EPJ Web of Conferences **137**, 07023 (2017) DOI: 10.1051/epjconf/201713707023
XIIth Quark Confinement & the Hadron Spectrum

Early hydrodynamisation, energy loss and small systems in holographic heavy ion collisions

Wilke van der Schee[1,2,a]

[1] *Center for Theoretical Physics, MIT, Cambridge, MA 02139, USA*
[2] *Institute for Theoretical Physics,Utrecht University, Leuvenlaan 4, 3584 CE Utrecht, The Netherlands*

Abstract. There has been much recent progress in describing heavy ion collisions using strongly coupled methods. We start by introducing a few caveats in applying holography to QCD, after which we will highlight recent progress on the initial stage, by simulating shock collisions in the dual gravitational theory. Secondly, we describe an attempt at describing jet evolution through strongly coupled plasma, where we specifically also included the width of the jet, and the interplay of this width with the energy lost as the jet propagates through the plasma. Lastly, there has been much recent discussion about the formation of quark-gluon plasma in small systems at low multiplicies, such as produced in proton-nucleus or even proton-proton collisions. At infinitely strong coupling it is possible to estimate that this may be possible, provided that the particle multiplicity satisfies $dN_{tot}/dy \gtrsim 4$.

1 Outline

Care must always be taken when interpreting results from holography in relation to heavy ion collisions. Especially around the QCD confinement/deconfinement transition pure $\mathcal{N} = 4$ SYM is different from QCD, although much progress can be made by considering more realistic holographic models including a running of the coupling constant and a confinement/deconfinement transition [1] . At high energy scales a direct comparison is also complicated, since asymptotic freedom implies that the coupling becomes weak at (very) high energies. The coupling, however, runs only logarithmically with energy, and hence at LHC energy scales there is always a need of including some strongly coupled physics.

It is hence a sensible approach to model heavy ion collisions at weak coupling, using perturbative QCD, as well as to model the collisions at strong coupling, using holography. Combining insights from both approaches can then lead to a realistic understanding of heavy ion collisions. Here, we will give a recent overview of holographic attempts to describe the initial stage as well as the propagation of jets. Lastly, we comment on recent excitement where flow-like signals were observed in small systems at relatively low multiplicities.

[a] e-mail: wilke@mit.edu

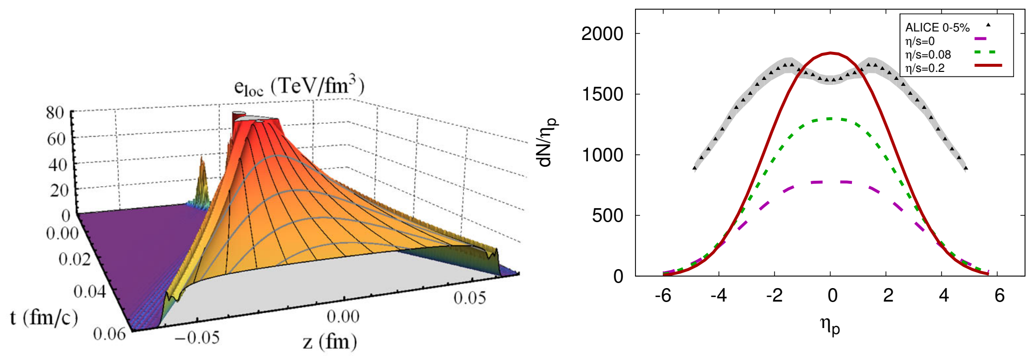

Figure 1. We show the energy density in the local rest frame as produced by high energy shock wave collisions (left). On the right we show the pseudo-rapidity distribution following from the holographic universal rapidity distribution. The experimental profile is about 50% wider than the data, which is perhaps natural at infinitely strong coupling. We note that both experimental and theoretical profiles are Gaussian to a good approximation when plotted as a function of rapidity, instead of pseudo-rapidity. (Fig. from [4])

2 Fast hydrodynamisation and a Gaussian rapidity profile

The work in this section was done in collaboration with Jorge Casalderrey-Solana, Paul Chesler, Michal Heller, Niki Kilbertus, David Mateos, Bjorn Schenke and Miqual Triana, as presented in [2–5].

