Scattering amplitudes in QCD exhibit a definite RG flow with energy towards the unitarity limit. In this paper we put forward an evolution equation which allows one to modify continuously the pre-asymptotic RG flow towards “saturation” of Wilson line correlators. It preserves the linearly unstable zero field fixed point and the unitarity limit attractor. We present the evolution equation in a form suitable for efficient numerical solution. Hence, the proposed evolution equation in principle permits phenomenological comparisons of our incorrect pre-asymptotic RG flows to the BK/JIMWLK flow of QCD. Ultimately, it may lead to observational constraints on the approach to unitarity, constraining it to be specifically the one of QCD. However, it appears that single-inclusive particle spectra in the forward region of p+A collisions at RHIC alone do not provide stringent constraints on the evolution.

The density of soft gluons in strongly boosted hadrons or nuclei is non-perturbatively high [1–4] thereby saturating the unitarity limit of scattering amplitudes from such an object. QCD evolution in rapidity towards the unitarity limit is described at leading logarithmic (LL) accuracy by the B-JIMWLK renormalization group equations [5–14]. These can be cast in the form of a Langevin process in the space of Wilson line correlators of Wilson lines [14–16] over the transverse plane of the collision.

Our approach here is to construct a new RG flow with energy of all correlators of Wilson lines, i.e. for the entire Balitsky hierarchy, directly by modifying the noise correlator of JIMWLK. Expanding the conventional JIMWLK evolution equation for a small step dy in rapidity and using a noise-noise correlator
\[
\langle \xi^a_y \xi^b_y \rangle = \sigma^2 \delta^{ij} \delta^{ab} \delta y^2
\]
with an arbitrary variance \( \sigma^2 \), we obtain the modified Langevin process
\[
V_{\vec{x}}(y+dy) = V_{\vec{x}}(y) \left[ 1 + \frac{\sqrt{\alpha_s}}{\pi} dy \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \cdot \left( \xi_{\vec{z}} - V_{\vec{x}}^\dagger V_{\vec{x}} \xi_{\vec{z}} V_{\vec{x}} \right) - \frac{\alpha_s}{2\pi^2} \sigma^2 \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \left( 2 - U_{\vec{x}}^\dagger U_{\vec{x}} - U_{\vec{x}} U_{\vec{x}}^\dagger \right)^{ab} t^a t^b + \frac{\alpha_s}{2\pi^2} \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \left( U_{\vec{x}}^\dagger U_{\vec{x}} - U_{\vec{x}} U_{\vec{x}}^\dagger \right)^{ab} t^a t^b \right].
\]
We can rewrite adjoint Wilson lines in terms of fundamental ones, \( U_{\vec{x}}^{c_{ij}} = (V_{\vec{x}}^\dagger c_{ij} V_{\vec{x}}) \) to obtain the equivalent form
\[
\begin{align*}
V_{\vec{x}}(y+dy) &= V_{\vec{x}}(y) \left[ 1 + \frac{\sqrt{\alpha_s}}{\pi} dy \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \cdot \left( \xi_{\vec{z}} - V_{\vec{x}}^\dagger V_{\vec{x}} \xi_{\vec{z}} V_{\vec{x}} \right) \\
&\quad - \frac{\alpha_s}{2\pi^2} \sigma^2 \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \left( \sigma^2 N_c - \sigma^2 V_{\vec{x}}^\dagger V_{\vec{x}} V_{\vec{x}}^\dagger V_{\vec{x}} (V_{\vec{x}} V_{\vec{x}}^\dagger) + \frac{\sigma^2}{2} V_{\vec{x}}^\dagger V_{\vec{x}} V_{\vec{x}}^\dagger V_{\vec{x}} (V_{\vec{x}} V_{\vec{x}}^\dagger) - \frac{\sigma^2}{2} V_{\vec{x}}^\dagger V_{\vec{x}} V_{\vec{x}}^\dagger V_{\vec{x}} (V_{\vec{x}} V_{\vec{x}}^\dagger) \right) \right] \end{align*}
\]
Here, \( \tilde{K}_{\vec{x}} = \vec{x}/x^2 \) is the “square root” of the BFKL kernel [17, 18]. This equation suffices for the derivation of evolution equations of specific Wilson line operators. However, for numerical solution the evolution should be cast in a form that preserves \( V_{\vec{x}} \in SU(N_c) \) exactly. With some algebra one can show that this is achieved by writing
\[
V_{\vec{x}}(y+dy) = \exp \left\{ -i \sigma^{-2} \frac{\sqrt{\alpha_s}}{\pi} dy \int_{\vec{x}} \tilde{K}_{\vec{x} \rightarrow \vec{z}} \cdot \left( V_{\vec{x}}^\dagger V_{\vec{x}} \right) \right\}
\]

