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CONSTRAINTS FOR NUCLEAR GLUON SHADOWING FROM DIS DATA

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ABSTRACT

The Q^2 dependence of the ratios of the cross sections of deep inelastic lepton–nucleus scattering is studied in the framework of leading twist, lowest order perturbative QCD. The $\log Q^2$ slope of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ is computed by using the DGLAP evolution equations, and shown to be sensitive to the nuclear gluon distribution functions. Four different parametrizations for the nuclear effects of parton distributions are studied. We show that the NMC data on the Q^2 dependence of $F_2^{\rm Sn}/F_2^{\rm C}$ rule out the case where nuclear shadowing (suppression) of gluons at $x \sim 0.01$ is much larger than the shadowing observed in the ratio $F_2^A/F_2^{\rm D}$. We also show that the possible nonlinear correction terms due to gluon fusion in the evolution equations do not change this conclusion. Some consequences for computation of RHIC multiplicities, which probe the region $x \gtrsim 0.01$, are also discussed.

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1 Introduction

Nuclear parton distributions (nPDF) are needed in the computation of inclusive cross sections of hard, factorizable, processes in high energy nuclear collisions. In the framework of collinear factorization and leading twist, it is possible to extract universal nuclear parton distributions $f_{i/A}(x, Q^2)$ from the measurements of deeply inelastic lepton-nucleus scattering (DIS) and hard processes in pA collisions such as the Drell-Yan process. The power corrections in the cross sections [1] and in the evolution equations [2, 3] can be neglected if the scales Q^2 and momentum fractions x involved are large enough. The dependence of the nPDF on the scale Q^2 is then given by the DGLAP evolution equations [4]. Analogously to the global analyses of the parton distributions of the free proton, the nPDF can be determined based on the DGLAP evolution, sum rules and fits to the data. Sets of nPDF like EKS98 [5, 6] (the code can be found in [7, 8]) and HKM [9] (Hirai, Kumano and Miyama, the code in [10]) have become available.

In the DGLAP analyses of parton distributions, the problem boils down to fixing the initial distributions for the DGLAP evolution at an initial scale Q_0^2. For the nPDF in particular, there are some uncertainties in the initial distributions due to the lack of experimental data. For instance, in the EMC region ($x \gtrsim 0.2$), gluons and sea quarks are so far not well constrained [11] (see also [12, 9] for more discussion). Also, it seems to be a common belief that the nuclear gluon distributions are constrained only very weakly by the DIS data at small values of x. In this letter our aim is to emphasize that this is not the case but very valuable constraints for the nuclear gluon distributions in the shadowing region ($x \sim 0.02$) can be obtained from the Q^2 dependence of the structure function ratio $F_2^{\rm Sn}/F_2^{\rm C}$ measured by the New Muon Collaboration (NMC) [13]. As first discussed in [14, 5], this is a consequence of the fact that in the lowest-order DGLAP evolution $\partial F_2(x, Q^2)/\partial \log Q^2 \sim \alpha_s x g(2x, Q^2)$ [15] at small x. One of the tasks in the present paper is to compare the results of the HKM analysis [9], where the constraint from the measured Q^2 dependence of $F_2^{\rm Sn}/F_2^{\rm C}$ has not been applied, with the results obtained with EKS98 [5, 6], which makes use of this constraint.

The Relativistic Heavy Ion Collider (RHIC) initiated the collider era for the search of the Quark-Gluon Plasma in ultrarelativistic AA collisions. From the point of view of the nPDF, the measurements of charged particle multiplicities $dN_{\rm ch}/d\eta$ in Au+Au collisions [16, 17, 18, 19] have been truly exciting, since production of semihard gluons at 1...2 GeV scales is expected to dominate particle and entropy production at $\sqrt{s} \gtrsim 200$ GeV. The measured multiplicities are thus a probe of the nuclear gluon distributions at $x \sim 0.01$ at these scales: in the models employing saturation [2, 20, 21, 22, 23], the multiplicity scales as $dN_{\rm ch}/d\eta \sim x g_A(x_{\rm sat}, Q^2_{\rm sat})$ (with possible powers of α_s added). In the two-component (hard+soft) models, such as HIJING [24, 25], the nPDF are probed through the perturbative (hard) minijet component.

Recently, in [25], it has been suggested based on the HIJING model that the RHIC data on multiplicities would indicate that the gluons were more strongly shadowed than the sea quarks. To challenge this interesting suggestion, we shall perform a DGLAP analysis of the nPDF based on the initial nuclear effects for quark and gluon distributions as given in [25]. As the second and main point of this paper, we shall show that in the leading twist, lowest order DGLAP framework the NMC data on the Q^2 dependence of $F_2^{\rm Sn}/F_2^{\rm C}$ [13] rules out very strongly shadowed gluons. In relation to HIJING, consequences of this observation for the parameter p_0, which determines the division into soft and hard components, will be discussed.

