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Pre-thermalization dynamics: initial conditions for QGP at the LHC and RHIC from perturbative QCD [1]

K.J. Eskola[2]

CERN, Theory Division,
CH-1211 Geneva 23, Switzerland
and
Department of Physics, University of Jyväskylä,
P.O. Box 35, FIN-40351 Jyväskylä, Finland

Abstract

I discuss how the initial conditions for QGP-production in ultrarelativistic heavy ion collisions at the LHC and RHIC can be computed from perturbative QCD.

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[2] kari.eskola@cern.ch

1 Introduction

Particle and transverse energy production in the central rapidity region of ultrarelativistic heavy ion collisions (URHIC) can be viewed as a combination of perturbative (hard and semihard) parton production and non-perturbative (soft) particle production. The key feature in the division is that one should be able to *compute* the hard and semihard parton production from perturbative QCD (pQCD), provided that the parton distributions of colliding nuclei are known [1, 2, 3]. On top of this, the non-perturbative particle production can be modelled *e.g.* through strings [4, 5], perhaps also through a decaying strong background colour field [6].

The semihard QCD-processes are expected to become increasingly important with increasing cms-energies $\sqrt{s}$, particularly in URHIC, due to the following reasons: Firstly, already in $p\bar{p}(p)$ collisions the rapid rise of the total and inelastic cross sections can be explained by copious production of semihard partons, "minijets," with transverse momenta $p_{\rm T} \geq p_0 \sim 1...2$ GeV [7]. In other words, the events tend to become more "minijetty" towards higher $\sqrt{s}$. This is also expected to happen in AA collisions, and even more so, because of the $A^{4/3}$-scaling of hard collisions.

Secondly, the rapid rise of the structure function F_2 observed at HERA at $x \lesssim 0.01$ and $Q \gtrsim 1$ GeV [8] implies that minijet production at the LHC gets a further boost [9]: the mid-rapidity minijets with $p_{\rm T} \sim 2$ GeV typically probe the gluon distributions at $x \sim 2p_{\rm T}/\sqrt{s} \sim 7 \times 10^{-4}$ for nuclear collisions at $\sqrt{s}/A = 5.5$ TeV.

Thirdly, the time scale for producing partons and transverse energy into the central rapidity region by semihard collisions is short, typically $\tau_{\rm h} \sim 1/p_0 \sim 0.1$ fm$/c$, where $p_0 \sim 2$ GeV is the smallest transverse momentum included in the computation. The soft processes are completed at later stages of the collision, typically at $\tau_{\rm s} \sim 1/\Lambda_{\rm QCD} \sim 1$ fm$/c$. If the density of partons produced in the hard and semihard stages of the heavy ion collision becomes high enough - as will be the case - a saturation in the initial parton production can occur already in the perturbative region $p_{\rm T} \gg \Lambda_{\rm QCD}$ [1, 10, 11, 12], and, consequently, softer particle production will be screened. The fortunate consequence of this is that a larger part of parton production in the central rapidities can be *computed* from perturbative QCD (pQCD) at higher energies, and the relative contribution from soft collisions with $p_{\rm T} \lesssim 2$ GeV becomes smaller. Typically, the expectation is that at the SPS (Pb+Pb at $\sqrt{s} = 17$ AGeV), the soft component dominates, and at the LHC (Pb+Pb at $\sqrt{s} = 5.5$ ATeV) the semihard component is the dominant one. At RHIC (Au+Au at $\sqrt{s} = 200$ AGeV) one will be in an intermediate region where both components should be taken into account.

The semihard processes have also been implemented in several event generators for URHIC [13]. For example, HIJING [5] and Parton Cascade Model (and VNI) [14] rely on the dominance of initial semihard particle production at high cms-energies.

In what follows, I am going to focus on how to compute the initial minijet and transverse energy production in an AA-collisions from pQCD, and what are the implications of such a computation for the initial conditions of QGP-formation at the LHC and RHIC. I will also discuss the uncertainties in the computation.

