The properties of vector mesons in nuclear matter are discussed. Results for the momentum dependence of the $\rho$-meson self energy in matter due to $\rho N$ interactions are presented, and consequences for the production rate of $e^+ e^-$ pairs in hot and dense matter are discussed. I also examine the constraints imposed by elementary processes on the widths of $\rho$ and $\omega$ mesons in nuclear matter, and conclude that in recent models the in-medium widths of these mesons are overestimated.

1 Introduction

In 1991 Gerry Brown and Mannque Rho proposed a universal scaling law for effective hadron masses in matter, which connects the hadron masses with the quark condensate, an order parameter for the spontaneous breaking of chiral symmetry. According to the scaling law, often referred to as Brown-Rho or BR scaling, the in-medium hadron masses, except for those of the pseudo-scalar mesons, are reduced when the quark condensate is reduced, i.e., when the chiral symmetry is restored. Thus, if BR scaling is correct, it opens the possibility to experimentally explore the restoration of chiral symmetry in dense and hot matter, provided one can identify experimental signatures for the modification of the effective masses of hadrons. This idea has triggered a lively discussion on the properties of hadrons in matter.

The electromagnetic decay of the vector mesons into $e^+ e^-$ and $\mu^+ \mu^-$ pairs makes them particularly well suited for exploring the conditions in dense and hot matter in nuclear collisions. The lepton pairs provide virtually undistorted information on the mass distribution of the vector mesons in the medium. Because of this, the discussion has concentrated on the in-medium properties of the $\rho$ and $\omega$ mesons and the corresponding signatures in the lepton-pair spectrum.

The universal scaling conjecture is supported by the QCD sum rule calculations of Hatsuda and Lee, who find a strong reduction of the $\rho$- and $\omega$-meson masses in nuclear matter. However, recently the results of Hatsuda and Lee...
were questioned in the work of Klingl, Kaiser and Weise. Klingl et al. find that the vector meson widths, which were ignored in the work of Hatsuda and Lee, play an important role in the QCD sum rules and that the QCD sum rule for the $\rho$ meson is satisfied by a model, where its width is strongly enhanced in nuclear matter, while its energy remains almost unchanged. For the $\omega$ meson they find a strong downward shift of the energy and a relatively large enhancement of its width. The tradeoff between the vector-meson energy and width in the QCD sum rules was then explored in more detail by Leupold, Peters and Mosel.

The lepton-pair spectrum in nucleus-nucleus collisions at SPS energies exhibits a low-mass enhancement compared to proton-proton and proton-nucleus collisions. A quantitative interpretation of the lepton-pair data can be obtained within a scenario, where the effective vector-meson masses are reduced in a hadronic environment. On the other hand, attempts to interpret the low-mass enhancement of lepton pairs in terms of many-body effects also yield good agreement with the data. In these calculations the broadening of the $\rho$ meson in nuclear matter due to the interactions of its pion cloud with the medium and the momentum dependence of the $\rho$-meson self energy due to the coupling with baryon-resonance–nucleon-hole states are taken into account. At least superficially, this mechanism seems to be completely decoupled from the BR-scaling scenario.

In this talk I discuss the many-body effects in the light of recent theoretical developments. I also examine the constraints on the imaginary part of the vector meson self energy in nuclear matter that can be derived from elementary reactions, like $\pi N \rightarrow \rho/\omega N$.

2 The pion cloud

The $\rho$ and $\omega$ mesons couple strongly to two and three pion states, respectively. Consequently, both are surrounded by a pion cloud. Since pions interact strongly with nucleons, the pion clouds of the vector mesons are modified in nuclear matter. For a $\rho$ meson at rest in nuclear matter this effect was explored several years ago by Herrmann et al., Chanfray and Schuck and by Asakawa et al. In these calculations conservation of the isospin current, which acts as a source for the $\rho$-meson field, was a guiding principle. The main effect considered was the dressing of the pion by the excitation of $\Delta(1232)$-hole states. In spite of differences in the details of the models, all groups came to similar conclusions, namely that the $\rho$-meson width is enhanced in matter, but that its energy does not change appreciably. Furthermore, at high densities ($\rho_N = 2 - 3\rho_0$), a new branch appears at an energy of roughly 450 MeV. This
branch corresponds to the decay of the $\rho$ meson into a pion and a $\Delta$-hole state. In Fig. 1 I show the resulting $\rho$-meson spectral function of ref. 11 and in Fig. 2 the corresponding lepton pair production rate from $\pi^+\pi^-$ annihilation in hadronic matter at a temperature of $T = 100$ MeV, and different densities. Note that the broadening of the $\rho$ and the new branch leads to an enhancement of the production rate for low-mass lepton pairs.

