Nuclear Drell-Yan effect in a covariant model

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(Dated: December 11, 2013)

We investigate effects of nuclear medium on antiquark distribution in nuclei applying the results of a recently developed relativistically-covariant self-consistent model for the pion and the isobar. We take into account Fermi motion including Pauli blocking and binding effects on the nucleons and medium effects on the isobar and pion leading to modest enhancement of the pion light-cone-momentum distribution in large nuclei. As a consequence the Drell-Yan cross-section ratio with respect to the deuteron exceeds one only for small values of the light-cone momentum.

I. INTRODUCTION

Improving our understanding of the quark and gluon degrees of freedom of nucleons bound in nuclei necessitates further efforts in spite of numerous investigations and successes [1]. The planned new nuclear Drell-Yan (DY) scattering experiment SeaQuest at Fermilab (E-906) [2–4] strongly motivates recent advancements in analysis of nuclear parton distributions [5] as well as attempts directed at more reliable and accurate calculation of experimentally observable quantities. In line with the second objective is our aim to update and extend analyses of nuclear parton distributions [5] as well as attempts directed at more reliable and accurate calculation of experimentally observable quantities. In line with the second objective is our aim to update and extend analyses of nuclear parton distributions as well as attempts directed at more reliable and accurate calculation of experimentally observable quantities. In line with the second objective is our aim to update and extend analyses of nuclear parton distributions as well as attempts directed at more reliable and accurate calculation of experimentally observable quantities.

In conventional models based upon meson exchange nuclear binding of nuclear matter comes for about 50% from (virtual) pions present in the nucleus. Can we observe these? Some indication for pions is present in EMC effect enhancement around $x = 0.1$, which can be ascribed to the fact that pions (and heavier mesons) carry a fraction of the momentum sum rule. Can one see these pions more explicitly, e.g. in the form of an enhancement of anti-quarks in the nucleus? Anti-quarks can be probed directly in Drell-Yan scattering but previous experiments [6] within the experimental uncertainty of about 10% did not show a nuclear enhancement; results of calculations varied strongly.

In practice one can distinguish two main types of theoretical interpretations of the classical EMC effect: (i) in terms of nucleon as constituents which are bound but not modified in the medium, and (ii) in terms of off-shell nucleons with medium modified structure functions, e.g. through scalar and vector fields acting on the quarks. In the first category one has the non-relativistic models which use a computed spectral function that accounts for large removal energy (50 MeV in nuclear matter) and Fermi motion due to correlations. This approach can reproduce the observed slope of the reduction for $x < 0.5$ in the EMC ratio $2F_2^A(x)/AF_2^N(x)$ but not the behavior around the minimum in the ratio at $x = 0.8$. The latter seems to require ad-hoc off-shell effects [7]. In the second category [10, 11] one usually starts from the Walecka model in the mean-field approximation which has a small net binding effect (8 MeV per nucleon) and hence yields a very small EMC effect, and then one adds the effect of external scalar and vector fields. Since in the present study we are interested in the antiquark distribution for $x < 0.4$ we rely on the conventional convolution approach using a parameterized nucleon distribution, which has the two above mentioned parameters related to the removal energy $\eta$ and the Fermi momentum.

In Sec. II we study the antiquark distributions in the free nucleon and determine the off-shell $\pi NN$ and $\pi N\Delta$ form-factors which lead to good description of the isovector part of the proton antiquark distribution by the pion cloud. In Sec. III we turn to general discussion of the nuclear effects consisting of binding and Fermi motion of nucleons and modification of the pion cloud. Detailed consideration of the medium effects on the nucleon’s pion cloud follows in Sec. IV where expressions are derived for the pion light-cone distributions originating from the $\pi N$ and $\pi \Delta$ states. We emphasize the careful treatment of the in-medium delta baryon based on a complete relativistically covariant basis for its dressed propagator. Numerical results for the in-medium pion distribution and DY cross-section ratios for nuclear targets relative to the deuteron are presented and discussed in Sec. V. Finally, Sec. VI contains a summary of our results.

II. ANTIQUARKS IN FREE NUCLEONS

Before turning to the nuclear case we want to investigate whether the pion cloud approach we will use in the medium can reproduce the observed flavor asymmetry in the free nucleon. The distribution of anti-
quarks in the nucleon can be decomposed into a flavor-symmetric isoscalar part (originating from gluon splitting and possibly meson cloud) and a non-perturbative isovector meson-cloud contribution. The latter can be considered to be the source of the $\bar{u} - \bar{d}$ asymmetry. In addition to the pion cloud of the nucleon we include also the isobar with its pion cloud since it was shown to give substantial contributions \[12\]. In this way one can constrain the $\pi NN$ and $\pi N\Delta$ form factors using the empirical antiquark $u$-$d$ flavor asymmetry. The physical nucleon state is expressed approximately as

$$|N\rangle = \sqrt{Z}|N\rangle_{\text{bare}} + \alpha|N\pi\rangle + \beta|\Delta\pi\rangle.$$  \hspace{1cm} (1)

Neglecting off-shell effects the light-cone momentum distribution of a quark with flavor $f$ in a proton can be written as ($B = N, \Delta$):

$$q_f(x) = Z q_{f,\text{bare}}(x) + \sum_{B,f} c_i \left[ \int_x^1 \frac{dy}{y} f_{B,f,\text{bare}}(y) q_{f,B,\text{bare}}(x/y) + \int_x^1 \frac{dy}{y} f_{NN}(y) q_{\pi,\text{bare}}(x/y) \right],$$  \hspace{1cm} (2)

where $c_i$ ($i$ labels the charge states) are the appropriate isospin Clebsch-Gordan coefficients, $q_{f,B,\text{bare}}(x)$ is the parton distribution in the bare $B_i$ baryon and $q_{\pi}(x)$ is the pion parton distribution function.

