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Gluon shadowing in the Glauber–Gribov model at HERA

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Abstract

We calculate shadowing using new data on the gluon density of the pomeron recently measured with high precision at HERA. The calculations are made in a Glauber–Gribov framework and pomeron tree-diagrams are summed up within a unitarity-conserving procedure. The total cross section of $\gamma^* A$ interaction is then found in a parameter-free description, employing gluon diffractive and inclusive distribution functions as input. A strong shadowing effect is obtained, in a good agreement with several other models. Impact parameter dependence of gluon shadowing is also presented.

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1. Introduction

Nuclear shadowing is a well-established phenomenon, which attracts attention of both experimentalists and theoreticians. It was found (see [1,2] and references therein) that the inclusive nuclear structure function is smaller in nuclei than in a free nucleon at small values of the Bjorken variable $x \leqslant 0.01$. The nature of shadowing, emerging, e.g. in deep inelastic scattering (DIS), can be well understood in terms of multiparticle scattering in the target rest frame. The incoming photon is represented as a superposition of gluons, quarks, anti-quarks and their bound states. At high bombarding energies the photon converts into $q\bar{q}$-pair long before the target, and its hadronic component interacts coherently with several nucleons of the target nucleus. This process leads to absorption and, therefore, to nucleon shadowing (for a review see e.g. [3]). Among the important consequences of the phenomenon is, for instance, severe reduction of particle multiplicity in heavy-ion collisions at LHC energies ($\sqrt{s} = 5.5$ TeV), since multiple scattering is connected to diffraction [4–6]. The effect can be further decomposed onto the quark shadowing and the gluon shadowing; the latter provides largest uncertainties in the theory and is a subject of our present study.

In recent years a lot of interest has been generated about the possibility of parton saturation in the nuclear wave function at the smallest x accessible at HERA, and there is an ongoing discussion if the same effects can be observed for hadron–nucleus and nucleus–nucleus collisions at RHIC. In this Letter, we will focus mostly on the low-x effects mentioned above. Understanding the so-called cold nuclear effects in hadron–nucleus collisions serves also as a baseline for the correct treatment of possible final state effects in nucleus–nucleus collisions, e.g. high-p_T particle suppression and heavy-flavor production at RHIC.

Our starting point is noticing that a significant change in the underlying dynamics of a hadron–nucleus collision takes place with growing energy of the incoming particles. At low energies, the total cross section is well described within the probabilistic Glauber model [7], which only takes into account elastic rescatterings of the incident hadron on the various nucleons of

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the target nucleus. Elastic scattering is described by pomeron exchange. At higher energies, $E > E_{\mathrm{crit}} \sim m_N \mu R_A$ (μ is a characteristic hadronic scale, $\mu \sim 1$ GeV, and R_A is the radius of the nucleus) corresponding to a coherence length

$$l_C = \frac{1}{2m_N x}, \tag{1}$$

the typical hadronic fluctuation length can become of the order of, or even bigger than, the nuclear radius and there will be coherent interaction of constituents of the hadron with several nucleons of the nucleus. The sum of all diagrams was calculated by Gribov [8,9], which corrected the Glauber series by taking into account the diffractive intermediate states in the sum over subsequent rescatterings. The space–time picture analogy to the Glauber series is nevertheless lost, as the interactions with different nucleons of the nucleus occurs nearly simultaneous in time. The phenomenon of coherent multiple scattering is referred to as shadowing corrections.

An additional effect which comes into play at high energies, is the possibility of interactions between soft partons of the different nucleons in the nucleus. In the Glauber–Gribov model this corresponds to interactions between pomerons. These diagrams are called enhanced diagrams [10], and can also be understood as interactions between strings formed in the collision. Actually, the necessity to include such diagrams at high energies can be related to unitarization of the total cross section. There is a connection between these effects and saturation effects already mentioned earlier.

The Glauber–Gribov model is described and a unitarity-conserving procedure for finding the total cross section of γ^*-nucleus ($\gamma^* A$) interaction, which corresponds to summing up pomeron fan-diagrams, is presented in Section 2. Further we will concentrate on new and interesting data on gluon diffractive distribution function, which we will describe in Section 3. The results for gluon shadowing are presented in Section 4, and our conclusions are drawn in Section 5.

2. The model

We consider the nucleus as a set of nucleons, in the spirit of the Glauber model. The elastic $\gamma^* A$ scattering amplitude can then be written as the sum of diagrams shown in Fig. 1, i.e. as multiple γ^*-nucleon ($\gamma^* N$) scattering diagrams with pomeron exchange [4,11]. The contribution from $1, 2, \ldots$ scatterings

$$\sigma_{\gamma^* A} = A\sigma_{\gamma^* N} + \sigma^{(2)}_{\gamma^* A} + \cdots, \tag{2}$$

should be summed up to obtain the total cross section. In Eq. (2), the first term simply equals to the Glauber elastic contribution and subsequent terms describe multiple interactions of the incoming probe with the nucleons in the target nucleus.

