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Diffraction at HERA, Color Opacity and Nuclear Shadowing in DIS off nuclei

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Abstract. The QCD factorization theorem for diffractive processes in DIS is used to derive formulae for the leading twist contribution to the nuclear shadowing of parton distributions in the low thickness limit (due to the coherent projectile (photon) interactions with two nucleons). Based on the current analyzes of diffraction at HERA we find that the average strength of the interactions which govern diffraction in the gluon sector at $x \leq 10^{-3}, Q_0 = 2GeV$ is $\sim$ 50mb. This is three times larger than in the quark sector and suggests that applicability of DGLAP approximation requires significantly larger Q_0 in the gluon sector. We use this information on diffraction to estimate the higher order shadowing terms due to the photon interactions with $N \geq 3$ nucleons which are important for the scattering of heavy nuclei and to calculate nuclear shadowing and Q^2 dependence of gluon densities. For the heavy nuclei the amount of the gluon shadowing: $G_A(x,Q_0^2)/AG_N(x,Q_0^2)_{|\, x\leq 10^{-3}} \sim 0.25-0.4$ is sensitive to the probability of the small size configurations within wave function of the gluon "partonometer" at the Q_0 scale. At this scale for $A \sim 200$ the nonperturbative contribution to the gluon density is reduced by a factor of $4-5$ at $x \leq 10^{-3}$ unmasking PQCD physics in the gluon distribution of heavy nuclei. We point out that the shadowing of this magnitude would strongly modify the first stage of the heavy ion collisions at the LHC energies, and also would lead to large color opacity effects in eA collisions at HERA energies. In particular, the leading twist contribution to the cross section of the coherent J/ψ production off $A \geq 12$ nuclei at $\sqrt{s} \geq 70$ GeV is strongly reduced as compared to the naive color transparency expectations. The Gribov black body limit for $F_{2A}(x,Q^2)$ is extended to the case of the gluon distributions in nuclei and shown to be relevant for the HERA kinematics of eA collisions. Properties of the final states are also briefly discussed.

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1 Introduction

It has been realized by Gribov already before the advent of QCD that there exists a deep relation between the phenomenon of high-energy diffraction and the nuclear shadowing phenomenon [1]. In particular, the nuclear shadowing due to the interaction of a virtual photon with two nucleons can be unambiguously calculated in terms of the $\gamma^* + N \to X + N$ diffractive cross section if the coherence length $l_c = \frac{2q_o}{Q^2+M^2}$ is much larger than the nucleus radius, R_A. Here M^2 is the invariant mass squared of the quark-gluon system to which a virtual photon is transformed [1]. In the case of the charged parton structure functions ($F_{2A}(x,Q^2)$) connection between shadowing and diffraction has been explored for a long time, see [2–11] and references therein. The importance of the color fluctuations-weakly interacting configurations in the shadowing phenomenon was first understood in [2] where this effect has been estimated based on the QCD aligned jet model and included in the calculation of $F_{2A}(x,Q^2)$.

Additional contributions to the nuclear structure function are related to the piece of the photon wave function for which the coherence length l_c is of the order of the average internucleon distance $r_{NN} \approx 1.7Fm$. These important nuclear effects have been estimated and explored in [2,5,12] using constraints which follow from the QCD momentum and baryon sum rules. Account for these effects leads to a more complicated QCD evolution which mixes shadowing region and the region of larger x.

In the recent paper [13] we started analysis of the implications of the information which is now available from HERA on the role of the gluon degrees of freedom in the diffractive processes in DIS for the gluon nuclear shadowing. We were able to study shadowing for $x \leq 10^{-3}$ and $Q^2 \sim 20-50\mathrm{GeV}^2$ and predict a factor $\sim 2-3$ larger shad-

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owing for the gluon channel than for the quark channel. This is in line with expectations of [14], though it differs from the pattern assumed in a number of the models, see e.g. for the recent summary [15].

In this paper we will extend this analysis to a broad range of x and Q^2. The main tool we will use is the QCD factorization theorem for the hard diffractive scattering [16], see also [17]. Application of the QCD factorization theorem makes it possible to establish correspondence between the deuteron(nucleus) infinite momentum frame (IMF) and the rest frame descriptions and therefore to explore advantages of both descriptions.

An evident advantage of the IMF description is the simple interpretation of the momentum and the baryon sum rules. On the other hand space-time development of high energy processes and nuclear shadowing phenomenon have a more clear interpretation within the nucleus rest frame approach. Using the QCD factorization theorem [16] and the Gribov analysis of nuclear shadowing we will derive the model independent expressions for the leading twist nuclear shadowing of parton densities in the case of coherent interactions with two nucleons (Sect. 2). [1] However the shadowing due to the interaction with two nucleons cannot diminish the parton density by more than a factor of 0.75 without introducing ghosts into the theory (see discussion in the end of Sect. 3). Hence in Sect. 3 we use the recent analysis [18] of the HERA diffractive data to extract the information on the S-channel dynamics of diffraction which is necessary to calculate the effects of coherent interactions with $N \geq 3$ nucleons. We find out that this analysis implies that the average strength of the interaction responsible for the diffraction in the gluon channel at the resolution scale $Q_0 \sim 2$GeV and $x \leq 10^{-3}$ is very large: $\sigma_{eff} \sim 50-60mb$. Large value of the interaction strength could be related to the large cross section of the small color octet dipole interaction with a nucleon which is given by

$$\sigma^{inel}_{\text{“}color\ octet\ dipole\text{''},N}(E_{inc}) = \frac{3\pi^2}{4} b^2 \alpha_s(Q^2) x G_N \times \left(x, Q^2 \equiv \frac{\lambda}{b^2}\right), \quad (1)$$

where $x = \frac{Q^2}{2m_N E_{inc}}$. This is a factor of 9/4 larger than for the case of “color triplet dipole” [17]. If we take for λ the value we estimated before for the color triplet case: $\lambda(x \approx 10^{-3}) \approx 9$, we find that the cross section is close to the S-channel unitarity limit for the range of applicability of DGLAP approximation corresponding to $\sigma_{inel} \geq \sigma_{el}$ for $x \sim 10^{-4}$ and $Q^2 \sim 10$ GeV^2.

For the scattering of a system with a radius $r \ll r_N$ the slope of the elastic scattering is given by the square of the two-gluon form factor $F^2_{2g}(t) \approx \exp(B_{2g}t)$ with $B_{2g} \sim 4GeV^{-2}$. For this situation condition $\sigma_{inel} = \sigma_{el}$ corresponds to the effective cross section of $8\pi B_{2g} = 40mb$ ($\simeq 35mb$ if the correction due to the real part of the amplitude is taken into account). This value is close to the one which emerges from the analysis of the diffractive data where the size of diffractive system is smaller though not negligible as compared to the nucleon size. Note also that the smallness of the shrinkage of the diffractive cone for the J/ψ elastic photoproduction ($\Delta B \leq 1GeV^{-2}$ for $\sqrt{s}$ between 5 and 200 GeV as compared to $\Delta B \approx 3GeV^{-2}$ expected in the soft regime) indicates that perturbative physics occupies most of the rapidity range for $Q^2 \geq 4GeV^2$ for HERA energy range.

It is worth emphasizing that inapplicability of the DGLAP evolution equation and possible closeness to the unitarity limit we discuss here are due to the growth of $xG_N(x, Q^2)$ generated predominantly by the $\log Q^2$ terms in the DGLAP evolution equations rather than solely by the $\ln(1/x)$ terms which would be the BFKL approximation. Hence the pattern discussed here is qualitatively different from the BFKL scenario of high-energy dynamics. Indeed, large values of $xG_N(x, Q^2)$ ($\sim 10-20$ in the small x HERA kinematics and growing with Q as $\sim \sqrt{Q}$) emerge not because of long ladders in rapidity - the ladders contain no more than 2-3 gluons in the multi-Regge kinematics, but rather due to a large number of emitters at the lower resolution scale. Possible closeness to the unitarity limit makes it likely that for moderate $Q^2 \leq 10GeV^2$ corrections to the DGLAP predictions for the nuclear shadowing would be rather large. This would primarily affect our predictions for moderate Q^2 since the information about the gluon induced diffraction is obtained predominantly at larger Q^2 and extrapolated to lower Q^2 via the DGLAP equations.

In Sect. 4 we first analyze the dynamics of the fluctuations of the interaction strength (color coherence - color opacity and color transparency phenomena) and explain that significant fluctuations of the strength of interaction should be present in particular due to the QCD evolution. Next we study the nuclear shadowing originating from the interactions with $N \geq 3$ nucleons. We point out that the eikonal type approximation seems reasonable in the soft QCD regime when the projectile wave function contains a large number of constituents. On the contrary in the PQCD regime where photon wave function is given by a $q\bar{q}$ dipole not more than two inelastic collisions are allowed by energy conservation law. Otherwise the energy released in the inelastic collisions calculated through the cuts of exchanged parton ladders will be larger than the sum of the energies of the colliding particles. Evaluation of a larger number of rescatterings in PQCD is beyond the scope of the naive semiclassical approximation and requires an accurate account of the space-time evolution of the scattering process, in particular a calculation of the NLO approximation to the photon wave function. We demonstrate that the $N \geq 3$ interactions are sensitive to the existence of the fluctuations of the interaction strength. The sensitivity is rather small for $A \sim 12$. For such A we predict significantly larger shadowing for gluons: $G_A(x, Q_0^2)/AG_N(x, Q_0^2)_{|x\leq 10^{-3}, Q_0^2=4GeV^2} \sim 0.7$ than for quarks: $F_{2A}(x, Q_0^2)/AF_{2N}(x, Q_0^2)_{|x\leq 10^{-3}, Q_0^2=4GeV^2} \sim$

[1] For an early discussion of the general arguments for the presence of the nuclear shadowing in the leading twist and references see [2].