2 Production of semihard $g, q, \bar{q}$ in AA

Hadronic jets originating from high-$p_{\rm T}$ parton-parton scatterings have been observed in $p\bar{p}$ collisions from $p_{\rm T} \gtrsim 5$ GeV [15] up to $p_{\rm T} \sim 440$ GeV [16]. Below 5 GeV it becomes very difficult to distinguish the individual transverse energy clusters [15] from the underlying background, so that minijets with $p_{\rm T} \sim 2$ GeV cannot be observed directly [3]. In AA collisions, where hundreds (RHIC) or thousands (LHC) of minijets are expected to be produced within the central rapidity unit, it becomes certainly impossible to distinguish them individually. However, due to their abundance, minijets are expected to dramatically contribute to QGP-formation and its further evolution at the LHC and RHIC.

Unlike in a jet measurement (or next-to-leading order (NLO) calculation) we do not define any jet size for an individual (mini)jet, but rather, our goal will now be to compute minijet and transverse energy production in the *whole* central rapidity unit $|y| \leq 0.5$. The key quantity is the integrated minijet cross section,

$$\sigma_{\rm jet}(\sqrt{s}, p_0) = \frac{1}{2} \int_{p_0^2} dp_{\rm T}^2 dy_1 dy_2 \sum_{\substack{ij<kl>= \\ q,\bar{q},g}} x_1 f_{i/A}(x_1, Q)\, x_2 f_{j/A}(x_2, Q) \times$$
$$\times \left[\frac{d\hat{\sigma}}{d\hat{t}}^{ij\to kl}(\hat{s}, \hat{t}, \hat{u}) + \frac{d\hat{\sigma}}{d\hat{t}}^{ij\to kl}(\hat{s}, \hat{u}, \hat{t}) \right] \frac{1}{1+\delta_{kl}}. \qquad (1)$$

Collinear factorization is assumed, and $f_{i/A}(x, Q)$ are the number densities (per nucleon) of partons i in a nucleus A at a fractional momentum x and a factorization scale Q, chosen here as $Q = p_{\rm T}$. From conservation of energy and momentum in simple $2 \to 2$ kinematics with negligible initial transverse momenta it follows that

$$x_1 = \frac{p_{\rm T}}{\sqrt{s}}({\rm e}^{y_1} + {\rm e}^{y_2}) \quad \text{and} \quad x_2 = \frac{p_{\rm T}}{\sqrt{s}}({\rm e}^{-y_1} + {\rm e}^{-y_2}), \qquad (2)$$

where $y_{1,2}$ are the rapidities of the outgoing partons k, l and $p_{\rm T}$ is the magnitude of their transverse momentum. Note that in this case k and l are back-to-back in transverse momentum. The kinematic invariants on the parton level become

$$\hat{s} = x_1 x_2 s = 2p_{\rm T}^2(1 + \cosh Y), \;\; \hat{t} = -p_{\rm T}^2(1 + {\rm e}^{-Y}), \quad \hat{u} = -p_{\rm T}^2(1 + {\rm e}^{Y}), \qquad (3)$$

where $Y = y_1 - y_2$.

[3] In this sense minijet is not a "jet" at all.

In the lowest order pQCD there are eight types of subprocesses $\hat{\sigma}^{ij\to kl}$:

$$gg \to gg, q_i\bar{q}_i \quad gq \to gq \quad q_i\bar{q}_i \to q_j\bar{q}_j, q_i\bar{q}_i, gg \quad q_iq_j \to q_iq_j \quad q_iq_i \to q_iq_i. \tag{4}$$

The cross sections can be found in the literature [17]. In $\sigma_{\rm jet}$ the contribution from $gg \to gg$ is clearly dominant. In the computation below, I will use (4), but as a little side-remark it is perhaps worth noticing that the "single effective subprocess" -approximation [18] works also nicely: it is straightforward to show that

$$\frac{[d\hat{\sigma}(\hat{s},\hat{t},\hat{u})/d\hat{t}]^{gg\to gg}}{[d\hat{\sigma}(\hat{s},\hat{t},\hat{u})/d\hat{t}]^{gq\to gg}+[d\hat{\sigma}(\hat{s},\hat{u},\hat{t})/d\hat{t}]^{gq\to gg}} = \frac{9}{4} + \left\{ \begin{array}{ll} \mathcal{O}(\chi), & \chi \ll 1 \\ \mathcal{O}(1/\chi), & \chi \gg 1 \end{array} \right. \tag{5}$$

where $\chi \equiv \hat{u}/\hat{t}$. Considering similarly the process $qq \to qq$, one obtains an approximate relation between the different main subprocesses:

$$(gg \to gg) : (gq \to gq) : (qq \to qq) = 1 : \frac{9}{4} : (\frac{9}{4})^2, \tag{6}$$

so that Eq. (1) becomes simply [3]

$$\sigma_{\rm jet}(\sqrt{s}, p_0) = \frac{1}{2}\int_{p_0^2} dp_{\rm T}^2 dy_1 dy_2 x_1 F(x_1, Q) x_2 F(x_2, Q) \frac{d\hat{\sigma}}{d\hat{t}}^{gg\to gg} \tag{7}$$

with $F(x,Q) \equiv xg(x,Q) + \frac{4}{9}\sum_q x[q(x,Q)+\bar{q}(x,Q)]$.

