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Diffractive production of charm in DIS and gluon nuclear shadowing

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Abstract. We evaluate nuclear shadowing of the total cross section of charm particles production in DIS within the framework of Gribov theory of nuclear shadowing generalized to account for the QCD evolution. We use as an input the recent QCD Pomeron parton density analysis of the HERA diffractive data. Assuming that the QCD factorization theorem is applicable to the charm production off nuclei we also calculate shadowing of the gluon densities in nuclei and find it sufficiently large for heavy nuclei: $G_{A\sim 200}(x,Q^2)/AG_N(x,Q^2) \sim 0.45 - 0.5 \cdot (A/200)^{-0.15}$ for $x \sim 10^{-3 \div -4}, Q^2 \sim 20 \div 40 GeV^2$ to influence significantly the physics of heavy ion collisions at LHC. We evaluate also suppression of minijet and hidden charm production in the central AA collisions. We also discuss some properties of the final states for $\gamma^* A$ processes dominated by the scattering off small x gluons like the high p_t jet and charm production.

PACS. 13.60.-r Photon and charged-lepton interactions with hadrons – 24.85.+p Quarks, gluons, and QCD in nuclei and nuclear processes

1 Introduction

The aim of this paper is to evaluate gluon shadowing in nuclei in the HERA kinematics using generalised Gribov theory of nuclear shadowing and information on diffraction processes in DIS obtained at HERA.

The paper is organized as follows. In Sect. 2 we review the Gribov theory of nuclear shadowing for $F_{2A}(x,Q^2)$ [1] which established a connection of this phenomenon with the process of the diffraction in the elementary process of the scattering of a projectile off a single nucleon (4). In essence, this connection is due to a possibility of very low momentum transfer to a nucleon in the diffraction processes comparable to internal momenta of nucleons in nuclei which leads to a significant interference of the amplitudes of diffractive scattering off different nucleons of the nucleus. We explain also that the Gribov theory takes into account only contribution of large longitudinal distances in deep inelastic scattering. Hence one has to take into account an additional QCD effect of the change of the average longitudinal distances with increase of Q^2 at fixed x. This ammounts to taking into account the leading twist QCD evolution of the parton densities at $x \geq 0.02$ at $Q^2 \geq Q_0^2$. Numericaly this effect is very small in the kinematics which we study in this paper.

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In Sect. 3 we focus on shadowing in the charm production off nuclei and gluon nuclear shadowing. We are able to extend the Gribov theory to these processes due to the recent systematic studies at HERA of the diffractive phenomena in DIS, and in the real photon induced hard processes, see [2]. One important observation beared by all analyses of the HERA data as well as of the hard hadronic diffraction data is the dominance of the gluon degrees of freedom as compared to quark degrees of freedom (by a factor $\sim 5-6$). This dominance at $Q^2 \leq 20 GeV^2$ is much stronger than in the case of small x parton densities at $x \leq 10^{-4}$ (a factor ~ 2). Hence we come to *a qualitative conclusion that shadowing in the gluon channel should be larger ($\sim 2-3$) that in the charge parton channel.*

To quantify this conclusion we use the global analysis of the diffractive HERA data [3] wich has demonstrated that the scaling violation in $\sigma_{diff}(x,Q^2)$, diffractive production of jets in the γp scattering, diffractive production of charm in DIS can be described based on the leading twist factorization approximation for the parton densities of diffraction which are often referred to as the Pomeron structure functions. In this paper we focus on the Q^2 range covered by the HERA diffractive experiments sensitive to the gluon intiated diffraction: $Q^2 \sim 20-50$ GeV^2 and the $x \leq 10^{-3}$ range where average longitudinal distances in the $\gamma^* - A$ interactions are much larger than the nucleus size. We will discuss the Q^2 dependence of shadowing for lower Q^2 and the whole range of x in the subsequent paper [4]. Our results for the gluon shadowing are given by (6), (9), (10). They provide a direct connection between the gluon shadowing due to the interaction with two nucleons and hard diffraction processes with nucleons dominated by the gluon degrees of freedom. The Gribov theory does not produce unique predictions for the contribution of the interactions where three or more nucleons are involved. As a starting point we treat these interactions neglecting effects of fluctuations of the strength of interaction of the projectile with a target. This corresponds to the quasieikonal approximation in the hadron-nucleus scattering. The results of numerical studies in this approximation presented in Fig. 2 confirm that shadowing in the gluon channel is indeed larger (by about a factor of two for $Q^2 \sim 20GeV^2$). than the quark shadowing.

Next we analyse effects of the fluctuation of the strength of interaction. Already simple dimensional arguments indicate that these fluctuations in our case should be large. We analyse this effect and find that indeed it becomes important in the case of suffiently large A which ammounts to a violation of the quasieikonal approximation and filtering out of the contribution of the strongly interacting configurations. However we find that even for central impact parameters for heavy nuclear targets ($A \leq 250$) this effect remains a correction. One needs $A \geq 1000$ to reach the region where these effects dominate. In the end of the paper we analyse the final states in the charm production. We predict that for the case of scattering of heavy nuclei coherent diffractive photoproduction of charm would constitute nearly half of the total charm production cross section. We also predict strong screening effects for coherent exclusive production of vector mesons in DIS $Q^2 \leq 20GeV^2$ and $x \leq 10^{-3}$.