Recently Klingl et al. constructed a vector-meson–baryon effective interaction of finite range based on a chirally symmetric Lagrangian. In this model the vector-meson–nucleon interaction is to a large extent mediated by the meson cloud of the vector mesons, much like in the earlier work 11, 12, 13. The resulting in-medium spectral functions satisfy the QCD sum rules. In agreement with the earlier calculations, Klingl et al. find that a $\rho$ meson at rest in nuclear matter is broadened, but that its energy remains almost unchanged. On the other hand, the energy of the $\omega$ meson is shifted down appreciably and its width is enhanced substantially. Finally, for the $\phi$ meson they find almost no change of its energy and a moderate enhancement of its width. In this work the vector-meson–nucleon scattering amplitude is renormalized in the limit corresponding to Compton scattering of long-wavelength photons. Thus, the model involves an extrapolation from the photon point $q^2 = 0$ to

Figure 1: The spectral function of a $\rho$ meson at rest in nuclear matter at density $\rho = 0$ (full), $\rho_0$ (dotted), $2\rho_0$ (dashed) and $3\rho_0$ (dash-dotted line) (from ref. 11).
Figure 2: The production rate for lepton pairs in nuclear matter at $T = 100$ MeV corresponding to the spectral function shown in Fig. 1. The densities are $\rho = 0$ (full), $\rho_0$ (dotted), $2\rho_0$ (dashed) and $3\rho_0$ (dash-dotted line) (from ref. 11).

$q^2 = m_V^2$. As we shall see below, this extrapolation, which is governed by the form factors of the model, turns out to be hazardous.

The imaginary parts of the $\rho$-nucleon and $\omega$-nucleon T-matrices obtained by Klingl et al. are shown in Fig. 3 and 4. They define a complex scattering length through

$$ a_{VN} = \frac{M_N}{4\pi(M_N + m_V)} T_{VN}(\omega = m_V). \tag{1} $$

The imaginary parts of the scattering lengths are related to the existence of inelastic channels, the $\pi N$ and $\pi\pi N$ channels, which are open at the vector-meson–nucleon threshold. The complex $\rho N$ and $\omega N$ scattering lengths quoted by Klingl et al. are $a_{\rho N \rightarrow \rho N} = (0.04 + i 1.62)$ fm and $a_{\omega N \rightarrow \omega N} = (3.34 + i 2.1)$ fm, respectively. As discussed in section 4, one can relate the scattering lengths to the vector-meson self energies in nuclear matter at low density. Thus, the imaginary parts of the scattering lengths correspond to very large in-medium enhancements of the $\rho$ and $\omega$ widths. The values of the scattering lengths imply that on the free mass shell $\Delta \Gamma_\rho \simeq 300$ MeV and $\Delta \Gamma_\omega \simeq 390$ MeV.
at normal nuclear matter density. The real part of the $\omega N$ scattering length indicates that the $\omega$-meson potential in nuclear matter is strongly attractive, while the small real part of the $\rho N$ scattering length corresponds to an almost vanishing shift of the $\rho$-meson energy in a nuclear medium. If the energy of the $\omega$ meson is reduced in matter due to the attractive potential, its in-medium width is of course smaller than the value quoted above. Both the real and imaginary parts of the $\phi N$ scattering length are small in this model.

