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Mutual boosting of the saturation scales in colliding nuclei

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ABSTRACT

Saturation of small-x gluons in a nucleus, which has the form of transverse momentum broadening of projectile gluons in pA collisions in the nuclear rest frame, leads to a modification of the parton distribution functions in the beam compared with pp collisions. The DGLAP driven gluon distribution turns out to be suppressed at large x, but significantly enhanced at $x \ll 1$. This is a high twist effect. In the case of nucleus–nucleus collisions all participating nucleons on both sides get enriched in gluon density at small x, which leads to a further boosting of the saturation scale. We derive reciprocity equations for the saturation scales corresponding to a collision of two nuclei. The solution of these equations for central collisions of two heavy nuclei demonstrate a significant, up to several times, enhancement of Q^2_{sA}, in AA compared with pA collisions.

1. Introduction

The transverse momentum distribution of gluons in nuclei is known to be modified compared with a free nucleon. The mean transverse momentum squared increases up to a value called saturated scale, Q^2_{sA}, which depends on the nuclear profile. This phenomenon, called color glass condensate [1], is related to parton saturation at small x [2], and can be also understood in terms of the Landau–Pomeranchuk principle [3] as a consequence of coherent gluon radiation from multiple interactions in the nucleus [4]. The value of the saturation momentum was calculated and compared with data on broadening in [4] and has been modeled recently in [5–12].

The saturation scale can be measured as p_T-broadening of a parton propagating through the nucleus in its rest frame [4],

$$Q^2_{sA} = \Delta p^2_T. \tag{1}$$

Although in leading order both sides of this relation rise linearly with nuclear profile T_A [13–17], this dependence slows down by gluon shadowing. These phenomena, broadening and suppression of gluons, are closely related, since both result from coherence of gluon radiation in multiple interactions. Solving the corresponding equation derived in [4], one arrives at a saturation scale considerably reduced compared to the leading order. The T_A dependence is slower than linear, and at very large (unrealistic) nuclear thicknesses the saturation scale saturates, becoming independent of T_A. Notice that the solution found in [4] is similar to the result of numerical solution [22] of the Balitsky–Kovchegov equation [18,19]. Broadening of gluons radiated in heavy ion collisions was studied with numerical simulations in [20,21].

2. Modification of the beam PDF by a nuclear target

Due to broadening a nuclear target probes the parton distribution in the beam hadron with a higher resolution. Therefore, the effective scale Q^2 for the beam PDF drifts to a higher value $Q^2 + Q^2_{sA}$. At first glance this seems to contradict causality, indeed, how can the primordial parton distribution in the hadron depend on the interaction which happens later? However, there is nothing wrong. The interaction performs a special selection of Fock states in the incoming hadron. The same phenomenon happens when one is measuring the proton parton distribution in DIS. The proton PDF “knows” in advance about the virtuality of the photon which it is going to interact with.

The shift in the scale also can be interpreted as a manifestation of the Landau–Pomeranchuk principle [3]: at long coherence times gluon radiation (which causes the DGLAP evolution) does not depend on the details of multiple interactions, but correlates only with the total momentum transfer, $\vec{q} + \Delta\vec{p}_T$, which after squaring and averaging over angles results in $Q^2 + \Delta p^2_T$.

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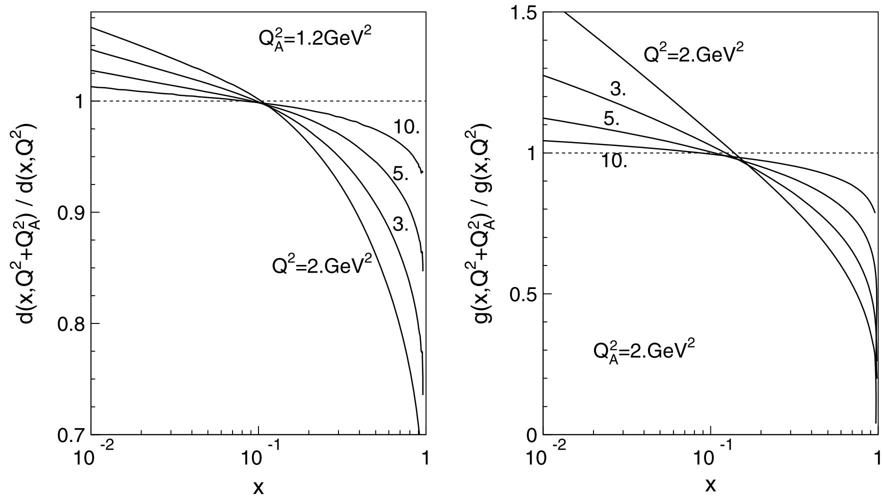

Fig. 1. Ratio of parton distribution functions in a reaction characterized by a hard scale $Q^2 = 2$, 3, 5, 10 GeV^2 on a nuclear ($A = 200$) and proton targets. *Left panel*: ratio of the d-quark distributions for the quark saturation momentum $Q^2_{sA} = 1.2$ GeV^2. *Right panel*: ratio for gluons with the gluon saturation momentum $Q^2_{sA} = 2$ GeV^2.

