Hadrons in Medium

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After a short motivation I outline a consistent treatment of hadronic spectral functions based on transport theory. As examples I discuss nucleon spectral functions and the observable effects of changes of the properties of vector mesons inside nuclei.

1. Introduction

Studies of in-medium properties of hadrons are driven by a number of - partly connected - motivations. A first motivation for the study of hadronic in-medium properties is provided by our interest in understanding the structure of large dense systems, such as the interior of stars. This structure obviously depends on the composition of stellar matter and its interactions (for a recent review see [1]). On a smaller scale, this also holds for the structure of nuclei which is influenced by the properties of their building blocks, the baryons, and that of the exchange mesons that provide the binding.

The second motivation is based on the expectation that changes of hadronic properties in medium can be precursor phenomena to changes in the phase structure of nuclear matter. Here the transition to the chirally symmetric phase, that exhibits manifestly the symmetries of the underlying theory of strong interactions, i.e. QCD, is of particular interest. Present day’s ultrarelativic heavy-ion collisions explore this phase boundary in the limit of high temperatures \( T \approx 170 \text{ MeV} \) and low densities. The other limit (low temperatures and high densities) is harder to reach, although the older AGS heavy-ion experiments and the planned CBM experiment at the new FAIR facility [2] may yield insight into this area. However, even in these experiments the temperatures reached are still sizeable \( T \approx 120 \text{ MeV} \). At very low temperatures the only feasible method seems to be the exploration of the hadronic structure inside ordinary nuclei, at the prize of a low density. Here the temperature is \( T = 0 \) and the density at most equals the equilibrium density of nuclear matter, \( \rho_0 \). It is thus of great interest to explore if such low densities can already give precursor signals for chiral symmetry restoration.

1.1 Phenomenology

That hadrons can indeed change their properties and couplings in the nuclear medium has been well known to nuclear physicists since the days of the Delta-hole model that
dealt with the changes of the properties of the pion and Delta-resonance inside nuclei [3]. Due to the predominant $p$-wave interaction of pions with nucleons one observes here a lowering of the pion branch with increasing pion-momentum and nucleon-density. This effect can be seen in optical model analyses of pion scattering on nuclei, but the absorptive part of the $\pi$-nucleus interaction limits the sensitivity to small densities. More recently, experiments at the FSR at GSI have shown that also the mass of a pion at rest in the nuclear medium differs from its value in vacuum [4]. This is interesting since there are also recent experiments [5] that look for in-medium changes of the $\sigma$ meson, the chiral partner of the pion. Any comparison of scalar and pseudoscalar strength could thus give information about the degree of chiral symmetry restoration in nuclear matter.

In addition, experiments for charged kaon production at GSI [6] have given some evidence for the theoretically predicted lowering of the $K^-$ mass in medium and the (weaker) rising of the $K^+$ mass. State-of-the-art calculations of the in-medium properties of kaons have shown that the usual quasi-particle approximation for these particles is no longer justified inside nuclear matter where they acquire a broad spectral function [7,8].

At higher energies, at the CERN SPS and most recently at the Brookhaven RHIC, in-medium changes of vector mesons have found increased interest. The interest in vector meson production stems from the fact that these mesons couple strongly to the photon so that electromagnetic signals can yield information about properties of hadrons deeply embedded into nuclear matter. Indeed, the CERES experiment [9] has found a considerable excess of dileptons in an invariant mass range from $\approx 300$ MeV to $\approx 700$ MeV as compared to expectations based on the assumption of freely radiating mesons.

This result has found an explanation in terms of a shift of the $\rho$ meson spectral function down to lower masses, as expected from theory (see, e.g., [10–13]). However, the actual reason for the observed dilepton excess is far from clear. This is so partly because any signals from heavy-ion collisions are time-integrals over the often quite complex collision history with very different densities and temperatures.

I have therefore already some years ago proposed to look for the theoretically predicted changes of vector meson properties inside the nuclear medium in reactions on normal nuclei with more microscopic probes [14,15]. Of course, the nuclear density felt by the vector mesons in such experiments lies much below the equilibrium density of nuclear matter, $\rho_0$, so that naively any density-dependent effects are expected to be much smaller than in heavy-ion reactions.

