Proton and anti-proton production in the forward region of d+Au collisions at RHIC from the color glass condensate

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Abstract

The power-law tail of high-$p_t$ $\pi^\pm$ spectra observed in forward d+Au collisions at RHIC can be attributed to the power-law decrease of the dipole forward scattering amplitude appearing in the color glass condensate (CGC) approach. Forward particle production probes the small-$x$ gluon distribution of the target nucleus where its anomalous dimension is rather flat ($\gamma = 0.6 \sim 0.8$) for moderately high $p_t$ ($\lesssim 5$ GeV), and where the leading-twist DGLAP approximation is not valid. In the same framework, we examine $p$ and $\bar{p}$ production using baryon fragmentation functions parameterized in the Lund fragmentation scheme. This provides a good description of the forward $p$ spectrum while it underestimates the $p$ data by as much as a factor of $2 \sim 3$ at $p_t \lesssim 4$ GeV. Part of this anomalous baryon excess can be attributed to surviving constituent diquarks from the deuteron projectile. Thus, the contribution from diquark scattering may play an essential role for forward baryon formation.
# 1 Introduction

One mysterious problem in high-energy p(d)+A collisions is a baryon excess seen in the $p/\pi^+$ ratio, observed at semi-hard transverse momenta ($p_t \gtrsim 1$ GeV) over a wide range of rapidity. Measurements include the central rapidity region at Fermilab ($\sqrt{s} = 27.4, 38.8$ GeV) \[1\] and RHIC ($\sqrt{s} = 200$ GeV) \[2\] as well as forward rapidities ($y \sim 3$) at RHIC \[3\]. In a phenomenological picture, the standard string models assume a diquark-quark string structure of the proton beam and dominance of soft processes for leading baryon production as in the inside-outside cascade picture. This leads to baryon production only near beam rapidity and with small $p_t$. To explain the baryon production near midrapidity two additional nonperturbative models have been proposed: the diquark breaking \[4\] and gluon junction \[5\] mechanisms. However, these models are also restricted to small momentum transfer ($p_t \lesssim 2$ GeV).

Next-to-leading order (NLO) perturbative QCD calculations with standard fragmentation functions cannot be reconciled with the forward-rapidity data from d+Au collisions at RHIC, either. The data clearly show a large baryon excess \[6\]. Meanwhile, a pointlike diquark picture involving the diquark form-factor gives a good description of the FNAL data near midrapidity over $2 \lesssim p_t \lesssim 10$ GeV \[7\]. Thus, the study of high-$p_t$ baryon spectra in high-energy p(d)+A collisions is very challenging and provides an opportunity for revealing soft successive processes from proton breakup to their parton fragmentations.

It is also known that physics toward forward rapidity of the incident proton from midrapidity is accompanied by a change of particle production mechanism: transition from string breaking region, where the particle production is dominated by quark-antiquark pairs created from vacuum, to projectile fragmentation region, where initial charge and isospin of the projectile are most likely conserved for the particle production. A signal of such a transition appears, for instance, in antiparticle-to-particle ratios from p+p collisions at large transverse momentum, which are independent of rapidity at less than $y \sim 1.5$ but above this point decrease with rapidity, irrespective of the detected hadron species \[8\]. In the forward-rapidity region, therefore, diquarks carrying away a large fraction ($\sim 2/3$) of the incoming proton momentum will increase the importance of its role for leading baryon production in the projectile fragmentation region \[9\].

In very high-energy p+A collisions, in particular for the kinematics realized at forward rapidity, a dense target close to the black-body (or unitarity) limit should destroy completely the coherence of the projectile partons since all of them experience large transverse kicks on the order of the saturation scale (onset of gluon saturation phenomenon) and hence fragment independently, mainly into pions \[10\]. This leads to a strong suppression of the forward baryon number. Also, according to this idea the $p_t$ spectrum is expected to be rather flat up to the saturation scale. Although recent BRAHMS d+Au data at $y \sim 3$ \[3\] are almost flat over the range $1 < p_t \lesssim 3$ GeV, these data do not seem to support this idea of forward baryon number suppression at RHIC energy. Rather, the ratio $p/\pi^+ \sim 0.8$ is substantially larger than for p+p collisions \[1\] in the range $1 < p_t \lesssim 3$ GeV. Therefore, additional quantitative theoretical studies of forward baryon production in p(d)+A collisions,

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1Pythia simulations \[11\] also underpredict this ratio by a factor $2 \sim 3$ as compared to the data \[3\].
and its relation to saturation physics is of importance in detailed comparisons with the data.

Another important experimental observation is that the ratio of hadron $p_t$-distributions in $d+Au$ versus $p+p$ collisions is suppressed at large rapidity \cite{12}, while showing a (species-dependent) Cronin enhancement at midrapidity \cite{13, 14}. Within the Color Glass Condensate (CGC) formalism \cite{15}, the disappearance of the Cronin peak at moderate $p_t$ is induced by quantum evolution of the dense coherent gluons in the target \cite{16, 17, 18}, therefore signaling the emergence of saturation physics\footnote{Other scenarios also were proposed to explain the phenomenon \cite{19, 20}.}. In the CGC picture, the overall growth of the gluon density is decelerated with rapidity because the rate of gluon fusion becomes comparable with that of gluon emission. With increasing rapidity, thus the cross section of the gold target at fixed impact parameter grows less rapidly than that of a proton target \cite{21}. Along the lines of the CGC formalism, Ref. \cite{22} developed an asymmetric DGLAP$_{\text{proj}}$ \cite{23} $\otimes$ BFKL$_{\text{targ}}$ \cite{24} factorization scheme which accounts for recoil effects due to collinear gluon radiation in the projectile, because radiation becomes important with higher $p_t$ \cite{25}. Furthermore, in Ref. \cite{26} a parameterization of the anomalous dimension $\gamma$ describing the quantum $x$-evolution away from the Glauber-Mueller \cite{27} or McLerran-Venugopalan saturation models \cite{28} (which do not include evolution due to fixed $\gamma = 1$) over a wide range of rapidity was provided. Indeed, these formulations give a good description of charged hadron or neutral pion production. Then, the importance of recoil effects and the appropriate use of the $p_t$ and $y$-dependent anomalous dimension in the high $p_t$ region were discussed quantitatively \cite{22, 26}.

