega_k}\frac{1}{E^2 - 4\omega_k^2 + i\eta}, \quad \omega_k^2 = m_\pi^2 + k^2. \tag{7}$$

The pseudopotential generated by the s-channel ρ pole graph is given by

$$V_{\pi\pi}(E,q_1,q_2) = v(q_1)D^0_\rho(E)v(q_2), \tag{8}$$

where

$$D^0_\rho(E) = [E^2 - (m_\rho^{(0)})^2]^{-1},$$
$$v(q) = \sqrt{\frac{2}{3}\frac{g}{2\pi}}\,2qF(q) \tag{9}$$

are the bare ρ propagator and the vertex functions, respectively. The hadronic form factor is chosen to be of dipole form normalized to one at the physical resonance energy $m_\rho = 0.77$ GeV:

$$F(q) = \left(\frac{2\Lambda_\rho^2 + m_\rho^2}{2\Lambda_\rho^2 + 4\omega_q^2}\right)^2. \tag{10}$$

The parameters g, Λ_ρ, and $m_\rho^{(0)}$ are fitted to the experimental p-wave $\pi\pi$ phase shift and the pion electromagnetic form factor

$$|F_\pi(E)|^2 = \frac{(m_\rho^{(0)})^4}{[E^2 - (m_\rho^{(0)})^2 - \mathrm{Re}\Sigma^0_\rho(E)]^2 + \mathrm{Im}\Sigma^0_\rho(E)^2}$$
$$\equiv (m_\rho^{(0)})^4 |D^0_\rho(E)|^2, \tag{11}$$

with

$$\Sigma^0_\rho(E) = \bar{\Sigma}^0_\rho(E) - \bar{\Sigma}^0_\rho(0),$$
$$\bar{\Sigma}^0_\rho(E) = \int k^2 dk\, v(k)^2 G^0_{\pi\pi}(E,k). \tag{12}$$

The subtraction for Σ^0_ρ at $E = 0$ is necessary to ensure the normalization $F_\pi(0) = 1$. We obtain a satisfactory fit with $g^2/4\pi = 2.7$, $\Lambda_\rho = 3.1$ GeV, and $m_\rho^{(0)} = 0.829$ GeV as indicated in Fig. 1.

A proper calculation of the in-medium ρ propagator

$$D_\rho(E) = [E^2 - (m_\rho^{(0)})^2 - \Sigma_\rho(E)]^{-1} \tag{13}$$

must include constraints from gauge invariance. A detailed derivation for the self-energy in homogeneous matter is given in Refs. [5,6], and we do not repeat it here. The result can be expressed in terms of the spin-isospin longitudinal and transverse response functions at a given density and temperature

$$\Pi_L(k_0,k) = (k_0^2 - k^2 - m_\pi^2)\tilde{\Pi}^0(k_0,k)D_L(k_0,k),$$
$$\Pi_T(k_0,k) = (k_0^2 - k^2 - m_\rho^2)\tilde{\Pi}^0(k_0,k)D_T(k_0,k), \tag{14}$$

where D_L and D_T are the longitudinal and transverse propagators which are expressed in terms of the pion self-energy as

