Color Glass Condensate at RHIC?

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Heavy-ion collisions at the BNL-RHIC collider can probe whether gluon saturation effects in nuclei at small $x$ have set in, or whether leading-twist perturbative estimates of particle production are still applicable. I discuss that soon to come centrality dependence of hadron production from RHIC collisions at RHIC may provide more systematic insight into this problem than present data on the energy and centrality dependence of hadron production from Au+Au collisions. Results from pA collisions at RHIC could also help to understand some controversial and puzzling results from ion-ion collisions, for example the large azimuthal asymmetry at high $p_\perp$, as measured by the STAR collaboration.

I. INTRODUCTION

Perturbative (short-distance) QCD predicts a double-logarithmic distribution of gluons in a hadron or nucleus at small $x$ (i.e., at rapidities far away from the valence quarks): $dN_g \sim dk_\perp^2/k_\perp^2 dx/x$. This result, however, is derived in the dilute limit where some effects, like gluon recombination, have been neglected. Therefore, as $k_\perp$ decreases, the DGLAP approximation should eventually become invalid, and the predicted rise of the gluon density should flatten out.

Namely, at large distances (though still much smaller than the confinement scale) the non-linear terms in the Yang-Mills equations tame the growth of the color charge density in hadrons and nuclei as compared to the power-law growth predicted by the linearized DGLAP theory. The gluon density should saturate at distances where the color charge density $\rho_s$ per unit transverse area becomes of order $1/g^2$, where $g$ is the coupling, and where therefore QCD-evolution becomes genuinely non-perturbative. Such large occupation numbers are characteristic of a condensate, and so classical methods may be useful in determining the color field in a hadron or nucleus at small $x$, i.e. at high density (the “saturation effect” is expected to show up in large nuclei much “earlier”, at larger $x$, than for protons). The local color charge density is a stochastic variable (fluctuating randomly as one moves around in the transverse plane) and observables eventually have to be averaged over some distribution of $\rho_s$. This is analogous to spin glasses, and hence the name “Color Glass Condensate” for the structure of a nucleus at small $x$.

The Relativistic Heavy-Ion Collider (RHIC) at BNL is dedicated to the study of high-energy hadronic interactions, from $p+p$ over $p+A$ to $A+B$ collisions. One of the important issues to settle is whether the above-mentioned gluon saturation scale $Q_s$ at the values of $x$ that are relevant at RHIC energy is in fact (substantially) larger than the QCD confinement scale $\Lambda$. Optimistic estimates do indeed yield $Q_s^2(y=0) \sim 2$ GeV$^2$ for central Au+Au collisions at RHIC, that is $Q_s^2/\Lambda^2 \sim 50$. If so, then perturbative estimates of particle production at leading twist can not be employed to compute the dominant part of the inelastic cross section because one needs to resum all powers of $\alpha_s$ times the charge density squared per unit of transverse area (“higher twists”). In other words, particle production will be dominated by the decay of the classical non-abelian field, except at very short distances (truly high $p_t$) where gluon occupation numbers are small and one enters the region of applicability of perturbation theory. The initial conditions for the subsequent evolution, which possibly proceeds through a so-called Quark-Gluon Plasma, then have to be obtained from the Color Glass Condensate model. On the other hand, if it turns out that at RHIC energy $Q_s$ is not substantially larger than $\Lambda$, as some conservative estimates indicate, then this means that leading-twist perturbation theory is safe at RHIC, and that it can be employed to describe the bulk of particle production.

To understand particle multiplicities and energy densities in the central region one does not need to consider the soft classical field then.

Naturally, as RHIC started operation with collisions of gold ions, theoretical studies focused on the energy and centrality dependence of particle production from the CGC model versus that from leading twist perturbation theory with fixed infrared cutoff $p_0 \sim 2$ GeV. Within the CGC model, a rather mild increase of $dN_{ch}/d\eta$ both with centrality (or the number of participants) and with energy is expected. In contrast, phase space for particle production at $p_\perp > p_0$ rapidly opens up in perturbation theory, and so the growth is faster. The energy dependence of particle production at RHIC turns out to be quite weak, approximately consistent with saturation models. However, to my opinion it is hard for both saturation models and soft-hard two-component models to make absolute predictions on the 10-20% level (at present), which would be required to falsify one or the other model using just the limited energy range of RHIC. Similar concerns apply to the centrality dependence of particle production. This observable is plagued in particular by the uncertainties in determining experimentally the number of participants for large
impact parameter collisions. In that respect, it would be
cleaner to determine $dN_{ch}/dy/N_{part}$ by using only cen-
tral collisions but varying the mass number of the collid-
ing nuclei. Of course, this will not be available within
the near future, although it is doable in principle at RHIC,
which can collide essentially any ions from $p$ to $Au$.