Let us first neglect all nuclear effects in parton densities and take $f_{i/A} = f_{i/p} \equiv f_i$, which is a reasonable first approximation to start with. Then, from the *pp*-level quantity $\sigma_{\rm jet}$, we obtain the *average* number of produced minijets with $p_{\rm T} \geq p_0$ in an AA collision with an impact parameter $\mathbf{b}$ simply by multiplying $\sigma_{\rm jet}$ by the nuclear overlap function T_{AA} [2, 3]: $\bar{N}_{AA}(\mathbf{b}) = 2T_{AA}(\mathbf{b})\sigma_{\rm jet}$, where

$$T_{AA}(\mathbf{b}) = \int d^2\mathbf{s}\, T_A(\mathbf{s}) T_A(\mathbf{b}-\mathbf{s}), \tag{8}$$

where the thickness function is

$$T_A(\mathbf{s}) = \int dz\, n_A(\sqrt{s^2+z^2}) \tag{9}$$

with normalizations $\int d^2\mathbf{b} T_{AA}(\mathbf{b}) = A^2$ and $\int d^2\mathbf{s} T_A(\mathbf{s}) = A$, *i.e.* the hard cross sections in AA scale as A^2, as expected in absence of any nuclear effects in $f_{i/A}$. For central collisions with Woods-Saxon nuclear densities n_A we get $T_{AA}(b=0) \approx A^2/(\pi R_A^2)$, so clearly $N_{AA} \sim A^{4/3}$. Here one should notice that this approach is based on *independent* binary parton collisions, and is valid when each semihard sub-collision consumes only a negligible fraction of the beam-energy, and when there are sufficiently many partons available for scattering. These conditions are fulfilled for minijet production at central rapidities even up to the LHC-energies and $A \sim 200$. However, at the LHC for large y together with a very large A, one eventually runs into trouble with energy and baryon number conservation [12], which is an indication that in this regime one should consider coherence effects in parton production. This is easy to understand simply by noticing that the available energy and valence quark number scale as $\sim A$, while $N_{AA} \sim A^{4/3}$. For RHIC energies, however, there does not seem to be a problem in the above approach for any y. [12]

We want to focus on the central rapidity unit, so we will have to do some book-keeping of where the partons are produced in rapidity, and define our acceptance cuts. Especially, we want to compute the average transverse energy carried by the partons in the central rapidity unit. Furthermore, we would like to keep track of contributions from gluons, quarks and antiquarks separately. Let us therefore define the $E_{\rm T}$-distribution for each flavour f in our acceptance window Δy as [3, 12]:

$$
\begin{aligned}
\frac{d\sigma^f}{dE_{\rm T}} = \frac{1}{2}\int dp_{\rm T}^2 dy_1 dy_2 \sum_{\substack{ij\\ \langle kl\rangle}} x_1 f_i(x_1,Q)\, x_2 f_j(x_2,Q) \frac{1}{1+\delta_{kl}} \times \\
\times\Big\{ \frac{d\hat\sigma}{d\hat t}^{ij\to kl}(\hat s,\hat t,\hat u)\,\delta(E_{\rm T}-[\delta_{fk}\epsilon(y_1)+\delta_{fl}\epsilon(y_2)]p_{\rm T}) + \\
+\frac{d\hat\sigma}{d\hat t}^{ij\to kl}(\hat s,\hat u,\hat t)\,\delta(E_{\rm T}-[\delta_{fl}\epsilon(y_1)+\delta_{fk}\epsilon(y_2)]p_{\rm T})\Big\},
\end{aligned}
\tag{10}
$$

with a normalization $\sum_f \int dE_{\rm T}\frac{d\sigma^f}{dE_{\rm T}} = \sigma_{\rm jet}(\sqrt{s},p_0)$. Above, our acceptance is defined through a "measurement function" $\epsilon(y)$, which will in the following be chosen as a step function

$$
\epsilon(y) = \begin{cases} 1, & \text{if } |y| \le 0.5 \\ 0, & \text{otherwise.} \end{cases}
\tag{11}
$$