In this paper, we discuss $\pi^\pm$, $p$, $\bar{p}$ production with high transverse momentum ($p_t \gtrsim 1$ GeV) in the deuteron fragmentation region ($y = 3.0$) of $d+Au$ collisions at RHIC energy. At such rapidity the light-cone momentum fractions of the partons participating in the scattering becomes very asymmetric between projectile (d) and target (Au) \cite{29}; large-$x$ quarks with $x_p = O(10^{-1})$ from the deuteron collide with (many) small-$x$ gluons, $x_A = O(10^{-3} \sim 10^{-4})$ from the gold target (see figures in \cite{22, 26}). Therefore, one should treat the deuteron as a dilute and the gold nucleus as a dense object involving saturation effects. Its saturation scale is given by $Q^2_s \sim A^{1/3} e^{\lambda y}$ with $A \simeq 200$ the mass number and $y$ the rapidity of the target gluons. $\lambda$ is a constant describing the growth rate of the saturation momentum with $1/x$. Thus, the saturation effects in the target are most evident at large $y$ and the black-body limit of the target discussed in \cite{10} is most likely expected there. On the other hand, hadron production close to the deuteron beam rapidity is dominated by its valence quarks (or their diquark bound state). Energy loss of these leading quarks traversing the cold nucleus might occur beyond the leading log $Q^2$ approximation \cite{30}.

We will basically use the same expressions for the single-inclusive hadron production cross sections as discussed in Refs. \cite{22, 26} in this paper, but will reformulate them slightly to include mass corrections (massless particles have been assumed in Refs. \cite{22, 26}); cf. Sec. 2. The collinearly factorized form derived here represents our basic equation. In subsequent sections we discuss details of the three "building blocks" (i.e. parton distribution function (PDF), fragmentation function (FF) and dipole cross section), which are universal
(or process-independent) objects. As emphasised in [7], the scattered constituent diquarks represent a possible source of the baryon excess measured at semi-hard \( p_t \). This forces us to incorporate the diquark component into the first building block, i.e. into the PDF of a proton. We shall employ the existing diquark PDF from Ref. [31], which is based on the extreme (scalar) diquark picture with \( Q^2 \)-dependent form-factor of the diquark, to discuss the \( p_t \) spectra of hadron cross sections including scaling violation. Here we will introduce the breakup probability of the diquark scattered off the CGC assuming that the cross section for diquark scattering is damped by a \( Q^2 \)-dependent form-factor [7]. Regarding diquark fragmentation into hadrons, phenomenological considerations using counting rules have been applied to the data, and also more theoretical models without scaling violations have been proposed [32, 7]. We will, however, not rely on these models but rather model the diquark FFs in terms of the standard FFs with scaling violations. As our standard baryon FFs we employ parameterizations extracted from JETSET simulations [33], rescaling the \( Q^2 \)-dependence such as to obtain a good fit to the KKP [34] or AKK [35] FFs.

Although considerable model-dependences originate from mainly the scaling violation of baryon or diquark FFs, the inclusion of constituent diquarks at large \( x \) in the deuteron proves essential to explain the measured baryon excess relative to standard PDF and FF sets. For \( \pi^\pm \) and \( \bar{p} \) production, on the other hand, the diquark contribution is not so large.

This paper is organized as follows. In the next section, we formulate the single-inclusive hadron production cross sections in high-energy \( d+Au \) collisions including mass corrections. In Secs. 3 and 4 the predictions from the the model are confronted with experimental data for forward \( \pi^\pm \) and \( p/\bar{p} \) production, respectively. A detailed discussion is given in Sec. 5. Finally, we summarize in Sec. 6.

2 Forward hadron production in proton-nucleus collision within the CGC formalism

It is well known that collinear factorization theorem is generally broken by higher-twist (multiple scattering) effects in hadron-hadron collisions, although the validity of the theorem is proved up to the twist-4 level [36]. Within the CGC formalism, however, all higher-twist effects can be included into one building block, i.e. the dipole forward scattering amplitude. Hence, even if one considers higher-twist effects in \( p+A \) collisions, the predictive power of the theorem is still preserved [37]. It was verified also for the case with one-loop radiative corrections to projectile partons [22]. This expression [22] is valid only under the condition that the target is in the saturation or so-called “extended geometric scaling” regime while the projectile is dilute [38]. Therefore, this approach does hold for forward rapidity kinematics which we are interested in here, unlike the usual \( k_t \)-factorization approach [39]. In this section, we shall briefly derive collinearly factorized forms describing the single-inclusive hadron production cross section in \( p+A \) collisions.

To describe the \( p + A \rightarrow h + X \) process depicted in Fig. 1 we work with light-cone momenta in a frame where the proton has large plus component \( P_p = (\sqrt{s/2}, 0, \vec{0}_t) \) and
Figure 1: Kinematics of the $p + A \rightarrow h + X$ process. The incoming parton $i$ from the proton interacts with the target CGC field of small-$x$ gluons to all orders via a $2 \rightarrow 1$ process. The outgoing parton $f$ hadronizes into a massive hadron $h$ with momentum fraction $z_h$.

The nucleus a large minus component $P_A = (0, \sqrt{s/2}, \vec{0}_t)$. $\sqrt{s}$ denotes the center of mass energy. The incoming parton $i$ carries a fraction $x_p$ of the proton momentum $P_p^+$, interacting with the incoming small-$x$ gluons of the dense nucleus (the CGC) through a $2 \rightarrow 1$ process. After the eikonal scattering process, the outgoing parton $f$ has momentum $p_f = (q_t e^{y_f}/\sqrt{2}, q_t e^{-y_f}/\sqrt{2}, \vec{q}_t)$ with rapidity $y_f$ and produces a massive hadron $h$ with the momentum fraction $z_h$ and mass $m_h$. The detected hadron has a momentum $p_h = (m_t e^{y_h}/\sqrt{2}, m_t e^{-y_h}/\sqrt{2}, \vec{p}_t)$, where $m_t$ is the transverse mass $m_t = \sqrt{m_h^2 + p_t^2}$ and $y_h$ is the hadron rapidity. The Feynman-$x$ of the produced hadron is given by $x_F = p_h^+ / P_p^+ = m_t e^{y_h}/\sqrt{s}$. In the eikonal approximation we have $z_h = p_h^+ / p_f^+ = x_p / x_F$ because of $x_p = q_t e^{y_f}/\sqrt{s}$. The internal variables of the parton $f$, i.e. $q_t$ and $y_f$, therefore, can be expressed in terms of $p_t$, $m_t$, $y_h$, and $x_p$: $q_t = p_t / z_h = p_t x_p / x_F$ and $y_f = y_h + \log(m_t / p_t)$. Another important internal variable is the momentum fraction $x_A$ carried by the small-$x$ gluons in the nucleus. It is related to that of the impinging projectile parton $x_p$ in the eikonal approximation by $x_A = x_p e^{-2y_f}$ (see Appendix B in [22]). This leads to the rapidity of the gluons,