2 Shadowing of quark distributions within a nucleus

First let us summarize necessary elements of the Gribov theory in the case of DIS scattering of a deuteron where calculation can be performed in a model independent way. Theoretical description of nuclear shadowing is simplified in the nucleus rest frame. The starting point is the presence of a large coherence length in the interaction of a γ^* with a nucleon with mass m_N: a virtual photon with the "mass^2" $\sim -Q^2$ fluctuates into a quark-gluon configuration of characteristic invariant mass M_X^2 over the average longitudinal distances (in the nucleus rest frame) (see [5] and for further development see [6] and references wherein)

$$l_{coh} = \frac{1}{m_N x(1 + M_X^2/Q^2)}. \qquad (1)$$

l_{coh} corresponds to the average distance over which the light-cone wave function of a photon which includes radiation of partons is build. In preQCD period it has been assumed that $M_X^2 \approx Q^2$. In this case the expression for coherence length obtaines famous form

$$l_{coh} = \frac{1}{2m_N x} \ , \qquad (2)$$

which by far exceed the nuclear radius even for heavy nuclei. $l_{coh} \gg R_A$.

For this limit Gribov has established the connection between the phenomenon of diffraction in the scattering off a nucleon and the presence of the nuclear shadowing [1]. The key elements of the derivation are the dominance of the vacuum exchange in the t-channel of the amplitude of the aN scattering, which is well established now experimentally, and the possibility to describe the deuteron and heavier nuclei as loosely bound systems build of nucleons. The difference in the scales characterizing a nuclear bound state (average internucleon distances in nuclei $r_{NN} \approx 1.7Fm$ and somewhat larger within the deuteron) and the nucleon radius, $r_N \approx 0.8Fm$, justifies taking into account only the contribution of the nearest pole in the nuclear vertex. A similar assumption is routinely made in the Glauber theory of the intermediate energy hadron-nucleus scattering which is known to work with a few percent accuracy.

In the case of the deep inelastic scattering off the deuteron the only quantity which enters into consideration is the simultaneous interaction of the virtual photon with two nucleons. In this case

$$\sigma^D_{shad} \equiv \sigma_{\gamma^* p} + \sigma_{\gamma^* n} - \sigma_{\gamma^* D}, \qquad (3)$$

is directly expressed through the cross section of diffraction in the $\gamma^* N$ processes. (To simplify notations we do not write explicitly in (3) differential in x and Q^2.) One has [1]

$$\sigma^D_{shad} = \eta \frac{1}{4\pi} \int S(4t) \frac{d^2\sigma^{\gamma^*+N\to X+N}}{dM^2 dt} dM^2 dt. \qquad (4)$$

Here $d^2\sigma/dM^2dt$ is the cross section for producing a state X with mass M and four-momentum transfer squared t in the inclusive diffractive reaction $\gamma^* + N \to X + N$. $S(t)$ is the electromagnetic form factor of the deuteron, $-t = k_t^2 + (x m_N(1 + M_X^2/Q^2) + k_t^2/2m_N)^2$. For simplicity we neglect spin effects in the wave function of the deuteron. The factor η is due to nonzero value of the real part of the amplitude A of the diffractive processes, [7]: $\eta = \frac{(1-\lambda^2)}{(1+\lambda^2)}$, and

$$\lambda = ReA/ImA \approx \frac{\pi}{2} \frac{\partial \ln(A/s)}{\partial \ln 1/x}. \qquad (5)$$

The current data for cross section of inclusive diffraction at HERA appear to indicate that $\lambda = 0.2 - 0.3$, leading to $\eta = 0.8 \div 0.9$. Note that in (4) there is no separation between leading and nonleading twist effects. This shortcoming is practically unimportant to the extent that we restrict ourselves by the leading twist effects only.

The account of the hard physics requires certain modifications of the Gribov formulae for the limit of small but fixed x and large Q^2. The reason is that the hard QCD physics – emission of gluons leads to the change of the relation between the coherence length and x with increase of Q^2. For given x average l_{coh} decreases with increase of Q^2

so that at extreemly large Q^2 contribution of $l_{coh} \leq 2fm$ would dominate leading to disappearance of nuclear shadowing. As a result in this limit one becomes sensitive to the A-dependence of the nuclear parton densities at $x \geq 0.02$: an enhancement of gluon and valence quark parton densities at $x \sim 0.1$ and the EMC effect at $x \geq 0.4$. As soon as the nuclear parton densities at $x \geq 0.02$ are known at a low Q^2 normalization point this effect can be taken into account using QCD DGLAP evolution equations [8, 9]. This leads to the QCD generalization of the Gribov formulae. For heavier nuclei interactions with $N \geq 3$ nucleons become important and hence one needs a more detailed information about diffraction to determine the shadowing effects. Considerations of structure function of nuclei F_{2A} in the framework of the Gribov theory were performed in a number of papers, see e.g. [8,6,9–12,7,13–15] though effects of enhancement of gluon densities were neglected in [10–15]. Note in passing that the eikonal approximation maybe a good starting approximation for multiparton configurations. On the other hand if initial configuration consists of two bare partons only inelastic interactions with two nucleons may contribute.

