Gauge theory plasmas from AdS/CFT

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We review the AdS/CFT description of gauge theory plasmas. We discuss the low shear viscosity, the jet quenching, and the J/ψ-suppression, which are three major signatures for the quark-gluon plasma observed at RHIC experiments.

1 Introduction

In this note, we review the connection between AdS/CFT duality and quark-plasma experiments at RHIC (see Ref. [1] for a review of QGP physics). RHIC stands for Relativistic Heavy Ion Collider at Brookhaven National Laboratory. The name "heavy ion" comes from the fact that it collides heavy ions such as gold nuclei ^{197}Au instead of usual e^+e^-, pp or $p\bar{p}$. The goal of the experiment is to realize the deconfinement transition and form the quark-gluon plasma (QGP). In principle, it should be possible to form QGP if one has high enough temperature or high enough density. However, it is not an easy job to confirm QGP formation because of the following problems: First, what one observes is not QGP itself but only the by-products after hadoronization, and one has to infer what had happened from the by-products. Second, those secondary particles are mostly strongly-interacting, and the perturbative QCD is not very reliable for the current and near-future experimental temperatures.

To resolve these problems, many attempts are made to identify the generic signatures of QGP. Some of the generic signatures discussed to date are as follows:

1. The elliptic flow which may be the consequence of very low viscosity of QGP
2. The jet quenching
3. J/Ψ-suppression

All of these signatures have been discussed in the AdS/CFT duality, so I review recent developments explaining these phenomena.

Let me first remind you of the duality. The original AdS/CFT duality is [2, 3, 4] (see Ref. [5] for a review)

$$\mathcal{N}=4 \text{ SYM} \quad \leftrightarrow \quad \text{type IIB string theory on } \mathrm{AdS}_5 \times S^5. \tag{1}$$

We will use the finite temperature version of the duality and its cousins:

$$\begin{aligned}&\mathcal{N}=4 \text{ SYM at finite temperature} \leftrightarrow \\ &\quad \text{type IIB string theory} \\ &\quad \text{on Schwarzschild-AdS}_5 \text{ black holes} \\ &\quad (\mathrm{SAdS}_5) \times S^5.\end{aligned}$$

The black hole appears on the right-hand side because a black hole is also a thermal system due to the Hawking radiation. As with the original AdS/CFT duality, this duality can be motivated from the near-horizon limit of the D3-brane.

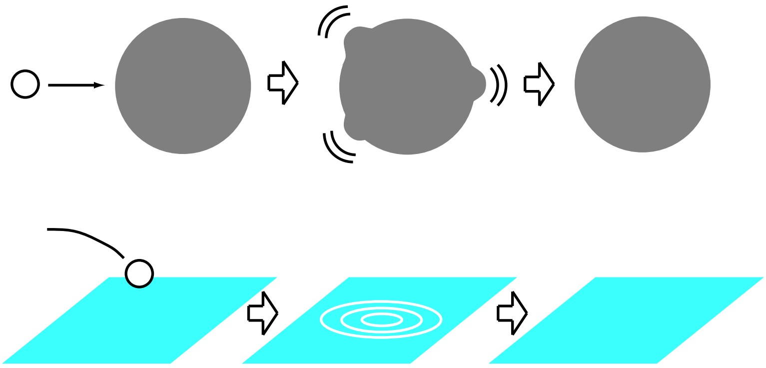

Figure 1: When one adds a perturbation to a black hole, the black hole behavior is similar to a hydrodynamic system. In hydrodynamics, this is a consequence of the viscosity.

2 Black holes and hydrodynamics

According to the RHIC experiment, QGP behaves like a liquid. AdS/CFT then implies that a black hole also behaves like a liquid. Then, the plasma viscosity should be calculable from black holes. In fact, black holes and hydrodynamic systems behave similarly. Consider adding a perturbation to a black hole, *e.g.*, drop some object (Fig. 1). Then, the shape of the black hole horizon becomes irregular, but such a perturbation decays quickly, and the black hole returns to the original symmetric shape. The no-hair theorem is one way to see this. According to the theorem, the stationary black hole is unique and symmetric. Thus, the perturbed black hole cannot be stable. If you regard this as a diffusion, the diffusion occurs since the perturbation is absorbed by the black hole.

