I. INTRODUCTION

An account of jet-quenching in AdS/CFT based on classical string trajectories has been developed in [1,2] and a number of subsequent works has been published with a parallel line of development starting in [3]. The approach of [1] was to model an energetic gluon as a string in AdS5−S5—Schwarzschild with both ends passing through the horizon. Light quarks could be modeled analogously by having a string with one endpoint ending on a D7-brane in the bulk of AdS5. The complementary approach of [2] was to model a light quark-anti quark pair by an initially pointlike open string created close to the boundary with endpoints that are free to fly apart. In either case, the string extends in a direction parallel to the boundary as it falls toward the black hole horizon. In all three works [1,2,3], it was found that the maximum distance that an energetic probe can travel for a fixed energy $E$ in a thermal $\mathcal{N} = 4$ SYM plasma at a temperature $T$ scales as $\Delta r_{\text{max}} \propto E^{1/3}T^{-4/3}$. The constant of proportionality is important for phenomenological applications as it determines the overall strength of jet quenching.

To compute the observables such as the nuclear modification factor $R_{AA}$ and the elliptic flow parameter $v_2$ of light hadrons, we need to know the details of the instantaneous energy loss of light quarks. To tackle this problem, a general formula for computing the instantaneous energy loss in non-stationary string configurations was developed [4], using methods related to the earlier work [5] but with somewhat different results. The application of this formula of [4] to the case of falling strings requires a precise definition of the energy loss (i.e., roughly speaking, what part of the string is to be considered as the “jet” and what part as the thermalized energy to which the jet energy is being lost) and is susceptible to the details of the initial conditions. Nevertheless, studies have shown some rather universal qualitative features of the light quark energy loss [3], including a modified Bragg peak and a seemingly linear path dependence which made it look very similar to the radiative pQCD energy loss.

Using these results for the stopping distance and the path dependence of the energy loss, one can perform the simplest constructions of the nuclear modification factor $R_{AA}$ [6]. The comparison with the LHC pion suppression data showed that, although it had the right qualitative structure, the overall magnitude was too low, indicating that the predicted jet quenching was too strong. Introducing higher derivative corrections to the AdS5 showed substantial increase in $R_{AA}$, but this effect alone was not enough to get close to the data. This is all illustrated in Fig. 1, where the dashed red line shows how far below the data this best-case scenario ($\lambda = 1$ and including the higher derivative corrections) is, and the dashed green curve shows how much we would effectively need to decrease the coupling (all the way to about $\lambda = 0.01$) to come close to the data.

A novel step towards addressing this as well as some other issues, and ultimately towards a more realistic description of energetic quarks, was taken recently by introduction of finite momentum at the endpoints of the string [5]. Typically, standard boundary conditions for open strings require vanishing endpoint momentum, but if one thinks of the endpoints as representing energetic quarks and the string between them representing the color field they generate, then a string with most of its energy packed into its endpoints provides a more natural holographic dual to a pair of quarks that have undergone a hard scattering event. In this way, one also obtains a clear distinction between the energy in the hard probe and energy contained in the color fields surrounding it, hence offering a clear definition of the instantaneous jet energy loss that was missing in earlier accounts. Another important feature of this proposal is that the distance the finite momentum endpoints travel for a given energy is greater than in the previous treatments of the falling strings. In other words, quark jets that these strings
II. ENERGY LOSS

In this section we will develop a phenomenologically usable form of the instantaneous energy loss \(dE/dx\) based on the finite endpoint momentum framework. As shown in [8], a direct consequence of having finite endpoint momentum is that the trajectories of the endpoints are piecewise null geodesics in \(AdS_5\) along which the endpoint momentum evolves according to equations that do not depend on the bulk shape of the string:

\[
\frac{dE}{dx} = -\frac{\sqrt{\lambda}}{2\pi} \frac{\sqrt{f(z^*)}}{z^4},
\]

(2.1)

where \(\sqrt{\lambda} = L^2/\alpha'\) is the 't Hooft coupling, \(f(z) = 1 - z^4/z_H^4\) (in this coordinate system the boundary is at \(z = 0\)), \(z_H = 1/(\pi T)\) and \(z_*\) is the minimal (inverse) radial coordinate the geodesic reaches and which hence completely determines the motion of the endpoint.

