A Microscopic Calculation of Photoabsorption Cross Sections on Protons and Nuclei

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Abstract

A recently developed model for ρ-meson propagation in dense hadronic matter is applied to total photoabsorption cross sections in γ-proton and γ-nucleus reactions. Within the vector dominance model the photon coupling to the virtual pion cloud of the nucleon, two-body meson-exchange currents, as well as γ-nucleon resonances are included. Whereas the γp reaction is determined by the low-density limit of the model, higher orders in the nuclear density are important to correctly account for the experimental spectra observed on both light and heavy nuclei over a wide range of photon energies, including the region below the pion threshold. In connection with soft dilepton spectra in high-energy heavy-ion collisions we emphasize the importance of photoabsorption to further constrain the parameters of the model.

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Recent measurements of dilepton spectra in heavy-ion collisions at both intermediate and high bombarding energies have shown a strong enhancement of the pair yield in the low-invariant-mass region \((M \approx 0.2-1.5 \text{ GeV})\) as compared to expectations based on free hadronic sources \([1–3]\). So far, the most promising approaches to explain these data are based on medium modifications in \(\pi^+ \pi^- \rightarrow \ell^+ \ell^-\) annihilation occurring during the interaction phase of the colliding nuclear system. In particular, the assumption of a dropping \(\rho\)-meson mass \([4]\) in the hot and dense medium has been shown to give a good description of the HELIOS-3 and CERES data \([5,6]\). On the other hand, the inclusion of in-medium hadronic interactions in \(\rho\) - and \(\pi\pi\) propagation also gives reasonable agreement with these experiments \([7,8]\). It is evident that any model of dilepton enhancement has to be in accordance with a wide variety of related data, e.g. the free \(\rho\) meson in the time-like region has to be accounted for by properly describing p-wave \(\pi\pi\) scattering and the pion electromagnetic form factor. Obviously another important constraint is provided by photoabsorption experiments \([9,10]\), which represent the limit of vanishing invariant mass of the (virtual) photon, \(M^2 \rightarrow 0\). In this note we will extend our model for \(\rho\)-meson propagation in hadronic matter \([7]\) to analyze photoabsorption spectra on both protons and nuclei, thereby further constraining the model parameters.

The starting point is the general expression for the total absorption cross section of a photon on a volume element \(d^3x\) of cold nuclear matter

\[
\frac{d\sigma}{d^3x} = -\frac{4\pi\alpha}{q_0} \epsilon_\mu(q, \lambda) \epsilon_\nu(q, \lambda) \text{Im} G^{\mu\nu}(q_0, \vec{q}; \rho_N). \tag{1}
\]

Here, \(\epsilon_\mu\) and \((q_0, \vec{q})\) denote the photon polarization vector and four-momentum, respectively, while \(G^{\mu\nu}\) represents the electromagnetic current correlator of the hadronic source at a given nuclear density \(\rho_N\). Invoking the vector dominance model (VDM) and neglecting small contributions from isoscalar vector mesons the correlator is determined by the \(\rho\)-meson propagator as

\[
G^{\mu\nu}(q_0, q; \rho_N) \equiv \frac{(m_\rho^{(0)})^4}{g^2} D^{\mu\nu}_\rho(q_0, q; \rho_N) \tag{2}
\]

where

\[
D^{\mu\nu}_\rho(q_0, q; \rho_N) = \frac{P^{\mu\nu}_L}{M^2 - (m_\rho^{(0)})^2 - \Sigma^L_\rho(q_0, q; \rho_N)} + \frac{P^{\mu\nu}_T}{M^2 - (m_\rho^{(0)})^2 - \Sigma^T_\rho(q_0, q; \rho_N)} \\
+ \frac{q^\mu q^\nu}{(m_\rho^{(0)})^2 M^2}. \tag{3}
\]

Here \(P^{\mu\nu}_L\) (\(P^{\mu\nu}_T\)) is the standard longitudinal (transverse) projection operator and \(\Sigma^L_\rho\) (\(\Sigma^T_\rho\)) the corresponding scalar part of the selfenergy. For real photons with \(M^2 = q_0^2 - \vec{q}^2 = 0\) Eq. (1) can be rewritten as the total photoabsorption cross section normalized to the number of nucleons, \(A\), as

\[
\frac{\sigma^{\text{abs}}_A}{A} = -\frac{4\pi\alpha}{q_0} \frac{(m_\rho^{(0)})^4}{g^2} \frac{1}{\rho_N} \text{Im} D^T_\rho(q_0, \vec{q}; \rho_N) \tag{4},
\]
\[ D^T \rho \] denotes the transverse \( \rho \)-meson propagator, defined through \( \Sigma^T_\rho \).