\*Electronic address: adrian.dumitru@baruch.cuny.edu
\*Electronic address: vskokov@ncsu.edu
\[ V_{\bar{x}} \exp \left\{ \frac{1}{\pi} \frac{\alpha_s \, dy}{s} \int_{\zeta} K_{\bar{x} - \zeta} \cdot \left[ \langle \xi_{\zeta} \rangle - (1 - \sigma^2) V_{\bar{x}}^{\dagger} V_{\bar{x}} \xi_{\zeta} V_{\bar{x}}^{\dagger} \right] \right\} . \] (4)

This reproduces eq. (3) to linear order in \( d_y \). Also, just like in standard JIMWLK evolution the implementation of eq. (4) requires just two (matrix valued) FFTs per rapidity step to compute the arguments of the exponentials for all \( \bar{x} \) (which is crucial for numerical solutions to be feasible). In the limit, \( \sigma \to 1 \), eq. (4) reproduces LL JIMWLK evolution.

From eq. (3) we can derive the evolution equation for the expectation value of the dipole scattering matrix \( S_{\bar{x} \bar{y}} = \text{tr} V_{\bar{x}} V_{\bar{y}}^{\dagger} / N_c \):

\[
\frac{\partial}{\partial Y} \langle S_{\bar{x} \bar{y}} \rangle = \frac{\bar{\alpha}_s \, \sigma^2}{2\pi} \int_{\zeta} \left\{ \frac{(\bar{x} - \bar{y})^2}{(\bar{x} - \zeta)^2 (\bar{y} - \zeta)^2} \left[ \langle S_{\bar{x} \bar{y}} \rangle - \langle S_{\bar{x} \zeta} S_{\zeta \bar{y}} \rangle \right] + \frac{1}{\bar{x} - \zeta} \frac{1 - \sigma^2}{2\sigma^2} \left[ \langle S_{\bar{x} \zeta} S_{\bar{y}} \rangle - \langle S_{\bar{x} \zeta} Q_{\bar{y} \bar{x} \bar{y}} \rangle \right] \right\} ,
\] (5)

where \( \bar{\alpha}_s = \alpha_s N_c / \pi \). For \( \sigma^2 = 1 \) this, of course, reduces to the B-JIMWLK evolution equation for the dipole. However, for \( \sigma^2 \neq 1 \) the r.h.s. involves the quadrupole \( Q_{\bar{x} \bar{y} \bar{z} \bar{w}} = \text{tr} V_{\bar{x}} V_{\bar{y}} V_{\bar{z}} V_{\bar{w}}^{\dagger} / N_c \). Note that UV divergences for \( \bar{z} \to \bar{x}, \bar{y} \) cancel for any value of \( \sigma^2 \). Diagrammatically [19, 20], the modification to the dipole evolution equation is due to a factor of \( \frac{1 + \sigma^2}{2\sigma^2} \) for the “real emission” diagrams where one of the quarks emits and reabsors a gluon; and due to a new contribution from this diagram corresponding to dipole \( \to \) dipole + quadrupole splitting. On the other hand, the virtual corrections and the real emission diagrams where the quark and the anti-quark exchange a gluon are unmodified. Therefore, we conjecture that no theory (defined in terms of Feynman diagrams) exists which corresponds to the evolution equation (5) for the dipole, when \( \sigma^2 \neq 1 \).