On the other hand, there is a big advantage to these experiments: they proceed with the spectator matter being close to its equilibrium state. This is essential because all theoretical predictions of in-medium properties of hadrons are based on an equilibrium model in which the hadron (vector meson) under investigation is embedded in cold nuclear matter in equilibrium and with infinite extension. The properties so calculated are then, in a second step, being locally inserted into a time-dependent event simulation. In actual experiments these hadrons are observed through their decay products and these have to travel through the surrounding nuclear matter to the detectors. Except for the case of electromagnetic signals (photons, dileptons) this is connected with often sizeable final state interactions (FSI) that have to be treated as realistically as possible. For a long period the Glauber approximation which allows only for absorptive processes along a straight-line path has been the method of choice in theories of photonuclear reactions on nuclei. This may be sufficient if one is only interested in total yields of strongly absorbed particles. However, it is clearly insufficient when one aims at, for example, reconstructing the spectral function
of a hadron inside matter through its decay products. Rescattering and sidefeeding through coupled channel effects can affect the final result so that a realistic description of such effects is absolutely mandatory [16].

2. Theory

2.1 Chiral Symmetry

A large part of the current interest in in-medium properties of hadrons comes from the hope to learn something about quarks in nuclei. More specifically, one hopes to see precursors of a restoration of the original symmetries of the theory of strong interactions, i.e. QCD, which are spontaneously broken in our world. In [17] and in particular in [18] I have discussed this point at some length. Here I just state that the so-called NJL model that is manifestly chirally invariant indeed leads to a linear relationship between the fermion mass and the chiral condensate, the order parameter of chiral symmetry breaking. In this model this condensate in turn drops with baryon density and temperature.

In the NJL model the dropping of the chiral condensate with density and/or temperature directly causes a drop of the mass because both are linearly proportional. This is no longer the case in complex hadronic systems. How the drop of the scalar condensate there translates into observable hadron masses is not uniquely prescribed. The only rigorous connection is given by the QCD sum rules that relate an integral over the hadronic spectral function to a sum over combinations of quark- and gluon-condensates with powers of $1/Q^2$ [19].

Since the spectral function appears under an integral the information obtained is, however, not very specific. However, Leupold et al. have shown [20,21] that the QCDSR provides important constraints for the hadronic spectral functions in medium, but it does not fix them. Recently Kaempfer et al have turned this argument around by pointing out that measuring an in-medium spectral function of the $\omega$ meson could help to determine the density dependence of the chiral condensate [22].

2.2 QCD Sum Rules and Collisional Broadening

Since QCD sum rules do not fix the hadronic properties models are needed for the hadronic interactions. The quantitatively reliable ones can at present be based only on ‘classical’ hadrons and their interactions. Indeed, in lowest order in the density the mass and width of an interacting hadron in nuclear matter at zero temperature and vector density $\rho_v$ are given by (for a meson, for example)

\[
m^* = m^2 - 4\pi \imag \Re f_{mN}(q_0, \theta = 0) \rho_v
\]

\[
m^* \Gamma^* = m \Gamma^0 - 4\pi \imag \Im f_{mN}(q_0, \theta = 0) \rho_v.
\]

Here $f_{mN}(q_0, \theta = 0)$ is the forward scattering amplitude for a meson with energy $q_0$ on a nucleon. The width $\Gamma^0$ denotes the free decay width of the particle. For the imaginary part this is nothing other than the classical relation $\Gamma^* - \Gamma^0 = v \sigma \rho_v$ for the collision width, where $\sigma$ is the total cross section. This can easily be seen by using the optical theorem.
Actually evaluating mass and width from (1) requires knowledge of the scattering amplitude which can only be obtained from very detailed analyses of experiments. The $s$-channel contributions to this scattering amplitude are determined by the properties of nucleon resonances and these are often not very well known yet. Here, resonance physics meets in-medium physics.