Attributing the asymmetry in the $\bar{u} - \bar{d}$ antiquark distributions to the nucleon meson cloud we are concerned with the pion light-cone distribution in the nucleon which gets contributions from final states with either nucleon or isobar:

$$f_{\pi/N}(y) = f_{\pi/N}(y) + f_{\pi/\Delta}(y).$$  \hspace{1cm} (3)

The nucleon term was calculated by Sullivan \[13\]:

$$f_{\pi NN}(y) = \frac{g_{\pi NN}}{16\pi^2} \int_{M^2 y^2(1-y)}^\infty \frac{dt}{(t + m_\pi^2)^2} F_{\pi NN}(t)^2,$$  \hspace{1cm} (4)

with $y = (k_0 + k_3)/M$ being the pion-light cone momentum fraction, with $M$ the physical mass of the nucleon (as a convenient scale), $F_{\pi NN}(t)$ is the off-shell form-factor of the $\pi NN$ vertex, while $g_{\pi NN}$ is the $\pi^0 N N$ coupling. The free-pion propagator, $D_0^\pi$, appears in the above expression in the form $(t + m_\pi^2)^{-1}$, where $t \equiv -q^2$ where $q$ denotes the pion four-momentum. The isobar contribution also plays an important role \[12\] despite the kinematical suppression coming from the isobar-nucleon mass difference. In Ref. \[12\] it was calculated using the free isobar propagator, i.e. neglecting its width. A complete relativistically covariant treatment of the isobar in vacuum and nuclear medium was introduced in Ref. \[14\] and we use that formalism to take into account the vacuum width consistent with the measured pion-nucleon scattering phase shift in the spin-3/2 isospin-3/2 channel. The full Lorentz structure of the vacuum propagator of the Rarita-Schwinger field can be expressed in terms of 10 Lorentz scalar functions \[14,15\] which contains both spin-3/2 and spin-1/2 sectors \[16\]. However, using the convenient basis from Ref. \[14\] it turns out that a single term, namely the (on mass shell) positive energy spin-3/2 contribution gives the dominant contribution and all others (some terms in the propagator are identically zero) are completely negligible. In the notation of Ref. \[14\] this is the coefficient of the projector sum $Q_{\pi NN}^{\mu\nu} \equiv Q_{\pi NN}^{\mu\nu} + P_{\pi NN}^{\mu\nu}$ which we denote by $G_{\pi NN}(p)$. The pion light-cone distribution originating from the $\Delta\pi$ state then can be expressed as:

$$f_{\pi/\Delta}(y) = \frac{y M g_{\Delta NN}^2}{6\pi^2} \int_{-\infty}^{-0} df \int_0^\infty dp_\perp \int_0^\infty p_\perp dp_\perp \cdot F_{\pi NN}(t)^2 F_{\pi \Delta}(p_\perp)^2 \frac{(M + p_\perp)}{(t + m_\Delta^2)^2}.$$  \hspace{1cm} (5)

where $p$ and $p'$ are the four-momenta of the nucleon and isobar, $t \equiv -(p - p')^2$ and $F_{\pi NN}$ are the form factors of the $\pi NN$ vertex. The form factor

$$F_{\pi NN}(p) = \exp \left[ -\frac{p^2}{2} - \frac{(M + m_\pi)^2}{\Lambda^2}\right],$$  \hspace{1cm} (6)

with $\Lambda = 0.97$ GeV (and $g_{\pi NN} = 20.2$ GeV$^{-1}$) was used in Ref. \[15\] and shown to give a good fit to the relevant pion-nucleon phase shift. For the $\pi NN$ and $\pi N\Delta$ off-shell form factors which take into account the off-shell pion we take a dipole form:

$$F_{\pi NN}(t) = \left( \frac{\Lambda_{\pi NN}^2 - m_\pi^2}{\Lambda_{\pi NN}^2 + t}\right)^2,$$  \hspace{1cm} (7)

with $X$ standing for $N$ or $\Delta$.

In order to calculate the $d - \bar{u}$ distribution for the free proton we assume that the pion sea is isospin symmetric leaving only the contribution of valence distributions. The final state with nucleon contributes through the presence of $\tau^+ \tau^- \pi^0$ with distribution $2 f_{\pi NN}$, while the isobar final state can have also a $\tau^+ \tau^- \pi^0$ with isospin weight $1/3$ or a $\tau^+ \tau^- \pi^0$ with minus sign (because of the pion valence $\bar{u}$ distribution) with respect to $f_{\pi/\Delta}$ giving in total:

$$(\bar{d} - \bar{u})_{\rho}(x) = \int_x^1 \frac{dy}{y} \left( 2 f_{\pi NN}(y) - \frac{2}{3} f_{\pi/\Delta}(y) \right) q_{\pi}(x/y),$$  \hspace{1cm} (8)

with $q_\pi(x)$ denoting the valence parton distribution of charged pion. In Fig. \[1\] we show the pion distributions with nucleon final state (solid line) and isobar final state (dashed line). For the form-factor cutoff we used the following values: $\Lambda_{\pi NN}^\tau = 0.95$ GeV and $\Lambda_{\pi NN}^\tau = 0.75$ GeV. The bare-nucleon probability then takes the value $Z = 0.69$ which suggests that higher-order terms with more than one pion do not contribute significantly. The calculated value for the $d - \bar{u}$ asymmetry in the free proton is shown in Fig. \[2\] by solid line and compared to the result using
the $\bar{u}$ and $\bar{d}$ fits CT10 \cite{CT10} (dot line). Also shown are separately the contributions from the nucleon final state

$$R_{A/d} = \frac{\sum_f c_f^2 \left\{ q_f^p(x_1) [q_f^p(y/A)(x_2) + q_f^n(y/A)(x_2)] + q_f^n(x_1) \left[ q_f^p(y/A)(x_2) + q_f^n(y/A)(x_2) \right] \right\}}{\sum_f c_f^2 \left\{ q_f^p(x_1) [q_f^p(y/A)(x_2) + q_f^n(y/A)(x_2)] + q_f^n(x_1) \left[ q_f^p(y/A)(x_2) + q_f^n(y/A)(x_2) \right] \right\}}$$