0.85. For larger A sensitivity to fluctuations steadily increases. However we find that the average interaction strength in the gluon channel is large at the normalization scale of $Q_0 = 2$ GeV so a significant nuclear shadowing of average and larger than average interaction strengths into the cross section is determined by the geometry of collisions and rather insensitive to the structure of the distribution over the strengths. As a result of shadowing of strongly interacting (nonperturbative ?) configurations, the relative contribution of the interactions with small σ is strongly enhanced in the parton distributions in heavy nuclei. We estimate possible effects of the weakly interacting configurations and find that they may contribute up to 50 % to $G_{A\sim 200}(x \leq 10^{-3}, Q_0 \sim 2GeV)$. [2] At the same time the fraction of the cross section due to weakly interacting configurations should diminish with decrease of x.

We want to stress here that the use of information on the diffraction in DIS at HERA allows us to take into account the nonperturbative effects in the gluon nuclear parton densities at the boundary of the QCD evolution. In the previous studies the gluon shadowing either was treated purely perturbatively as for example in the IMF model of McLerran and Venugopalan [19] or it was introduced in a phenomenological way assuming similarity between the shadowing in the gluon and quark channels, see e.g. [5, 15]. Overall a currently popular scenario which is used in the discussion of the heavy ion collisions assumes that reduction of gluon densities is a relatively small correction, for the recent review and references see [20].

Next, we introduce the constraints on the gluon densities which follow from the momentum sum rule and imply presence of the gluon enhancement at $x \sim 0.1$. Combing this effect with the quark and gluon shadowing for small x we calculate the x, Q^2 dependence of the leading twist nuclear densities. In the end of the section we also consider nuclear structure functions in the limit when the nucleus thickness is large enough so the black disk approximation is applicable.

Obviously, the predicted large gluon shadowing has many implications for the various high-energy processes of scattering off nuclei. In Sect. 5 we calculate the impact parameter dependence of the gluon shadowing and briefly analyze two phenomena: the emergence of the color opacity in the coherent production of J/ψ and Υ-mesons from nuclei in the HERA kinematics, and the suppression on the minijet production in AA collisions at the LHC energies. We find both the color opacity effect and minijet suppression to be very large. For example, for the lead-lead collisions we predict a suppression of the minijet production at $p_t = 2(3)$ GeV/c by a factor $\geq 7 (\geq 4)$.

In Sect. 6 we briefly discuss properties of final states and predict a dip in the ratio of the spectra of leading hadrons produced in the current fragmentation region in eA and in eN collisions for rapidities shifted from the maximum rapidity by $\ln \left[\left\langle M^2_{diff}\right\rangle/\mu^2\right]$ where $\left\langle M^2_{diff}\right\rangle$ is the average diffractive mass^2 produced in eN scattering.

In Sect. 7 we compare our approach with some of the recent studies of the nuclear shadowing.

[2] We are indebted to A.Mueller who stressed the effect of filtering of the PQCD physics in the parton distributions in nuclei.

2 The QCD factorization theorem and the leading twist shadowing for the parton densities

The studies of the diffraction production in hard processes lead to the introduction of the diffractive parton densities $f^D_{j/B}(\beta, Q^2, x_{I\!P}, t)$ with $\beta = \frac{x}{x_{I\!P}}$, which represent the number densities of partons in the initial hadron, but conditional on the detection of the diffracted outgoing hadron B in the target fragmentation region with light-cone fraction $1 - x_{I\!P}$ and fixed momentum transfer t. For example in the case of the diffractive process $e+p \rightarrow e+p+X$ the diffractive structure function F^D_2 which is introduced via

$$\frac{d^4\sigma_{diff}}{d\beta dQ^2 dx_{I\!P} dt} = \frac{2\pi\alpha^2}{\beta Q^4}\left(\left[1+(1-y)^2\right]F^D_2 - y^2 F^D_L\right), \quad (2)$$

can be written as

$$F^D_2(\beta, Q^2, x_{I\!P}, t) = \sum_a e^2_a \beta f_{a/p}(\beta, Q^2, x_{I\!P}, t) + HT\ corrections. \quad (3)$$

In the case of the proton production this structure of the hard diffractive processes was first suggested in the framework of the Ingelman-Schlein model [21]. Recently it was demonstrated [16] that the QCD factorization theorem is valid for the x, Q^2 evolution of these parton densities at fixed $x_{I\!P}, t$. The evolution is governed by the same DGLAP equations as for the inclusive processes. The HERA data on diffraction in DIS are consistent with the dominance of the leading twist contribution except near the edge of the phase space (see discussion below).

For the processes dominated by the vacuum channel the Gribov theory [1] unambiguously relates diffractive processes in the scattering of a projectile off a single nucleon to the process of nuclear shadowing due to the interaction of the projectile with two nucleons. The simplest way to visualize this connection for example in the case of the scattering off the deuteron is to consider $\gamma^* d$ scattering in the deuteron rest frame in the kinematics where $l_c \gg R_d$ (R_d is the radius of the deuteron). Due to the difference of the spatial scales characterizing the deuteron and the soft QCD strong interactions, the dominant contribution is given by the diagrams where the photon dissociates into a hadron component before deuteron and then this component interacts with both nucleons. Let us use the AGK theorem [22] and consider the cut of the double scattering diagram corresponding to the diffractive final state (Fig. 1). This corresponds to the scattering off one nucleon in the $|in\rangle$ state and off the second nucleon in the $\langle out|$ state. The final state interaction between nucleons is accounted for as usual within the closure approximation. The interference between two diagrams results

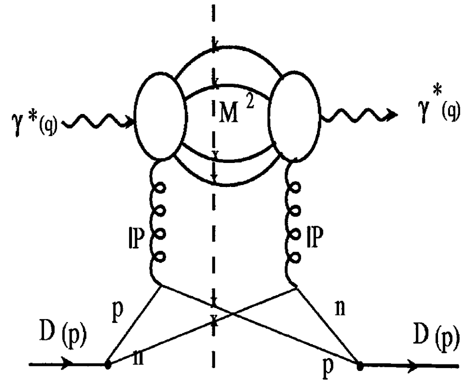

Fig. 1. Double scattering diagram for the $\gamma^* D$ scattering corresponding to a diffractive final state

from the Fermi motion of nucleons in the deuteron since the spectator nucleon in the $|in\rangle$ state has to have a momentum equal to the momentum of the diffracted nucleon in the $\langle out|$ state. The screening effect is expressed ultimately through $-Ref^2$ where f is the diffractive amplitude of the interaction of the probe with the nucleon as compared to $|f|^2$ in the case of diffractive scattering off the nucleon. The real part of the diffractive amplitude is rather small and can be calculated from the information about the imaginary part of the amplitude. Thus the difference between $|f|^2$ and $-Ref^2$ is small and easy to deal with.

Hence we can apply the Gribov results for the scattering off the deuteron and nuclei to evaluate the shadowing contribution to the deuteron parton density of flavor j in terms of the corresponding nucleon diffractive densities (we consider only the Pomeron type contribution, so we do not distinguish diffraction of protons and neutrons)

$$f_{j/^2H}(x,Q^2) = f_{j/p}(x,Q^2) + f_{j/n}(x,Q^2) - \eta\frac{1}{4\pi}\int dx_{I\!P} dt S(4t) f^D_{j/N}\left(\beta, Q^2, x_{I\!P}, t\right). \quad (4)$$

Here $S(t)$ is the electromagnetic form factor of the deuteron, and $-t = (k_t^2 + (x_{I\!P} m_N)^2)/(1 - x_{I\!P})$, and $\eta = (1 - (ReA_{dif}/ImA_{dif})^2)/1 + (ReA_{dif}/ImA_{dif})^2)$.

Similarly, in the approximation when only scattering off two nucleons in the nucleus is taken into account one can similarly deduce the expression for the shadowing term in terms of the parton densities

$$\begin{aligned} &f_{j/A}(x,Q^2)/A \\ &= f_{j/N}(x,Q^2) - \frac{1}{2}\int d^2b \int_{-\infty}^{\infty} dz_1 \int_{z_1}^{\infty} dz_2 \int_x^{x_0} dx_{I\!P} \cdot \\ &\cdot f^D_{j/N}\left(\beta, Q^2, x_{I\!P}, t\right)_{|k_t^2=0}\, \rho_A(b,z_1)\rho_A(b,z_2) \\ &\cdot \cos(x_{I\!P} m_N(z_1 - z_2)). \quad (5)\end{aligned}$$

Here $\rho_A(r)$ is the nucleon density in the nucleus normalized according to the equation $\int \rho_A(r) d^3r = A$. For simplicity we gave the expression for the limit when the slope of the dependence of diffractive amplitude on the momentum transferred to target nucleon, t, is much smaller than the one due to the nucleus form factor so that impact parameters of two nucleons are equal. Note that (5) is similar to the corresponding expression for the shadowing in the vector dominance model, see (5.4) in [23] since the space-time evolution of the interaction is the same in both cases. This leads to the same structure of the nuclear block, provided one substitutes the VDM expression for the longitudinal momentum transfer, $q_z = M_V^2/2\nu$ by the Bjorken limit value: $q_z = x_{I\!P} m_N$.

