Pions in the nuclear medium and Drell-Yan scattering

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Abstract

We investigate the modification of the pion-cloud in the nuclear medium and its effect on the nuclear Drell-Yan process. The pion’s in-medium self-energy is calculated in a self-consistent delta-hole model, with particle-hole contribution also included. Both the imaginary and real part of the pion’s and delta’s self-energy are taken into account and related through a dispersion relation assuring causality. The resulting in-medium pion light-cone momentum distribution shows only a slight enhancement compared to the one of the free nucleon. As a consequence the ratio of the cross-section for Drell-Yan scattering on nuclear matter and nucleonic target is close to unity in agreement with experiment.
I. INTRODUCTION

Recently there has been renewed interest in the role of the pion propagator in the nuclear medium for several reasons: First the ratio of spin-longitudinal and spin-transverse response functions below the quasi-elastic peak in a naive model is predicted to be much larger than unity, while experiment finds a ratio close to unity in $(\vec{p}, \vec{\ell})$ polarization transfer \[1\]. Second the observed ratio of cross sections for Drell-Yan scattering on nuclear and nucleonic targets is consistent with no excess pions present in the nucleus, i.e. no enhancement of the sea-quarks \[4\].

In the past the effect of medium modifications of the pion propagator has been investigated in detail by Ericson and Thomas \[3\] (in connection with the EMC effect) and Bickerstaff et al. \[4\]. The apparent absence of the pion-cloud enhancement in DY scattering was explained by Brown et al. \[3\] in terms of partial restoration of chiral symmetry and the associated decrease in masses of the nucleons and vector mesons in the nuclear medium. It was also pointed out \[3\] that, in a definite model considered, the correct normalization of the physical nucleon state, consisting of the bare nucleon and pion-cloud term, considerably reduces the effect of the pion-cloud enhancement.

The sensitivity of the pion-cloud to the pion-nucleon-nucleon ($\pi NN$) vertex cut-off $\Lambda$ was emphasized by Thomas \[7\]. Not only is the sea-quark content of the nucleon very sensitive to the value of the cut-off, the enhancement of the pion light-cone-momentum distribution in the nuclear medium also follows that pattern, even if the value of the cut-off is kept the same in the medium as in free space \[8\]. Another uncertainty concerns the values of the Migdal $g'$ parameters ($g'_{NN}, g'_{N\Delta}, g'_{\Delta\Delta}$), describing the effects of short-range repulsion. Taking large enough values of these parameters it is possible to reduce the medium enhancement of the pion distribution. However, it has not been reliably established that using such large values of the $g'$ parameters is physically justified.

Our approach in reexamining the pion-cloud enhancement is to perform a more ambitious computation of the pion light-cone-momentum distribution in the nuclear medium, using the
energy and momentum dependent pion self-energy and propagator. The latter are calculated in a recently developed self-consistent delta-hole model, including particle-hole states, and allowing for a width of the delta [9]. This model takes into account both the real and the imaginary parts of the pion’s and delta’s self-energy, in a way that assures causality and absence of unphysical poles. The Schwinger-Dyson equations (without vertex corrections) for the delta and the pion are solved self-consistently, while the nucleon is treated in a mean-field approximation.

The framework of the present approach (and that of Ref. [9]) is the effective field theory of hadrons, which is used consistently to calculate physical processes. The basic difference with previous approaches is the self-consistent treatment of the delta-hole contribution to the pion self-energy, including both its real and imaginary part. As already noted in Ref. [3] much of the pion-distribution enhancement came from the delta-hole term through its coupling to particle-hole term by $g'_{N\Delta}$. However, that result was based on a crude approximation for the delta-hole contribution, including only its real part (violating causality), with a momentum dependence coming only from the p-wave nature of the coupling and the form factor.

For the practical calculation of the pion light-cone momentum distribution in isospin symmetric nuclear matter (at temperature $T = 0$) we compare two approaches. First in section II we present a calculation of the pion light-cone-momentum distribution, based on explicit summation of the relevant diagrams involving the dressed pion propagator, and using an explicit integration over the initial and final nucleon momenta.

Second, in section III we consider the computation of the same quantity, but based on the nuclear-matter response function. While somewhat simpler, this approach gives only an approximate (however, for typical nuclear densities very close) expression for the pion distribution. This happens since the calculation of the imaginary part of the pion self-energy involves the full phase-space factors for the two nucleons [10], while such momentum dependent factor is not present for the incoming nucleon in the calculation involving explicit summation of relevant diagrams. However, the response-function approach takes into account also the process when a nucleon becomes a delta plus a pion, if the pion self-energy
includes the delta-hole contribution. We checked numerically that the two approaches give practically the same result. The ratio for Drell-Yan scattering on an isospin-symmetric nuclear target and the deuteron is presented in section IV. Numerical results, a discussion and comparison with other computations and experimental results is the subject of section V.

II. PION DISTRIBUTION IN THE NUCLEAR MEDIUM

It was noticed long time ago that the pion cloud gives a scaling contribution to the deep inelastic lepton scattering on the free nucleon [11]. The nucleon wave function can be schematically expressed as

\[
|N{>}_{\text{phys}} = \sqrt{Z}|N{>}_{\text{bare}} + \alpha|N{}\pi{>} + \beta|\Delta{\pi{>} + \cdots}
\]  

(1)

In this approach the light-cone momentum distribution of a quark with flavor \( f \) in a proton can be written as \((B = N, \Delta)\)

\[
q_f(x) = Zq_f,bare(x) + \sum_{B,i} c_i \left[ \int_x^1 \frac{dy}{y} f_{B_i/N}(y) q_{f,B_i,bare}(x/y) + \int_x^1 \frac{dy}{y} f_{\pi_i/N}(y) q_{\pi_i}(x/y) \right].
\]  

(2)