The phenomenon of nuclear shadowing for the quark sea $F_{2A}(x,Q^2)$, as the leading twist effect is well established experimentally, see e.g. [16] though in a rather limited range of DIS kinematics. Further studies of this phenomenon which is one of the key elements of the physics of AA collisions at LHC may be possible at HERA in a eA mode [17].

It is possible to describe quantitatively these data by applying the Gribov theory of nuclear shadowing [1] which connects the shadowing for $F_{2A}(x,Q^2)$ and diffraction in DIS scattering off a nucleon for the same x, Q^2 [8,6,9, 11,7,13,15]. In the early papers one had to start with modeling the cross section of diffraction in DIS which was done based essentially on the QCD generalization of the Bjorken aligned jet model [8]. More recently one could employ information on diffraction in the HERA kinematics to reduce the model element of considerations [7,12–15]. Overall uncertainties involved in the estimates of higher order terms appear to be rather small - see discussion below. At the same time the effects leading to an enhancement of the parton densities for $x \sim 0.02-0.2$ which so far could be treated only on the basis of the momentum and baryon charge sum rules [8] prevents quantitative predictions for $x \geq 0.01$.

It is worth emphasizing, that in the case of shadowing in $\gamma^* A$ reaction one expects that there will be no smooth transition from the real phton limit to the limit of small $x \leq 10^{-3}$ and $Q^2 \sim$ few GeV^2 [18] since the small t diffraction for the real photon case in $2 \div 3$ times larger that in the DIS limit. At the same time for larger Q^2 rapidity gap probability weakly depends on Q^2 leading to a weak dependence of shadowing on Q^2.

Note in passing that several approaches to nuclear shadowing have been developed based on the light-cone treatment of the nucleus wave function, for a recent review and references see [19]. If one imposes within such an approach a condition that the nucleus is build of weakly bound nucleons, this approach could be considered as complementary to the approach outlined above. However so far the light-cone approaches discussed in [19] did not deal with the diffractive physics and hence it is not clear whether approximations currently used in these approaches are consistent with the HERA diffractive data.

3 Nuclear shadowing of charm quark and gluon distributions in nuclei

3.1 Diffractive production of charm as a gluon partonometer

Obviously if one would have the information on diffraction in the scattering of the hard probe coupled directly to gluons – "a gluon partonometer" – one would be able to repeat the same program for the gluon shadowing.

The above mentioned observation of the operational validity of the factorization approximation for the range of x, Q^2, M_X^2 covered by HERA [3] implies that we can actually select the processes dominated by scattering off the gluon field of the target and use them for calculating the gluon shadowing in nuclei. We choose as such a process production of charm because in this case the QCD evolution equation describes well the inclusive cross section while the factorization approximation describes the diffractive charm production. So we will apply the Gribov theory to calculate nuclear shadowing for charm diffractive production. As far as the charm production of nuclei can be described in terms of the QCD factorization theorem we can interpret this shadowing as the shadowing of the nuclear gluon density.

To get a better insight into the picture of nuclear shadowing it is convenient to consider it in the S-channel picture in the nucleus rest frame where the projectile is considered as a superposition of different quark-gluon configurations. In particular, in the case of the small x charm production the virtual photon can be represented as $c\bar{c}$, $c\bar{c}g$, ... configurations which scatter off the target gluon field. Since the $c\bar{c}$ configuration has a small size $\leq 1/m_c$, diffraction in this picture occurs predominantly due to the scattering of the colorless dipole built of $c\bar{c} - g$ in a color octet configuration off the target [20,21]. In many respects these configurations are analogous to the $q\bar{q}$ aligned jet configurations. The cross section of the interaction of the color octet dipole is larger by a factor of 9/4 than that for the color triplet dipole (for the same transverse separation). Hence the relative importance of the diffraction for the probes coupled to gluons and to quarks is determined by the interplay of two factors - larger cross section interaction for color octet dipoles versus the different probability of the production of the large size configurations.

