MANIFESTATION OF THE COLOR GLASS CONDENSATE IN PARTICLE PRODUCTION AT RHIC*

Kirill Tuchin

Physics Department, Brookhaven National Laboratory
Upton, NY 11973-5000, USA

(Received October 4, 2004)

We discuss general properties of the Color Glass Condensate. We show that predictions for particle production in $p(d)A$ and AA collisions derived from these properties are in agreement with data collected at RHIC.

PACS numbers: 13.85.Hd, 13.60.Hb, 13.85.Ni, 24.85.+p

In this paper we discuss the experimental signatures of the new form of nuclear matter — the Color Glass Condensate (CGC) in particle production at RHIC. Let us first see how the notion of the CGC arises in pA collisions at high energy. Consider a process of inclusive particle production in pA collisions in a nucleus rest frame. At high energies the typical values of the Bjorken x are small. It is well-known that at small x hard processes develop over large 'coherence length' $l_{\rm c}$. In particular, a gluon production is coherent over $l_{\rm c} \simeq 1/(Mx)$, where M is the proton's mass. For instance, at midrapidity at RHIC the coherence length of 2 GeV gluon is $l_{\rm c} \simeq 20$ fm (at $\sqrt{s} = 200$ GeV). It is much bigger than the size of the target $\simeq 6.5$ fm. This allows formal separation of the gluon production process into two parts: slow gluon emission described by the proton's light cone wave function, and almost instantaneous interaction with the target at given impact parameter $\underline{b}$ described by the amplitude $N_{\rm G}(\underline{r}, \underline{b}, x)$, where $\underline{r}$ is the variable Fourier-conjugated to the gluon's transverse momentum $\underline{k}$. In the one gluon exchange approximation, assuming that scattering on different nucleons is independent, one arrives at the formula [1–3]

$$N_{\rm G}(\underline{r}, \underline{b}, x) = 1 - \exp\left[-\frac{1}{4}\underline{r}^2 Q_{\rm s}^2 S(\underline{b}) \ln(1/r\mu)\right], \qquad (1)$$

* Presented at the XXXIV International Symposium on Multiparticle Dynamics, Sonoma County, California, USA, July 26–August 1, 2004.

similar to the Glauber formula for the low energy hadron-nucleus scattering. Here $Q_{\rm s}$ is a parameter with dimension of mass, $S(\underline{b})$ is a nuclear profile function and $r \equiv |\underline{r}|$. $N_{\rm G}(\underline{r},\underline{b},x)$ can be interpreted as a forward scattering amplitude of a gluon dipole off a heavy nucleus. One can see from (1) that for hard gluons, such that $1/r \sim k \gg Q_{\rm s}$ the scattering amplitude coincides with the usual perturbative expression

$$N_{\rm G}^{\rm pert}(\underline{r},\underline{b},x) = \underline{r}^2\,\pi^2\,\alpha_{\rm s}\,\rho\,T(\underline{b})\,xG(r,x)/(2\,C_{\rm F})\,. \tag{2}$$

In the opposite limit $k \ll Q_{\rm s}$ the scattering amplitude (1) becomes independent of its variables $\underline{r}$, $\underline{b}$ and x as it approaches its unitarity limit. This is the phenomenon of *saturation* in nuclear and hadronic reactions [4]. The scale $Q_{\rm s}$ is called *the saturation scale*. Eq. (2) implies that $Q_{\rm s}^2 \propto A^{1/3}$, where A is the atomic number. Thus, for a very big nucleus $A \gg 1$ the saturation scale becomes a perturbative scale $Q_{\rm s} \gg \Lambda_{\rm QCD}$ which in turn means that $\alpha_{\rm s}(Q_{\rm s}^2) \ll 1$. It can be argued that the total multiplicity of produced gluons is dominated by gluons with the typical momentum $k \simeq Q_{\rm s}$. Therefore, gluon production in pA collisions at high energies is calculable in the perturbation theory.