Using Eq. (1) and the results shown in Fig. 3 and 4 one can extract the contribution from the $\pi N$ channel to the imaginary parts of the scattering lengths. I find $\text{Im} a^{(\pi N)}_{\rho N \rightarrow \rho N} \simeq 0.35$ fm and $\text{Im} a^{(\pi N)}_{\omega N \rightarrow \omega N} \simeq 1.1$ fm, which correspond to $\Delta \Gamma^{(\pi N)}_{\rho} = 65$ MeV and $\Delta \Gamma^{(\pi N)}_{\omega} = 200$ MeV on the free mass shell (at $\rho_N = \rho_0$). In this model, the $\pi N$ channel is responsible for the major part of the in-medium width of the $\omega$ meson at masses below the free one, while for the $\rho$ meson the $\pi \Delta$ channel dominates. As I discuss below, the $\pi N$ contributions can be related to the experimentally accessible reactions $\pi N \rightarrow VN$. Consequently, such reactions provide constraints on models for the vector-meson self energy in nuclear matter at low baryon densities.
Figure 4: Same as Fig. 3 for the ω-nucleon T-matrix.

The fact that the QCD sum rules are satisfied by the in-medium ρ-meson spectral function of Hatsuda and Lee and by the one of Klingl et al. indicates that the ρ-meson width plays an important role in this context. Indeed, Leupold et al. find that there is a trade-off between the in-medium ρ-meson energy and its in-medium width. For a small ρ width, the QCD sum rule requires an appreciable downward shift of the ρ energy, while for a large width the sum rule is satisfied with a small energy shift. Note that these results are obtained with the vacuum saturation assumption for the 4-quark condensate. Thus, the uncertainty connected with this approximation remains.

3 ρN interactions at finite \( \vec{q} \)

So far I have discussed only vector mesons at rest in nuclear matter. In a heavy-ion collision, the leptonic decay of such a vector meson gives rise to a lepton pair of vanishing total momentum in the rest frame of the participants. However, so far the experimental invariant mass distributions are integrated over the total momentum of the lepton pair. Consequently, in a theoretical analysis of the lepton-pair spectrum in nucleus-nucleus collisions the properties of vector mesons moving with respect to the medium must be considered. In a
recent paper H.J. Pirner and I. estimated the interactions of $\rho$ mesons with nucleons at finite momenta and explored the consequences for the production of lepton pairs in relativistic nucleus-nucleus collisions.

As discussed in section 4 the $\rho$-meson self energy in low-density nuclear matter can be constructed from the $\rho$-nucleon scattering amplitude $f_{\rho N}$. In this work we approximate the scattering amplitude with baryon resonance contributions. There are two resonances in the mass range of interest ($m \approx m_N + m_\rho$) which couple strongly to the $\rho N$ channel, namely $N^*(1720)$ and $\Delta(1905)$. The $N^*(1720)$, which is the more important one, decays into a $\rho$ meson and a nucleon in a relative p-wave with more than 70% probability, while the $\Delta(1905)$ decays into the $\rho N$ channel with a branching ratio above 60% in a relative p- or f-wave. Thus, the corresponding $\rho$ self energy is momentum dependent.

In order to describe the coupling of a virtual photon to hadrons we adopt the formulation of the vector meson dominance (VMD) model due to Kroll, Lee and Zumino. This formulation of VMD has the advantage of being explicitly gauge invariant, and it allows us to adjust the $\gamma N$ and $\rho N$ partial widths independently.

The free parameters of the model are the $N^*N\gamma$ and $N^*N\rho$ transition matrix elements together with the corresponding ones for the $\Delta(1905)$. These are fixed by fitting the partial widths for the resonance decays into the $\gamma N$ and $\rho N$ channels. In the static approximation, the corresponding $\rho$-meson self energy in nuclear matter is, to lowest order in density, given by

$$\Sigma_{\rho}(\omega, \vec{q}) - \Sigma_{\rho}^{(0)}(\omega, \vec{q}) = 4 f_{N^*N\rho}^2 (\frac{m_N^2}{m_\rho^2}) F(q^2) q^2 \rho_B \left( \frac{\varepsilon_{N^*} - m_N}{\omega^2 - (\varepsilon_{N^*} - m_N)^2} \right) + (N^* \rightarrow \Delta),$$

where

$$\varepsilon_{q^2}^{N^*} = \sqrt{q^2 + m_{N^*}^2} - \frac{i}{2} \Gamma_{N^*},$$

$F(q^2) = \Lambda^2/(\Lambda^2 + q^2)$ is a form factor with $\Lambda = 1.5$ GeV and $\Sigma_{\rho}^{(0)}(\omega, \vec{q})$ denotes the $\rho$-meson self energy in vacuum. In Eq. $\Gamma_{N^*}$ is the full width of the resonance, modified by the phase space appropriate for a resonance embedded in the $\rho$-meson self energy. This is important for the correct threshold behaviour of the in-medium self energy. Note that in the static approximation the self energy vanishes for a $\rho$ meson at rest in nuclear matter.