Here \( c_i \) (\( i \) labels the charge states) are the appropriate isospin Clebsch-Gordan coefficients. Note that following [12] we use the \( Z \) factor, the bare nucleon probability, as a renormalization of the bare nucleon only, not as an overall normalization of the whole right-hand side. In Refs. [5,6] the Sullivan contribution was also multiplied with the wave function renormalization factor. To see the relation we note that in quantum field theory one can write

\[
|N{>}_{\text{phys}} = \sqrt{Z'}(|N{>}_{\text{bare}} + g_0^0 |N{\pi{>}),}
\]

where \( g^0 \) denotes the bare coupling. The latter is related to the physical one (in lowest order) by \( g = \sqrt{Z'} g^0 \), i.e. \( Z' = (1 + \int dy f_{\pi/N}(y))^{-1} \). The prescription used in the present paper (see also Ref. [13]) is consistent with the standard nuclear physics definition of the \( \pi NN \) coupling constant derived from NN interaction at large distances. Of course the basic assumption is that we can restrict ourselves to only one pion in the air. For the form-factors we used in present calculation (whose conclusions are not affected by some variation of the cut-off) the value of \( Z \approx 0.6 - 0.7 \) suggests that this is a good approximation.
The pion light-cone momentum distribution, \( f^{\pi/N}(y) \), may contain either a nucleon or a delta final state i.e. \( f^{\pi/N}(y) = f^{\pi N/N}(y) + f^{\pi \Delta/N}(y) \). Let us first consider the nuclear final state. For the free nucleon the well-known result \[1\] for the pion distribution is

\[
f^{\pi N/N}(y) = \frac{3g_{\pi NN}^2}{16\pi^2} y \int_{M^2 y^2/(1-y)}^{\infty} dt \frac{|F_{\pi NN}(t)|^2 t}{(t + m_\pi^2)^2},
\]

(3)

Here \( y = \frac{k_0 + k_3}{M} \) gives the pion light-cone momentum fraction, with \( M \) being the physical mass of the nucleon (as an arbitrary but convenient scale), \( F_{\pi NN}(t) \) is the \( \pi NN \) vertex form-factor, while \( g_{\pi NN} \) is the \( \pi^0 NN \) coupling. The free-pion propagator, \( D_{\pi 0}^0 \), appears in the above expression in the form \((t + m_\pi^2)^{-1}\).

In the literature various prescriptions for the \( \pi NN \) form-factor have been proposed. Instead of the covariant formalism, with the pion taken off-shell and the final nucleon on-shell as used above, in Ref. \[12\] and also in work of the Jülich group, Ref. \[3\], old-fashioned perturbation theory in the infinite momentum frame was used. In that formulation the pion and final nucleon are taken to be on the mass shell (however, energy is not conserved in the interaction vertex) and the form-factor is conveniently taken a function of the center-of-mass energy of the intermediate \( \pi N \) system, \( s = (p' + k)^2 \). In this case the probability to find a baryon in nucleon with momentum fraction \( y \) is equal to the probability of finding a meson in the nucleon with momentum fraction \( 1 - y \), which is not the case in the covariant approach with the form-factor depending only on \( t \equiv -k^2 \).

The impracticality of making in-medium (especially self-consistent) calculations in the infinite-momentum frame makes us adopt the covariant approach. Since the covariant and the old-fashioned perturbation approach are equivalent (apart from the effect of form factors), we assume the conditions which should be satisfied by the first and second moments of the two distributions for flavor-charge and momentum conservation can be assured (at least approximately), also in the former scheme by a suitable \( p^2 \) dependence in the form-factor. In that case (see section IV) it is enough to consider only the change of \( f^\pi(y) \) due to medium effects, which is the main subject of the present work. Furthermore, the pion distributions obtained in the two schemes are very similar if in the covariant approach a
monopole form-factor is used as in our calculation.

It was confirmed [14] that dressing the pion in vacuum gives practically negligible contribution. That is not the case, however, in nuclear medium where delta-hole and particle-hole states play an important role. Consequently, we want to base our calculation on the dressed-pion self-energy and propagator, depending on $k_0$ and $k ≡ |\vec{k}|$. The types of diagrams which contribute are shown in Fig. 1. They can be summed in RPA in two steps:

(i) first, the pion rescattering (without the $g'$ coupling to external nucleon) leads to the dressed pion propagator $D_\pi(k_0, k)$, $D_\pi^{-1} = (D_\pi^0)^{-1} - k^2\Pi$, where the pion self-energy $k^2\Pi$ contains both the delta-hole and particle-hole contributions [15]:

$$\Pi = \Pi_{NN} + \Pi_{N\Delta} + 2g'_{N\Delta}\Pi_{NN}\Pi_{N\Delta}$$

Here $\Pi_{NN}$ and $\Pi_{N\Delta}$ are the iterated particle-hole and delta-hole self-energies divided by the common factor $k^2$ and the arguments $k_0, k$ are left out for brevity:

$$\Pi_{NN} = \frac{\Pi_{0NN}(k_0, k)}{1 - g'_{NN}\Pi_{0NN}(k_0, k)}$$

$$\Pi_{N\Delta} = \frac{\Pi_{0N\Delta}(k_0, k)}{1 - g'_{N\Delta}\Pi_{0N\Delta}(k_0, k)}$$

with $\Pi_{0NN}$ and $\Pi_{0N\Delta}$ being the one-loop particle-hole and delta-hole self-energies divided by $k^2$.

(ii) second, the inclusion of diagrams in which the pion first couples with $g'$ has the effect of replacing $D_\pi$ by

$$\tilde{D}_{\pi N}(k_0, k) = D_\pi(k_0, k) \left(1 + \frac{X_1(k_0, k) + X_2(k_0, k)}{X_3(k_0, k)}\right)$$

with

$$X_1(k_0, k) = g'_{NN}\Pi_{NN} \left(1 + g'_{N\Delta}\Pi_{N\Delta}\right),$$

$$X_2(k_0, k) = g'_{N\Delta}\Pi_{N\Delta} \left(1 + g'_{NN}\Pi_{NN}\right),$$

$$X_3(k_0, k) = \left[1 - (g'_{N\Delta})^2\Pi_{NN}\Pi_{N\Delta}\right].$$

As a result in nuclear matter $D_0^\pi$ in Eq. (3) is replaced by $\tilde{D}_{\pi N}$ and in addition the integration limits are adjusted. The nuclear matter ground state is approximated by a mean field, i.e.
the nucleon that emits the pion has an initial momentum \( p \leq p_F \) and to satisfy the Pauli principle a final momentum \( p' > p_F \). The effect of binding is included in an effective mass and a shift of energy approximation, i.e. \( E = \sqrt{p^2 + M_*^2} + c_0 \). The values of the parameters in the above expression are determined from nuclear matter models and the requirement for leading to a reasonable value of the Fermi energy.