When the density of gluons in the wave functions of the colliding nuclei (or hadrons) becomes large enough, gluon fusion starts to play a role. This induces non-linearities into the QCD evolution equations [2, 3, 26, 27, 28] at small values of x and Q^2. As the last subject to study in this letter, we add the non-linear terms (GLRMQ) into the DGLAP equations for Sn and C nuclei. Comparison with the NMC data [13] shows, that a very strong gluon shadowing is ruled out even more clearly when the non-linearities are included.

2 DGLAP analysis and different parametrizations

Following the notation in our previous works [5, 6, 12], we define the nPDF through the nuclear effects,

$$R_i^A(x,Q^2) \equiv \frac{f_i^A(x,Q^2)}{f_i(x,Q^2)}, \qquad i = g,\ u,\ d,\ s,\ \bar{u},\ \bar{d},\ \bar{s}, \ldots, \tag{1}$$

where $f_i^A \equiv f_i^{p/A}$ is the number density distribution of a flavour i in a bound proton of a nucleus A, and f_i is the corresponding distribution in the free proton. The parton distributions of bound neutrons in isoscalar nuclei are obtained through isospin symmetry, $f_{u(\bar{u})}^{n/A} = f_{d(\bar{d})}^{p/A}$ and $f_{d(\bar{d})}^{n/A} = f_{u(\bar{u})}^{p/A}$, and we expect this to be a good approximation for non-isoscalar nuclei as well. Nuclear effects in the structure function F_2^A we define through

$$R_{F_2}^A(x,Q^2) = \frac{\frac{1}{A}F_2^A}{\frac{1}{2}F_2^{\mathrm{D}}} \approx \frac{\frac{1}{2}(F_2^{p/A}+F_2^{n/A}) + \frac{1}{2}(\frac{2Z}{A}-1)(F_2^{p/A}-F_2^{n/A})}{\frac{1}{2}(F_2^p+F_2^n)}, \tag{2}$$

where F_2 of the protons (neutrons) of the nucleus A is $F_2^{p(n)/A}(x,Q^2) = \sum_q e_q^2[xf_q^{p(n)/A}(x,Q^2) + xf_{\bar{q}}^{p(n)/A}(x,Q^2)]$, and where the small nuclear effects in D have been neglected.

In order to explicitly show the effect of nuclear gluon shadowing to the scale dependence of the ratio $F_2^{\mathrm{Sn}}/F_2^{\mathrm{C}}$, we shall study different parametrizations for the nuclear effects. On one hand, we will directly make use of the available results from the global DGLAP analyses of nPDF, EKS98 [5, 6] and HKM [9]. On the other hand, we perform the DGLAP evolution for nPDF by taking the initial modifications $R_i^A(x,Q_0^2)$ from two different parametrizations in which the Q^2 dependence is assumed to be negligible:

- HPC parametrization [29] which is based on [30]:

$$\begin{aligned} R_i^A(x) = R_{F_2}^A(x) \quad &= \quad R_{\mathrm{sh}}\frac{1+c_Dc_A(1/x-1/x_{\mathrm{sh}})}{1+c_AA^{p_A}(1/x-1/x_{\mathrm{sh}})}\Theta(x_{sh}-x) + \\ &+(a_{\mathrm{emc}}-b_{\mathrm{emc}}x)\Theta(x-x_{sh})\Theta(x_f-x) + R_f\left(\frac{1-x_f}{1-x}\right)^{p_{\mathrm{f}}}\Theta(x-x_f) \end{aligned} \tag{3}$$

  where the different regions are matched together by setting $R_{\mathrm{sh}} = a_{\mathrm{emc}} - b_{\mathrm{emc}}x_{\mathrm{sh}}$, $R_f = a_{\mathrm{emc}} - b_{\mathrm{emc}}x_f$ and $a_{\mathrm{emc}} = 1 + b_{\mathrm{emc}}x_{\mathrm{emc}}$. The A dependence of b_{emc} is $b_{\mathrm{emc}} = p_{\mathrm{emc}}[1 - A^{-1/3} - 1.145A^{-2/3} + 0.93A^{-1} + 0.88A^{-4/3} - 0.59A^{-5/3}]$ from Ref. [31]. A fit to DIS data results in $p_A = 0.10011$, $c_A = 0.0127343$, $c_D = 1.05570$ $x_{\mathrm{sh}} = 0.154037$, $x_{\mathrm{emc}} = 0.275097$, $p_{\mathrm{emc}} = 0.525080$, $x_{\mathrm{f}} = 0.742059$, and $p_{\mathrm{f}} = 0.320992$.

