

The charmonium dissociation in an “anomalous wind”

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ABSTRACT: We study the charmonium dissociation in a strongly coupled chiral plasma in the presence of magnetic field and axial charge imbalance. This type of plasma carries “anomalous flow” induced by the chiral anomaly and exhibits novel transport phenomena such as chiral magnetic effect. We found that the “anomalous flow” would modify the charmonium color screening length by using the gauge/gravity correspondence. We derive an analytical expression quantifying the “anomalous flow” experienced by a charmonium for a large class of chiral plasma with a gravity dual. We elaborate on the similarity and *qualitative* difference between anomalous effects on the charmonium color screening length which are *model-dependent* and those on the heavy quark drag force which are fixed by the second law of thermodynamics. We speculate on the possible charmonium dissociation induced by the chiral anomaly in heavy ion collisions.

KEYWORDS: Quark-Gluon Plasma, Holography and quark-gluon plasmas, Anomalies in Field and String Theories

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Contents

1	Introduction and summary	1
2	The holographic setup	3
3	The screening length l_s in the anomalous flow	5
3.1	String stretched by “anomalous flow”	8
4	Representative results for $\mathcal{N} = 4$ SYM chiral plasma	9
5	Discussion and phenomenological consequences	11
5.1	Comparison with anomalous contribution to heavy quark drag force	11
5.2	Phenomenological implications	12
A	Derivation of (3.10)	13

1 Introduction and summary

The work of Matsui and Satz [1] introduced the idea of using a quarkonium to probe quark gluon plasma (QGP). In a deconfined QGP, quarkonium bound states such as charmonium $c\bar{c}$ will dissociate because of color screening and thereby exhibit a suppression relative to the confined phase. Due to its importance, the problem of charmonium dissociation in QGP has been a focus of many recent studies (see refs. [2, 3] for reviews and references).

In this paper, we consider the problem of charmonium dissociation in a chiral (parity-violating) plasma with a finite chiral (axial) charge density¹ and in the presence of an external magnetic field. This environment is pertinent to the QGP, which is approximately chiral, created in heavy-ion collisions. First, a very strong magnetic field is generated from the incoming nuclei that are positively charged and move at nearly the speed of light. Such magnetic field has a magnitude of the order of $eB \sim m_\pi^2$ and its lifetime can be significant when medium’s effect is taken into consideration [5, 6]. Meanwhile, QCD as a non-Abelian gauge theory has topologically nontrivial gluonic configurations such as instantons and sphalerons. These configurations couple to quarks through the chiral anomaly and translate topological fluctuations into the chiral imbalance for quarks.

The focus of our study is on the effects of the chiral anomaly on the color screening length l_s , which is an important parameter quantifying charmonium dissociation. In heavy-ion collisions, the produced charmonium is moving relative to QGP and the relative velocity v (or rapidity $\eta = \tanh^{-1}(v)$) can be significant. We therefore also take dependence of l_s on rapidity η into consideration.

¹For a recent discussion of lattice QCD with chiral chemical potential see e.g. [4].

Anomaly-induced effects in a chiral medium has attracted much interests recently (see refs. [7–10] for reviews). One familiar example is the chiral magnetic effect (CME) [11–16], the generation of a vector current \mathbf{j}_V along an external magnetic field \mathbf{B} . In particular and closely related to the current work, those anomalous effects modify hydrodynamics of chiral fluids [17] (see also refs. [18–20]). For such fluid in the frame that energy density is at rest (i.e. in the Landau frame), the entropy density is not at rest and is moving opposite to the direction of chiral magnetic current $\mathbf{j}_{\text{CME}} = C_A \mu_A \mathbf{B}$:

$$\mathbf{s}_{\text{anom}} = s v_{\text{anom}} \hat{B}, \quad v_{\text{anom}} \equiv -\frac{C_A \mu_A \mu}{(\epsilon + p)} B, \quad C_A = \frac{N_c}{2\pi^2}, \quad (1.1)$$

where ϵ , s , p , μ (μ_A) denotes the energy density, entropy density, pressure and vector (axial) chemical potential respectively and the coefficient C_A is fixed by the chiral anomaly.² In this work, we will use charmonium (or in general quarkonium) to probe such an anomalous chiral fluid and ask how its screening length $l_s(\eta; v_{\text{anom}})$ would be influenced by the presence of the “anomalous flow” v_{anom} .

To compute the rapidity-dependent color screening length $l_s(\eta; v_{\text{anom}})$, we will use the gauge/gravity correspondence following the general formalism of ref. [21]. Previously, $l_s(\eta)$ and the dissociation of a moving charmonium has been studied in the framework of holographic correspondence from both top-down [21, 22] and bottom-up [23] approaches. Quarkonia dissociation in the presence of magnetic field has also been addressed previously (see for example refs. [24–28]). To the extent of our knowledge, the effects of the chiral anomaly on the charmonium dissociation have not been reported in literature before.

The main finding of this paper is that the charmonium color screening length $l_s(\eta; v_{\text{anom}})$ receives contributions from the chiral anomaly: a charmonium finds itself in a wind induced by anomalous flow (1.1). Let us quantify the “anomalous flow” felt by a charmonium introducing η_{anom} with following properties:

$$l_s(\eta; v_{\text{anom}}) = l_s(\eta + \eta_{\text{anom}}; v_{\text{anom}} = 0). \quad (1.2)$$

In other words, the color screening length of a charmonium moving at rapidity η in the *presence* of the anomalous flow v_{anom} equals to that of a charmonium moving at rapidity $\eta + \eta_{\text{anom}}$ in the *absence* of the anomalous flow. For small $v_{\text{anom}} \ll 1$, we obtain an analytical expression within the current holographic model at linear order in v_{anom} for η_{anom} . We observe that the magnitude of η_{anom} is proportional to v_{anom} . However, its value is *model-dependent*. Very recently, anomalous contributions to the heavy quark drag force were studied for a holographic chiral fluids [29]. Our study here provides further insights on the “anomalous flow” felt by heavy probes of the chiral plasma.