As a result the pion light-cone distribution (per nucleon) in isospin symmetric nuclear matter can be expressed as

\[
    f^{\pi/A}_N(y) = \frac{9yg_{\pi NN}^2M_*}{4(2\pi p_F)^3}\int_{-p_F}^{p_F} dp_3 \sqrt{\gamma_F-p_3^2} \int_0^{p'_{\perp\max}} dp'_{\perp} \int_0^{p'_{\perp}} dp'_{\perp} \times \int_0^{2\pi} d\theta k^2 F_{\pi NN}(k) \frac{1}{z'} |\tilde{D}_{\pi N}|^2.
\]  

(11)

In the above \( p'_{\perp\max} = \sqrt{2z'M_*E_F - (z^2M_*^2 + M_*^2)} \), \( \theta \) is the angle between \( \vec{p}_{\perp} \) and \( \vec{p}'_{\perp} \), \( \vec{k} = \vec{p} - \vec{p}' \) (the three-momentum of the pion), and

\[
    z' = \frac{1}{M_*} \left( -M y + p_3 + \sqrt{M_*^2 + p_3^2 + p_{\perp}^2} \right) = \frac{1}{M_*} \left( p_3' + \sqrt{M_*^2 + p_*^2} \right),
\]  

(12)

\[
    k_0 = \sqrt{M_*^2 + p_{3}^2 + p_{\perp}^2} - \sqrt{M_*^2 + p_{\perp}^2 + p_{3}^2},
\]  

(13)

We note that \( f^{\pi/A}_N(y) \) has the correct non-relativistic low-density limit, \( \lim_{p_F \to 0} f^{\pi/A}_N(y) = f^{\pi/N}_N(y) \), which is not the case in Ref. [4].

We also consider diagrams with a delta replacing the nucleon in the final state, in which case the relativistic analog of expression (3) reads

\[
    f^{\pi\Delta/N}(y) = \frac{g_{\pi NN\Delta}M}{24\pi^2 M_\Delta^2} y \int dp_3 p_{\perp} dp_{\perp} \frac{1}{p_0} (M_*^2 + t)^2 (M_*^2 + t) \frac{|F_{\pi\Delta}(t)|^2}{(t + M^2)^2} \rho_{\Delta}(p_0, p),
\]  

(14)

where \( M_\pm = M_\Delta \pm M \). The above expression takes into account the width of the delta through its (vacuum) spectral-function \( \rho_{\Delta} \). The energy of the delta is determined by the light-cone-momentum of the pion as \( p_0 = M(1 - y) - p_3 \) and \( t = p_{3}^2 - M_*^2 y^2 - 2M yp_3 \).

We used the nonrelativistic limit of this expression, which amounts to replacing the term \( (M_*^2 + t)^2 (M_*^2 + t) \) by the first term of low-momentum expansion: \( 16M^3M_\Delta k^2 \), where \( k \) is the magnitude of the pion’s three-momentum. In isospin-symmetric nuclear medium the pion light-cone distribution (per nucleon) takes a form analogous to (11):
\[ f_{\pi/A}^\Delta(y) = \frac{y M g_{\pi NN}^2 g_{\pi N\Delta}^2}{2 (\pi p_F)^2 M_A^2} \int_{-p_F}^{p_F} dp_3 \int_0^{\sqrt{p_F^2 - p_3^2}} p_\perp dp_\perp \int_{-\infty}^{\infty} dp'_3 \int_0^{\infty} p'_\perp dp'_\perp \times \int_0^{2\pi} d\theta k^2 F_{\pi N\Delta}^2(k) \frac{1}{p_0} \rho_{\Delta}(p'_0, p') |\bar{D}_{\pi\Delta}|^2, \] 

where

\[ \bar{D}_{\pi\Delta}(k_0, k) = \frac{D_{\pi}(k_0, k)}{1 - (g'_{\pi N\Delta})^2 \Pi_{NN}\Pi_{\pi N\Delta}} \sum_{i=1}^{3} X_{\Delta i}(k_0, k), \]

with

\[ X_{\Delta 1}(k_0, k) = g'_{\pi N\Delta} \Pi_{NN} (1 + g'_{\pi N\Delta} \Pi_{\pi N\Delta}), \]

\[ X_{\Delta 2}(k_0, k) = g'_{\pi N\Delta} \Pi_{\pi N\Delta} (1 + g'_{\pi N\Delta} \Pi_{NN}), \]

\[ X_{\Delta 3}(k_0, k) = \left[ 1 - (g'_{\pi N\Delta})^2 \Pi_{NN}\Pi_{\pi N\Delta} \right]. \]

In Eq. (15) \( p' = \sqrt{p_\perp^2 + p'_3}, p_0' = \sqrt{M_\pi^2 + p_3^2 + c_0 - M_y + p_3 - p'_3} \) and \( k_0 = \sqrt{M_\pi^2 + p_3^2 + p_\perp^2 + c_0 - p'_0} \). Obviously the total pion light-cone momentum distribution in the medium is the sum of the contributions with a nucleon and with a \( \Delta \) in the final state,

\[ f_{\pi/A}^\pi(y) = f_{\pi/N}^\pi(y) + f_{\pi/\Delta}^\pi(y). \]  

### III. PION DISTRIBUTION AND THE NUCLEAR RESPONSE FUNCTION

The spin-isospin nuclear response function essentially involves the imaginary part of the (iterated) pion self-energy and since diagrams of the same type appear in deep inelastic scattering off the pion emitted by the nucleon, the pion distribution in the medium can be expressed in terms of the response function [4].

