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Broken boost invariance in the Glasma via finite nuclei thickness

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ABSTRACT

We simulate the creation and evolution of non-boost-invariant Glasma in the early stages of heavy ion collisions within the color glass condensate framework. This is accomplished by extending the McLerran–Venugopalan model to include a parameter for the Lorentz-contracted but finite width of the nucleus in the beam direction. We determine the rapidity profile of the Glasma energy density, which shows deviations from the boost-invariant result. Varying the parameters both broad and narrow profiles can be produced. We compare our results to experimental data from RHIC and find surprising agreement.

1. Introduction

Heavy ion collisions at the Relativistic Heavy Ion Collider (RHIC) and the Large Hadron Collider (LHC) provide insight into the properties of nuclear matter under extreme conditions. The evolution of the Quark–Gluon Plasma (QGP) that is created in such collisions is well described by relativistic viscous hydrodynamics [1,2]. A first principles description of the initial state of heavy-ion collisions is provided by the Color Glass Condensate (CGC) framework [3–5]. The CGC is a classical effective field theory for nuclear matter at ultrarelativistic energies. Models such as the IP-Glasma [6,7] in combination with hydrodynamics are able to correctly reproduce azimuthal anisotropies and event-by-event multiplicity distributions [8,9]. Furthermore, the CGC can explain long-range rapidity correlations like the ridge [10,11].

A Gaussian shaped rapidity profile of particle multiplicity can be found in experiments covering various energy ranges, from LHC [12] to RHIC Beam Energy Scan [13,14]. This shape is well explained by the Landau model [15] up to RHIC energies [13], which assumes full stopping of the colliding nuclei. The Landau model is in contrast to the Bjorken model [16] which relies on approximate boost invariance. A Gaussian profile has also been found in holographic calculations of colliding shock waves [17–19].

In its original formulation collisions in the CGC picture are assumed to be boost-invariant [20–23] and were thus only understood as an approximation valid close to midrapidity. This approach implicitly assumes infinitely thin Lorentz-contracted nuclei and entails a classical, boost-invariant evolution of the Glasma at leading order.

Only at the next-to-leading order the boost invariance is broken by a change of the initial conditions through JIMWLK evolution [24–27], and non-boost-invariant rapidity profiles can be obtained. Recently such rapidity dependencies have been found to agree reasonably well with experimental data where observables like charged particle multiplicities show a Gaussian rapidity profile [28,29]. On the other hand it has been suggested that if one considers nuclei with finite extent in the beam direction, deviations from boost invariance may arise already at the classical level [30]. In the case of proton–nucleus collisions methods have been developed to systematically include finite width corrections of the nucleus [31–33]. However, so far there has been no consistent simulation of the subsequent three-dimensional evolution for heavy-ion collisions even at the classical level.

In this letter, we show that Gaussian rapidity profiles of energy density can arise already from 3+1 dimensional purely classical CGC simulations, if incoming nuclei have a finite extent in the beam direction. As one would expect, the Gaussian profiles become broader at higher collision energy. In principle, we can cover the wide range from very thin nuclei with almost boost-invariant behavior to thick nuclei at low collision energies and narrow Gaussian profiles. We simulate the collision in the laboratory frame, see Fig. 1, which makes it necessary to include the propagating nuclei already before and during the collision.

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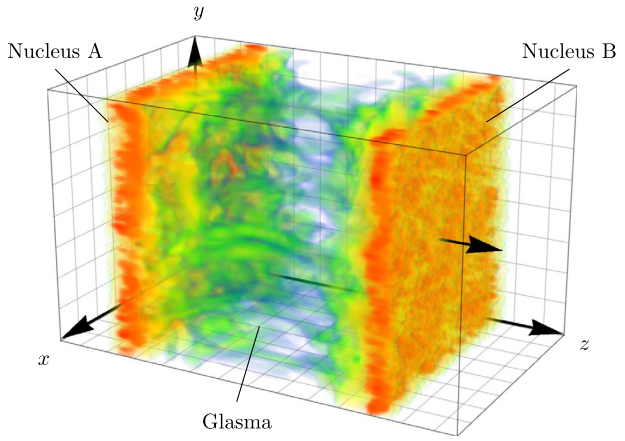

Fig. 1. A 3+1 dimensional colored particle-in-cell simulation of the collision of two thick sheets of relativistic nuclear matter. The simulation box covers a small part of the full transverse extent of the nuclei in the x, y plane. This figure shows a density plot of the energy density of both nuclei A and B, and the three-dimensional Glasma that is created in the collision. An animated version of this figure can be found at [34].

2. Theoretical framework

The nuclei in the CGC picture consist of hard partons that are surrounded by soft gluons. The hard partons can be described as classical color charges moving at the speed of light, while the soft gluons form a highly occupied coherent non-Abelian gauge field. The collision of two such infinitely thin condensates produces the Glasma whose evolution can be described classically by solving the Yang–Mills equations for early proper times. At finite nuclei width, the collision region is not pointlike anymore, and the nuclei, the collision, and the evolution of the Glasma can not be described separately and require one consistent simulation that covers all these steps.