This paper is organized as follows. In section 2, we describe our holographic set-up. In section 3, we derive the analytic formula which determines the anomalous contribution to the color screening length $l_s(\eta; \eta_{\text{anom}})$. At this point, our results are valid for a large class of holographic chiral fluids. In section 4, we take $\mathcal{N} = 4$ SYM theory as an example and present η_{anom} as well as anomalous contributions to $l_s(\eta; v_{\text{anom}})$. We compare our results

²We set charge $e = 1$ here but will recover e -dependence in (5.2).

with anomalous contributions to the heavy quark drag force found in refs. [29, 30] and speculate on phenomenological implications for heavy-ion collisions in section 5.

2 The holographic setup

We start with the 4+1-dimensional asymptotic AdS bulk metric which is dual to the dense strongly coupled plasma of $\mathcal{N} = 4$ SYM theory with a nonzero axial chemical potential (as a manifestation of the chiral imbalance) and with a homogeneous external magnetic field \mathbf{B} . Such metric can be determined by solving the bulk Einstein-Maxwell-Chern-Simons equations. For analytical transparency, we will consider the situation that v_{anom} given by (1.1) is small. Consequently, the bulk metric, at leading order in v_{anom} , can be obtained by solving the linearized bulk equation of motion around AdS Reissner-Nordstrom (RN) black-brane solution, see refs. [17, 31, 32] for details. The resulting metric, in the Landau frame of the fluid, is of the form:

$$ds^2 = g_{xx}(r)\{-f(r)dt^2 + d\mathbf{x}^2 - 2v_{\text{anom}}Q(r)dzdt\} - 2\frac{v_{\text{anom}}Q(r)}{f(r)}dzdr + [g_{xx}(r)f(r)]^{-1}dr^2. \quad (2.1)$$

Here r denotes holographic coordinates with the boundary at $r \rightarrow \infty$ and the horizon at $r \rightarrow r_H$, where $f(r_H) = 0$. Without losing generality, we take the chiral magnetic current to be along the z direction (cf. figure 1).

To keep our discussion as general as possible, we will defer writing down explicitly r -dependence of the metric (2.1) to section 4. Our discussion presented in this section and section 3 can be applied to a large class of chiral fluids in the presence of the anomalous flow with holographic dual of the form (2.1). In (2.1), $g_{tz}(r)$ and $g_{zr}(r)$ which are proportional to $v_{\text{anom}}Q(r)$ are induced by the chiral anomaly.

We now consider a “dipole” moving through a thermal plasma at the velocity \mathbf{v} . To be specific, \mathbf{v} here denotes the velocity with respect to the Landau rest frame. To concentrate on the effects related to the chiral anomaly, we take the velocity \mathbf{v} to be along z direction (see figure 1 for a schematic view). To proceed, it is convenient to consider the gravity background (2.1) in the frame that the “dipole” is at rest while energy density of the plasma is moving. This metric can be determined by boosting the metric (2.1) using:

$$dt = \cosh \eta dt' - \sinh \eta dz', \quad dz = -\sinh \eta dt' + \cosh \eta dz'. \quad (2.2)$$

Here we have also introduced the rapidity:

$$v \equiv \tanh(\eta). \quad (2.3)$$

As a result, we have (dropping the prime to save notations):

$$ds^2 = G_{tt}(r)dt^2 + G_{zz}(r)dz^2 + g_{xx}(r)(dx^2 + dy^2) - \left(\frac{1}{g_{xx}(r)f}\right)dr^2 + 2[G_{tz}(r)dt dz - v_{\text{anom}} \cosh \eta Q(r)f^{-1}dzdr + v_{\text{anom}} \sinh \eta f^{-1}(r)Q(r)dt dr], \quad (2.4)$$

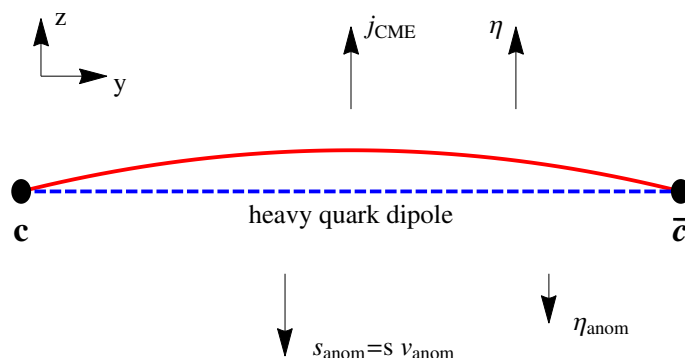


Figure 1. A schematically view of the heavy-quark “dipole” configuration and the corresponding string world-sheet (projected into $y - z$ plane) in the holographic set-up considered in this paper. The chiral magnetic current is along z -direction and the “dipole” is placed along y -direction. The entropy flow v_{anom} in (1.1) is opposite to the chiral magnetic current. We will show η_{anom} (1.2) is also along v_{anom} direction. Here the “dipole” is moving at rapidity η along v_{anom} direction. Nevertheless, η_{anom} is invariant under $\eta \rightarrow -\eta$. Without the anomalous flow v_{anom} , the string profile is constant along z -direction (dashed blue curve). However, the anomalous flow will stretch the string profile along z -direction (red curve, see (3.1) for details).

where:

$$G_{tt}(r) = (\cosh \eta)^2 g_{xx}(r) \left[(-f(r) + \tanh^2 \eta) + 2v_{\text{anom}} \tanh \eta Q(r) \right], \quad (2.5)$$

$$G_{zz}(r) = (\cosh \eta)^2 g_{xx}(r) \left[(1 - \tanh^2 \eta f(r)) + 2v_{\text{anom}} \tanh \eta Q(r) \right],$$

$$G_{tz}(r) = -(\cosh \eta)^2 g_{xx}(r) \left[\tanh \eta (1 - f(r)) + v_{\text{anom}} Q(r) (1 + \tanh^2 \eta) \right]. \quad (2.6)$$

To describe the interaction potential energy E of a heavy quark “dipole” with quarks separated by a distance L and moving at rapidity η , we consider the Wilson loop in the frame (2.4) whose contour \mathcal{C} is given by a rectangle with large extension \mathcal{T} in the t -direction and short sides of length L along some spatial direction. We take short sides of \mathcal{C} to lie in the y direction that is transverse to both \mathbf{v} and “anomalous flow” \mathbf{v}_{anom} . (cf. figure 1). The generalization to the arbitrary angle between L and \mathbf{v} is straightforward.