One of the prime accomplishments of holography is the convincing demonstration that at strong coupling there are well understood settings whereby the system starts out very far-from-equilibrium, but nevertheless is well described by viscous hydrodynamics within a time of approximately the inverse temperature *at that time* [6–8]. This process is now called hydrodynamisation, and it is important to stress that this normally happens much before the system thermalises, i.e.much before the pressures are equal and the hydrodynamic viscous contribution is small. Much of this intuition has been gained by sourcing the system in a specific way [6], or by starting with generic initial conditions [7, 8].

In terms of heavy ion collisions more can be learnt by starting as an initial condition with two colliding gravitational shock waves [9, 10]. In the dual gauge theory these correspond to bits of plasma boosted to the speed of light, while keeping their energy fixed. Before the collision these are exact solutions of the Einstein equations, both for planar and non-planar shock waves, and they can have an arbitrary energy profile in the longitudinal direction. It is appealing that when restricting to pure Einstein gravity the initial conditions in the bulk are completely fixed when the initial energy distribution at a time before the collision is known.

Here we focus on initial planar shock waves with a Gaussian energy profile in the longitudinal direction: $T_{\pm\pm} = \frac{N_c^2}{2\pi^2} m^4 e^{-(t\pm z)^2/2w^2} / \sqrt{2\pi w^2}$, with w the width of the shock and m^3 the rescaled energy density per transverse area. Most prominently, these collisions hydrodynamise even more rapidly than the previous studies, within a time of $0.25/T$ at mid-rapidity [2, 11]. Perhaps more interestingly, it is also possible to compute the rapidity profile of the hydrodynamic initial conditions, which led to a universal Gaussian rapidity, with width 0.96 in the limit of high energy collisions [5]. Perhaps unsurprisingly, although indeed both at RHIC [12] and LHC [13] approximately Gaussian rapidity profiles are found (with width 2.27 and 3.86 at $\sqrt{s_{NN}}$ of 0.2 and 2.76 TeV respectively), even after hydrodynamical evolution and freeze-out the experimental widths of both experimental RHIC and

DOI: 10.1051/epjconf/201713707023

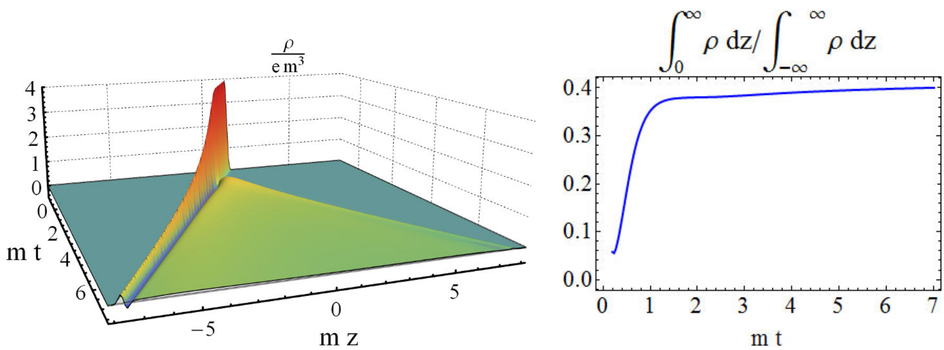

Figure 2. In a similar set-up as in Fig. 1 we gave the shock moving towards negative z a globally conserved charged, while keeping the other shock neutral. The transparency and decay of the charge on the lightcone behaves analogous to the energy density. On the right panel we show the fraction of the charge at $z > 0$ versus time, all in the centre-of-mass frame. Interestingly, due to the strong interactions present we see that 41% of the charge changes direction and ends up at positive z at late times. (Fig. from [3])

LHC profiles are roughly 50% wider than the profile obtained by the shock collisions (Fig. 1 (right), as a function of pseudorapidity). Of course this indicates that infinitely strong coupling produces too much stopping, something we will come back to in the discussion. Recently, this framework has been extended to a full five dimensional numerical computation, without assuming any symmetry [14], which allows to study much richer dynamics.