$$y_A = \log(1/x_A) = \log(1/x_p) + 2y_h + \log(m_t^2 / p_t^2).$$  \hspace{1cm} (1)

Here, complying with standard practice, we have defined the direction of the final hadron produced in the “proton fragmentation region” as the positive $z$-direction, i.e. $y_h > 0$.

In the framework of collinear factorization, the single-inclusive hadron production cross section can be written as a convolution of the inclusive parton $f$ production cross section with the PDF of the initial parton $i$, $f_{i/p}$, and with the FF of $f$ into a hadron $h$, $D_{h/f}$:

$$\frac{d\sigma(pA \rightarrow hX)}{\pi dy_h dp_t^2 d^2 b} = \frac{d\sigma(pA \rightarrow hX)}{\pi dy_h dm_t^2 d^2 b} \sum_{i,f} \int dx_p dz_h dy_f dq_t^2 f_{i/p}(x_p, Q_f^2) \frac{d\sigma(iA \rightarrow fX)}{\pi dy_f dq_t^2 d^2 b} D_{h/f}(z_h, \mu_f^2) \times \delta(m_t^2 - M(z_h, q_t)^2) \delta(y_h - Y(y_f, z_h, q_t)), \hspace{1cm} (2)$$
where \( M(z_h, q_f)^2 = m_h^2 + (z_h q_f)^2 \) and \( Y(y_f, z_h, q_f) = y_f - \log(m_t/z_h q_f) \) as derived above. \( \vec{b} \) is the impact parameter with respect to the center of the nucleus. The factorization scale of the PDF is denoted by \( Q_f \) and the fragmentation scale of the FF by \( \mu_f \), both of which are hereafter set to \( Q_f = \mu_f = p_t \).

After integration over \( y_f, q_f^2 \) and \( z_h \), we obtain the impact-parameter averaged single-inclusive hadron production cross section,

\[
\frac{d\sigma(pA \rightarrow hX)}{\pi dY_h dp_t^2} = \frac{1}{2\pi S_A} \int_{x_F}^{1} dx_p \int_{0}^{R_A} db \times \left[ f_{q/p}(x_p, Q_f^2) N_F \left( \frac{x_p}{x_F}p_t, y, b \right) D_{h/q} \left( \frac{x_F}{x_p}, \mu_f^2 \right) + f_{g/p}(x_p, Q_f^2) N_A \left( \frac{x_p}{x_F}p_t, y, b \right) D_{h/g} \left( \frac{x_F}{x_p}, \mu_f^2 \right) \right].
\]

(3)

Here we used \( d\sigma(iA \rightarrow fX)/\pi dy_f dq_f^2 db = q_f^2 \delta(q_f^2 - x_p p_t^+)^N(q_t, y_A, b)/(2\pi)^2 \), where \( N(q_t, y_A, b) \) is the scattering probability of dipoles from the nucleus. \( N_F \) corresponds to a projectile quark impinging with the target small-\( x \) gluons while \( N_A \) is for a projectile gluon. The indices \( q \) in \( f_{q/p} \) and \( D_{h/q} \) are summed over all quark species. For simplicity we assume a spherical nucleus of radius \( R_A \) with a sharp edge, where its profile function is given by \( T(b) = 2\sqrt{R_A^2 - b^2} \). \( S_A \) is the transverse area of the nucleus (\( = \pi R_A^2 \)).

The same expression as Eq. (3) was derived in Ref. [22] assuming massless hadrons and hence this equation is consistent with that of [22] in the massless limit. To leading \( \log p_t^2 \) accuracy it was verified in [22] that the recoil of the projectile parton by hard gluon radiation can be converted into the \( Q^2 \)-evolution of the PDF and FF according to the full DGLAP [23] evolution equations. The recoil effect are important for the interpretation of the forward-rapidity data from RHIC.

The dipole profiles of transverse size \( r_t \) at an impact parameter \( b \) are defined in the fundamental and adjoint representations of \( SU(N_c) \), respectively, as

\[
N_F(\vec{r}_t, y_A, \vec{b}) \equiv \frac{1}{N_c} \text{Tr}_c \langle 1 - V^\dagger(\vec{b} - \vec{r}_t/2)V(\vec{b} + \vec{r}_t/2) \rangle,
\]

\[
N_A(\vec{r}_t, y_A, \vec{b}) \equiv \frac{1}{N_c^2 - 1} \text{Tr}_c \langle 1 - U^\dagger(\vec{b} - \vec{r}_t/2)U(\vec{b} + \vec{r}_t/2) \rangle.
\]

(4)

where \( V \) and \( U \) denote Wilson lines along the light cone [40] in the corresponding representation, and their correlators are averaged over the color source in the nucleus.

For the actual description of the dipole profiles we rely on the KKT model [17], which is a phenomenologically reasonable parameterization facilitating comparisons with experimental data, rather than on solutions of the intricate JIMWLK equations [40], which is not yet feasible. For a dipole in the adjoint representation,

\[
N_A(\vec{r}_t, y_A, \vec{b}) = 1 - \exp \left[ -\frac{1}{4}(r_t^2 Q_s^2(y_A, b))^{\gamma(r_t, y_A, b)} \right],
\]

(5)

where \( Q_s(y_A, b) \) is the saturation scale of the target nucleus and \( \gamma(r_t, y_A, b) \) the anomalous dimension of its gluon distribution with saturation boundary condition. \( \gamma = 1 \) reduces [34]
to the Golec-Biernat-Wüsthoff saturation model [11] or the classical MV model [28]. The dipole in the fundamental representation, \( N_F \), differs by a factor of \( Q_s^2 \to Q_s^2 C_F / C_A = 4/3 Q_s^2 \) [17].