In Fig. 1, the integration region in the (y_1, y_2)-plane (with a fixed $p_{\rm T}$),

$$
|y_1| \le \log\Big[\frac{\sqrt{s}}{2p_{\rm T}} + \sqrt{(\frac{\sqrt{s}}{2p_{\rm T}})^2-1}\,\Big],\ -\log\Big[\frac{\sqrt{s}}{p_{\rm T}} - {\rm e}^{-y_1}\Big] \le y_2 \le \log\Big[\frac{\sqrt{s}}{p_{\rm T}} - {\rm e}^{y_1}\Big],
\tag{12}
$$

is illustrated together with the cuts corresponding to our acceptance $\Delta y : |y| \le 0.5$. After this, it is straightforward to compute the first $E_{\rm T}$-moment for each flavour f in the semihard collision,

$$
\begin{aligned}
\sigma_{\rm jet}(\sqrt{s},p_0)\langle E_{\rm T}^f\rangle_{\Delta y} &= \int dE_{\rm T} E_{\rm T}\frac{d\sigma^f}{dE_{\rm T}} \\
&= \int dp_{\rm T}^2 dy_1 dy_2 \sum_{\substack{ij\\ \langle kl\rangle}} x_1 f_i(x_1,Q)\, x_2 f_j(x_2,Q) \times \\
&\times\frac{1}{1+\delta_{kl}}\Big[\delta_{fk}\frac{d\hat\sigma}{d\hat t}^{ij\to kl}(\hat s,\hat t,\hat u) + \delta_{fl}\frac{d\hat\sigma}{d\hat t}^{ij\to kl}(\hat s,\hat u,\hat t)\Big]p_{\rm T}\epsilon(y_1)
\end{aligned}
\tag{13}
$$

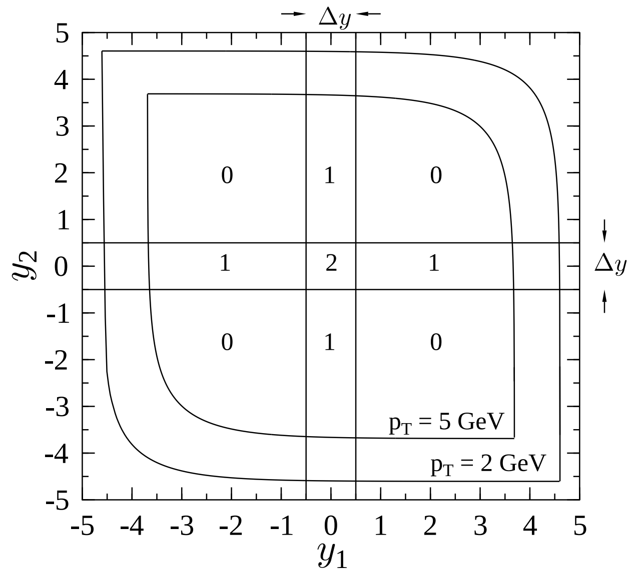

Figure 1: Kinematic cuts in rapidity. The oval regions show the rapidity phase space (12) for fixed $p_{\rm T} = 2$ GeV and 5 GeV, and $\sqrt{s} = 200$ GeV. The values of $\epsilon(y_1) + \epsilon(y_2)$ corresponding to our acceptance region, $\Delta y : |y| \leq 0.5$, are indicated.

In an AA collision with an impact parameter $\mathbf{b}$, the average transverse energy carried by semihard partons of flavour f to the acceptance region becomes then [3, 12]

$$\bar{E}_{\rm T}^f(\mathbf{b}, \sqrt{s}, p_0, \Delta y) = T_{AA}(\mathbf{b})\sigma_{\rm jet}(\sqrt{s}, p_0)\langle E_{\rm T}^f\rangle_{\Delta y}, \tag{14}$$

where $T_{AA}(\mathbf{b})\sigma_{\rm jet}(\sqrt{s}, p_0)$ is the average number of semihard collisions (all y) and $\langle E_{\rm T}^f\rangle_{\Delta y}$ is the average transverse energy carried by the flavour f at $|y| \leq 0.5$ in each semihard collision.