A new quantitative insight can be gained by giving the shock wave charge, as can be done by adding a $U(1)$ gauge field to the bulk spacetime [3]. This could be thought of as modelling baryon charge, and we find rather similar results regarding the hydrodynamisation time as well as the rapidity profile. Interestingly, by including charge it is possible to collide a charged shock with a neutral shock, and thereby determine which part of the charge in the plasma came from either the left- or right-moving shock wave. We find that at strong coupling in the centre-off-mass frame 41% percent of the charge on the right of the collision point is actually from the left-moving shock, which clearly demonstrates the strong interactions at play (see Fig. 2).

3 Interplay between jet width and jet energy loss

The work in this section was done in collaboration with Krishna Rajagopal and Andrey Sadofyev, as presented in [15].

A complementary approach to study the dynamics of quark-gluon plasma is to trace highly energetic partons as they move through the plasma. While the initial production of these partons is typically well described by methods within perturbative QCD, the propagation through the plasma involves softer energy scales, such that the evolution of those virtual partons requires strongly coupled physics. It is therefore interesting to gain intuition what would happen to jets when they propagate through an infinitely strongly coupled plasma in $\mathcal{N} = 4$ SYM-theory, which can be reliably computed using holography.

Within holography it is indeed possible to introduce quarks transforming in the fundamental representation of $SU(N_c)$, which is done by introducing (probe) D7 branes in the bulk AdS spacetime (either space-filling or upto a finite AdS coordinate for massless or massive quarks respectively).

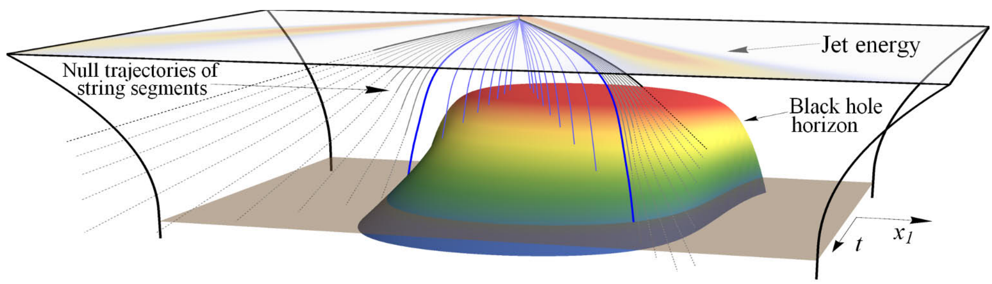

Figure 3. This figure illustrates a dijet traversing a cooling and expanding plasma, as modelled with bits of string following null geodesics. The left-moving jet starts narrower and traverses a shorter distance so it loses less energy (the blue bits of string). It is also clear that the right-moving string curves down more strongly, which will translate in a wider final jet.

Quark-antiquark pairs then correspond to fundamental strings as described by the Nambu-Goto action, which in the planar limit at strong coupling propagate classically without breaking [16]. Conveniently, after a short initial time the strings are described by null strings: each part of the string moves along an independent null geodesics through the bulk space-time. It is hence only necessary to specify the direction of each part of the string as an initial condition after which the future evolution can be easily determined by solving the geodesic equation. Perhaps the most interesting part of this exercise is to determine the part of the string that falls into the black hole, which is interpreted as energy loss.

There have been several studies employing such a scheme in order to compute observables such as the ratio of jets produced in nucleus-nucleus compared to proton-proton collisions, as measured by R_{AA} [17–19][1]. These studies, however, only matched the initial energy of the produced jets, but lacked a more detailed description, such as the width of a jet. This is perhaps particularly unrealistic, since most studies used initial string configurations constructed to minimise the energy loss. More diverse ensembles of string configurations would hence have more varied properties, which will also affect the energy loss.

In this new study we tried to ameliorate this problem, by introducing a more realistic ensemble of string configurations, with a jet width distribution $C_1^{(1)}$, defined through the variable

$$C_1^{(\alpha)} \equiv \sum_{i,j} z_i z_j \left(\frac{|\theta_{ij}|}{R}\right)^{\alpha}. \tag{1}$$

We obtained the initial distribution of these jet widths from perturbative QCD [21]. In this first model we assumed that $C_1^{(1)} = a\,\sigma_0$, with a an unknown constant and σ_0 the angle determining how steeply the endpoint of the string falls into the bulk, which agrees with the intuition that as matter moves deeper in the bulk it leaves an imprint in the field theory over a wider region (scale-radius duality).