The saturation scale \( Q_s \) at \( b = 0 \) is given for a nucleus of mass number \( A \) (Au(197)) as [22]

\[
Q_s^2(y_A, b = 0) = A^{1/3} Q_0^2 \left( \frac{x_0}{x_A} \right)^3 = A^{1/3} Q_0^2 x_0^3 e^{\lambda y_A},
\]

where \( Q_0 \approx 1 \) GeV, \( \lambda \approx 0.3 \) and \( x_0 \approx 3.0 \times 10^{-4} \) are fixed by the DIS data [11]. The energy dependence of \( Q_s \) is controlled through the constant growth rate \( \lambda = \partial \log(\Lambda_{QCD}^2) / \partial \log(1/x_A) \), which is obtained from fixed-coupling LO BFKL evolution [12]. Since the squared saturation scale has dependence on the impact parameter, for instance, through the nuclear profile \( T(b) = T(b = 0) \sqrt{1 - (b/R_A)^2} \) in the hard sphere approximation of the nuclear target and the pointlike proton projectile, we define the saturation scale as [11]

\[
Q_s^2(y_A, b) = Q_s^2(y_A, b = 0) \sqrt{1 - (b/R_A)^2}.
\]

This naive approximation for the nuclear surface is sufficient for minimum bias observables.

The anomalous dimension \( \gamma \) is parameterized as [26]

\[
\gamma(r_t, y_A, b) = \gamma_s + (1 - \gamma_s) \frac{|\log(1/r_t^2 Q_s^2(y_A, b))|}{\lambda y_A + |\log(1/r_t^2 Q_s^2(y_A, b))| + d \sqrt{y_A}},
\]

where \( \gamma_s \approx 0.627 \) is the anomalous dimension for BFKL evolution [24] with saturation boundary condition, i.e. for evolution along the saturation line [42], and \( d \) is a free parameter which is fitted to experimental data. Throughout this paper, we will make replacement \( \gamma(r_t, y_A, b) \to \gamma(1/q_t, y_A, b) \). As indicated in [26], this parameterization of \( \gamma \) stays within \( \gamma = 0.6 \sim 0.8 \) at large rapidity \( y_b \) over a comparatively wide range of \( p_t = 1 \sim 5 \) GeV.

Below we focus on the minimum-bias cross section obtained by impact-parameter averaging of Eq. (3). Since in the integrand of (3) the impact-parameter dependence is carried only by the saturation scale, we take the average value of \( Q_s^2(b) \) with respect to \( b \) instead of integrating (3) over \( b \):

\[
\langle Q_s^2(b) \rangle = \frac{\pi}{S_A} \int_0^{R_A} db^2 Q_s^2(y_A, b) = \frac{2}{3} Q_s^2(y_A, b = 0).
\]

We have used Eq. (7) where this approximation is valid to good accuracy [14] 4.

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3As emphasized in [26], this remains inside the (at least, extended) geometric scaling regime. When going beyond the scaling regime like higher-\( p_t \) or lower rapidities, Eq. (3) assumes that crossover from the scaling regime to the perturbative one is very slow and smooth. A similar behavior of \( \gamma \) was discussed in Ref. [33].

4Refs. [22, 29] employ an effective saturation scale for minimum bias collisions such as \( Q_s^2(b) = A_{eff} Q_0(x_0/x_A)^\lambda \) with \( A_{eff} = 18.5 \), where the factor \( A_{eff}^{1/3} \) is by \( \sim 30\% \) smaller than the corresponding factor \( (2/3) A^{1/3} \) in [10].
Then, Eq. (3) reads

\[
\frac{d\sigma_{m.b.}(pA \rightarrow hX)}{\pi dy_h dp_t^2} = \frac{1}{(2\pi)^2} \int_{x_F}^1 dx_p \frac{x_p}{x_F} \left[ f_{q/p}(x_p, Q_f^2)N_F \left( \frac{x_p}{x_F}p_t, y_A, \langle Q_s^2(b) \rangle \right) D_{h/q} \left( \frac{x_F}{x_p}, \mu_f^2 \right) \right] + f_{g/p}(x_p, Q_f^2)N_A \left( \frac{x_p}{x_F}p_t, y_A, \langle Q_s^2(b) \rangle \right) D_{h/g} \left( \frac{x_F}{x_p}, \mu_f^2 \right),
\]

where the Fourier transform of the dipole profile functions is given by

\[
N_{A,F}(q_t, y_A, \langle Q_s^2(b) \rangle) = -\int d^2 r_t e^{i\vec{q} \cdot \vec{r}_t} N_{A,F}(r_t, y_A, \langle Q_s^2(b) \rangle) = -2\pi \int_0^\infty dr_t J_0(r_t q_t) N_{A,F}(r_t, y_A, \langle Q_s^2(b) \rangle).
\]

3 Forward $\pi^\pm$ production in d+Au collision at RHIC

We are now in the position of applying our results (10) to deuteron-gold collisions at RHIC energy $\sqrt{s} = 200$ GeV, at large rapidity ($y_h = 3.0$). We first address minimum bias $\pi^\pm$ production observed by BRAHMS collaboration [15]. This process is dominated by valence quarks since $x_p \sim 0.1$ on the deuteron side and hence might be affected by initial quark correlation (like diquarks) with large $x$ momentum inside deuteron.

In this section, our interest is twofold: First, we check the validity of our leading order CGC formalism using standard parameterizations of the PDF and the FF, and determine the free parameter $d$ as well as the $K$-factor prescribed to NLO corrections to reproduce well the experimental data. Second, using those very same parameters ($d$ and $K$) we investigate the diquark contribution.

For a deuteron projectile, we treat its parton distributions as a simple superposition of those in a proton and a neutron without any nuclear modification, which indeed has negligible effects (less than 5%) for the deuteron. The parton distributions in the neutron are obtained through isospin symmetry, $f_{u,\bar{u}/p} = f_{d,\bar{d}/n}$ and $f_{d,\bar{d}/p} = f_{u,\bar{u}/n}$, and other parton species have the same distribution as those in the proton [16]. This isospin symmetry may affect the final pion productions composed of flavor $SU(2)$ light quarks, where it leads to no difference between light $\pi^+(u\bar{d})$ and $\pi^-((\bar{u}d)$ productions to good accuracy. Actually, such behavior can be seen for $p_t \lesssim 4$ GeV in the BRAHMS data plotted in Fig. 2.