To characterize the dispersion over the strengths of the interaction it is convenient to introduce the notion of cross section fluctuations for the projectile nucleon interactions. In our case we consider fluctuations of the $c\bar{c}$ component of the virtual photon wave function. The relevant quantity is the distribution over the cross section strength $P(\sigma)$. The cross section of diffraction at $t = 0$ is expressed through $P(\sigma)$ using the Miettenen-Pumplin relation [22] generalized to account for the real part of the amplitude:

$$16\pi \frac{\frac{d\sigma_{diff}}{dt}}{\sigma_{tot}} = (1+\lambda^2)\frac{<\sigma^2>}{<\sigma>}\exp Bt, \qquad (6)$$

where $\lambda = Re/Im$ was defined in 5. By definition:

$$\langle \sigma^n \rangle = \int d\sigma \sigma P(\sigma)\sigma^n \qquad (7)$$

The quantity which can be extracted from the analysis of the data within the model of [3]

$$R_{diff}(charm) = \frac{\int_0^{x_0} dt dx_{I\!P} \frac{d\sigma_{diff}^{charm}(M^2,s,x,Q^2)}{dt}}{\sigma_{tot}^{charm}(x,Q^2)}, \qquad (8)$$

where $x_P = M^2/s$, M^2 is the mass[2] of diffractively produced hadron system, and s is the square of the invariant energy of the $\gamma^* N$ collision. We fix $x_o = 0.05$ since the nucleus form factor naturally cuts the contribution of larger values of $x_{I\!P}$. For the fit **D** of [3] which provides the best description of the charm data and dijet production in the real photon scattering we find $R_{diff}(charm)(x,Q^2)$ to be a very weak function of x, Q^2 for $x \leq 10^{-3}$ and $Q^2 \geq 10GeV^2$. For the $10^{-4} \leq x \leq 2.5 \cdot 10^{-3}$ range the ratio changes between from 0.20-0.21 at $Q^2 \sim 8GeV^2$ to 0.16-0.17 at $Q^2 \sim 70GeV^2$. Note that $R_{diff}(charm)$ is larger than the corresponding quantity for the total cross section of diffraction in DIS where it is close to 0.10. Qualitatively this is consistent with above mentioned larger cross section of interaction for the color octet dipole. Recently the t dependence of the inclusive diffractive cross section was measured for $x_L \geq 0.97$ and $Q^2 \geq 3GeV^2$ [23]. The data were fitted assuming an exponential dependence of the cross section on t: $\frac{d\sigma_{diff}}{dt} \propto \exp Bt$, finding $B = 7.1 \pm 1 \pm 1.2GeV^{-2}$. This is consistent with the dominance of strong interaction physics in diffraction[1]. Since the diffraction forms a significant fraction of total cross section of the charm production the large size configurations should give large contribution which is consistent with the analysis of [3] which assumes the Ingelman-Schlein model and the QCD evolution, i.e assumes that diffraction is due to a soft QCD-pomeron exchange. Hence it is natural to assume that the slope B of the diffractive production of charm is the same as that for the inclusive diffraction. Taking $R_{diff} = 0.2$ and neglecting the real part of the amplitude (we will explore sensitivity to the value of λ below) we find[2]:

$$\sigma_{eff} \equiv \langle \sigma^2 \rangle / \langle \sigma \rangle = 16\pi R_{diff} B = 28mb \cdot (1 \pm 0.23) \quad (9)$$

To illustrate the magnitude of the uncertainties we give here errors due to uncertainties in the value of the t-slope. We do not include an error due to uncertainties in the absolute cross section of diffraction which is more difficult to estimate. In the following analysis we will neglect a small decrease of R_{diff} with increase of Q^2 which leads to a slow decrease of the shadowing with increase of Q^2. We will consider this effect as well as the x dependence of shadowing elsewhere [4].

The large value of σ_{eff} for charm diffraction is another indication that relatively soft QCD physics determines the cross section of the charm diffraction.[3] It is worth emphasizing that there is no contradiction between this statement and the applicability of the QCD factorization model. In this model this statement reflects the properties of the Pomeron parton density at the Q^2 boundary for the evolution. This is similar to the situation for $F_{2N}(x,Q^2)$.

3.2 Quasieikonal approximation

To evaluate the shadowing for the scattering of an arbitrary nucleus we have to sum the diagrams where the projectile interacts with 2,3,4... nucleons. The nuclear wave function enters via many-nucleon form factors, which suppress large momentum transfer in the transition amplitudes.

First, let us make a simplifying assumption that all configurations interact with the same cross section σ_{eff}. This corresponds to neglecting the dispersion of the distribution over σ in $P(\sigma)$ - we will explore this assumption below. We also assume that the mass of the system propagating between the points where the photon has transformed into a hadron configuration and where the hadron configuration collapsed back into a photon does not change by a large factor, say a factor larger than 2, or smaller than 0.5, (see discussion of this point below). In this case we can write an expression for $R = \sigma(eA \rightarrow e + charm + X)/A\sigma(eN \rightarrow e + charm + X)$ in the form analogous to the case of vector dominance model - (5.4) of [24]. The only difference is the substitution $m_\rho^2/s \rightarrow (Q^2 + M^2)/s \equiv x_{I\!P}$ [6] and the integration

[1] Note that one may expect that the t-dependence of the diffractive cross section should be better approximated for $-t \leq 1GeV^{-2}$ by $\frac{d\sigma_{diff}}{dt} \propto G_N^2(t)\exp \tilde{B}t$, where $G_N(t)$ is the nucleon electric form factor. Use of such a form would lead to an increase of $\frac{d\sigma_{diff}}{dt}_{|t=0}$ and hence to a larger value of σ_{eff}, cf. Eq.9.