We can also compute the number distribution $d\sigma^f/dn$ for each flavour in a similar way (in Eq.(10) replace $E_{\rm T}$ by n, and in the delta-functions $p_{\rm T}$ by 1) to obtain the average number of flavour f produced within our acceptance window:

$$\bar{N}_{AA}^f(\mathbf{b}, \sqrt{s}, p_0, \Delta y) = T_{AA}(\mathbf{b})\sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{\Delta y}, \tag{15}$$

where $\sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{\Delta y}$ can be obtained directly from Eq. (13) simply by removing the weight $p_{\rm T}$. The normalization $\sum_f \sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{\Delta y\to\infty} = 2\sigma_{\rm jet}(\sqrt{s}, p_0)$ can be verified from Eqs. (1) and (13).

3 Initial conditions at $\tau = 0.1$ fm/c

The subprocess cross sections $d\hat{\sigma}^{ij}/d\hat{t}$ diverge as $\sim p_{\rm T}^{-4}$ when $p_{\rm T} \to 0$, which is why we cannot reliably extend our computation below $p_0 = 1...2$ GeV. The QCD-evolution of the parton densities does not improve the situation enough because of two competing effects: towards smaller p_T both x and Q decrease; the downwards QCD-evolution reduces the densities (for fixed, small x) but the simultaneous small-x increase of the densities has an effect in the other direction. We would like to extend our computation down to as low values of $p_{\rm T}$ as possible but making sure we can still believe that we are dealing with perturbative quantities. Of course, only after computing the NLO minijet cross sections, the question of convergence of the perturbation series can be better answered. For large-$p_{\rm T}$ jet measurements the NLO-calculation has been done [19] but for the few GeV $p_{\rm T}$-range and for the rapidity acceptance we have in mind this has not been completed yet.

So, how to fix the parameter p_0? In Fig. 2, I have plotted the behaviour of $\sigma_{\rm jet}(\sqrt{s}, p_0)\langle E_{\rm T}^f\rangle_{\Delta y}$ and $\sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{\Delta y}$ for the LHC and RHIC energies as functions of p_0. In the computation, I have assumed that $f_{i/A} \approx f_i$. I have used the GRV94 LO-distributions [22], which reproduce the observed small-x rise of the partons densities reasonably well. Note that there is no *ad hoc* "K-factor" included in the computation.

One way to fix p_0 is to formulate a saturation criterion [1, 10, 12] for the produced parton system in URHIC: let us assign a transverse area $\pi/p_{\rm T}^2$ for each parton within our rapidity acceptance. Most of the partons have $p_{\rm T} \sim p_0$, so that the effective total transverse area occupied by the partons is $\bar{N}_{AA} \times \pi/p_0^2$. The available nuclear transverse area can be estimated as πR_A^2, so that the saturation of parton production in $|y| \leq 0.5$ should occur when $\bar{N}_{AA} \times \pi/p_0^2 \gtrsim \pi R_A^2$, *i.e.* when

$$\sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{|y|\leq 0.5} \gtrsim 134\left(\frac{p_0}{2\,{\rm GeV}}\right)^2 {\rm mb}, \tag{16}$$

where $R_{\rm Pb} = 6.54$ fm and $T_{\rm PbPb}(0) = 32/$mb has been used. This procedure results in $p_0 \approx 2$ GeV for the LHC, and $p_0 \approx 1$ GeV for RHIC, as illustrated in Fig. 2.

Saturation of parton production can also be studied in a more self-consistent - but still phenomenological - manner by computing a screening mass (electric, static) generated by the partons produced by a fixed time $\tau < 1/p_{\rm T}$ [11]. This mass then screens further parton production at smaller $p_{\rm T}$ and later in time. The more partons are produced the stronger the screening becomes, and, the more damped further parton production is. In this way, saturation of transverse energy production seems again to take place at $p_{\rm T} \sim 2$ GeV for LHC and $p_{\rm T} \sim 1$ GeV for RHIC [11].