As mentioned before, considering endpoints as energetic quarks themselves and the string as the color field they generate, we will identify the rate at which the energy gets drained from the endpoint with the energy loss of an energetic quark. It is worth pointing out that (2.1) is a unique answer, independent of the initial conditions: it does not depend on the energy stored in the endpoint [9] and it is a function of only the radial coordinate \(z\) at which the endpoint is located and weakly dependent on the \(z_*\) of the geodesic along which the endpoint is moving.

To express \(dE/dx\) as a function of \(x\), we need to solve the null geodesic equation. Assuming that initially, at \(x = 0\), the endpoint is at \(z = z_0\) going towards the boundary, we have:

\[
x_{\text{geo}}(z) = \frac{z_H^2}{z} 2F_1 \left( \frac{1}{4}, \frac{5}{4}, \frac{z^4}{z_H^4} \right) - \frac{z_H^2}{z_0} 2F_1 \left( \frac{1}{4}, \frac{5}{4}, \frac{z_*^4}{z_0^4} \right),
\]

(2.2)

where \(2F_1\) is the ordinary hypergeometric function. One could now numerically invert this relation (for given \(z_*\) and \(z_0\)) to obtain \(z(x)\) and plug it in (2.1) to obtain \(dE/dx\) as a function of \(x\) which would result in a characteristic bell-shaped curve for energy loss [8]. However, (2.2) has a particularly simple and universal form for small \(z_*\):

\[
x_{\text{geo}}(z) = \frac{z_H^2}{z_0^4} \left[ \left( \frac{1}{z} - \frac{1}{z_0} \right) + \mathcal{O} \left( \frac{z_*^4}{10^2 z^5}, \frac{z_*^4}{10^5 z_0^8} \right) \right].
\]

(2.3)

The reason we are interested in this expansion is phenomenological: from (2.1) we see that if we start at \(z\) close to the boundary, the energy loss will be large,
which means that the jets dual to these endpoints will be quenched quickly and hence won’t be observable. Therefore, for observable jets, we need to start rather close to the horizon, and for $z_* < z$ we see that the expansion \(^{(2.3)}\) is strongly convergent, resulting in an interesting novel universal form for energy loss:

$$\frac{dE}{dx} = -\frac{\pi}{2} \sqrt{xT^2} \left( \frac{1}{\tilde{z}_0} + \pi T x \right)^2, \tag{2.4}$$

where $\tilde{z}_0 \equiv \pi T z_0 \in [0, 1]$. This form of energy loss has an interesting physical interpretation: at small $x$, it looks like a pure $\sim T^2$ energy loss, similar to the pQCD elastic energy loss (with a running coupling); for intermediate $x$, it looks like $\sim x T^3$ with a path dependence (but not the energy dependence) similar to the pQCD radiative energy loss; and, finally, for large $x$, it has a novel $\sim x^2 T^4$ behavior. The size of $\tilde{z}_0$ (i.e. how much above the horizon the endpoint starts) dictates at what $x$ each of these regimes becomes relevant. This is an interesting (and a very specific) generalization of the simpler “abc” models of energy loss \([10]\), where $dE/dx \propto E^a x^b T^c$.