Following holographic dictionary, the interaction potential energy E of the “dipole”, measured by the thermal expectation value $\langle W(\mathcal{C}) \rangle$, is given by [33]:

$$e^{-iE\mathcal{T}} = \langle W(\mathcal{C}) \rangle = e^{iS[\mathcal{C}]}, \quad (2.7)$$

where $S[\mathcal{C}]$ is related to the corresponding Nambu-Goto action:

$$S[\mathcal{C}] = \frac{1}{2\pi\alpha} \int d\sigma d\tau \mathcal{L} = \frac{1}{2\pi\alpha} \int d\sigma d\tau \sqrt{-\det h_{\alpha\beta}}, \quad (2.8)$$

with the induced metric given by:

$$h_{\alpha\beta} = G_{MN} \partial_\alpha X^M \partial_\beta X^N. \quad (2.9)$$

Here, M, N run over t, \mathbf{x}, r and α, β run over τ, σ and G_{MN} can be read from (2.4).

The action $S(\mathcal{C})$ is invariant under the choice of τ, σ and we choose $\tau = t$ and $\sigma = y$ for convenience. Since $\mathcal{T} \gg L$, we can assume that the surface is translationally invariant along τ direction and therefore $X^M(\sigma)$ depends on σ only. The ‘‘Lagrangian’’ in (2.8) reads:

$$\mathcal{L}(r', z'; r) = \sqrt{A(r) \left[1 + (B(r))^{-1} (r')^2 \right] + C(r)(z')^2 + 2v_{\text{anom}}D(r)z'r'}, \quad (2.10)$$

where $A(r), B(r), C(r)$, to linear order in v_{anom} , are given by:

$$\begin{aligned} A(r) &= \cosh^2 \eta g_{xx}^2(r) [f(r) - \tanh^2 \eta - 2v_{\text{anom}} \tanh \eta Q(r)] + \mathcal{O}(v_{\text{anom}}^2), \\ B(r) &= f(r)g_{xx}^2(r) + \mathcal{O}(v_{\text{anom}}^2), \quad C(r) = \cosh^2 \eta g_{xx}^2(r)f(r) + \mathcal{O}(v_{\text{anom}}^2), \\ D(r) &= -\cosh \eta g_{xx}(r)Q(r) + \mathcal{O}(v_{\text{anom}}). \end{aligned} \quad (2.11)$$

Here and hereafter, we use prime to denote the derivative with respect to σ (or equivalently to y). As $\mathcal{L}(r', z'; r)$ is independent of z , the momentum π_z conjugated to z' is a constant of the motion:

$$\pi_z \equiv \frac{\partial \mathcal{L}}{\partial z'} = \frac{C(r)z' + D(r)r'}{\mathcal{L}} = \text{const.} \quad (2.12a)$$

Moreover, as $\mathcal{L}(r', z'; r)$ has no explicit dependence on σ , the corresponding Hamiltonian H :

$$H \equiv \mathcal{L} - r' \frac{\partial \mathcal{L}}{\partial r'} - z' \frac{\partial \mathcal{L}}{\partial z'} = \frac{A(r)}{\mathcal{L}} = \text{const.}, \quad (2.12b)$$

is also a constant of motion. From (2.12) and boundary conditions,

$$z \left(\frac{L}{2} \right) = z \left(-\frac{L}{2} \right) = 0, \quad r \left(\frac{L}{2} \right) = r \left(-\frac{L}{2} \right) = \Lambda \rightarrow \infty, \quad (2.13)$$

$r(\sigma)$ and $z(\sigma)$ can be determined for each given set of integration constants. Here Λ is a cut-off along holographic direction. We will show in section 3.1 that $\pi_z = 0$ for the given setup. Therefore the solution only depends on the integration constant H . In fact, for a given H , one could determine L by solving (2.12) and (2.13). Consequently, one might consider $L(H)$ as a function of H . Generically, $L(H)$ will reach a maximum at some H , say H_m (see for example ref. [21]). Following ref. [21], we will interpret:

$$l_s \equiv L(H_m), \quad (2.14)$$

as the screening length of the quarkonium potential. The physical picture of such definition of the screening length is clear: for $L > l_s$, (2.12) has no solutions. Therefore there is no L -dependent potential between the quark and antiquark and attractive force between a quark and anti-quark pair separated by $L > l_s$ will be screened.

