PHYSICAL REVIEW C, VOLUME 61, 051901(R)

Matter-induced ρ-δ mixing: A source of dileptons

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(Received 8 February 2000; published 5 April 2000)

We study the possibility of ρ-δ mixing via N-N excitations in dense nuclear matter. This mixing induces a peak in the dilepton spectra at an invariant mass equal to that of the δ. We calculate the cross section for dilepton production through the mixing process and we compare its size with that of π-π annihilation. In-medium masses and mixing angles are also calculated.

PACS number(s): 25.75.Dw, 24.10.Cn

Heavy ion physics has recently seen a considerable effort being devoted to the study of the properties of hadrons in a hot and/or dense nuclear medium. Those activities were stimulated in part by the suggestion that in the nuclear medium, the vector meson masses would drop from their values in free space and that this could be interpreted as a precursor phenomenon of chiral symmetry restoration [1]. Several attempts have been made to highlight and understand the in-medium behavior of vector mesons, both in theory and experiment [1–4]. In this respect electromagnetic signals constitute valuable probes, especially lepton pairs. This owes to the fact that the leptons couple to hadrons via vector mesons and, therefore, hadronic processes involving e^+e^- in the final channel are expected to reveal their properties in the dilepton spectra. Furthermore, the e^+e^- pairs suffer minimum final state interactions and are thus likely to bring information to the detectors essentially unscathed. Several experiments have measured, or are planning to measure, the lepton pairs produced in nucleus-nucleus collisions. They have been carried out by the Dilepton Spectrometer (DLS) at Lawrence Berkeley Laboratory (LBL) [5], and by HELIOS [6] and CERES [2] at CERN. Two new initiatives that will focus on electromagnetic probes will be PHENIX at Relativistic Heavy Ion Collider (RHIC) [7] and HADES at Gesellschaft für Schwerionenforschung (GSI) [8]. The density-dependent characteristics of vector mesons can also be highlighted through experiments performed at Thomas Jefferson National Accelerator Facility (TJNAF) [9]. The last two projects will involve measurements performed in environments where the possible modifications from vacuum properties will mostly be density driven. It is with those in mind that we have performed the theoretical estimates which we report in this Rapid Communication.

While several theoretical studies have sought to investigate the in-medium properties of the vector mesons (mainly the ρ), their possible mixing with other mesons has only started to receive attention in the context of dense baryonic matter. An exception is the case of ρ-ω [10,11]. This specific mixing can be omitted when dealing with symmetric nuclear matter, as we will here. The popularity of the ρ meson resides in the fact that in nuclear collisions a substantial contribution to the dilepton spectra comes from π-π annihilation, which proceeds through ρ as an intermediate state. This fact can also be stated as the dilepton spectrum sampling the in-medium vector meson spectral function [12–14].

We explore here the possibility of ρ-δ (or a_0 as listed in [15]) mixing via nucleon(n)-nucleon(n) excitations in nuclear matter. Such a mixing, in effect, is similar to the known ω-σ mixing [16–19]. This is a pure density-dependent effect and is forbidden in free space on account of Lorentz symmetry. We will show that such a mixing opens up a new channel for the dilepton productions and induces an additional peak in the ϕ mass region.

The interaction Lagrangian we will use can be written as

$$\mathcal{L}_{\rm int}=g_\sigma\bar{\psi}\phi_\sigma\psi+g_\delta\bar{\psi}\phi_{\delta,a}\tau^a\psi+g_{\omega NN}\bar{\psi}\gamma_\mu\psi\omega^\mu +g_\rho\left[\bar{\psi}\gamma_\mu\tau^\alpha\psi+\frac{\kappa_\rho}{2m_n}\bar{\psi}\sigma_{\mu\nu}\tau^\alpha\partial^\nu\right]\rho^\mu_\alpha, \tag{1}$$

where ψ, ϕ_σ, ϕ_δ, ρ, and ω correspond to nucleon, σ, δ, ρ, and ω fields, and τ_a is a Pauli matrix. The values used for the coupling parameters are obtained from Ref. [20].