The resulting spectral function is shown in Fig. at $|\vec{q}| = 750$ MeV and $\rho_B = 2\rho_0$ together with the spectral function of a $\rho$ meson in vacuum. Note that at finite momenta strength is moved down to energies below the $\rho$-meson
Figure 5: The $\rho$-meson spectral function at $|\vec{q}| = 750$ MeV in vacuum (full line) and in nuclear matter at $\rho_B = 2\rho_0$ (from ref. 14). The arrows show the positions of the unperturbed levels in the zero width limit.

peak in vacuum. Due to the energy dependence of the self energy and the large widths of all particles involved, the spectral strength is quite fragmented at large momenta. The peaked structure at low energies, which one can identify with a quasiparticle, carries only about 20% of the strength at $q = 750$ MeV. Nevertheless, it may have interesting consequences for the lepton-pair spectrum in nucleus-nucleus collisions.

In Fig. 6 the resulting production rate for $e^+e^-$ pairs is shown for various densities, at a temperature of $T = 140$ MeV. Also shown is the rate due to $\pi^+\pi^-$ annihilation computed without any medium modifications. The rate includes an integral over all pair momenta, but no corrections for experimental acceptance. The p-wave interaction leads to a strong enhancement of the population of lepton pairs with invariant masses around 300-500 MeV. In the mass region of interest, the pion annihilation and the $N^*(1720)$ contributions dominate. Because the p-wave interactions lead to a characteristic momentum dependence of the lepton-pair spectrum, an analysis of the data in terms of total momentum and invariant mass may shed light on the mechanism responsible for the low-mass enhancement.

This calculation was recently extended by Peters et al., who include all baryon resonances of a mass below 1.9 GeV and attempt a self consistent
treatment of the $\rho$-meson spectral function. In the resulting spectral function, the $\rho$-meson quasi-particle has acquired a very large width. Similar results have been obtained by Rapp et al. \cite{9,10}, who also compute the corresponding lepton-pair spectrum for $S + Au$ and $Pb + Au$ collisions at CERN energies, including the broadening of the $\rho$ meson in a nuclear medium discussed in the previous section and the p-wave interactions discussed above. This model provides a quantitative description of the data of the CERES collaboration.

4 Constraints from elementary processes

The low-density theorem states that the self energy of e.g. a vector meson $V$ in nuclear matter is given by

$$\Sigma_V(\rho_N) = -4\pi \langle 1 + \frac{m_V}{m_N} \rangle \langle f_{VN} \rangle \rho_N + \ldots,$$

(4)

where $m_V$ is the mass of the vector meson, $m_N$ that of the nucleon, $\rho_N$ the nucleon density and $\langle f_{VN} \rangle$ denotes the $VN$ forward scattering amplitude $f_{VN}$, appropriately averaged over the nucleon Fermi sea. Note that when there is
a narrow resonance close to threshold, the validity of the low-density theorem may be limited to very low densities (see e.g. ref.\(\text{20}\)). For the vector mesons \(\rho, \omega\) and \(\phi\) the elastic scattering amplitudes have to be extracted indirectly, from e.g. vector-meson production experiments, like \(\gamma N \to V N\) and \(\pi N \to V N\).

In ref.\(\text{21}\) a simple model for the photo-induced production of \(\rho\) and \(\omega\) mesons was studied. Using the low density theorem and vector meson dominance to extrapolate the amplitude from the photon point (\(q^2 = 0\)) to an on-shell \(\rho\) meson (\(q^2 = m_{\rho}^2\)), one finds that the data are consistent with a strongly attractive \(\rho\)-meson self energy in nuclear matter. However, the extrapolation over a wide range in mass introduces a strong model dependence. Although it may be possible to eliminate this model dependence to some extent\(\text{22}\), the safest approach is clearly to constrain the \(V N\) scattering amplitude only with data in the relevant kinematic range. As an example, I shall discuss the implications of the data on pion-induced vector-meson production for the in-medium width of \(\rho\) and \(\omega\) mesons.

