The shear viscosity of gauge theory plasma with chemical potentials

Paolo Benincasa\textsuperscript{1}, Alex Buchel\textsuperscript{1,2} and Roman Naryshkin\textsuperscript{1,3}

\textsuperscript{1}Department of Applied Mathematics  
University of Western Ontario  
London, Ontario N6A 5B7, Canada  
\textsuperscript{2}Perimeter Institute for Theoretical Physics  
Waterloo, Ontario N2J 2W9, Canada  
\textsuperscript{3}Physics Department  
Taras Shevchenko Kiev National University  
Prosp.Glushkova 6, Kiev 03022, Ukraine

Abstract

We consider strongly coupled gauge theory plasma with conserved global charges that allow for a dual gravitational description. We study the shear viscosity of the gauge theory plasma in the presence of chemical potentials for these charges. Using gauge theory/string theory correspondence we prove that at large ’t Hooft coupling the ratio of the shear viscosity to the entropy density is universal.

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1 Introduction

The gauge theory/string theory correspondence of Maldacena [1,2] provides a valuable insight into the nonperturbative dynamics of strongly coupled gauge theory plasma. Originally, the string theory correspondence was formulated for static properties of strongly coupled gauge theories. It was pointed out in [3] that computing equilibrium two-point correlation functions of stress-energy tensor one can extract the physics of the near-equilibrium (hydrodynamic) description of strongly coupled hot gauge theory plasma. The computations of the finite temperature transport properties of the $\mathcal{N} = 4 SU(N)$ supersymmetric Yang-Mills (SYM) theory plasma at large t’ Hooft coupling\(^1\) [3, 6] were extended to various non-conformal gauge theory plasmas in [7–15]. One of the most surprising results of these computations was the discovery of the universality of the shear viscosity of strongly coupled gauge theory plasma. Specifically, for all gauge theories which allow for a dual supergravity description, and without chemical potentials for conserved global charges (if such charges are present), it was shown that the ratio of the shear viscosity $\eta$ to the entropy density $s$ is a universal constant [16–18]

\[ \frac{\eta}{s} = \frac{1}{4\pi}, \tag{1.1} \]

in units $\hbar = k_B = 1$.

$\mathcal{N} = 4$ SYM has a maximum $U(1)^3 \subset SO(6)_R$ abelian subgroup of the R-symmetry for which one can introduce (at most three) different chemical potentials. At strong coupling and finite temperature, the dual supergravity description of this gauge theory plasma is given by a system of near extremal D3 branes with generically three different angular momenta along the five-sphere $S^5$ [19]\(^2\). This supergravity solution allows for a consistent Kaluza-Klein reduction to five dimensions, where it is described within $\mathcal{N} = 2$ supergravity coupled to two abelian vectors commonly referred to as the STU model [21]. Within this effective five dimensional description, finite temperature $\mathcal{N} = 4$ SYM plasma is dual to non-extremal black holes carrying generically three different $U(1)$ charges corresponding to three different $S^5$ angular momenta of the near-extremal D3 branes in the type IIB supergravity description. Even though $\mathcal{N} = 4$ SYM plasma with chemical potentials violates the assumptions of the universality theorem [16–18], explicit computation of shear viscosity leads to (1.1) [22–24].

\(^1\)Finite t’ Hooft coupling corrections to $\mathcal{N} = 4$ hydrodynamics were discussed in [4, 5].

\(^2\)The misprints in the five-form expression for this supergravity solution were fixed in [20].
Expectation that (1.1) might in fact be more universal than anticipated in [16–18] was strengthened\(^3\) by explicit construction of new models of finite temperature gauge/string theory correspondence with an \(R\)-charge chemical potential [26]. Specifically, it was found that shear viscosity of all strongly coupled \(Y^{p,q}\) quiver gauge theory [27–30] plasmas with a \(U(1)\) R-charge chemical potential satisfies the universal relation (1.1). It was further conjectured in [26] that (1.1) holds true even in the presence of chemical potentials.

In this paper we prove that (1.1) is indeed satisfied for generic strongly coupled gauge theory plasma with global charge chemical potentials. We begin with spelling out explicit assumptions (and their implications) under which we obtain (1.1).