[2] In this estimate we use the parameterization of [3] for all $\beta = Q^2/(Q^2+M^2)$, though for $\beta \sim 1$ higher twist effects are important. However this region gives very small contribution to the total diffractive cross section.

[3] σ_{eff} may be overestimated within the above evaluation since one of characteristic features of hard diffraction is a smaller slope B. But a smaller slope according to (9) would lead to smaller σ_{eff}. At the same time even for hard diffraction the slope is $\approx 4.5 - 5GeV^{-2}$ as measured in the exclusive processes of J/ψ and ρ-meson production, one would obtain a value of σ_{eff} close to the lower limit of (9). Moreover a better fit is given in this case by $f(t) = F_N^2(t)$, where $F_N(t)$ is the nucleon dipole form factor. Using this fit one gets the values of σ_{eff} within the range of (9).

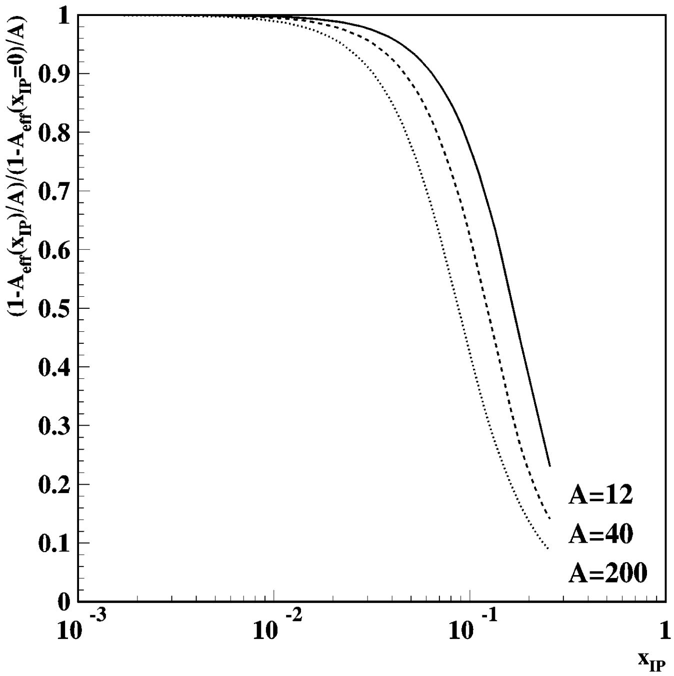

Fig. 1. Suppression due to the longitudinal momentum transfer calculated as a function of $x_{I\!P}$

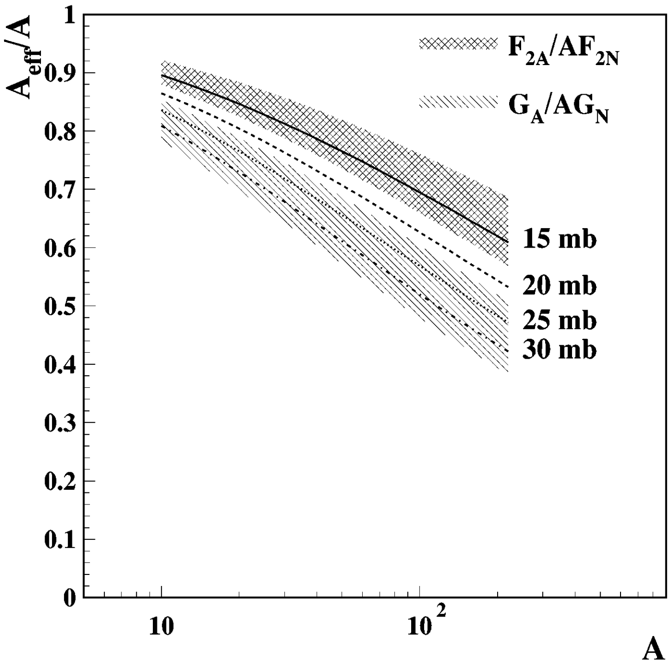

Fig. 2. Dependence of A_{eff}/A on σ_{eff} for $\lambda = 0$. The shaded regions are expected values of F_{2A}/AF_{2N} and G_A/AG_N for the range of σ_{eff} values consistent with the HERA diffractive data

over $x_{I\!P}$, i.e. over M^2-the mass of the diffractively produced system[4]:

$$R = 1 - \frac{1}{2} Re \left[\sigma_{eff}(1 - i\lambda) \int d^2b \int_{-\infty}^{\infty} dz_1 \int_{z_1}^{\infty} dz_2 \right. \\ \times \int_x^{x_0} dx_{I\!P} \eta(x_{I\!P}) \rho_A(b, z_1) \rho_A(b, z_2) \\ \left. \exp(-iq_{\parallel}(z_1 - z_2)) \exp(-\left(\frac{1}{2}\sigma_{eff} \int_{z_1}^{z_2} \rho_A(b, z') dz' \right) \right]. \tag{10}$$