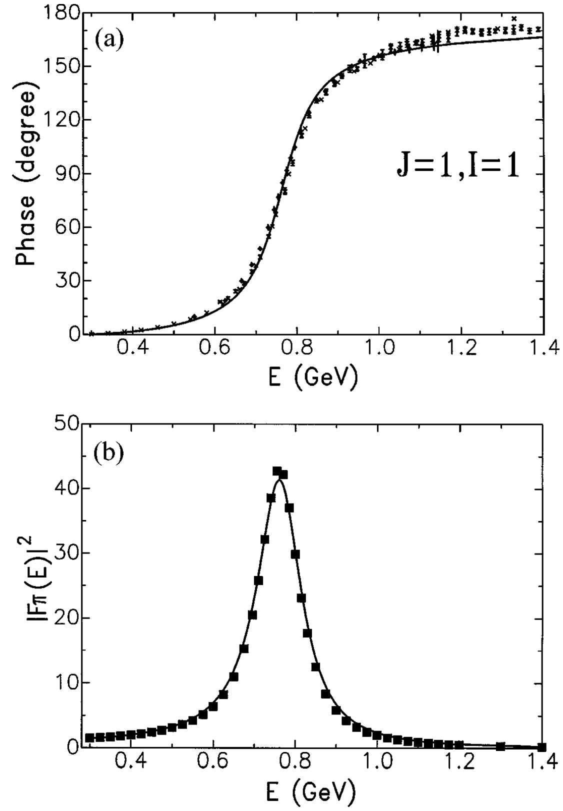

FIG. 1. Our fit to the p-wave $\pi\pi$ phase shifts (upper panel) and the pion electromagnetic form factor (lower panel); the squares in the lower panel are the values from the Gounaris-Sakurai formula [14], which itself gives an accurate description of the data.

$$D_L(k_0,k) = [k_0^2 - k^2 - m_\pi^2 - \Sigma_\pi(k_0,k)]^{-1},$$
$$D_T(k_0,k) = [k_0^2 - k^2 - m_\rho^2 - C_\rho\Sigma_\pi(k_0,k)]^{-1},$$
$$\tilde{\Pi}^0(k_0,k) = \frac{1}{k^2}\Sigma_\pi(k_0,k). \tag{15}$$

Obviously D_L is equal to the single pion propagator D_π. Making use of derivations in Ref. [5] the imaginary part of the self-energy for a ρ meson at rest then takes the form

$$\mathrm{Im}\Sigma_\rho(q_0,\vec{0}) = \int_0^\infty k^2\,dk\,v(k)^2 \int_0^{q_0}\frac{dk_0}{\pi}\,\mathrm{Im}D_\pi$$
$$\times\ (q_0 - k_0,k)\mathrm{Im}[\alpha(q_0,k_0,k)D_\pi(k_0,k)$$
$$+\ \tfrac{1}{2}(1/k^2)\Pi_L(k_0,k) + (1/k^2)\Pi_T(k_0,k)], \tag{16}$$

where the function α is given by

$$\alpha(q_0,k_0,k) = [1 + \tilde{\Pi}^0_R(k_0,k) + \tilde{\Pi}^0_R(q_0 - k_0,k)$$
$$+ \tfrac{1}{2}\tilde{\Pi}^0_R(k_0,k)\tilde{\Pi}^0_R(q_0 - k_0,k)],$$
$$\tilde{\Pi}^0_R(k_0,k) = \mathrm{Re}\,\tilde{\Pi}^0(k_0,k). \tag{17}$$

The real part of Σ_ρ is obtained via a dispersion integral:

$$\mathrm{Re}\Sigma_\rho(q_0) = -\mathcal{P}\int_0^\infty \frac{dE'^2}{\pi}\,\frac{\mathrm{Im}\Sigma_\rho(E')}{q_0^2 - E'^2}\,\frac{q_0^2}{E'^2}. \quad (18)$$

The factor q_0^2/E'^2 results from a subtraction at $q_0 = 0$, which is necessary to ensure gauge invariance [5]; the imaginary part is unaffected by this subtraction since $\mathrm{Im}\Sigma_\rho(q_0 = 0) = 0$.

It remains to specify the model for the single-pion self-energy Σ_π. We employ the standard model of p-wave particle-hole excitations extended to finite temperature [8]. Here the pions are dressed by NN^{-1}, ΔN^{-1} as well as $N\Delta^{-1}$, and $\Delta\Delta^{-1}$ excitations, the latter appearing as a consequence of a thermally excited Δ abundance in the gas. One has

$$\Sigma_\pi(\omega, k) = -k^2 \sum_\alpha \chi_\alpha(\omega, k), \quad (19)$$

where the summation is performed over all excitation channels $\alpha = ab^{-1}$, $a, b = N, \Delta$. The susceptibilities χ_α contain short-range correlations taken into account by Landau-Migdal parameters $g'_{\alpha\beta}$ such that

