Elliptic Flow from Non-equilibrium Initial Condition with a Saturation Scale

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A current goal of relativistic heavy ion collisions experiments is the search for a Color Glass Condensate (CGC) as the limiting state of QCD matter at very high density. In viscous hydrodynamics simulations, a standard Glauber initial condition leads to estimate $4\pi \eta/s \sim 1$, while employing the Kharzeev-Levin-Nardi (KLN) modeling of the plasma leads to at least a factor of 2 larger $\eta/s$. Within a kinetic theory approach based on a relativistic Boltzmann-like transport simulation, our main result is that the out-of-equilibrium initial distribution reduces the efficiency in building-up the elliptic flow. At RHIC energy we find the available data on $v_2$ are in agreement with a $4\pi \eta/s \sim 1$ also for KLN initial conditions. More generally, our study shows that the initial non-equilibrium in p-space can have a significant impact on the build-up of anisotropic flow.

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Ultra-relativistic heavy-ion collisions (uRHICs) at the Relativistic Heavy-Ion Collider (RHIC) and the Large Hadron Collider (LHC) create a hot and dense system of strongly interacting matter. In the last decade it has been reached a general consensus that such a state of matter is not of hadronic nature and there are several signatures that it is a strongly interacting quark-gluon plasma (QGP) \[1,2\]. A main discovery has been that the QGP has a very small shear viscosity to density entropy, $\eta/s$, which is more than one order of magnitude smaller than the one of water \[4,5\], close to the lower bound of $1/4\pi$ conjectured for systems at infinite strong coupling \[13\]. A key observable to reach such a conclusion is the so-called elliptic flow, $v_2 = \langle \cos(2\varphi_p) \rangle = \langle (p_x^2 - p_y^2)/(p_x^2 + p_y^2) \rangle$, with $\varphi_p$ being the azimuthal angle in the transverse plane and the average meant over the particle distribution. In fact, the expansion of the created matter generates a large anisotropy of the emitted particles that can be primarily measured by $v_2$. Its origin is the initial spatial eccentricity, $\epsilon_2 = (x^2 - y^2)/(x^2 + y^2)$, of the overlap region in non-central collisions. The observed large $v_2$ is considered a signal of a very small $\eta/s$ because it means that the system is very efficient in converting $\epsilon_2$ into an anisotropy in the momentum space $v_2$, a mechanism that would be strongly damped in a system highly viscous that dissipates and smooths anisotropies \[1,2,11\]. Quantitatively both viscous hydrodynamics \[1,2,12,13,14\], and transport Boltzmann-like approaches \[16,20\] agree in indicating an average $\eta/s$ of the QGP lying in the range $4\pi \eta/s \sim 1 - 3$.

The uRHIC program offers the tantalizing opportunity to explore the existence of an exotic state, namely the Color Glass Condensate (CGC) \[21,22\], see \[23,24\] for reviews. Such a state of matter would be primarily generated by the very high density of the gluon parton distribution function at low $x$ (parton momentum fraction), which triggers a saturation of the gluon distribution function at a $p_T$ below the saturation scale, $Q_g$ \[25\]. Even if at first sight surprisingly, the study of the shear viscosity $\eta/s$ of the QGP and the search for the CGC are related. In fact, the main source of uncertainty for $\eta/s$ comes from the unknown initial conditions of the created matter \[8\] and confirmed later by further works \[10,12,26\].