We want to clarify the relation of these two quantities for the relativistic approach. Even though we are using nonrelativistic nucleon and delta propagators (as well as \( \pi NN \) and \( \pi N\Delta \) vertices), our use of relativistic kinematics necessitates the discussion, especially for establishing the correct low-density limit, as well as for getting the factors of nucleon’s physical and effective mass in the expression correctly (the latter does not appear in previous treatments [4][6], while in Ref. [3] a peculiar rescaling is introduced).
As a first step we consider the relation of the imaginary part of the particle-hole pion self-energy and the relevant part of the deep inelastic diagram shown in Fig. 2. Iterating the pion self-energy (with free-pion propagators in between) to get the full response function will not change the established relation (in the kinematic region of interest), since the off-shellness of the emitted pion means that cutting the free-pion propagator gives zero contribution. The square of the absolute value of the diagram shown in Fig. 2b gives:

\[
L = -2g_{\pi NN}^2 \int \frac{d^3p \, d^3p'}{(2\pi)^6 2\sqrt{M_*^2 + \vec{p}'^2}} \text{Tr}[\gamma_5(\not{p_*} + M_*) \gamma_5(\not{p'_*} + M_*)] \Theta(p_F - |\vec{p}|) \Theta(|\vec{p}'| - p_F)
\]

\[
= -4g_{\pi NN}^2 \int \frac{d^3p \, d^3k}{(2\pi)^6 2\sqrt{M_*^2 + (\vec{p} - \vec{k})^2}} (k_0^2 - \vec{k}^2) \Theta(p_F - |\vec{p}|) \Theta(|\vec{p} - \vec{k}| - p_F),
\]  

(21)

where the isospin degeneracy factor of 2 is included, \( p_* \) denotes a \( p \) whose zeroth component has a mean-field shift, \( \Theta \) is the step function and \( k_0 \equiv \sqrt{M_*^2 + \vec{p}'^2 - \sqrt{M_*^2 + (\vec{p} - \vec{k})^2}} \).

On the other hand, the contribution of the mean-field nucleons to the imaginary part of the pion self-energy in isospin symmetric nuclear matter can be written as \[ 17 \]

\[
\text{Im}\Pi(k_0, k) = -2 \times 2g_{\pi NN}^2 \int \frac{d^4p}{(2\pi)^4} \text{Tr}[\gamma_5(\not{p_*} + M_*) \gamma_5(\not{p_*} - \vec{k} + M_*)] \delta((p_0 - c_0)^2 - \vec{p}'^2 - M_*^2)
\]

\[
\times \delta((p_0 - c_0 - k_0)^2 - (\vec{p} - \vec{k})^2 - M_*^2) \Theta(p_F - |\vec{p}|) \Theta(|\vec{p} - \vec{k}| - p_F)
\]

\[
= -4g_{\pi NN}^2 \int \frac{d^3p}{(2\pi)^6 2\sqrt{M_*^2 + \vec{p}'^2}} (k_0^2 - \vec{k}^2) \delta((p_0 - k_0)^2 - (\vec{p} - \vec{k})^2 - M_*^2)
\]

\[
\times \Theta(p_F - |\vec{p}|) \Theta(|\vec{p} - \vec{k}| - p_F).
\]

(22)

Integrating the above expression over \( k_0 \) and \( \vec{k} \) we obtain

\[
\int \frac{d^4k}{(2\pi)^4} \text{Im}\Pi(k_0, k) = -2g_{\pi NN}^2 \int \frac{d^3p \, d^3k}{(2\pi)^6 2\sqrt{M_*^2 + (\vec{p} - \vec{k})^2}} \frac{k_0^2 - \vec{k}^2}{2\sqrt{M_*^2 + \vec{p}'^2}}
\]

\[
\times \Theta(p_F - |\vec{p}|) \Theta(|\vec{p} - \vec{k}| - p_F),
\]

(23)

with \( k_0 \equiv \sqrt{M_*^2 + \vec{p}'^2 - \sqrt{M_*^2 + (\vec{p} - \vec{k})^2}} \) on the right-hand side. Comparison of expressions (21) and (23) shows that in general one cannot express (21) in terms of (23) because of presence of the momentum-dependent factor \( 1/\sqrt{M_*^2 + \vec{p}'^2} \) in Eq. (23). However, if we approximate \( \sqrt{M_*^2 + \vec{p}'^2} \) with a \( p \)-independent constant (say \( M_* \)), which is an excellent approximation for baryon densities not much larger than the saturation density, the following relationship emerges:
\[ L = -4M_\ast \int \frac{d^4k}{(2\pi)^4} \text{Im} \, \Pi(k_0, k). \] (24)

Apart from the factor \(1/\sqrt{M^2_\ast + \vec{p}^2}\) whose momentum dependence in the present context does not play an important role, there is another momentum-dependent factor in expressions (21) and (23): \(1/\sqrt{M^2_\ast + (\vec{p} - \vec{k})^2}\), whose presence is vital for correct low-density limit. In our derivation of expression (22) this factor appears automatically since we use the fully relativistic nucleon propagators in the particle-hole loop. However, by using the non-relativistic limit (as, for example in Ref. [16]), this term would be replaced by \(1/M_\ast\). This leads then to an incorrect low-density limit of the pion distribution, since the Sullivan expression (3) for the free-nucleon is based on relativistic calculation and contains the momentum dependence of the mentioned term (with \(\vec{p} = 0\)).

Taking into account the more complicated diagrams means replacing \(\Pi\) by \(\Pi + \Pi D^0_\pi k^2 \Pi + \Pi D_\pi^0 k^2 D^0_\pi k^2 \Pi + \cdots = \Pi(1 - D^0_\pi k^2 \Pi)^{-1}\), where \(D^0_\pi\) is the free-pion propagator. Also, when integrating over the pion momentum we have to take into account the constraint on the pion light-cone momentum fraction \(y = (k_0 + k_3)/M\). This can be imposed by an appropriate delta-function, which should be inserted in both expressions (21) and (23). Performing the \(k_3\) integration and dividing by the total number of nucleons in unit volume we obtain the expression for the pion light-cone momentum distribution

\[ \tilde{f}^{\pi/A}(y) = -\frac{9yMM_\ast}{4\pi p_F^3} \int k_\perp dk_\perp d\omega \text{Im} \left( \frac{k^2 \Pi(\omega, k)}{1 - D^0_\pi k^2 \Pi(\omega, k)} \right) |D^0_\pi|^2, \] (25)

where (in the general case) the pion self-energy which contains both the delta-hole and particle-hole contributions is given by Eq. (4). In \(\Pi\) and \(D^0_\pi\) the argument \(k\) is given by \(k = \sqrt{k_\perp^2 + k_3^2}\), where \(k_3 = My + \omega\). The integration region here is from zero to infinity for both variables, since conditions on incoming nucleon (below the Fermi sea) and outgoing nucleon (above the Fermi sea) are automatically satisfied by the correct calculation of the imaginary part of \(\Pi\), if the nucleons are treated as a mean field.