We first use the leading order (LO) CTEQ5 PDFs [17]. Then, we have several choices for the FFs, e.g. the Kretzer [18], KKP [31] and AKK [35] sets. The FFs by Kretzer assume that the charge-conjugation symmetry $D_{q/\pi^\pm} = D_{q/\pi^\pm}$ holds at low input scale. Both KKP and AKK FFs (the latter provide flavor-dependent FFs of light quarks as an update of KKP) provide only the average of charged pion FFs. Therefore, all three FFs give the same neutral pion FF and should show almost the same $p_t$-spectra at LO.

We checked that the result with the LO Kretzer FFs gives a very similar curve to that of the LO KKP FFs (but different $K$-factor) over the range $1 \lesssim p_t \lesssim 5$ GeV, with a pion mass of $m_\pi = 0.14$ GeV. The result with the LO Kretzer FFs is in good agreement with the data at $d = 0.6$ and $K = 1.4$, while the LO KKP FFs require $d = 0.6$ and $K = 1.0$. 

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Figure 2: $p_t$ spectra of $\pi^+$ and $\pi^-$ from $d+Au$ collisions compared to BRAHMS preliminary minimum bias data ($y_h = 3.0$). The lines show the CGC results with the LO Kretzer FFs and $K = 1.4$, where either the LO CTEQ5 or diquark PDFs are used.

The main difference is due to the way that gluons fragment to pions; its contribution in the KKP set is larger than that in the Kretzer one [2]. The sensitivity of the results to the value $d$ is not so large in the range of $d = 0.6 \sim 1.2$ and is visible only at high $p_t$, where the data however have large errors. In Fig. 2 we show the LO Kretzer result with $K = 1.4$ as the dotted line.

Next, in order to investigate the diquark contribution to this process, we make use of the PDFs given by the Stockholm diquark model [31]. It assumes that only scalar diquarks are genuine bound states and other axial-vector diquarks are negligible. As is well known, the diquark is classified into $\bar{3}_c$ and $6_c$ channels in color $SU(3)$. The scalar diquark belongs to the $\bar{3}_c$ channel while the axial-vector diquark is in the $6_c$ representation. The reason why we adopt only the former is that the one-gluon exchange between quarks is attractive in the $\bar{3}_c$ channel while repulsive in the $6_c$ channel. Thus, for a first estimate it is rather plausible to focus on the more tightly bound scalar-diquark. This parameterization is obtained from a fit to the proton structure function $F^p_2$ observed in high energy $e+p$ collisions at SLAC, BCDMS and EMC, over the wide range $1 < Q^2 < 200 \text{ GeV}^2$.

Ref. [31] has two possible parameterizations which differ by the choice of a mass scale in the diquark form-factor. We choose the parameter set with $M^2 = 10 \text{ GeV}^2$ and a dipole
form-factor

\[ F^2(Q_f^2) = \frac{1}{(1 + Q_f^2/M^2)^2}. \]

It ensures that with larger \( Q_f^2 \) all diquarks are resolved and the usual quark PDFs are recovered. The scale \( M \) is related to the internal binding energy of the diquark and \( M^2 = 10 \text{ GeV}^2 \) corresponds to a diquark radius of about 0.2 fm, which is obtained from the mean-square radius \( \langle r^2 \rangle = 6dF^2(Q_f^2)/dQ_f^2 |_{Q_f^2=0} \). It yields larger contribution from diquark component due to its tightly bound state than that of \( M^2 = 3 \text{ GeV}^2 \). The parameterization for each parton, then, is given as [31]

\[
\begin{align*}
x_p u(x_p, Q_f^2) &= \frac{(0.52 + 0.37s)x_p^{0.45+0.12s} (1-x_p)^{2.2+2.39s} - x_p^{0.71-0.02s} (1-x_p)^{7.3-0.51s}}{(0.52 + 0.37s)B(0.45 + 0.12s, 3.2 + 2.39s) - B(0.71 - 0.02s, 8.3 - 0.51s)} \\
&\quad + \frac{x_p^{0.9-0.83s} (1-x_p)^{5.0-1.88s}}{B(0.9 - 0.83s, 6.0 - 1.88s)} [1 - F^2(Q_f^2)], \\
x_p d(x_p, Q_f^2) &= \frac{x_p^{0.9-0.83s} (1-x_p)^{5.0-1.88s}}{B(0.9 - 0.83s, 6.0 - 1.88s)} [1 - F^2(Q_f^2)], \\
x_p s(x_p, Q_f^2) &= x_p \bar{u}(x_p, Q_f^2) - x_p \bar{d}(x_p, Q_f^2) = x_p \bar{s}(x_p, Q_f^2) \\
&= (0.35 - 0.06s)(1 - x_p)^{6.89+0.75s} \\
x_p f_{DQ}(x_p, Q_f^2) &= \frac{x_p^{0.93-0.52s} (1-x_p)^{1.5-1.1s}}{B(0.93 - 0.52s, 2.5 - 1.1s)} F^2(Q_f^2),
\end{align*}
\]  

(13)

where we assumed flavor \( SU(3) \) symmetry for the sea quark distributions, and \( f_{DQ} \) denotes the diquark distribution, \( B(\mu, \nu) \) Euler’s Beta function and \( s = \log[\log(Q^2_f/\Lambda^2)/\log(Q_0^2/\Lambda^2)] \) with \( \Lambda = 0.2 \text{ GeV} \) and \( Q_0^2 = 4 \text{ GeV}^2 \).

When one views the d+Au collisions in this diquark picture, one faces other theoretical uncertainties besides modelling of the diquark distribution in the deuteron: the breakup of diquarks scattered off the CGC and the diquark fragmentation into hadrons. The former depends on the collision dynamics between diquarks with a finite size and the CGC. To take into account this breakup probability in a simple way, we express it as the dipole form-factor of [12] with \( M^2 = 10 \text{ GeV}^2 \) as done in Ref. [7]. This probability goes to zero in the high-\( p_t \) limit (\( Q_f = p_t \)), and then the diquark bound state is completely broken inside the CGC, where the constituent \( u \) and \( d \) quarks fragment independently.