The spin-averaged cross section for the reaction \(\pi^- p \to \omega n\) is given by

\[
\sigma_{\pi^- p \to \omega n} = \frac{1}{2} \sum_{\text{spins}} \int |f_{\pi^- p \to \omega n}|^2 \frac{k_\omega}{k_\pi} d\Omega, \quad (5)
\]

where \(k_\omega\) and \(k_\pi\) are the c.m. momenta in the \(\omega N\) and \(\pi N\) channels. Detailed balance implies a relation between cross sections for time reversed reactions\(\text{23}\), which for the case at hand means that

\[
3k_\omega^2 \sigma_{\omega n \to \pi^- p} = k_\pi^2 \sigma_{\pi^- p \to \omega n}. \quad (6)
\]

The factor 3 is the ratio of the spin degeneracies in the \(\omega N\) and \(\pi N\) channels. Furthermore, it follows from (two-body) unitarity that

\[
\text{Im} f_{\omega n \to \omega n}(\theta = 0) = \sum_\nu k_\nu \int |f_{\omega n \to \nu}|^2 \frac{d\Omega}{4\pi}, \quad (7)
\]

where \(\text{Im} f_{\omega n \to \omega n}(\theta = 0)\) is the imaginary part of the \(\omega\)-nucleon forward scattering amplitude and the sum runs over all possible two-body final states \(|\nu\rangle\).

The contribution of the \(\pi^- p\) channel is given by

\[
\frac{4\pi}{k_\pi} \text{Im} f_{\omega n \to \omega n}^{(\pi^- p)}(\theta = 0) = \sum_{\text{spins}} \int |f_{\omega n \to \pi^- p}|^2 d\Omega, \quad (8)
\]

where the sum over spins applies only to the final state \(|\pi^- p\rangle\). Thus, using unitarity and detailed balance, one finds
where \( \tilde{f} \) denotes the spin-averaged scattering amplitude. Close to the \( \omega n \) threshold, the scattering amplitude can be expanded in powers of the relative momentum in the \( \omega n \) channel \( k_\omega \). An excellent fit to the data from threshold up to \( s = 3.1 \) GeV\(^2\) is obtained with \( \text{Im} \tilde{f}^{(\pi^- p)}_{\omega n \to \omega n}(\theta = 0) = a + bk_\omega^2 + ck_\omega^4 \), where \( a = 0.013 \) fm, \( b = 0.10 \) fm\(^3\) and \( c = -0.08 \) fm\(^5\) (see Fig. 7). The coefficient \( a \) is the imaginary part of the scattering length, while \( b \) may be due to a p-wave contribution or a generalized effective range term. Note that if the \( k \)-dependent terms are dropped, i.e., if only the scattering length is kept, the energy dependence of the cross section is too weak, and one cannot describe the data at energies beyond the first two data points.

The contribution of the \( \pi \)-nucleon channel to the width of the \( \omega \) meson at rest in nuclear matter can now be obtained by using the low-density theo-

\[ \sigma_{\pi^- p \to \omega n} = 12\pi \frac{k_{\omega}}{k_\pi^2} \text{Im} f^{(\pi^- p)}_{\omega n \to \omega n}(\theta = 0), \]
\[ \Delta \Gamma_\omega = 4\pi(1 + \frac{m_\omega}{m_N}) \frac{3}{2} \frac{\langle \text{Im} f_{\omega_n \rightarrow \omega_n}^{(\pi^-p)} \rangle \rho_N}{m_\omega} \]  

(10)

The factor \( \frac{3}{2} \) accounts for the \( \pi^0n \) channel. At nuclear matter density (\( \rho_N = 0.16 \text{ fm}^{-3} \)) the average over the Fermi sea yields an effective scattering amplitude \( \langle \text{Im} f_{\omega_n \rightarrow \omega_n}^{(\pi^-p)} \rangle = 0.031 \text{ fm} \). This implies that the width of the \( \omega \) meson in nuclear matter is increased by 9 MeV due to the \( \pi^-nucleon \) channel.