(dash line) and isobar final state (dash-dot line). These results are quite similar to the $\bar{d} - \bar{u}$ obtained in Ref. \cite{18}, although there the infinite-momentum-frame formalism was used with suitably adjusted values of the $\pi N$ and $\pi \Delta$ form factors.

III. NUCLEAR EFFECTS

First calculations of the nuclear Drell-Yan process \cite{19} suggested an enhancement coming from the medium modification of the pion cloud. However, the experimental data \cite{8} did not show that enhancement within a 10% uncertainty. Later on other groups reported more detailed calculations of the Drell-Yan ratio with a large variation in results as shown in Refs. \cite{22}. Here we consider the ratio of the cross sections of proton-nucleus and proton-deuteron scattering,

$$R_{A/d} = \frac{2 d\sigma^{pA}/dx_1 dx_2}{A d\sigma^{pd}/dx_1 dx_2}$$

where $A$ denotes the nucleus and its nucleon number. We specialize for the case of isoscalar targets for which the cross-section ratio becomes:

$$2 q_f^p(x/y) + \int_x^y \frac{dy}{y} f^{N/A}(y) \left[ q_f^p(x/y) + 2 q_f^p(x/y) \right]$$

where isospin-symmetric nuclear medium was assumed. Adding the difference of the left-hand side and right-hand side of (11) to the right-hand side of (14) and repeating the same procedure for the neutron we obtain

$$2 q_f^n(x/y) + \int_x^y \frac{dy}{y} f^{N/A}(y) \left[ q_f^n(x/y) + 2 q_f^n(x/y) \right]$$

which isospin-symmetric nuclear medium was assumed. The (anti-)quark distribution in the medium can be modified in two ways, (i) through Fermi motion and binding of the nucleon, and (ii) modification of the nucleon’s pion cloud. To establish the connection to the (anti-)quark distribution of the free nucleon we use Eq. (2) which for the free proton gives (with isobar terms not written out for brevity):

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

with

$$Z = 1 - \int_0^1 dy f^{N/A}(y) - 3 \int_0^1 dy f^{nN/N}(y),$$

where the last equality expresses flavor-charge conservation. Similarly, the quark distribution for the nuclear proton can be written:

$$q_f^p(x) = Z_A q_{f,bare}^p(x) + \frac{1}{3} \int_x^y \frac{dy}{y} f^{N/A}(y) \left[ q_{f,bare}^p(x/y) + 2 q_{f,bare}^p(x/y) \right]$$
The presence of the proton–antiproton asymmetry in free proton (solid line) compared to the difference parameter values: $\Lambda_{\pi N} = 0.95$ GeV and $\Lambda_{\pi N} = 0.75$ GeV.

Eq. (11) and its analog for the neutron. A much simpler, though approximate procedure is just to subtract the meson-cloud contribution form the antiquark distribution of the physical nucleon. Indeed, using the fact that antiquark distributions at small $x$ behave as $1/x$ one can confirm that the first and second term on the right-hand side of (11) combine to give the bare distribution if one adds to (11) its neutron analog. Using the same argument about the small $x$ behavior of antiquark distributions one can establish an approximate cancellation of the second and third as well as the fourth and fifth terms on the right-hand side of Eq. (11) and of the corresponding terms involving the neutron. This simplification was used in our previous work [8], but in the present calculation we want to take into account these contributions with the bare antiquark distributions determined by the above mentioned subtraction of the pion contribution from the physical distribution. In Fig. 3 we show the proton-neutron average of the sum of the second, third, fourth and fifth terms in Eq. (11) devided by the antiquark distribution of the free “isoscalar” nucleon. Solid line is for the pion parameter set (1) and dash line for parameter set (2).

We observe that this contribution is indeed quite small as expected from the form of the nucleon antiquark distribution and pion (as well as related nucleon) light-cone-momentum distributions. The parameter set (1) is: $M_\pi = 0.89$ GeV, $\Sigma_N^e = 0$, $\Sigma_A^e = -0.1$ GeV, $\Sigma_N^{12} = 0$, $g_{11} = 0.9$, $g_{12} = 0.3$, $g_{22} = 0.3$, while the set (2) is given by: $M_\pi = 0.89$ GeV, $\Sigma_N^e = 0$, $\Sigma_A^e = -0.05$ GeV, $\Sigma_N^{12} = 0$, $g_{11} = 1.0$, $g_{12} = 0.4$, $g_{22} = 0.3$; where $M_\pi$ is the mean-field shifted nucleon mass, $\Sigma_N^e$ the energy shift of the nucleon, $\Sigma_A^e$ and $\Sigma_A^{12}$ the delta’s mean-field shifts and $g_{ij}$ the Migdal four-fermion interaction parameters.