For the actual calculations we have to specify a model for the in-medium \( \rho \)-meson propagator. It consists of a bare \( \rho \) with mass \( m^{(0)}_\rho \) renormalized through various selfenergy contributions involving two-pion and single-baryon interactions

\[ \Sigma^L,T_\rho = \Sigma^L,T_\rho^{\pi\pi} + \Sigma^L,T_\rho^N. \]  

(5)

In free space \( (\rho_N=0) \), only \( \Sigma^{0}_{\rho\pi\pi}(M) \equiv \Sigma_{\rho\pi\pi}(q_0, \vec{q}; \rho_N = 0) \) survives, representing the coupling of the \( \rho \) to vacuum two-pion states. Aside from the two-pion loop we include a pion-tadpole contribution rendering \( \Sigma^{0}_{\rho\pi\pi}(M) \) transverse and zero at the photon point, once properly regularized \[12\]. The parameters are fixed by ensuring a good description of the p-wave \( \pi\pi \) phase shifts as well as the pion electromagnetic form factor in vacuum. At finite density \( \Sigma_{\rho\pi\pi} \) is modified through pion interactions with the surrounding nucleons \[7\]. As is well known, the dominant contribution to the in-medium pion propagator arises from p-wave nucleon-hole and delta-hole polarizations. The corresponding selfenergies contain coupling constants \( f_{\pi\alpha} \) \((\alpha = NN^{-1}, \Delta N^{-1})\) related via \( f_{\pi N\Delta} = 2f_{\pi NN} \) (Chew-Low factor \[11\]), a monopole form factor

\[ F_{\pi\alpha}(k) = \left( \frac{\Lambda^2_\pi - m^2_\pi}{\Lambda^2_\pi + k^2} \right), \]  

(6)

spin-isospin factors \( SI(\pi\alpha) \) and a Lindhard function \( \phi_{\alpha}(\omega, k) \) for the loop integration over the nucleon Fermi sea (see below). The precise value of the cutoff parameter \( \Lambda_\pi \) will be determined from the fit to the photoabsorption data. Furthermore, the pion selfenergies have to be corrected for short-range correlation effects, conveniently parameterized in terms of Migdal parameters \( g'_{\alpha\beta} \). The calculations of ref. \[4\] were done in back-to-back kinematics \((\vec{q} = 0)\) for simplicity. When going to the photon point, however, one necessarily has to allow for finite 3-momentum of the \( \rho \) meson relative to the rest frame of the nuclear medium. Recently this has been achieved by Urban et al. maintaining exact conservation of the hadronic vector current and will be discussed separately \[12\]. In the low-density limit, the in-medium \( \rho \)-meson propagator per nucleon reduces to the forward Compton amplitude on the proton. The model specified above describes the coupling of the photon to the virtual pion cloud of the nucleon via an intermediate \( \rho \) meson and yields non-resonant ‘background contributions’ to the Compton amplitude.

The second piece of the in-medium \( \rho \) selfenergy, \( \Sigma_{\rho N} \), in Eq. (5) arises from direct coupling of the \( \rho \) meson to the surrounding nucleons leading to nucleonic resonances. We assume the \( \rho N \) amplitude to be governed by s-channel pole graphs. This was first discussed in ref. \[13\] for the case of the \( N(1720) \) and \( \Delta(1905) \), which both show a large branching ratio \((> 60\%)\) into the \( \rho N \) channel. However, the photoabsorption data (especially for the free proton) require the inclusion of additional, lower-lying resonances. We account for the most important states which will allow us to saturate the experimental spectra. They can be divided into two groups:

(i) positive parity states, which exhibit a predominant p-wave coupling to \( \rho N \). In the non-relativistic limit, suitable interaction Lagrangians are given by

\[ \mathcal{L}^p_{\rho BN} = \frac{f_{\rho BN}}{m_\rho} \Psi_B (\vec{s} \times \vec{q}) \cdot \vec{t}^a \Psi_N + h.c., \]  

(7)
where the summation over \( a \) is in isospin space. The spin operators \( \vec{s} = \vec{\sigma}, \vec{S} \) for \( J=1/2,3/2 \) and the isospin operators \( \vec{I} = \vec{T}, \vec{T} \) for \( I=1/2,3/2 \), respectively, are chosen in accordance with the quantum numbers of baryon \( B \) for \( B=N(939), \Delta(1232) \) and \( N(1720) \), where \( \sigma_j \) and \( \tau_a \) are the usual Pauli matrices and \( S_j, T_a \) the isospin operators.