The polarization vector through which the δ couples to ρ via the n-n loop is given by

$$\Pi_\mu(q_0,|\vec{q}|)=2ig_\delta g_\rho\int\frac{d^4k}{(2\pi)^4}{\rm Tr}[G(k)\Gamma_\mu G(k+q)], \tag{2}$$

where 2 is an isospin factor and the vertex for ρ-nn coupling is

$$\Gamma_\mu=\gamma_\mu-\frac{\kappa_\rho}{2m_n}\sigma_{\mu\nu}q^\nu. \tag{3}$$

$G(k)$ is the in-medium nucleon propagator given by [21]

$$G(k_0,|\vec{k}|)=G_F(k)+G_D(k_0,\vec{k}) \tag{4}$$

with

$$G_F(k)=\frac{(\not{k}+m_n^*)}{k^2-m_n^{*2}+i\epsilon}, \tag{5}$$

and

$$G_D(k_0,|k|)=(\not{k}+m_n^*)\frac{i\pi}{E_k^*}\delta(k_0-E_k^*)\theta(k_F-|\mathbf{k}|), \tag{6}$$

where $E_k^*=\sqrt{k^2+m_n^{*2}}$. The second term ($G_D$) deletes on-mass shell propagation of nucleons having momenta below the Fermi momentum k_F. In Eq. (4), the subscripts F and D refer to the free and density-dependent part of the propaga-

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tor. In the subsequent equations m_n^* denotes the effective nucleon mass evaluated at the mean field level [21].

With the evaluation of the trace, and after a little algebra, Eq. (2) could be cast into a suggestive form:

$$\Pi_\mu(q_0,|q|)=\frac{g_\rho g_\delta}{\pi^3}2q^2\left(2m_n^*-\frac{\kappa q^2}{2m_n}\right)\times\int_0^{k_F}\frac{d^3k}{E^*(k)}\frac{k_\mu-\frac{q_\mu}{q^2}(k\cdot q)}{q^4-4(k\cdot q)^2}. \tag{7}$$

This leads immediately to two conclusions. First, it respects the current conservation condition, *viz.*, $q^\mu\Pi_\mu=0=\Pi_\nu q^\nu$. Secondly, there are only two components which would survive after the integration over azimuthal angle. In fact this guarantees that it is only the longitudinal component of the ρ meson which couples to the scalar meson while the transverse mode remains unaltered. Furthermore, current conservation implies that out of the two nonzero components of Π_μ, only one is independent. It should be noted here that the tensor interaction, as evident from Eq. (7), inhibits the mixing.

In the presence of mixing the combined meson propagator might be written in a matrix form where the dressed propagator would no longer be block-diagonal:

$$\mathcal{D}=\mathcal{D}^0+\mathcal{D}^0\Pi\mathcal{D}. \tag{8}$$

It is to be noted that the free propagator is block-diagonal and, following the notation used in [16], has the form

$$\mathcal{D}^0=\begin{pmatrix} D^0_{\mu\nu} & 0\\ 0 & \Delta_0\end{pmatrix}. \tag{9}$$

In Eq. (9) the noninteracting propagator for the δ and ρ are given, respectively, by

$$\Delta_0(q)=\frac{1}{q^2-m_\delta^2+i\epsilon}, \tag{10}$$

$$D^0_{\mu\nu}(q)=\frac{-g_{\mu\nu}+\frac{q_\mu q_\nu}{q^2}}{q^2-m_\rho^2+i\epsilon}. \tag{11}$$

The mixing is characterized by the polarization insertion in Eq. (8), which is a 5×5 matrix containing nonblock-diagonal elements:

$$\Pi=\begin{pmatrix}\Pi^\rho_{\mu\nu}(q) & \Pi_\nu(q)\\ \Pi_\mu(q) & \Pi^\delta(q)\end{pmatrix}. \tag{12}$$