Introducing a shorter notation for the bare nucleon contribution to the in-medium antiquark distribution:

$$q^p_{f,\text{bare}}(x) = q^p_f(x) + q^p_{f,\text{bare}}(x) \int_0^1 f^{N/N}(y)dy - \int_x^1 \frac{dy}{y} f^{N/N}(y)q^p_{f,\text{bare}}(x/y)$$

$$- q^p_{f,\text{bare}}(x) \int_0^A f^{N/A}(y)dy + \int_x^A \frac{dy}{y} f^{N/A}(y)q^p_{f,\text{bare}}(x/y) + (p \rightarrow n),$$

FIG. 2. (Color online) The pion-cloud result for the $\bar{d} - \bar{u}$ asymmetry in free proton (solid line) compared to the difference of the proton $\bar{d}$ and $\bar{u}$ distributions (dot line) from the fit CT10 [17] and data points by Fermilab E866/NuSea Collaboration [20]. The isobar contribution (dash-dot line) is negative and much smaller than the nucleon term (dash line). Used parameter values: $\Lambda_{\pi N} = 0.95$ GeV and $\Lambda_{\pi N} = 0.75$ GeV.

FIG. 3. (Color online) Sum of the second, third, fourth and fifth terms in Eq. (11) devided by the antiquark distribution of the free “isoscalar” nucleon. Solid line is for the pion parameter set (1) and dash line for parameter set (2).
we can finally write the sum of in-medium proton and neutron antiquark distribution as

\[ q_f^{p/A}(x) + q_f^{n/A}(x) = \int_x^A \frac{dy}{y} f_{N}^{N}(y) q_f^{p+n}(x/y) + 2 \int_x^A \frac{dy}{y} \left[ f_{p}^{p/A}(y) - f_{p}^{N/A}(y) \right] \left[ q_{\pi}^f(x/y) + q_{\pi}^{f+}(x/y) + q_{\pi}^{f-}(x/y) \right]. \]  

(16)

The convolution with \( f_{N}^{N}(z) \) takes into account Fermi motion and binding effects on the in-medium nucleons. For the function \( f_{N}^{N}(z) \) we take the result of Birse [21]:

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

(17)

where \( \epsilon \equiv p_F/M, \eta \) is a parameter with value slightly below one which takes into account the nuclear binding and \( \Theta(x) \) is the unit step function. In Fig. 4 we show the ratio of the \( F_2(x) \) structure functions for the isospin symmetric nuclear matter and the deuteron. We assume negligible medium effects in the deuteron and for the nuclear medium we use the convolution model with light-cone distribution \( (17) \) and parameters \( p_F = 250 \text{ MeV} \) and \( \eta = 0.97 \). In this way one can reproduce the negative slope in the classical EMC effect for \( 0.1 < x < 0.5 \) as shown for example in Ref. 22. The experimental enhancement observed around \( x = 0.1 \) can be attributed to the pion enhancement as shown in the figure for the two parameter sets (1) and (2) used also for the plots in Fig. 3 and given above. The pion enhancement term was calculated by the convolution of the in-medium pion light-cone distribution enhancement relative to the free nucleon and the pion \( F_2 \) distribution. Note that expected shadowing effects would lead to decrease of the nuclear cross section for \( x \leq 0.05 \).

FIG. 4. (Color online) The ratio of \( F_2(x) \) per nucleon for isospin symmetric nuclear medium with parameters: Fermi momentum \( p_F = 250 \text{ MeV}, \eta = 0.97 \), and deuteron. Pion contribution with parameter sets (1) and (2) is included in the results shown by solid and dash-dot lines.

FIG. 5. Types of diagrams contributing to the in-medium pion distribution with outgoing nucleon. The same types appear also with outgoing delta baryon. The dash line denotes the dressed pion propagator, while the solid and the double line correspond to nucleon and delta.

IV. ANTIQUARKS IN BOUND NUCLEONS: PION CONTRIBUTION

We now turn to consideration of the pion contribution to antiquark distributions in nucleons bound in large nuclei which we model by an infinite system with appropriate average nuclear density.

For corresponding pion properties in the nuclear matter we use the results of a fully covariant self-consistent model developed in Ref. 23. Compared to the case of the free nucleon the nuclear environment changes the pion propagator appearing in the Sullivan formula (11) and renormalizes the \( \pi N N \) as well as the \( \pi N \Delta \) vertices through nucleon-nucleon correlations modeled by the Migdal four-fermion interactions [24]. Nucleon properties are also affected and we take into account the binding effects through mean-field mass and energy shifts consistent with approach in Ref. 22.

The inclusion of the dressed pion propagator is straightforward but the dressing of the \( \pi NN \) and the \( \pi N \Delta \) vertices requires summation of nucleon-hole and delta-hole bubbles. The types of relevant diagrams are shown in Fig. 5. For resummation of these diagrams we use the relativistically covariant formalism introduced in Ref. 25 and applied for pion and delta self energy calculation in Ref. 20. In the present case it concerns a different type of contribution which for the nucleon in
where \( q = p - p' \) is the pion 4-momentum, \( u(p) \) is the nucleon in-medium spinor (with mean-field shifts of mass and energy) and \( \Pi_{\mu
u}(q) \) is the resummed contribution of nucleon-hole and delta-hole loops. Using the decomposition of nucleon-hole and delta-hole loops\(^{23}\):

\[
\Pi^{(NN)}_{\mu\nu}(q) = \sum_{i,j=1}^{2} \Pi^{(nn)}_{ij}(q) L^{(ij)}_{\mu\nu}(q) + \Pi^{(Nh)}_{\mu\nu}(q) T_{\mu\nu}(q),
\]

\[
\Pi^{(Nh)}_{\mu\nu}(q) = \sum_{i,j=1}^{2} \Pi^{(Nh)}_{ij}(q) L^{(ij)}_{\mu\nu}(q) + \Pi^{(Nh)}_{T\mu\nu}(q) T_{\mu\nu}(q),
\]

where \( L^{(ij)}_{\mu\nu}(q) \) and \( T_{\mu\nu}(q) \) have projector properties, the nonzero contribution involving one particle-hole (nucleon-hole or delta-hole) loop comes from

\[
\Pi^{(1)}_{\mu\nu} = g'_{11} \left[ \Pi^{(Nh)}_{11}(q) L^{(11)}_{\mu\nu}(q) + \Pi^{(Nh)}_{21}(q) L^{(21)}_{\mu\nu}(q) \right] + g'_{12} \left[ \Pi^{(Nh)}_{12}(q) L^{(12)}_{\mu\nu}(q) + \Pi^{(Nh)}_{22}(q) L^{(22)}_{\mu\nu}(q) \right],
\]

where \( g'_{11}, g'_{12}, g'_{22} \) are the usual Migdal parameters\(^{23}\).