Equation (5) has two fixed points: one corresponds to the zero field limit \( S, Q = 1 \), the other to the strong field/unitarity limit \( S = Q = 1 \). We require the former to be linearly unstable, and the latter to be attractive. This restricts \( \sigma^2 > \frac{1}{2} \). In the limit \( \sigma^2 \to \infty \), for example, the evolution of \( \langle S_{\bar{x} \bar{y}} \rangle \) with rapidity is described by

\[
\frac{\partial}{\partial Y} \langle S_{\bar{x} \bar{y}} \rangle = \frac{\bar{\alpha}_s \, \sigma^2}{2\pi} \int_{\zeta} \left\{ \frac{(\bar{x} - \bar{y})^2}{(\bar{x} - \zeta)^2 (\bar{y} - \zeta)^2} \left[ \langle S_{\bar{x} \bar{y}} \rangle - \langle S_{\bar{x} \zeta} S_{\zeta \bar{y}} \rangle \right] - \frac{1}{2(\bar{x} - \zeta)} \frac{1}{2(\bar{y} - \zeta)} \left[ \langle S_{\bar{x} \zeta} S_{\bar{y}} \rangle - \langle S_{\bar{x} \zeta} Q_{\bar{y} \bar{x} \bar{y}} \rangle \right] \right\} ,
\] (6)

which involves the effective coupling \( \bar{\alpha}_s \, \sigma^2 \).

Applying a large-\( N_c \) mean field approximation [21, 22] \( \langle S S \rangle \to \langle S \rangle \langle S \rangle, \langle S Q \rangle \to \langle S \rangle \langle Q \rangle \), and assuming \( \langle Q \rangle \sim
FIG. 2: The average dipole $S$-matrix as a function of dipole size $r$ at different evolution times $t = \bar{\alpha}_s Y = 0, \ldots, 6$. Solid lines correspond to BK-JIMWLK ($\sigma^2 = 1$) evolution while dashed lines correspond to modified flow at $\sigma^2 = 3$ (here with the coupling rescaled as $\bar{\alpha}_s \rightarrow \bar{\alpha}_s/6$). These curves are numerical solutions of eq. (4). The horizontal line at $S(r) = 1 - 1/e \approx 0.63$ indicates the transition to the non-linear regime. Units for $r$ are given by the RMS color charge density $g^2 \mu$ of the MV model.

$\langle S \rangle^2$ 1, we find that the zero field fixed point is linearly unstable just like for standard BK evolution [5–7, 21, 22] (however, the eigenvalue is different). On the other hand, the unitarity limit is asymptotically (linearly) stable and attractive. However, while the BK/JIMWLK flow to the $\langle S \rangle = 0$ fixed point exhibits “repulsion” at quadratic order in $\langle S \rangle$, in eq. (6) the repulsion is pushed to cubic order. Figure 1 compares the evolution of the dipole $S$-matrix at fixed $\vec{r} = \vec{x} - \vec{y}$ for BK-JIMWLK vs. our modified flow, starting from MV model [25–27] initial conditions where $\langle S \rangle$ is real.

We have rescaled the coupling constant, which is a free parameter in the LL evolution equation, to approximately match the two curves. Hence, increasing $\sigma^2$ from its value $\sigma^2 = 1$ in QCD one is able to modify the pre-asymptotic RG flow to the fixed point corresponding to the unitarity limit.

It is also interesting to note the correlated contribution to dipole pair evolution,

$$\frac{\partial}{\partial Y} \langle S_{\vec{x}\vec{y}} S_{\vec{r}\vec{s}} \rangle - \langle S_{\vec{x}\vec{y}} \rangle \frac{\partial}{\partial Y} \langle S_{\vec{r}\vec{s}} \rangle - \langle S_{\vec{r}\vec{s}} \rangle \frac{\partial}{\partial Y} \langle S_{\vec{x}\vec{y}} \rangle =$$

$$- \bar{\alpha}_s \sigma^2 \int_{\vec{z}} \left( \vec{K}_{\vec{x} \vec{z}} \cdot (\vec{K}_{\vec{s} \vec{r}} - \vec{K}_{\vec{r} \vec{z}}) \right) \left[ Q_{\vec{x}\vec{y}\vec{r}\vec{z}} + Q_{\vec{r}\vec{y}\vec{r}\vec{z}} - S_{\vec{x}\vec{y}\vec{r}\vec{z}} - S_{\vec{x}\vec{y}\vec{r}\vec{z}} \right] ,$$

where $S_{\vec{x}\vec{y}\vec{r}\vec{z}} = \text{tr} V_{\vec{x}} V_{\vec{y}} V_{\vec{r}} V_{\vec{z}} / N_c$ is a trace over six Wilson lines. The r.h.s. does not involve any new operators as compared to the JIMWLK equation; it has just been rescaled to the effective coupling $\bar{\alpha}_s \sigma^2$. This is due to the fact that this contribution arises from the second term in eq. (3), i.e. the one linear in the noise $\xi_{\vec{y}}$, which has the same structure as in JIMWLK evolution. Therefore, the correlated contribution to dipole pair scattering is affected differently by our modification of the flow than the evolution of the $S$-matrix of a single dipole. Two-particle or dijet correlations [28–37] could provide means to constrain evolution at small-$x$.