A suitable numerical method was developed in our previous publication [35] based on the colored particle-in-cell method (CPIC) [36–39], which is a non-Abelian extension of the particle-in-cell method for the simulation of Abelian plasmas [40,41]. In contrast to the traditional approach of simulating the Glasma, where the field equations are solved in the forward light-cone parametrized by proper time τ and space–time rapidity η_s, we describe the collision in the laboratory frame using the lab-frame time t and the longitudinal beam direction z. The most striking difference of this approach is the explicit inclusion of the nuclei in the simulation, whereas in a boost-invariant simulation the information about the nuclei and their color currents is completely encoded in the initial conditions at the boundary of the light-cone, i.e. $\tau = 0$. To solve this problem numerically we simulate the continuous color charge densities of the nuclei with a large number of color-charged point-like particles, mimicking the dynamics of the continuous cloud of color charges on a lattice. This enables us to describe the full 3+1 dimensional collision and the subsequent evolution of the Glasma beyond the boost-invariant approximation. For a more detailed description we refer the reader to [35].

2.1. Initial conditions

The initial conditions in our simulation differ from the traditional approach as well. Instead of starting at $\tau = 0$, our simulation begins before the collision with the nuclei well-separated in the longitudinal direction. Here we quickly review how to solve the Yang–Mills equations in the covariant gauge and the transformation to the temporal gauge in the laboratory frame for a single nucleus. We base our model of the initial state on the McLerran–Venugopalan (MV) model [42,43], extended by a thickness parameter in longitudinal direction. The transverse charge density $\rho^a(x_T)$ as a function of the transverse coordinate x_T is a random variable following the usual gauge-invariant Gaussian probability functional $W[\rho]$ with the two-point correlation function

$$\left\langle \rho^a(x_T)\rho^b(y_T)\right\rangle = g^2\mu^2\delta^{ab}\delta^{(2)}(x_T - y_T), \tag{1}$$

where μ is the MV model parameter controlling average color charge density and g is the Yang–Mills coupling constant. For a single nucleus moving in the positive z direction, we embed this two-dimensional charge density into the three-dimensional laboratory frame via $\rho^a(x_T, x^-) = f(x^-)\rho^a(x_T)$ with a longitudinal profile function $f(x^-)$, where $x^\pm \equiv (t \pm z)/\sqrt{2}$ are the usual light-cone coordinates. For the longitudinal profile we choose a Gaussian

$$f(x^-) = \frac{1}{\sqrt{2\pi}L}\exp\left(-(x^-)^2/L^2\right), \tag{2}$$

where we introduce the thickness parameter L. In the limit of $L \to 0$ we have $f(x^-) \propto \delta(x^-)$ and restore the boost-invariant limit of the original MV model. Note that this model explicitly neglects non-trivial longitudinal color structure [44]. The only non-vanishing component of the light-like color current of the nucleus is then given by

$$J^+_{\text{cov}}(x_T, x^-) = \sqrt{2}f(x^-)\rho^a(x_T)t^a, \tag{3}$$

where t^a are the generators of the gauge group SU(N). The subscript "cov" denotes that this defines the color current in the covariant gauge $\partial_\mu A^{\mu,a} = 0$. Using this ansatz we can solve the Yang–Mills equations

$$D_\mu F^{\mu\nu} = J^\nu \tag{4}$$

in the covariant gauge by finding a solution to the two-dimensional Poisson equation

$$-\Delta_T A^+(x_T, x^-) = J^+_{\text{cov}}(x_T, x^-), \tag{5}$$

which is solved by

$$\phi^a(x_T) = \int_0^\Lambda \frac{d^2k_T}{(2\pi)^2}\frac{\tilde\rho^a(k_T)}{k_T^2 + m^2}e^{-ik_T\cdot x_T}, \tag{6}$$

$$A^+(x_T, x^-) = \sqrt{2}f(x^-)\phi^a(x_T)t^a, \tag{7}$$

where $\tilde\rho^a(k_T)$ is the Fourier transform of $\rho^a(x_T)$. We introduced an infrared regulator m and an ultraviolet cutoff Λ, since the MV model is both infrared and UV divergent. The regularization in (6) should be read as a modification of the charge densities $\rho^a(x_T)$ while the field equations remain unchanged.

Our numerical method requires the gauge fields to satisfy the temporal gauge condition $A^a_0 = 0$. Switching to this gauge from the covariant gauge renders the fields purely transverse. The transverse field components and the color current are given by

$$A_i(x_T, x^-) = \frac{i}{g}V(x_T, x^-)\partial_i V^\dagger(x_T, x^-), \tag{8}$$

$$J^+(x_T, x^-) = V(x_T, x^-)J^+_{\text{cov}}(x_T, x^-)V^\dagger(x_T, x^-), \tag{9}$$

with the temporal Wilson line

$$V(x_T, x^-) = \mathcal{T}\exp\left(-ig\int_{-\infty}^{t} dt' f(x'^-)\phi(x_T)\right). \tag{10}$$