Is it possible then to verify or falsify the gluon sat-
uation in nuclei at RHIC energies? I think a good
opportunity for that exists in $pA$ since protons are likely
quite dilute at RHIC energy, while saturation may occur
for large nuclei \[ 13,14 \]. In that case, distinct signatures
emerge in the transverse momentum and rapidity distri-
butions of secondaries, which are correlated (and so can
be tested experimentally without necessarily relying on
absolute predictions). Also, one might be able to test
various scalings of observables like $dN/\hat{y}$, $\langle p_{\perp}\rangle$ and so
on. More detailed observables like the $p_{\perp}$-broadening of
minijets at forward rapidities \[ 13 \], or diffractive single-jet
production \[ 16 \] may as well provide deeper insight.

II. PROTON-NUCLEUS COLLISIONS

$pA$ collisions at high energy are very interesting be-
cause one can actually compute the total multiplicity
analytically \[ 13,14 \], and also make some predictions for
transverse momentum spectra and the scaling of $\langle p_{\perp}\rangle$
with the produced multiplicity $dN/\hat{y}$.

The key point is that the color charge density per
unit transverse area is much smaller in a proton than
in a big nucleus. Therefore, the saturation momentum
of the proton is much smaller than that of the nucleus,
$Q_s^{(1)} < Q_s^{(2)}$. At $p_{\perp} > Q_s^{(2)}$, the fields of both the proton
and of the nucleus are weak, and so one expects that or-
dinary perturbation theory is applicable. In particular,
to compute single-inclusive gluon production at high $p_{\perp}$
the well-known expressions from collinear factorization
(with DGLAP evolution) can be employed.

On the other hand, the inelastic cross section (particle
production) is dominated by the kinematic region $Q_s^{(1)} <
 p_{\perp} < Q_s^{(2)}$. In that regime, the field of the proton is still
weak, and so can be treated within perturbation theory;
but the field of the nucleus is in the non-linear regime,
and one must resum interactions of the radiated gluons
with the background field of the nucleus to all orders.
A simple way to understand why is to realize that while
each additional gluon “line” from the produced gluon to
the nuclear field comes with an additional power of the
coupling $\alpha_s$, the occupation numbers of the nuclear field
at (or below) the scale $Q_s^{(2)}$ are of order $1/\alpha_s$, and so
additional “rescatterings” are not suppressed by powers of
the coupling constant.

The calculation of gluon production essentially consists
then of solving the Yang-Mills equations in the forward
light-cone to all orders in the background field of the nu-
cleus, but to leading order in the field of the weak proton.
From the Fourier transform of the field at asymptotic
times one reads off the amplitudes of the modes, which are
then squared and averaged over the color charge distri-
butions in the proton and in the nucleus, using a Gaus-
sian weight \[ 1 \]. Finally, this leads to the gluon distribu-
tion \[ 1 \]

\[
\frac{dN}{d^2 b d^2 k_{\perp} dy} = g^2 \gamma(1, y) \frac{N_c^2 - 1}{(2\pi)^4} \int d^2 u_{\perp} \int_{\Lambda_{QCD}} d^2 p_{\perp} e^{i(p_{\perp} - k_{\perp}) \cdot u_{\perp}} \chi_1(y) \chi_2(y) \frac{k_{\perp}}{Q_s^{(2)}} \exp \left\{ \frac{y}{\Lambda_{QCD}^2} \right\},
\]

with

\[
\gamma(x_{\perp}) = \int_{\Lambda_{QCD}^2} d^2 q_{\perp} \left( \frac{4}{(2\pi)^2} q_{\perp}^2 \right) \frac{e^{i q_{\perp} \cdot x_{\perp}}}{q_{\perp}^4} \frac{k_{\perp}}{Q_s^{(2)}} \chi_1(y).
\]