Using the expression for the nucleon density \(\rho = 2p_F^3/3\pi^2\) and introducing the longitudinal response-function through
expression (23) can be written in the usual form

\[
\tilde{f}_{\pi/A}(y) = \frac{3g_{\pi NN}^2y}{16\pi^2\rho} \int dk^2 k^2 F_{\pi NN}(k)^2 \int d\omega \frac{R_L(\omega, k)}{(\omega^2 - k^2 - m_\pi^2)^2},
\]

(27)

Note that the form factor and the coupling have been removed from the imaginary part of the (iterated) pion self-energy in the definition (26) of the response function.

We confirmed by numerical computation that Eqs. (11) and (25) give practically identical results (thus validating the approximation used to obtain eq.(24)), if the pion self-energy contains only the particle-hole contribution (assuring the final state with only nucleon and not delta, corresponding to diagrams summed in Eq. (11)). Inclusion of the self-consistent delta-hole contribution to the pion self-energy takes into account also the diagrams where a delta replaces the nucleon in the final state. This corresponds to taking into account the contribution of diagrams summed in eq. (15), but also some other higher-order diagrams, which appear because the delta is itself dressed by pion-nucleon loops (where the pion is also dressed). Numerically this shows up as a very small difference (at values \( y < 0.2 \)) of the distribution (27) and the sum of eqs.(11) and (15). The smallness of this difference indicates that it is reasonable to expect that contributions of other higher-order diagrams left out from the calculation (for example, the pion-nucleon loop term of the nucleon self-energy) will not affect the pion distribution significantly.

IV. DRELL-YAN SCATTERING

The cross-section of the Drell-Yan process \( a + b \rightarrow \bar{l}l \) can be expressed as

\[
\frac{d\sigma^{ab}}{dxf_1dx_2} = \frac{4\pi\alpha^2}{9sx_1x_2K(x_1, x_2)} \sum_f e_f^2 \left[q_f^a(x_1)\bar{q}_f^b(x_2) + \bar{q}_f^a(x_1)q_f^b(x_2)\right],
\]

(28)

where \( s \) is the center-of-mass energy squared and the summation of products of quark and antiquark distribution functions is over flavors. The factor \( K(x_1, x_2) \) takes into account higher order QCD corrections and is of the order 1.5. The values of \( x_1, x_2 \) are extracted...
from experiment via the invariant mass of the lepton pair. We are interested in the ratio of the cross-sections for proton-nucleus and proton-deuteron scattering:

\[ R_{A/d} \equiv \frac{2 \, d\sigma^{pA}/dx_1dx_2}{A \, d\sigma^{pd}/dx_1dx_2}, \]  

(29)

where \( A \) denotes both the nucleus and its nucleon number.

Since our in-medium computation is performed for the case of isospin symmetric nuclear medium, we specialize to the ratio of differential cross-sections on a nucleus consisting of equal numbers of protons and neutrons, and on a deuteron (for which we consider medium effects negligible)

\[ R_{A/d} = \frac{\sum_f e_f^2 \left\{ q_f^p(x_1) \left[ q_{f/A}^p(x_2) + q_{f/A}^n(x_2) \right] + q_f^n(x_1) \left[ q_{f/A}^p(x_2) + q_{f/A}^n(x_2) \right] \right\}}{\sum_f e_f^2 \left\{ q_f^p(x_1) \left[ q_{f/A}^p(x_2) + q_{f/A}^n(x_2) \right] + q_f^n(x_1) \left[ q_{f/A}^p(x_2) + q_{f/A}^n(x_2) \right] \right\}}. \]  

(30)

To relate the quark distribution of the bound nucleon to that of the free nucleon we start from (2), where for brevity we write out only nucleon terms:

\[ q_f^p(x) = Z q_{f/bare}^p(x) + \frac{1}{3} \int_x^1 dy f^{N/N}(y) \left[ q_{f/bare}^{p,N}(x/y) + 2 q_{f/bare}^{n,N}(x/y) \right] 
+ \frac{1}{3} \int_x^1 dy f^{\pi/N}(y) \left[ q_{f/0}^{p,\pi}(x/y) + 2 q_{f/0}^{n,\pi}(x/y) \right], \]  

(31)

with

\[ Z \equiv 1 - \int_0^1 dy f^{N/N}(y) = 1 - \int_0^1 dy f^{\pi/N}(y), \]  

(32)

where the last equality in (32) is a requirement for flavor-charge conservation. Similarly, the starting expression for the quark distribution in a nuclear proton is

\[ q_f^p(x) = Z_A q_{f/bare}^{p,A}(x) + \frac{1}{3} \int_x^A dy f^{N/A}(y) \left[ q_{f/bare}^{p,A}(x/y) + 2 q_{f/bare}^{n,A}(x/y) \right] 
+ \frac{1}{3} \int_x^1 dy f^{\pi/A}(y) \left[ q_{f/0}^{p,\pi}(x/y) + 2 q_{f/0}^{n,\pi}(x/y) \right], \]  

(33)

where we specialized to the case of isospin-symmetric nuclear matter. Adding and subtracting the expression (31) to expression (33) and repeating the same procedure for the neutron (since only the sum of these two terms is relevant in isospin-symmetric medium) we get
\[\bar{q}_f(x) + \bar{q}_f^p(x) = q_f^p(x) + q_f^{p,bare}(x) \int_0^1 f^{N/N}(y) dy - \int_x^1 \frac{dy}{y} f^{N/N}(y) q_f^{p,bare}(x/y) - q_f^{p,bare}(x) \int_0^A f^{N/A}(y) dy + \int_x^A \frac{dy}{y} f^{N/A}(y) q_f^{p,bare}(x/y) + (p \to n) + 2 \frac{\int_x^A dy}{y} [f^{\pi/A}(y) - f^{\pi/N}(y)] [q_f^{\pi^0}(x/y) + q_f^{\pi^+}(x/y) + q_f^{\pi^-}(x/y)]. \tag{34}\]

To proceed without approximation a fit of the bare structure functions \(q_f^{p,bare}(x)\) would be required, based on expression (31) and the experimentally extracted \(q_f^p(x)\). This is clearly beyond the scope of present work and in the following we present arguments that the expression used in Refs. [3,18,8] for the pion contribution to the change of parton distribution:

\[\delta q_f^p(x) = \int_x^A dy \frac{dy}{y} \delta f^{\pi}(y) q_f^p(x/y), \tag{35}\]

(where \(\delta f^{\pi}(y) \equiv f^{\pi/A}(y) - f^{\pi/N}(y)\) and isospin factors are not shown) is a good approximation for antiquark distributions, the change of which is probed by the Drell-Yan pair production.