- New HIJING parametrization [25] which replaces the obsolete one in HIJING [24] (see also the discussion in [12]). In the new parametrization, the quark sector is fitted to modern DIS data, while the gluon sector, especially shadowing and antishadowing are constrained by the requirement that (with the updated $\sqrt{s}$ dependence of the cut-off parameter p_0 for the transverse momentum of minijets) HIJING reproduces the measured charged-particle multiplicities in Au+Au collisions at RHIC. In this way, gluon shadowing is suggested to be much stronger than that of quarks:

$$\begin{aligned} R_q^A(x) = R_{\bar{q}}^A = R_{F_2}^A(x) &= 1.0 + 1.19 \log^{1/6} A \, (x^3 - 1.2x^2 + 0.21x) \\ &\quad - s_q \, (A^{1/3} - 1)^{0.6} (1 - 3.5\sqrt{x}) \exp(-x^2/0.01) \end{aligned} \tag{4}$$

$$\begin{aligned} R_g^A(x) &= 1.0 + 1.19 \log^{1/6} A \, (x^3 - 1.2x^2 + 0.21x) \\ &\quad - s_g \, (A^{1/3} - 1)^{0.6} (1 - 1.5x^{0.35}) \exp(-x^2/0.004), \end{aligned} \tag{5}$$

with $s_q = 0.1$ and $s_g = 0.24 \ldots 0.28$. Below, we shall use the value 0.24.

In Fig. 1, in the left panels, we plot the nuclear effects for gluons in tin ($A = 117$, top row) and carbon ($A = 12$, third row) nuclei at a scale $Q_0^2 = 2.25$ GeV^2, as given by EKS98 and HKM, and by the HIJING and HPC parametrizations above. For further discussion, we also show $R_{F_2}^A(x, Q_0^2)$ from Eq. (2) (2nd and 4th rows, correspondingly). With the EKS98, HIJING and HPC parametrizations, we have used the MRST (central gluon) LO PDFs of the free proton [32, 8]. The HKM results are obtained directly from the HKM code [10], based on [9] and where the LO MRST (c-g) PDFs are also used.

As seen in Fig. 1, for the gluons there are quite distinctive differences between the sets used: in EKS98 there is strong gluon antishadowing, which originates from the requirement of conservation of momentum combined with the constraint obtained from the measured Q^2 dependence of $F_2^{\rm Sn}/F_2^{\rm C}$. In HKM in turn, the gluons are less shadowed at small values of x but no EMC-effect appears at $x \sim 0.3 \ldots 0.7$. Momentum is conserved also in HKM, so the deficit of momentum at $x \lesssim 0.2$ is compensated by a rapid increase of R_g^A at $x \gtrsim 0.2$. Of the four cases studied here, the HIJING parametrization has the strongest gluon shadowing. Since no antishadowing appears, the HIJING parametrization underestimates the momentum sum at Q_0^2 by about 10 %. As discussed in [12], the HPC parametrization underestimates the momentum sum at Q_0^2 by 5 %.

Regarding the ratio $R_{F_2}^A$, all parametrizations are quite similar at $x \gtrsim 0.02$, as they are based on the fits to the DIS data. At $x \lesssim 0.005$, however, due to the lack of DIS data in the perturbative region, some differences arise. Note also that in Fig. 1 the ratio $R_{F_2}^A$ from EKS98 is plotted for isospin symmetrized $A = 117$ (in order to compare with the NMC data [13]) but the HKM-results are for non-isospin symmetric tin, and with the small effects for D included, as obtained directly from the numerical code [10] by HKM. At the region of small values of x, where the focus of the present paper is, the isospin effects can in any case be safely neglected.

Next, we consider the (lowest order) DGLAP evolution of the nPDF and explicitly show the consequences of different assumptions of gluon shadowing. In EKS98 and HKM the scale evolution has already been done but for the other two cases it needs to be performed separately. For the HPC and HIJING nuclear effects, we do this by choosing Q_0^2 as the initial scale, and