The crucial feature of (4,5) is that the parton densities which enter in the shadowing term evolve according to the leading twist evolution equations. When they are folded with a function of $x_{I\!P}$ which does not depend on Q^2 they retain this property. Since the QCD evolution of real and imaginary parts of hard amplitude is governed by the same evolution equation at sufficiently small x we investigate in the paper the fact that real parts enter into diffraction and into shadowing in a different way does not influence the QCD evolution. This proves that (4,5) correspond to the leading twist contribution to the nuclear parton densities. *In the limit of the low nuclear densities (4,5) provide a complete description of the leading twist nuclear shadowing.*

Obviously the derived equations could not provide a complete picture of the deviations of nuclear parton densities from the sum of the nucleon densities for all x. This is because the derived equations take into account the contributions related to the distances $l_c \gg R_A$ but not the ones related to the configurations with much smaller coherence lengths. The simplest way to estimate the corresponding additional piece is to apply the energy-momentum and baryon sum rules which are exact in QCD for the leading twist parton densities. Therefore to satisfy these sum rules the shadowing should be accompanied by an enhancement of some parton densities at higher x. This enhancement term has to be added to (4,5). If we introduce this term at a scale Q_0^2 for $x \geq x_0$ it would contribute for large Q^2 for much smaller x. Hence the Gribov type approximation becomes inapplicable for fixed x and $Q^2 \to \infty$ [3]. Below, to deal with the enhancement effects we will adopt the procedure of [5] in which these effects are estimated at a low normalization point and the subsequent evolution is dealt with by solving the DGLAP evolution equations.

The range of the validity of approximation where interactions with $N \geq 3$ nucleons are neglected strongly depends on the strength of the corresponding diffraction channel. Hence in the next two sections we review the results of the recent analysis of the HERA diffractive data and build approximation for treating interactions with several nucleons.

[3] In principle one should also take into account the effects of nonnucleonic degrees of freedom (the large x EMC effect) but for any practical purposes this effect is negligible.

3 Diffraction at HERA and shadowing in the low nuclear thickness limit

3.1 Gap probability for the gluon induced hard diffraction

First, let us briefly summarize the results of the studies of diffraction in DIS which were performed over the last few years at electron-proton and proton-antiproton colliders and recast them in the form necessary for the studies of the nuclear shadowing phenomena. The data obtained at HERA include studies of the diffractive structure functions in $\gamma^* + p \to X + p$ scattering, production of dijets and charm in $\gamma^* + p$ scattering. The leading twist contribution appears to describe the data well except very close to the edge of the phase space where higher twist effects are important. So the factorization theorem for these processes [16] seems to hold for the studied Q^2 range, see [18] for the recent analysis.

In the practical applications an assumption is usually made that the semiinclusive parton densities at small values of $x_{I\!P}$ can be written as a product of a function of $x_{I\!P}$ and a parton density which depends on $\beta = x/x_{I\!P}$ and Q^2.

Hence for the sake of brevity we will refer to these densities as the parton densities in the "Pomeron". [4] As a result one can define a diffractive parton density at given x as a convolution of the so called Pomeron flux factor, $f_{p/I\!P}(x_{I\!P})$ and corresponding Pomeron parton density for example for gluons:

$$xg_{dif}(x,Q^2) = \int_x^{x_{max}} f_{p/I\!P}(x_{I\!P}) x/x_{I\!P} \, f_{g/I\!P}(x/x_{I\!P}, Q^2) dx_{I\!P}, \qquad (6)$$

where x_{max} is the maximal value of $x_{I\!P}$ for which diffractive picture still holds.

The important finding of the HERA diffractive studies is that $f_{g/I\!P}(\beta) \gg f_{q/I\!P}(\beta)$ for a wide range of β (Similar trend was observed in $p\bar{p}$ collisions, see review in [26]). For example, in the best global fit of the HERA diffractive data [18] (fit **D**):

$$\beta f_{g/I\!P}(x, Q_0^2) = (9.7 \pm 1.7)\beta(1-\beta),$$
$$\frac{\Sigma_q \beta f_{q/I\!P}(x,Q_0^2)}{\beta f_{g/I\!P}(x,Q_0^2)} \approx 0.13. \qquad (7)$$

Let us consider probability of diffractive events where the proton remains intact for the hard leading twist processes coupled solely to the gluons. It can be defined as

$$P^g_{dif}(x,Q^2) = \frac{xg_{dif}(x,Q^2)}{xg_N(x,Q^2)}. \qquad (8)$$

Since this definition includes only the leading twist contribution into diffraction it effectively excludes the contribution of small masses to the diffraction which could originate from the higher twist effects. Since at small $x \sim 10^{-3}$ and $Q^2 \sim$ few GeV^2 the ratio of the quark and gluon densities in a nucleon is $\sim \frac{1}{2}$, and $f_{g/I\!P}(\beta) \gg f_{q/I\!P}(\beta)$ one obviously expects

$$P^g_{dif}(x,Q^2) = \frac{g_{dif}}{q_{dif}} \frac{q(x,Q_0^2)}{g(x,Q_0^2)} P^q_{dif}(x,Q^2) \gg P^q_{dif}(x,Q^2). \qquad (9)$$

We can quantitatively estimate $P^g_{dif}(x,Q^2)$ using the fit **D** of [18]. The analysis of [18] chooses the initial conditions for the DGLAP evolution at $Q_0^2 = 4GeV^2$ - see (7). We also take $x_{max} = 0.02$ which is the highest $x_{I\!P}$ for which the one has enough sensitivity to the gluon density in the "Pomeron". However for $x \ll x_{max}$ the diffraction probability practically does not depend on the choice of x_{max}. We find that

$$P^g_{dif}(10^{-4} \leq x \leq 3 \cdot 10^{-3}, Q_0^2) \approx 0.34 \cdot (1 \pm 0.15), \qquad (10)$$

which is much larger than $P^q_{gap}(x,Q_0^2) \sim 0.12$. Total probability of rapidity gap which includes double diffractive events (proton dissociation) is larger by a factor ~ 1.4. This factor can be estimated assuming the Regge factorization for $t = 0$: $\frac{d\sigma(\gamma^*+p\to X_1+X_{rec}/dt}{d\sigma(\gamma^*+p\to X_1+p)/dt} |_{t=0} \sim 0.2$ independent of the diffraction state X_1, and taking into account that the slope of the t dependence in the double dissociation should be smaller by about a factor of two due to almost complete disappearance of the proton form factor in the proton vertex. However the cross section given by the HERA groups includes a small contribution of the proton dissociation of about 15% [27]. So effectively the scaling factor is smaller ~ 1.25.

Thus in the gluon channel the ratio of total diffraction to total cross section reaches the value ~ 0.4 for $Q = 2$ GeV. Thus we conclude the ratio of single diffraction to total cross section in the gluon channel is close to that for pp collision (for the soft hadronic processes analogous quantity is the ratio of sum of the elastic and the single diffraction cross sections to the total cross sections).

The QCD evolution leads to a decrease of this probability since at larger Q^2 many small x partons originate

[4] Note that in difference from the usual parton densities which are process independent the "Pomeron" parton densities may depend on the target, on the mass of diffractively produced system etc. In particular, for small masses $M^2 \ll Q^2$ contributions of the higher twist to diffraction become important which are proportional to $xG_N(x,Q^2)^2$ and hence lead to intercept of the effective "Pomeron" $\alpha_{I\!P}(0) \geq 1.20$ [24]. In the analysis of [18] this kinematics was excluded from the fit. At the same time for $M^2 \gg Q^2$ intercept should be more close to $\alpha_{I\!P}(0) = 1.08$ familiar from the soft QCD interactions. These are particular illustrations of the deep difference between the QCD factorization theorem and the Regge pole factorization [25]. Note also that the energy dependence of diffraction is different from that for soft hadronic processes. This is not surprising since the coherence length for the soft hadronic processes is significantly larger than that for soft hadronic processes in DIS.

from "ancestors" at Q_0^2 with $x \geq x_{max}$ which cannot produce protons with small $x_{I\!P} \leq x_{max}$. So at $Q^2 = 25GeV^2$, P^g_{dif} drops to about 0.2. [5]

3.2 Implications for the S-channel picture of hard diffraction

The large probability of diffraction in the gluon channel comparable to that in soft hadron interactions indicates that configurations involved have large interaction cross section (here we effectively switch to the S-channel language of description of diffraction [28,29]). We can quantify this by using the optical theorem $\frac{d\sigma_{dif}}{dt}\mid_{t=0} = \frac{\sigma^2}{16\pi}$ to introduce the strength of interaction σ_{eff} as [6], [7]

$$\sigma_{eff}(x,Q^2) \equiv \frac{16\pi d\sigma^{dif}/dt_{|t=0}}{\sigma} = P^g_{dif}(x,Q^2)16\pi B \quad (11)$$

The effective cross section $\sigma_{eff}(x,Q^2)$ characterizes within the Gribov theory the diffractive rescatterings of the produced quark-gluon system, cf. (12). The results of the calculation are presented in Fig. 2 for $Q = 2$GeV and $x_{max} = 0.02$ for the quark and gluon channels and show that σ_{eff} for the gluon channel is of the order 55 mb for small x and about 3 times larger than for the quark channel. [8] Large value of σ_{eff} can be interpreted as an indication that the interactions in the gluon channel is remaining strong up to much larger virtualities than in the quark channel. This matches rather naturally with the perturbative QCD pattern of a factor of 9/4 larger cross sections for color octet dipole-nucleon interaction than for the color triplet dipole-nucleon interaction [14], [17], see discussion in the introduction after (1).