## III. CONFORMAL $R_{AA}$

In this section we will use the proposed formula \((2.4)\) for the energy loss to compute $R_{AA}$ for pions at RHIC and LHC. There are several steps that will be taken in order to compute a more realistic $R_{AA}$, the purpose of which is to imitate some of the features of QCD:

- The first step is to express all the energy variables in GeV’s and length variables in femtometers.
- To account for roughly three times more degrees of freedom in $\mathcal{N} = 4$ SYM than in QCD, we will relate the temperatures via \([11]\):

$$T_{\text{SYM}} = 3^{-1/4} T_{\text{QCD}}. \tag{3.1}$$

- The next step is to promote a constant $T_{\text{QCD}}$ to a Glauber-like $T_{\text{QCD}}(\vec{x}_\perp, t, \phi)$.
- We will introduce the transverse expansion via a simple blast wave dilation factor \([12]\):

$$r_{\text{bl}}(t) = \sqrt{1 + \left( \frac{v_T t}{R} \right)^2}, \tag{3.2}$$

where $R$ is the mean nuclear radius. We will take the transverse velocity $v_T = 0.6$. The effect of this dilation factor will be to replace $\rho_{\text{part}}(\vec{x}_\perp) \rightarrow \rho_{\text{part}}(\vec{x}_\perp/r_{\text{bl}})/(v_T^2 r_{\text{bl}}^2)$ in the Glauber model.
- Finally, we will use the fragmentation functions \([13]\) to obtain the pionic $R_{AA}$ from the partonic one (neglecting the gluon contribution) \([14]\).

We will use the standard optical Glauber model to compute the participant and binary collisions densities, include the effects of longitudinal expansion and model the spacetime evolution of the temperature. If a jet is created at position $\vec{x}_\perp$ in the transverse plane at time $t = 0$ and moves radially at angle $\phi$, then the local temperature that it sees at some later time $t$ is given by

$$T(\vec{x}_\perp, t, \phi) = \left[ \frac{\pi^2}{\zeta(3)} \frac{1}{16 + n_f N_{\text{part}}} \frac{dN/dy}{\rho_{\text{part}}(\vec{x}_\perp + t \hat{c}(\phi))} \right]^{1/3}, \tag{3.3}$$

where $n_f$ is the number of active flavors (which we will take to be 3), $dN/dy$ is the multiplicity and $t_i$ is the plasma equilibration (formation) time (which will be typically between $0.5$ and $1$ fm/$c$). For this jet we can find, using the energy loss formula \((2.4)\), its initial energy $p_{T,i}$, provided it has a fixed final energy $p_{T,f}$ at the time when the temperature reaches the freezeout temperature $T_{\text{freeze}}$. Averaging the ratio of the initial production spectra $d\sigma/dp_T$ \([15,16]\] at final and initial energies over the transverse plane we get the nuclear modification factor:

$$\left\langle R_{AA}^\text{p} \right\rangle_{p_{T,f}} = \int d^2\vec{x}_\perp \frac{T_{AA}(\vec{x}_\perp, t, \phi)}{N_{\text{bin}}} \frac{d\sigma/dp_T(p_{T,i}(p_{T,f}))}{d\sigma/dp_T(p_{T,f})}, \tag{3.4}$$

where $T_{AA}$ is the number density of binary collisions and $N_{\text{bin}}$ is the total number of binary collisions. For $\phi = 0$ we obtain $R_{AA}^\text{p}$ from \((3.4)\), while for $\phi = \pi/2$ we get $R_{AA}^\text{p}$. For nucleon-nucleon inelastic cross sections we will use $\sigma_{NN} = 42$ mb and $\sigma_{NN} = 63$ mb \([17]\). We will also use the mixed participant-binary number scaling of the multiplicities in the non-central collisions with $85\%$ of participant scaling and $15\%$ of binary number scaling. For charged multiplicities we use $dN_{\text{RHIC,ch}}/d\eta = 700$ and $dN_{\text{LHC,ch}}/d\eta = 1584$.

The results are shown in Fig. 2 where we chose (as in the rest of the plots in the paper) $\tilde{z}_0 = 1$. There we see that, first of all, qualitatively, as expected, our $R_{AA}$ calculations seem to match the data well. To obtain a satisfactory quantitative fit, with a reasonable choice of parameters, we needed to choose $\lambda = 3$ at RHIC; however, using the same parameters and $\lambda$ at the LHC shows that the data is severely underpredicted. Lowering $\lambda$ to 1 for LHC data is not enough: one seems to need a radically small $\lambda = 0.25$ to obtain a satisfactory fit. Of course, with that $\lambda$, RHIC is then severely overpredicted. This is precisely the “surprising transparency” of the LHC \([10]\), where the effects of temperature increase from RHIC to LHC affect the $R_{AA}$ much more than the competing increase of the production spectra \([18]\).