Interestingly, as was previously found in [22], at strong coupling there is a rich interplay between the energy loss of a jet and the width of the jet: wide jets lose more energy. This is intuitive, since as strings fall down steeper into the black hole space time, more energy falls into the black hole and hence more energy is lost. Secondly, it is also apparent that all jets individually become broader, since the curvature of the geometry increases the angle σ_0 and hence the jet width.

[1]The study [17] employs a specific and perhaps natural initial condition where all the energy is initially in the end-point [20], akin to the Lund string model. These strings however take a much longer time to be accurately described as a null string.

The black hole geometry we evolved our ensemble of strings through is specified by the temperature profile, where we take [17]

$$T = b\left[\frac{dN_{\rm ch}}{dy}\frac{\rho_{\rm part}(\vec{x_\perp}/r_{\rm bl}(t))}{N_{\rm part}(t+t_i)\, r_{\rm bl}(t)^2}\right]^{1/3}, \tag{2}$$

where $\rho_{\rm part}(\vec{x_\perp})$ is the participant density, as given by an optical Glauber model, $N_{\rm part}$ gives the total participants, $dN_{\rm ch}/dy \approx 1870$ [13] the particle multiplicity at mid-rapidity at LHC and $r_{\rm bl}(t) = \sqrt{1+(v_T t/R)^2}$, with $v_T = 0.6$ and $R = 6.7$ fm. For QCD the constant b is a measure of the multiplicity per entropy S, which for $S/N_{\rm ch} = 7.5$ [23, 24] would give $b = 0.782$. Here, however, we take b as a free parameter, which can be thought of as incorporating the fact that QCD has fewer degrees of freedom and weaker coupling than infinitely strongly coupled SYM.

Fig. 3 shows R_{AA} for five different sets of our free parameters (a, b), chosen such that the R_{AA} is in rough agreement with data (shown to guide the eye). We stress that the computed R_{AA} can not be directly compared with data, since the model does not include effects such as hadronisation and jet finding. Interestingly, even though the R_{AA} is very similar for all combinations (a, b) we see that the $C_1^{(1)}$ distribution, as shown in Fig. 4 (right) show different qualitative behaviour. For larger a we see that narrow jets get significantly wider as compared with vacuum jets, which is expected since each individual jet gets wider. For smaller a, on the other hand, it is apparent that we find fewer wide jets. This second effect can be explained by realising that the jet spectrum is a steeply falling function of energy, falling roughly as E_{jet}^{-6}, and that wide jets lose more energy. The combination means that wider jets end up in a lower energy bin, where they do not contribute significantly due to the steeply falling energy spectrum. As can be seen, in this model either effect can dominate depending on the free parameters (a, b). Interestingly, it is hence entirely possible that the mean value of $C_1^{(1)}$ in a specified energy range decreases (as experiments indicate [25]), even though every individual jet in the ensemble widens in angle.

Of particular interest related to this work is an similar study at weak coupling [26], where it was also found that the width of a jet crucially influences the jet energy loss. Secondly, an exciting measurement [27, 28] of the modification of jet substructure in quark-gluon plasma indicates that individual partons within a jet lose a varying amount of energy, even when their path length is identical. This is indeed consistent with our picture that a fluctuating jet width or virtuality is essential for energy loss.

4 Small systems

Recently there has been a lot of excitement about measurements in proton-nucleus and even proton-proton collisions [29, 30], where similar (flow) signals were found that were previously found to be convincing evidence for the formation of quark-gluon plasma in nucleus-nucleus collisions. Whereas the community has gotten used to the surprisingly successful description by hydrodynamics for plasmas of about 10 fm diameter, it was never expected that a similar hydrodynamic picture would emerge from systems only 1 fm in size. Given these new signals this now seems nevertheless a reasonable avenue to explore, as it is hard to explain long-range rapidity correlations at equal transverse angle (near-side ridge) without hydrodynamics since they originate from very early times by causality (see however [31] for a quantum explanation within the glasma picture).