Note, that the scattering amplitude (1) was calculated in a quasi-classical approximation. The quasi-classical approximation corresponds to a quantum system with high occupation numbers. In terms of the QCD action $S_{\rm QCD}$ this implies $S_{\rm QCD} \gg 1$. At small x gluons dominate over quarks. Therefore, we have

$$\frac{1}{g^2}\int d^4x\,{\rm tr}\,\tilde{G}_{\mu\nu}(x)\,\tilde{G}^{\mu\nu}(x) \gg 1\,, \tag{3}$$

where the rescaled gluon field is defined as $\tilde{A}^a_\mu(x) = g\,A^a_\mu(x)$. On the other hand, $\alpha_{\rm s}(Q_{\rm s}^2) \ll 1$. Therefore, the typical gluon field of nucleus is of order of $A^a_\mu \sim 1/g$ [1]. This strong gluon field at small coupling is called *the Color Glass Condensate*. This configuration is very much different from the perturbation theory where both the gluon (and quark) field and the coupling are small, and from the non-perturbative regime where both gluon (and quark) field and the coupling are large [5]. Phenomenologically, the CGC in a quasi-classical approximation manifests itself as a saturation of the scattering amplitude $N_{\rm G}(\underline{r},\underline{b},x)$ at small transverse momenta [4].

The CGC in a quasi-classical approximation can be thought of as a model of multiple rescatterings of a hadron in a heavy nucleus at high energy. As such it has much in common with many other models of multiple rescatterings. In particular, their common prediction is the Cronin effect, *i.e.* enhancement of particle production at intermediate transverse momenta k in pA collisions as compared with pp scaled by the atomic number A. The

origin of this effect is simple: a gluon traversing a heavy nucleus gains additional transverse momentum due to multiple rescatterings. On the other hand, in a quasi-classical approximation the total number of particles is conserved. Therefore, if there are less particles with low transverse momentum, then there are more particles with high transverse momentum [6–9]. Of course, this effect is predicted to increase for heavier nuclei or more central collisions.

However, a quasi-classical approximation breaks down at high energy since quantum evolution becomes an important process. Indeed, additional gluon production is parametrically of order $\alpha_{\rm s}\ln(1/x)$. Therefore, at $x \ll e^{1/\alpha_{\rm s}}$ a quasi-classical approximation is no longer valid. One might attempt to take the evolution into account using collinear factorization, which basically means incoherent production of gluons in the proton's wave function. As a result, proton will suffer more scatterings in a nucleus and the Cronin effect will increase with energy/rapidity. However, this expectation contradicts the experimental data as we discuss later, see Fig. 2. The reason is that the collinear factorization scheme and the OPE break down as soon as multiple rescatterings are important, see *e.g.* [11]. This is because each additional scattering is a higher-twist effect. The effect of coherence of the parton evolution at high energies can be taken into account in the *nonlinear evolution equations* of QCD [4, 12, 13]. These equations describe the high energy quantum evolution of the CGC. That is, if the scattering amplitude $N_{\rm G}(\underline{r}, \underline{b}, x)$ is known at some initial value of x_0, *e.g.* as given by (1), the evolution equations allow calculation of the scattering amplitude at any $x < x_0$.

In the large N_c approximation the differential cross section for a gluon production can be written in the $k_{\rm T}$-factorized form [3, 14, 15]

$$\frac{d\sigma_{\rm G}^{pA}}{d^2k\ dy} = \frac{C_{\rm F}\, S_A\, S_d}{\alpha_{\rm s}\,\pi\,(2\pi)^3}\,\frac{1}{\underline{k}^2}\int d^2r\,\nabla_z^2\, n_{\rm G}(\underline{z}, Y-y)\, e^{-i\underline{k}\cdot\underline{r}}\,\nabla_r^2\, N_{\rm G}(\underline{r}, y), \quad (4)$$