## IV. HIGHER DERIVATIVE CORRECTIONS

A possible way to make our setup more realistic is to add higher derivative $R^2$-corrections to the gravity sector of $AdS_5$, which are the leading $1/N_c$ corrections in the presence of a $D7$-brane. It has been shown \([9]\) that these
type of corrections can increase \( R_{AA} \) significantly and in this section we will explore their effect in the context of finite endpoint momentum strings.

We will model the \( R^2 \) corrections by a Gauss-Bonnet term, i.e., we will consider the action of the form:

\[
S = \frac{1}{2k_5^2} \int d^4x \sqrt{-G} \left[ R + \frac{12}{L^2} + L^2 \frac{\lambda_{GB}}{2} \left( R^2 - 4R_{\mu\nu}^2 + R_{\mu\nu\rho\sigma}^2 \right) \right],
\]

where \( \lambda_{GB} \) is a dimensionless parameter, constrained by causality \cite{20} and positive-definiteness of the boundary energy density \cite{21} to be:

\[
-\frac{7}{36} < \lambda_{GB} \leq \frac{9}{100}.
\]

A black hole solution in this case is known analytically \cite{22}:

\[
ds^2 = \frac{L^2}{z^2} \left( -a^2 f_{GB}(z) dt^2 + dx^2 + \frac{dz^2}{f_{GB}(z)} \right),
\]

where

\[
f_{GB}(z) = \frac{1}{2\lambda_{GB}} \left( 1 - \sqrt{1 - 4\lambda_{GB}(1 - z^4/z_H^4)} \right),
\]

\[
a^2 = \frac{1}{2} \left( 1 + \sqrt{1 - 4\lambda_{GB}} \right).
\]

The 't Hooft coupling and the temperature are given by

\[
\lambda = a^2 \frac{L^2}{\alpha'}, \quad T = \frac{a}{\pi z_H}.
\]

Using the same procedure as before, we can easily find the energy loss from the finite endpoint momentum in geometry \cite{3} and obtain a formula similar to \cite{2}:

\[
d\mathcal{E}_{GB} = \frac{\sqrt{\lambda}}{2\pi^2} \frac{f_{GB}(z)}{z} \frac{dx_{geo}(z)}{dz}.
\]

To express this as a function of \( x \), we need to solve for null geodesics:

\[
dx_{geo} = \pm \frac{1}{\sqrt{f_{GB}(z)-f_{GB}(z)}}.
\]

To obtain the generalization of the formulas from the previous section, we will send \( z \to 0 \) here and then, for a given \( \lambda_{GB} \), we can easily numerically integrate \cite{4} and invert to obtain \( z(x) \), which can then be plugged in \cite{4} to get \( d\mathcal{E}/dx \) as a function of \( x \) and \( T \).

However, since according to \cite{4}, \( \lambda_{GB} \) is constrained to be small, these expressions are suitable for a perturbative expansion in \( \lambda_{GB} \), allowing for a more practical analytic expression. To do this, we will expand \cite{4} in \( \lambda_{GB} \) up to some order \( n \) and neglect all terms higher than \( 1/z^2 \), as they are \( O(z^4) \) subleading. Of course, we will be able to check how accurate this is by comparing with the full numerical solution. We define a polynomial in \( \lambda_{GB} \):

\[
P_n(\lambda_{GB}) = \frac{2}{z_H^2} \lim_{z \to 0} z^2 \left( \frac{dx_{geo}(n)}{dz} \right)_{z=0}.
\]

where \( n \) denotes the order of expansion in \( \lambda_{GB} \). In this case, we can easily solve the geodesic equation:

\[
z_n(x, \lambda_{GB}) = \frac{z_H^2 z_0 P_n(\lambda_{GB})}{z_H^2 P_n(\lambda_{GB}) - 2x z_0}.
\]