A simple geometrical description through the Glauber model \[28\] predicts a $\epsilon_2$ smaller at least 25-30% than the eccentricity of the CGC, for most of the centralities of the collisions, see for example results within the Kharzeev-Levin-Nardi (KLN) model \[27,28,30\], factorized KLN (fKLN) model \[31\], Monte Carlo KLN (MC-KLN) model \[31,32\] and dipole model \[27,33\]. The uncertainty in the initial condition translates into an uncertainty on $\eta/s$ of at least a factor of two as estimated by mean of several viscous hydrodynamical approaches \[8,10,12,26\]. More explicitly, the experimental data of $v_2(p_T)$ at the highest RHIC energy are in agreement with a fluid at $4\pi \eta/s \sim 1$ according to viscous hydrodynamics simulation, assuming a standard Glauber initial condition. Assuming an initial fKLN or MC-KLN space distribution the comparison favors a fluid at $4\pi \eta/s \sim 2$. The reason is the larger initial $\epsilon_2$ of the fKLN, which leads to larger $v_2$ unless a large $\eta/s$ is considered. However, in \[10,26\] it has been shown that viscous hydrodynamics fails to reproduce both $v_2$ and $v_3 = \langle \cos(3\varphi_p) \rangle$ if the same $4\pi \eta/s \sim 2$ is assumed. At variance a Glauber initial condition can account for both with the same $4\pi \eta/s = 1$. However the indirect effect on $\epsilon_2$ is not a unique and solid prediction of the CGC modellings; for example, the approach based on the solution of the Classical Yang-Mills (CYM) equations predict a somewhat smaller initial eccentricity \[34,36\]. Very recently employing a $x - space$ distribution inspired to the CYM approach in a viscous hydrodynamical approach \[53\] it has been shown that not only $v_2$ but also higher harmonics can be correctly predicted with a $4\pi \eta/s \sim 1.5$ instead of $\sim 2$, which is in
qualitative agreement with the fact that CYM tend to predict quite smaller \( \epsilon_s \) with respect to fKLN. However, our present studies focus on the effect of the initial nonequilibrium in \( p \)-space, an issue discarded in all previous studies including the recent ones \[35,36\].

In this Letter, we point out that the implementation of the melted CGC in hydrodynamics takes into account only the different space distribution with respect to a geometric Glauber model, discarding the key and more peculiar feature of the CGC that is below the \( Q_s \) saturation scale. We have found by mean of kinetic theory that this has a pivotal role on the build-up of \( v_2 \).

We adopt the model which was firstly introduced by Kharzeev, Levin and Nardi \[29\] (KLN model) even if in the regime of over saturation in \( A + A \) collisions some aspects are better caught by a CYM approach \[48\]. In particular, to prepare the initial conditions of our simulations we refer to the factorized-KLN (fKLN) approach as introduced in \[31\] \[32\]. This will allow for a direct comparison with viscous hydrodynamics results, in which the coordinate space distribution function of gluons arising from the melted CGC is assumed to be

\[
\frac{dN_g}{dy d^2x_\perp} = \int d^2p_T p_A(x_\perp) p_B(x_\perp) \Phi(p_T, x_\perp, y),
\]

where \( \Phi \) corresponds to the momentum space distribution in the \( k_T \) factorization hypothesis \[37\] \[38\].

\[
\Phi(p_T, x_\perp, y) = \frac{4\pi^2 N_c}{N_c - 1} \frac{1}{p_T^2} \int d^2k_T \alpha_s(Q^2)
\times \phi_A(x_1, k_T^2, x_\perp)
\times \phi_B(x_2, (p_T - k_T)^2, x_\perp).
\]

Here \( x_{1,2} = p_T \exp(\pm y)/\sqrt{s} \) and the ultraviolet cutoff \( p_T = 3 \text{ GeV}/c \) assumed in the \( p_T \) integral in Eq. (1); \( \alpha_s \) denotes the strong coupling constant, which is computed at the scale \( Q^2 = \max(k_T^2, (p_T - k_T)^2) \) according to the one-loop \( \beta \) function but frozen at \( \alpha_s = 0.5 \) in the infrared region as in \[30\] \[33\] \[39\]. In Eq. (1) \( p_{A,B} \) denote the probability to find one nucleon at a given transverse coordinate, \( p_A(x_\perp) = 1 - (1 - \sigma_{in} T_A(x_\perp)/A)^\Lambda \), where \( \sigma_{in} \) is the inelastic cross section and \( T_A \) corresponds to the usual thickness function of the Glauber model.