Here $\rho_A(r)$ is the nucleon density in the nucleus normalized according to the equation $\int \rho_A(r) d^3r = A$. The quantity $\eta(x_{I\!P})$ is the differential diffractive cross section normalized to the total diffractive cross section

$$\eta(x_{I\!P}) = \frac{\frac{d\sigma_{diff}(x_{I\!P})}{dx_{I\!P}}}{\int_x^{x_0} \frac{\sigma_{diff}(x_{I\!P})}{dx_{I\!P}}} \tag{11}$$

The parameter $q_{\parallel} = x_{I\!P} m_N$ is the longitudinal momentum transfer to the nucleon in the $\gamma^* \to M_X$ transition. It is simply related to the coherence length parameter defined in 1 as $q_{\parallel} = \frac{1}{l_c}$.

The factor $\exp(-iq_{\parallel}(z_1 - z_2))$ in (10) suppresses the contribution of large $x_{I\!P}$. To illustrate this point we present in Fig. 1 the result of the calculation of the ratio $\frac{(1-R(x_{I\!P}))}{(1-R(x_{I\!P}=0))}$ for the value of $\sigma_{eff} = 28mb$, where instead of integrating over $x_{I\!P}$ we fix its value in 10. One can see from this plot that the suppression depends weakly on $x_{I\!P}$, which makes the neglect of the change of the mass in the rescattering quite safe.[5] One can see also that the contribution of the region of $x_{I\!P} \geq 0.05$ is strongly suppressed.

To explore the sensitivity of the result to the value of σ_{eff} we present in Fig. 2 the value of A_{eff}/A calculated for different values of σ_{eff}, and $\lambda = 0$ in the ranges relevant for shadowing in the gluon channel and in inclusive eA scattering. One can see that the estimated uncertainties in σ_{eff} of $\sim 23\%$ do not lead to a significant change of the shadowing effect. We will explore details of the Q^2 and x dependence of gluon shadowing elsewhere.

Note also that the cross section of the diffractive charm production rapidly increases with increase of energy. Hence (5) implies that the real part of the amplitude is not small. Based on the fits of [3] and (5) we estimate $\lambda \sim 0.4$. Hence we need to explore the sensitivity of A_{eff}/A to the value of λ. In Fig. 3 we present the result of calculation for A_{eff}/A for A=12,240 for a range of the values of λ and for the fixed value of $\sigma_{eff}(1 + \lambda^2)$ which is determined from the diffractive ratio of (6). One can see from the figure that for a realistic range of λ (shaded area) the sensitivity to real part is rather small, especially for heavy nuclei.

[4] In the discussed approximation $\langle \sigma^2 \rangle / \langle \sigma \rangle$ is fixed and further effects of dispersion are neglected. Hence it is close to the enhanced eikonal approximation used in [25]

[5] This approximation is similar to the one used in hadron-nucleus scattering in the model [25] which describes well available data on $\sigma_{tot}(hA)$ for the incident hadron energies $E^h_{inc} \leq 400GeV$ where inelastic shadowing effects play a noticeable role.

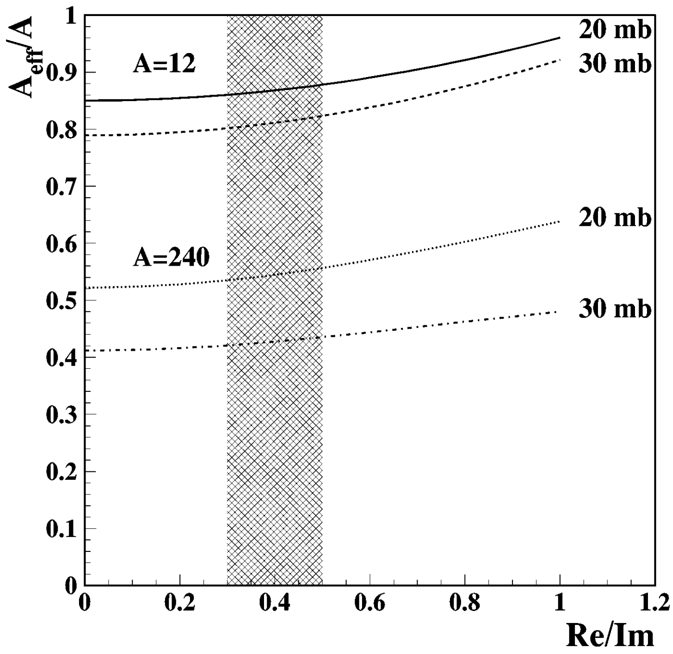

Fig. 3. Dependence of A_{eff}/A on $\lambda \equiv Re/Im$ for A=12 and A=240 and $\sigma_{eff}(1+\lambda^2)$ equal to 20 and 30mb. The shaded region indicates a realistic range of λ for charm diffraction