- First, we consider strongly coupled gauge theory plasma that allow for a dual string theory description. This allows us to use the gauge theory/string theory Minkowski-space prescription [31] for computing two-point correlators of the stress-energy tensor. We will work in the regime of large ’t Hooft coupling, where the supergravity approximation to string theory is valid.

- The gauge theory/string theory correspondence relates a \((D+1)\)-dimensional strongly coupled gauge theory on \(\mathbb{R}^{D,1}\) space-time to a particular background of ten-dimensional type IIB supergravity. In all known examples of gauge theory/string theory correspondence it is possible to do a Kaluza-Klein reduction along a compact \((8-D)\)-dimensional manifold and relate \((D+1)\)-dimensional strongly coupled gauge theory to an effective \((D+2)\)-dimensional gauged supergravity, obtained from a consistent truncation of the full ten-dimensional supergravity. However, we do not believe that this is always possible to do. In fact, the strongest form of the universality (1.1) proven in [18] does not rely on this assumption\(^4\). In this paper we assume that our gauge theory plasma is described by a type IIB gravitational background that allow for a consistent truncation to a \((D+2)\)-dimensional gauged supergravity. However, we will not assume that the geometry is asymptotically flat, or that the asymptotically flat region can be “re-attached”. Thus, our universality class of (1.1) is necessarily weaker than that of

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\(^3\)Shear viscosity of the M2-brane plasma also appears to satisfy the universal relation (1.1) [25].

\(^4\)Both proofs [16,17] implicitly assume the existence of a consistent truncation. The proof [17] further implicitly assumes that the background is asymptotically flat, as it relies on the universality of the cross section for the graviton scattering from the black hole horizon [32,33]. The latter universality was derived only for asymptotically flat space-times, and in general it is not known how to “re-attach” the asymptotically flat region in gravitational dual to non-conformal gauge theories (other than those described by flat Dp branes).
the zero chemical potential case considered in [18]. Since we would like to study gauge theory plasma with finite chemical potentials, these gauge theories must have a set of abelian conserved charges — otherwise we simply would not be able to introduce a chemical potential. In the dual supergravity picture, for each conserved $U(1)$ charge, there must be a corresponding $U(1)$ isometry. Introducing a chemical potential for a particular $U(1)$ charge on a gauge theory side results in gauging corresponding $U(1)$ isometry on the supergravity side. In the effective $(D + 2)$-dimensional gravitational description, besides Einstein-Hilbert term and scalar fields, one gets a set of Maxwell fields, one for each gauged isometry (or a chemical potential from the gauge theory perspective). To summarize, we consider strongly coupled $(D + 1)$-dimensional gauge theory plasmas which have a following dual effective $(D + 2)$-dimensional gravitational description

$$S_{D+2} = \frac{1}{16\pi G_{D+2}} \int_{M_{D+2}} d\xi^{D+2} \sqrt{-g} \left\{ R - \mathcal{K}_{\alpha\beta}(\phi)\partial_\mu \phi^\alpha \partial_\nu \phi^\beta - \tau_{ab}(\phi) F^{(a)}_{\mu\nu} F^{\mu\nu(b)} - V(\phi) \right\},$$

(1.2)

where $V$, $\mathcal{K}_{\alpha\beta}$ and $\tau_{ab}$ are arbitrary functions of scalar fields $\phi^\alpha$, and the index $(a)$ in Maxwell fields $F^{(a)}_{\mu\nu} = \partial_\mu A^{(a)}_\nu - \partial_\nu A^{(a)}_\mu$ runs over the set of nonzero chemical potentials. Corresponding to finite temperature $(D + 1)$-dimensional gauge theory with chemical potentials, the effective action (1.2) must admit a black $D$-dimensional brane solution, electrically charged under vector fields $A^{(a)}_\mu$. The $SO(D)$ invariance of the solution implies that it is of the form

$$ds^2_{D+2} = -c_1^2(r) (dt)^2 + c_2^2(r) \sum_{i=1}^D (dx^i)^2 + c_3^2(r) (dr)^2,$$

(1.3)

$$A^{(a)}_\mu = \delta^t_\mu \Phi^{(a)}(r), \quad \phi^a = \phi^a(r),$$

for some scalar potentials $\Phi^{(a)}$.  