In order to perform the summation involving arbitrary number of nucleon-hole or delta-hole loops it is convenient to introduce the following matrices\(^{25}\):

\[
g^{(L)} = \begin{pmatrix}
-\frac{g'_{11}}{1} & \frac{g'_{12}}{0} & 0 & 0 \\
\frac{g'_{12}}{0} & \frac{g'_{11}}{0} & 0 & 0 \\
0 & \frac{g'_{22}}{0} & \frac{g'_{22}}{1} & 0 \\
0 & \frac{g'_{22}}{0} & 0 & \frac{g'_{22}}{1}
\end{pmatrix},
\]

\[
\Pi^{(L)} = \begin{pmatrix}
\Pi^{(Nh)}_{11} & 0 & 0 & 0 \\
0 & \Pi^{(Nh)}_{12} & \Pi^{(Nh)}_{21} & \Pi^{(Nh)}_{22} \\
0 & \Pi^{(Nh)}_{21} & 0 & 0 \\
0 & \Pi^{(Nh)}_{22} & 0 & 0
\end{pmatrix}.
\]

The lowest order contribution\(^{26}\) can then be written as:

\[
\Pi^{(1)}_{\mu\nu} = \left[ (\Pi^{(L)} g^{(L)})_{11} + (\Pi^{(L)} g^{(L)})_{31} \right] L^{(11)}_{\mu\nu}(q) + \left[ (\Pi^{(L)} g^{(L)})_{21} + (\Pi^{(L)} g^{(L)})_{41} \right] L^{(21)}_{\mu\nu}(q).
\]

Higher order terms are accounted for by taking appropriate matrix elements of products of \((\Pi^{(L)} g^{(L)})\) matrices and the summation of terms with arbitrary number of loops is simply achieved by replacing \((\Pi^{(L)} g^{(L)})\) in\(^{27}\) by \((\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)} )^{-1} \) leading to:

\[
\Pi_{\mu\nu} = \left[ (\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)} )^{-1})_{11} + (\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)} )^{-1})_{31} \right] L^{(11)}_{\mu\nu}(q) + \left[ (\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)} )^{-1})_{21} + (\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)} )^{-1})_{41} \right] L^{(21)}_{\mu\nu}(q).
\]

Adding \( g_{\mu\nu} \) to \( \Pi_{\mu\nu} \) in\(^{28}\) we obtain the full contribution of the diagram with nucleon final state. Squaring its absolute value and performing summation over the spin projections of the nucleon in the final state and averaging for the nucleon in the initial state we obtain:

\[
\frac{1}{2} \sum_{s,s'} |K_N|^2 = A_{qq} \left( 2(p - \Sigma^u_N u) \cdot q (p' - \Sigma^u_N u) \cdot q - q^2 [M^2 + (p' - \Sigma^u_N u) \cdot (p - \Sigma^u_N u)] \right) + 2A_{qu} \cdot \left( (p - \Sigma^u_N u) \cdot u (p' - \Sigma^u_N u) \cdot q + (p' - \Sigma^u_N u) \cdot u (p - \Sigma^u_N u) \cdot q \right. \\
- q \cdot u [M^2 + (p' - \Sigma^u_N u) \cdot (p - \Sigma^u_N u)] + A_{uu} \cdot \left. (2(p - \Sigma^u_N u) \cdot u (p' - \Sigma^u_N u) \cdot u - M^2 - (p' - \Sigma^u_N u) \cdot (p - \Sigma^u_N u), \right)
\]

where \( u \) is the 4-velocity of the medium (implicitly present also in Eqs.\(^{18}-20\)), \( M_s = M_N + \Sigma^s_N, \Sigma^s_N \) and \( \Sigma^u_N \) are the nucleon mean-field mass and energy shifts. The factors \( A_{qq}, A_{qu}, A_{uu} \) are given by:

\[
A_{qq} = 2 \left[ 1 + A + q \cdot u B / \sqrt{q^2 - (q \cdot u)^2} \right]^2,
\]
\[
A_{qu} = -2 \text{Re} \left[ q^2 B / \sqrt{q^2 - (q \cdot u)^2} (1 + A - q \cdot u B / \sqrt{q^2 - (q \cdot u)^2}) \right],
\]
\[
A_{uu} = 2 q^2 B / \sqrt{q^2 - (q \cdot u)^2},
\]

(26)
where the bar denotes complex conjugation,

\[ A = [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{11} + [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{31}, \]

\[ B = [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{21} + [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{41}, \]

and we used that:

\[ f^{(1)}_{\mu
u}(q) q^\nu = q_\mu, \]

\[ f^{(21)}_{\mu
u}(q) q^\nu = \frac{q \cdot u}{\sqrt{q^2 - (q \cdot u)^2}} q_\mu - \frac{q^2}{\sqrt{q^2 - (q \cdot u)^2}} u \mu \]