This behavior is very similar to a liquid. Suppose that one drops a ball in a water pond. Then, you generate surface waves, but they decay quickly, and the water pond returns to a state of stable equilibrium. In hydrodynamics, this is a consequence of the viscosity. Thus, one can consider the notion of viscosity for black holes as well. And the "viscosity" for black holes should be calculable by considering the above process.

Let me remind you of freshman physics of viscosity. As a simple example, consider a fluid between two plates and move the upper plate with velocity v (Fig. 2). Then, the fluid is dragged and the lower plate experiences a force. This force is the manifestation of the viscosity. In this case, the force F the lower plate experiences per unit area is given by

$$\frac{F}{A} = \eta \frac{v}{L} \ . \tag{2}$$

The proportionality constant η is called the (shear) viscosity.

Microscopically, the viscosity arises due to the momentum transfer between molecules. Figure 2 shows a close-up view of the fluid and I put an artificial boundary to divide the fluid into two parts. The molecules collide each

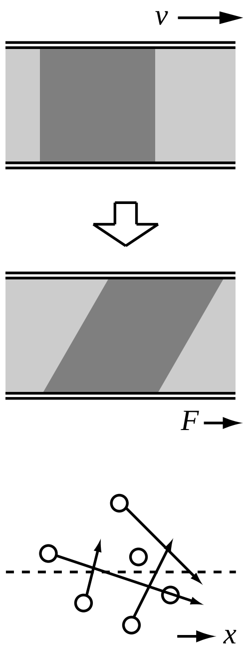

Figure 2: When one moves the upper plate, the fluid is dragged due to the viscosity and the lower plate experiences a force. The bottom figure is the close-up view of the fluid.

other and are exchanged randomly through the boundary. But in the situation where you move the upper plate, the molecules in the upper-half part, on average, have more momentum in the x-direction than the ones in the lower-half part. These molecules are exchanged, which means that momentum in the x-direction is transported through the boundary.

Going back to the black hole, how can one calculate the plasma viscosity? I first give a brief explanation and its implications. Then, I justify the claim more in detail.

2.1 Quick argument

There are many ways to compute the viscosity and I explain one simple method. To do so, let me go back to the D3-brane from the SAdS. In the gravity side, the diffusion occurs by black hole absorption. So, it is natural to associate the shear viscosity with the absorption cross section by black holes. (I explain this point more in detail later. The detailed argument suggests that this is the cross section for the graviton of particular polarization.) Now, there is a general theorem on black holes [7] to state that the cross section $\sigma_{\rm BH}$ is equal to the horizon area A for a broad range of black holes[1] :

$$\eta \propto \lim_{\omega \to 0} \sigma_{\rm BH} = A \ . \tag{3}$$

But the horizon area is the famous quantity, namely it represents the black hole entropy, so it must be the plasma entropy [2]

$$S_{\rm BH} = \frac{A}{4G\hbar} k_{\rm B} \tag{4}$$

($k_{\rm B}$: Boltzmann constant). Then, the shear viscosity divided by the entropy becomes constant.[3] The constant can be determined from the argument later and the result is

$$\frac{\eta}{s} = \frac{\hbar}{4\pi k_{\rm B}} \ . \tag{5}$$

This value is very small. In comparison, η/s for water is about 3×10^3 under normal circumstances.

[1] As a simple example, this is true for the usual Schwarzschild black hole as well. Precisely speaking, the theorem applies only to the low energy limit $\omega \to 0$.

[2] We use the unit $c = 1$ in this lecture.

[3] Precisely speaking, one divides by the entropy density s. We consider black holes with infinite extension, so the entropy itself diverges. The area for the absorption cross section should be understood in a similar way.

Now, the point is that all the relations we used (cross section versus horizon area, black hole entropy versus horizon area) are generic, so the result must be universal as well. Namely, it does not depend on the details of gauge theories. So, the claim is [13]

Gauge theory plasmas which have gravity duals have a universal low value of η/s at large 't Hooft coupling (at zero chemical potential).[4]

This is a rather indirect argument, but this claim has indeed been checked for many known gravity duals.