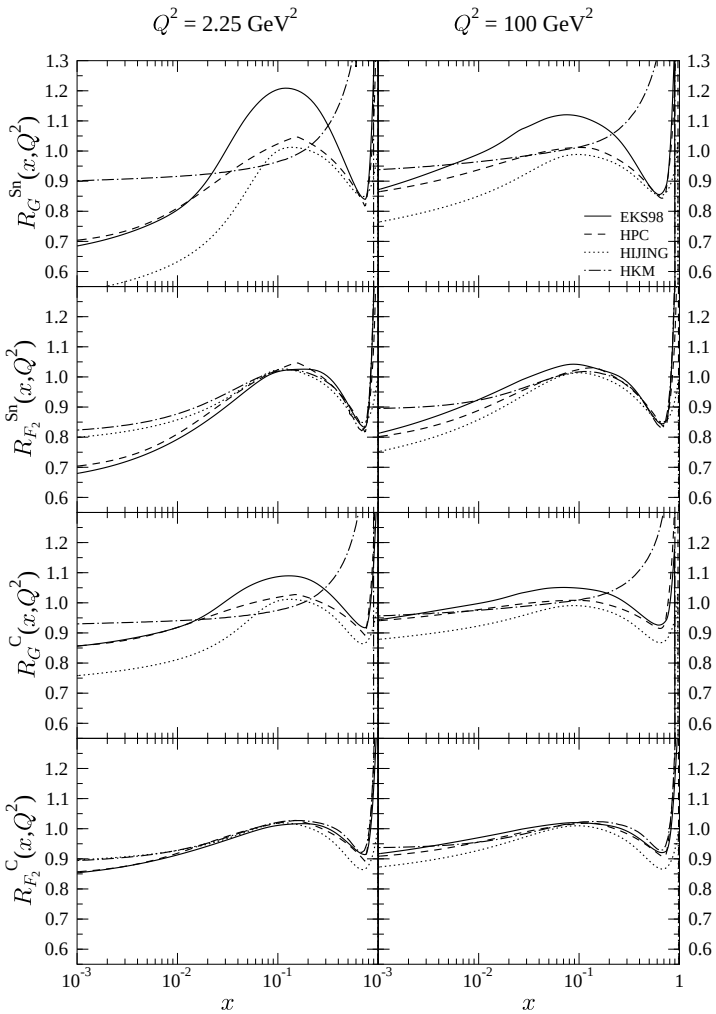

Figure 1: $R_g(x,Q^2)$ and $R_{F_2}(x,Q^2)$ for Sn ($A = 117$) and C ($A = 12$) as function of x for scales $Q^2 = Q_0^2 = 2.25$ GeV^2 and $Q^2 = 100$ GeV^2. The difference of the ratios $R_g^A(2x,Q^2)/R_{F_2}^A(x,Q^2)$ for Sn and C is proportional to the slope $\partial(F_2^{\rm Sn}/F_2^{\rm C})/\partial \log Q^2$, as shown in Eq. (8).

computing the intial distributions with the MRST distributions, taking the initial nuclear effects from Eqs. (3) and (4)-(5), correspondingly. The absolute nPDF are then evolved from Q_0^2 to higher scales with the DGLAP equations. The results for $R_G^A(x,Q^2)$ and $R_{F_2}^A(x,Q^2)$ at a scale $Q^2 = 100$ GeV^2 from all cases studied, are shown in the right panels of Fig. 1. Again, with the EKS98 we use the MRST distributions in plotting the ratio $R_{F_2}^A$, and the HKM results are directly from the HKM code [10].

Two observations can be immediately made: first, the pQCD scale evolution is a sizable effect for the gluon ratios at $x \lesssim 0.3$, especially in the small-x region. Second, at small values of x, due to the very strong gluon shadowing, the HIJING parametrization predicts the slope $\partial R_{F_2}^A(x,Q^2)/\partial \log Q^2$ to be *negative*, contrary to the other three cases studied.

The comparison between the calculated results and the NMC data [13] for the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ is shown in Fig. 2. The EKS98 parametrization reproduces well the experimental data, the

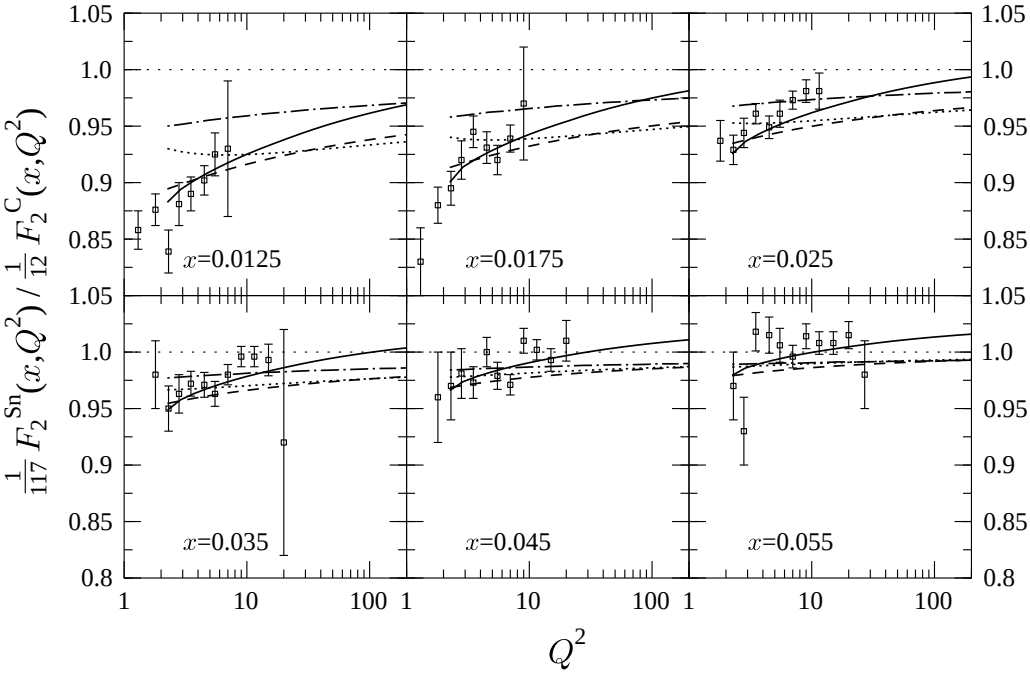

Figure 2: Comparison of the calculated and measured Q^2 dependence of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$. The NMC data [13] are shown with statistical errors only. The results for EKS98 [5, 6] (solid lines) and HKM [9] (dotted-dashed) are from the corresponding global DGLAP analyses. The Q^2 dependence of the HPC (dashed) and HIJING (dotted) cases is obtained from the DGLAP equations by taking the initial conditions for the nuclear effects from Eqs. (3) and (4)-(5).