The main ingredient to specify in Eq. (2) is the unintegrated gluon distribution function (uGDF) for partons coming from nucleus \( A \), which is assumed to be:

\[
\phi_A(x_1, k_T^2, x_\perp) = \frac{k_T^2}{\alpha_s(Q^2)} \left[ \frac{\theta(Q_s - k_T^2)}{Q_s^2 + \Lambda^2} + \frac{\theta(k_T^2 - Q_s^2)}{k_T^2 + \Lambda^2} \right],
\]

where we see the peculiar feature of the CGC that is the saturation of the distribution for \( p_T < Q_s \); a similar equation holds for partons belonging to nucleus \( B \). Following \[31\] we take the saturation scale for the nucleus \( A \) as

\[
Q_{s,A}(x, \perp) = 2 \text{GeV}^2 \left( \frac{T_A(x_\perp)}{1.53 p_A(x_\perp)} \right) \left( \frac{0.01}{x} \right)^{\lambda},
\]

with \( \lambda = 0.28 \), and similarly nucleus \( B \). This choice is the one adopted in fKLN or MC-KLN and in hydro simulations \[12\] \[31\] to study the dependence of \( v_2(p_T) \) on \( \eta/s \). Using Eqs. (1) and (11) we find that \( \langle Q_s \rangle \approx 1.4 \text{ GeV} \) where the average is understood in the transverse plane.

We employ transport theory as a base of a simulation code of the fireball expansion created in relativistic heavy-ion collision \[10\] \[20\] \[40\] \[41\], therefore the time evolution of the gluons distribution function \( f(x, p, t) \) evolves according to the Boltzmann equation:

\[
p_{\ell} \partial_x f_1 = \int d^2\Gamma_2 d\Gamma_3 f_2 (f_1 f_2 - f_1 f_2) \times |\mathcal{M}|^2 \delta^4(p_{1} + p_{2} - p_{1} - p_{2}),
\]

where \( d^3p_{k} = 2E_{k} (2\pi)^3 d\Gamma_{k} \) and \( \mathcal{M} \) corresponds to the transition amplitude.

At variance with the standard use of transport theory, we have developed an approach that, instead of focusing on specific microscopic calculations or modelings for the scattering matrix, fixes the total cross section in order to have the wanted \( \eta/s \). In Ref.\[42\] it has been shown in 1+1D such an approach is able to recover the Israel-Stewart viscous hydrodynamical evolution when \( \eta/s \) is sufficiently small. In 3+1D some of the authors has studied the analytical relation between \( \eta/s \), temperature,
cross section and density and as shown in [41, 44], the 
Chapmann-Enskog approximation supplies such a rela-
tion with quite good approximation [43], in agreement 
with the results obtained using the Green Kubo formula. 
Therefore, we fix \( \eta/s \) and compute the pertinent total 
cross section by mean of the relation 
\[
\sigma_{\text{tot}} = \frac{1}{15} \rho g(a) \frac{1}{\eta/s},
\]
which is valid for a generic differential cross section 
\( d\sigma/dt \sim \alpha_s^2 (t - m_D^2)^3 \) as proved in [44]. In the above 
equation \( a = T/m_D \), with \( m_D \) the screening mass regul-
ating the angular dependence of the cross section, while 
\[
g(a) = \frac{1}{50} \int dyy^9 \left[ (y^2 + \frac{1}{3})K_3(2y) - yK_2(2y) \right] h \left( \frac{a^2}{y^2} \right),
\]
with \( K_n \) the Bessel function and \( h \) corresponding to the 
ratio of the transport and the total cross section. The maximum value of \( g \), namely \( g(m_D \rightarrow \infty) = h(m_D \rightarrow \infty) = 2/3 \), is reached for isotropic cross section; a smaller 
value of \( g(a) \) means that a higher \( \sigma_{\text{tot}} \) is needed to repro-
duce the same value of \( \eta/s \). However, we notice that in 
the regime were viscous hydrodynamic applies (not too 
large \( \eta/s \) and \( p_T \)) the specific microscopic detail of the 
cross section is irrelevant and our approach is an effective 
way to employ transport theory to simulate a fluid at a 
given \( \eta/s \). From the operative point of view, keeping \( \eta/s \) 
constant in our simulations is achieved by evaluating lo-
caley in space and time the strength of the cross section 
by means of Eq. [45], where both parton densities and temperature are computed locally in each cell. To real-
ize a realistic freeze-out, when the local energy density 
reaches the cross-over region, the \( \eta/s \) increases linearly 
to match the estimated hadronic viscosity, as described in [20, 41]; this affects in the same way all the cases con-
sidered in the following.