The gluon propagator. $\chi_1(y)$ denote the effective den-
sities of color charge as “seen” by the gluon produced
at rapidity $y$. It is the integrated charge density of ei-
ther one of the sources, from its beam rapidity down to
$y$. Note that $\chi_1$ is linear in $\hat{y}$ but resums all orders
in $\chi_2$. A numerical evaluation of the above result has
not yet been performed, mainly because one needs to de-
terminate the dependence of $\chi_2$ on rapidity \[ 3 \]. One can
nevertheless gain some insight by considering two lim-
its. Namely, at asymptotically large $k_{\perp}$, the exponential
can be expanded to first non-trivial order, leading to the
well-known pQCD result

\[
\frac{dN}{d^2 b dy} = g^2 N_c(N_c^2 - 1) \frac{\chi_1(y)\chi_2(y)}{k_{\perp}^4} \log \frac{k_{\perp}}{Q_s^{(2)}},
\]

where we introduced the shorthand notation $Q_s^{(1,2)}(y) \equiv \sqrt{g^2 N_c \chi_1,2(y)/8\pi \Lambda_{QCD}}$. Thus, one recovers the standard perturbative $\alpha_s^2/k_{\perp}^4$ behavior at very high $k_{\perp}$, with a logarithmic correction analogous to DGLAP evolu-
tion. Note that $\chi_1, \chi_2$ scale as $A_{1/3}^{1/3} A_2^{1/3}$ \[ 16 \], while the integral over $d^2 b$ gives a factor
of $\pi R_s^2 \propto A_2^{2/3}$. Therefore, in this kinematic region
$dN/d^2 b dy$ scales like $A_{1/3}^{1/3} A_2$, up to logarithmic
corrections. This holds also for the integrated distribution
$dN(k_{\perp} > p_0)/dy$ above some fixed $A_2$-independent scale
$p_0$, as in the soft-hard two-component model. On the
other hand, when integrating over $k_{\perp}$ from $Q_s^{(2)}$ to in-
finity, the contribution from large $k_{\perp}$ to the rapidity
density is

\[
\frac{dN}{d^2 b dy} = \frac{g^2 (N_c^2 - 1)}{\pi^2} \chi_1(y).
\]

Again, the integral over $d^2 b$ gives a factor $\pi R_s^2 \sim A_2^{2/3}$, and so $dN/dy$ scales like $A_{1/3}^{1/3} A_2^{2/3}$.

On the other hand, when $Q_s^{(1)}(y) \ll k_{\perp} \ll Q_s^{(2)}(y)$, one finds that
\[
\frac{dN}{d^2bdk_\perp dy} \approx 2g^2\chi(y)(\frac{N_c^2}{2\pi^3})^{\frac{1}{2}} \frac{1}{k_\perp} \log \frac{k_\perp}{Q_s^{(1)}(y)}.
\]

This form \( \sim \alpha_s \chi_1 / k_\perp^2 \) is to be compared with that from eq. (3), \( \sim \alpha_s^2 \chi_1 \chi_2 / k_\perp^4 \), valid at high \( k_\perp \). A schematic distribution in transverse momentum is shown in Fig. 1.

![FIG. 1. Schematic \( k_\perp \) distribution for gluons produced in high-energy \( A_1 + A_2 \) collisions at rapidity \( y \) such that \( Q_s^{(1)}(y) \ll Q_s^{(2)}(y) \). In the perturbative regime, \( dN/dk_{\perp}^2 dy \sim 1/k_\perp^4 \). In between the saturation scales for the two sources, \( dN/dk_{\perp}^2 dy \sim 1/k_\perp^2 \). Although the figure does not provide a quantitative prediction for experiments at RHIC, one should notice non-trivial features. It is predicted that when strong-field effects set in around \( Q_s^{(2)} \), the \( k_\perp \) distribution of secondaries flattens. Not only does the scaling with \( A_2 \) differ in the two regimes, as mentioned above, but more importantly the “turnover” point \( Q_s^{(2)}(y) \) is a function of rapidity! Experimentally, one can thus take the rapidity dependence of the saturation momentum from a parametrization of HERA data [1], which also seems to fit the observed rapidity dependence of \( dN/dy \) from \( Au + Au \) at RHIC [2], and test whether the turnover in the transverse momentum distribution from Fig. 1 moves as one changes the rapidity in a way consistent with those parametrizations.