In the above expression, Π^δ and $\Pi^\rho_{\mu\nu}$ refer to the diagonal self-energies of the δ and ρ meson induced by the *n-n* polarization:

$$\Pi^\delta(q_0,|\vec{q}|)=-2ig_\delta^2\int\frac{d^4k}{(2\pi)^4}\mathrm{Tr}[G(k)G(k+q)], \tag{13}$$

$$\Pi^\rho_{\mu\nu}(q_0,|\vec{q}|)=-2ig_\rho^2\int\frac{d^4k}{(2\pi)^4}\mathrm{Tr}[G(k)\Gamma_\mu G(k+q)\tilde{\Gamma}_\nu]. \tag{14}$$

It should be mentioned that, unlike mixing, Eqs. (13) and (14) also involve a free part stemming from the $G_F(k)G_F(k)$ combination which is divergent. This therefore needs to be regularized. The regularization condition we employ is $\partial^n\Pi^F(q^2)/\partial(q^2)^n|_{m_n^*\to m,q^2=m_v^2}=0\quad(n=0,1,2,\ldots,\infty)$ [10]. For ρ the results may be found in Ref. [10], which we do not present here and for δ the free part of the self-energy is given by

$$\Pi^\delta(q^2)=\frac{3g_\delta^2}{2\pi^2}[3(m_n^{*2}-m_n^2)-4(m_n^*-m_n)m_n-(m_n^{*2}-m_n^2)]\int_0^1dx\ln\left[\frac{m_n^{*2}-x(1-x)q^2}{m_n^2}\right]-\int_0^1dx(m_n^2-x(1-x)q^2)\ln\left[\frac{m_n^{*2}-x(1-x)q^2}{m_n^2-x(1-x)q^2}\right]. \tag{15}$$

It might be worthwhile to say here that being a vector, the collective oscillations set by the ρ meson propagation through matter would have longitudinal (L) and transverse (T) components depending upon whether its spin is aligned along or perpendicular to the direction of propagation. Accordingly, with a special choice of z axis along the direction of the momenta ($\vec{q}$), one can define the longitudinal and transverse polarization as $\Pi_L=-\Pi_{00}+\Pi_{33}$ and $\Pi_T=\Pi_{11}=\Pi_{22}$, respectively [16]. To determine the collective modes, one defines the dielectric function as [16]

$$\epsilon(q_0,|\vec{q}|)=\det(1-\mathcal{D}^0\Pi)=\epsilon_T^2\times\epsilon_{\mathrm{mix}}, \tag{16}$$

where ϵ_T corresponds to two identical transverse (T) modes and ϵ_{mix} correspond to the longitudinal mode with the mixing. The latter, of course, also characterizes the mode relevant for the δ meson propagation:

$$\epsilon_T=1-d_0\Pi_T,\quad d_0=\frac{1}{q^2-m_\delta^2+i\epsilon}$$

$$\epsilon_{\mathrm{mix}}=(1-d_0\Pi_L)(1-\Delta_0\Pi_s)-\frac{q^2}{|\vec{q}|^2}\Delta_0 d_0(\Pi_0)^2. \tag{17}$$

The zeros of the dielectric functions define the dispersion relation for meson propagation. Figure 1 shows the relevant dispersion curves with and without mixing at density $\rho=2.5\rho_0$. As only the L mode mixes with the scalar mode, we do not consider the T mode. The latter in fact is the same as

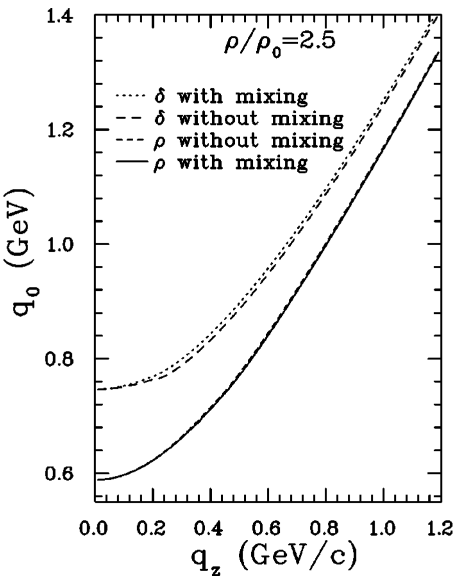

FIG. 1. The dispersion curve for ρ and δ meson with and without mixing at $\rho=2.5\rho_0$.