(ii) negative parity states, which exhibit a predominant \( s \)-wave coupling to \( \rho N \). In the non-relativistic limit, the interaction Lagrangians can be chosen as \[ L_{\rho BN}^{s-wave} = \frac{f_{\rho BN}}{m_\rho} \Psi_B \left( \bar{q}_0 \vec{s} \cdot \vec{r}_a - \rho_a^0 \vec{s} \cdot \vec{q} \right) t_a \Psi_N + h.c. \] for \( B=N(1520), \Delta(1620), \Delta(1700) \).

From these interaction vertices we derive in-medium self-energy tensors for \( \rho \)-like \( BN^{-1} \) excitations. In close analogy to the pionic case one obtains for the purely transverse \( p \)-wave contributions

\[ \Sigma_{\rho_{a,pw}}^{(0),T}(q_0, q) = - \left( \frac{f_{\rho a} F_{\rho a}(q)}{m_\rho} \right)^2 SI(\rho a) q^2 \phi_{\rho a}(q_0, q), \]  

while for \( s \)-waves

\[ \Sigma_{\rho_{a,sw}}^{(0),T}(q_0, q) = - \left( \frac{f_{\rho a} F_{\rho a}(q)}{m_\rho} \right)^2 SI(\rho a) q_0^2 \phi_{\rho a}(q_0, q), \]

with a monopole form factor \( F_{\rho a}(q) = \Lambda_\rho^2/(\Lambda_\rho^2 + q^2) \) and spin-isospin factors \( SI(\rho a) \) summarized in table \( \text{II} \) (note that for \( M^2 = 0 \) we have \( q^2 = q_0^2 \) and thus the expressions for \( s \)-wave and \( p \)-wave coupling become identical). In analogy to pion-induced excitations we include short-range correlation effects in the particle-hole bubble through Migdal parameters \( g' \), which also induce a mixing between the various excitations of a given partial wave \( \text{[7]} \). The explicit form of the Lindhard functions reads

\[ \phi_{\rho a}(q_0, \vec{q}) = - \int_0^{p_f} \frac{p^2 dp}{(2\pi)^2} \int_{-1}^{+1} dx \sum_{+,-} \frac{1}{\pm q_0 + E_p^N - E_{p0}^B(x) \pm \frac{1}{2} \Gamma_{tot}} , \]  

which is equivalent to the pionic case \( E_p^N = \sqrt{m_N^2 + p^2}, E_{p0}^B(x) = \sqrt{m_B^2 + q^2 + p^2 + 2pqx} \). The free baryon widths are each taken as the sum of \( \rho N \) channel and \( \pi N \) channel in the appropriate partial wave,

\[ \Gamma_{B}^0(s) = \Gamma_{B \to \rho N}^0(s) + \Gamma_{B \to \pi N}^0(s) . \]  

In the nuclear medium we account for a density-dependent correction as

\[ \Gamma_B(s; \rho) = \Gamma_{B}^0(s) + \Gamma_{B \to \rho N}^{med} \frac{\rho}{\rho_0} , \]
where possible energy dependencies in $\Gamma_{B}^{med}$ have been neglected for simplicity. It remains to fix the coupling constants $f_{\rho BN}$. For $B=\Lambda (939)$ and $\Delta (1232)$ we take values close to the Bonn potential; all other resonances considered have a sizable branching ratio into the $\rho N$ channel. These branching ratios are used to obtain an estimate for the $\rho BN$ coupling constants via

$$\Gamma_{0 B \rightarrow N \rho}^{0} (\sqrt{s}) = \frac{f_{\rho BN}^{2}}{4\pi m_{\rho}^{2}} \frac{2m_{N}}{\sqrt{s}} \frac{(2I_{\rho} + 1)}{(2J_{B} + 1)(2I_{B} + 1)} SI(\rho BN) \sqrt{s-m_{N}} \int_{2m_{\pi}}^{M d M} \frac{A_{\rho}^{0} (M) q_{cm} F_{\rho} (q_{cm})^{2}}{\pi} f (M),$$

(14)

where the kinematic factor $f (M)$ is given by $f (M) = q_{cm}^{2}$ for p-wave coupling and $f (M) = (2q_{0}^{2} + M^{2})$ for s-wave coupling with

$$q_{cm}^{2} = \frac{(s - M^{2} - m_{N}^{2})^{2} - 4m_{N}^{2}M^{2}}{4s}$$

(15)

being the $\rho / N$ decay momentum in the resonance rest frame. $A_{\rho}^{0} (M) = -2 \text{Im} D_{\rho}^{0} (M)$ denotes the $\rho$-meson spectral function in the vacuum. The $\pi BN$ coupling constant used for $\Gamma_{0 B \rightarrow \pi N}$ in Eq. (12) is then chosen such that the total width matches its experimental value at the resonance mass $s = m_{B}^{2}$.