Although so far various models for the diquark fragmentation have been proposed [32], they have focused mainly on how the FFs behave as a function of \( x \), leaving aside scaling violation by the \( Q^2 \) evolution. For our present investigation of \( p_t \)-spectra of pions and baryons from diquarks it is important to account for the scaling violation. For this purpose, as a naive model estimate we express the diquark FFs in terms of the LO Kretzer FFs, which describe fragmentation into charged hadrons separately. We approximate the

\[ \text{footnote text} \]
(ud)-diquark FFs into charged pions as

\[D_{\pi^+/(ud)/DQ(ud)}(z_h, \mu_f^2) = D_{\pi^+/d}^{Kr-LO}(z_h, \mu_f^2)F^2(\mu_f^2)\]

\[+ [D_{\pi^+/(ud)/DQ(ud)}^{Kr-LO}(z_h, \mu_f^2) + D_{\pi^+/d}^{Kr-LO}(z_h, \mu_f^2)](1 - F^2(\mu_f^2)),\]

\[D_{\pi^-(ud)/DQ(ud)}(z_h, \mu_f^2) = D_{\pi^-/u}^{Kr-LO}(z_h, \mu_f^2)F^2(\mu_f^2)\]

\[+ [D_{\pi^-/(ud)/DQ(ud)}^{Kr-LO}(z_h, \mu_f^2) + D_{\pi^-/u}^{Kr-LO}(z_h, \mu_f^2)](1 - F^2(\mu_f^2)),\] (14)

where we multiplied the FFs by the breakup probability. The first term is the contribution where the valence quarks of the produced pions do not include either one of the quarks from the scattered diquark, and thus the pion production occurs via quark-antiquark pair creation from the vacuum, like \(D_{\pi^+/d}\) or \(D_{\pi^-/u}\). The second contribution corresponds to diquark breakup followed by independent fragmentation of each quark into pions.

In Fig. 2 we plot the diquark result, using Kretzer FFs both for the diquarks as well as for the single partons, as the dashed line (with \(d = 0.6\) and \(K = 1.4\)). This curve, however, has a problem, since it should approach the result with CTEQ5+Kretzer (dashed line) at high \(p_t\) due to \(F^2(p_t^2) \rightarrow 0\); that is, the diquark contribution should vanish in this limit and the single parton picture should be restored again, but apparently the former is much harder than the latter. To improve this unfavorable behavior, we simply change the dipole form-factor \(F^2(p_t^2)\) to \(F^4(p_t^2)\) for both the PDFs and the breakup probabilities. This result is shown as the solid line, which is softer at high \(p_t\) and is in good agreement with the CTEQ5+Kretzer result, especially at high \(p_t\). Since this change of the form-factor increases the diquark radius from 0.2 fm to 0.3 fm \(^6\), the new form-factor becomes more sensitive to \(p_t\) and the diquarks are broken up more easily. The solid line is still somewhat harder than the dotted one at low \(p_t\), but it remains a good fit to the data. Thus, the result in the diquark picture can reproduce well the data within error bars. As expected, this implies that the diquark contribution to pion production is not very large.

4 Forward \(p\) and \(\bar{p}\) production

In this section we examine forward proton or anti-proton production in d+Au collisions at RHIC and apply Eq. (10) to obtain the cross sections within the same setup as for pion production (discussed above). At the RHIC energy \(\sqrt{s} = 200\) GeV and in the deuteron fragmentation region \((y_h = 3.0)\), we first show the BRAHMS preliminary data \(^{[45]}\) for minimum bias cross sections of \(p\) (solid circle) and \(\bar{p}\) (open circle) production in Fig. 3.

Unlike for pions, charge symmetry is broken more strongly for \(p\) and \(\bar{p}\) production. The proton cross section exceeds that of anti-protons by one order of magnitude for \(1.5 < p_t < 3.5\) GeV. This feature can also be seen in p+p collisions at large rapidity and high \(p_t\) \(^{[45]}\). Such a large asymmetry between them, and in particular the high abundance of protons, may be associated with the production mechanism involving the fragmentation of diquarks into protons.

\(^6\)This radius is still smaller than the other parameter set of \(^{[7]}\) with \(M^2 = 3\) GeV\(^2\), where \(\sqrt{\langle r^2 \rangle} \approx 0.4\) fm.
Figure 3: $p_t$ spectra of $p$ and $\bar{p}$ in $d+Au$ collisions from BRAHMS preliminary minimum bias data ($y_h = 3.0$) compared to computations within the CGC formalism ($K = 1.0$). The lines show the results with the common LO CTEQ5 PDFs and the Lund FFs rescaled to either the KKP or AKK FFs with $K = 1.0$. The lines with symbols show the results with the diquark PDFs.

To our knowledge, the best parameterization of baryon FFs handling the $Q^2$ evolution to some extent is given in Ref. [33], where the FFs are obtained by parameterizing results computed by Monte Carlo simulations in the Lund string model, evaluated at invariant mass $W = 2Q_f$. The advantage of using this parameterization is that we can treat proton and anti-proton production separately while the KKP and AKK FFs only provide averages. Moreover, the Lund string model JETSET also models mass effects in the fragmentation which are important when the momentum of the outgoing parton is not much larger than the mass of its daughter hadrons. However, the $Q^2_f$ dependence of these Lund FFs $D_{B/q}(z_h, Q^2_f)$ has not been implemented very carefully: Ref. [33] only requested that the average multiplicity $\int_0^1 dz_h D_{B/q}(z_h, Q^2_f)$ satisfies the $Q^2_f$ dependence of the experimental data. This treatment would not describe correctly scaling violations, especially in the high-$z_h$ regions, where the FFs change steeply.

To check the $Q^2_f$ behavior, we computed the average $p_t$-distribution of protons and anti-protons, $(p + \bar{p})/2$, and compared with the results obtained with the NLO KKP and NLO AKK FFs, using the LO CTEQ5 PDF for all three FF sets.\footnote{The reason why we employ the NLO KKP or NLO AKK FFs is as follows: Recent STAR data show...}
and AKK-NLO/Lund are shown in Fig. 4, where the baryon mass is set to $m_{p,\bar{p}} = 0.938$ GeV and we choose $d = 0.6$ and $K = 1.0$.