[5] Note that experimental studies of the "Pomeron" gluon densities are performed either at large virtualities of $Q \geq 5GeV$ or via scaling violation of f^q_{dif} for $Q \geq 2GeV$. So they cannot directly measure the large value of P^g_{dif}.

[6] Here and below we neglect $\leq 5\%$ corrections due to the real part of the amplitude since other uncertainties in the input are of the order $15-20\%$.

[7] For the sake of simplicity we parameterize the t dependence of diffractive cross section as $d\sigma^{dif}/dt = d\sigma^{dif}/dt_{|t=0}\exp Bt$.

[8] Determination of σ_{eff} requires the knowledge of the t-dependence of the diffraction. Experimentally it was measured for the process $\gamma^* + p \to X + p$ only [30] and for relatively large $\langle x_{I\!P}\rangle \sim 0.01$. For this kinematics the fit of [18] describes the t-dependence well. The fit also assumes the rate of the diffractive cone shrinkage $\propto \exp(2\alpha' \ln(1/x_{I\!P}))$ with α' from the soft processes. Due to a larger value of diffraction and hence a larger value of σ_{eff} for the gluon channel the assumption of [18] that the slope for single diffraction in the gluon channel is at least as large as for the quark channel seems also very natural. One should however remember that the lack of direct measurements of the t-dependence of the gluon induced diffraction introduces an additional uncertainty in the results of calculations. Overall our guess for the uncertainty in the value of the parameter σ_{eff} for $10^{-4} \leq x \leq 10^{-3}$ is about 20%.

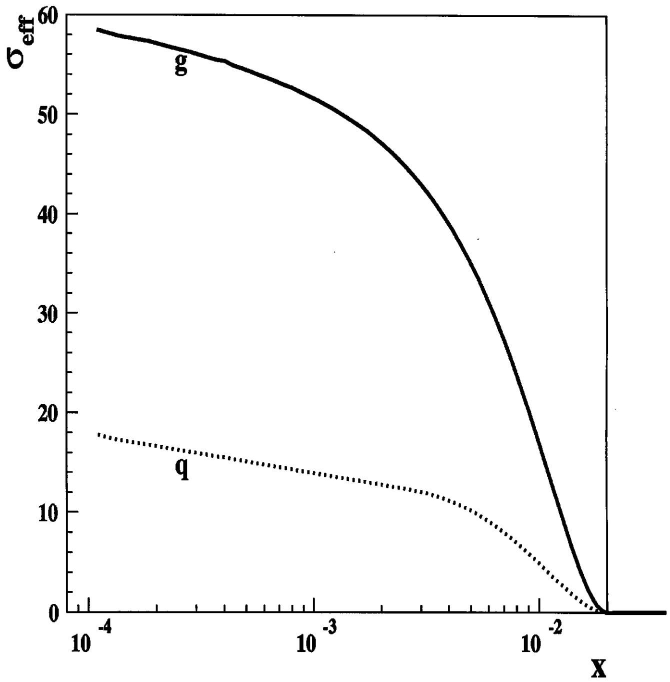

Fig. 2. Dependence of σ_{eff} as defined in (11) for gluon and quark channels on x for $Q = Q_0 = 2$ GeV

3.3 σ_{eff} and shadowing for small nuclear densities

Now we are in a position to rewrite the results of Sect. 2 for the parton shadowing in the limit of small nuclear densities using the notion of σ_{eff}. This would allow us as a next step to go beyond the two nucleon approximation for the shadowing effects. For $x \ll x_{max}$, and not too large A such that $l_c \gg R_A$ the $\cos(x_{I\!P}M_N(z_1 - z_2))$ factor in (5) can be substituted by one. Hence in this limit the amount of shadowing is directly proportional to the differential probability of diffraction at $t = 0$:

$$xG_A(x,Q^2) = AG_N(x,Q^2) - 4\pi \int d^2bT_A^2(b) \cdot \int_x^{x_{max}} \frac{df(x_{I\!P})}{dt}\mid_{t=0} f_g(x/x_{I\!P},Q^2)dx_{I\!P}, \quad (12)$$

where

$$T_A(b) = \int_{-\infty}^{\infty} dz\rho_A(z,b), \quad (13)$$

is the usual nuclear thickness function and $\rho_A(r)$ is the nuclear density. For larger $x_{I\!P}$ one has to take into account the damping factor due to the $\cos(x_{I\!P}M_N(z_1 - z_2))$ factor which originates from the nuclear form factor due to the longitudinal momentum transfer in the transition to the diffractive mass equal to $m_N x_{I\!P}$. However with our choice of $x_{max} = 0.02$ this effect is small even for $A \sim 200$. Equation (12) leads to the shadowing proportional to σ_{eff}:

$$1 - xG_A(x,Q^2)/AxG_N(x,Q^2) = \frac{\sigma_{eff}}{4A}\int d^2bT_A^2(b) \quad (14)$$

and hence predicts a factor ~ 3 larger shadowing for the gluon channel than for the quark channel for $Q = 2GeV$. This is in line with expectations of [14], though it differs from the pattern assumed in a number of the models, see e.g. for the recent summary [15].

Note also that if one does not separate leading and higher twist contributions in the diffraction off a nucleon one can still use the experimental data about the total cross section of diffraction off a nucleon to calculate using (12) the total amount of shadowing in the corresponding channel for the scattering off the deuteron (and nuclei in the approximation when the interactions with $N \geq 3$ nucleons are neglected). The simplest way to see this is to apply the the AGK cutting rules [22] which are valid for the scattering off nuclei. In particular, if the the higher twist effects due to interactions with two nucleons described by the Mueller and Qiu model [31] were important in the nuclear shadowing at the normalization point they should be manifested as well in the diffraction off a nucleon. So, as far as the diffraction is described by the leading twist factorization approximation, (12) leads to the DGLAP evolution of the nuclear shadowing.

Note also that in the approximation when a probe (photon) may interact not more than with two nucleons there exists a relation between the shadowing for the total cross section and the partial cross section of inelastic processes with the multiplicity similar to the one in the inelastic ep scattering - σ_1:

$$\sigma_{tot} = \sigma_{imp} - \sigma_{double}, \quad \sigma_1 = \sigma_{imp} - 4\sigma_{double}, \qquad (15)$$

where σ_{imp} is the impulse approximation cross section and σ_{double} is the screening cross section due to the interaction with two nucleons [22]. One can see from (15) that shadowing due to the photon interaction with two nucleons can diminish the total cross section by not more than a factor of 0.75 as compared to the impulse approximation without introducing ghosts into the theory: for $A_{eff}/A \leq 0.75$ the partial cross section σ_1 would become negative [32].[9] This implies that for $A_{eff}/A \leq 0.75$ shadowing interactions with a larger number of nucleons could not be ignored.

4 Fluctuations of the interaction strength and multinucleon shadowing

4.1 Modeling effects of cross section fluctuations

Due to the large value of $\sigma_{eff}(x, Q^2)$ for the gluon channel, deviations from (12) due to interactions with $N \geq 3$ nucleons become large already for $A \sim 10$. To account for these effects we address the Q^2 dependence of $\sigma_{eff}(x, Q^2)$. Within the DGLAP approximation it basically reflects an influx to small x of configurations which at a lower resolution Q' correspond to configurations with larger $x \equiv x_{parent}$ and hence with smaller $\sigma_{eff}(x_{parent}, Q')$. Configurations which interacted strongly at $Q \sim Q_0$ interact strongly at large Q as well, but they contribute smaller and smaller fraction of the total cross section relevant for the nuclear shadowing phenomenon at a fixed x. This pattern is the same as in the QCD aligned jet model [2]. Since the gluon shadowing strongly reduces gluon densities already $Q \sim Q_0$ the deviations from the DGLAP equations for $Q \sim Q_0$ due higher order order terms in nuclear parton density originating from the average masses of the diffractively produced system should be significantly smaller than in the model of [31] where shadowing at the starting scale is neglected. One can speculate of course that these effects have already occurred between $Q^2 \sim 1GeV^2$ and $Q^2 = Q_0^2$. Higher twist effects are enhanced for the contribution of diffractively produced system with the masses $M^2 \ll Q_0^2$. This interesting question is beyond of the scope of this paper.

It is straightforward to take into account effects of the longitudinal momentum transfer in the diffraction [4,13]. However as we demonstrated in [13] that these effects are important only for $x_{I\!P} \geq 0.03$ for $A \sim 200$ and even for larger $x_{I\!P} \sim 0.05$ for light nuclei. Since we have chosen $x_{max} = 0.02$ we can safely neglect this effect in the following discussion. In this approximation to account for the fluctuation effects it is convenient to introduce the probability distribution over the strength of interaction in the gluon channel - $P_g(\sigma)$. σ_{eff} is expressed in terms of $P_g(\sigma)$ as [34]

$$\sigma_{eff} = \int d\sigma \sigma^2 P_g(\sigma) / \int d\sigma \sigma P_g(\sigma). \qquad (16)$$

We obtain in the generalized eikonal approximation:

$$\frac{G_A(x, Q_0^2)}{G_N(x, Q_0^2)} = \frac{\int d^2b d\sigma P_g(\sigma)(2 - 2\exp(-T(b)\sigma/2))}{A \int d\sigma \sigma P_g(\sigma)}, \qquad (17)$$

where $T(b)$ was defined in (13). Fluctuations lead to a decrease of shadowing effect as compared to the quasieikonal approximation where $P(\sigma) \propto \delta(\sigma - \sigma_{eff})$[10]. We will study effects of fluctuations at length elsewhere. However for characteristic $\sigma \sim \sigma_{eff} \sim 55$ mb the exponential factor in the numerator of (17) is very small for $A \sim 200$ and small enough b a wide range of σ, leading to cross section $\approx 2\pi R_A^2$. This suppresses the contribution of large σ -nonperturbative QCD physics and therefore enhances the contribution of small σ-PQCD physics.