$\log Q^2$ slopes in particular, which is not surprising as these data have been taken into account in the analysis [5, 6]. The DGLAP-evolved HPC also reproduces the data reasonably well. The HKM results miss the absolute normalization of the data at the smallest values of x but have the right kind of curvature. The data would also suggest faster evolution at small scales than that in HKM. The HIJING parametrization results in a negative Q^2-slope for the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ at small values of x and Q^2. This clearly is in contradiction with the data. This behaviour is caused by the strong gluon shadowing in the HIJING parametrization at small x, as will be discussed next.

At small values of x, the structure function F_2 (of both p and n) is dominated by sea quarks, whose DGLAP evolution in turn is dominated by gluons. The $\log Q^2$ slope of F_2 can be approximated at lowest order as [15]

$$\frac{\partial F_2^{p(n)}(x,Q^2)}{\partial \log Q^2} \approx \frac{10\alpha_s}{27\pi} x g(2x,Q^2). \tag{6}$$

This leads to

$$\frac{\partial R_{F_2}^A(x,Q^2)}{\partial \log Q^2} \approx \frac{10\alpha_s}{27\pi} \frac{x g(2x,Q^2)}{\frac{1}{2}F_2^{\rm D}(x,Q^2)} \Big\{ R_g^A(2x,Q^2) - R_{F_2}^A(x,Q^2) \Big\}, \tag{7}$$

and

$$\frac{\partial(\frac{1}{117}F_2^{\rm Sn}/\frac{1}{12}F_2^{\rm C})}{\partial \log Q^2} \approx \frac{10\alpha_s}{27\pi}\frac{xg(2x,Q^2)}{\frac{1}{2}F_2^{\rm D}(x,Q^2)}\frac{R_{F_2}^{\rm Sn}(x,Q^2)}{R_{F_2}^{\rm C}(x,Q^2)}\left\{\frac{R_g^{\rm Sn}(2x,Q^2)}{R_{F_2}^{\rm Sn}(x,Q^2)}-\frac{R_g^{\rm C}(2x,Q^2)}{R_{F_2}^{\rm C}(x,Q^2)}\right\}, \quad (8)$$

where xg is the gluon distribution in the free proton, and $F_2^{\rm D}=F_2^p+F_2^n$.

In the HIJING parametrization gluons are much more strongly shadowed than quarks, so $R_g^A(2x,Q_0^2)<R_q^A(x,Q_0^2)$. Eq. (7) thus directly shows why a negative $\log Q^2$ slope for the ratio $R_{F_2}^A$ seen in Fig. 1 is bound to follow. Based on Eq. (8) we can also understand the origin for the differences between the computed $\log Q^2$ slopes of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$. The second term in the curly brackets in Eq. (8) is obviously always closer to unity than the first one, since the nuclear effects in smaller nuclei are smaller. On the other hand, the NMC data in Fig. 2 indicates a clearly positive $\log Q^2$-slope. These facts imply that R_g^A is bounded from below as $R_g^A(2x,Q^2)>R_{F_2}^{\rm A}(x,Q^2)$. Consequently, $R_g^A(x,Q_0^2)$ cannot be much smaller than $R_{F_2}^{\rm A}(x,Q_0^2)$. The values of $R_g^A(2x,Q_0^2)$ and $R_{F_2}^{\rm A}(x,Q_0^2)$ at $x=0.0125$ for Sn and C are directly readable off from Fig. 1. Both the sign and the relative order in the magnitude of the computed slopes in the first panel of Fig. 2 can be understood by substituting these values in Eq. (8). The fact that the $\log Q^2$ slopes from the HIJING parametrization of nuclear effects are opposite to the measured ones leads us to the conclusion that very strongly shadowed gluons are ruled out within the leading twist DGLAP framework.

The observation above has an interesting consequence for the HIJING model. As discussed in [25], particle production in AA collisions in HIJING is due to contributions from soft and hard components,

$$\frac{dN_{ch}^{AA}}{d\eta}=\frac{1}{2}\langle N_{\rm parts}^{AA}\rangle\langle n_{\rm soft}\rangle+\langle N_{\rm binary}^{AA}\rangle\langle n_{\rm hard}\rangle\frac{\sigma_{\rm jet}^{AA}(\sqrt{s},p_0)}{\sigma_{\rm in}^{NN}(\sqrt{s})}, \quad (9)$$