For $\sigma^2 = 3$ we have checked the “travelling wave” solutions [38, 39] for $\langle S_{\vec{x}\vec{y}} \rangle(Y)$ which lead to geometric scaling of the cross section for scattering of a virtual photon from a proton [40]. In fig. 2 we have again rescaled the coupling constant ($\bar{\alpha}_s \rightarrow \bar{\alpha}_s/\kappa$) so that the speed at $S = 1 - 1/e$ is similar to that for BK-JIMWLK evolution. We do observe a modification of the shape of the travelling waves at intermediate rapidities, deep in the non-linear regime at small $S$. For $\sigma^2 < 1$, the rescaling factor $\kappa$ is < 1. In particular, we confirmed numerically that the evolution of $\langle S_{\vec{x}\vec{y}} \rangle$

1 This does not refer to a naive factorization $\langle Q_{\vec{x}\vec{y}\vec{u}\vec{v}} \rangle \rightarrow \langle S_{\vec{x}\vec{y}} \rangle \langle S_{\vec{u}\vec{v}} \rangle$. Rather, here we use that in the approach to unitarity the typical magnitude of the quadrupole is of order the squared magnitude of the dipole. See refs. [23, 24] for detailed discussions.
“stalls” ($\kappa \to 0$) for a value of $\sigma^2$ slightly less than 1/2.

The modified evolution equation (3, 4) permits phenomenological comparisons of our incorrect evolution (towards unitarity) to the BK/JIMWLK flow of QCD. Ultimately, this would provide observational constraints on the RG flow (bounds on $|1 - \sigma^2|$), i.e. on the approach towards “saturation” of QCD scattering amplitudes. The question we ask is not whether or not unitarity is achieved. Rather, we are interested in constraining from observation the flow towards unitarity to be the one of QCD.

In fig.3 we compare single-inclusive $p_T$-distributions obtained from fixed-coupling ($\alpha_s = 0.1$) evolution at various values for $\sigma^2$; we employ the “hybrid formalism” of ref. [42] to compute the $p_T$ spectrum. Realistic phenomenology of QCD evolution requires one to account for the running of the coupling [43, 44]. Its implementation in B-JIMWLK evolution has been discussed in ref. [45], and references therein. One should also account for NLO corrections to particle production [46–49]. However, our main goal here is not to obtain a good fit to the data but rather to check the sensitivity of the spectra to $\sigma^2$. As $\sigma^2$ was varied we performed no adjustment of the coupling in the evolution equation, $\alpha_s = 0.1$, of the initial saturation scale, $Q_s(0) = 1$ GeV on average over impact parameters, of the form of the initial ensemble of Wilson lines at $Y = 0$ (MV model [25–27]), or of the scales for the proton parton distribution functions [50] and the fragmentation functions [51], $Q^2 = p_T^2$. The figure shows that the single-inclusive spectra in the forward region of pA collisions at RHIC alone do not constrain the evolution to be close to that for QCD.

The modification of the approach to the asymptotic fixed point should affect suitable observables. The wealth of data from HERA, RHIC, LHC, and in the future from the EIC, may make it possible to set quantitative constraints on the flow towards the asymptotic limit of QCD. A challenge for small-$x$ QCD phenomenology is to provide a bound $|1 - \sigma^2|/\sigma^2 \ll 1$.

\[2\] The running coupling prescription suggested in ref. [45] also introduces a quadrupole operator in the evolution of the dipole. However, those are specific NLO corrections associated with running ($\beta(\alpha_s) \sim -\alpha_s^2$) of $\alpha_s$. They do not correspond to a modification of the proper QCD evolution equation like in our scenario.
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**Appendix A: Reweighting**

Here we outline briefly how operator expectation values in the standard B-JIMWLK ensemble can be reconstructed from a simulation with $\sigma^2 \neq 1$ through reweighting.