As a result of large absorption for $\sigma \sim \sigma_{eff}$ fluctuations near the **average** value of σ practically do not change the shadowing (for $A \leq 250$) provided σ_{eff} is kept fixed. If, for example, we assume that $P(\sigma) = a\theta(\sigma - \sigma_0)$ the G_A/G_N would change even for heavy nuclei by less than 20%.

Besides, as we mentioned in the introduction the diffraction in the gluon induced hard processes without proton break up constitutes about 35% of the events at $x \leq 10^{-3}$. Large contribution of small σ in $\int d\sigma \sigma P_g(\sigma)$

[9] Note that in [33] where shadowing was calculated in the Mueller and Qiu model the values of A_{eff}/A as low as 0.5 were obtained.

[10] Note that since the value of σ_{eff} is fixed by the cross section of the diffractive scattering this approximation differs from the eikonal approximation often used for the hadron-hadron scattering in which only elastic rescatterings are included.

would imply that for scattering in a significant fraction of configurations the gap probability is much smaller than average. Hence some configurations would have to generate the gap events with a probability exceeding 35%. This is one of indications for the problems for the applicability of DGLAP to describe ep scattering at HERA at low normalization point Q_0^2. At the same time there is a natural mechanism in PQCD for the generation of a contribution of very small σ in $\int d\sigma \sigma P_g(\sigma)$. It comes from the QCD evolution. The partons at given x which originated from ancestors at lower resolution scale with $x_{initial} \geq x_{max}$ do no contribute to diffraction and hence effectively correspond to $\sigma \sim 0$ (In other words this effect reflects the contribution of small coherence lengths to the small x physics due to the QCD evolution). This effect is especially pronounced in the models of QCD evolution with low normalization point like the GRV model [35]. To estimate the sensitivity to this effect we will use two models. The first one is the quasieikonal which neglects higher order fluctuations. The second model is the fluctuation two-component model (to which we will refer to as a fluctuation model) which implements an extreme assumption that a fraction λ of the total cross section at given x originates from configurations with small cross section and the rest from the average ones. The second model is similar to the QCD aligned jet model of [2,4]. It corresponds to

$$\sigma P_g(\sigma) \propto \lambda \delta(\sigma) + (1-\lambda)\delta(\sigma - \sigma_0), \qquad (18)$$

where σ_0 is fixed by (16) to

$$\sigma_0 = \frac{\sigma_{eff}}{1-\lambda}. \qquad (19)$$

The requirement that the gap probability for this component (including dissociation of the nucleon) does not exceed 50% puts an upper limit on λ. Taking into account uncertainties in the value of the total rapidity gap probability we estimate that $\lambda \leq 0.2$. So we will use $\lambda = 0.2$ in this model. Note that for small values of $\sigma_{eff} \leq 20mb$ and moderate values of λ the screening weakly depends on λ. Hence we will ignore this effect for the quark channel. For simplicity we will also assume that λ does not depend on x. This should be considered as a rather rough approximation since for small enough $x \leq 0.005$ one may expect the contribution of small l_c to decrease with decrease of x. Also for larger x the drop of the parameter σ_{eff} can be due to an increase of λ. However in this region we anyway have significant uncertainties due to the contribution of $x_{I\!P} \sim x_{max}$. Thus the results in this x range can be considered simply as a smooth interpolation between the region of large shadowing and the region where shadowing disappears.

Obviously, small σ's in (17) would give the dominant contribution for $A \to \infty$. However for large σ_{eff} and $\lambda \leq 0.2$ for $A \sim 240$ a contribution of the weakly interacting component could at most become comparable to the soft contribution, see Figs. 3,4 below.

4.2 Models for shadowing at Q_0^2 and numerical results for Q^2 dependence

Based on the above discussion of the fluctuation effects we adopt the following prescription for the calculation of the gluon and quark nuclear shadowing for $A \geq 10$:

(i) For $Q = Q_0$ we use two models: the quasieikonal model and the fluctuation model with $\lambda = 0.2$ analogous to the ones we used in [4,13]. The relation between the amount of shadowing in the two models is given by

$$\begin{aligned} & G_A(x,Q_0)/G_N(x,Q_0)_{fluct.mod.}(\sigma_{eff}(x,Q_0)) \\ & \quad = \lambda + (1-\lambda) G_A(x,Q_0)/G_N(x,Q_0)_{quasieik.mod.} \\ & \quad \cdot \left(\frac{\sigma_{eff}(x,Q_0)}{1-\lambda}\right) \end{aligned} \qquad (20)$$

(ii) To constrain the behavior of the gluon density at $x \geq 0.02$ in the normalization point we use the analysis of [5] which indicates that gluons in nuclei carry approximately the same fraction of the momentum as in a free nucleon:

$$\int_0^1 dx x G_A(x,Q^2) \approx \int_0^1 dx x G_N(x,Q^2). \qquad (21)$$

This allows to estimate the amount of the gluon enhancement at $x \geq x_{max}$ assuming that it should be concentrated at $x \leq 0.2$ where average longitudinal distances (the Ioffe distances) contributing to the parton density are comparable to the internucleon distances - this procedure is similar to one we introduced in [5].

(iii) Based on the rational presented above we use the DGLAP evolution equations to calculate the nuclear shadowing for larger Q^2.

Since in this paper we are interested primarily in the behavior of the nuclear gluon and quark densities at $x \leq 0.01$ we are not sensitive to details of the enhancement pattern. So we do not try to introduce an A dependent shape for the enhancement and assume that the enhancement is present for $0.02 \leq x \leq 0.2$ and can be approximated by $G_A(x,Q_0)/G_N(x,Q_0) = C(A)(x-0.02)(0.2-x)$. We also do not model small enhancement for $F_{2A}(x,Q^2)$ and $V_A(x,Q^2)$ at $x \sim 0.1$. In the calculations we use the standard Fermi step fit to the nuclear densities: $\rho(r) = C/(1+\exp((-R+r)/b))$, where $R = 1.1fm \cdot A^{1/3}, b = 0.56fm$. First, in Fig. 3 we present a comparison of the gluon shadowing calculated in two models at $Q_0 = 2$GeV and $x = 10^{-3}, 10^{-4}$. One can see that shadowing for $A \sim 12$ is already very significant, though the effect of fluctuations is still small. With a further increase of A the effect of fluctuations becomes larger and it reaches a factor of 1.4 for A$\sim$ 200. For these A in the fluctuation model weakly interacting configurations contribute approximately half of the cross section. Note however that we expect that in a more realistic model of fluctuations a relative contribution of these configurations is likely to decrease with decrease of x, see discussion in Sect. IV.A. In Figs. 4, 5 we present results for the Q^2 dependence of shadowing for gluons and for quarks for A=12, 40, 100, 200 calculated in the quasieikonal model. One can see that the gluon shadowing is large already for A=12 and for heavy nuclei reaches the

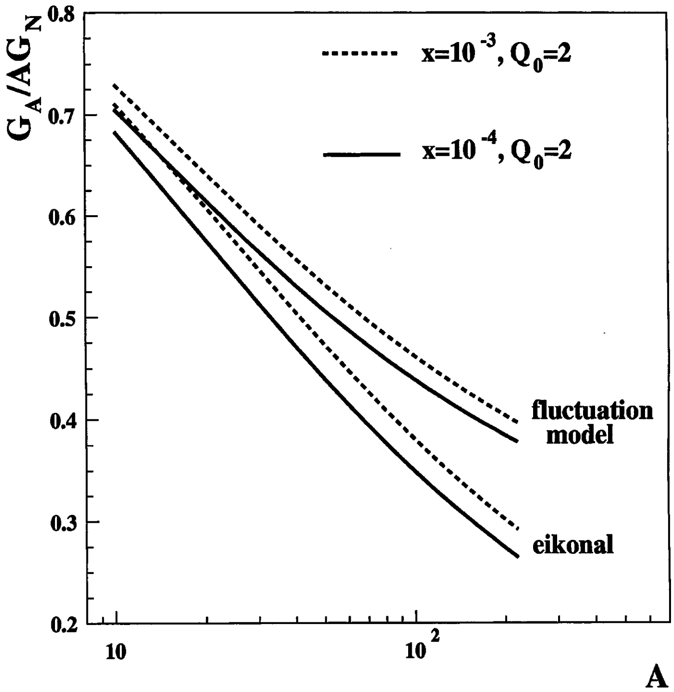

Fig. 3. Comparison of the gluon shadowing calculated in the quasieikonal model and the fluctuation model for $x = 10^{-3}, 10^{-4}$

level of $A_{eff}/A \sim 0.2$. (Variations of the parameter σ_{eff} within a factor of 1 ± 0.2 allowed by the uncertainties in the value of the gluon diffractive density at $t = 0$ lead to the similar variations of G_A/AG_N at $x \leq 10^{-3}$ and of $\left[1 - G_A(x, Q^2)/AG_N(x, Q^2)\right]$ for $0.02 \leq x \leq 0.2$.) Shadowing decreases with increase of Q^2 but remains large up to very large Q^2. The shadowing in the charged parton channel is much smaller and rather weakly decreases with Q^2. In fact, at $x \sim 10^{-4}$ shadowing for $F_{2a}(x, Q^2)$ first increases with increase of Q^2 due to a larger shadowing in the gluon channel. Note also that (21) leads to a rather large enhancement of G_A/G_N for $0.05 \leq x \leq 0.2$. However the enhancement is rather large already for $A = 12$. As a result a further growth of the enhancement between carbon and tin is of the order of 10% and it is well consistent with the analysis of the current data in [36]. The decrease of shadowing with increase of Q^2 is due to the feeding of the small x by partons which originated from $x \geq 0.02$ at Q_0^2. Effectively, as we discussed above, the QCD evolution leads to fluctuations in the value of interaction strength due to the mixing of the contributions of small and large l_c. Therefore the use of (17) in the quasieikonal model with $\sigma \propto \sigma_{eff}$ defined at a high resolution Q^2 would lead to an overestimate of the nuclear shadowing since the cross section fluctuations lead to a decrease of the shadowing for the fixed value of σ_{eff}, see discussion in [13].