We would also like to point out that this evidence is not conclusive for the existence of quark-gluon plasma and in fact there is the observation that the 4-particle cumulant, called $c_{2(4)}$, can have a positive sign for systems with low multiplicity, although this depends somewhat on the method used [32]. This is thought to be incompatible with a standard hydrodynamic + freeze-out prescription. Also,

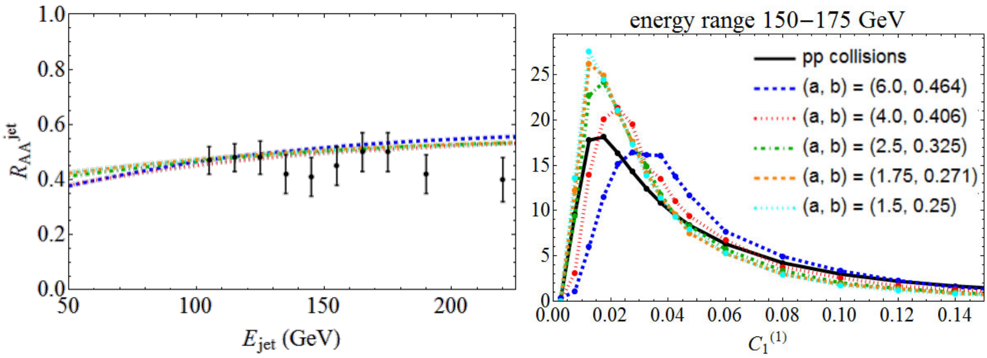

Figure 4. For five sets of our free parameters a and b (see text) we show the jet ratio R_{AA} (left), as well as the $C_1^{(1)}$ distributions, which are to be compared with the proton-proton input from [21]. There are two effects visible, where either effect can dominate depending on our choice of (a, b). The first effect is that each jet gets wider and hence we find fewer narrow jets. The second effect is that wider jets lose more energy, and with a steeply falling jet versus energy spectrum this means that fewer wide jets are found in each energy bin.

if a plasma really forms it would be expected that jets, as described in section 3, suffer an energy loss comparable with the energy loss observed in nucleus-nucleus collisions at similar multiplicity. Instead, the data seems consistent with zero energy loss in these small systems [33].

Nevertheless, it is an interesting possibility to work out if it is theoretically possible that plasma is indeed formed in those small systems. At infinitely strong coupling we will argue that this is indeed feasible, though we see three (related) important theoretical constraints:

1. The plasma should have enough time to hydrodynamise
2. The plasma should behave collectively, so it should be bigger than the thermal wavelength
3. Both non-hydrodynamic modes and higher order hydrodynamic modes that have unknown transport coefficients should be small

The first constraint is perhaps the easiest, as we have already reviewed that shock collisions lead to hydrodynamisation times of $0.25/T$. At strong coupling the second constraint has first been studied in [34], where in a boost-invariant context it was asked how small a fluctuation could be while still being described by hydrodynamics. It was concluded that a size of 0.5 fm was feasible, which was later refined to $1/T$ by an explicit computation using 5-dimensional numerical general relativity [35]. The second and third constraint are studied in [36–39], where it was concluded that hydrodynamics breaks down for momentum modes of $|\mathbf{k}| > 4.5T$, which translates into a size $\lambda \simeq 2\pi/|\mathbf{k}| > 1.4/T$. [2]

So perhaps unsurprisingly all three constraints give a similar outcome. Interestingly, this can easily be converted to a multiplicity estimate, using that $N_{tot} \approx S_{tot}/5.0$ [23, 24], and that entropy is approximately conserved after hydrodynamics becomes valid. Since $S_{tot} = s\,V \approx 16T^3V$ (using lattice [40]), it follows that for a transverse size of $R \gtrsim 1/T_{\rm ini}$ (note that we take this condition initially since R is expected to grow faster than $1/T$) and volume $R^2\tau$ that $dN_{tot}/dy \gtrsim 3.2\,\tau_{\rm ini}T_{\rm ini} \approx 4$, where we estimated that $\tau_{\rm ini}T_{\rm ini} \approx (0.5\,{\rm fm}/c)(500\text{ MeV}) \approx 1.25$.