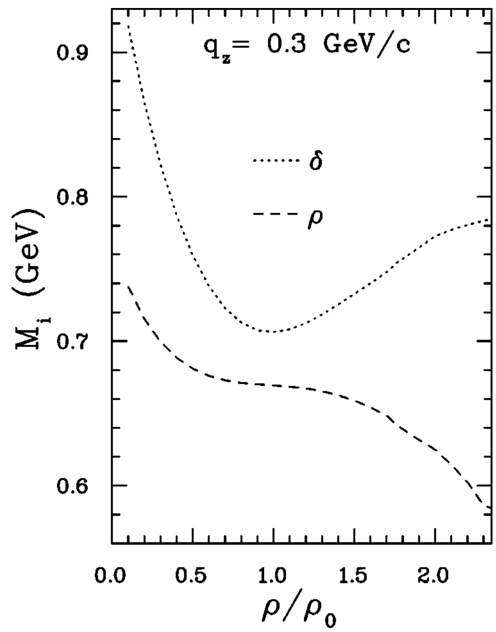

FIG. 2. The ''invariant mass'' as a function of the relative nuclear density.

presented in Ref. [10] for ρ meson. The effect of mixing on the pole masses, as evident from Fig. 1, are found to be very small. However, the mixing could be large when the mesons involved go off-shell. It should be noted that the modes with mixing move away from each other compared to what one obtains without mixing. This can be understood in terms of ''level-level'' repulsion driven by the off-diagonal terms of the dressed propagator [18].

The ''invariant mass'' ($M_i=\sqrt{q_{0(i)}^2-|\vec{q}_i|^2}$) ($i=\rho,\delta$) is determined by finding the solutions (q_0) of the equation $\epsilon_{\rm mix}=0$ for a fixed value of $|\vec{q}|$. Figure 2 shows the dependence of M_i on nuclear densities for $|\vec{q}|=0.3$. It is evident that the difference of the ''invariant masses'' first decreases with density reaching a minimum, and then again starts increasing. This behavior arises from the nonmonotonic density dependence of the polarization functions. This trend is the same for the whole relevant domain of $|\vec{q}|$.

To calculate the mixing angle, one diagonalizes the mass matrix [10] with the mixing and obtains

$$\theta_{\rm mix}=\frac{1}{2}\arctan\left(\frac{2\Pi_{\rm mix}^{\rho\delta}}{m_\delta^2-m_\rho^2-\Pi_L^\rho+\Pi^\delta}\right). \tag{18}$$

In Eq. (18) $\Pi_{\rm mix}^{\rho\delta}=M_i/|\vec{q}|\Pi_0$, which increases with density. Π_0 is the zero component of Eq. (7). Equation (18) clearly