Unfortunately it is not a-priori known up to which densities the low-density expansion (1) is useful. Post et al. [10] have recently investigated this question in a coupled-channel calculation of selfenergies. Their analysis comprises pions, $\eta$-mesons and $\rho$-mesons as well as all baryon resonances with a sizeable coupling to any of these mesons. The authors of [10] find that already for densities less than $0.5\rho_0$ the linear scaling of the selfenergies inherent in (1) is badly violated for the $\rho$ and the $\pi$ mesons, whereas it is a reasonable approximation for the $\eta$ meson. Reasons for this deviation from linearity are Fermi-motion, Pauli-blocking, selfconsistency and short-range correlations. For different mesons different sources of the discrepancy prevail: for the $\rho$ and $\eta$ mesons the iterations act against the low-density theorem by inducing a strong broadening for the $D_{13}(1520)$ and a slightly repulsive mass shift for the $S_{11}(1535)$ nucleon resonances to which the $\rho$ and the $\eta$ meson, respectively, couple. The investigation of in-medium properties of mesons, for example, thus involves at the same time the study of in-medium properties of nucleon resonances and is thus a coupled-channel problem.

2.3 Coupled Channel Treatment of Incoherent Particle Production

In order to avoid the intrinsic difficulties connected with using equilibrium hadronic properties in a non-equilibrium situation such as a heavy-ion reaction, we have looked for possible effects in reactions that proceed closer to equilibrium, i.e. reactions of elementary probes such as protons, pions, and photons on nuclei. The densities probed in such reactions are always $\leq \rho_0$, with most of the nucleons actually being at about $0.5\rho_0$. On the other hand, the target is stationary and the reaction proceeds much closer to (cold) equilibrium than in a relativistic heavy-ion collision. If any observable effects of in-medium changes of hadronic properties survive, even though the densities probed are always $\leq \rho_0$, then the study of hadronic in-medium properties in reactions with elementary probes on nuclei provides an essential baseline for in-medium effects in hot nuclear matter probed in ultra-relativistic heavy-ion collisions.

With the aim of exploring this possibility we have over the last few years undertaken a number of calculations for proton- [23], pion- [24,25] and photon- [26] induced reactions. All of them have one feature in common: they treat the final state incoherently in a coupled channel transport calculation that allows for elastic and inelastic scattering of, particle production by and absorption of the produced hadrons. A new feature of these calculations is that hadrons with their correct spectral functions can actually be produced and transported consistently. This is quite an advantage over earlier treatments [27,28] in which the mesons were always produced and transported with their pole mass and their spectral function was later on folded in only for their decay. The method is summarized in the following section, more details can be found in [26].

We separate the photonuclear reaction into three steps. First, we determine the amount of shadowing for the incoming photon; this obviously depends on its momentum transfer $Q^2$. Second, the primary particle is produced and third, the produced particles are propagated
through the nuclear medium until they leave the nucleus.

a Shadowing. Photonnuclear reactions show shadowing in the entrance channel, for real photons from an energy of about 1 GeV on upwards [29]. This shadowing is due to a coherent superposition of bare photon and vector meson components in the incoming photon and is handled here by means of a Glauber multiple scattering model [16]. In this way we obtain for each value of virtuality $Q^2$ and energy $\nu$ of the photon a spatial distribution for the probability that the incoming photon reaches a given point; for details see [16,30,31].

b Initial Production. The initial particle production is handled differently depending on the invariant mass $W = \sqrt{s}$ of the excited state of the nucleon. If $W < 2$ GeV, we invoke a nucleon resonance model that has been adjusted to nuclear data on resonance-driven particle production [26]. If $W > 2$ GeV the particle yield is calculated with standard codes developed for high energy nuclear reactions, i.e. FRITIOF or PYTHIA; details are given in [32]. We have made efforts to ensure a smooth transition of cross sections in the transition from resonance physics to DIS.

c Groundstate Correlations and Final State Interactions. The groundstate correlations and the final state interactions are consistently described by a semiclassical coupled channel transport theory that had originally been developed for the description of heavy-ion collisions and has since then been applied to various more elementary reactions on nuclei with protons, pions and photons in the entrance channel.