where $y = \ln(1/x)$ and $n_{\rm G}(\underline{r}, \underline{b}, y)$ is a forward gluon dipole scattering amplitude off a proton. S_A and S_p are cross sectional areas of the gold nucleus and proton correspondingly and Y is the total rapidity interval. The evolution effects in a nucleus are enhanced by a factor of $A^{1/3} \gg 1$ as compared to those in proton (deuteron). Therefore, $n_{\rm G}(\underline{r}, \underline{b}, y)$ approximately satisfies the linear BFKL evolution equation [16] (this is correct at not very high energies, when the Pomeron loops are small). The gluon dipole scattering amplitude can be related to the unintegrated gluon distribution function $\phi(\underline{k}, x)$ as [14]

$$\phi(\underline{k}, x) = \frac{C_{\rm F}}{\alpha_{\rm s}(2\pi)^3}\int d^2b\, d^2r\, e^{-i\underline{k}\cdot\underline{r}}\,\nabla_r^2\, N_{\rm G}(\underline{r}, \underline{b}, x)\,. \quad (5)$$

The main property of $\phi(\underline{k}, x)$ which follows directly from the nonlinear evolution equations is the *geometric scaling* which means that $\phi(\underline{k}, x) = \phi(k/Q_{\rm s}(x))$, *i.e.* the gluon distribution becomes a function of only one variable at low x [4]! Here $Q_{\rm s}(x)$ is the same saturation scale as in (1). Being the only dimensional parameter at low x, $Q_{\rm s}(x)$ sets the scale for the gluon field. Eq. (3) implies $A^a_\mu \sim Q_{\rm s}(x)/g$. In course of evolution A^a_μ increases as x decreases due to increase of number of color sources. Hence, $Q_{\rm s}(x)$ is increasing function of $1/x$. It follows from the nonlinear evolution equation that [17]:

$$Q_{\rm s}^2(x) = \left(\frac{x_0}{x}\right)^\lambda A^{1/3}\,{\rm GeV}^2, \tag{6}$$

where $\lambda \approx 0.3$ [18]. The same gluon distribution function $\phi(\underline{k}, x)$ (5) enters expressions for the structure functions in Deep Inelastic Scattering. It allows to fit the initial value of x using experimental data collected at HERA. In (6): $x_0 = 3 \times 10^{-4}$ and $\lambda = 0.28$. For RHIC and LHC it is convenient to write (6) in the center-of-mass frame

$$Q_{\rm s}^2 = \left(\frac{\sqrt{s}}{3.3\,{\rm TeV}}\right)^\lambda e^{\pm\lambda y}\, A^{1/3}\,{\rm GeV}^2. \tag{7}$$

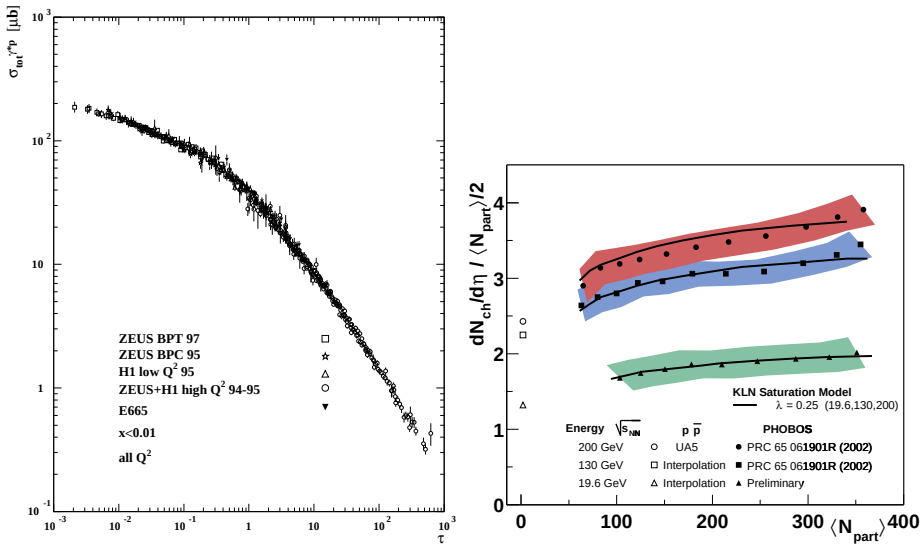
The geometric scaling of the gluon distribution holds as long as the logarithms of energy gained in course of the BFKL evolution are bigger than the logarithms of transverse momentum gained in course of the DGLAP evolution:

$$\alpha_{\rm s} \log Q_{\rm s}^2/\,\Lambda^2 \;\sim\; \alpha_{\rm s}\, y \;\gg\; \alpha_{\rm s}\, \log \underline{k}^2/Q_{\rm s}^2, \tag{8}$$

which implies the geometric scaling in a wide kinematical region $k < k_{\rm geom} = Q_{\rm s}^2/\Lambda$ [19]. The experimental evidences of the geometric scaling in DIS and heavy ion collisions are shown in Fig. 1.

Another consequence of the quantum evolution on the unintegrated gluon distribution (5) is that its anomalous dimension γ acquires strong dependence on the scaling variable $k/Q_{\rm s}$. In the perturbative regime $k \gg k_{\rm geom}$ we get the usual leading-twist expression $\phi(\underline{q}, x) \propto S_A\, Q_{\rm s}^2/\underline{q}^2$ modulo DGLAP corrections. However, in the saturation $k < Q_{\rm s}$, $\phi(\underline{q}, x) \propto S_A$, *i.e.* $\gamma \to 0$. As we have already noted this signals the breakdown of the OPE. In the intermediate region $Q_{\rm s} < k < k_{\rm geom}$ the saddle point of the BFKL amplitude is located at $\gamma \approx 1/2$, which implies $\phi(\underline{q}, x) \propto S_A Q_{\rm s}/q$. Recalling, that $Q_{\rm s} \propto A^{1/3}$ we find that at $k < k_{\rm geom}$ the gluon distribution in a nucleus of atomic number A is less than A times the gluon distribution in a proton:

$$\frac{\phi_A(k, x)}{A\,\phi_p(k, x)} = \frac{1}{A^\rho}\,, \tag{9}$$



(a) (b)

Fig. 1. (a) Geometric scaling of the total DIS cross section $\sigma(Q^2, x)$ as a function of $\tau = Q^2/Q_{\rm s}^2$ [20]. (b) Geometric scaling of the total charged hadron multiplicity in Au–Au collisions at RHIC. Solid lines: prediction of the saturation model of [5]. Prediction based on the collinear factorization is quite different $dN/d\eta\, N_{\rm part}^{-1} \propto N_{\rm part}^{1/3}$.

with $\rho = 1/3$ at $k \ll Q_{\rm s}$ and $\rho \approx 1/6$ at $Q_{\rm s} < k < k_{\rm geom}$. Let me emphasize that the CGC takes into account two nuclear shadowing effects. First one is a quasi-classical effect of multiple rescatterings. It necessarily requires higher twist effects to be included in a calculation. It predicts suppression at $k < Q_{\rm s}$ followed by enhancement at $k \sim Q_{\rm s}$. Second one is a quantum evolution effect. It predicts suppression in wide kinematical region $k < k_{\rm geom}$ both in the linear evolution region at $k > Q_{\rm s}$ ('leading twist shadowing') and the nonlinear evolution one at $k < Q_{\rm s}$ (saturation). Using (4) and (5) it is easily seen that nuclear shadowing in $\phi_A(k, x)$ translates into suppression of particle production in deuteron-gold collisions [6–10]. In Fig. 2 recent RHIC data for charged particle production at different rapidities and centralities [21] is shown. We see that at $\eta = 0$ there is the Cronin enhancement of the particle production in dA as compared to pp in central collisions as predicted by the CGC as well as by many multiple rescatterings models [22]. This implies that quasi-classical approximation is valid. At pseudo-rapidity $\eta = 3.2$ corresponding to 25 times smaller x's than at $\eta = 0$ for the same transverse momentum, the evolution becomes essential. It manifests itself as suppression of particle production at large transverse momenta in $p(d)A$ as compared to pp at higher rapidities and centralities. Not only that none