Finally, we assume that the black brane horizon of (1.3) is nonsingular. This implies that there is a choice of the radial coordinate such that as $r \to r_{\text{horizon}}$

$$c_1 \to \alpha_1 (r - r_{\text{horizon}})^{1/2}, \quad c_2 \to \alpha_2, \quad c_3 \to \alpha_3 (r - r_{\text{horizon}})^{-1/2},$$

(1.4)

where constants $\alpha_i$ satisfy

$$\alpha_1 \alpha_2 \alpha_3 \neq 0.$$
Notice that given (1.4) the black brane temperature $T$ and the entropy density $s$ are correspondingly

$$T = \frac{\alpha_1}{4\pi\alpha_3}, \quad s = \frac{\alpha^D}{4G_{D+2}}. \quad (1.6)$$

2 The proof

In this section, we examine the hydrodynamics of the gauge theory plasma at finite chemical potential dual to the generic black hole solution (1.3). In particular, using prescription [31], we compute the retarded Green’s function of the boundary stress-energy tensor $T_{\mu\nu}(t, x^i)$ ($\mu = \{t, x^i\}$) at zero spatial momentum, and in the low-energy limit $\omega \to 0$:

$$G_{12,12}^R(\omega, 0) = -i \int dt d^Dx \ c^{\omega t} \theta(t) \langle [T_{12}(t, x^i), T_{12}(0, 0)] \rangle. \quad (2.1)$$

Computation of this Green’s function allows for a determination of the shear viscosity $\eta$ through the Kubo relation

$$\eta = \lim_{\omega \to 0} \frac{1}{2\omega} \left[ G_{12,12}^A(\omega, 0) - G_{12,12}^R(\omega, 0) \right], \quad (2.2)$$

where the advanced Green’s function is given by $G^A(\omega, 0) = (G^R(\omega, 0))^*$. We find

$$G_{12,12}^R(\omega, 0) = -\frac{i\omega s}{4\pi} \left( 1 + O\left(\frac{\omega}{T}\right) \right), \quad (2.3)$$

where $s$ is the entropy density. Inserting this expression into (2.2) yields the universal ratio (1.1).

We begin the computation of (2.1) by recalling that the coupling between the boundary value of the graviton and the stress-energy tensor of a gauge theory is given by $\delta g_{12}^{\mu \nu} T_{\mu \nu}^2/2$. According to the gauge/gravity prescription, in order to compute the retarded thermal two-point function (2.1), we should add a small bulk perturbation $\delta g_{12}(t, r)$ to the metric (1.3)

$$ds_{D+2}^2 \to ds_{D+2}^2 + \delta g_{12}(t, r) \ dx^1 dx^2, \quad (2.4)$$

and compute the on-shell action as a functional of its boundary value $\delta g_{12}^b(t)$. Symmetry arguments [34] guarantee that for a perturbation of this type in the background (1.3) all the other components of a generic perturbation $\delta g_{\mu \nu}$, along with the gauge potentials perturbations $\delta A_{\mu}^{(a)}$ and scalar perturbations $\delta \phi^a$, can be consistently set to zero.
Instead of working directly with $\delta g_{12}$, we find it convenient to introduce the field $\psi = \psi(t, r)$ according to
\[
\psi = \frac{1}{2} g^{11} \delta g_{12} = \frac{1}{2} c_2^{-2} \delta g_{12}.
\] (2.5)

The retarded correlation function $G_{12,12}^R(\omega, 0)$ can be extracted from the (quadratic) boundary effective action $S_{\text{boundary}}$ for the metric fluctuations $\psi^b$ given by
\[
S_{\text{boundary}}[\psi^b] = \int \frac{d^{D+1}k}{(2\pi)^{D+1}} \psi^b(-\omega) F(\omega, r) \psi^b(\omega) \bigg|_{\partial M_{D+2}^r},
\] where
\[
\psi^b(\omega) = \int \frac{d^{D+1}k}{(2\pi)^{D+1}} e^{-i\omega t} \psi(t, r) \bigg|_{\partial M_{D+2}^r}.
\] (2.6)