To compute the pion light-cone distribution per nucleon in the medium we integrate over incoming nucleons in the Fermi sea and outgoing ones above the Fermi sea, restricting the pion light-cone momentum fraction to the specified value by inserting a delta function and finally divide by the nucleon density. The final expression obtained is:

\[
f_N(y) = 3 M y \left( \frac{f_N}{m_\pi} \right)^2 \frac{1}{32 \pi^3 p_F^4} \int_{-p_F}^{p_F} dp_3 \int_0^{(p_F - p_3)^2} p_\perp dp_\perp \int_0^\infty p_\perp' dp_\perp' \int_0^{2\pi} d\vartheta \frac{1}{2b} \sum_{s,s'} |K_N|^2 |F_{\pi NN}^{(p)}(-q^2)|^2, \]

where

\[ D_\pi(q) \]

is the in-medium dressed pion propagator, \( b \equiv M y - p_3 - \sqrt{M_\pi^2 + p_3^2 + p_\perp^2} \), \( p_\perp^\text{min} = \sqrt{2b \sqrt{M_\pi^2 + p_3^2} - M_\pi^2 - b^2} \), \( \vartheta \) is the angle between \( \vec{p}_\perp \) and \( p_\perp' \), and the \( \pi N N \) form factor \( F_{\pi NN}^{(\pi)}(-q^2) \) was also included.

The contribution coming from the delta baryon in the final state is made more involved by the complicated structure of the in-medium delta propagator \([14]\). However, considerable simplification can be achieved by including only the two dominant contributions in the convenient relativistically covariant decomposition since the imaginary part of the other components is typically two orders of magnitude smaller at nuclear saturation density \([20]\). The dominant contributions come from the \( Q_{11}^{(\pi)} \) and \( P_{55}^{(\pi)} \) terms which were degenerate in the free delta case, but are different in the medium \([14]\). Summation of particle-hole loops dressing the \( \pi N \Delta \) vertex is analogous to the \( \pi NN \) case, with only difference being in the relevant matrix elements of \( \Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1} \), replacing the expression \([24]\) with:

\[
\Pi_{\mu\nu}^\Delta = \left( [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{13} + [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{33} \right) L^{(1)}_{\mu
u}(q) \\
+ \left( [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{23} + [\Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1}]_{43} \right) L^{(21)}_{\mu
u}(q). \]

The expression analogous to \([25]\) in this case takes the form:

\[
\frac{1}{2} \sum_{s,s'} |K_\Delta|^2 = \frac{1}{2} \text{Tr} \left[ (\mathbf{p} - \Sigma_N^{\mu} \sigma^\mu + M_\pi) \text{Im} G_{\mu\nu}^{\mu\nu}(p') \right] (g_{\mu\alpha} g_{\nu\beta} + \Pi_{\mu\alpha} \Pi_{\nu\beta}) q^\alpha q^\beta, 
\]

where \( \text{Im} G_{\mu\nu}(p') \) denotes the imaginary part of the in-medium \( \Delta \) propagator for which we take the dominant contribution given in the basis used in Ref. \([14]\) by just two terms:

\[
G_{\mu\nu}(p') = Q_{11}^{(\pi)}(p') G_{11}^{(Q)}(p') + P_{55}^{(\pi)}(p') G_{55}^{(P)}(p'). 
\]

In this way the expression \([31]\) takes the form:

\[
\frac{1}{2} \sum_{s,s'} |K_\Delta|^2 = \left[ A_{qq}^{(\Delta)} c_{qq}^{(Q)} + 2 A_{qq}^{(\Delta)} c_{q\alpha}^{(Q)} + A_{qq}^{(\Delta)} c_{q\alpha}^{(Q)} \right] \text{Im} G_{11}^{(Q)}(p') \\
+ \left[ A_{qq}^{(\Delta)} c_{qq}^{(P)} + 2 A_{qq}^{(\Delta)} c_{q\alpha}^{(P)} + A_{qq}^{(\Delta)} c_{q\alpha}^{(P)} \right] \text{Im} G_{55}^{(P)}(p'), 
\]

where the expressions for \( A_{qq}^{(\Delta)}, A_{q\alpha}^{(\Delta)}, A_{\alpha q}^{(\Delta)}, c_{qq}^{(Q)}, c_{q\alpha}^{(Q)}, c_{\alpha q}^{(Q)}, c_{qq}^{(P)}, c_{q\alpha}^{(P)}, c_{\alpha q}^{(P)} \) are given in the Appendix. The pion light-cone momentum distribution stemming from the process with the nucleon emitting a pion and a delta baryon is analogous to expression \([29]\) and reads:

\[
f_\Delta(y) = 3 M y \left( \frac{f_\Delta}{m_\pi} \right)^2 \frac{1}{32 \pi^4 p_F^4} \int_{-p_F}^{p_F} dp_3 \int_0^{(p_F - p_3)^2} p_\perp dp_\perp \int_0^\infty dp_\perp' \int_0^{2\pi} d\vartheta \frac{1}{2} \sum_{s,s'} |K_\Delta|^2 \left| F_{\pi NN}^{\Delta}(p') F_{\pi NN}^{(\pi)}(-q^2) D_\pi(q) \right|^2. 
\]
We checked by explicit numerical calculation that both Eqs. (29) and (34) have the correct low-density limit, i.e. reproduce the free nucleon and delta results.