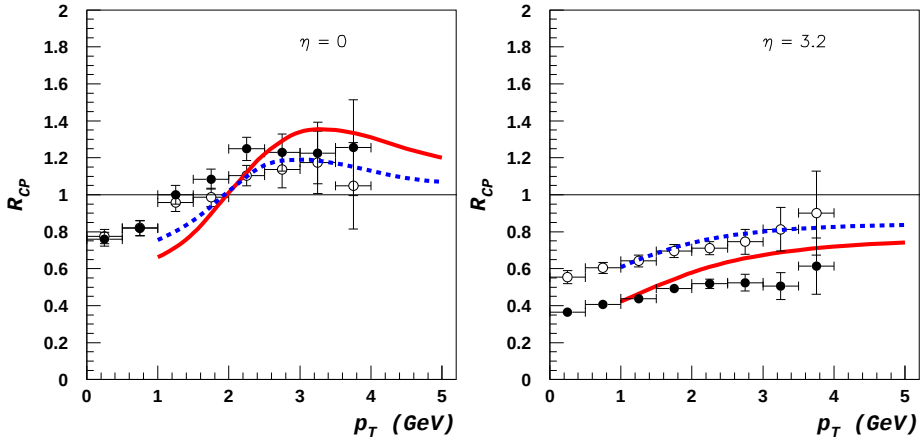


Fig. 2. The nuclear modification factor: ratio of the charged hadron multiplicity in central (full dotes, $b \approx 3$ fm) and semi-central (open dots, $b \approx 5$ fm) dA collisions to those in peripheral ones ($b \approx 7$ fm) rescaled by the ratio of corresponding numbers of participated nucleons. Lines: result of a simple CGC model of [6].

of the existing conventional nuclear shadowing models can explain *both* the Cronin effect and the suppression in the deuteron fragmentation region, but also none of those shadowing models can explain the large value of that suppression [23]. In the framework of CGC both effects are *predicted* to follow from the nonlinear evolution equation [6–8, 10]. The value of suppression factor comes naturally as a consequence of (9).

Since the suppression of charged particles in $p(d)A$ and AA at forward rapidities at RHIC originates in a gluon shadowing (9) there should be similar suppression in open charm production. In that case the region of the geometric scaling, and hence of suppression, is $m_t < k_{\rm geom}$, where $m_t^2 = \underline{k}^2 + m_c^2$ [24]. Another important signature of CGC is weakening of jet–jet correlations [25]. Indeed, since at low x a lot of particles with a typical momentum $Q_{\rm s}$ can be produced in single nucleon–nucleon subcollision any two of them need not to be correlated back-to-back unless their transverse momenta are very large. Once all these pieces of evidence all collected together they will become a strong evidence for the Color Glass Condensate at RHIC.

REFERENCES

[1] L.D. McLerran, R. Venugopalan, *Phys. Rev.* **D49**, 2233 (1994); *Phys. Rev.* **D49**, 3352 (1994); *Phys. Rev.* **D50**, 2225 (1994).

[2] Y.V. Kovchegov, *Phys. Rev.* **D54**, 5463 (1996); *Phys. Rev.* **D55**, 5445 (1997); J. Jalilian-Marian, A. Kovner, L.D. McLerran, H. Weigert, *Phys. Rev.* **D55**, 5414 (1997).

[3] Yu.V. Kovchegov, A.H. Mueller, *Nucl. Phys.* **B529**, 451 (1998).

[4] L.V. Gribov, E.M. Levin, M.G. Ryskin, *Phys. Rep.* **100**, 1 (1983).

[5] D. Kharzeev, E. Levin, *Phys. Lett.* **B523**, 79 (2001).