Note also that it is often stated in the literature that nuclear shadowing for the total cross section of $\gamma^* A$ should be practically the same for the real photon and for virtual photons with moderate (few GeV^2) virtualities. However our analysis predicts a significant drop of the nuclear shadowing in the total cross sections of $\gamma^* A$ scattering between $Q^2 = 0$ and $Q^2 = 4GeV^2$ due to the strong Q^2 dependence of the diffraction contribution to $\sigma_{\gamma^* N}$. The results of our calculation using data on the diffraction in the γp scattering at $W \sim 14GeV$ [37] and at HERA [38] and a smooth interpolation between two energy ranges is presented in Fig. 6. For a rough estimate we use the quasieikonal approximation which leads to a slight overestimate of shadowing for the heavy nuclei.

4.3 Structure functions of nuclei in the the black disk limit

It is of interest to consider also the structure functions in the limit Q^2=const, $x \ll \frac{1}{m_N R_A}$, $A \to \infty$ first analyzed by Gribov for σ_T, σ_L [1]. It was argued in this paper that in such a limit interactions of all essential configurations in the virtual photon wave function with a nucleus can be treated in the black body (S-channel unitarity) limit. Since in the black disc limit the dispersion over the strengths of interactions can be neglected one finds [1]

$$\frac{1}{Q^2} F_{2A}(x, Q^2) = \frac{\pi R_A^2}{12\pi^2} \int_{m_o^2}^{\frac{W^2}{2m_N R_A}} \frac{m^2 \rho(m^2) dm^2}{(m^2 + Q^2)^2}, \quad (22)$$

and

$$\frac{1}{Q^2} F_{L\ A}(x, Q^2) = \frac{\pi Q^2 R_A^2}{12\pi^2} \int_{m_o^2}^{\frac{W^2}{2m_N R_A}} \frac{\rho(m^2) dm^2}{(m^2 + Q^2)^2}, \quad (23)$$

where $\rho(m^2) = \sigma(e^+e^- \to hadrons)/\sigma(e^+e^- \to \mu^+\mu^-)$.

We can generalize the Gribov formulae deduced for the sea quark distribution [1] to the case of the gluon channel:

$$xG_A(x, Q^2)/Q^2 = \frac{3}{2} \frac{\pi R_A^2}{12\pi^2} \int_{m_o^2}^{\frac{W^2}{2m_N R_A}} \tilde{\rho}(m^2) \frac{m^2 dm^2}{(m^2 + Q^2)^2} \quad (24)$$

Here $\tilde{\rho}(m^2) = \sigma(\hat{J} \to hadron)/\sigma(\hat{J} \to 2g)$ is the ratio of the cross section of the gluon hadronic processes initiated by the local operator $\hat{J} = \frac{1}{\sqrt{-\Box}} F^{\alpha}_{\mu\lambda} F^{\mu\lambda}_{\alpha}$ introduced in [19] to the perturbative cross section of the annihilation of the gluon source into gluons. For $m^2 \to \infty$, $\tilde{\rho}(m^2) = 1$. So in the black disk limit

$$xG_A(x, Q^2)/Q^2 = \frac{\pi R_A^2}{8\pi^2} \ln\left(\frac{W^2}{Q^2 2m_N R_A}\right) \equiv \frac{\pi R_A^2}{8\pi^2} \ln\left(\frac{x_0}{x}\right), \quad (25)$$

where $x_0 = \frac{1}{2m_N R_A}$. Equation (25) provides a solution for the problem of the gluon nuclear shadowing in the theoretical limit Q^2=const, $x \ll x_0$, $A \to \infty$. It also illustrates that investigation of the Q^2 dependence of the nuclear structure functions will provide an effective method to establish whether fluctuations of strengths play a significant role in the gluon structure functions.

It is useful to compare (25) with the impulse approximation value of $G_A = AG_N$. For example for $A = 200$ this equation leads to $xG_A(10^{-4}, 10GeV^2)/A =$

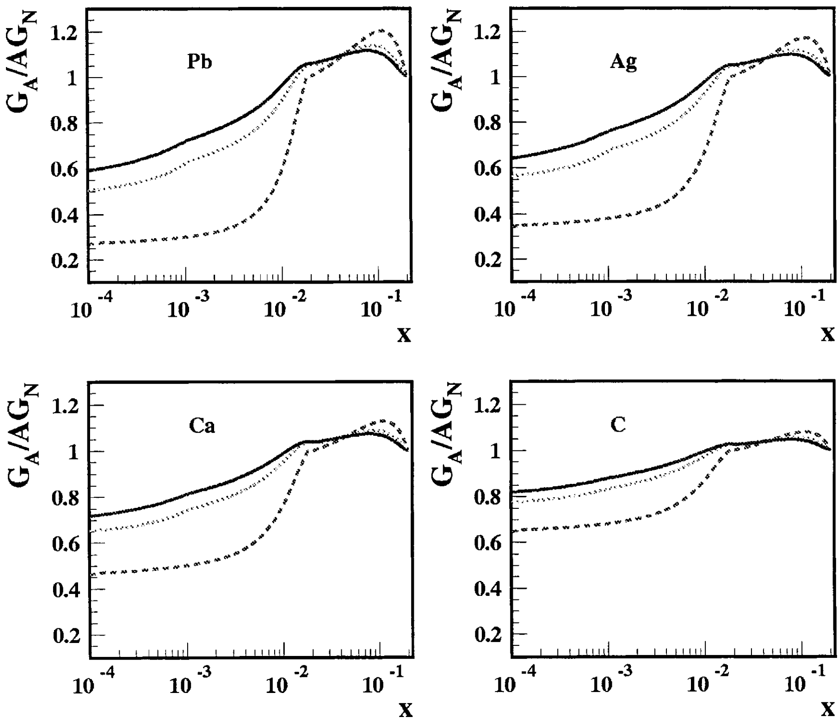

Fig. 4. Dependence of G_A/AG_N on x for Q=2,5,10 GeV (dashed, dotted, solid curves) calculated in the quasieikonal model

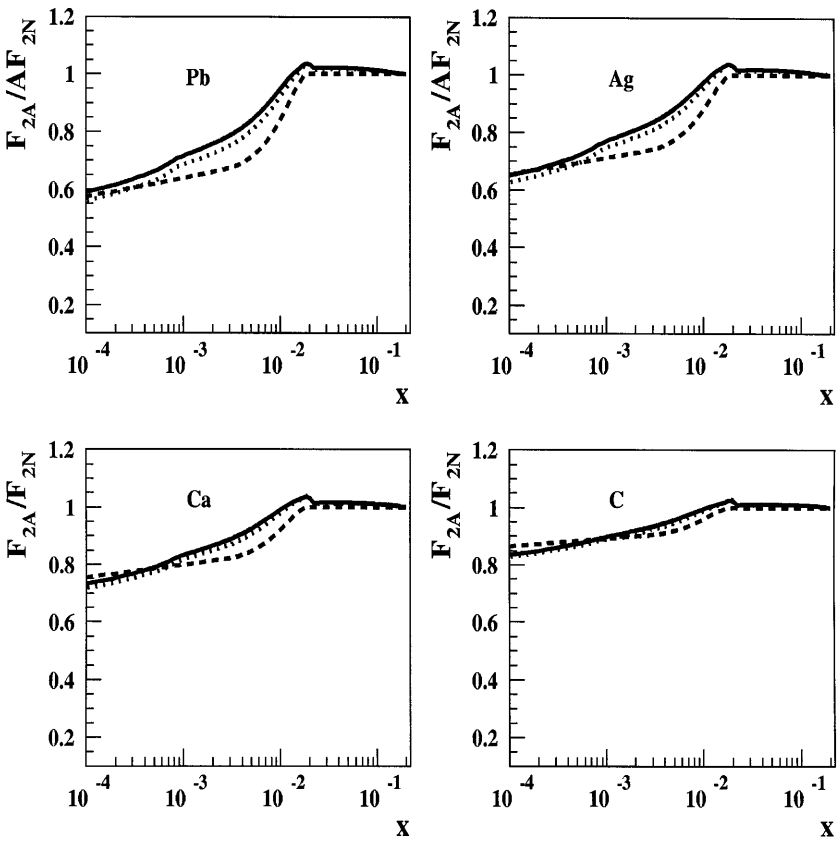

Fig. 5. Dependence of F_{2A}/AF_{2N} on x for Q=2,5,10 GeV (dashed, dotted and solid curves) calculated in the quasieikonal model