![Figure 4: Summed $p + \bar{p}$ yields obtained with the NLO KKP or NLO AKK FFs divided by the Lund FFs.](image)

We see that the ratios change steeply as a function of $Q_f = p_t$ and that the $Q_f^2$ dependence of the Lund FFs is significantly harder than those of KKP and AKK sets. In more detail, the cross section in the Lund model is larger than that with the KKP set for $p_t > \sim 1.5$ GeV, while the behavior is opposite for the AKK set at $p_t < \sim 3.5$ GeV. This forces us to improve the Lund model FFs to enforce a $Q_f^2$-dependence that matches that of the KKP or AKK FFs. In this figure we can also see that the KKP FFs are smaller by a factor of $5 \sim 12$ than the AKK FFs in the region $1.5 < p_t < 5$ GeV and the deviation becomes larger at higher $p_t$.

For that purpose, we fit the ratios using each function via

\[
\begin{align*}
\text{KKP(NLO)/Lund} &= \frac{1}{[-0.01 + (p_t/1.56)^{1.36}]^{2.01}}, \\
\text{AKK(NLO)/Lund} &= \frac{1}{[0.12 + (p_t/3.70)^{1.91}]^{1.38}}. 
\end{align*}
\]

(15)

that high-$p_t$ $p + \bar{p}$ yields around midrapidity in $d+Au$ collisions at $\sqrt{s} = 200$ GeV agree better with NLO pQCD calculations using the AKK rather than the KKP FFs [2]. The latter is known to underestimate the data by one order of magnitude and so it would be best to adopt the AKK FFs for baryon production. Unfortunately, the AKK FFs only provide a NLO parameterization. The use is inconsistent within our LO formalism, but despite such inconsistency, the use of AKK FFs for baryon productions is of importance to reproduce the yields. We also checked that the NLO KKP FFs are quite similar to the LO KKP FFs, but the former is a bit softer than the latter at high $p_t$. 

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These functions provide rather reasonable fits. Multiplying the Lund FFs by the fitting functions (15), we then rescale the $Q_f^2$ dependence of the FFs. We shall employ the same rescaling functions for both $p$ and $\bar{p}$ FFs.

For the diquark fragmentation into baryons we construct the following model in the same way as for pion production:

$$
D_{p(\bar{u}\bar{d})}/DQ(\bar{u}\bar{d})(z_h, \mu_f^2) = D_{Kr-LO}^{\pi^+} + \frac{1}{2}[D_{p/u}^{Lund}(z_h, \mu_f^2) + D_{p/d}^{Lund}(z_h, \mu_f^2)](1 - F^2(\mu_f^2)),
$$

$$
D_{\bar{p}(u\bar{u})}/DQ(\bar{u}\bar{d})(z_h, \mu_f^2) = \frac{1}{2}[D_{\bar{p}/u}^{Lund}(z_h, \mu_f^2) + D_{\bar{p}/d}^{Lund}(z_h, \mu_f^2)](1 - F^2(\mu_f^2)),
$$

(16)

where $D_{Kr-LO}$ and $D_{Lund}$ denote the LO Kretzer and Lund sets respectively. The first terms express the diquark fragmentation as a single entity and the second ones do as a system consisting of independent valence quarks. The first term of the proton FF is based on the conjecture that the fragmentation of the $(ud)$-diquark proceeds via pick-up of a $u$-quark from the vacuum, similar to that of $d$ into $\pi^+$, if the finite size of diquark is neglected. The first term of the anti-proton production treats a diquark as a single entity (like $u$ or $d$ quarks) and hence takes their average.

In Fig. 3 we plot the results of Eq. (10) rescaled to the KKP or AKK FFs, and using the form-factor $F^4$. For anti-protons, we show a total of four curves with the CTEQ5 or diquark PDFs, rescaled to either the KKP (dashed) or AKK (dotted) FFs. As indicated above, the KKP set significantly underestimates the anti-proton data by about a factor 10, independent of the choice of PDFs. The AKK FFs fit better although the data appears somewhat softer than our results. The difference between the CTEQ5 and diquark PDFs is generically small but increases at low $p_t$, similarly to the pion case discussed in Sec. 3. Both these curves are in good agreement with the data within the large error-bars.

For protons, we first show the CTEQ5 result rescaled to the AKK set as the solid line, which underestimates the data by a factor $2 \sim 3$ in the low-$p_t$ region, but approaches the data at high $p_t$. Next, using the same FFs, the diquark result is plotted as the square+solid line, which enhances the cross section at low $p_t$, but still underestimates the data. This feature is, however, unlike the cases of anti-protons or pions and the enhancement of the cross section at low $p_t$ is favored by the data. As a somewhat extreme case we also plot a curve corresponding to “complete diquark survival” by setting the breakup probability to zero (i.e. $F^4(\mu_f^2) = 1$) as the open-circle+solid line. This result is very close to the data in the low-$p_t$ region.

The $p/\pi^+$ and $\bar{p}/\pi^-$ ratios will be relevant for investigating the $p_t$ dependence of baryon production more precisely. In Fig. 4 we plot the data extracted by taking the ratios of the corresponding data points from Figs. 2. Here, as seen in Fig. 2 there are no data points of the $\pi^-$ in the range $1.5 \lesssim p_t \lesssim 2.1$ GeV. To interpolate the lack in this interval (6 points), we substitute the corresponding data points of the $\pi^+$. This is reasonable because for a deuteron projectile isospin symmetry seems to work well at low $p_t$ as discussed above. This figure shows that the $p/\pi^+$ ratio is almost $0.6 \sim 0.9$ which is almost flat, while the
Figure 5: Ratios of $p/\pi^+$ and $\bar{p}/\pi^-$. The experimental points were obtained by taking the ratios of the corresponding data from Figs. 2 and 3. The lines show the results with the common LO CTEQ5 PDFs for both baryons and pions, where the former uses the Lund FFs rescaled to the AKK set with $K = 1.0$ and the latter the LO Kretzer FFs with $K = 1.4$. The lines with symbols show the results with the diquark PDFs.