where $\langle N_{\rm parts}\rangle\sim 2A$ is the average number of participants in AA for the centrality selection considered, $\langle n_{\rm soft}\rangle=1.6$ is the average multiplicity from soft processes to the rapidity interval considered, $\langle N_{\rm binary}\rangle\sim A^{4/3}$ is the average number of binary nucleon-nucleon collisions, $\langle n_{\rm hard}\rangle=2.2$ represents particle production from (mini)jet hadronization, and $\sigma_{\rm in}^{NN}(\sqrt{s})$ is the inelastic nucleon-nucleon cross section. The integrated minijet cross section $\sigma_{\rm jet}^{AA}(\sqrt{s},p_0)$ is computed perturbatively by using the nPDF. The parameter p_0 for the minimum transverse momentum of minijets defines the division into soft and hard components. The values of p_0 are determined based on fits to the inelastic cross sections measured in pp and $p\bar{p}$ collisions. The $\sqrt{s}$ dependence of p_0 in HIJING has now been updated [25], and p_0 is an increasing function of $\sqrt{s}$. Unlike the saturation scale in models with parton saturation [2, 20, 21, 22, 23], the scale p_0 does not depend on A.

We write $\sigma_{\rm jet}^{AA}(\sqrt{s},p_0)=[R_g^A(\langle x\rangle,\langle Q^2\rangle)]^p\sigma_{\rm jet}(\sqrt{s},p_0)$, where $\langle x\rangle\approx 2p_0/\sqrt{s}$ and $\langle Q^2\rangle=ap_0^2$ with $a\sim 2$ for the minijet production at central rapidities. The effective power $p=1\ldots 2$ describes the net effect of gluon shadowing. From Eq. (9) we then obtain

$$[R_g^A(\frac{2p_0}{\sqrt{s}},ap_0^2)]^p\sigma_{\rm jet}(\sqrt{s},p_0)=\frac{\sigma_{\rm in}^{NN}(\sqrt{s})}{\langle N_{\rm binary}^{AA}\rangle\langle n_{\rm hard}\rangle}\left[\frac{dN_{ch}^{AA}}{d\eta}-\frac{1}{2}\langle N_{\rm parts}^{AA}\rangle\langle n_{\rm soft}\rangle\right]. \quad (10)$$

Requiring that the model reproduces the measured multiplicity in Au+Au collisions at RHIC, the r.h.s. of Eq. (10) is a fixed number for each $\sqrt{s}$ and A, and thus independent of p_0.

As discussed above, within the DGLAP framework the NMC data implies that R_g^A should be larger than that of the HIJING parametrization in Eq. (5) (this is due to both weaker gluon shadowing and due to scale evolution). Correspondingly, $\sigma_{\rm jet}(\sqrt{s}, p_0)$ should be smaller. This in turn implies that p_0 should be larger than in the pp case. We are thus lead to the conclusion that the same multiplicities as are currently obtained from the HIJING model with an A-independent $p_0(\sqrt{s})$ and very strong gluon shadowing, can be obtained with weaker gluon shadowing by introducing a scale p_0 which is an increasing function of both $\sqrt{s}$ and A.

3 Non-linear effects in Q^2 evolution

The DGLAP scale evolution discussed above is linear in the parton densities. At very small values of x the density of gluons increases to the extent that contributions from non-linear phenomena due to gluon fusion may start to play a role. The first non-linear correction terms to pQCD evolution equations have been computed in [2, 3], let us call them "GLRMQ terms". One could argue that perhaps these corrections would change the scale evolution in such a way that stronger gluon shadowing could be allowed. In order to study this possibility, we next add the GLRMQ terms into the DGLAP equations. In connection with the DGLAP analysis of nPDF, the non-linear effects have been numerically studied e.g. in [33] (see also [34]).

The valence quark evolution remains unmodified, but for the gluons and sea quarks the generic form of the equations reads

$$\frac{\partial x g_A}{\partial \log Q^2} = \left(\frac{\partial x g_A}{\partial \log Q^2}\right)_{\rm DGLAP} - \frac{1}{Q^2}\mathcal{R}_{ggg}(x, Q^2) \tag{11}$$

$$\frac{\partial x \bar{q}_A}{\partial \log Q^2} = \left(\frac{\partial x \bar{q}_A}{\partial \log Q^2}\right)_{\rm DGLAP} - \frac{1}{Q^2}\left[\mathcal{R}_{\bar{q}gg}(x, Q^2) - \mathcal{R}^{\rm HT}_{\bar{q}gg}(x, Q^2)\right] \tag{12}$$

$$\frac{\partial x g_{HT}}{\partial \log Q^2} = -\mathcal{R}_{ggg}(x, Q^2) \tag{13}$$