This can be plugged in \cite{4} to get explicitly the form of \( d\mathcal{E}/dx \) for a given order \( n \), yielding an expression similar
to (2.4):
\[
\frac{dE_{\text{GB}}}{dx} = -\sqrt{\lambda}T^2F_{\text{n}}(\lambda_{\text{GB}}) \left( \frac{G_n(\lambda_{\text{GB}})}{z_0} + \pi T x \right)^2.
\]  
(4.10)

The functions \(F_n\) and \(G_n\) are functions of \(\lambda_{\text{GB}}\) only and do not have a particularly illuminating explicit form, even for small \(n\). For \(\lambda_{\text{GB}}\) as large as \(-7/36\), by comparing to the all-order numerical result, we found that it is enough to go to \(n = 5\) order in expansion.

In the left plot of Fig. 3 we compare this energy loss to the energy loss without the Gauss-Bonnet term, where we can see that, at a maximally negative \(\lambda_{\text{GB}}\), the energy loss with the Gauss-Bonnet corrections can be up to two times smaller. This, expectedly, has noticeable consequences for the \(R_{AA}\) results in greater stopping distances. Application of this formula, using endpoints that start close to the horizon, offers a clear definition of the energy loss that is independent of the details of the string configuration and results in greater uncertainties. Note that our energy loss formula (2.4) has a strong sensitivity to the temperature, \(dE/dx \sim \sqrt{\lambda}T^3\) or even \(\sqrt{\lambda}T^4\). Hence, even a small change in the temperature, \(T \to \kappa T\) can have the same effect as a large change in the coupling, \(\lambda \to \kappa^2 \lambda\) or \(\kappa \lambda\).

We cannot offer at the moment a concrete physical reason that would justify the possibility of temperature uncertainties, but we can speculate based on some very general arguments. From the perspective of the temperature formula (3.3) we see that one would expect the LHC to be roughly 30% hotter than RHIC, based on the ratio of the multiplicities. However, if the initial time \(t_i\) in the two cases is different, then the jet effectively feels a cooler or a hotter medium, according to (3.3). This is precisely what was suggested in [23], where the authors used a bigger initial time at the LHC than at RHIC where \(t_t = 0.6\) fm/c [25], based on the requirements of the hydrodynamic simulations to fit the low \(p_T\) elliptic flow data. We see the effect of this in the left plot of Fig. 3, where changing \(t_i\) from 1 fm/c (blue) to 0.6 fm/c (red) for the case of \(\lambda = 1\) and \(\lambda_{\text{GB}} = -0.2\) (which fits the central LHC data) leads to a significant decrease in \(R_{AA}\). Additionally, if we allow \(\eta/s\) to decrease (relative to LHC, where the temperature range is higher), which means increasing \(\lambda_{\text{GB}}\) (as per (4.11)), we can approach the RHIC data even more (yellow curve). We should note that this is just an illustration of the effect of the decrease of \(\eta/s\) on \(R_{AA}\), as the same hydrodynamic calculations of [23] and [24] suggest that this decrease is not so strong. If we want to keep the same \(\eta/s\) at RHIC as it was at the LHC (meaning keeping \(\lambda_{\text{GB}} = -0.2\) then we can get close to the data by increasing the coupling approximately 4 times (green curve). If we keep these parameters of the green curve and pass onto LHC (right plot of Fig. 5), where we set \(t_i = 1\) fm, we see that the curve is below the data (blue curve), but lowering the overall LHC temperature by about 10% we are able to approach the data (red curve).

VI. CONCLUSIONS

In this Letter we have proposed a novel formula (2.4) for the instantaneous energy loss of light quarks in a strongly coupled SYM plasma. This formula was derived in the framework of finite endpoint momentum strings, where the jet energy loss is identified with the energy flux from the endpoint to the bulk of the string, offering a clear definition of the energy loss that is independent of the details of the string configuration and results in greater stopping distances. Application of this formula, using endpoints that start close to the hori-
FIG. 3. **Left:** Ratio of the instantaneous energy loss in pure AdS \(^4\) and the energy loss with the Gauss-Bonnet corrections \([4,10]\) as a function of \(x\), for \(z_0 = 1\) and for several different values of \(\lambda_{GB}\). **Right:** Nuclear suppression factor \(R_{AA}\) at the LHC for \(\lambda = 1\), with and without the higher derivative corrections. All other parameters are the same as in Fig. 2.