In this method the spectral phase space distributions of all particles involved are propagated in time, from the initial first contact of the photon with the nucleus all the way to the final hadrons leaving the nuclear volume on their way to the detector. The spectral phase space distributions $F_h(\vec{r}, \vec{p}, \mu, t)$ give at each moment of time and for each particle class $h$ the probability to find a particle of that class with a (possibly off-shell) mass $\mu$ and momentum $\vec{p}$ at position $\vec{r}$. Its time-development is determined by the BUU equation

$$\left( \frac{\partial}{\partial t} + \frac{\partial H_h}{\partial \vec{p}} \frac{\partial}{\partial \vec{r}} - \frac{\partial H_h}{\partial \vec{r}} \frac{\partial}{\partial \vec{p}} \right) F_h = G_h a_h - L_h F_h.$$  

(2)

Here $H_h$ gives the energy of the hadron $h$ that is being transported; it contains the mass, the selfenergy (mean field) of the particle and a term that drives an off-shell particle back to its mass shell. The terms on the lhs of (2) are the so-called drift terms since they describe the independent transport of each hadron class $h$. The terms on the rhs of (2) are the collision terms; they describe both elastic and inelastic collisions between the hadrons. Here the term inelastic collisions includes those collisions that either lead to particle production or particle absorption. The former is described by the gain term $G_h a_h$ on the rhs in (2), the latter process (absorption) by the loss term $L_h F_h$. Note that the gain term is proportional to the spectral function $a$ of the particle being produced (see discussion below), thus allowing for production of off-shell particles. On the contrary, the loss term is proportional to the spectral phase space distribution itself: the more particles there are the more can be absorbed. The terms $G_h$ and $L_h$ on the rhs give the actual strength of the gain and loss terms, respectively. They have the form of Born-approximation collision integrals and take the Pauli-principle into account. The free collision rates themselves are taken from experiment or are calculated [26].

The collision term on the lhs of (2) is responsible for the collision broadening that all particles experience when they are embedded in a dense medium. Collisions either change energy and momentum of the particles or absorb them altogether. Both processes contribute to collisional broadening.
The detailed structure of the gain and loss terms can be obtained from quantum transport theory [33,34]. To see this I start by summarizing briefly the known fundamental relations of quantum transport theory for the description of non-stationary processes in an interacting quantum system, following closely the presentation in [35]. In quantum transport theory the non-stationary processes which introduce a coupling between causal and anti-causal single particle propagation are described by the one-particle correlation functions

$$ g^{>}(1, 1') = -i \langle \Psi(1) \Psi^\dagger(1') \rangle \quad g^{>}(1, 1') = i \langle \Psi^\dagger(1') \Psi(1) \rangle . $$

(3)

where $\Psi$ are the nucleon field operators in Heisenberg representation. Correspondingly, in an interacting quantum system the single particle self-energy operator includes correlation self-energies $\Sigma^{<>}$ which couple particle and hole degrees of freedom [33,34]. Clearly, $g^{<>}$ and $\Sigma^{<>}$ are closely related. The wanted relation is obtained from transport theory. After a Fourier transformation to energy-momentum representation the corresponding self-energies are found as [33]

$$ \Sigma^{>}(\omega, p) = g \int \frac{d^3p_2}{(2\pi)^4} \frac{d\omega_2}{(2\pi)^2} \frac{d^3p_3}{(2\pi)^4} \frac{d\omega_3}{(2\pi)^2} \frac{d^3p_4}{(2\pi)^4} \frac{d\omega_4}{(2\pi)^4} \delta^4(p + p_2 - p_3 - p_4) |\tilde{M}|^2 \times g^{<}(\omega_2, p_2)g^{>}(\omega_3, p_3)g^{>}(\omega_4, p_1) $$

$$ \Sigma^{<}(\omega, p) = g \int \frac{d^3p_2}{(2\pi)^4} \frac{d\omega_2}{(2\pi)^2} \frac{d^3p_3}{(2\pi)^4} \frac{d\omega_3}{(2\pi)^2} \frac{d^3p_4}{(2\pi)^4} \frac{d\omega_4}{(2\pi)^4} \delta^4(p + p_2 - p_3 - p_4) |\tilde{M}|^2 \times g^{>}(\omega_2, p_2)g^{<}(\omega_3, p_3)g^{<}(\omega_4, p_4) . $$

(4)

Here, $g = 4$ is the spin-isospin degeneracy factor and $|\tilde{M}|^2$ denotes the square of the in-medium nucleon-nucleon scattering amplitude, averaged over spin and isospin of the incoming nucleons and summed over spin and isospin of the outgoing nucleons. Note that energy $\omega$ and three-momentum $p$ are not related by a dispersion relation thus allowing for a description of the full off-shell behavior of the self-energies.