In particular, the Green’s function is given simply by
\[
G_{12,12}^R(\omega, 0) = \lim_{\partial M_{D+2}^r \to \partial M_{D+2}} 2 F(\omega, r),
\] (2.8)

where $\mathcal{F}$ is the kernel in (2.6). The boundary metric functional is defined as
\[
S_{\text{boundary}}[\psi^b] = \lim_{\partial M_{D+2}^r \to \partial M_{D+2}} \left( S_{\text{bulk}}^r[\psi] + S_{\text{GH}}[\psi] + S_{\text{counter}}[\psi] \right),
\] (2.9)

where $S_{\text{bulk}}^r$ is the bulk Minkowski-space effective supergravity action (1.2) on a cut-off space $M_{D+2}^r$ (where $M_{D+2}$ in (1.3) is regularized by the compact manifold $M_{D+2}^r$ with a boundary $\partial M_{D+2}^r$). Also, $S_{\text{GH}}$ is the standard Gibbons-Hawking term over the regularized boundary $\partial M_{D+2}^r$. The regularized bulk action $S_{\text{bulk}}^r$ is evaluated on-shell for the bulk metric fluctuations $\psi(t, r)$ subject to the following boundary conditions:

(a) : $\lim_{\partial M_{D+2}^r \to \partial M_{D+2}} \psi(t, r) = \psi^b(t),$

(b) : $\psi(t, r)$ is an incoming wave at the horizon.

The purpose of the boundary counterterm $S_{\text{counter}}$ is to remove divergent (as $\partial M_{D+2}^r \to \partial M_{D+2}$) and $\omega$-independent contributions from the kernel $\mathcal{F}$ of (2.6).

The effective bulk action for $\psi(t, r)$ is derived in Appendix A. It takes the form
\[
S_{\text{bulk}}[\psi] \equiv \frac{1}{16\pi G_{D+2}} \int_{M_{D+2}} d^{D+2} \mathcal{L}_{D+2} = \frac{1}{16\pi G_{D+2}} \int_{M_{D+2}} d^{D+2} \xi \left[ c_1 c_2 c_3 \left\{ \frac{1}{2c_1^2} (\partial_t \psi)^2 - \frac{1}{2c_2^2} (\partial_r \psi)^2 \right\} \right. \right.
\]
\[
\left. + \left\{ -\partial_t \left( \frac{2c_2 c_3}{c_1} \psi \partial_t \psi \right) + \partial_r \left( \frac{2c_2}{c_1} c_3 \psi \partial_r \psi + \frac{c_1 c_2}{c_3} \left( \partial_t \psi \right)^2 \right) \right\} \right].
\] (2.11)
The second line in (2.11) is the effective action for a minimally coupled scalar in the geometry (1.3), while the third line is a total derivative. Thus the bulk equation of motion for $\psi$ is that of a minimally coupled scalar in (1.3). Decomposing $\psi$ as

$$\psi(t, r) = e^{-i\omega t} \psi_\omega(r),$$

we find that the equation of motion reduces to

$$\frac{\omega^2 \psi_\omega}{c_1^2} + \frac{\psi_\omega''}{c_2^2} + \left[ \ln \frac{c_1 c_2^3}{c_3} \right] \frac{\psi_\omega'}{c_2^2} = 0,$$

where primes denote derivatives with respect to $r$.

Let's understand the characteristic indices of $\psi_\omega$ as $r \to r_{\text{horizon}}$. Consider the following ansatz

$$\psi_\omega(r) \sim (\Lambda(r - r_{\text{horizon}}))^{\lambda},$$

where $\Lambda$ is a characteristic, finite, $\omega$-independent energy scale associated with the background geometry (1.3). For background geometries dual to conformal gauge theory plasmas the role of $\Lambda$ is being played by the temperature. More generally, $\Lambda$ could be a scale at which conformal invariance is broken (either explicitly or spontaneously). It is crucial for $\Lambda$ to be finite — otherwise the hydrodynamic approximation to a strongly coupled gauge theory plasma would not applicable\(^5\). Substituting (2.14) into (2.13) and using the near horizon asymptotics (1.4), we find to leading order in $(\Lambda(r - r_{\text{horizon}})) \ll 1,$