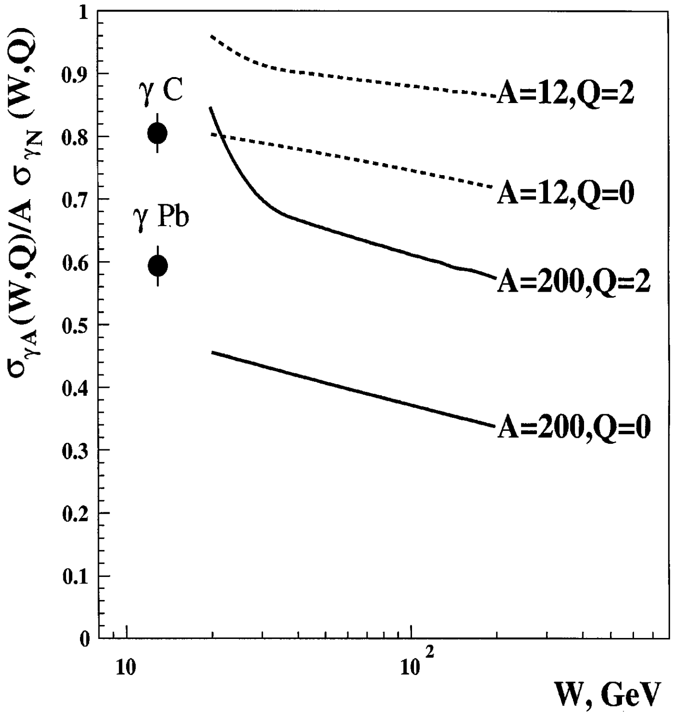

Fig. 6. Dependence of $\frac{\sigma_{\gamma^* A}}{A\sigma_{\gamma^* N}}$ on W for Q=0,2 GeV (dashed and solid curves). The data points are from [39]. They are corrected for the small effect of the nuclear shadowing in the deuteron

8 while the current fits to the nucleon data lead to $xG_N(10^{-4}, 10GeV^2) \approx 20$. Hence the model independent unitarity constrain implies a large gluon shadowing for this kinematics. This is consistent with our model calculations presented in the previous subsection.

5 Onset of color opacity regime in hard diffraction and suppression of minijet production in nucleus-nucleus collisions

Let us now briefly discuss some of the consequences of the found magnitude of the gluon shadowing.

The production of minijets is often considered as an effective mechanism of producing high densities in the head on heavy ion collisions. However in the LHC kinematics for the central rapidities minijets are produced due to collisions of partons with $x_{jet} = \frac{2p_t}{\sqrt{s_{NN}}}$. For heavy ion collisions $s_{NN} \geq 4\ TeV$ and the gluon-gluon collisions are responsible for production of most of the minijets. Therefore the gluon nuclear shadowing would lead to a reduction of the rate of the jet production due to the leading twist mechanism by a large factor up to $p_t \sim 10GeV/c$, see Fig. 7 where we give results of calculation in the quasieikonal and fluctuation models. The nuclear gluon shadowing leads to a similar very strong reduction of the heavy onium production in pA and AA collisions at LHC energies for $y_{c.m.} \sim 0$ and small p_t.

For the central impact parameters the reduction is even larger. Using generalized eikonal approximation of

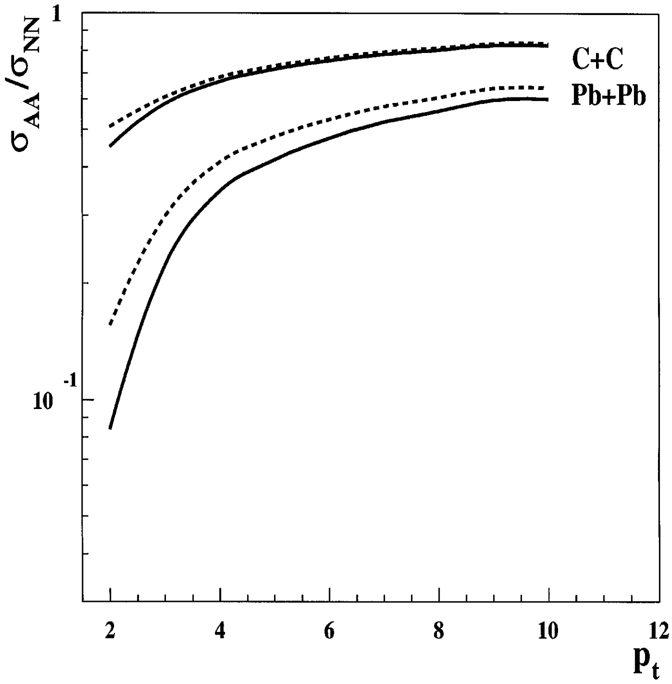

Fig. 7. Suppression of the jet production in AA collisions due to gluon shadowing at $y = 0$ calculated in the quasieikonal and fluctuation models (solid and dashed curves)

(17) we can calculate also the suppression of the parton densities at a given impact parameter b as [40]:

$$G_A(x, Q_0^2, b)_{shadowed} = \frac{\int d\sigma P_g(\sigma)(2 - 2\exp(-T(b)\sigma/2)G_N(x, Q_0^2)}{\int P_g(\sigma)\sigma d\sigma}. \quad (26)$$

Since in the impulse approximation $G_A(x, Q^2, b)_{imp} = G_N(x, Q^2)T_A(b)$ we finally obtain

$$G_A(x, Q_0^2, b)/G_N(x, Q_0^2) = \frac{\int d\sigma P_g(\sigma)(2 - 2\exp(-T(b)\sigma/2)}{T(b)\int d\sigma\sigma P_g(\sigma)}. \quad (27)$$

The results of calculation using (27) are presented in Fig. 8.

So we conclude that the shadowing effects are likely to reduce very substantially the parton densities generated at the first stage of heavy ion collisions at the LHC energies. However theoretical uncertainties related to the role of point-like configurations lead to rather large uncertainties in the estimate of the suppression for the case of the heavy ion collisions. Hence it would be very important to perform a direct measurement of the gluon nuclear shadowing at HERA in this kinematics.

Presence of a large gluon shadowing leads to large effects in the diffractive eA collisions at HERA energies. Here we consider the simplest example - coherent diffractive production of vector mesons at large Q^2 by the longitudinally polarized photons - $\gamma_L^* + A \rightarrow V + A$, and photo(electro) production of heavy onium states. Since these processes are dominated by production of $q\bar{q}$ in a

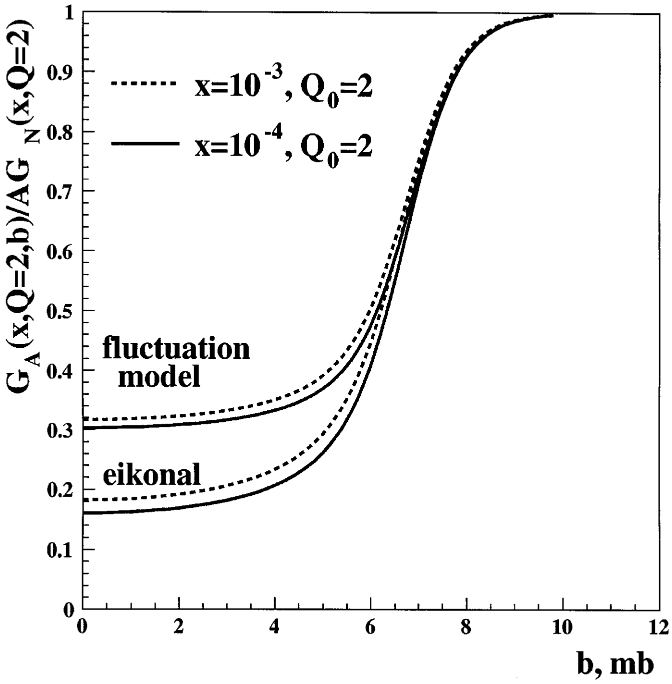

Fig. 8. Impact parameter dependence of the gluon shadowing for the scattering of *Pb* for $Q = 2\text{GeV}$ and $x = 10^{-3}, 10^{-4}$ calculated in the quasieikonal and fluctuation models

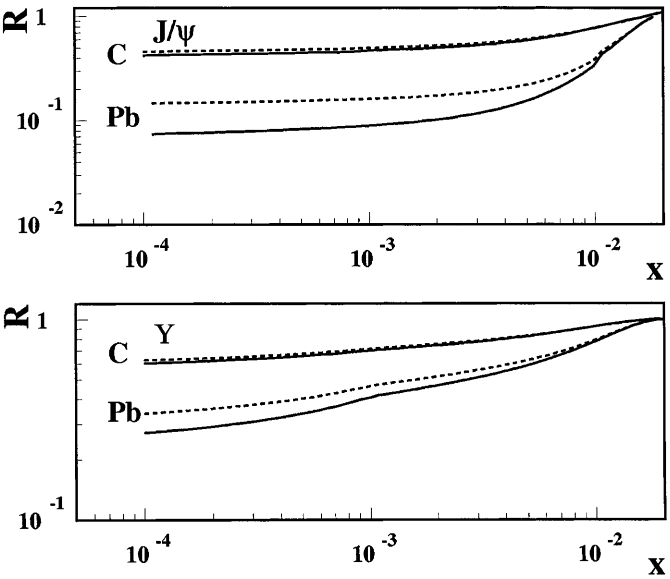

Fig. 9. Color opacity effect for the ratio of the coherent production of J/ψ and Υ from carbon(lead) and a nucleon normalized to the value of this ratio at $x = 0.02$ calculated in the quasieikonal and fluctuation models (solid and dashed curves)

small size configuration one may naively expect that color transparency should hold for such processes and the amplitudes of these processes should be proportional to A. However the QCD factorization for exclusive processes leads to the amplitude of this process been proportional to $G_A(x, Q^2)$ [24]. For small x and t

$$\frac{d\sigma^{\gamma_L^*+A\to V+A}}{dt} = F_A^2(t)\frac{G_A^2(x,Q^2)}{G_N^2(x,Q^2)}\frac{d\sigma^{\gamma_L^*+N\to V+N}}{dt}. \quad (28)$$

So we expect that the color transparency regime for $x \geq 0.02$ (with a small enhancement at $x \sim 0.1$ due to enhancement of $G_A(x, Q^2)$ for these x) would be followed by the color opacity regime for $x \leq 0.01$. As an illustration in Fig. 9 we present the ratio of the cross section of the J/ψ and Υ production off nuclei with $A = 12, 200$ and nucleon calculated under the assumption that the leading twist gives the dominant contribution in these processes. It is plotted as a function of x and normalized to the value of the ratio at $x = 0.02$. In the calculation we use the analysis of [41] which indicates that $Q^2_{eff} \approx 5(40)GeV^2$ for $J/\psi(\Upsilon)$ production. One can see from the figure that we expect onset of the Color Opacity regime for J/ψ starting at $x \sim 0.01$. The Color Opacity effect remains quite significant for production of Υ.