3 The screening length l_s in the anomalous flow

In this section, we will determine the anomalous contribution at the linear order in v_{anom} . In particular, we will expand $L(H; \eta)$ and $l_s(\eta)$ as:

$$\begin{aligned} L(H; \eta) &= L_0(H; \eta) + v_{\text{anom}}L_1(H; \eta) + \mathcal{O}(v_{\text{anom}}^2), \\ l_s(\eta; v_{\text{anom}}) &= l_0(\eta) + v_{\text{anom}}l_1(\eta) + \mathcal{O}(v_{\text{anom}}^2), \end{aligned} \quad (3.1)$$

and obtain expressions for $L_1(H; \eta)$ and $l_1(\eta)$. The zeroth order results (i.e. results in the absence of anomalous effects) such as string profile $r_0(\sigma), z_0(\sigma)$ and $l_0(\eta)$ are extensively discussed in literature [21, 22]. Using the results of ref. [21], we have (see also section 3.1):

$$z_0(\sigma) = \text{const.} \quad (3.2)$$

Eq. (3.2) implies that $z' \sim v_{\text{anom}}$ and hence (2.12b) becomes

$$H = \frac{A(r)}{\sqrt{A(r)[1 + (r')^2/B(r)]}} + \mathcal{O}(v_{\text{anom}}^2). \quad (3.3)$$

This leads to the expression:

$$\frac{dr}{d\sigma} = -\frac{1}{H} \sqrt{B(r)[A(r) - H^2]}. \quad (3.4)$$

We note that with given H , r.h.s. of (3.4) will vanish at $r = r_c$ with r_c satisfying:

$$A(r_c) = H^2. \quad (3.5)$$

Due to $r'(\sigma) = 0$ at this point, r_c would be the minimum r the world sheet would reach. In other words, the world sheet stretches from $r = \Lambda \rightarrow \infty$ down to $r = r_c$. Since $r(\sigma)$ is an even function of σ , using boundary condition (2.13), we have:

$$L(H; \eta) = 2H \int_{r_c}^{\Lambda} dr \frac{1}{\sqrt{B(r)[A(r) - H^2]}}. \quad (3.6)$$

It is clear from the above derivation that zeroth order results $L_0(H)$ can be recovered by the replacement $A(r) \rightarrow A_0(r)$, $r_c \rightarrow r_c^0$ in (3.6):

$$L_0(H; \eta) = 2H \int_{r_c^0}^{\Lambda} dr \frac{1}{\sqrt{B(r)[A_0(r) - H^2]}}, \quad (3.7)$$

where r_c^0 is fixed by:

$$A_0(r_c^0) = H^2. \quad (3.8)$$

Here we define A_0 and A_1 by expanding A in powers of v_{anom} : $A(r) = A_0(r) + v_{\text{anom}} A_1(r)$ and one finds using (2.11) that

$$A_0(r) = \cosh^2 \eta g_{xx}^2(r) [f(r) - \tanh^2 \eta], \quad A_1(r) = -\sinh(2\eta) g_{xx}^2(r) Q(r). \quad (3.9)$$

As the anomalous contribution to B is already at $\mathcal{O}(v_{\text{anom}}^2)$, we will not distinguish B from its zeroth order part B_0 (similar for C and D used in section 3.1).

At the linear order in v_{anom} , one can determine $L_1(H; \eta)$ defined in (3.1) by expanding (3.6) in powers of v_{anom} . As a result (see appendix A for details), we have:

$$L_1(H; \eta) = -\frac{2H A_1(r_c^0)}{V_A(r_c^0)} \langle V_A(r); H \rangle, \quad (3.10)$$

where we introduce

$$V_A(r) \equiv \frac{v_{\text{anom}} Q(r)}{1 - f(r)}. \quad (3.11)$$

We will discuss the physical interpretation of $V_A(r)$ shortly. In (3.10), we have defined the “average” over holographic coordinate for any function of $F(r)$:

$$\begin{aligned} \langle F(r); H \rangle \equiv & \int_{r_c^0}^{\infty} dr \left\{ v \left[\frac{(1-f(r))g_{xx}^2(r)}{(1-f(r_c^0))g_{xx}^2(r_c^0)} F(r) - \frac{A_0'(r)r}{A_0'(r_c^0)r_c^0} F(r_c^0) \right] \times \frac{1}{2[B(r)]^{1/2}[A_0(r)-A_0(r_c^0)]^{3/2}} \right. \\ & \left. + \frac{F(r_c^0)}{A_0'(r_c^0)r_c^0} \left[\frac{2B(r)-B'(r)r}{2[B(r)]^{3/2}[A_0(r)-A_0(r_c^0)]^{1/2}} \right] \right\}, \end{aligned} \quad (3.12)$$

to make our notations compact. Obviously, due to (3.11), $L_1(H; \eta)$ determined from (3.10) is independent of v_{anom} .

With (3.10), we now ready to determine $l_s(\eta)$. Let H_m^0 and H_m be the H corresponding to the maximum of $L_0(H; \eta)$ and $L(H; \eta)$ respectively, i.e. $l_0(\eta) = L_0(H_m^0; \eta)$, $l_s(\eta) = L(H_m; \eta)$. Note that $L_0(H_m; \eta) = L_0(H_m^0; \eta) + \mathcal{O}(v_{\text{anom}}^2)$ due to $\partial_H L_0(H_m^0; \eta) = 0$. Hence we have $L(H_m; \eta) = L_0(H_m^0; \eta) + v_{\text{anom}} L_1(H_m; \eta) = L_0(H_m^0; \eta) + v_{\text{anom}} L_1(H_m^0; \eta) + O(v_A^2)$ where we noticed that the difference between H_m^0 and H_m is of the order v_{anom} . We therefore have

$$l_s(\eta) = l_0(\eta) + v_{\text{anom}} L_1(H_m^0(\eta); \eta), \quad (3.13)$$

Now one can compute η_{anom} , the “anomalous flow” felt by the charmonium moving in a chiral plasma using (1.2). From (3.13), we easily find:

$$\eta_{\text{anom}} = \frac{\langle V_A(r); H_0^m \rangle}{\langle 1; H_0^m \rangle}. \quad (3.14)$$

Eq. (3.14) is one of the main results of this paper. We note from (3.12) and (2.11) that A_0, B is invariant under $\eta \rightarrow -\eta$, therefore η_{anom} is an even function of η .