V. NUMERICAL RESULTS AND DISCUSSION

For the computation of in-medium pion and isobar properties we rely on the recently developed relativistically covariant self-consistent model presented in Ref. [23] and used for the nuclear photoabsorption calculation in the isobar region in Ref. [27]. For the medium computation we use the same $\pi NN$ and $\pi N\Delta$ form factors as in the vacuum one and include them in the model of Ref. [23]. The values of the Migdal $g'$ parameters which model the short-range nucleon and isobar correlations were taken in the range preferred by the results of Ref. [27]; in this work a good description of the nuclear photo-absorption cross section in the isobar region was obtained. Binding effects for the nucleon are taken into account by the effective (mean-field) mass $M_\star$ and the energy shift $\Sigma_\star$. A consequence of the use of the mean-field approximation is a reduction of the in-medium pion distribution coming from the nucleon final state. Namely the dominant contribution to it is the term proportional to $A_{qq}$ in Eq. (25) with:

$$
(2(p - \Sigma_\star u) \cdot (p' - \Sigma_\star u) - q^2 \left[ M_\star^2 + (p' - \Sigma_\star u) \cdot (p - \Sigma_\star u) \right]) = -2M_\star^2q^2,
$$

(35)

which is the same expression as for the free nucleon, except that $M_\star$ appears instead of $M$.

Since $M_\star/M < 1$ a further suppression in addition to that from the Pauli blocking is obtained, depending on the actual value of $M_\star/M$. The latter is difficult to constrain since observables generally are only sensitive to the combination $M_\star + \Sigma_\star$. Since our aim is to make comparison with experiments on finite nuclei (rather than nuclear matter) with an average density smaller than the saturation density we assume small values for the energy shift in the range zero to $\Sigma_\star = 0.04\text{GeV}$, corresponding to effective mass values in the range of $0.85\text{ GeV}$ and $0.89\text{ GeV}$. These values are close to the ones used in more elaborate treatments of nuclear matter [28, 29] where values of $0.8 - 0.85\text{ GeV}$ at saturation density give good agreement with observables. The mean-field shifts of the isobar mass and energy are chosen in such a way that they reproduce the isobar-nucleon mass difference used in Ref. [27]. This means $\Sigma_\Delta = -0.05\text{GeV}$ and $-0.1\text{ GeV}$ and zero for the energy shift.

In Figs. 6 and 7 we show the pion distributions $f^{\pi N/A}(y)$ and $f^{\pi \Delta/A}(y)$ for in-medium nucleons for different parameter sets. For the nucleonic distribution one observes a reduction coming partly from the Pauli blocking of the nucleons in the medium and partly from the $M_\star$ effect (which leads to a suppression roughly by the factor $(M_\star/M)^2$). The pion broadening in the nuclear medium only partly compensates these effects and a net reduction is the result. This is not completely surprising since the computations of Ref. [23] do not lead to appreciable softening of the pion spectrum in the medium which would result in enhanced pion distribution. In this respect the pion dressing of Ref. [23] is not significantly different from an older calculation [30] which used a nonrelativistic treatment of the isobar and a softer pion-nucleon-delta form factor. On the other hand a significant enhancement is observed for the contribution originating from the transition $N \to \pi \Delta$ which is not Pauli suppressed. These results emphasize the importance of careful treatment of the in-medium isobar self energy and propagator, which is made possible by the convenient complete basis introduced in this context in Ref. [14]. As a consequence in the nuclear medium the combined effects from the pion and from the isobar can produce sizeable increase in the pion light-cone distribution. The latter is constrained to smaller light-cone-momentum ratio $y$ values because of kinematical effect of the isobar-nucleon mass difference but can still have significant effects on the DY cross-section ratio.

Since we are considering isospin symmetric nuclear medium and make a comparison with the deuteron it is advantageous to consider the pion distribution in an “isoscalar” nucleon, i.e. to consider a proton-neutron average. Taking into account pions of all charges gives the complete pion distribution of an “isoscalar” nucleon:

$$
f^{\pi/A}(y) = 3f^{\pi N/A}(y) + 2f^{\pi \Delta/A}(y).
$$

(36)

In Fig. 8 we show the function $f^{\pi/A}(y)$ for different input parameter values compared to the pion distribution of the free “isoscalar” nucleon. The probability $Z_A$ of the bare nucleon in the medium takes the values from 0.6 to 0.65, i.e. just slightly smaller than in the free nucleon case.
Before examining the DY cross-section we show the ratio of the antiquark distribution in the in-medium proton and the same distribution in the free proton. The up and down antiquark distributions experience different in-medium modification due to different weights of nucleon and delta contributions even in isospin-symmetric nuclear medium. In Fig. [9] we show the ratios of antiquark distributions for an in-medium proton relative to the free proton for parameter sets (1) and (2).

due to the delta-baryon final state as compared to quite modest enhancement and even suppression for the down antiquark as a consequence of larger weight of nucleon final state and smaller weight of delta-baryon final state as compared to the up antiquark. This difference points to the possibility of distinguishing between effects coming from the medium modification of the nucleon and delta baryon by examining observables to which up and down antiquarks contribute with different weights.

We now turn to the DY cross-section ratio (10). In Figs. [10] and [11] we show the cross-section ratio (10) as a function of $x_2$ for fixed values of $x_1$. The input parameter values are given in the figure caption. We observe an enhancement only for small values of $x_2$, typically less than 0.2, and for $x_2 > 0.1$ a decreasing trend as a result of the convolution with nucleon distribution (17).

For comparison with the measurements of Ref. [8], we computed the ratio of the nuclear and deuteron cross sections for given $x_2$ and integrating over $x_1$ satisfying the condition $x_1 > x_2 + 0.2$ corresponding to the experimental cut-off. Fig. [12] shows the measured values with error bars and the calculated curves for different input parameters. We consider the lowest curve in Fig. [12] with $M_π = 0.8$ GeV and corresponding rather pronounced suppression of the order $(M_π/M)^2$ probably exaggerating the effect of the nucleon mean-field approximation and regard the other two curves as representing better our results based on the preferred parameter sets.