This result is very important, so let me rephrase in a different way (Fig. 3). Gauge theories of which we can actually compute the shear viscosity are supersymmetric gauge theories, not the real QCD. We compute the shear viscosity from black holes, but the gravity dual of QCD is not known. So, one cannot use AdS/CFT directly to compute QCD properties. However, as we saw, the quantity corresponding to the shear viscosity is universal on the black hole side, so one can immediately apply the $\mathcal{N}=4$ result to the real QCD even though these two theories are completely different.

In fact, RHIC suggests that QGP has a very low viscosity and the estimated value [9, 10]

$$\frac{\eta}{s} \sim 0.1 \times \frac{\hbar}{k_{\mathrm{B}}}? \tag{6}$$

is very close to the above AdS/CFT value. One important point is that the temperature in question is still order of Λ_{QCD}. At this range of temperature, QCD is still strongly coupled and pQCD is not very reliable. (In fact, the naive extrapolation of the weak coupling result gives a larger value for η/s.) Thus, the AdS/CFT duality which predicts the strong coupling behavior may be useful to analyze QGP.

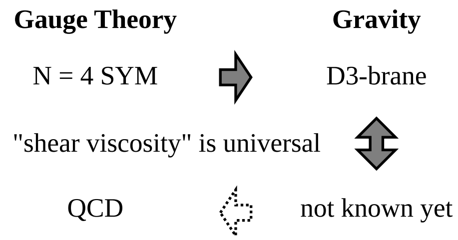

Figure 3: Gauge theories described by D-branes are mostly super-Yang-Mills theories, and not the real QCD. However, the black hole quantity corresponding to the "shear viscosity" is universal, so probably the results for super-Yang-Mills theories are directly applicable to the real QCD.

2.2 More in detail

I now describe the relation between the shear viscosity and the absorption cross section. To do so, one first has to understand the interactions of bulk and boundary fields. These fields can interact at the boundary. For example, the graviton produces the back-to-back scattering of gluons (Fig. 4). The relevant interactions are

$$S_{\mathrm{int}} \sim \int d^4x\, \delta\phi F_{\mu\nu}^2 + \delta h^{\mu\nu} T_{\mu\nu} + \cdots \tag{7}$$

(ϕ: dilaton, $h^{\mu\nu}$: graviton, $F_{\mu\nu}$ and $T_{\mu\nu}$: the field strength and the energy-momentum tensor of the gauge theory). Such interactions can be obtained by expanding the D-brane action (DBI-action) around the expectation values of

[4] The assumption of zero chemical potential arises from a somewhat technical assumption I did not mention.

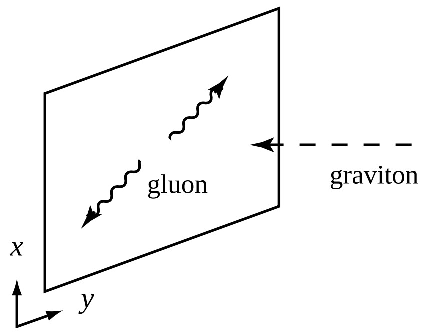

Figure 4: The bulk graviton produces the back-to-back scattering of gluons on the boundary. In black hole picture, it is natural to regard the graviton decay rate as the absorption cross section by the black hole.

the bulk fields. The DBI-action with at most two derivatives contains the following term:

$$\mathcal{L} \sim e^{-\phi} F_{\mu\nu}^2 + \cdots \ . \tag{8}$$

If ϕ is constant, one gets the standard gauge theory Lagrangian with $g_{\rm YM}^2 \sim g_s$, where $g_s := e^{\phi}$. If ϕ fluctuates, the action is expanded as

$$\mathcal{L} \sim e^{-\langle\phi\rangle} F_{\mu\nu}^2 + \delta\phi F_{\mu\nu}^2 \ , \tag{9}$$

where $\phi = \langle\phi\rangle - \delta\phi$. The second term is nothing but the first term in Eq. (7).

This fact could be stated as

Bulk field fluctuations act as sources of boundary fields

I just rephrase the same statement, but this leads to the so-called "GKP-Witten relation," which is the definition of the AdS/CFT.