A comparison of the imaginary part of the empirical scattering length \( (1.5 \times 0.02 \text{ fm} = 0.03 \text{ fm} \) or the effective scattering amplitude \( (1.5 \times 0.034 \text{ fm} = 0.05 \text{ with that extracted from the calculation of Klingl et al.} \) in section 2 (1.1 fm) reveals a very large discrepancy. In this model the imaginary part of the \( \omega N \) scattering length and the in-medium width of the \( \omega \) meson are overestimated by more than an order of magnitude! This is a striking illustration of how dangerous an extrapolation over a wide range in energy can be, without guidance from experiment.

For the \( \rho \) meson the situation is a bit more complicated. First of all the experimentally accessible \( \pi N \) channel is subdominant. Second, both isospin 1/2 and 3/2 are allowed. Thus, three independent reactions are needed to pin down the amplitudes of the two isospin channels and their relative phase. For the \( \omega \) meson the situation is simpler, since only isospin 1/2 is allowed. The available data on the reactions \( \pi^-p \rightarrow \rho^0n, \pi^+p \rightarrow \rho^+p \) and \( \pi^-p \rightarrow \rho^-p \) up to \( s = 4 \text{ GeV}^2 \) are shown in Fig. \[ \text{These cross sections involve linearly independent combinations of the isospin amplitudes and would, if measured down to threshold, be sufficient to determine the amplitudes. The enhancement of the \( \rho \)-meson width in nuclear matter is then given by}

\[ \Delta \Gamma_\rho = \frac{4\pi}{2m_\rho}(1 + \frac{m_\rho}{m_N}) \left( \langle \text{Im} f_{\rho^0n \rightarrow \rho^0n}^{(\pi^-p)} \rangle + \langle \text{Im} f_{\rho^+p \rightarrow \rho^+p}^{(\pi^+p)} \rangle \right) \]

(11)

Unfortunately only the first one, \( \pi^-p \rightarrow \rho^0n \), is well known close to threshold. For the second one, \( \pi^+p \rightarrow \rho^+p \), only two points are available in the relevant energy range, while the data on \( \pi^-p \rightarrow \rho^-p \) start only at \( s > 4 \text{ GeV}^2 \). Thus, in order to extract numbers, we must make assumptions on the cross section.

\[ \text{Note that the imaginary part of the } \omega N \text{ scattering amplitude extracted from the data is surprisingly small. The corresponding amplitude for } \eta N \text{ scattering, extracted from the } \pi^-p \rightarrow \eta n \text{ data, is about an order of magnitude larger: } \text{Im} a_{\eta n \rightarrow \eta n}^{(\pi^-p)} = 0.2 - 0.3 \text{ fm (see e.g. ref. 25, 26).} \]
Figure 8: The data for the reactions $\pi^- p \rightarrow \rho^0 n$ (squares), $\pi^+ p \rightarrow \rho^+ p$ (circles) and $\pi^- p \rightarrow \rho^- p$ (triangles) (ref. 24).

for the last reaction close to threshold. Clearly data on $\pi^+ p \rightarrow \rho^+ p$ and $\pi^- p \rightarrow \rho^- p$ close to threshold would be very useful.

Because the $\rho$-meson width is large, the detailed-balance relation is more complicated ref. 27. One must take into account the fact that at a given c.m. energy $\sqrt{s}$ part of the $\rho$-meson spectral strength may be energetically unavailable. Thus, the cross section for production of $\rho$ mesons is of the form

$$\sigma_{\pi^- p \rightarrow \rho^0 n} = 12\pi \int_{2m_\pi}^{\sqrt{s} - m_N} \frac{k_\rho}{k^2_\pi} \text{Im} \tilde{f}_{\rho^0 n \rightarrow \rho^0 n} A_\rho(m^2) \frac{d m^2}{\pi}. \quad (12)$$

Here $A_\rho(m^2)$ is the $\rho$-meson spectral function, and $k^2_\rho = ((s - m^2 - m^2_N)^2 - 4m^2_N m^2)/4s$. For a narrow resonance, where the spectral function can be approximated by a delta function $A_R(m^2) \simeq \delta(m^2 - m^2_R)$, Eq. (12) reduces to
Figure 9: The data for the reaction $\pi^- p \to \rho^0 n$ near threshold and the fit described in the text.