[2] Perhaps complementary, Ref. [38] argued that hydrodynamics cannot be said to work if it is sensitive to 2nd order transport coefficients, which were assumed to be poorly constrained. This was tested phenomenologically, and it was found that for $dN_{ch}/d\eta < 2$ the results depended sensitively on higher order transport coefficients.

Of course we just provided what we think is an absolute lower bound: if $dN_{tot}/dy < 4$ then we should never expect hydrodynamic behaviour. It is however a difficult problem to convincingly distinguish hydrodynamic from non-hydrodynamic behaviour, especially as far-from-equilibrium evolution can also develop flow [34, 41, 42]. This type of flow even builds up during free streaming, even though in free streaming one naturally has no multi-particle correlations. To be able to distinguish a hydrodynamic and non-hydrodynamic evolution it is hence necessary to accurately understand the hydrodynamic initial conditions and evolution to see if this matches the experimental results.

It is important to stress that the discussion above in principle does not involve the actual numbers of particles to be large in order for hydrodynamics to be valid; instead, at high temperature the plasma is strongly coupled and has no quasi-particles, such that the relevant question is if the thermal wavelength is smaller than the size of the plasma. On the other hand, even at strong coupling, a constraint on the number of particles is expected, since the estimate of a minimum size of $1/T$ only holds for a large number of colours N_c, which is equivalent to a large entropy since the entropy scales as N_c^2. It would therefore be interesting to see how the estimate $1/T$ would be affected by $1/N_c$ corrections.

5 Discussion

We hope to have conveyed an updated status on a strongly coupled approach in heavy ion collisions, having particularly reviewed the hydrodynamisation of strongly coupled plasma, and the evolution of jets through such a plasma. Strongly coupled systems hydrodynamise quickly, perhaps within a time of $0.25/T$, with T the temperature at hydrodynamisation. For collisions this gives a Gaussian rapidity profile, which has width 0.96 in the high energy infinitely strongly coupled limit.

Naturally, our computations have not been performed in QCD, which in particular does not have infinitely strong coupling and is neither scale invariant. It would therefore be very interesting to see how all the results above change, at least at a qualitative level, when running coupling effects or a breaking of scale invariance are included [3]. Also interesting would be a better understanding when including finite coupling corrections. Promising starts have recently been made by considering higher-derivative (curvature squared) theories [37, 44, 45], indicating slightly longer hydrodynamisation times, as well as more energy on the lightcone and less entropy production.

Acknowledgments

This overview owes many thanks to all collaborators of the work presented above, in particular Jorge Casalderrey-Solana, Paul Chesler, Michal Heller, David Mateos, Scott Pratt, Paul Romatschke, Krishna Rajagopal, Andrey Sadofyev, Bjorn Schenke and Miquel Triana. This work is supported by the U.S. Department of Energy under grant Contract Number DE-SC0011090.

References

[1] U. Gursoy, E. Kiritsis, JHEP **0802**, 032 (2008), 0707.1324

[2] J. Casalderrey-Solana, M.P. Heller, D. Mateos, W. van der Schee, Phys. Rev. Lett. 111, **181601** (2013), 1305.4919

[3] J. Casalderrey-Solana, D. Mateos, W. van der Schee, M. Triana, JHEP **09**, 108 (2016), 1607.05273

[3] The running coupling is expected to make a particularly important difference for the scaling of the entropy versus the collision energy, we thank Elias Kiritsis for stressing this again during the conference [43].