FIG. 4. Nuclear modification factor at RHIC in non-central collisions (left plot) and the elliptic flow parameter at the LHC (right plot). The experimental data for both RHIC \([19]\) and LHC \([24]\) are for the 20-30 % centrality class. In the left plot we compare the \(R_{AA}\) in central collisions as well as the in and out \(R_{AA}\) in non-central collisions to our calculations for \(b = 3\) fm and \(b = 7\) fm, respectively. In the right plot, the band corresponds to the \(v_2\) calculations for \(b\) between 7 and 9 fm. All the other parameters in these plots are the same as in Fig. 2.

...zerson ("shooting" strings) and including the higher derivative \(R^2\)-corrections showed, independently, a very good match with the RHIC and LHC central \(R_{AA}\) data for light hadrons, and even partially for the elliptic flow \(v_2\). A consistent simultaneous match of both the RHIC and the LHC central \(R_{AA}\) data remains challenging, but, as we argue, the temperature sensitivity of our formula coupled with the uncertainties in the formation time \(\xi\) and the shear viscosity \(\eta/s\) at RHIC and LHC may enable us to reconcile these differences. In particular, we have shown that, using a smaller formation time at RHIC and perhaps allowing for a slightly smaller \(\eta/s\) than at the LHC, one can significantly reduce the RHIC-LHC splitting. Further inclusion of non-conformal effects, which are known to moderately increase the energy loss at lower temperatures (and hence affect RHIC more than the LHC), may provide an additional reduction of the splitting.

**ACKNOWLEDGMENTS**

The work of A.F. and M.G. was supported by U.S. DOE Nuclear Science Grant No. DE-FG02-93ER40764. The work of S.S.G. was supported in part by the Department of Energy under Grant No. DE-FG02-91ER40671.
FIG. 5. **Left:** Nuclear modification factor at RHIC in central collisions for different choices of the 't Hooft coupling $\lambda$, dimensionless $\lambda_{GB}$ and the initial time $t_i$. **Right:** Nuclear modification factor at the LHC in central collisions with and without the temperature adjustment $T \to \kappa T$.

[1] S. S. Gubser, D. R. Gulotta, S. S. Pufu, and F. D. Rocha, JHEP 0810, 052 (2008), arXiv:0803.1470 [hep-th].
[2] P. M. Chesler, K. Jensen, and A. Karch, Phys. Rev. D79, 025021 (2009), arXiv:0804.3110 [hep-th].
[3] Y. Hatta, E. Iancu, and A. H. Mueller, JHEP 05, 037 (2008), arXiv:0803.2481 [hep-th].
[4] A. Ficnar, Phys. Rev. D86, 046010 (2012), arXiv:1201.1780 [hep-th].
[5] P. M. Chesler, K. Jensen, A. Karch, and L. G. Yaffe, Phys. Rev. D79, 125015 (2009), arXiv:0810.1985 [hep-th].
[6] A. Ficnar, J. Noronha, and M. Gyulassy, Nucl. Phys. A910-911, 252 (2013), arXiv:1208.0305 [hep-ph].
[7] S. Chatrchyan et al. (CMS Collaboration), Eur. Phys. J. C72, 1945 (2012), arXiv:1202.2554 [nucl-ex].
[8] A. Ficnar and S. S. Gubser, (2013), arXiv:1306.6648 [hep-th].

[9] In [8], $z_e$ was related to the initial energy of the endpoint $E_0$ by saying that the energy at the moment when the endpoint reaches the horizon must be zero. The reason behind this proposal was that this prescription maximized the stopping distance for a given energy, without the string performing a “snap-back”. However, in general, it only matters that the string does not perform a snap-back within some phenomenologically relevant distance $L$ so for a given $z_e$, we can have a multitude of energies.