To analyze (3.14), we first discuss the physical meaning of $V_A(r)$ defined in (3.11). We note that the chiral anomaly will induce a non-zero $g_{tz}(r) = g_{xx}(r)Q(r)$ in the metric (2.1). The presence of this $g_{tz}(r)$ in the metric is equivalent to a boost of the non-anomalous metric with r -dependent velocity $V_A(r)$ (keeping terms only up to the linear order in v_{anom}):

$$dt = dt' - V_A(r)dz', \quad dz = -V_A(r)dt' + dz'. \quad (3.15)$$

Therefore the bulk velocity $V_A(r)$ describes an “anomalous wind” flowing in the bulk. The “anomalous flow” η_{anom} probed by a charmonium is given by an average of $V_A(r)$ over the holographic coordinate r with an appropriate weight as given by (3.12). It is clear that the resulting η_{anom} will depend on the details of the bulk profiles thus will be *non-universal*.

In figure 2, we plot a representative $V_A(r)/v_{\text{anom}}$ vs r for $\mathcal{N} = 4$ SYM chiral plasma at finite chemical potential. Note that the asymptotics of $V_A(r)$ is model-independent. As one can check, the energy flow of the chiral fluid dual to (2.1) is proportional to $V_A(r \rightarrow \infty)$. As we are working in the Landau frame, the energy flow $V_A(r \rightarrow \infty)$ must vanish. On the other hand, at the horizon $r \rightarrow r_H$ one finds $V_A(r \rightarrow r_H) \rightarrow v_{\text{anom}}$. This is because $V_A(r \rightarrow r_H)$ is related to the entropy flow (1.1) induced by the anomaly which is model-independent. Indeed, As figure 2 indicates, $V_A(r)$ decreases from v_{anom} to 0 from boundary towards the horizon.

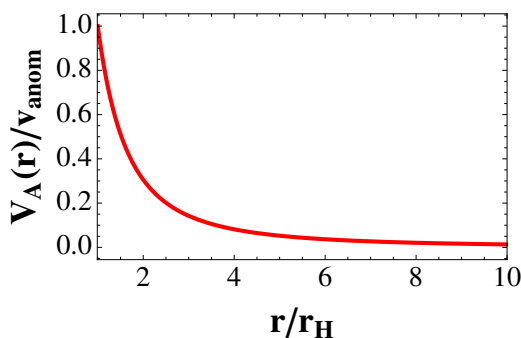


Figure 2. $V_A(r)/v_{\text{anom}}$, the bulk “anomalous velocity”, introduced in (3.11) as a function of r/r_H for $\mathcal{N} = 4$ SYM chiral plasma at finite chemical potential (cf. section 4). In numerical computations, we consider right-handed fermions with $\mu_R/T = 0.1$ (to the leading order in μ_R/T).

3.1 String stretched by “anomalous flow”

As another manifestation of the “anomalous flow” experienced by a moving charmonium, we now consider string profile along z -direction $z(\sigma)$. As mentioned before, we expand string profile $z(\sigma)$ as $z(\sigma) = z_0(\sigma) + v_{\text{anom}}z_1(\sigma)$.

First, let us briefly recall the derivation of (3.2) that $z_0(\sigma) = 0$ following the argument presented in ref. [21]. In the absence of the anomaly, (2.12a) becomes

$$\pi_z^0 = \frac{C(r)z'}{\mathcal{L}}. \tag{3.16}$$

If π_z^0 is non-zero then $z'_0 \neq 0$. The boundary condition (2.13) implies that there is at least a point at which $z'_0 = 0$. On the other hand, the Lagrangian \mathcal{L} has to be positive definite. Therefore r.h.s. of (3.2) vanishes at this point where $z' = 0$. That is in contradiction with the assumption that $\pi_z^0 \neq 0$ and thus we conclude that $\pi_z^0 = 0$ and $z_0(\sigma) = \text{const}$.

The situation is different due to anomaly. The condition (2.12a) at the leading order in v_{anom} can be expressed in terms of the zeroth-order profile:

$$\pi_z^1 = \frac{C[r_0(\sigma)]z'_1(\sigma) + D[r_0(\sigma)]r'_0(\sigma)}{\sqrt{A_0[r_0(\sigma)]\{1 + B_0[r_0(\sigma)]\}[r'_0(\sigma)]^2}} = \frac{H}{A_0[r_0(\sigma)]}\{C[r_0(\sigma)]z'_1(\sigma) + D[r_0(\sigma)]r'_0(\sigma)\}, \tag{3.17a}$$

where we have used the zeroth order relation for (2.12b):

$$H = \frac{A_0(r)}{\sqrt{A_0(r)[1 + B(r)(r')^2]}}. \tag{3.17b}$$

Solving (3.17a) and using boundary condition $z_1(-L/2) = 0$, we then have:

$$z_1(\sigma) = \left(\frac{\pi_z^1}{H}\right) \int_{-L/2}^{\sigma} d\tilde{\sigma} \frac{A_0[r_0(\tilde{\sigma})]}{C[r_0(\tilde{\sigma})]} - \int_{-L/2}^{\sigma} d\tilde{\sigma} \frac{D[r_0(\tilde{\sigma})]}{C[r_0(\tilde{\sigma})]} r'_0(\tilde{\sigma}). \tag{3.18}$$

The integration constant π_z^1 can be fixed by taking $\sigma \rightarrow L/2$ in (3.18) and imposing the boundary condition $z_1(L/2) = 0$. Noting $r_0(\sigma)$ is an even function of σ and $r'_0(\sigma)$ is an

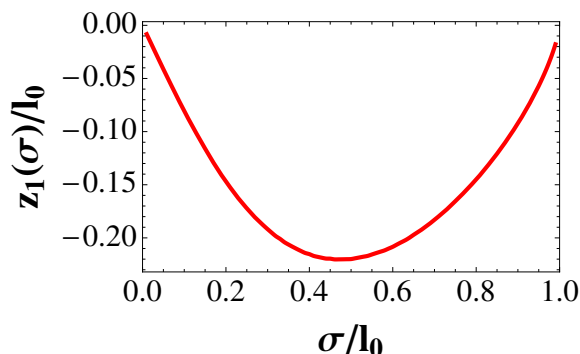


Figure 3. The anomalous contribution to the string profile $z_1(\sigma)$ as given by (3.20) for $\mathcal{N} = 4$ SYM chiral plasma at finite chemical potential (cf. section 4). The string is “dragged” by the anomalous flow. Both $z_1(\sigma)$ and σ are rescaled by the corresponding zeroth order screening length l_0 . In numerical computations, we consider right-handed fermions only with $\mu_R/T = 0.1$ and $v = 0.5$ (to the leading order in μ_R/T).