We argue that for antiquark distributions at small \(x\) it is a good approximation to neglect the difference of the second and third term on the right-side of Eq. (34), as well as the difference of the fifth and fourth term. The reason is that \(f^{N/N}(y)\) and \(f^{N/A}(y)\) are very small for \(y < 0.3\) (since the pion distributions are negligible for \(y > 0.7\)), thus if \(x < 0.3\) (the region where antiquark distributions are significant), the zero lower limit of integrals can be safely shifted to \(x\). Taking into account that antiquark distributions at small \(x\) behave as \(1/x\) (and that \(x/y\) is also small in the \(y\) region of most significant contribution), cancellation of the considered terms follows. A much simpler argument in favor of expression (35) is based on the assumption that the nucleon’s entire antiquark sea can be attributed to its virtual meson cloud which, however, has not been confirmed [13]. Since valence-quark distributions increase with decreasing \(x\) slower than \(1/x\) in the small \(x\) region, the above approximation may be less good in that case, but the effect on the Drell-Yan process of a relatively small change in quark distribution is not significant. The sum of the nuclear proton and neutron (anti)quark distribution of flavor \(f\) thus becomes
\[ \tilde{q}_f^p(x) + \tilde{q}_f^n(x) = q_f^p(x) + q_f^n(x) + \frac{2}{3} \int_x^A \frac{dy}{y} \left[ f^{\pi/A}(y) - f^{\pi/N}(y) \right] \left[ q_f^{\pi^0}(x/y) + q_f^{\pi^+}(x/y) + q_f^{\pi^-}(x/y) \right]. \] (36)

Performing a convolution on the nucleon part to take into account the Fermi motion of the nucleons we obtain

\[ q_f^{p/A}(x) + q_f^{n/A}(x) = \int_x^A \frac{dz}{z} f_{\text{Fermi}}^N(z) \left[ q_f^p(x/z) + q_f^n(x/z) \right] \]

\[ + \frac{2}{3} \int_x^A \frac{dy}{y} \left[ f^{\pi/A}(y) - f^{\pi/N}(y) \right] \left[ q_f^{\pi^0}(x/y) + q_f^{\pi^+}(x/y) + q_f^{\pi^-}(x/y) \right]. \] (37)

We note that the above expression satisfies flavor-charge conservation by construction and the momentum conservation sum-rule (with only pions included)

\[ \int_0^A zdz f_{\text{Fermi}}^N(z) + \int_0^A ydy \left[ f^{\pi/A}(y) - f^{\pi/N}(y) \right] = 1. \] (38)

turns out to be satisfied with values of \( M_* \) and \( c_0 \) based on nuclear matter models with accuracy better than 1\% (see next section). Heavier mesons in general give much smaller contributions than the pion \[12\] and would not affect significantly the sum-rule.

One can treat the contribution of the delta in the same way, an important difference being that even in the isospin symmetric case the structure function of the delta appears in the structure function of the in-medium nucleon.

To be consistent with mean-field treatment of nucleons which was used for calculating the pion’s self-energy, for the nucleon light-cone momentum distribution \( f_{\text{Fermi}}^N \) we use the Fermi gas model with mean-field corrections for mass and energy, \( E(p) = \sqrt{M_*^2 + p^2 + c_0} \).

The expression, which takes into account the flux factor, is the one obtained by Birse \[19\]:

\[ f_{\text{Fermi}}^N(z) = \frac{3}{4e^3} \left[ \epsilon^2 - (z - \eta)^2 \right] \Theta(\epsilon - |z - \eta|), \] (39)

where \( \epsilon \equiv p_F/M, \eta \equiv (\sqrt{M_*^2 + p_F^2} + c_0)/M \) and \( \Theta(x) \) is the step function.

V. NUMERICAL RESULTS AND DISCUSSION

First we examine the effect of Pauli blocking and Fermi motion on the pion light-cone momentum distribution, \( f^{\pi A}(y) \) by comparing in Fig. 3 the free-nucleon case (full line) to the
in-medium one, but the latter calculated with the free-pion propagator. The pion-nucleon-nucleon coupling has the usual value $g_{\pi NN} = 13.5$.

For the actual form of the form factor various parametrizations have been used, e.g. $F(t) = (\Lambda^2 - m^2)^n$ with $n = 1$ (monopole) or $n=2$ (dipole), and also exponential forms. The corresponding values of the cut-off, $\Lambda$, can be related, see Ref. [13]. Previous studies [5,12,13,20,21] of observables related to antiquark distributions obtained cut-offs in the range $0.6-0.8$ GeV, when translated to monopole form, which is much smaller than the one used in the Bonn potential. We performed computations with a monopole form using $\Lambda_{\pi NN} = 0.7$ and $0.8$ GeV, and also exponential form with $1$ GeV. While the effect of changing the cut-off $\Lambda_{\pi NN}$ is significant for both the free nucleon case and in the medium, their difference shows little sensitivity to this change. This is in contrast to Ref. [8] where a decrease in the cut-off produced a decrease in the enhancement too. In the following for all presented results we used a monopole form-factor with $\Lambda_{\pi NN} = 0.8$ GeV, both in vacuum and in the medium.

For the values of $M_*$ and $c_0$ we turn to nuclear matter models. While the simplest form of the Walecka model [23] gives at saturation density $M_* \approx 0.7$ GeV, more elaborate treatments tend to increase the effective mass to $0.8-0.85$ GeV at saturation density, leading to better agreement with observables [24,25]. Since we use for the Fermi momentum $p_F = 0.256$ GeV (i.e. slightly below the saturation density), we take $M_* = 0.85$ GeV. Adopting a smaller value for the nucleon’s effective mass decreases the in-medium pion enhancement, as noticed in Ref. [6]. The value of the energy shift $c_0$ should be such that a reasonable value of the Fermi energy is obtained [19], thus we take $c_0 = 0.04$ GeV, giving $E_F - M = -11$ MeV. For these values of $M_*$ and $c_0$ the momentum-conservation sum-rule (38) is satisfied with accuracy better than $1\%$ (for values of the $g'$ parameters in the region from 0 to 0.4), still leaving a little space for the effect of other mesons.