[10] W. Horowitz and M. Gyulassy, Nucl. Phys. A872, 265 (2011), arXiv:1104.4958 [hep-ph].
[11] S. S. Gubser, Phys. Rev. D76, 126003 (2007), arXiv:hep-th/0611272 [hep-th].
[12] B. Betz and M. Gyulassy, (2013), arXiv:1305.6458 [nucl-th].
[13] B. A. Kniehl, G. Kramer, and B. Potter, Nucl. Phys. B582, 514 (2000), arXiv:hep-ph/0010289 [hep-ph].
[14] Another more natural feature of shooting strings that differs from the falling strings is that the virtuality of the endpoint, $Q^2 \equiv p_{\perp}^2 + p_z^2$, is proportional to the endpoint’s energy squared, $Q^2 = E^2 z_t^2 / (1 - z_t^2)$. This energy, and hence virtuality as well, decreases even during the “ascending” phase in the geodesic trajectory, in a unique way given by (2.1) and (2.2). However, at finite $N_c$, one should bear in mind that even bulk constructions can be off-shell; thus some $N_c$ suppressed contribution to virtuality may be significant compared to the rather suppressed classical expression just given. For this reason, we will use the usual prescription of $Q = p_T$, and note that, due to $a$, for our purposes, low sensitivity of the fragmentation functions to the $Q^2$-evolution, we do not expect that a choice of a different prescription would affect our results significantly.

[15] M. Cacciari, P. Nason, and R. Vogt, Phys. Rev. Lett. 95, 122001 (2005), arXiv:hep-ph/0502203 [hep-ph].
[16] M. L. Mangano, P. Nason, and G. Ridolfi, Nucl. Phys. B373, 295 (1992).
[17] W.-T. Deng, X.-N. Wang, and R. Xu, Phys. Rev. C83, 014915 (2011), arXiv:1008.1841 [hep-ph].
[18] This may be partially resolved by considering duals of non-conformal field theories where the coupling gets an effective temperature running. For example, by introducing a specific potential for the dilaton, one can construct a bottom-up non-conformal deformation of $N = 4$ SYM [26, 27] that has the same thermodynamics (and Polyakov loops) as the one provided by lattice QCD. In those models the running of the dilaton causes the energy loss at low temperatures to increase relative to the conformal limit, which in turn affects the $R_{AA}$ at RHIC more than at the LHC, but this effect was not strong enough to resolve the problem entirely.

[19] A. Adare et al. (PHENIX Collaboration), (2012), arXiv:1208.2254 [nucl-ex].
[20] M. Brigante, H. Liu, R. C. Myers, S. Shenker, and S. Yaida, Phys. Rev. D77, 126006 (2008), arXiv:0712.0805 [hep-th].
[21] D. M. Holman and J. Maldacena, JHEP 0805, 012 (2008), arXiv:0803.1467 [hep-th].
[22] R.-G. Cai, Phys.Rev. D65, 084014 (2002), arXiv:hep-th/0109133 [hep-th].

[23] H. Song, S. Bass, and U. W. Heinz, (2013), arXiv:1311.0157 [nucl-th].

[24] S. Chatrchyan et al. (CMS Collaboration), Phys.Rev.Lett. 109, 022301 (2012), arXiv:1204.1850 [nucl-ex].

[25] H. Song, S. A. Bass, U. Heinz, T. Hirano, and C. Shen, Phys.Rev. C83, 054910 (2011), arXiv:1101.4638 [nucl-th].

[26] S. S. Gubser, A. Nellore, S. S. Pufu, and F. D. Rocha, Phys.Rev.Lett. 101, 131601 (2008) arXiv:0804.1950 [hep-th].

[27] A. Ficnar, J. Noronha, and M. Gyulassy, J.Phys. G38, 124176 (2011), arXiv:1106.6303 [hep-ph].