Given the interaction term, one can easily calculate the graviton decay rate (with polarization h_{xy}, where x and y are directions along the brane) from the standard field theory formula:

$$\begin{aligned}
&\sigma_{\rm QFT} \\
&\quad = \frac{1}{2\omega} \sum_{\text{final states}} \int \frac{d^3p_1}{(2\pi)^3 2\omega_1} \frac{d^3p_2}{(2\pi)^3 2\omega_2} \\
&\qquad \times (2\pi)^4 \delta^4(p_f - p_i) |\mathcal{M}|^2 \\
&\quad = \frac{8\pi G}{\hbar\omega} \int d^4x\, e^{i\omega t} \langle [T_{xy}(t,x), T_{xy}(0,0)] \rangle \ .
\end{aligned}$$

The first equality is just the Fermi's golden rule with matrix element $\mathcal{M}$. The matrix element is proportional to T_{xy}, so the formula is written by a correlator of the energy-momentum tensor. (This is an optical theorem.)

In black hole description, it is natural to regard this decay rate as the absorption cross section of the graviton. This has been checked for the D3-brane at zero temperature [11], so [5]

$$\sigma_{\rm BH} = \frac{8\pi G}{\hbar\omega} \int d^4x\, e^{i\omega t} \langle [T_{xy}(t,x), T_{xy}(0,0)] \rangle \ . \tag{10}$$

We can use this relation to compute the shear viscosity since the viscosity is given by a Kubo

[5] There is a caution on Eq. (10). The relation is known to hold for the D3-brane at zero temperature but is unclear for generic AdS black holes. The problem is that a black hole is the strong coupling limit of a gauge theory, so Eq. (10) is the strong coupling statement. But, of course, it is not easy to compute the right-hand side in gauge theory. One can compute it in the weak coupling, but as we saw, weak coupling results differ from strong coupling results in general. In the special case of the zero-temperature D3-brane, such a comparison actually makes sense due to the nonrenormalization theorem, and this is the case where Eq. (10) has been checked. But it is an open question if this is also true for generic AdS black holes. Here, we simply assume that the relation holds at strong coupling.

formula microscopically:

$$\eta = \lim_{\omega \to 0} \frac{1}{2\hbar\omega} \int d^4x\, e^{i\omega t} \langle [T_{xy}(t, x), T_{xy}(0, 0)] \rangle \ . \tag{11}$$

This formula has the same form as the absorption cross section, so one can immediately write the viscosity in terms of the cross section:

$$\eta = \frac{\lim_{\omega \to 0} \sigma_{\text{BH}}}{16\pi G} \ . \tag{12}$$

We have written all plasma quantities in terms of black hole quantities. Using Eqs. (3) and (4), one obtains

$$\frac{\eta}{s} = \frac{\frac{A}{16\pi G}}{\frac{A}{4G\hbar} k_{\text{B}}} = \frac{\hbar}{4\pi k_{\text{B}}} \ . \tag{13}$$

This is the previous formula (5).

2.3 Chemical potential issue

We saw that gauge theories at strong coupling have a universality of shear viscosity. However, there is an important restriction. There are many proofs of the universality, but all fail in the presence of a chemical potential [8, 12, 13, 14]. So, the natural question is what happens to the universality at finite chemical potential.

Actually, it is not easy to realize the realistic finite density in AdS/CFT, *i.e.*, baryon number density. But there is a simple alternative. Namely, consider charged AdS black holes instead of neutral black holes. A black hole is known to obey thermodynamic-like laws and its first law is written as

$$dM = T dS_{\text{BH}} + \Phi dQ \tag{14}$$

(T: black hole temperature, Φ: electromagnetic potential, Q: black hole charge). As one can see, the electromagnetic field Φ plays a role of a chemical potential.

In AdS/CFT, such a charge arises as follows. As in Eq. (1), the full geometry involves S^5. One can add an angular momentum along S^5, which is known as the "spinning" D3-brane solutions [15, 16]. The angular momentum becomes a Kaluza-Klein charge after the S^5 reduction.

What is the gauge theory interpretation of the charge? The symmetry of S^5 corresponds to an internal symmetry of SYM, R-symmetry $SO(6)$. This $SO(6)$ rotates adjoint scalars in the $\mathcal{N} = 4$ supermultiplet. (These scalars correspond to string oscillations in the transverse directions to the brane. Thus, there are $D - p = 6$ degrees of freedom, and they label S^5.) The R-symmetry group $SO(6)$ is rank 3, so one can add at most three independent charges. The three-charge solution is known as the STU solution [17]. When all charges are equal, the STU solution is the well-known RN-AdS$_5$ (Reissner-Nordström-AdS) black hole. Thus, the charge in question is a $U(1)_R$ charge.