the standard form (9). A good fit to the $\pi^- p \to \rho^0 n$ data near threshold is obtained with $\text{Im} \bar{f}_{\rho^0 n \to \rho^0 n} = a + bk^2$, where $a = 0.021$ fm and $b = 0.006$ fm$^3$ (see Fig. 9). By averaging over the Fermi sea, I find an effective scattering amplitude $\langle \text{Im} \bar{f}_{\rho^0 n \to \rho^0 n} \rangle = 0.022$ fm.

Similarly, a fit to the low-energy $\pi^+ p \to \rho^+ p$ data can be obtained. Because there are only two data points in the energy range of interest, it makes no sense to extract a momentum dependent scattering amplitude. A reasonable fit is obtained with $\text{Im} \bar{f}_{\rho^+ p \to \rho^+ p} = 0.06 - 0.08$ fm.

One can use the isospin decomposition of the $\pi N \to \rho N$ scattering amplitude and the two measured cross sections to set limits on the magnitude of the unknown $\pi^- p \to \rho^- p$ amplitude $^4$. The absolute value of the isospin $3/2$ $\pi N \to \rho N$ amplitude is fixed by the $\pi^+ p \to \rho^+ p$ data, while the $\pi^- p \to \rho^0 n$ data give a relation between the isospin $1/2$ amplitude and the relative phase of the amplitudes. This relation implies limits on the magnitude of the isospin

$^4$In this argument I neglect the momentum dependence of the amplitudes, which for the reaction $\pi^- p \to \rho^0 n$ gives rise to at most a $10\%$ correction at the momenta of interest here.
1/2 amplitude, and consequently also constraints on the $\pi^-p \rightarrow \rho^-p$ amplitude. One thus finds that the absolute value of the isospin 1/2 amplitude must be less than 0.35 fm, which implies for the $\langle \text{Im} f^{(\pi^-p)}_{\rho^-p \rightarrow \rho^-p} \rangle < 0.25$ fm and that the effective “scattering length” due to the $\pi N$ channel fulfills

$$\text{Im} a^{(\pi N)}_{\text{eff}} = \frac{\langle \text{Im} f^{(\pi^-p)}_{\rho^0 n \rightarrow \rho^0 n} \rangle + \langle \text{Im} f^{(\pi^+p)}_{\rho^-p \rightarrow \rho^-p} \rangle + \langle \text{Im} f^{(\pi^+p)}_{\rho^-p \rightarrow \rho^-p} \rangle}{2} < 0.18 \text{ fm.} \quad (13)$$

This constraint is not very restrictive, since it allows the unknown scattering amplitude near threshold to be an order of magnitude larger than that for the reaction $\pi^-p \rightarrow \rho^0 n$. To proceed further one must make an assumption for $\langle \text{Im} f^{(\pi^-p)}_{\rho^-p \rightarrow \rho^-p} \rangle$. At energies, where all reaction channels are measured ($s > 4 \text{ GeV}^2$), the cross sections agree within a factor $\sim 2$. Hence, a reasonable guess is that this amplitude is not very different from the ones extracted for the other reactions. If one assumes that the unknown scattering amplitude is between those obtained for the reactions $\pi^-p \rightarrow \rho^0 n$ and $\pi^+p \rightarrow \rho^+p$, one finds that the $\pi N$ contribution to $\text{Im} a^{(\pi N)}_{\text{eff}} = 0.05 - 0.09$ fm. This corresponds to an enhancement of the $\rho$-meson width of $\Delta \Gamma^{(\pi N)}_{\rho} = 9 - 17 \text{ MeV}$ at nuclear matter density. A comparison with the value extracted for the model of Klingl et al. (0.35 fm), reveals a discrepancy of a factor $4 - 7$. In order to reproduce their value, the unknown scattering amplitude would have to be $\langle \text{Im} f^{(\pi^-p)}_{\rho^-p \rightarrow \rho^-p} \rangle = 0.61 \text{ fm}$, which clearly violates the upper limit derived above.