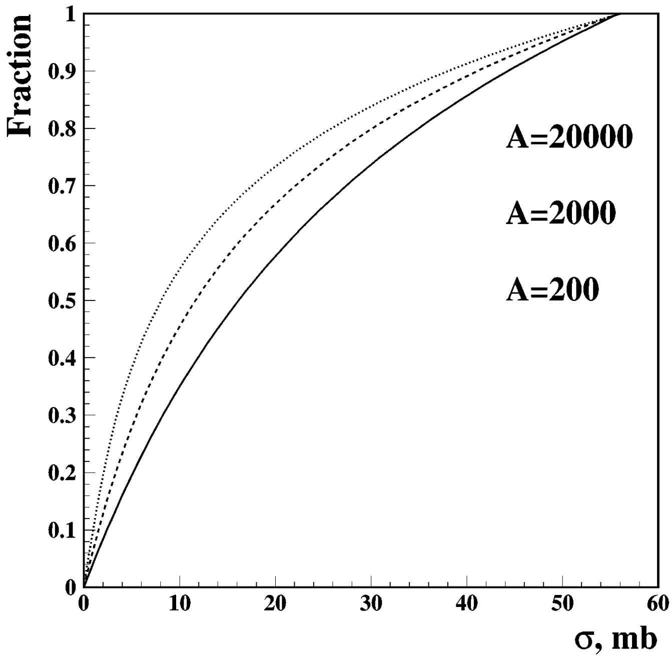

Fig. 4. Comparison of the quasieikonal based model and the fluctuation model for the value of shadowing for A=240, and $b=0$, as a function of σ_{eff}

3.3 Shadowing for central impact parameters

For the heavy ion applications it is interesting to consider shadowing of DIS processes for the central impact parameters. It is natural to define the parton density per unit area and to consider the suppression factor $R(b)$ for the value of this density as compared to that in the uncorrelated nucleons placed at the same impact parameter. The same equation 10 is valid for $R(b)$ with the integral over d^2b removed and an extra factor $\int dz\rho(b,z) \equiv T(b)$ in the denominator [26].

$$R(b) = 1 - \frac{A}{2T_A(b)} Re\Bigg[\sigma_{eff}(1-i\lambda) \int_{-\infty}^{\infty} dz_1 \int_{z_1}^{\infty} dz_2 dx_{I\!P} \int_x^{x_0} \eta(x_{I\!P}\rho_A(b,z_1)\rho_A(b,z_2) \exp(-iq_{\|}(z_1-z_2)) \exp(-\left(\frac{1}{2}\sigma_{eff}\int_{z_1}^{z_2}\rho_A(b,z')dz'\right)\Bigg]. \quad (12)$$

For large A and $b=0$ neglecting the edge effects (that is taking the nuclear density constant and equal to $\rho_0 \approx 0.16 fm^{-3}$ we obtain neglecting the $q_{\|}$ effects but accounting for the fluctuations of the interaction strengths as:

$$A_{eff}(b=0)/A = \frac{\int d\sigma P(\sigma)(2-2\exp(-\sigma L\rho_0/2))}{\rho_0 L \int d\sigma\sigma P(\sigma)}, \quad (13)$$

where $L \approx 2R_A$. If we neglect the dispersion in σ this would lead, for $\sigma_{eff}L\rho_0/2 \gg 1$, to

$$A_{eff}(b=0)/A \approx \frac{2}{\rho_0 L\sigma_{eff}} \quad (14)$$

For $A=208$ this results in $A_{eff}(b=0)/A = 0.32$ which is significantly smaller than the average A_{eff}/A for lead.

3.4 Fluctuations of the interaction strength

Let us investigate the sensitivity of this estimate to dispersion of the interaction strength. Motivated by the general structure of $P(\sigma)$ for a photon, where for small σ - $P(\sigma) \propto \frac{1}{\sigma}$ [27] (this behaviour is essentially consequence of the dimentional counting for $P(\sigma)$ for a point-like object), we take a model:

$$P(\sigma) = c\frac{1}{\sigma}\theta(\sigma_0 - \sigma) \quad (15)$$

The parameter $\sigma_0 = 2\sigma_{eff}$ is fixed based on (9). In this case (13) leads to

$$A_{eff}(b=0)/A = \frac{\int_0^{\sigma_0}(1-\exp(-\sigma\rho_0 L/2))/\sigma d\sigma}{\sigma_0\rho_0 L/4} \quad (16)$$