Because the charge corresponds to the $U(1)_R$ charge, this is by no means realistic. However, the theory has interesting features which are common to the real QCD. For instance, the phase diagram is qualitatively similar to the QCD diagram [18, 19].[6] Our system does not represent a realistic chemical potential, but it may mimic the realistic case and one may learn an interesting lesson for gauge theory plasmas at finite density.

The shear viscosity for charged AdS black holes was computed by 4 groups, and the re-

[6]This is the case of the RN-AdS black hole with compact horizon or compact SYM. In this review, we consider the RN-AdS black hole with noncompact horizon, which is always in plasma phase.

sult turns out to be $\eta/s = 1/(4\pi)$ again [20]-[23]. So, the universality seems to hold even at nonzero chemical potential. If this is true for generic chemical potential, $\eta/s = 1/(4\pi)$ may be true even at finite baryon number density.

3 Heavy quarks in medium

Recently, there has been much discussion on heavy quark dynamics in plasma medium. Two applications have been discussed: the issues related to the J/ψ suppression and the issues related to the jet quenching. In this section, I quickly summarize the discussion.

3.1 J/Ψ-suppression

Since J/Ψ is heavy, charm pair production occurs only at the early stages of the nuclear collision. However, if the production occurs in the plasma medium, charmonium formation is suppressed due to the Debye screening. One technical difficulty is that the $c\bar{c}$ pair is not produced at rest relative to the plasma. Therefore, the screening length is expected to be velocity-dependent. Such a computation has been done only for the Abelian plasma [24].

In the AdS/CFT framework, a heavy quark may be realized by a fundamental string which stretches from the asymptotic infinity (or from a "flavor brane") to the black hole horizon. This string transforms as a fundamental representation; In this sense, the string represents a "quark." The fundamental string has an extension and the tension, so the string has a large mass, which means that the string represents a heavy quark.

Such a string has been widely studied in the past to measure the heavy quark potential [25, 26]. For a $q\bar{q}$ pair, two individual strings extending to the boundary is not the lowest energy configuration. Instead, it is energetically favorable to have a single string that connects the pair (Fig. 5). The energy difference is interpreted as a $q\bar{q}$ potential, and one gets a Coulomb-like potential for $\mathcal{N} = 4$.

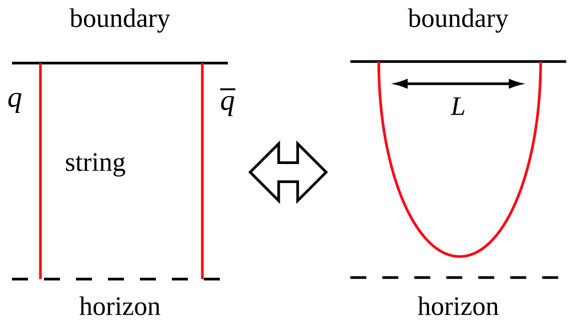

Figure 5: For a $q\bar{q}$ pair, the left configuration is not the lowest energy configuration. Instead, the right configuration is energetically favorable. The energy difference is interpreted as a $q\bar{q}$ potential. At finite temperature, this is no longer true. For large enough separation of the pair, the left configuration becomes favorable. This phenomenon is the AdS/CFT description of the Debye screening.

At finite temperature [27, 28], it is no longer true that a string connecting the $q\bar{q}$ pair is always the lowest energy configuration; for large enough separation of the pair (L_s), isolated strings are favorable energetically. This phenomenon is the AdS/CFT description of the Debye screening and L_s may be interpreted as the screening length.

References [29, 30] proposed how to compute the velocity dependence of the screening length. They computed the screening length in the $q\bar{q}$ rest frame, *i.e.*, they considered the plasma flowing at a velocity v. Such a "plasma wind" is obtained by boosting a black hole.