In the following, we will consider three different types 
of initial distribution function in the phase-space, two of 
which are the one employed till now for the investigation 
of the \( \eta/s \), while the third one is the genuine novelty of the 
present study. Furthermore, we refer to \( Au + Au \) col-
losion at \( \sqrt{s} = 200 \text{AGeV} \) and \( b = 7.5 \text{fm} \). In this case, 
our result for initial eccentricity in the fKLN model is 
\( \epsilon_x = 0.357 \) (which is in agreement with MC-KLN [31] 
result used in hydro simulations). The standard initial 
condition for simulations of the plasma fireball created 
at RHIC is a \( x \)-space distribution given by the Glauber 
model and a \( p \)-space thermalized spectrum in the tran-
verse plane at a time \( \tau = 0.6 \text{fm}/c \) with a maximum initial 
temperature \( T_0 = 340 \text{MeV} \). In this case, for a standard 
mixture of \( N_{\text{part}} \) and \( N_{\text{coll}} \) we find \( \epsilon_x = 0.284 \). We will 
refer to this case as Th-Glauber. Instead the study of the 
impact of an initial CGC state has been performed con-
sidering an \( x \)-space distribution given by the fKLN (or 
MC-KLN), while in the momentum space the spectrum 
has been considered thermalized at \( \tau_0 \sim 0.6 \text{fm}/c \); we refer 
to this case as Th-fKLN and it is represents the case 
implemented in hydrodynamics, that has lead to the con-
clusion that the CGC suggests a \( 4\pi \eta/s \sim 2 \) [8, 10, 12, 24]. 
The third initial conditions is the full fKLN initial condi-
tions where, beyond the \( x \)-space, the saturated dis-
tribution in \( p \)-space is implemented as well, see Fig. [4] 
solid thick line. As initial time we take \( \tau_0 = 0.15 \text{fm}/c \) be-
cause in this case there is no pre-assumption of thermal-
ization. This is not usually considered in hydrodynamics 
because there it is implicitly assumed a distribution function in \( p \)-space in local equilibrium, at least in the transverse plane. The choice of \( \tau_0 \) for the plasma-like ini-
tial condition is inspired by the recent results of [50, 52] 
where it is discussed that, even if at \( \tau = 0^+ \) the longi-
tudinal pressure of the glasma is negative, within a time 
\( \tau \approx 1/3Q_s \) it becomes positive thus making a description of the 
expanding system in terms of a partonic distribution 
function quite reliable. For all the previous cases, 
as usually done, a Bjorken scaling at the initial time is 
assumed, identifying momentum rapidity \( y_L \) with space 
rapidity \( \eta_L = arctg q^{-1}((z/\tau) \) For all the case considered 
the multiplicity \( dN/dy \) at mid rapidity has been fixed 
initially equal to 400 to approximately match the experi-
mental data that for the impact parameter considered 
corresponds to about \( N_{\text{part}} = 150 \) in [45].

In Fig. [4] we plot the initial spectra for the IKLN 
(thick solid line), Th-fKLN (dashed line) and Th-Glauber 
(thin solid line) at their respective initial times \( \tau_0 \), 
and the spectrum of the fKLN model after a time evolution 
\( \Delta \tau = 0.8 \text{fm}/c \) (dashed green line). We notice that ini-
tially the fKLN spectrum is quite far from a thermalized 
spectrum; in fact, it embeds the saturation effects which 
are proper of the melted CGC. Nevertheless, the spec-
trum evolves to a thermalized one within 1 \text{fm}/c. Such a 
feature is confirmed by the inset of Fig. [4] where the 
quantity \( T^* \cdot \tau^{1/3} \) is shown, with \( T^* = E/3N \) represent-
ing, the temperature in the case of a thermalized system. 
It is known that in the case of 1D expansion a thermal-
ized system should keep such a ratio constant. We find 
that in the case of the IKLN (solid green line) the product 
\( T^* \cdot \tau^{1/3} \) is strongly dependent on time because the system 
is quite far from equilibrium; however at \( \tau \approx 0.8 \text{fm}/c \) 
both the value and the time evolution become indistin-
guishable from the thermal cases represented by the Th-
Glauber and Th-fKLN. We notice a little adjustment also 
for these cases that we have indicated as thermal. The 
reason is that the initial spectra are thermal only in the 
transverse plane, but there is a boost invariance in the 
longitudinal direction. This causes a little re-adjustment 
that would disappear assuming a thermal spectral also 
in the longitudinal direction.