Since both $g^{<>}$ and $\Sigma^{<>}$ describe the correlation dynamics, the spectral function can be obtained from either of the two quantities as the difference over the cut along the energy real axis. In terms of the correlation propagators, the spectral density is defined by

$$ a(\omega, p) = i (g^{>}(\omega, p) - g^{<}(\omega, p)) . $$

(5)

Non-relativistically, the single particle spectral function is explicitly found as

$$ a(\omega, p) = \Gamma(\omega, p) \frac{\Gamma(\omega, p)}{\left(\omega - \frac{p^2}{2m_N} - \text{Re}\Sigma^{<}(\omega, p) + \frac{1}{2}\text{Re}\Sigma^{>}(\omega, p)\right)^2 + \frac{1}{4}\text{Im}^2(\omega, p)} , $$

(6)

including the particle and hole nucleon self-energy $\Sigma$. The width $\Gamma$ is given by the imaginary part of the retarded self-energy

$$ \Gamma(\omega, p) = 2 \Im\Sigma^{<}(\omega, p) = i(\Sigma^{>}(\omega, p) - \Sigma^{<}(\omega, p)) . $$

(7)

In the limiting case of vanishing correlations, i.e. $\Im\Sigma \rightarrow 0$, the usual deltaliike quasiparticle spectral function is recovered.

The correlation functions $g^{<>}$ can be re-written in terms of the phasespace distribution function $f(\omega, p)$.
\[ g^>(\omega, p) = -ia(\omega, p)(1 - f(\omega, p)) \]
\[ g^<(\omega, p) = ia(\omega, p)f(\omega, p). \] (8)

For a system at \( T = 0 \), \( f \) reduces to
\[ f(\omega, p) = \Theta(\omega_F - \omega). \] (9)

As a result, we obtain for the self-energies the conditions
\[ \Sigma^>(\omega, p) = 0 \quad \Gamma(\omega, p) = -i\Sigma^<(\omega, p) \quad \text{for} \quad \omega \leq \omega_F \]
\[ \Sigma^<(\omega, p) = 0 \quad \Gamma(\omega, p) = i\Sigma^>(\omega, p) \quad \text{for} \quad \omega \geq \omega_F. \] (10)

The correlation function \(-ig^<(\omega, p)\) is nothing else than the Fourier-transform of the spectral phase space density \( F_h \) in Eq. (2), with the variable \( \mu \) in (2) being an ‘off-shell mass’ \( \mu = \sqrt{\omega^2 - p^2} \). Also the gain and loss terms in (2) can be expressed in terms of the correlations functions defined here. As a result, we have
\[ G = i\Sigma^<(1 - f) \quad L = i\Sigma^>. \] (11)

Due to the dependence of the width and the single particle spectral function upon each other the calculation of \( a(\omega, p) \) requires a self-consistent treatment. Therefore, the transport theoretical approach leads to single particle propagators including correlation self-energies to all orders. The collision rates embedded in \( G \) and \( L \) determine the collisional broadening of the particles involved and thus their spectral function \( a \). The widths of the particles, resonances or mesons, thus evolve in time away from their vacuum values. In addition, broad particles can be produced off their peak mass and then propagated. The extra ‘potential’ in \( H \) already mentioned ensures that all particles are being driven back to their mass-shell when they leave the nucleus. The actual method used is described in [26]. It is based on an analysis of the Kadanoff-Baym equation that has led to practical schemes for the propagation of off-shell particles [36,37]. The possibility to transport off-shell particles represents a major breakthrough in this field. For further details of the model see Ref. [26] and [32] and references therein.