6 Nuclear effects in the inclusive leading hadron spectra in eA collisions

Numerous data on hadron-nucleus scattering at fixed target energies indicate that the multiplicities of the leading hadrons $N_A(z) \equiv \frac{1}{\sigma_{tot}(aA)}\frac{d\sigma(z)^{a+A\to h+X}}{dz}$ decrease with increase of A. Here z is the light-cone fraction of the projectile "a" momentum carried by the hadron "h". On the contrary, the QCD factorization theorem for the inclusive hadron production in DIS implies that in the case of DIS electron-nucleus scattering no such dependence should be present. This indicates that there should be a interesting transition from the soft physics dominating in the interactions of real photons with nuclei to the hard physics in the inclusive hadron production in the DIS kinematics. It would be manifested in the disappearance of the A-dependence of the leading spectra at large z:

$$N_A(z, Q^2) = N_N(z, Q^2),\ for\ z \geq 0.2, Q^2 \geq\ few\ GeV^2, \quad (29)$$

At small x a new interesting phenomenon should emerge due to the presence of diffraction and nuclear shadowing for smaller z. Indeed, the diffraction originates from the presence in the wave function of γ^* of partons with relatively small virtualities which screen the color of the leading parton(partons) with large virtuality and can rescatter elastically from a target (several target nucleons in the case of nuclear target). Inelastic interactions of these soft partons with several nucleons should lead to a plenty of new revealing phenomena in small x DIS eA scattering, which resemble hadron-nucleus scattering but with a shift in rapidity from $y_{max}(current)$ related to the average rapidities of these soft partons. This shift can be expressed through the average masses of the hadron states produced in the diffraction:

$$y_{soft\ partons} \sim y_{max} - \ln(\langle M^2_{dif}\rangle/\mu^2), \quad (30)$$

where $\mu \sim 1GeV$ is the soft scale. Partons with these rapidities will interact in multiple collisions and loose their

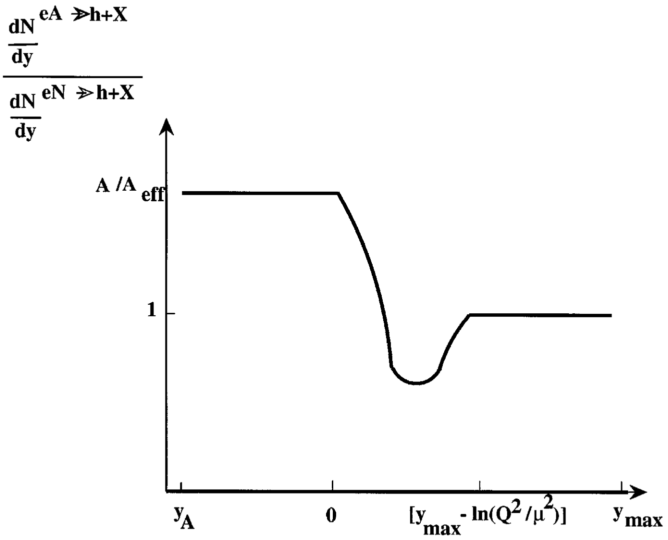

Fig. 10. Sketch of the A dependence of the hadron multiplicity $\frac{dN}{dy}$ for small x plotted as a function of hadron rapidity

energy leading to a dip in the ratio $\eta_A(y) \equiv N_A(y)/N_p(y)$. At the same time these multiple interactions should generate larger multiplicities at smaller rapidities. Application of the AGK rules indicates that for $y \leq y_{soft\ partons} - \Delta$, where $\Delta = 2 - 3$ the hadron multiplicity in the case of nuclei will be enhanced by the factor:

$$\eta_A(y) = \frac{AF_{2p}(x, Q^2)}{F_{2A}(x, Q^2)}. \tag{31}$$

At the rapidities close to the nuclear rapidities a further increase of $\eta_A(y)$ is possible due to formation of hadrons inside the nucleus. A sketch of the expected rapidity dependence of $\eta_A(y)$ is presented in Fig. 10.

One also expects a number of phenomena due to long range correlations in rapidity. This includes: (a) Local fluctuations of multiplicity in the central rapidity region, e.g. the observation of a broader distribution of the number of particles per unit rapidity, due to fluctuations of the number of wounded nucleons [9].

These fluctuations should be larger for the hard processes induced by gluons, for example the direct photon production of two high p_t dijets. (b) Correlation of the central multiplicity with the multiplicity of neutrons in the forward neutron detector, etc.

Another important manifestation of nuclear shadowing is a large probability of diffractive final states for small x. In the case of the generic $e + A$ scattering this probability would reach $\sim 35\%$ for $A \sim 200$ [9]. Due to a larger effective cross section of interaction in processes induced by the hard interactions with gluons, this effect should be even more pronounced in the lepton scattering processes where a hard process corresponds to hard $\gamma^* - g$ interaction. Hence for example in the charm electroproduction of a heavy nucleus we expect about a half of the events to originate from the coherent diffraction where the nucleus remains intact [13]. Also one expects more strong filtering out of the gluon dominated diffraction than the diffraction dominated by the coupling to the quarks.

7 Comparison with other approaches

Several approaches were developed over last decade to the dynamical calculation of nuclear shadowing phenomenon in DIS [11]

First group of approaches is based on the Gribov work [1] which established connection between the diffraction and nuclear shadowing. Among these considerations [2–11] the one of [11] is the most detailed and comes closest to our analysis in the case of F_{2A}. The analysis is based on the fit to the HERA diffractive data on the $e+p \to e+X+p$ reaction within the model developed by the authors and a set of assumptions about higher order screening effects which are anyway rather small for F_{2A}. They demonstrate that the model can well describe the NMC data at $x \sim 0.01$ and give predictions for the HERA kinematics. The main differences from our approach are the use of the model for diffraction which does not explicitly satisfy the QCD factorization theorem for diffraction in DIS and neglect by the effects of enhancement of gluon distributions in nuclei at $x \sim 0.1$. Besides the gluons appear to play a rather small role in their model of ep diffraction leading to the expectation of the gluon shadowing smaller than in the case of F_{2A} as compared to the larger shadowing for gluons expected in our analysis.

In the case of hadron-nucleus scattering both total cross section and inelastic diffraction can be described based on the idea of the fluctuations of the interaction strength in the projectile treating interaction of each component in the eikonal approximation, for the review and references see [42]. There are a several models where a similar approach has been applied to the calculation of the nuclear shadowing by introducing the impact parameter $q\bar{q}$ virtual photon wave function $\Psi_\gamma(b)$ and introducing the cross section of the $q\bar{q}$-N interaction for the fixed b, see e.g. [43,44] and references therein. So far it was assumed in these models that $\sigma_{q\bar{q}-N} = cb^2$. Hence the QCD evolution which leads to a fast increase of $\sigma_{q\bar{q}-N}$ with incident energy was neglected. Moreover in this approximation shadowing for the small b configurations is a higher twist effect, leading to an expectation of lack of the leading twist shadowing for σ_L and lack of gluon shadowing.

Another group of approaches uses the infinite momentum picture and treats all the process within the perturbative QCD. To avoid problems with positivity of the cross section for $A_{eff}/A \leq 0.75$ (see discussion in section XX) one has to include interactions with $N \geq 3$ nucleons. The current models which include such interactions (see [45–47] and references therein) assume that all shadowing is generated perturbatively and do not include information about diffractive processes. Qualitative expectations of these models are a rather large gluon shadowing

[11] A phenomenological approach based on fitting the existing nuclear data and imposing the momentum and baryon sum rules in the spirit of [5] was pursued in [12,15].

and significant nonlinear effects in the evolution of the parton densities. It would be interesting to compare two approaches after the leading twist shadowing effects are implemented and constraints following from the HERA diffractive data are taken into account.

8 Conclusions

We have demonstrated that dominance of gluons in the "Pomeron" diffractive parton densities leads to a large enhancement of the nuclear gluon shadowing in a wide range of x, Q. Gluon shadowing of this magnitude will strongly affect the first stage of the heavy ion collisions at LHC, lead to a number of Color Opacity phenomena in the HERA kinematics for the eA collisions. Study of the gluon shadowing for heavy nuclei may allow to enhance contribution of the small interaction strengths, allowing to unmask PQCD physics in the eA collisions at HERA. The studies of coherent diffraction off nuclei and hadron production in the inclusive eA scattering will provide complementary handles for studying the small x dynamics.

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