[4] W. van der Schee, B. Schenke, Phys. Rev. **C92**, 064907 (2015), 1507.08195
[5] P.M. Chesler, N. Kilbertus, W. van der Schee, JHEP **11**, 135 (2015), 1507.02548
[6] P.M. Chesler, L.G. Yaffe, Phys.Rev.Lett. **102**, 211601 (2009), 0812.2053
[7] M.P. Heller, R.A. Janik, P. Witaszczyk, Phys.Rev.Lett. **108**, 201602 (2012), 1103.3452
[8] M.P. Heller, D. Mateos, W. van der Schee, D. Trancanelli, Phys.Rev.Lett. **108**, 191601 (2012), 1202.0981
[9] R.A. Janik, R.B. Peschanski, Phys.Rev. **D73**, 045013 (2006), hep-th/0512162
[10] D. Grumiller, P. Romatschke, JHEP **0808**, 027 (2008), 0803.3226
[11] P.M. Chesler, L.G. Yaffe, Phys.Rev.Lett. **106**, 021601 (2011), 1011.3562
[12] I. Bearden et al. (BRAHMS Collaboration), Phys.Rev.Lett. **94**, 162301 (2005), nucl-ex/0403050
[13] E. Abbas et al. (ALICE), Phys. Lett. **B726**, 610 (2013), 1304.0347
[14] P.M. Chesler, L.G. Yaffe, JHEP **10**, 070 (2015), 1501.04644
[15] K. Rajagopal, A.V. Sadofyev, W. van der Schee, Phys. Rev. Lett. **116**, 211603 (2016), 1602.04187
[16] A. Karch, E. Katz, JHEP **0206**, 043 (2002), hep-th/0205236
[17] A. Ficnar, S.S. Gubser, M. Gyulassy, Phys. Lett. **B738**, 464 (2014), 1311.6160
[18] J. Casalderrey-Solana, D.C. Gulhan, J.G. Milhano, D. Pablos, K. Rajagopal, JHEP **1410**, 19 (2014), 1405.3864
[19] R. Morad, W.A. Horowitz, JHEP **11**, 017 (2014), 1409.7545
[20] A. Ficnar, S.S. Gubser, Phys.Rev. **D89**, 026002 (2014), 1306.6648
[21] A.J. Larkoski, S. Marzani, G. Soyez, J. Thaler, JHEP **05**, 146 (2014), 1402.2657
[22] P.M. Chesler, K. Rajagopal, Phys.Rev. **D90**, 025033 (2014), 1402.6756
[23] B. Muller, K. Rajagopal, Eur.Phys.J. **C43**, 15 (2005), hep-ph/0502174
[24] S.S. Gubser, S.S. Pufu, A. Yarom, Phys.Rev. **D78**, 066014 (2008), 0805.1551
[25] S. Chatrchyan et al. (CMS), Phys. Lett. **B730**, 243 (2014), 1310.0878
[26] J.G. Milhano, K.C. Zapp, Eur. Phys. J. **C76**, 288 (2016), 1512.08107
[27] Collaboration (CMS) (2016), CMS-PAS-HIN-16-006
[28] G. Aad et al. (ATLAS), Phys. Lett. **B751**, 376 (2015), 1506.08656
[29] V. Khachatryan et al. (CMS), Submitted to: Physics Letters B (2016), 1606.06198
[30] M. Aaboud et al. (ATLAS) (2016), 1609.06213
[31] K. Dusling, R. Venugopalan, Phys. Rev. **D87**, 094034 (2013), 1302.7018
[32] Collaboration (ATLAS) (2016), ATLAS-CONF-2016-106
[33] Collaboration (ATLAS) (2016), ATLAS-CONF-2016-108
[34] W. van der Schee, Phys.Rev. **D87**, 061901 (2013), 1211.2218
[35] P.M. Chesler, JHEP **03**, 146 (2016), 1601.01583
[36] P. Romatschke, Eur. Phys. J. **C76**, 352 (2016), 1512.02641
[37] S. Grozdanov, N. Kaplis, A.O. Starinets (2016), 1605.02173
[38] P. Romatschke (2016), 1609.02820
[39] M.P. Heller, A. Kurkela, M. Spalinski (2016), 1609.04803
[40] M. Cheng, N. Christ, S. Datta, J. van der Heide, C. Jung et al., Phys.Rev. **D77**, 014511 (2008), 0710.0354
[41] J. Vredevoogd, S. Pratt, Phys.Rev. **C79**, 044915 (2009), 0810.4325
[42] W. van der Schee, P. Romatschke, S. Pratt, Phys.Rev.Lett. **111**, 222302 (2013), 1307.2539

[43] E. Kiritsis, A. Taliotis, JHEP **1204**, 065 (2012), 1111.1931
[44] T. Andrade, J. Casalderrey-Solana, A. Ficnar (2016), 1610.08987
[45] S. Grozdanov, W. van der Schee (2016), 1610.08976