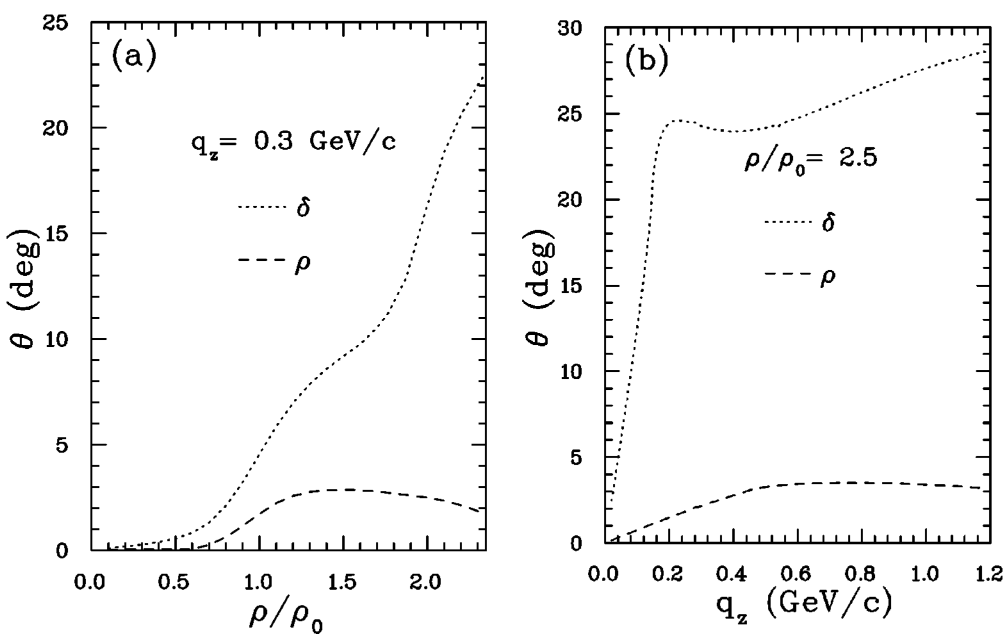

FIG. 3. The mixing angle as a function of the relative nuclear density (a) and momentum (b).

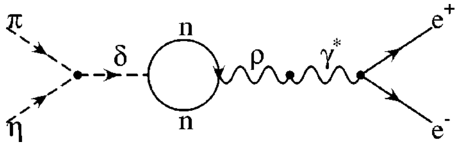

FIG. 4. The Feynman diagram for the process $\pi+\eta\to e^+ + e^-$.

shows that the mixing angle depends not only on the mixing amplitude (Π_0) but also on the ''energy denominator.'' The latter, as seen in Fig. 2, first decreases as a function of density then again shows an increase characterizing the density dependence of the mixing angle as presented in Fig. 3. It should be noted here that the mixing angle is a function of q_0 and $|\vec{q}|$. Therefore estimated values presented here correspond to the two poles, δ and ρ, determined by the zeros of ϵ_{mix} [22]. The mixing angle in Fig. 3 corresponds to $|\vec{q}| = 0.3$ GeV/c. The momentum dependence for a density 2.5 times higher than the normal nuclear matter density is shown in the right panel of the same figure. This shows that for momenta beyond $|\vec{q}|\approx 0.2$ GeV/c the mixing is quite appreciable, which affects the dilepton yield substantially as shown later. It should also be noted that the mixing angle vanishes at $|\vec{q}| = 0$ or at $\rho = 0$, as it should.

The ρ-δ mixing opens a new channel, *viz.*, $\pi+\eta\to e^+ + e^-$ in dense nuclear matter through $n-n$ excitations. The Feynman diagram of the process is depicted in Fig. 4.

The cross section for this process might be expressed in terms of the mixing amplitude (Π_0)

$$\sigma_{\pi\eta\to e^+e^-} = \frac{4\pi\alpha^2}{3q_z^2 M}\frac{g_{\delta\pi\eta}^2}{g_\rho^2}\frac{m_\rho^4}{(M^2-m_\rho^2)^2+m_\rho^2\Gamma_\rho^2(M)} \times \frac{1}{(M^2-m_\delta^2)^2+m_\delta^2\Gamma_\delta^2(M)}\frac{1}{\sqrt{M^2-4m_\pi^2}}|\Pi_0|^2. \tag{19}$$

For the decay widths we consider the invariant mass dependence as presented below:

$$\Gamma_\rho(M) = \frac{g_{\rho\pi\pi}^2}{6\pi}\frac{\left(\frac{M^2}{4}-m_\pi^2\right)^{\frac{3}{2}}}{M^2}, \tag{20}$$

and

$$\Gamma_\delta(M) = \frac{g_{\delta\pi\eta}^2}{16\pi}\frac{\sqrt{(M^2-(m_\pi+m_\eta)^2)(M^2-(m_\pi-m_\eta)^2)}}{M^3}. \tag{21}$$

To describe the $\pi\delta\eta$ vertex we use

$$\mathcal{L}_{\delta\pi\eta} = f_{\delta\pi\eta}\frac{m_\delta^2-m_\eta^2}{m_\pi}\phi_\eta\vec{\phi}_\pi\cdot\vec{\phi}_\delta, \tag{22}$$

where for later convenience we define $g_{\pi\delta\eta} = f_{\delta\pi\eta}(m_\delta^2 - m_\eta^2)/m_\pi$.