where GLRMQ terms $\mathcal{R}_{ggg} \sim \alpha_s^2 \int_x^1 (dy/y) y^2 g_A^{(2)}(y, Q_0^2)$, $\mathcal{R}_{\bar{q}gg} \sim \alpha_s^2 x^2 g_A^{(2)}(x, Q_0^2)$ and $\mathcal{R}^{\rm HT}_{\bar{q}gg} \sim \alpha_s \int_x^1 (dy/y)(x/y)\bar{\gamma}_{\rm FG}(x/y) y g_{\rm HT}(y, Q^2)$. The detailed form of these terms can be found in [3]. Following [3] and [33], for the 2-gluon density we take $x^2 g_A^{(2)}(x, Q^2) = \frac{A}{\pi R_A^2}[x g_A(x, Q^2)]^2$, where $\frac{A}{\pi R_A^2}$ and the nuclear radius $R_A = 1.12 A^{1/3} - 0.86 A^{-1/3}$ are based on the Woods-Saxon parametrization of nuclear densities. For the gluon higher-twist term, we assume $x g_{HT}(x, Q_0^2) = x^2 g_A^{(2)}(x, Q_0^2)$ and that $g_{HT}(x, Q^2) \geq 0$.

We obtain the initial conditions for the actual nPDF as before, from the EKS98, HIJING and HPC parametrizations and with the MRST distributions for the free proton. We do not make an attempt to include the GLRMQ terms to the HKM analysis. The results for the scale evolution of the ratio $\frac{1}{117}F_2^{\rm Sn}/\frac{1}{12}F_2^{\rm C}$ are shown in Fig. 3 against the NMC data. We observe that the effects of the GLRMQ corrections remain fairly modest (as they should, in order to stay as corrections) but that they make the positive $\log Q^2$-slopes (EKS98, HPC) flatter, and the negative slopes (HIJING) even more negative than in the case of pure DGLAP evolution. Our conclusion therefore is that the GLRMQ terms do not give support to the strong shadowing of gluons, either.

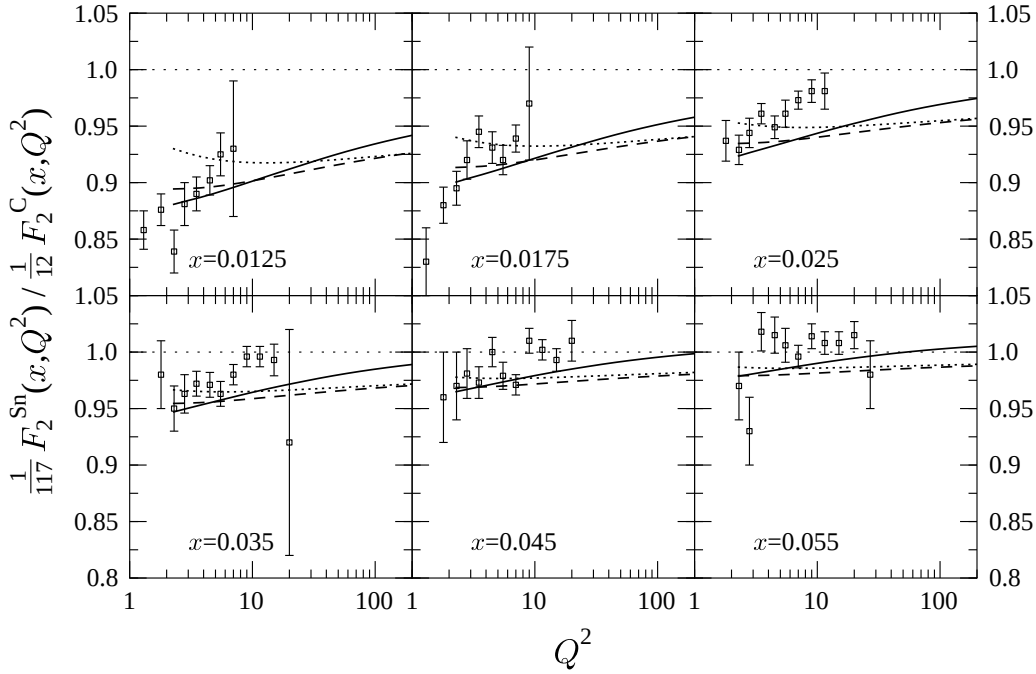

Figure 3: The scale dependence of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ calculated using DGLAP evolution with MQ corrections, and compared to NMC data [13]. Initial conditions for nuclear effects are taken from EKS98 (solid lines), HPC (dashed) and HIJING (dotted) parametrizations.

The systematics of the change in the $\log Q^2$ slopes at the initial scale Q_0^2 can also be easily understood from Eqs. (12). The evolution equation for F_2^A can be written as $F_2^{A\prime} = F_{2\,\rm DGLAP}^{A\prime} + F_{2\,\rm GLRMQ}^{A\prime}$, where the prime stands for $\partial/\partial \log Q^2$, and where the second term contains only the GLRMQ corrections for the sea quark evolution from Eq. (12). The net effect of the GLRMQ corrections to the $\log Q^2$ slope of sea quarks is negative, and dominated by the term $\mathcal{R}_{\bar{q}gg} \sim A^{1/3}(R_g^A)^2(xg)^2$. The $\log Q^2$ slope of the the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ can then be expressed as