[6] D. Kharzeev, Y.V. Kovchegov, K. Tuchin, *Phys. Rev.* **D68**, 094013 (2003).

[7] J.L. Albacete, N. Armesto, A. Kovner, C.A. Salgado, U.A. Wiedemann, *Phys. Rev. Lett.* **92**, 082001 (2004).

[8] R. Baier, A. Kovner, U.A. Wiedemann, *Phys. Rev.* **D68**, 054009 (2003).

[9] E. Iancu, K. Itakura, D.N. Triantafyllopoulos, *Nucl. Phys.* **A742**, 182 (2004).

[10] D. Kharzeev, E. Levin, L. McLerran, *Phys. Lett.* **B561**, 93 (2003).

[11] E. Gotsman, E. Levin, U. Maor, L.D. McLerran, K. Tuchin, *Phys. Lett.* **B506**, 289 (2001); *Nucl. Phys.* **A683**, 383 (2001).

[12] I. Balitsky, *Nucl. Phys.* **B463**, 99 (1996); Y.V. Kovchegov, *Phys. Rev.* **D60**, 034008 (1999).

[13] J. Jalilian-Marian, A. Kovner, A. Leonidov, H. Weigert, *Nucl. Phys.* **B504**, 415 (1997); *Phys. Rev.* **D59**, 014014 (1999); E. Iancu, A. Leonidov, L. D. McLerran, *Nucl. Phys.* **A692**, 583 (2001); E. Ferreiro, E. Iancu, A. Leonidov, L. McLerran, *Nucl. Phys.* **A703**, 489 (2002);

[14] Yu.V. Kovchegov, K. Tuchin, *Phys. Rev.* **D65**, 074026 (2002).

[15] M.A. Braun, *Phys. Lett.* **B483**, 105 (2000).

[16] E.A. Kuraev, L.N. Lipatov, V.S. Fadin, *Sov. Phys. JETP* **45**, 199 (1977) [Zh. Eksp. Teor. Fiz. **72**, 377 (1977)]; I.I. Balitsky, L.N. Lipatov, *Sov. J. Nucl. Phys.* **28**, 822 (1978) [Yad. Fiz. **28**, 1597 (1978)].

[17] E. Levin, K. Tuchin, *Nucl. Phys.* **B573**, 833 (2000).

[18] D.N. Triantafyllopoulos, *Nucl. Phys.* **B648**, 293 (2003).

[19] E. Iancu, K. Itakura, L. McLerran, *Nucl. Phys.* **A708**, 327 (2002); E. Levin, K. Tuchin, *Nucl. Phys.* **A691**, 779 (2001).

[20] A. M. Stasto, K. Golec-Biernat, J. Kwiecinski, *Phys. Rev. Lett.* **86**, 596 (2001).

[21] I. Arsene *et al.* [BRAHMS Collaboration], `nucl-ex/0403005`.

[22] A. Accardi, M. Gyulassy, *Phys. Lett.* **B586**, 244 (2004); X.N. Wang, *Phys. Rev.* **C61** (2000) 064910; I. Vitev, *Phys. Lett.* **B562**, 36 (2003); B.Z. Kopeliovich, A.V. Tarasov, A. Schafer, *Phys. Rev.* **C59**, 1609 (1999).

[23] V. Guzey, M. Strikman, W. Vogelsang, `hep-ph/0407201`.

[24] D. Kharzeev, K. Tuchin, *Nucl. Phys.* **A735**, 248 (2004); K. Tuchin, `hep-ph/0401022`; J.P. Blaizot, F. Gelis, R. Venugopalan, *Nucl. Phys.* **A743**, 57 (2004).

[25] Y.V. Kovchegov, K.L. Tuchin, *Nucl. Phys.* **A708**, 413 (2002); *Nucl. Phys.* **A717**, 249 (2003); K. Tuchin, *AIP Conf. Proc.* **698**, 698 (2004); D. Kharzeev, E. Levin, L. McLerran, `hep-ph/0403271`.