As has been noted long ago the most simple version of the VDM, Eq. (2), typically results in an overestimation of the $B \rightarrow N \gamma$ branching fractions when using the hadronic coupling constants deduced from the $B \rightarrow N \rho$ partial widths. However, one can correct for this by employing an improved version of the VDM [15], which allows to adjust the $BN \gamma$ coupling $\mu_{B}$ (the transition magnetic moment) at the photon point independently [13]. It amounts to replacing the combination $(m_{\rho}^{0})^{4} \text{Im} D_{\rho}^{T} (q_{0}, \bar{q}; \rho N)$ entering Eq. (4) by the following 'transition form factor':

$$\bar{F} (q_{0}, q; \rho N) = -\text{Im} \Sigma^{T}_{\rho \pi \pi} [d_{\rho} - 1]^{2} - \text{Im} \Sigma^{T}_{\rho \pi N} [d_{\rho} - r_{B}]^{2}$$

$$d_{\rho} (q_{0}, q; \rho N) = \frac{M^{2} - \Sigma^{T}_{\rho \pi \pi} - r_{B} \Sigma^{T}_{\rho \pi N}}{M^{2} - (m_{\rho}^{(0)})^{2} - \Sigma^{T}_{\rho \pi \pi} - \Sigma^{T}_{\rho \pi N}}$$

(16)

where

$$r_{B} = \frac{\mu_{B}}{f_{\rho BN} \frac{(m_{\rho}^{(0)})^{2}}{g}}$$

(17)

denotes the ratio of the photon coupling to its value in the naive VDM. In principle each resonance state $B$ can be assigned a separate value for $r_{B}$ but, as will be seen below, reasonable fits to the photoabsorption spectra can be achieved with a common value for both $r_{B}$ and $\Lambda_{\rho}$, making use of some freedom in the hadronic couplings $f_{\rho BN}$ within the experimental uncertainties of the partial widths, Eq. (14). The final formula to be used for the photoabsorption calculations then reads

$$\sigma_{abs}^{\gamma A} = -\frac{4\pi \alpha}{g^{2}q_{0}^{0}} \frac{1}{\rho_{N}} \bar{F} (q_{0}, q; \rho N).$$

(18)
In taking the low-density limit, \( \rho_N \rightarrow 0 \), only terms linear in density contribute to \( \bar{F} \), representing the absorption process on a single nucleon, as mentioned above. Fig. 1 shows the resulting cross section with the free parameters taken as \( \Lambda_\pi=550 \text{ MeV} \), \( \Lambda_\rho=600 \text{ MeV} \) \((\Lambda_{\rho NN}=1 \text{ GeV})\) and \( r_B=0.5 \) (the actual hadronic coupling constants \( f_{\rho BN} \) are given in table 1). The various contributions from \( \rho \)-meson coupling to the virtual pion cloud of the nucleon (the 'background' described by \( \Sigma_{\rho\pi\pi} \)) as well as from the \( \rho N \) resonances (contained in \( \Sigma_{\rho N} \)) add incoherently.

For absorption spectra on nuclei, one experimentally observes an almost independent scaling with the atomic number \( A \) of different nuclei \((cp. \text{ Fig. 2})\). This justifies to perform the calculations in infinite nuclear matter at an average density, which we take to be \( \bar{\rho}_N=0.8\rho_0 \) (we have checked that the results for the normalized cross section, Eq. (18), only weakly depend on density within reasonable limits). As compared to the free proton two additional features appear in the nuclear medium: short-range correlation effects in the resummation of the particle-hole bubbles and in-medium corrections to the resonance widths. Due to the rather soft form factors involved, the \( p \)-wave excitations turn out to favor rather small Landau-Migdal parameters of \( g'_{NN}=0.6 \) and \( g'_{\rho\rho}=0.25 \) for all other transitions, whereas the \( s \)-wave bubbles show no significant evidence for short-range correlations (therefore we set \( g'_{s-\text{wave}}=0 \)). Note that the rather large in-medium correction to the \( N(1520) \) width, \( \Gamma_{\text{med}}^{N(1520)}=250 \text{ MeV} \), can be understood microscopically in a selfconsistent treatment of the \( \rho \) spectral function in nuclear matter \([14]\). On the other hand, the net in-medium correction to the \( \Delta(1232) \) width is quite small. This reflects the fact that a moderate in-medium broadening is largely compensated by Pauli blocking effects on the decay nucleon. The sensitivity of our results with respect to the in-medium widths of the higher lying resonances is comparatively small. As can be seen from Fig. 2, a reasonable fit is obtained up to incident photon energies of about 1 GeV. Beyond that the inclusion of further baryon resonances in both the \( \pi N \) and \( \rho N \) interactions as well as higher partial waves seems to be required. It is noteworthy that below the pion threshold some strength appears. This is nothing but the well-known ‘quasi deuteron’ tail above the giant dipole resonance, arising from pion-exchange currents. These are naturally included in our model.