3. In-medium Hadrons – Observables

3.1 Nucleon Spectral Functions

Even the elementary building blocks of nuclei, the nucleons, are medium-modified inside nuclei. Due to correlations they pick up a width and a mass-shift, the latter on top of that caused by the long range mean field. This becomes evident by considering the case of elastic collisions only in Eq. (2). Nucleons moving in a mean field can still collide, if the final states are above the Fermi-surface thus acquiring a collisional width. This is illustrated in Fig. 1.

The nucleon spectral functions in Fig. 1 show the distinct quasiparticle peak at the energy \( \omega = \sqrt{p^2 + m^2} \), broadened by collisions with the other nucleons. The width gets larger with increasing momentum due to the phase space opening up at the higher
momenta. In addition to the quasiparticle peak the spectral functions exhibit a distinct zero at the Fermi energy, caused by the Pauli principle.

Observable effects of this broadening of nucleon spectral functions in medium are well known: they show up in $(e, e'p)$ reactions on nuclei where an analysis of the momentum and energy of the outgoing proton allows to determine its spectral function inside the nucleus. The method has recently been used to calculate the spectral function of nucleons also at finite temperatures and higher densities [38], as well as for isospin-asymmetric nuclear matter [39].

3.2 $\omega$ Production

Many of the early studies of hadronic properties in medium concentrated on the $\rho$-meson [11,40], partly because of its possible significance for an interpretation of the CERES experiment. It is clear by now, however, that the dominant effect on the in-medium properties of the $\rho$-meson is collisional broadening that overshadows any possible mass shifts [10] and is thus experimentally hard to observe. The emphasis has, therefore, shifted to the $\omega$ meson. An experiment measuring the $A(\gamma, \omega \rightarrow \pi^0\gamma')X$ reaction is presently being analyzed by the TAPS/Crystal Barrel collaborations at ELSA [41]. The varying theoretical predictions for the $\omega$ mass (640-765 MeV) [40] and width (up to 50 MeV) [24,42] in nuclear matter at rest encourage the use of such an exclusive probe to learn about the $\omega$ spectral distribution in nuclei.
Figure 2. Mass differential cross section for $\pi^0\gamma$ photoproduction off $^{93}\text{Nb}$. Shown are results both with and without a mass-shift as explained in [44]. The quantity $p$ gives the escape probability for one of the four photons in the $2\pi^0$ channel. The two solid curves give results of calculations with $p = 26\%$ with and without mass-shift, the two dashed curves give the same for $p = 5\%$ (from [43]).

Simulations have been performed at 1.2 GeV and 2.5 GeV photon energy, which cover the accessible energies of the TAPS/Crystal Barrel experiment. After reducing the combinatorial and rescattering background by applying kinematic cuts on the outgoing particles, we have obtained rather clear observable signals for an assumed dropping of the $\omega$ mass inside nuclei [44]. Therefore, in this case it should be possible to disentangle the collisional broadening from a dropping mass.

Our calculations represent complete event simulations. It is, therefore, possible to calculate these background contributions and to take experimental acceptance effects into account. An example is shown in Fig. 2 which shows the effects of a possible misidentification of the $\omega$ meson. This misidentification can come about through the $2\pi^0 \rightarrow 4\gamma$ channel if one of the four photons escapes detection and the remaining three photons are identified as stemming from the $\pi^0\gamma \rightarrow 3\gamma$ decay channel of the $\omega$-meson. The calculations show that the misidentification does not affect the low-mass side of the omega spectral function.

Fig. 2 shows a good agreement between the data of the TAPS/CB@ELSA collaboration [45] for a photon escape probability of 5\% and a mass shift $m_\omega = m^0_\omega - 0.18 \rho/\rho_0$. In [44] we have also discussed the momentum-dependence of the $\omega$-selfenergy in medium and have pointed out that this could be accessible through measurements which gate on different three-momenta of the $\omega$ decay products.
Dileptons, i.e. electron-positron pairs, in the outgoing channel are an ideal probe for in-medium properties of hadrons since they experience no strong final state interaction. A first experiment to look for these dileptons in heavy-ion reactions was the DLS experiment at the BEVALAC in Berkeley [46]. Later on, and in a higher energy regime, the CERES experiment has received a lot of attention for its observation of an excess of dileptons with invariant masses below those of the lightest vector mesons [9]. Explanations of this excess have focused on a change of in-medium properties of these vector mesons in dense nuclear matter (see e.g. [47,48]). The radiating sources can be nicely seen in Fig. 3 that shows the dilepton spectrum obtained in a low-energy run at 40 AGeV together with the elementary sources of dilepton radiation.