$$\frac{\partial(F_2^{\rm Sn}/F_2^{\rm C})}{\partial \log Q^2} = \frac{F_2^{\rm Sn}}{F_2^{\rm C}} \left\{ \left[\frac{F_{2\,\rm DGLAP}^{\rm Sn\,\prime}}{F_2^{\rm Sn}} - \frac{F_{2\,\rm DGLAP}^{\rm C\,\prime}}{F_2^{\rm C}} \right] + \left[\frac{F_{2\,\rm GLRMQ}^{\rm Sn\,\prime}}{F_2^{\rm Sn}} - \frac{F_{2\,\rm GLRMQ}^{\rm C\,\prime}}{F_2^{\rm C}} \right] \right\} \tag{14}$$

where the latter term in brackets again contains only the GLRMQ contributions to the $\log Q^2$ slope from Eq. (12). Using $F_2^A \sim R_{F_2}^A(F_2^p + F_2^n)$, the GLRMQ part becomes

$$\frac{F_{2\,\rm GLRMQ}^{\rm Sn\,\prime}}{F_2^{\rm Sn}} - \frac{F_{2\,\rm GLRMQ}^{\rm C\,\prime}}{F_2^{\rm C}} \sim \frac{12^{1/3}(R_g^{\rm C})^2}{R_{F_2}^{\rm C}} \Big[1 - \left(\frac{117}{12}\right)^{1/3} \frac{R_{F_2}^{\rm C}}{R_{F_2}^{\rm Sn}} \left(\frac{R_g^{\rm Sn}}{R_g^{\rm C}}\right)^2 \Big]. \tag{15}$$

Reading the values for $R_{F_2}^A(x,Q_0^2)$ and $R_g^A(x,Q_0^2)$ at $x \sim 0.01$ off from Fig. 1, it is easy to see that the contribution from the GLRMQ terms to the $\log Q^2$ derivative of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ is indeed negative in all the three cases studied.

We emphasize that for the DGLAP+GLRMQ case we have not attempted to make a global analysis of the nPDF, where the initial conditions would be based on fits to various set of data. Such an analysis will require also detailed studies of the constraints for the magnitude of the

nonlinearities in the case the free proton. We note, however, that according to Fig. 3, in the framework of DGLAP+GLRMQ we could expect somewhat less gluon shadowing for $g_A(x, Q_0^2)$ than in the framework of pure DGLAP. This will also lead to slightly smaller excess of gluons at larger values of x (see EKS98 and HKM in Fig. 1). This analysis is left as a future task.

4 Conclusions

We have studied the Q^2 dependence of the ratio $F_2^{\rm Sn}/F_2^{\rm C}$ in the pQCD framework of lowest-order DGLAP evolution and leading twist. We emphasize that the NMC data on the Q^2 dependence of $F_2^{\rm Sn}/F_2^{\rm C}$ [13] is the most direct (although indirect) measurement of the nuclear gluon distributions that is currently available. These data provide a valuable constraint for pinning down the nuclear gluon shadowing in the global DGLAP analyses of nPDFs.

We have demonstrated the sensitivity of the $\log Q^2$ slopes of $F_2^{\rm Sn}/F_2^{\rm C}$ to the gluon shadowing by comparing four different approaches. The NMC data [13] implies that the nuclear effects in the gluon distributions (R_g^A), should lie close to those in $F_2^A/F_2^{\rm D}$ at $x \sim 0.01$ and $Q^2 \sim 2$ GeV^2. In particular, a very strong gluon shadowing, as suggested e.g. in [25], is ruled out by the NMC data, since it leads to a $\log Q^2$ slope whose sign is opposite to what is measured. Consequences for the HIJING model have been discussed in Sec. 2. We also suggest to use the NMC data [13] and [35] as further constraints in the DGLAP analysis of HKM [9].

The DGLAP approaches [5, 6] and [9] conserve momentum explicitly. The scale independent parametrizations from HIJING [24, 25] and HPC [29] do not do that, especially not if gluons are very strongly shadowed and no antishadowing appears. We emphasize that the strong antishadowing in the gluons of EKS98 results from the fact that the gluon shadowing is constrained by the NMC data at $x \sim 0.01$, and to compensate the loss of momentum there, antishadowing is needed.

We have also studied the effects of the nonlinear GLRMQ terms [2, 3] in the DGLAP equations. These terms decrease the slope $\partial(F_2^{\rm Sn}/F_2^{\rm C})/\partial \log Q^2$. For the very strongly shadowed gluons this results in even more negative slopes than without the GLRMQ corrections. Our conclusion therefore is that within the DGLAP and DGLAP+GLRMQ frameworks studied we cannot find any support from the DIS data for a much stronger gluon shadowing at $x \sim 0.01$ and $Q^2 \sim 2$ GeV^2 than what is observed in the ratio $F_2^A/F_2^{\rm D}$. To resolve the situation at smaller values of x, especially in the region relevant for the few-GeV scales at the LHC, more DIS data in the perturbative region would be needed.

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