One can see from the analysis of (16) that for large values of the nuclear thickness the shadowing effect is smaller when the fluctuations are included due to the contribution of small σ's. However numerically even for the heaviest nuclei the effect is an increase of A_{eff}/A by a factor of 1.3, see Fig. 4. The increase occurs due to the contribution of configurations with $\sigma \leq \sigma_{eff}$. For very large A this would lead to a dominance of the contribution of configurations with small σ - a kind of a filtering phenomenon. To estimate at what A small σ would dominate we present in Fig. 5 the fraction of the cross section due to configurations with σ smaller than a given one. One can see that

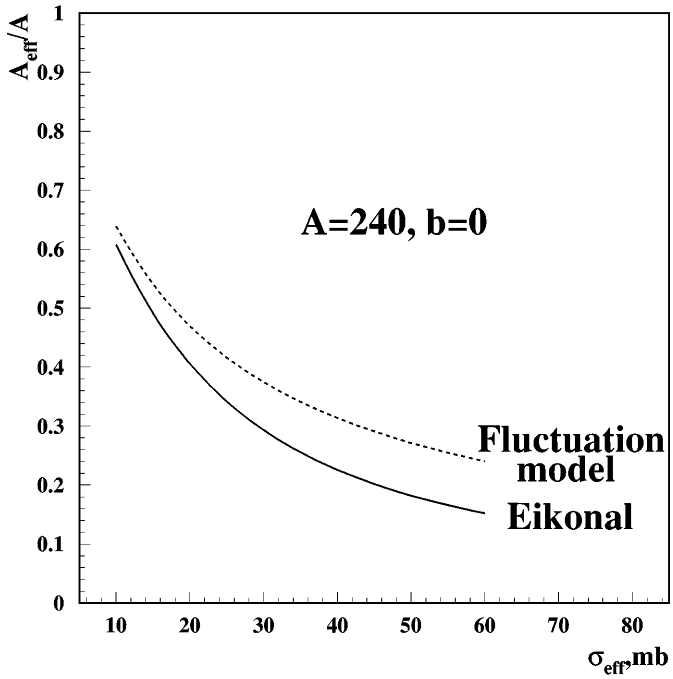

Fig. 5. Fraction of the total cross section due to the contributions of cross sections $\leq \sigma$ for $b = 0$ and A=200 (solid curve), A=2000 (dashed curve) and A=20000 (dotted curve)

for $\sigma_{eff} = 28$mb one would need extremely large A values to reach a situation where say $\sigma \leq 10$mb would give the dominant contribution. Due to the condition $l_{coh} \gg 2R_A$ this would require extremely small x.

The significant values of shadowing derived for small x and a wide range of Q^2 have obvious implications for AA collisions at LHC. For example, for production of J/Ψ and Υ at the central rapidities in Au-Au collisions we predict a suppression of the order of 1/5 ($\sim A^{-.3}$) for the inclusive cross section and $\sim 1/8(\sim A^{-.4})$ for the central collisions solely due to the gluon shadowing effects. Further suppression occurs due to the final state of the produced onium states $(\chi, J/\psi, ...)$. A similar suppression should occur for the minijet production for $x_T(jet) \leq 10^{-3}$.

3.5 Structure of the final states

Let us briefly discuss also the structure of the final states for the processes dominated by scattering off the gluons for $x \leq 10^{-3}$. In the case of inclusive scattering we have demonstrated in [7] using the AGK technique [28] that larger shadowing effects larger the rapidity gap fraction of events and the multiplicity fluctuations for production of soft particles for the central rapidity interval. Since in the case of the processes dominated by coupling to gluons σ_{eff} is a factor ~ 2 larger than for the inclusive case - application of the same technique indicates that for the case of charm production (i) the fraction of charm produced in the rapidity gap events for heavy nuclei should be in the range of 40-50%, (ii) the fraction of diffractive events with $\beta = Q^2/(Q^2 + M^2_{diff}) \geq 0.4$ should increase due to the enhancement of "elastic" scattering of the projectile configurations as compared to the "triple Pomeron type" processes [7], (iii) the distribution over the number of particles produced for a fixed rapidity interval around $y_{c.m.} \sim 0$ should be much broader than in ep collisions. Also, since the amplitude of hard diffractive process like $\gamma_L + A \to V(\rho, J/\psi, ...) + A$ is proportional to $x^2G^2_A(x, Q^2)$ for $t = 0$ – (the generalized color transparency [29]), the cross section of this process integrated over t should be $\sim A$, since $xG_A(x, Q^2)/A \propto A^{-0.15}$. Since $\sigma_{tot}(eA)/A \propto A^{-0.09}$ [7] we expect $\sigma(\gamma_L + A \to V(\rho, J/\psi, ...) + A)/\sigma_{tot}(eA) \propto A^{0.1}$.

In conclusion, we have demonstrated that the observation of the large diffractive cross section of charm production at HERA indicates that the shadowing for gluon densities at $x \leq 10^{-3}$ should be large and strongly affect the physics of nucleus-nucleus collisions at LHC. The observation of gluon shadowing of such magnitudes and associated effects in diffraction processes would be one of the green pastures for the eA collider program at HERA.

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