In summary, we have shown that our earlier model for rho-meson propagation in hadronic matter \([4]\) allows for a consistent application at the photon point \((M^2=0)\). With additional improvements on both the two-pion selfenergy \( \Sigma_{\rho\pi\pi} \) (including the full 3-momentum dependence) and resonant \( \rho N \) contributions (including \( s \)-wave contributions as well as an improved version of the VDM) an acceptable description of total photoabsorption cross sections on both the proton and nuclei has been achieved with a rather limited number of parameters, thereby further constraining their actual values. This clearly increases the confidence in the model when applying it to calculate dilepton production as measured in relativistic heavy-ion collisions at various bombarding energies. Indeed, employing our model in a transport theoretical analysis of the CERN experiments (CERES and HELIOS-3) gives good agreement with the observed dilepton spectra \([8]\).

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TABLE I. Properties of the \( \rho BN \) vertices as derived from the interaction lagrangians, Eqs. (7) and (8); table columns from left to right: baryon resonance, relative angular momentum in the \( \rho N \) decay, spin-isospin factor (note that in its definition we have absorbed an additional factor of \( 1/2 \) as compared to table 2 in ref. [7]), partial decay width into \( \rho N \) as extracted from ref. [16], coupling constant as estimated from \( \Gamma^0_{\rho N} \) (for \( N(939) \) and \( \Delta(1232) \) we have indicated the values from the BONN potential [17] which uses somewhat harder form factors), coupling constant as actually used in our photoabsorption fit, in-medium correction to the total decay width.

| B         | \( I_{\rho N} \) | \( SI(\rho BN^{-1}) \) | \( \Gamma^0_{\rho N} [\text{MeV}] \) | \( \left( \frac{f^2_{\rho BN}}{4\pi} \right)_{\text{est}} \) | \( \left( \frac{f^2_{\rho BN}}{4\pi} \right)_{\text{fit}} \) | \( \Gamma^\text{med} [\text{MeV}] \) |
|-----------|-----------------|---------------------|---------------------|---------------------|---------------------|---------------------|
| \( N(939) \) | \( p \)        | 4                   | –                   | 4.68                | 5.8                 | 0                   |
| \( \Delta(1232) \) | \( p \)        | 16/9               | –                   | 18.72               | 23.2                | 15                  |
| \( N(1520) \) | \( s \)        | 8/3                | 24                  | 6.95                | 5.5                 | 250                 |
| \( \Delta(1620) \) | \( s \)        | 8/3                | 22.5                | 1.01                | 0.7                 | 50                  |
| \( \Delta(1700) \) | \( s \)        | 16/9               | 45                  | 1.2                 | 1.2                 | 50                  |
| \( N(1720) \) | \( p \)        | 8/3                | 105                 | 8.99                | 9.2                 | 50                  |
| \( \Delta(1905) \) | \( p \)        | 4/5                | 210                 | 17.6                | 18.5                | 50                  |
Figure Captions

**Figure 1**: Total photoabsorption cross section on the proton: full result of our fit (solid line), $\pi\pi$ 'background' (dashed line) as well as the three dominant $\rho N$ resonances $\Delta(1232)$, $N(1520)$ and $N(1720)$ (dashed-dotted lines). The data are taken from ref. [18].

**Figure 2**: Total photoabsorption cross sections on different nuclei: the solid line represents our full result while the dashed line denotes the contribution of the $\pi\pi$ background only, both calculated at a nuclear density $\bar{\rho}_N=0.8\rho_0$. The data are compiled from refs. [19–22].
\[ \sigma_{\gamma A}^{\text{abs}} / A \ [\mu b] \]

- Sn (Frascati)
- Pb (Frascati)
- Pb (Saclay)
- U (averaged)