Lower limits for p_0 may also be studied in an eikonal approach [7], or, by convoluting the parton cross sections with the fragmentation functions of each parton flavour into pions and kaons [20]. Comparison of the fragmentation function approach with the measured $p_{\rm T}$-distributions of charged particles in $p\bar{p}$ collisions is successful at $\sqrt{s} \geq 200$

GeV at large $p_{\rm T}$ [21], whereas it seems that allowing parton scatterings with $p_{\rm T} < 1.5$ GeV would result in an overestimate of the charged particle $p_{\rm T}$-distributions. Therefore, let me fix $p_0 = 2$ GeV both for the LHC and for RHIC. Estimates for smaller p_0 can be easily obtained from Fig. 2.

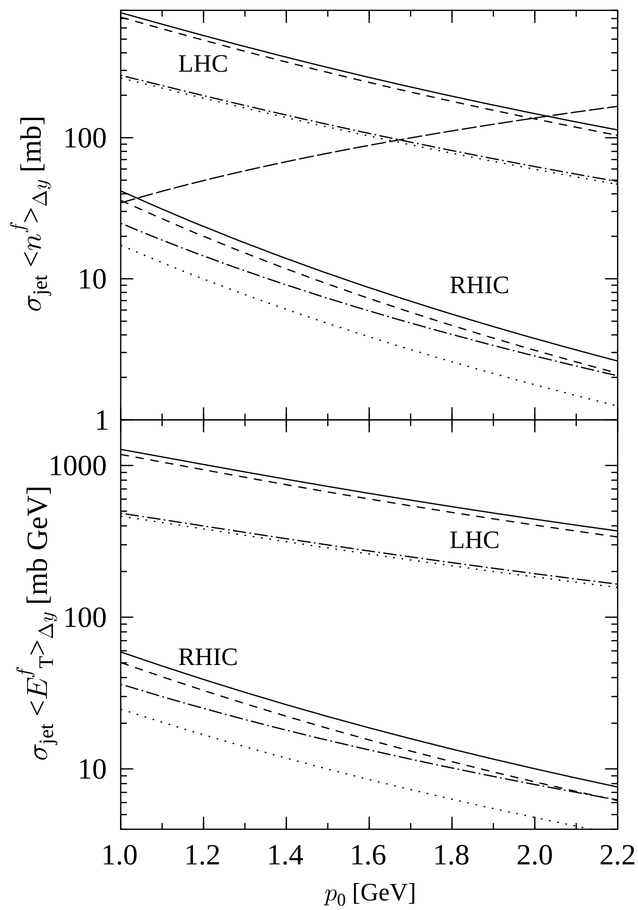

Figure 2: quantities $\sigma_{\rm jet}(\sqrt{s}, p_0)\langle n^f\rangle_{\Delta y}$ and $\sigma_{\rm jet}(\sqrt{s}, p_0)\langle E_{\rm T}^f\rangle_{\Delta y}$ from Eqs. (13-15) for $\sqrt{s} = 5500$ GeV (LHC) and $\sqrt{s} = 200$ GeV (RHIC) as functions of p_0 for $\Delta y : |y| \leq 0.5$. The solid curves show the total results, the dashed ones the gluon contributions. For the figure, quark and antiquark contributions have been multiplied by a factor 10 (7) for the LHC (RHIC) and they are shown with dot-dashed and dotted curves, respectively. The saturation limit from Eq. (16) is the increasing dashed line in the upper panel.

Table 1 shows the average numbers and transverse energies of semihard partons at $\tau = 0.1$ fm/c with $|y| \leq 0.5$ and $p_{\rm T} \geq 2$ GeV in central Pb–Pb collisions at the LHC and RHIC, as given by Eqs. (15) and (14).

We make the following five interesting observations:

1. Gluons clearly outnumber quarks and antiquarks and dominate $E_{\rm T}$-production at $\tau \sim 0.1$ fm/c both at the LHC and RHIC. The QGP is actually gluon plasma to a

	$\|y\| < 0.5$	total	g	q	$\bar{q}$
LHC:	$\bar{N}^f_{\rm PbPb}$	4741	4350	200	191
	$\bar{E}^f_{\rm T}$	14160	12950	619	590
RHIC:	$\bar{N}^f_{\rm PbPb}$	121	99.6	13.0	8.10
	$\bar{E}^f_{\rm T}$	321	263	36.1	21.9

Table 1: The initial conditions at $\tau = 0.1$ fm: $N^f_{\rm PbPb}$ and $\bar{E}^f_{\rm T}$ for $b = 0$, $\sqrt{s}/A = 5500$ GeV (LHC) and 200 GeV (RHIC), $p_0 = 2$ GeV and $|y| \leq 0.5$.

first approximation.