Consider an operator $O$ given by products of Wilson lines at different transverse points. Upon taking a step in rapidity its expectation value in the standard B-JIMWLK ensemble ($\sigma^2 = 1$) evolves as follows:

$$
\langle O \rangle(Y + dy) = \langle O \rangle(Y) + dy \langle O' \rangle(Y).
$$

(A1)

The operator $O'$ is obtained from $O$ by replacing all Wilson lines in $O$ by the r.h.s. of eq. (3), setting $\sigma^2 = 1$; this is followed by an expansion in $dy$ to linear order.

The brackets $\langle \cdot \rangle$ indicate an average over the noise (in the rapidity slice $Y$):

$$
\langle O' \rangle = \left[ \prod_{\vec{x},a,i} \int d\xi_{ia} P_{\sigma^2=1}(\xi_{ia}) \right] O'(\{\xi_{ia}\}) ; \quad P_{\sigma^2}(\xi_{ia}) = \frac{1}{\sqrt{2\pi \sigma^2}} e^{-\xi_{ia}/2\sigma^2}.
$$

(A2)

With $w(\xi_{ia}) = P_{\sigma^2=1}(\xi_{ia})/P_{\sigma^2}(\xi_{ia})$ we can also write this in the form

$$
\langle O' \rangle = \left[ \prod_{\vec{x},a,i} \int d\xi_{ia} w(\xi_{ia}) P_{\sigma^2}(\xi_{ia}) \right] O'(\{\xi_{ia}\}) = \langle O' w \rangle_{\sigma^2}.
$$

(A3)

Hence, $\langle O' \rangle$ can be computed by averaging over noise with variance $\sigma^2 \neq 1$ provided that one multiplies $O'$ by the weight $w = \prod_{\vec{x},a,i} w(\xi_{ia})$. However, we repeat that $O'$ must be obtained via the r.h.s. of eq. (3) with $\sigma^2 = 1$.

In general it is not possible to compute the averages over the noise exactly. Rather, one must employ Monte-Carlo importance sampling. Eq. (A2) will then lead to much more accurate results than eq. (A3) since the weight $w$ will fluctuate strongly from configuration to configuration [unless $|\sigma^2 - 1| \lesssim (A_L d_A)^{-1/2}$, where $d_A = N_c^2 - 1$ and $A_L$ is the area, i.e. the number of sites, of the transverse lattice].

Nevertheless, reweighted averages via eq. (A3) can be useful for some applications. For example, consider computing the expectation value of $O'$ over a biased ensemble, where one is interested in selecting rare evolution trajectories. In the standard B-JIMWLK ensemble this corresponds to evaluating $\langle O' b \rangle$ where $b$ denotes a bias such as high gluon multiplicity or mean transverse momentum, for which a simple model has been considered in refs. [52–54]. A simpler example in the present context would be a bias of the form $b = \prod_{\vec{x},a,i} \Theta(\Xi - |\xi_{ia}|)$. If $\Xi \gg 1$ ($\Xi \ll 1$) this bias evidently prefers a wider (narrower) distribution for the noise which is better sampled with $P_{\sigma^2=1}(\xi_{ia})$ than with $P_{\sigma^2}(\xi_{ia})$. A more physical example is

$$
b = \exp \left\{ \frac{\alpha_s}{N_c^4} \frac{1}{A_L} \int \frac{d^2 q}{(2\pi)^2} \left[ A_n^i(q) A_n^i(-q) - \frac{1}{q^2} q^i A_n^i(\vec{q}) q^j A_n^j(-\vec{q}) \right]^2 \right\},
$$

(A4)

which suppresses longitudinal and enhances transverse light-cone gauge fields (if $\lambda > 0$)

$$
A_n(x) = \frac{1}{ig} V_n^i \partial_i V_x.
$$

(A5)

Again, if $b$ selects rare evolution trajectories one should expect it to exhibit very large fluctuations across configurations. This would lead to large errors for $\langle O' b \rangle$ since this average would be completely dominated by a small subset of
configurations. Computing, instead, evolution trajectories for $\sigma^2 \neq 1$, and reweighted averages $\langle O^* b w \rangle_{\sigma^2}$, could in some cases increase the overlap with the desired ensemble, i.e. if the product $b w$ is approximately constant for a greater set of configurations. For some biases $b$ like the one in eq. (A4) it would be beneficial to employ correlated (in $\vec{x}$) and non-diagonal (in color) noise.

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