At the leading order in v, the screening length is obtained as [29, 30, 31, 32]

$$\begin{aligned} &(\text{screening length}) \\ &\propto \quad \frac{1}{\epsilon_0^{1/4}}(1-v^2)^{1/4} \\ &\sim \quad (\text{boosted plasma energy density})^{-1/4} \end{aligned} \tag{15}$$

for the $\mathcal{N}=4$ SYM, where ϵ_0 is the unboosted energy density. Thus, the screening effect becomes stronger than the $v=0$ case. One can also compute the screening length for the other gauge theories [32, 33]. The leading behavior in v seems universal. Namely, if one writes the screening length as

$$(\text{screening length}) \propto (1-v^2)^{\nu} \ , \tag{16}$$

the exponent ν is determined by the speed of sound c_s:

$$4\nu = 1 - \frac{3}{4}(1-3c_s^2) + \cdots \tag{17}$$

when the theory is nearly conformal, *i.e.*, $c_s^2 \sim 1/3$. One can make a simple estimate of the exponent for QCD. According to the lattice results cited in Ref. [34], all groups roughly predict $1/3 - c_s^2 \sim 0.05$ around $2T_c$. Bearing in mind that our results are valid to large-N_c theories and not to QCD, Eq. (17) gives $\nu \sim 0.22$. It would be interesting to compare this number with lattice calculations and experimental results.

One can also study the screening length at finite chemical potential [32]. At the leading order, the screening length at finite chemical potential is the same as the one at zero potential for a given energy density.

3.2 Jet quenching

Another interesting QGP phenomenon is the jet quenching. In the parton hard-scattering, jets are formed, but the jets have to travel in the QGP medium, so the jets are strongly suppressed. This phenomenon is known as the jet quenching. So, the interesting quantity is the energy loss rate of partons. The AdS/CFT descriptions of jet quenching have been proposed recently [35]-[38]. The proposals have been quickly extended to the other gauge theories.[7]

To discuss the jet quenching, one now moves the fundamental string with a velocity v along a brane direction. Then, the momentum carried by the string flows towards the horizon and one interprets the flow as the energy loss rate. For example, the energy loss rate for the $\mathcal{N}=4$ SYM becomes

$$\frac{dp}{dt} = -\frac{\pi}{2}\sqrt{\lambda}T^2\frac{v}{\sqrt{1-v^2}} \ , \tag{18}$$

where $\lambda := g_{\mathrm{YM}}^2 N$ is the 't Hooft coupling.

Unfortunately, the result obtained in this way has some drawbacks. First, the result is not universal and is model-dependent. Second, the black hole results become exact only in the $\lambda \to \infty$ limit, but the result does not have a finite large-λ limit. These two drawbacks are in contrast to the η/s case. One can still try to put the numerical values naively, but the value obtained in this way is not close to the experimentally favored value. It is still too early to draw conclusions, but this may suggest that one has to be careful to apply AdS/CFT to QGP. Namely, we should focus on the universality, and naive extrapolations of $\mathcal{N}=4$ results may not be a good idea.

[7] For an extensive list of literature, see, *e.g.*, Ref. [32].

4 Towards "AdS/QGP"

Hydrodynamic description of gauge theory plasmas using AdS/CFT is very powerful due to the universality. AdS/CFT may be useful to analyze experiments. Conversely, experiments or the other theoretical approaches (such as lattice) may be useful to confirm AdS/CFT. This approach is also important since there are many loose ends on the AdS/CFT derivation. However, one has to be careful to apply AdS/CFT to QGP if the universality does not hold (*e.g.*, jet quenching).

In this talk, I emphasized the universality approach, but of course finding the gravity dual of QCD is desirable. The biggest assumption of the universality argument is that such a dual indeed exists, so finding the dual is necessary for the universality approach as well. One well-known model of QCD is the Sakai-Sugimoto model [39]. This was a successful model in the confining phase, but unfortunately it is not a good model in the plasma phase. The Sakai-Sugimoto model involves the D4-brane compactified on a circle. Above the deconfinement temperature, the model looks as a true five-dimensional theory.

Another problem of the current approach concerns the large-λ limit; AdS/CFT and QGP are actually different limits. The large-λ limit of the AdS/CFT is $g_{\rm YM} \to 0$ and $N_c \to \infty$. But of course QGP has a large-λ since $g_{\rm YM} \sim O(1)$ and $N_c = 3$. We are not taking the same limits, so it is not clear why both give the close results.

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