$$\chi_\alpha = \chi_\alpha^{(0)} - \sum_\beta \chi_\alpha^{(0)} g'_{\alpha\beta}\chi_\beta,$$
$$\chi_\alpha^{(0)}(\omega, k) = \left(\frac{f_{\pi\alpha}\Gamma_\pi(k)}{m_\pi}\right)^2 \mathrm{SI}(\alpha)\phi_\alpha(\omega, k), \quad (20)$$

where $f_{\pi\alpha}$ denote the πNN, etc. coupling constants and $\mathrm{SI}(\alpha)$ is a spin-isospin factor. The form factor $\Gamma_\pi(k) = (\Lambda_\pi^2 - m_\pi^2)/(\Lambda_\pi^2 + k^2)$ accounts for the hadronic size of the pion-baryon vertex ($\Lambda_\pi = 1200$ MeV). The explicit form of the thermal Lindhard functions, ϕ_α, including the Δ width has been given in Ref. [8]. For the Migdal parameters we take $g'_{\alpha\beta} = 0.8$ for $\alpha\beta = aa^{-1}bb^{-1}$ and $g'_{\alpha\beta} = 0.5$ for all others. Finally, to account for a finite pion density, we supplement Eq. (16) with a two-pion Bose factor $[1 + f^\pi(k_0) + f^\pi(q_0 - k_0)]$, which can be derived rigorously within a Matsubara formalism [11]. The contribution of the gas pions to the pion self-energy $\Sigma_\pi(k_0, k)$ has been shown to be small [11] and we can safely neglect it here.

Expressions (2)–(20) specify our model for the dilepton rate in back-to back kinematics. Since we evaluate the ρ propagator for a homogeneous gas of nucleons and Δ's in thermal equilibrium, the rate at a given temperature T of the fireball is determined by the nucleon and Δ abundances at that T. To evaluate these abundances we make use of the transport results of Li, Ko, and Brown [2]. It is found that at the initial stage the total baryon density ρ_b is $3.5\rho_0$ ($\rho_0 = 0.16$ fm^{-3}) at a temperature of 170 MeV. Including besides the nucleon and Δ all baryon resonances with masses below 1.7 GeV as well as the lowest-lying hyperons [2] and assuming chemical equilibrium, the chemical potential $\mu_N = \mu_\Delta = \cdots = 0.448$ GeV is determined by the initial baryon density. The corresponding nucleon and Δ densities are $1.0\rho_0$ each, the sum being in good agreement with the results quoted in Ref. [2]. The time evolution of the abundances can be obtained from the time dependence of the temperature given in Ref. [2]. A reasonable fit of the transport results is obtained with

$$T(t) = (T^i - T^\infty)e^{-t/\tau} + T^\infty, \quad (21)$$

with an initial temperature $T^i = 170$ MeV, $T^\infty = 115$ MeV, and $\tau = 10$ fm$/c$. For chemical equilibrium the time evolution of ρ_b can then be calculated and the result is again in agreement with the transport calculations [2]. Thus we can extract ρ_N and ρ_Δ at each time for the integration of the time history.

For a direct comparison with experiment two further points have to be considered. The first is related to the fact that the rate is evaluated only in back-to-back kinematics ($\vec{q} = 0$). Minimally this can be corrected for by using [12]

$$\frac{dN_{l^+l^-}}{d^4x d^4q}(q_0, \vec{q}) = \left(\frac{dN_{l^+l^-}}{d^4x d^4q}\right)_{\vec{q}=0} F(M, q, T),$$
$$F(M, q, T) = \frac{\exp[-(\sqrt{\vec{q}^2 + M^2} - M)/T]}{\sqrt{1 + \vec{q}^2/M^2}}, \quad (22)$$