$$(r - r_{\text{horizon}})^{\lambda-1} \left\{ \frac{\omega^2}{\alpha_1^2} + \frac{\lambda(\lambda - 1)}{\alpha_3^2} + \frac{\lambda}{\alpha_2^2} \right\} = 0,$$

or

$$\lambda = \pm i \frac{\alpha_3 \omega}{\alpha_1} = \pm i \frac{\omega}{4\pi T},$$

where we used (1.6). Given (2.12), for the incoming wave at the horizon we must have

$$\psi_\omega(r)_{\text{incoming}} \sim (\Lambda(r - r_{\text{horizon}}))^{\lambda\omega/(4\pi T)}, \quad (\Lambda(r - r_{\text{horizon}})) \ll 1,$$

which to leading order in $\frac{\omega}{T}$ takes the form

$$\psi_\omega(r)_{\text{incoming}} \sim 1 - i \frac{\omega}{4\pi T} \ln((\Lambda(r - r_{\text{horizon}})),$$
valid when both inequalities are satisfied

\[ \left| \ln((\Lambda(\tau - r_{\text{horizon}}))) \right| \gg 1 \quad \text{and} \quad \frac{\omega}{T} \left| \ln((\Lambda(\tau - r_{\text{horizon}}))) \right| \ll 1. \tag{2.19} \]

Clearly, for finite \( \Lambda \) and sufficiently small \( \omega \) there is an overlap region in (2.19).

In what follows we will need a solution to (2.13), subject to (2.10), to linear order in \( \omega T \), i.e., in the low-frequency approximation. In this approximation one can neglect the first term in (2.13) and write down the most general solution as

\[ \psi_\omega(r) = A_1(\omega) + A_2(\omega) \int_r^\infty \frac{c_3(\rho)d\rho}{c_1(\rho)c_2(\rho)B}, \tag{2.20} \]

where \( A_i(\omega) \) are constants, depending at most linearly on \( \frac{\omega}{T} \). The first boundary condition in (2.10) implies that

\[ A_1(\omega) = 1. \tag{2.21} \]

Substituting (1.4) into (2.20) we obtain the leading near horizon behavior of \( \psi_\omega \)

\[ \psi_\omega(r) \approx 1 - \frac{\alpha_3 A_2(\omega)}{\alpha_1 \alpha_2^D} \ln(\Lambda(\tau - r_{\text{horizon}})), \quad \left| \ln(\Lambda(\tau - r_{\text{horizon}})) \right| \gg 1. \tag{2.22} \]

Comparing (2.22) with (2.18) in the hydrodynamic approximation, i.e., for sufficiently small \( \frac{\omega}{T} \), we conclude that

\[ A_2(\omega) = i \alpha_2^D \omega. \tag{2.23} \]

To summarize, we have to leading order in \( \frac{\omega}{T} \)

\[ \psi_\omega(r) = 1 + i \alpha_2^D \omega \int_r^\infty \frac{c_3(\rho)d\rho}{c_1(\rho)c_2(\rho)B}. \tag{2.24} \]

Notice that as \( r \to \infty \)

\[ \psi_\omega(r) \to 1, \quad \partial_r \psi_\omega \to -i \alpha_2^D \omega \frac{c_3(r)}{c_1(r)c_2(r)B}. \tag{2.25} \]

Once the bulk fluctuations are on-shell (i.e., satisfy equations of motion) the bulk gravitational Lagrangian becomes a total derivative. From (2.11) we find (without dropping any terms)

\[ L_{D+2} = \partial_t J^t + \partial_r J^r, \tag{2.26} \]

where

\[ J^t = -\frac{3c_2^D}{2c_1} \psi \partial_t \psi, \]

\[ J^r = \frac{3c_2^D c_1}{2c_3} \psi \partial_r \psi + \frac{c_2^{D-1}c_1^2 c_2}{c_3} \psi^2. \tag{2.27} \]
Additionally, the Gibbons-Hawking term provides an extra contribution so that

\[ J' \to J' - \frac{2c_2^D c_1}{c_3} \psi \partial_r \psi - \left( \frac{c_1 c_2^D}{c_3} \right)' \psi^2. \]  