In Fig. 3 we see a reduction due to Pauli blocking and practically no broadening for the effective mass of the nucleon $M_* = 0.85$ GeV (short-dashed line). We note, however, that the decrease of the effective mass of the nucleon contributes to the decrease and narrowing of the pion distribution (as pointed out in Ref. [6]), as can be seen also from Fig. 3, where
the results for $M_\pi = 0.939$ GeV and 0.7 GeV are also shown.

For the in-medium calculation we consider first the contribution from diagrams with a nucleon in the final state. The results for $f^\pi_A(y)$, Eq. (11), for different values of $g'$ parameters are shown and compared to the free-nucleon case in Fig. 4. For $g'_{NN}$ we used the values 0, 0.4 and 0.6, while the other two parameters are kept constant, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$. Increasing the $g'$ parameters in general leads to reduction in the enhancement of pion distribution, although for small values of $g'$ there can be some increase for certain $y$ values, as a consequence of corresponding enhancement in some kinematical regions [3]. For the in-medium calculation we keep the mass difference of the delta and the nucleon the same as in free space.

We also consider the contribution from the delta in the final state, $f^\pi_\Delta(y)$. In connection with it we make some remarks on the $\pi N\Delta$ coupling. In the past the cut-off for the $\pi N\Delta$ coupling was taken the same as for the $\pi NN$ (using SU(6)). More recently there is evidence that the former coupling is softer than the latter from several sources:

(i) The Juelich group has analyzed data for diffractive scattering $p + p \to \Delta^{++} + X$ in the one-pion exchange approximation,

(ii) Koepf et al. [13] pointed out that the empirical $\bar{u} - \bar{d}$ asymmetry can be explained in a meson cloud model only if $\Lambda_{\pi N\Delta}$ is softer than $\Lambda_{\pi NN}$ by about 100 MeV.

In the present approach [1] we have extracted $\Lambda_{\pi N\Delta}$ from a fit to $\pi N$ scattering phase-shift in the delta channel, using three parameters: mass of the delta, value of the coupling and cut-off, and obtained the value for the exponential form-factor of $\Lambda_{\pi N\Delta} = 0.38$ GeV (for dipole form-factor the value of the cut-off was 0.51 GeV). This value gives an excellent fit for the phase-shift for laboratory pion momentum up to 500 MeV. The $\pi N\Delta$ coupling used, $g_{\pi N\Delta}$, was 20 GeV$^{-1}$ (this implies an on-shell value of 14 GeV$^{-1}$, in good agreement with expectation based on the delta’s width), and the (bare) delta mass of 1.27 GeV. We remark that a similarly soft pion-nucleon-delta form-factor was extracted from pion-nucleon scattering in Ref. [22], where also a self-consistent dressing of the pion and the delta was performed. Since the $\pi N\Delta$ cut-off is rather small, we expect this contribution to be small.
This is indeed the case as one can see from Fig. 5. The pion distribution is nonzero only in the region of small $y$ ($y < 0.2$), and while for $g' = 0$ (here the $g'_{\Delta\Delta}$ plays the most prominent role) there is some enhancement over the free-nucleon case (the full line), its value is very small. To study the effect of the cut-off, we also performed computations with $\Lambda_{\pi N\Delta} = 0.5$ GeV (which gives a very poor fit to the pion-nucleon phase-shift) and confirmed the absence of medium enhancement in that case too. We note that using the spectral-function of the dressed delta in the medium presents numerical difficulties, since it is nonzero in the region where the real part of the pion propagator takes very large values. To obtain stable results we neglected the delta contribution if its spectral function was smaller than 0.2 GeV$^{-1}$ (this affects the spectral-function sum-rule negligibly).

The advantage of the expression (25) relating the response function to the pion light-cone-momentum distribution is that it incorporates both contributions with a nucleon as well as delta in the final state, if the pion self-energy contains the particle-hole and the delta-hole contributions. Numerical evaluation confirms that for $y > 0.2$ these results agree with those based on expression (11), while for small $y$ there is a small enhancement due to diagrams with the delta in the final state. The slight difference between results based on expression (25) and the sum of contributions (11) and (15) can be attributed to contributions of higher order diagrams present in the self-consistent delta self-energy. The results of free-space (which is approximated by a low-density calculation with $\rho = 0.1\rho_0$, where $\rho_0$ corresponds to nucleon density of $p_F = 0.256$ GeV) and in-medium computations, using the same parameters as for Fig. 4, are shown in Fig. 6.

As to the origin of the absence of significant in-medium enhancement of the pion cloud we point to the results of Ref. [9] for the pion spectral function. The basic feature of that model is the self-consistent treatment of the pion and the delta, with both real and imaginary parts of the self-energies taken self-consistently into account, thus assuring causality for self-energies and propagators. As a result the pion spectral function has a main maximum whose shift to smaller energies is much less pronounced than in previous, less elaborate treatments, an exception being only Ref. [22], whose results show similar behavior to those
of Ref. [9]. Since this shift forms the origin of the pion-cloud enhancement and comes from
the (negative) real part of the delta-hole contribution, we can conclude that this is where
the relevant difference between our pion’s self-energy and that of previous models’ rests.
The sensitivity of the pion enhancement to this term requires its careful treatment, which
is evidently missing in Refs. [3,16]. We emphasize once more that in Ref. [9] the real part
of the self-energy is always calculated from its imaginary part using a relevant dispersion
relation.

In the calculation of the structure functions of nuclear nucleons we do not take into
account the delta contribution. The reason is that we do not know the (bare) delta structure
functions which are needed as are the nucleon structure functions. The delta structure
functions are necessary even in the case of the isospin symmetric nuclear matter. However,
the above results for the pion light-cone distribution show that there can be no noticeable
medium enhancement.