The multiparticle content of subsequent diagrams in Fig. 1 is given by AGK cutting rules [12], where the intermediate states are on-shell. The cut contribution of the double rescattering diagram can be expressed in terms of diffractive deep inelastic scattering (DDIS). The usual variables for DDIS: Q^2, x, M^2 and t, or $x_{\mathbb{P}}$, are shown in Fig. 2. The variable $\beta = \frac{Q^2}{Q^2+M^2} = x/x_{\mathbb{P}}$ plays the same role for the pomeron as the Bjorken variable, x, for the nucleon. We assume that the amplitude of the process is purely imaginary; this is justified for a value of the pomeron intercept close to unity [13]. The contribution from the second term in Eq. (2) to the total $\gamma^* A$ cross section is given by [4]

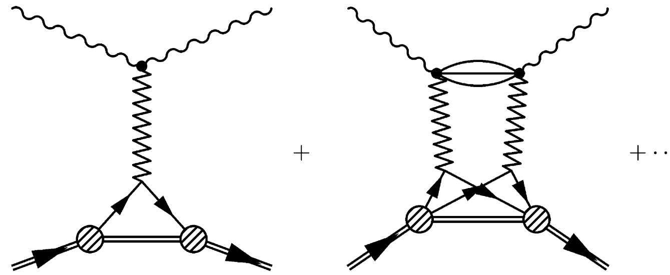

Fig. 1. The single and double scattering contribution to the total $\gamma^* N$ cross section.

$$\sigma^{(2)}_{\gamma^* A} = -4\pi A(A-1) \int \mathrm{d}^2 b\, T_A^2(b) \times \int_{M^2_{\min}}^{M^2_{\max}} \mathrm{d}M^2 \left[\frac{\mathrm{d}\sigma^{\mathcal{D}}_{\gamma^* N}(Q^2, x_{\mathbb{P}}, \beta)}{\mathrm{d}M^2\,\mathrm{d}t}\right]_{t=0} F_A^2(t_{\min}), \tag{3}$$

where $T_A(b) = \int_{-\infty}^{+\infty} \mathrm{d}z\, \rho_A(\boldsymbol{b}, z)$ is the nuclear normalized density profile, $\int \mathrm{d}^2 b\, T_A(b) = 1$. The form factor F_A is given by

$$F_A(t_{\min}) = \int \mathrm{d}^2 b\, J_0\big(\sqrt{-t_{\min}}\, b\big) T_A(b), \tag{4}$$

where $t_{\min} = -m_N^2 x_{\mathbb{P}}^2$, and $J_0(x)$ denotes the Bessel function of the first kind. Strictly speaking, Eq. (3) is valid for nuclear densities which depend separately on **b** and z, however we have checked that calculations with an exact expression lead to negligible corrections. Note that since Eq. (3) is obtained under very general assumption, i.e. analyticity and unitarity, it can be applied for arbitrary values of Q^2 provided x is very small. We have assumed $R_A^2 \gg R_N^2$, so that the t-dependence of the $\gamma^* N$ cross section has been neglected. For a deuteron, the double rescattering contribution has the following form

$$\sigma^{(2)}_{\gamma^* d} = -2 \int_{-\infty}^{t_{\min}} \mathrm{d}t \int_{M^2_{\min}}^{M^2_{\max}} \mathrm{d}M^2 \frac{\mathrm{d}\sigma^{\mathcal{D}}_{\gamma^* N}}{\mathrm{d}M^2\,\mathrm{d}t} F_D(t), \tag{5}$$

where $F_D(t) = \exp(at)$, with $a = 40\ \mathrm{GeV}^{-2}$ [4].

In Eqs. (3) and (5), $M^2_{\min}$ corresponds to the minimal mass of the diffractively produced hadronic system, $M^2_{\min} = 4m_\pi^2 = 0.08\ \mathrm{GeV}^2$, and $M^2_{\max}$ is chosen according to the condition: $x_{\mathbb{P}} \leqslant x_{\mathbb{P}}^{\max}$. The choice of $x_{\mathbb{P}}^{\max}$ is governed by the fact that the model is only valid for $x_{\mathbb{P}} \ll 1$, i.e. a large rapidity gap is required in the experimental data. We use the standard choice for $x_{\mathbb{P}}^{\max} = 0.1$ [14]. It is convenient as it guarantees the disappearance of nuclear shadowing at $x \sim 0.1$ as in experimental data. Coherence effects are taken into account through $F_A(t_{\min})$ in Eq. (4), which is equal to 1 at $x \to 0$ and decreases with increasing x due to the loss of coherence for $x > x_{\mathrm{crit}} \sim (m_N R_A)^{-1}$.