From (3), the \( k_\perp \)-integrated multiplicity in the nonperturbative regime \( Q_s^{(1)}(y) \leq k_\perp \leq Q_s^{(2)}(y) \) is

\[
\frac{dN}{d^2bdy} = g^2\chi(y)(\frac{N_c^2}{2\pi^3})^{\frac{1}{2}} \\log \frac{2Q_s^{(2)}(y)}{Q_s^{(1)}(y)}.
\]

Thus, at fixed impact parameter, the multiplicity scales as \( A_1^{1/3} \), up to the square of a logarithm of \( (A_2/A_1)^{1/3} \).

The average transverse momentum in the saturation regime is given by

\[
\langle k_\perp \rangle = 2Q_s^{(2)} \frac{\xi - 1 - \log \xi}{\log^2 \xi},
\]

where \( \xi(y) = Q_s^{(1)}(y)/Q_s^{(2)}(y) \). From dimensional considerations, it has been suggested [2] that in symmetric \( A + A \) collisions, and at central rapidity, \( \langle k_\perp \rangle^2 \) scales with the multiplicity per unit of transverse area and of rapidity, \( \langle k_\perp \rangle^2 \propto dN/d^2dy \). A similar scaling relation can be derived from eqs. (5, 7) for the asymmetric case,

\[
\langle k_\perp \rangle^2 \propto \frac{dN}{d^2bdy} \frac{g^2}{\xi^2} \frac{(\xi - 1 - \log \xi)^2}{\log^6 \xi}.
\]

Thus, \( \langle k_\perp \rangle^2 \) is proportional to the multiplicity per unit of rapidity and transverse area, times a function of the ratio of the saturation momenta. If source one is very much weaker than source two, i.e. in the limit \( |\log \xi| \gg 1 - \xi \), the third factor on the right-hand-side of (8) depends on \( \log \xi \) only. Neglecting that dependence, and assuming as before that \( \chi_{1,2} \) are proportional to \( A_{1,2}^{1/3} \), one has the approximate scaling relation

\[
\langle k_\perp \rangle^2 \propto \left( \frac{A_2}{A_1} \right)^{1/3} \frac{dN}{d^2bdy} \propto \left( \frac{1}{A_1A_2} \right)^{1/3} \frac{dN}{dy}.
\]

In practice though one expects significant corrections to the simple scaling relation (9), as given by eq. (8). \( pA \) collisions provide a natural testing ground for the scaling of \( \langle p_\perp \rangle \) with \( dN/dy \) since distortions of the above scaling relations from final-state interactions (“hydrodynamic expansion”) should not be a major issue.

**TABLE I.** Scaling of various quantities in the perturbative and saturation kinematic regimes.

| | sat. | pert. |
|---|---|---|
| \( dN/d^2bdk_\perp dy \) | \( A_1^{1/3} \) | \( A_1^{1/3} A_2^{1/3} \) |
| \( \langle k_\perp \rangle^2 \) | \( (A_1A_2)^{-1/3} \frac{dN}{dy} \) | 1 (for fixed \( p_0 \) scale) |
| \( \frac{d}{dy} \log \frac{dN}{d^2kd_\perp dy} \) | \( \frac{d}{dy} \log Q_s^{(1)}(y) \) | \( \frac{d}{dy} \left[ \log Q_s^{(1)}(y) + \log Q_s^{(2)}(y) \right] \) |
For rapidities far from the fragmentation region of the large nucleus, and for $Q_s^{(1)}(y) \ll k_\perp \ll Q_s^{(2)}(y)$, $\text{d}N/\text{d}^2k_\perp \text{d}y$ varies with rapidity like

$$\frac{\text{d} \text{d}N}{\text{d}y \text{d}^2k_\perp \text{d}y} \propto g^2 \chi_1^2(y) \cdot (10)$$

Thus, an experimental measure for the evolution of the CGC density parameter in rapidity is

$$\frac{\text{d} \log \text{d}N/\text{d}^2k_\perp \text{d}y}{\text{d}y} = \frac{\text{d} \log \chi_1(y)}{\text{d}y} \cdot (11)$$