The $\bar{p}/\pi^-$ ratio is 0.04 ~ 0.07, smaller by one order of magnitude, where at $p_t \lesssim 2.5$ GeV we see no significant $p_t$-dependence, but above this point the $p_t$-distribution starts to decrease largely with $p_t$.

For the $\bar{p}/\pi^-$ ratio, we plot two results: The dotted line uses the common LO CTEQ5 PDFs for both $\bar{p}$ and $\pi^-$, and the Lund FFs rescaled to the AKK set for $\bar{p}$ and the Kretzer FFs for $\pi^-$, with different $K$-factors: $K = 1.0$ for $\bar{p}$ and $K = 1.4$ for $\pi^-$. The cross-dotted line replaces the PDFs by the diquark set, with otherwise the same “ingredients”. These curves explain the data quantitatively, especially at low $p_t$, with small discrepancies between the CTEQ5 and diquark PDFs. The monotonic decrease with $p_t$ is supported by the data except the last data point.

For the $p/\pi^+$ ratio, we show the same curves. First, we employ CTEQ5 PDFs with the Lund FFs rescaled to the AKK set for $p$ and to the Kretzer FFs for $\pi^+$ (solid line). Second, we switch to the diquark PDFs with the same sets of FFs. The experimental data cannot be explained by the use of the CTEQ5 PDFs any more, which is too low by a factor

---

8This ratio is above unity around $p_t = 1.5$ GeV, because of the relatively large underestimate of our $\pi^+$ cross section as compared to the $\pi^+$ data.
of 2 ∼ 3 at low \( p_t \), as already seen in Fig. 3. The result is almost flat over a wide range of \( p_t \) and seems to work well only in the high-\( p_t \) region. Switching to the diquark PDFs may explain at least partly the large baryon excess seen in the data at low \( p_t \). Diquark effects disappear quite rapidly with increasing \( p_t \) which is not really supported by the data. The prediction of the “surviving diquark” picture is moderately larger at low \( p_t \) and perhaps offers a reasonable explanation of the data.

5 Discussion

The physical process of \( d+Au \) collisions illustrated in Fig. 1 shows that the cross section is factorized into three parts, i.e. into a convolution of the form \( f_{q,g/d} \otimes N_{F,A} \otimes D_{h/q,g} \) like in Eq. (10) (refer to ref. [22] for verifying this factorization to leading logarithmic accuracy 9). These three functions are universal, independent of the processes and associated with infrared hadron or nuclear structures, for example with respect to the fact that we employ the same forms of \( N_{F,A} \) for meson (\( D_{M/q,g} \)) and baryon (\( D_{B/q,g} \)) production. In other words, the \( p_t \) spectra may be determined by universal dipole profiles, irrespective of detected hadron species. These dipole profiles \( N_{F,A} \) describe the interaction of hard partons scattering off the fields of nucleus. Bremsstrahlung occurring either “before” or “after” the propagation through the nucleus yields logarithmic radiative corrections, which are absorbed into the DGLAP \( Q^2 \)-evolution of \( f_{q,g/d} \) and \( D_{h/q,g} \) [22].

As seen from Figs. 2, 3, the data displays a common power-law behavior over the range of \( 1 < p_t < 5 \) GeV, not an exponential form predicted by the parton recombination model [19]. In more detail, it is close to the exponential until \( p_t \sim 2 - 3 \) GeV, but above this it appears rather close to a power-law. In the CGC formalism, the dipole profile displays a power-law behavior for the \( p_t \)-distribution at high rapidity, whose specific shape is governed by the anomalous dimension \( \gamma \) [22, 26].

In fact, our dipole profiles give even slightly harder \( p_t \)-distributions than the data for both \( p \) and \( \bar{p} \) if we employ the CTEQ PDFs. If the factorization and the CGC framework work well, however, such a deviation from the data will be due mainly to the \( Q^2 \)-evolution FFs (here the Lund baryon FFs rescaled by the AKK FFs). If this is the case, we should determine more carefully the scaling violation of the baryon FFs in the Lund string scheme. Work on new parameterizations by means of Lund JETSET simulation will be reported elsewhere [51].

For \( p \) production, a remarkable point is that the diquark picture does have a noticeable effect on the spectra at large rapidity. In fact, for the specific parameterization employed here, the slope is even steeper than seen in the data. This comes from mainly the \( p_t \)-dependence in the form-factors, which are used both in the PDFs and in the breakup probability of diquarks. As a reference, we considered the case of completely pointlike diquarks, which are not broken by scattering off the CGC at all. At a first glance this result

9The process violating the factorization like the second rescattering of the hard parton with the target after the first rescattering and the succeeding gluon radiation vanishes in the high energy limit [29, 50], at least if we calculate in light-cone gauge.
is closer to the data compared to the results with breakup. However, this should be taken with caution, because this "surviving diquark" picture maximizes the contribution from diquarks and is inconsistent with the baryon suppression scenario in the deep saturation regime \[10\].

For a more quantitative comparison to experimental data we require more precise information of the $Q^2$-dependence of diquark distribution and fragmentation functions (which can not be measured in $e^+e^-$ annihilation like ordinary quark FFs). It might also be possible that the deuteron is composed of three-diquark system \[52\]. In the present formulation the third constituent diquark has not been considered. This contribution may lead to additional baryon formation.

6 Summary

We have focused on forward $p$ and $\bar{p}$ production in d+Au collisions using universal dipole profiles developed in \[22\] within the CGC formalism. This formalism is in a good agreement with the high $p_t$-spectrum of forward pions at $y = 3.0$, where its power-law behavior is described well by the dipole profile in the formalism. However, the abundance of protons can not be explained by the usual PDFs (like CTEQ5) and Lund FFs (rescaled by AKK FFs), while that of anti-protons is consistent with our results. Similarly, large deviations of the $p/\pi^+$ ratio from Pythia simulations is already observed for forward $p+p$ collisions \[15\]. One possible explanation for this extra production of protons at $p_t = 1 \sim 3$ GeV compared to standard PDFs is the direct formation of protons from (dominantly scalar) diquarks, whose contribution is incorporated into our formulation, including a diquark form-factor and fragmentation into hadrons. This diquark contribution is essential only for proton production but is less important for pions and anti-protons.

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