Our results on thermalization time are not in disagree-
ment with earlier studies showing that two-body collis-
sions are insufficient to achieve a fast thermal equilibrium 
[46]. In fact in that case a perturbative QCD two-body
cross section is employed, corresponding to $\eta/s$ about one order of magnitude larger than in our case $^{[44]}$, while here we normalize the cross section to get an $\eta/s$ corresponding to few fm/c if the coupling is strong enough, or in general the dynamics sufficiently nondissipative $^{[52]}$.  

In the left panel of Fig. 3 we plot the $v_2(p_T)$ for the case of Th-Glauber (thin solid line) and Th-fKLN (thick solid line) at a fixed $4\pi\eta/s = 1$. The Glauber initial condition reproduces quite well the data (circles); in the case of Th-fKLN (thick solid line) one gets a too large $v_2$ and for such initial conditions the agreement with the data is achieved only if the $\eta/s$ is increased by a factor of two (dashed line). These results are in agreement with the ones obtained from viscous hydrodynamics $^{[8, 10, 12, 26]}$, showing the solidity and consistency of our transport approach at fixed $\eta/s$.  

In the right panel of Fig. 3 we present our novel result for the fKLN model, when the CGC distribution function is implemented in both the $x$ and $p$ spaces. We find that fKLN with a $4\pi\eta/s = 1$ (thick solid line) gives a $v_2(p_T)$ quite similar to the Th-Glauber, while in such a case if $4\pi\eta/s = 2$, dashed line, the $v_2(p_T)$ would be too small. Our interpretation is that the initial larger $\varepsilon_x$ is compensated by the key feature of an almost saturated initial distribution in $p$-space below the saturation scale $Q_s$. In other words the initial out-of-equilibrium fKLN distribution reduces the efficiency in converting $\varepsilon_x$ into $v_2$. In fact, the elliptic flow can be understood as a larger slope of the momentum spectrum in the out of plane $\vec{x}$ direction with respect to the $\vec{y}$ one caused by a larger pressure in the $\vec{x}$ direction due to the elliptical shape. The net effect in terms of the difference of the particle yields between the two directions is larger if the spectra are decreasing exponentially with respect to the case in which they are nearly flat as a function of $p_T$. A detailed study is in preparation, but the result we present in this Letter shows that the initial out-of-equilibrium function implied by the isotropization itself. This result is not in disagreement with other studies $^{[52, 51]}$ which show how the expanding plasma becomes almost isotropic within few fm/c if the coupling is strong enough, or in general the dynamics sufficiently nondissipative $^{[52]}$.
distribution with a saturation behavior generates smaller $v_2$ with respect to the thermal one. This result is quite general, and we expect it should be valid, besides QGP in uRHICs, for systems like cold atoms in a magnetic trap which are characterized by a value of $\eta/s$ close to the QGP one \cite{17}. In the specific case of the KLN matter studied here, the effect of the initial non-equilibrium distribution affects the estimate of $\eta/s$ of about a factor of two. The relevance of our results is further enhanced by the fact that Th-fKLN with $4\pi\eta/s \sim 2$ would generate a low $v_2$ with respect to the available data, which is the main conclusion of \cite{28}. We notice that in the present kinetic approach the quantum nature of gluons has been discarded. This is justified at RHIC because in this Letter we have focused on an effect which is dominant at $p_T > 0.5$ GeV, where the $f(x,p)$ is still smaller than unity. At LHC, or anyway at small $p_T$, it would be necessary to include $(1+f)$ terms in Eq.\((\pi)\) that could drive the system toward a Bose-Einstein condensate \cite{55}.

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\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{example_graph.png}
\caption{Evolution of $v_2$ at $p_T = 1.5$ GeV as a function of the evolution time for all the different initial conditions and $\eta/s$ values. Calculations refer to Au-Au collisions at $\sqrt{s} = 200$ GeV, with an impact parameter $b = 7.5$ fm.}
\end{figure}
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