The figure exhibits clearly the rather strong contributions of the vector mesons – both direct and through their Dalitz decay – at invariant masses above about 500 MeV. The strong amplification of the dilepton rate at small invariant masses $M$ caused by the photon propagator, which contributes $\sim 1/M^4$ to the cross section, leads to a strong sensitivity to changes of the spectral function at small masses. Therefore, the excess observed in the CERES experiment can be explained by such changes as has been shown by various authors (see e.g. [27] for a review of such calculations).

In view of the uncertainties in interpreting these results discussed earlier we have studied the dilepton photo-production in reactions on nuclear targets. Looking for in-medium changes in such a reaction is not a priori hopeless: Even in relativistic heavy-ion reactions...
only about 1/2 of all dileptons come from densities larger than $2\rho_0$ [27]. In these reactions the pion-density gets quite large in the late stages of the collision. Correspondingly many $\rho$ mesons are formed (through $\pi + \pi \to \rho$) late in the collision, where the baryonic matter expands and its density becomes low again.

In [26] we have analyzed the photoproduction of dileptons on nuclei in great detail. After removing the Bethe-Heitler contribution the dilepton mass spectrum in a 2 GeV photon-induced reaction looks very similar to that obtained in an ultrarelativistic heavy-ion collision (Fig. 3). The radiation sources are all the same in both otherwise quite different reactions. The photon-induced reaction can thus be used as a baseline experiment that allows one to check crucial input into the simulations of more complicated heavy-ion collision.
A typical result of such a calculation for the dilepton yield – after removing the Bethe-Heitler component – is given in Fig. 4. The lower part of Fig. 4 shows that we can expect observable effect of possible in-medium changes of the vector meson spectral functions in medium on the low-mass side of the \( \omega \) peak. In [26] we have shown that these effects can be drastically enhanced if proper kinematic cuts are introduced that tend to enhance the in-medium decay of the vector mesons. There it was shown that in the heavy nucleus \( Pb \) the \( \omega \)-peak completely disappears from the spectrum if in-medium changes of width and mass are taken into account. The sensitivity of such reactions is thus as large as that observed in ultrarelativistic heavy-ion reactions.

An experimental verification of this prediction would be a major step forward in our understanding of in-medium changes. The ongoing g7 experiment at JLAB is presently analyzing such data [50]. This experiment can also yield important information on the time-like electromagnetic formfactor of the proton and its resonances [51] on which little or nothing is known.

4. Conclusions

In this talk I have first outlined the theoretical motivation for studies of in-medium properties of hadrons and their relation to QCD. I have then shown that transport theory can be used to calculate the in-medium properties of hadrons and their interactions in a consistent way. In particular photonuclear reactions on nuclei can give observable consequences of in-medium changes of hadrons that are as big as those expected in heavy-ion collisions which reach much higher densities, but proceed farther away from equilibrium. Special emphasis was put in these lectures not so much on the theoretical calculations of hadronic in-medium properties under simplified conditions, but more on the final, observable effects of any such properties. I have discussed that for reliable predictions of observables one has to take the final state interactions with all their complications in a coupled channel calculation into account; simple Glauber-type descriptions are not sufficient.

A first, well known example for the effects of in-medium interactions is given by the spectral functions of nucleons inside nuclei. As an example that is free from complications by FSI I have shown that in photonuclear reactions in the 1 - 2 GeV range the expected sensitivity of dilepton spectra to changes of the \( \rho \)- and \( \omega \) meson properties in medium is as large as that in ultrarelativistic heavy-ion collisions and that exactly the same sources contribute to the dilepton yield in both experiments. While the dilepton decay channel is free from hadronic final state interactions this is not so when the signal has to be reconstructed from hadrons present in the final state. Nevertheless, the \( \omega \) photoproduction, identified by the semi-hadronic \( \pi^0 \gamma \) decay channel, seems to exhibit a rather clean in-medium signal.

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