We now turn to a discussion of the DY ratio, Eq. (30). In applying Eq. (37) we use the
quark distributions in the nucleon and the pion from Refs. [26], and [27], respectively. The
data of Ref. [2] cover a rather limited region of $x_2$ values around $x_2 = 0.2$. That means that
in the expression (37) the convolution in the first term on the right-hand side is relatively
unimportant and that in the integral in the second term rather large $y$ values are sampled,
typically $y \geq 0.3$ and larger. Therefore one probes rather large 3-momenta in the response
function $D(k_0, k)$, Eq. (7). From $k^2 > k_3^2$ and $k_3 = -k_0 + my$ one has $|k| > |k_0| + my$, i.e.
typically $k > 400\text{MeV}/c$. As a consequence Pauli blocking (which is relevant for small $k$ )
is not very effective; also the delta-hole component at $k_{0,\Delta}$ is relatively unimportant since it
involves still higher momenta, $k > my + |k_{0,\Delta}|$. This is rather different from the situation
for the longitudinal spin response function probed in e.g. the quasi-free (p,n) reaction.

Results for the ratio (30) of Drell-Yan cross sections are shown in Fig. 7, for different
values of the Migdal $g'_\text{NN}$ parameters. A general feature of all these plots is the decreasing
trend for larger values of $x_2$ ($x_2 > 0.3$), leading to values below unity. The enhancement
for smaller values of $x_2$ even with no Migdal correction is quite modest. Values of the $g'_\text{NN}$
around 0.3 lead to practically no enhancement.

Compared to Ref. [5] we still obtain a net medium DY enhancement by 10% for \( g' < 0.4 \) in the region \( 0.1 < x_2 < 0.3 \). This difference can be ascribed to the fact that in [5] an *ad hoc* Z-factor (see discussion below Eq. (3)) occurs on right-hand side of Eq. (37) which is not present in our approach. In addition in [5] the assumed density dependent in-medium reduction of the cut-off \( \Lambda \) in the \( \pi NN \) form factor appears the main mechanism to reduce the pionic enhancement. However we feel there is no justification for such a medium vertex renormalization.

For direct comparison with measured results of Ref. [2] we calculated the ratio of nuclear and nucleon cross-section for given values of \( x_2 \) and the condition \( x_1 > x_2 + 0.2 \), corresponding to the experimental cut-off (here we assume that the factor \( K(x_1, x_2) \) is constant). In accordance with the above discussion we verified numerically that the ratio in the given \( x_2 \) region is mainly sensitive to the pion distribution in the region \( y > 0.3 \). Only the region of \( x_2 \) around 0.3 shows some sensitivity to smaller \( y \) values. This means that the drop of the ratio below one for \( x_2 \) around 0.05 (if the effect is real), probably cannot be explained by a change of the pion distribution.

Based on our results we can conclude that the treatment of the delta, although seemingly of secondary importance because of the kinematical region of the delta-hole contribution, does play a significant role. Including both the real and imaginary parts of its self-energy in a self-consistent way eliminates a large part of the pion-distribution in-medium enhancement and leads to agreement with experimental results on Drell-Yan scattering for modest values of the \( g' \) parameters. The reduction in the in-medium pion enhancement is achieved without introducing a medium reduction of the pion-nucleon-nucleon vertex cut-off parameter or a renormalization factor to suppress the pion-cloud contribution.

Finally we note that in the present paper we have restricted ourselves to isospin symmetric nuclear matter. On the other hand some of the data involve nuclei like Fe and W, which have a neutron excess. It is an interesting question whether there are additional medium effects to the asymmetry \( \bar{u}(x) - \bar{d}(x) \) in case of neutron excess. We plan to address this
question in the future.

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FIGURES

FIG. 1. Types of diagrams involving the pion self-energy and dressed propagator (dashed line), used for calculating the pion distribution in the medium. The double line denotes either a nucleon or delta.

FIG. 2. The particle-hole pion self-energy (a), and the pion emission part of the diagram showing the deep-inelastic scattering off the nucleon’s pion cloud (b).

FIG. 3. Effect of Pauli blocking and Fermi motion on the pion distribution. The full line corresponds to to the free nucleon, the short-dashed line to nuclear matter with $p_F = 0.256$ GeV and $M_* = 0.85$ GeV, but with free-pion propagator. The long-dashed line corresponds to $M_* = 0.7$ GeV, while the dot-dashed line is for $M_* = 0.939$ GeV.

FIG. 4. Pion distribution in the nuclear medium with Fermi momentum $p_F = 0.256$ GeV and for the free nucleon. Full line is for the free nucleon, long dashed line for medium with $g'_{NN} = 0$, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$, short dashed for $g'_{NN} = 0.4$, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$ and dot-dashed for $g'_{NN} = 0.6$, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$.

FIG. 5. Delta contribution to the pion distribution in the nuclear medium with Fermi momentum $p_F = 0.256$ GeV and for the free nucleon. Full line is for the free nucleon, long dashed line for medium with $g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0$, short dashed for $g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$ and dot-dashed for $g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0.5$. The effective mass of the delta in the medium is $M^*_\Delta = 1.18$ GeV (assuring the same effective mass difference for the delta and nucleon in the medium as in the free space), and the $\pi N\Delta$ vertex cut-off $\Lambda_{\pi N\Delta} = 0.38$ GeV.

FIG. 6. Pion distribution in the nuclear medium, based on computation using the response-function approach, expression (27). Full line is for the low-density case with $p_F = 0.119$ GeV, corresponding to 1/10 of the density used for in-medium calculations. Long dashed line is for medium with $g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0$, short dashed for $g'_{NN} = 0.4$, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$ and dot-dashed for $g'_{NN} = 0.6$, $g'_{N\Delta} = g'_{\Delta\Delta} = 0.3$. 23
FIG. 7. Ratio \((30)\) of the Drell-Yan cross-sections. Full line is for \(x_1 = 0.3\), dashed line for \(x_1 = 0.4\) and dot-dashed line for \(x_2 = 0.5\). The Migdal parameters are: a) \(g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0\), b) \(g'_{NN} = 0.4, g'_{N\Delta} = g'_{\Delta\Delta} = 0.3\).

FIG. 8. Experimental results from Ref. \([2]\) compared to our calculation for different \(g'\) values. Full line is for \(g'_{NN} = g'_{N\Delta} = g'_{\Delta\Delta} = 0.4\), dashed line for \(g'_{NN} = 0.4, g'_{N\Delta} = g'_{\Delta\Delta} = 0.3\), dot-dashed line for \(g'_{NN} = 0.6, g'_{N\Delta} = g'_{\Delta\Delta} = 0.3\).
Fig. 2
Fig. 3
Fig. 4

Pion distribution vs. $y$
Fig. 5

Pion distribution vs. y
Fig. 6
Fig. 7a
Fig. 7b
Fig. 8