odd function of σ , we find:

$$z_1(L/2) = \left(\frac{\pi_z^1}{H}\right) \int_{-L/2}^{L/2} d\tilde{\sigma} \frac{A_0[r_0(\tilde{\sigma})]}{C[r_0(\tilde{\sigma})]}. \quad (3.19)$$

The integration in r.h.s. of (3.19) is non-zero since $A_0(r_0(\sigma))/C(r_0(\sigma))$ is even function of σ . Thus to satisfy boundary conditions, we must have also $\pi_z^1 = 0$ and consequently (3.18) becomes:

$$z_1(\sigma) = - \int_{-L/2}^{\sigma} d\tilde{\sigma} \frac{D(r_0(\tilde{\sigma}))}{C(r_0(\tilde{\sigma}))} r_0'(\tilde{\sigma}). \quad (3.20)$$

A representative profile of $z_1(\sigma)$ is plotted in figure 3. As one can check from (3.20), $z_1(\sigma)$ is negative. This implies that the string profile along z -direction is dragged behind the “anomalous wind” v_{anom} . From (2.11), we note C and D are invariant under $\eta \rightarrow -\eta$, therefore z_1 is an even function of η .

Let us contrast the anomalous contribution (3.20) with zero order results (3.2). $z_0(\sigma) = 0$ suggests that even though there is a wind blowing in the z -direction, the string world sheet is not dragged at all by this wind. However, non-trivial $z_1(\sigma)$ given by (3.20) indicates that string will be stretched by the anomalous flow. It would be interesting to gain further physical insights in this aspect.

4 Representative results for $\mathcal{N} = 4$ SYM chiral plasma

Up to this point, our discussion is general and applicable to any metric of the form (2.1). We now take $\mathcal{N} = 4$ SYM theory as an example and report the anomalous contribution to color screening length as defined in (3.1).

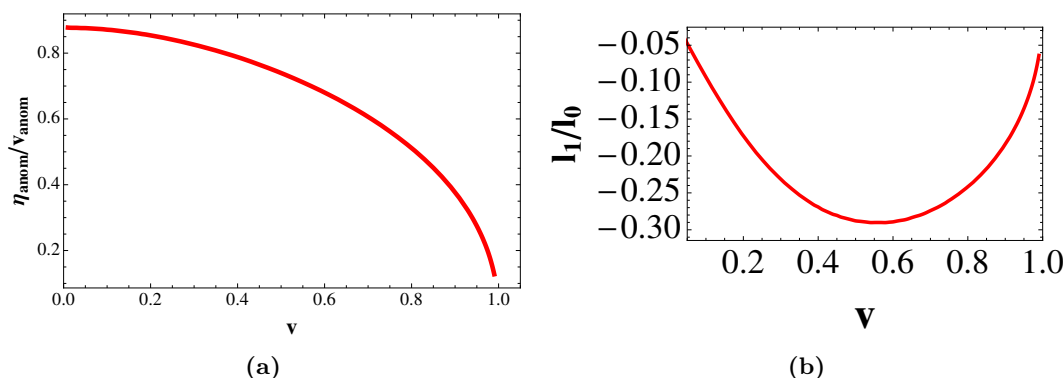


Figure 4. Left: anomalous velocity η_{anom} “felt” by a charmonium moving at velocity v in a chiral plasma with a anomalous entropy current $s v_{\text{anom}}$. Here, η_{anom} is determined from (3.14) and is along the direction of anomalous entropy current in Landau frame. Right: anomalous contribution to the color screening length l_s . As in (3.1), we expand the color screening length as $l_s = l_0 + v_{\text{anom}} l_1 + \mathcal{O}(v_{\text{anom}}^2)$. In numerical computations, we consider $\mathcal{N} = 4$ SYM chiral plasma at finite chemical potential with right-handed fermions only with $\mu_{\text{R}}/T = 0.1$ (to the leading order in μ_{R}/T). We have also translated the dependence on rapidity η into that of velocity v .

For $\mathcal{N} = 4$ SYM theory with a nonzero chemical potential and with a weak homogeneous external magnetic field \mathbf{B} , we have (see e.g. [34]):

$$g_{xx}(r) = r^2, \quad f(r) = 1 - \frac{M}{r^4} + \frac{q^2}{r^6},$$

$$Q(r) = \frac{2\pi^2 T^2}{r^2} - \frac{\pi^4 T^4}{r^4} + 2 \left(1 - \frac{\pi^4 T^4}{r^4} \right) \log \left(\frac{r^2}{r^2 + \pi^2 T^2} \right) + \mathcal{O}(\mu^2/T^2), \quad (4.1)$$

where for brevity we kept only the leading term in $Q(r)$ in powers of μ/T . In this section, we will consider one chirality, say right handed fermions, only and therefore μ should be interpreted as a chiral chemical potential μ_{R} . This is sufficient for our illustrative purpose.

Black hole parameters M , q could be easily related to physical quantities μ and T (see e.g. [29, 31]):

$$M = \frac{\pi^4 T^4}{16} \left[\sqrt{1 + \frac{8\mu^2}{3\pi^2 T^2}} + 1 \right]^3 \left[3\sqrt{\frac{8\mu^2}{3\pi^2 T^2}} + 1 - 1 \right], \quad q = \frac{\mu}{\sqrt{3}} \frac{\pi^2 T^2}{2} \left[\sqrt{1 + \frac{8\mu^2}{3\pi^2 T^2}} + 1 \right]^2. \quad (4.2)$$

As we will measure any dimensionful quantities in units of πT , the results then will depend on the ratio μ/T . In numerics, we choose a specific value $\mu/T = 0.1$ and use $Q(r)$ and $f(r)$ given by (4.1).