As far as the PDF of the projectile proton has a harder scale in pA collisions than in pp, the ratio of parton distributions should fall below one at forward and rise above one at backward rapidities. This may look like a breakdown of k_T-factorization, however, it is a higher twist effect.

Examples of pA to pp ratios $R_A(x, Q^2)$ calculated with MSTW2008 [23] are shown in Fig. 1 for d-quark and gluon distributions in a hard reaction (high-p_T, heavy flavor production, etc.). We see that the shift in the hard scale caused by saturation in the nucleus leads to a sizable suppression in the projectile parton distribution at large $x \to 1$ and enhancement at small $x \ll 1$. We also observe that the magnitude of nuclear modification quickly decreases with Q^2 confirming that this is a high twist effect.

Important for what follows is the observation of a considerably increased population of small-x partons in the projectile proton in pA compared with pp collisions.

3. Nucleus–nucleus collisions: reciprocity relations

Notice that in pA collisions the modification of the PDFs of the beam and target are not symmetric. Namely, the scale of the PDF of the beam proton gets a shift, $Q^2 \Rightarrow Q^2_{eff} = Q^2 + Q^2_{sA}$, while the PDFs of the bound nucleons, which do not undergo multiple interactions, remain the same as in pp collisions.

The situation changes in the case of a nucleus–nucleus collision: the bound nucleons in both nuclei participate in multiple interactions, therefore the scales of PDFs of all of them are modified. However, this modification goes beyond the simple shift $Q^2 \Rightarrow Q^2 + Q^2_{sA}$. Indeed, in an AB nuclear collision not only the two nucleons (one from A and one from B) participating in the hard reaction undergo multiple interactions, but also many other nucleons, the so-called participants, experience multiple soft interactions. For this reason their parton distributions are boosted from the soft scale μ^2 up to the saturation scale $\mu^2 \Rightarrow \mu^2 + Q^2_{sA(B)}$, which is usually much larger. Thus, the participant nucleons on both sides are boosted to a higher scale and get softer PDFs, with larger parton multiplicities at small x. This is illustrated on the cartoon in Fig. 2.

The next important observation is that the p_T-broadening on such "excited", or boosted nucleons, $\tilde{N}$ is larger than in pA collisions, $\Delta p^2_T|_{\tilde{N}} > \Delta p^2_T|_N$, since the density of target gluons is increased at small x. This should lead to a further mutual enhancement of broadening, i.e. a further increase of the saturation scales

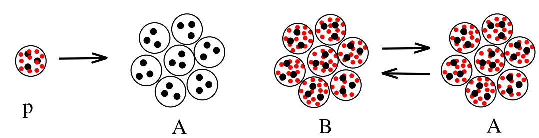

Fig. 2. *Left*: pA collision in which the colliding proton is excited by multiple interactions up to a saturated scale Q^2_{sA}, what leads to an increased multiplicity of soft gluons in the incoming proton. *Right*: nuclear collision in which participating nucleons on both sides are boosted to the saturation scales, Q^2_{sA} in the nucleus B, and Q^2_{sB} in the nucleus A. As a result, the low-x gluon population is enriched in both nuclei.

in both nuclei. Intuitively this seems to be clear, but a formal consideration below also supports this conclusion.

Broadening is predominantly a process based on many soft rescatterings of the projectile parton. It was found in [15,14] (see also [4]) that quark broadening is related to the dipole cross section,

$$Q^2_{sA} = \Delta p^2_T(E) = 2\frac{d\sigma(r,E)}{dr^2}\bigg|_{r=0} \int dz\, \rho_A(b,z), \tag{2}$$

where $\rho_A(b,z)$ is the nuclear density at impact parameter b and longitudinal coordinate z. Since the process is soft, Bjorken x is not a proper variable, but instead the parton energy E should be used. The energy dependent $\bar{q}q$ dipole cross section was parametrized in the saturated form and fitted to photoabsorption, low Q^2 DIS data, and πp total cross section in [28] (see also [15,24]). With that parametrization [15]

$$C_q(E) \equiv \frac{d\sigma(r,E)}{dr^2}\bigg|_{r=0} = \frac{1}{4}\sigma^{\pi p}_{tot}(E)\left[Q^2_{qN}(E) + \frac{3}{2\langle r^2_{ch}\rangle_\pi}\right], \tag{3}$$

where $\sigma^{\pi p}_{tot}(E)$ is the πp total cross section; $\langle r^2_{ch}\rangle_\pi \approx 0.44$ fm^2 is the mean pion charge radius squared; $Q_{qN}(E) = 0.19$ GeV $\times$ $(E/1\ \text{GeV})^{0.14}$ is the proton saturation momentum.

Notice, that Eq. (3) $C_q(E)$ has no scale dependence. It corresponds to the dipole–nucleon cross section $\sigma_{dip}(r) = C_q r^2$ and a soft scale characterizing the proton is implicitly contained in C_q. Strictly speaking, however, this coefficient is divergent at $r \to 0$, since it contains $\ln(r_T)$ [27]. This divergency originates from the ultraviolet behavior of the unintegrated gluon density $\mathcal{F}(x,k_T) \propto 1/k^4_T$ at large k_T. In reality this divergency is not harmful due to the natural cut-offs discussed in [15], and to a low sensitivity to their values. One should fix the $\ln r$ dependent factor term in C_q at