Now consider the case of high $k_\perp \gg Q_s^{(2)}(y)$ described by eq. (3). In that regime the rapidity distribution is proportional to $\chi_1(y)\chi_2(y)$, and so $\text{d}N/\text{d}y\text{d}^2k_\perp$ varies with rapidity like

$$\frac{\text{d} \log \text{d}N/\text{d}^2k_\perp \text{d}y}{\text{d}y} = \frac{\text{d} \log \chi_1(y)}{\text{d}y} + \frac{\text{d} \log \chi_2(y)}{\text{d}y} \cdot (12)$$

Subtracting (11) from (12) provides an experimental measure for the renormalization-group evolution of $\chi_2(y)$. This does not only provide an independent verification of the HERA fits [13], but also a consistency check on how the turnover point (where the distribution flattens) from fig. 1 moves as one goes to different rapidity! In table I I have listed a few scaling relations of various quantities.

In fact, the results obtained above may even be relevant for collisions of equal-size nuclei, like $Au + Au$. Namely, if we consider a collision at non-zero impact parameter, as shown in fig. 2, it is easy to realize that near point A in that figure the situation is precisely the one considered above for $pA$ collisions: a low-density projectile colliding with a high-density target.

From the discussion above, the $k_\perp$ distribution of gluons produced in the impact is harder near point A than near point B, due to saturation effects in the target. Now, if those gluons suffer from radiative energy loss as they interact with each other, it is clear that a detector positioned at the top of the page will detect less momentum flowing into it than a detector positioned at the right of the page. Thus, radiative energy loss will generate an azimuthal asymmetry

$$v_2(p_\perp) = \frac{\langle p_\perp^2 - p_\perp^2 \rangle}{p_\perp^2} \cdot (13)$$

at transverse momenta on the order of a few GeV [2]. The point here is that $v_2$ will be larger if the $p_\perp$-distribution of produced gluons is harder at A than at B as compared to the case when $\text{d}N/\text{d}^2p_\perp$ is the same at A and B. This might be important for understanding the large values of $v_2$ measured by STAR at RHIC at $p_\perp \sim 2-6$ GeV [2], as it has been pointed out that large energy loss by itself (assuming the same $p_\perp$-distribution throughout the overlap zone in fig. 1) does not lead to large enough $v_2$ [23] in particular for semi-central events (10% centrality class). Quantitative computations of $v_2$ from the gluon saturation model are presently under way [24].

So far, all attempts to understand the large $v_2$ at high $p_\perp$ invoke radiative parton energy loss [25]

$$- \Delta E \approx \frac{3\alpha_s}{\pi} E _{cr} \log \frac{2E}{L m_\perp^2} \cdot (14)$$

Here, $\lambda$ is the typical length of the medium traversed by the jet, $m_\perp$ is the IR cutoff provided by the mass of static electric gluons, $E > E _{cr}$ is the jet energy, and $E _{cr} = m_\perp^2 L_\perp^2 / \lambda$. $\Delta E / E$ depends on the density of the medium via the mean-free path $\lambda$ of the radiation. Going down in $\sqrt{s}$, presumably this should reduce $v_2$ at high $p_\perp$ (in the sense that the leading-twist prediction $v_2 = 0$ should be approached faster). Is this the case?

Finally, if indeed $v_2$ at $p_\perp = 2 - 6$ GeV is generated by final-state interactions of the minijets, it might be sensitive to nearly critical scattering at $T_c$ [27] which would evidently be very exciting. QCD with 3 colors is not an ideal gas near $T_c$; rather, the pressure $p(T)$ is significantly smaller than the ideal-gas pressure $p_{id}(T)$. The transition from confinement to deconfinement appears to be nearly a second order transition [27], in which case $(m_\perp / T)^2 \sim \sqrt{\partial p / \partial T_d}$ drops rapidly as $T_c$ is approached from above [28]. This behavior is quite different from a strong first-order transition, and can enhance energy loss (and $v_2$) near $T_c$.

**III. CONCLUSIONS**

$Au + Au$ collisions from RHIC may have offered a first glance at gluon saturation effects in large nuclei: exper-
mentally, the energy and centrality dependence of particle production is quite weak. However, in my opinion it is too early yet to draw definite conclusions from the \( \sqrt{s} = 130 \) AGeV and 200 AGeV. Collisions at lower energy and in particular with different projectile/target combinations should be analyzed in the future. Collisions of small projectiles on large targets, for which I use the generic acronym “pA”, could reveal distinct scaling laws predicted by the Color Glass Condensate model.

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