In figure 4a, we plot η_{anom} vs v as determined from (3.14). Observe as expected that η_{anom} is along the direction of the entropy flow v_{anom} . With increasing v , η_{anom} will decrease towards zero. This is of course not surprising as a ultra-relativistically moving charmonium should be insensitive to a small flow v_{anom} . Furthermore, In figure 4b, we plot

the (relative) anomalous contribution to the screening length, $l_1(v)/l_0(v)$ as a function of velocity v . Here, l_0 and l_1 are defined in (3.1) by expanding the color screening length in the presence of the anomaly to the leading order in v_{anom} . We observe that $l_1(v)/l_0(v)$ reaches its maximum at some intermediate velocity v . For small v , $l_s(v)/l_0(v)$ grows linearly in v . At high velocity v , $l_1(v)/l_0(v)$ is suppressed due to the vanishing of η_{anom} as shown in figure 4a. It is worthy noting that strictly speaking, in our small v_{anom} expansion, we have in-explicitly assumed that $v_{\text{anom}} \ll v$. Therefore small v asymptotic should be understood as taking $v \rightarrow 0$ limit while keeping $v_{\text{anom}} \ll v$.

5 Discussion and phenomenological consequences

In this work, we consider the charmonium color screening length l_s in a strongly coupled chiral plasma in the presence of flow induced by chiral anomaly. We use holographic correspondence to show l_s is influenced by the anomalous flow and establish an analytical formula (3.14) to quantify such influence. From the gravity side of the correspondence, such contributions stems from the modification of the bulk metric due to the chiral anomaly (cf. (2.1)). Consequently, it contributes to Nambu-Goto action (2.8) which determines the color screening length. While the anomalous effects on light quarks which are nearly chiral have been studied extensively, exploring the anomalous plasma using heavy probes is initiated very recently. Our results as a proof of principle, imply that those anomalous effects can be probed by charmonium. We will further discuss the physical interpretation and possible phenomenological implication below.

5.1 Comparison with anomalous contribution to heavy quark drag force

It is instructively to compare our results with the recent study of ref. [29] in which drag force $F(\eta; v_{\text{anom}})$ of a heavy quark moving at rapidity η in the presence of anomalous velocity v_{anom} (cf. (1.1)) is studied. Let us rephrase the results in ref. [29] by introducing $\eta_{\text{anom}}^{\text{drag}}$ which describes the ‘‘anomalous flow’’ experienced by a heavy quark. Similar to (1.2) where the ‘‘anomalous flow’’ experienced by a charmonium η_{anom} is introduced, we define $\eta_{\text{anom}}^{\text{drag}}$ by the condition:

$$F(\eta; v_{\text{anom}}) = F(\eta + \eta_{\text{anom}}^{\text{drag}}; v_{\text{anom}} = 0). \tag{5.1}$$

Ref. [29] indicates that η_{anom} is non-zero and is proportional to v_{anom} as well.

However, one should not overlook the *qualitative* difference between the case of the color screening length presented in here and that of the heavy quark drag coefficient. In fact, $\eta_{\text{anom}}^{\text{drag}}$ is independent of the microscopic details of the system while η_{anom} , as shown in this paper, is model-dependent. For example, in Landau frame and in $\eta \rightarrow 0$ limit (i.e. the heavy probe is at rest), one can see from ref. [29] that $\eta_{\text{anom}}^{\text{drag}} = V_A(r_H)$. Therefore $\eta_{\text{anom}}^{\text{drag}}$ only depends on the value of $V_A(r)$ at horizon and is uniquely fixed by v_{anom} . In contrast, η_{anom} in (3.14) is given by the bulk average over $V_A(r)$ and therefore depends on the bulk details of the holographic model. This remarkable difference is directly connected to the dissipation-less nature of the anomalous transport. As the presence of the drag force introduces dissipation, there should be a frame for a chiral plasma with anomalous flow

that drag force vanishes. In fact $-(\eta + \eta_{\text{anom}}^{\text{drag}})$ defines that frame as $F(0; v_{\text{anom}} = 0) = 0$. Therefore η_{anom} is constrained by the dissipation-less nature of the anomalous transport. Indeed, η_{anom} can be universally fixed by using anomalous hydrodynamics in ref. [30] and by the generalization of Landau’s criterion for superfluidity in ref. [35]. On the other hand, the color screening length is a static quantity and is not related to dissipative processes. This in turn implies that one can not constrain anomalous contributions to l_s directly from argument based on non-dissipationless nature of chiral effects. Our holographic study confirmed this difference.

5.2 Phenomenological implications

We now turn to phenomenological implications. As the anomalous contributions is controlled by the magnitude of the “anomalous flow” v_{anom} , let us begin our discussion by estimating the magnitude of v_{anom} in heavy-ion collisions. Recovering dependence on charge e , C_A in (1.1) we get:

$$C_A = \frac{C_{\text{EM}} e^2}{2\pi^2}. \tag{5.2}$$

We consider the case that only u, d will contribute to the CME current therefore $N_f = 2, C_{\text{EM}} = 5/9$. To estimate ϵ and p , we use the equation of state of a massless ideal quark-gluon gas $\epsilon = 3p$, with

$$p(T) = \frac{g_{\text{QGP}} \pi^2}{90} T^4, \tag{5.3}$$

where $g_{\text{QGP}} = g_g + 7g_q/8$ is the number of degrees of freedom with $g_g = (N_c^2 - 1)N_s$ and $g_q = 2N_c N_s N_f$; $N_c = 3$ and $N_f = 2$ are the numbers of colors and flavors and $N_s = 2$ is the number of spin states for quarks and (transverse) gluons.