VI. SUMMARY

In this work we presented an analysis of nuclear effects on the Drell-Yan process. The approach is based on the
We took into account the change of the pion cloud originating from both the pion-nucleon and pion-delta states and a small correction (neglected in previous work) attributed to the binding effect of bare nucleon.

Fermi motion and binding of nuclear nucleons were accounted for by the two-parameter light-cone-momentum distribution \[ F_{\pi NN}(x) \] which reproduces the negative slope of the classical EMC effect in the region \( 0.1 < x < 0.5 \) as shown in Fig. 10. Taking into account the pion enhancement which comes from the pion-delta state of the nucleon (which is significant only for small light-cone-momentum \( y \approx 0.2 \) values) leads to some enhancement of the \( F_2(x) \) ratio for \( x \leq 0.2 \).

Pion and delta properties in the nuclear medium are calculated in a recently developed fully covariant self-consistent model \[ \Sigma \] which consistently takes into account the \( \pi NN \) and \( \pi N\Delta \) vertex corrections due to Migdal short-range correlations. Pronounced softening of the in-medium pion spectrum present in simpler models does not appear in this approach and consequently Pauli blocking causes some suppression of the pion distribution coming from the pion-nucleon state for an in-medium nucleon. However, enhancement results from a careful treatment of the pion-delta state as a consequence of pion broadening and delta shift and broadening.

The net effect for preferred parameter values is a modest enhancement of pion light-cone-momentum distribution mostly concentrated around \( y \approx 0.2 \) value. As a consequence the DY cross-section ratio exceeds one for small values, typically less than 0.2, of the \( x_2 \) variable for fixed \( x_1 \) values or integration over it corresponding to some experimental cuts. The convolution with distribu-

\[ x_1 = 0.25 \]
\[ x_1 = 0.35 \]
\[ x_1 = 0.45 \]
tion (17) acts qualitatively on the antiquarks in the same way as on the quarks producing a negative slope which is less pronounced for the integrated properties of the pion and delta baryon. It would be very desirable to achieve measurements with considerably smaller uncertainties which could contribute to the resolution of some decades old issues of nuclear physics. We remark that another interesting possibility for studying sea quark distributions in nuclei would be the use of an Electron-Ion Collider as detailed in the joint report of the Brookhaven National Laboratory, the Institute for Nuclear Theory (Seattle, WA) and the Thomas Jefferson National Accelerator Facility 31. The proposed semi-inclusive deep-inelastic electron-nucleus scattering would provide new information about the structure of nuclei and quantum chromodynamics of nuclear matter and extend possibilities for studying effects of the transverse momentum distribution of partons 32.

Appendix

The terms of expression (33) necessary to calculate the contribution of the in-medium delta are given as follows:

\[
A_{qq}^{(\Delta)} = \left| A_\Delta + q \cdot u B_{\Delta} / \sqrt{q^2 - (q \cdot u)^2} \right|^2,
\]

\[A_{qu}^{(\Delta)} = -2 \text{Re} \left[ q^2 B_{\Delta} / \sqrt{q^2 - (q \cdot u)^2} (1 + A_\Delta - q \cdot u B_{\Delta} / \sqrt{q^2 - (q \cdot u)^2}) \right],\]

\[A_{uv}^{(\Delta)} = 2 \left| q^2 B_{\Delta} / \sqrt{q^2 - (q \cdot u)^2} \right|^2,\]

\[c_{qq}^{(Q)} = (M_\Delta + (p - \Sigma_N^u \cdot \hat{p}')) \left[ t + (q \cdot \hat{p}')^2 - (q \cdot X(p'))^2 \right],\]

\[c_{qu}^{(P)} = \frac{1}{3} (M_\Delta + (p - \Sigma_N^u \cdot \hat{p}')) \left[ t + (q \cdot \hat{p}')^2 + 3 (q \cdot X(p'))^2 \right],\]

\[c_{qu}^{(Q)} = -(M_\Delta + (p - \Sigma_N^u \cdot \hat{p}'))(q \cdot u - q \cdot \hat{p}' u \cdot \hat{p}' + q \cdot X(p') u \cdot X(p')) ,\]

\[c_{uu}^{(P)} = -\frac{1}{3} (M_\Delta + (p - \Sigma_N^u \cdot \hat{p}'))(q \cdot u - q \cdot \hat{p}' u \cdot \hat{p}' - 3q \cdot X(p') u \cdot X(p')) ,\]

\[c_{uu}^{(Q)} = (M_\Delta + (p - \Sigma_N^u \cdot \hat{p}')) \left[ -1 + (u \cdot \hat{p}')^2 - (u \cdot X(p'))^2 \right],\]

\[c_{uu}^{(P)} = \frac{1}{3} (M_\Delta + (p - \Sigma_N^u \cdot \hat{p}')) \left[ -1 + (u \cdot \hat{p}')^2 + 3(u \cdot X(p'))^2 \right],\]

\[A_\Delta \equiv \left[ \Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1} \right]_{13} + \left[ \Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1} \right]_{33},\]

\[B_\Delta \equiv \left[ \Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1} \right]_{23} + \left[ \Pi^{(L)} g^{(L)} (1 - \Pi^{(L)} g^{(L)})^{-1} \right]_{43},\]

\[X(p) \equiv \frac{(p \cdot u) p_\mu - p^2 u_\mu}{p^2 \sqrt{(p \cdot u)^2 / p^2 - 1}}.\]

where \( t = -q^2 \) and \( \hat{p}_\mu = p_\mu / \sqrt{p^2} \).

ACKNOWLEDGMENTS

This research was supported in part by the Hungarian Research Foundation (OTKA) grant 71989. CLK acknowledges the kind hospitality of the Kernfysisch Versneller Instituut in Groningen.

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