2. With $p_0 = 2$ GeV $\gg \Lambda_{\rm QCD}$ the glue is saturated at the LHC. This indicates that the pQCD-domain is dominant for $E_{\rm T}$-production. In other words, more gluons are produced into the system at $p < p_0$ but their transverse energy is relatively small. For RHIC I expect the soft component to be still important.

3. Estimating the volume of the system at $\tau_{\rm h} = 0.1$ fm/c and with $\Delta y = 1$ by $V_{\rm i} = \pi R^2_{\rm Pb}\Delta y\tau_{\rm h}$, we can convert the results into local energy and number densities: $\epsilon_f^{\rm pQCD} = \bar{E}^f_{\rm T}/V_{\rm i}$ and $n_f^{\rm pQCD} = \bar{N}^f_{\rm PbPb}/V_{\rm i}$. From the numbers in Table 1 for the LHC, we see that $E_{\rm T}$/gluon ≈ 3 GeV (naturally, since $p_0 = 2$ GeV). On the other hand, for an ideal gas of massless bosons $\epsilon_g^{\rm ideal}/n_g^{\rm ideal} \approx 2.7T_{\rm eq}$. For an ideal gas with $\epsilon_g^{\rm ideal} = \epsilon_g^{\rm pQCD}$ we get $T_{\rm eq} \approx 1.1$ GeV, which means that $(\epsilon_g/n_g)^{\rm pQCD} \approx (\epsilon_g/n_g)^{\rm ideal}$, and the glue seems to be initially thermalized at the LHC. [12] This result is essentially due to the small-x increase of the gluon distributions [9]. At RHIC the gluon thermalization will take place a little later. One should also keep in mind that a complete thermalization requires locally isotropic momentum distributions, and to study this more detailed modelling is needed.

4. Since $n_q, n_{\bar{q}} \ll n_g$, quarks and antiquarks are initially obviously far away from chemical equilibrium both at RHIC and the LHC.

5. By assigning a baryon number $1/3(-1/3)$ for each quark(antiquark) produced at $\tau = 0.1$fm/c, we observe that the initial net baryon number density at $|y| \leq 0.5$ becomes $n_{B-\bar{B}} = 0.21\,{\rm fm}^{-3}$ for the LHC and 0.12 fm^{-3} for RHIC. This density will naturally dilute quite fast ($\sim 1/\tau$) but it is interesting that for the LHC the initial number is actually larger than for RHIC, and even larger than the nuclear matter density 0.17 fm^{-3}. This behaviour is again due to the small-x increase of the gluon densities. Also, for the LHC, it is worth noticing that the transit time of the colliding nuclei is $2R_A/\gamma \sim 0.005$ fm/$c \ll 0.1$ fm/c. Finally, the initial net baryon number-to-entropy ratio can also be estimated [12], and we find that $(B-\bar{B})/S \sim 1/5000\,(1/2000)$ for the LHC (RHIC), so that we are far away from the conditions of the early Universe, where the inverse of the specific entropy is $\sim 10^{-9}$.

4 Discussion and outlook

I have shown how to compute initial conditions for QGP-production in URHIC at the LHC and RHIC from the lowest order pQCD. The analysis can obviously be sharpened in various ways: by computing the transverse energy production in NLO in α_s and by including nuclear effects in the parton distributions in a consistent, scale-dependent manner [23]. Resummation similar to the Drell-Yan case [24], if applicable, would hopefully help to reduce the uncertainty due to the parameter p_0. Also higher twist effects and more coherent scattering should be studied and formulated in more detail at large y. Connection to the BFKL-physics potentially relevant to the minijets involving small-x should be understood better [25] as well. Recently, a novel approach based on classical gluon fields has been developed [26] for semihard parton production. Applicability of this model at the energies of RHIC and the LHC has been studied lately [27]. Also, understanding the pre-thermal evolution of the newly produced QGP is a challenging task, involving space-time dependent phenomena in gauge theories.

To conclude, it is evident that there are several interesting open questions in pre-thermalization physics of URHIC, and much more precision-work is needed. Understanding the primary parton production mechanisms will be essential for understanding the further evolution and signals of the QGP, and more global variables measured at nuclear collisions at the LHC and RHIC.

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