As a results, we have:

$$v_{\text{anom}} \approx 0.003 \left(\frac{eB}{T^2} \right) \left(\frac{\mu_V \mu_A}{T^2} \right). \tag{5.4}$$

Keeping in mind the characteristic \sqrt{eB} and μ (is of order T) one can see that v_{anom} is numerically small. This implies that our results based on small v_{anom} expansion is applicable to many situations in heavy-ion collisions. However anomalous effects on charmonium dissociation are numerically small.

Moreover, we note that at early time of heavy-ion collisions, eB can be of the order $10 m_\pi^2$ at RHIC and of the order $100 m_\pi^2$ at LHC. Those numbers can be even larger from event by event fluctuations [36, 37]. This would lead to a significant v_{anom} . Therefore it would be interesting to consider the situation that B is large and explore if anomaly-induced η_{anom} would lead to a new mechanism for charmonium suppression at early time of heavy-ion collisions. We hope the results presented in this paper would encourage further study along this direction.

The current work can be extended in number of ways. In this exploratory study, we consider charmonium moving along the direction of anomalous flow and take the configuration of the heavy quark “dipole” to be perpendicular to the anomalous flow. The study for general angle between the anomalous flow and the “dipole” velocity would bring further

details on anomalous contributions to the screening length. We have restricted ourselves to the anomalous flow related to the chiral magnetic effects. One could extend it by incorporating chiral vortical effects [31, 32, 38] and possible contribution from gravitational anomaly [39]. We also assume both the magnetic field and axial charge to be static and homogeneous. In realistic situations such as the early times of heavy-ion collisions, both magnetic field and axial charge imbalance are dynamical. Those dynamical magnetic field and axial charge imbalance are shown to induce novel phenomena related to chiral anomaly (see for example refs. [40–44]). It would be interesting to understand their influences on the color screening. In the realistic situation the plasma is also anisotropic. Both the color screening length and the anomalous transport (see e.g. [45]) depend on the anisotropy of the medium (see e.g. [46]). Thus one would expect some interplay between them. However these directions are beyond the scope of this work and we leave them for future study.

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A Derivation of (3.10)

We now present the derivation of (3.10). From (3.6) and (3.7), we have:

$$\begin{aligned}
 v_{\text{anom}} L_1(\eta) &= L(H; \eta) - L_0(H; \eta) \\
 &= 2H \left\{ \int_{r_c}^r dr \frac{1}{\sqrt{B(r)[A(r) - H^2]}} - \int_{r_c^0}^r dr \frac{1}{\sqrt{B(r)[A_0(r) - H^2]}} \right\} \\
 &= 2H \left\{ \int_1^\infty d\tilde{r} \frac{r_c}{\sqrt{B(r_c \tilde{r})[A(r_c \tilde{r}) - H^2]}} - \int_1^\infty d\tilde{r} \frac{r_c^0}{\sqrt{B(r_c^0 \tilde{r})[A_0(r_c^0 \tilde{r}) - H^2]}} \right\} \\
 &= 2H \int_1^\infty d\tilde{r} \left\{ \frac{r_c}{\sqrt{B(r_c \tilde{r})[A(r_c \tilde{r}) - H^2]}} - \frac{r_c^0}{\sqrt{B(r_c^0 \tilde{r})[A_0(r_c^0 \tilde{r}) - H^2]}} \right\}. \quad (\text{A.1})
 \end{aligned}$$

In (A.1), we have introduced a rescaled variable \tilde{r} such that the integration is from 1 to ∞ . This would bring the convenience to expand (A.1) in power of v_{anom} ,

Let us now expand:

$$r_c = r_c^0 + v_{\text{anom}} r_c^1 + \mathcal{O}(v_{\text{anom}}^2), \quad (\text{A.2})$$

to obtain:

$$\begin{aligned}
 A(r_c \tilde{r}) &= A_0(r_c^0 \tilde{r}) + v_{\text{anom}} \left[A_1(r_c^0 \tilde{r}) + A_0'(r_c^0 \tilde{r}) (r_c^1 \tilde{r}) \right] + \mathcal{O}(v_{\text{anom}}^2), \\
 B(r_c \tilde{r}) &= B(r_c^0 \tilde{r}) + v_{\text{anom}} B'(r_c^0 \tilde{r}) (r_c^1 \tilde{r}) + \mathcal{O}(v_{\text{anom}}^2). \quad (\text{A.3})
 \end{aligned}$$

In the above expansion (A.3), we have assumed that $(rA'_0(r))/A_0(r)$ and $(rB'(r))/B(r)$ are finite for arbitrary r . As one can check, this is indeed the case for asymptotic AdS metric.

Substituting (A.3) into (A.1) and comparing terms proportional to v_{anom} in (A.1), we have:

$$L_1(\eta) = 2H \int_{r_c^0}^{\infty} dr \left\{ -\frac{[A_1(r) + A'_0(r)(r_c^1/r_c^0)r]}{2[B(r)]^{1/2}[A_0(r) - A_0(r_c^0)]^{3/2}} + \left(\frac{r_c^1}{r_c^0}\right) \left[\frac{2B(r) - B'(r)r}{2[B(r)]^{3/2}[A_0(r) - A_0(r_c^0)]^{1/2}} \right] \right\}, \quad (\text{A.4})$$

where we have replaced \tilde{r} with r/r_c^0 and H^2 with $A_0(r_c^0)$. Finally, from the definition of r_c (r_c^0) (3.5), i.e., $A(r_c) = A_0(r_c^0) = H^2$, we obtain:

$$r_c^1 = -\frac{A_1(r_c^0)}{A'_0(r_c^0)}. \quad (\text{A.5})$$

Putting (A.4) and (A.5) together, we arrived at (3.10).

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