(2.28)

We are now ready to extract the kernel \( F \) of (2.6). The regularized boundary effective action for \( \psi \) is

\[ S_{\text{boundary}}[\psi]^r = \frac{1}{16\pi G_{D+2}} \int_{\partial M_{D+2}} dt \, d^D x \left( -\frac{c_2^D c_1}{2c_3} \psi \partial_r \psi \right) + \text{c.t.}, \]  

(2.29)

where, as prescribed in [31], we need only to keep the boundary contribution. In (2.29) \( \text{c.t.} \) stands for (finite) contact terms that will not be important for computations. Substituting (2.12), (2.24) into (2.29), and using (2.25), we obtain \( F(\omega, r) \) in the limit \( r \to \infty \)

\[ \lim_{r \to \infty} F^r(\omega, r) = -\frac{i\alpha D \omega}{32\pi G_{D+2}} \left( 1 + O \left( \frac{\omega}{T} \right) \right) \]  

(2.30)

\approx -\frac{i\omega s}{8\pi},

where we have recalled the definition of \( s \) in (1.6). Using (2.8) to extract the Green’s function from \( F^r \) we find

\[ G_R^{12}(\omega, 0) \approx -\frac{i\omega s}{4\pi}, \]  

(2.31)

at least in the low frequency limit \( \omega \to 0 \). This is the result claimed in (2.3), giving rise to the universal ratio of shear viscosity to entropy density (1.1).

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A Effective bulk action for $\psi$

We begin with collecting some useful expressions. The trace of the Einstein equations derived from (1.2) takes the form

$$\mathcal{R} = \mathcal{K}_{ab} \partial \phi^a \partial \phi^b + \frac{D - 2}{D} \tau_{ab} F_{\mu \nu}^{(a)} F^{\mu \nu(b)} + \frac{D + 2}{D} \mathcal{V}. \tag{A.1}$$

Additionally we have

$$\sqrt{-g(0)} R_{x_1(0)} = - \left[ \frac{c_1 c_2^{D-1} c_2'}{c_3} \right] = \frac{1}{D} \sqrt{-g(0)} \left( \mathcal{V} - \tau_{ab} F_{\mu \nu}^{(a)} F^{\mu \nu(b)} \right)^{(0)}, \tag{A.2}$$

where we used superscript (0) to emphasize that the Ricci tensor component and the metric determinant are evaluated on the background (1.3). The second equality in (A.2) makes use of the Einstein equation for $R_{x_1(0)}$.

Expanding the effective action (1.2) to quadratic order in perturbation (2.4) and using (A.1), we find

$$S_{\text{bulk}}[\psi] = \frac{1}{16 \pi G_{D+2}} \int_{M_{D+2}} d^{D+2} \xi \sqrt{-g(0)} \left\{ R^{(2)} - \frac{1}{2} \psi^2 \frac{2}{D} \left( \mathcal{V} - \tau_{ab} F_{\mu \nu}^{(a)} F^{\mu \nu(b)} \right)^{(0)} \right\}$$

$$= \frac{1}{16 \pi G_{D+2}} \int_{M_{D+2}} d^{D+2} \xi \sqrt{-g(0)} \left\{ R^{(2)} - \psi^2 R_{x_1(0)}^{(0)} \right\}, \tag{A.3}$$

where in the second line we used identity (A.2). In (A.3), $R^{(2)}$ is a perturbation of the Ricci scalar, quadratic in $\psi$ and its derivatives

$$\sqrt{-g(0)} R^{(2)} = \sqrt{-g(0)} \left\{ \frac{1}{2c_1^2} (\partial_t \psi)^2 - \frac{1}{2c_3^2} (\partial_r \psi)^2 \right\}$$

$$+ \left\{ -\partial_t \left( \frac{2c_2^D c_3}{c_1} \psi \partial_t \psi \right) + \partial_r \left( \frac{2c_2^D c_1}{c_3} \psi \partial_r \psi + \frac{c_1 c_2^{D-1} c_2'}{c_3} \psi^2 \right) \right\} \tag{A.4}$$

where we again used (A.2). From (A.3) and (A.4) we obtain (2.11).

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