Of course, there is an uncertainty involved with the coupling parameter $f_{\delta\pi\eta}$ as discussed in Refs. [23,15]. This arises from the fact that δ (or a_0) lies close to the opening of the $K\bar{K}$ channel leading to a cusplike behavior in the resonant amplitude, therefore a naive Breit-Wigner form for the decay width is inadequate. Furthermore, as mentioned before, there is also uncertainty involved with the δNN coupling, which renders the precise extraction of δ- π-η coupling even more difficult [23]. We take a value for $f_{\delta\eta\pi}$=0.44 from Ref. [23] which gives $\Gamma_{\delta\to\pi\eta}(m_\delta)$=59 MeV, while the experimental vacuum width of δ is between 50−100 MeV [15].

One can notice in Fig. 5 that the process, $\pi+\eta\to e^+ + e^-$, at densities higher than $\boldsymbol{\rho}_0$, not only enhances the overall production of lepton pairs but also induces an additional peak near the ϕ mass region. The contribution at the δ mass is comparable to that of $\pi+\pi\to e^+ + e^-$ near the ρ peak, for densities higher than $\boldsymbol{\rho_0}$. Figure 5 also shows that

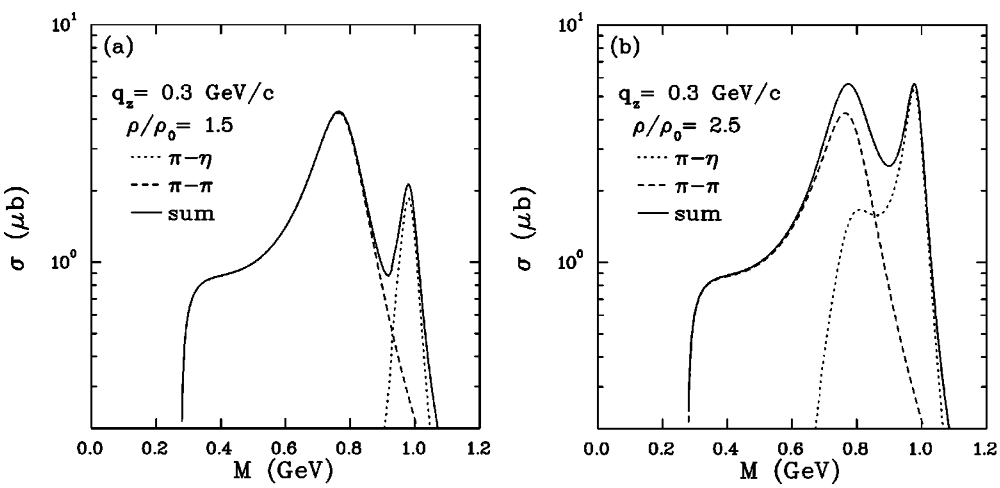

FIG. 5. Dilepton spectrum induced by $\pi+\pi\to e^+ + e^-$ and $\pi+\eta\to e^+ + e^-$ considering matter-induced $\rho-\delta$ mixing. (a) ρ=1.5ρ_0. (b) ρ=2.5ρ_0.

as the density goes even higher the dilepton yield arising out of the mixing also increases further. The cross section increases with increasing momenta of the mesons in keeping with the mixing angle as shown in Fig. 3.

We have highlighted the possibility of ρ-δ mixing in dense nuclear matter. We observe the appearance of an additional peak at a dilepton invariant mass that corresponds to that of the δ. With sufficient experimental resolution, this effect could be observable, probably not as an individual peak, because of the δ's vacuum width which is already not small, but more realistically as a shoulder in the ϕ spectrum. This feature is then exclusively density dependent. Our aim here was to establish the existence of the signal. Our calculation can, and will be, improved upon: further studies are in progress to assess finite temperature effects and to self-consistently incorporate the necessary many-body machinery. For example, the characteristics of the ρ can be modified in the nuclear medium [24] and the in-medium behavior of the δ needs to be addressed. We have verified that the inclusion of hadronic form factors does not change the conclusions we reach in this work. Detailed results will be presented elsewhere.

This work was supported in part by the Natural Sciences and Engineering Research Council of Canada and in part by the Fonds FCAR of the Québec Government.

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