which is exact for free pions. The second, more severe, point is the finite momentum acceptance of the detector. In the CERES S + Au experiment only dielectrons with opening angles $\Theta_{ee} > 35$ mrad and transverse momenta $p_t > 0.2$ GeV are detected. While the opening angle restriction is not severe the p_t cuts have to be included properly. To do so we statistically model e^+e^- ρ decays at given temperature and invariant mass and apply the acceptance cuts to each of the two lepton tracks [13] resulting in an acceptance function $A(M, T)$. The final dilepton yield can now be calculated by taking into account the above modifications and integrating over the time evolution of the fireball up to the freeze-out time t_f to obtain

$$\frac{dN_{e^+e^-}}{d^3x dM} = \int_0^{t_f} dt\, C(M, T(t)) \left(\frac{dN_{e^+e^-}(t)}{d^4x d^4q}\right)_{\vec{q}=0}, \quad (23)$$

where

$$C(M, T) = A(M, T) \int \frac{d^3q}{(2\pi)^3} F(M, q, T) \quad (24)$$

contains the acceptance function as well as the finite q correction to the back-to-back rate [Eq. (2)]. We take $t_f = 10$ fm$/c$, but our results are affected less than 10% when choosing $t_f = 15$ fm$/c$. Finally we fix the overall normalization in the measured rapidity window to the total yield from vacuum ρ decay obtained by Li, Ko, and Brown [2]. Figure 2 shows our final results supplemented with contributions from free Dalitz decay of π^0's, η's, and ω's (dashed-dotted line), which we extracted from Ref. [1]. We also implemented the contributions from free ω decay ($\omega \to e^+e^-$) according to Ref. [2], which, due to the small ω width, are expected to undergo only minor medium modifications [2]. The inclusion of the medium modifications of the ρ meson leads to an enhancement of the e^+e^- yield of about a factor of 3 for invariant masses around $M \sim 0.5$ GeV/c^2 (full line compared to the dotted line in Fig. 2). The main effect at work is the coupling of the rho meson to dressed pions and ΔN^{-1} excitations which was discussed in detail in Ref. [5]. Inclusion of NN^{-1}, $N\Delta^{-1}$, and $\Delta\Delta^{-1}$ excitations reduces the rho-meson peak further (~20%) by moving strength to invariant masses below 0.45 GeV/c^2.

In summary, we have presented a gauge invariant calculation of the ρ propagator at rest in a hot $\pi N \Delta$ gas taking into account the full off-shell dynamics of the intermediate pions when coupled to nucleons and Δ's. After correcting for a finite momentum of the dilepton pair as well as for the experimental acceptance of the CERES detector, our calculated yield can, to a large extent, account for the enhancement in the e^+e^- spectra observed in 200 GeV/u S + Au collisions in the CERES experiment.

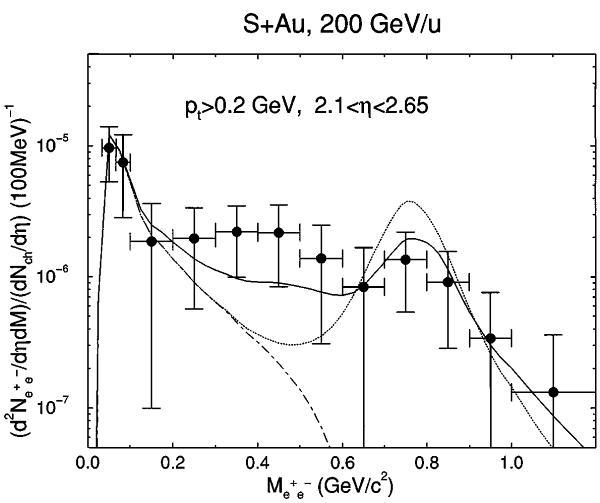

FIG. 2. Dielectron yield from free Dalitz decay (dashed-dotted line), free Dalitz + free ω+ free ρ (dotted line) and free Dalitz + free ω+ in-medium ρ decay (full line); the dots are the CERES data.

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