In ref. 3 it is found that the imaginary part of the $\rho N$ T-matrix is dominated by the experimentally not accessible $\Delta \pi$ channel. If this is the case, one cannot draw any firm conclusions on the in-medium width of the $\rho$ meson. Nevertheless, one can make qualitative statements. One does not expect the matrix element for $\rho N \rightarrow \pi \Delta$ to differ considerably from that for $\rho N \rightarrow \pi N$, since the $\pi \Delta N$ coupling constant is believed to be about a factor 2 larger than the $\pi NN$ one (cf. ref. 28). Furthermore, the spin and isospin factors favour the $\Delta \pi$ channel, while the phase space factor obviously is larger in the $\pi N$ channel. To the extent that the latter two effects cancel, one would expect a net factor of approximately 4. This agrees approximately with the relative strength of the two channels found by Klingl et al. for an on-shell $\rho$ meson (see Fig. 3). Consequently, if we assume that this ratio is correct, one finds for the total scattering length $\text{Im} a^{(\text{tot})}_{\text{eff}} \approx 5 \text{ Im} a^{(\pi N)}_{\text{eff}} = 0.25 - 0.45 \text{ fm}$ and for the enhancement of the $\rho$-meson width, including both channels, $\Delta \Gamma_{\rho} = 45 - 85 \text{ MeV}$ at $\rho_N = \rho_0$, while Klingl et al. find 1.62 fm, which implies $\Delta \Gamma_{\rho} = 300 \text{ MeV}$.

Note that if the estimate given above is correct and the total width of
the $\rho$ meson at nuclear matter density is $\sim 220$ MeV (on the free mass shell), the QCD sum rule requires a strong reduction of the $\rho$-meson energy at this density (modulo uncertainties due to the approximate treatment of the 4-quark condensate).

Also in the work of Herrmann et al., 11 where the broadening of the $\rho$ meson in nuclear matter is due to the $\Delta\pi$ channel, the enhancement of the width, $\Delta\Gamma_\rho \simeq 150$ MeV at nuclear matter density, is appreciably larger than the estimate given above. A similar model is used by Rapp et al. 9,10 in their calculations of lepton-pair production in nucleus-nucleus collisions. Part of the low-mass enhancement is in this model due to the broadening of the $\rho$ meson in nuclear matter. Clearly any serious attempt to understand the low mass lepton pairs should be confronted with the vector meson production data discussed here.

5 Summary

I have discussed recent theoretical developments on the many-body effects, which determine the properties of vector mesons in nuclear matter and the consequences for the lepton pair spectrum in hadronic collisions. The broadening of $\rho$ mesons in nuclear matter and the momentum dependence of the $\rho$-meson self energy leads to an enhancement of low-mass lepton pairs. In the work of Rapp et al. 9,10 these effects are responsible for most of the low-mass enhancement found by the CERES 5 and Helios-3 6 collaborations at CERN. Thus, the many-body processes offer an interpretation of the low-mass enhancement alternative to the dropping-mass scenario of Li, Ko and Brown. 7

However, elementary processes put constraints on the in-medium width of $\rho$ and $\omega$ mesons, which so far have not been implemented in the calculations of vector-meson properties in nuclear matter. The widths extracted from the data on pion-induced vector-meson production are much smaller than those obtained within various models. Since, in the many-body scenario, a substantial part of the low-mass enhancement is due to the in-medium broadening of the $\rho$ meson, this is a crucial point. How these constraints modify the lepton-pair spectrum in this approach is presently being investigated.

I argued that the root of the problem in some models is the potentially dangerous extrapolation from the photon point to $q^2 = m_V^2$. Consequently, a safer approach would be to exclude the photon data and to rely only on hadron scattering and production data in the relevant kinematic regime.

Finally, I note that Klingl et al. have now revised their calculation, so that the $\pi N$ channel is in agreement with the $\omega$-production data. 30 In the revised model they also include processes that were not included in their original cal-
calculation. These processes enhance the contribution of the \( \pi \pi N \) channel to the in-medium width of the \( \omega \) meson. Thus, Klingl \textit{et al.} now find that, at \( \rho = \rho_0 \), the \( \omega \) width is enhanced by \( \simeq 30 \) MeV and its mass is reduced by \( \simeq 120 \) MeV.

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