Holographic response from higher derivatives with homogeneous disorder

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Abstract

In this letter, we study the charge response from higher derivatives over the background with homogeneous disorder introduced by axions. We first explore the bounds on the higher derivatives coupling from DC conductivity and the anomalies of causality and instabilities. Our results indicate no tighter constraints on the coupling than that over Schwarzschild-AdS (SS-AdS) background. And then we study the optical conductivity of our holographic system. We find that for the case with γ1 < 0 and the disorder strength ˆα < 2/√3, there is a crossover from a coherent to incoherent metallic phase as ˆα increases. When ˆα is beyond ˆα = 2/√3 and further amplified, a peak exhibits again at low frequency. But it cannot be well fitted by the standard Drude formula and new formula for describing this behavior shall be called for. While for the holographic system with the limit of γ1 → 1/48, the disorder effect drives the hard-gap-like at low frequency into the soft gap and suppresses the pronounced peak at medium frequency.

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I. INTRODUCTION

The quantum critical (QC) dynamics described by CFT or its proximity physics is strongly coupled systems without quasi-particles descriptions [1] (also refer to [2–10]). In dealing with these problems, a controlled manner in traditional field theory can not usually be performed. The AdS/CFT correspondence [11–14] provides valuable lessons to understand such systems by mapping certain CFTs to higher dimensional classical gravity. Studying the optical conductivity $\sigma(\omega/T)$ by introducing the probe Maxwell field coupled to the Weyl tensor $C_{\mu\nu\rho\sigma}$ in the Schwarzschild-AdS (SS-AdS) black brane background, we find that the behavior of conductivity is similar with one in the superfluid-insulator quantum critical point (QCP) described by the boson Hubbard model [15–17]. It provides possible route to access this kind of problems. Further, to test the robustness of higher derivative (HD) terms, the author of Ref.[18] studies the charge response of a large class of allowed HD terms in the SS-AdS geometry and find some interesting results. Of particular interest is that the optical conductivity displays an arbitrarily sharp Drude-like peak and the bounds found in [15, 19] are violated.

Also, we explore the charge transport with Weyl term in a specific thermal state with homogeneous disorder, which is introduced by a pair of spatial linear dependent axionic fields in AdS geometry and is away from quantum critical point (QCP), and new physics is qualitatively found [20]. Of particular interesting is that for the positive Weyl coupling parameter $\gamma > 0$, the strong homogeneous disorder drives the Drude-like peak in QCP state described by Maxwell-Weyl system in SS-AdS geometry [15] into the incoherent metallic state with a dip, which is away from QCP. While an oppositive scenario is found for $\gamma < 0$. In addition, the particle-vortex duality in the dual field theory induced by the bulk electromagnetic (EM) duality related by changing the sign of $\gamma$ is still preserved. Nonetheless, there is still bound for the conductivity as in [15, 19]. In this letter, we study the charge response of a large class of HD terms in the specific thermal state with homogeneous disorder. The letter is organized as follows. We describe the holographic setup for a class of HD theory with homogeneous disorder in Section II. And then the constraints imposing on the HD coupling parameters in the Einstein-axions-AdS (EA-AdS) geometry are explored in Section III. In Section IV, we mainly study the optical conductivity from HD theory with homogeneous disorder. Finally the conclusion and discussion are presented in Section V.
II. HOLOGRAPHIC SETUP

A specific thermal excited state with homogeneous disorder can be holographically described by the EA theory [21],

\begin{equation}
S_0 = \int d^4x \sqrt{-g} \left( R + 6 - \frac{1}{2} \sum_{I=x,y} (\partial \phi_I)^2 \right),
\end{equation}

where \( \phi_I = \alpha x_I \) with \( I = x, y \) and \( \alpha \) being a constant. In the action above, there is a negative cosmological constant \( \Lambda = -6 \), which supports an asymptotically AdS spacetimes \(^{1}\). The EA action (1) gives a neutral black brane solution \(^{2}\)

\begin{equation}
ds^2 = \frac{1}{u^2} \left( -f(u) dt^2 + \frac{1}{f(u)} du^2 + dx^2 + dy^2 \right),
\end{equation}

where

\begin{equation}
f(u) = (1 - u) p(u), \quad p(u) = \sqrt{1 + 6\hat{\alpha}^2 - 2\hat{\alpha}^2 - 1} u^2 + u + 1.
\end{equation}

\( u = 0 \) is the asymptotically AdS boundary while the horizon locates at \( u = 1 \). Note that we have parameterized the black brane solution by \( \hat{\alpha} = \alpha/4\pi T \) with the Hawking temperature \( T = p(1)/4\pi \). Although the momentum dissipates due to the break of microscopic translational symmetry, the geometry is homogeneous and so we refer to this mechanism as homogeneous disorder and \( \hat{\alpha} \) denotes the strength of disorder \(^{2}\).

Now, we consider the following action beyond Weyl [18]

\begin{equation}
S_A = \int d^4x \sqrt{-g} \left( -\frac{1}{8g_F^2} F_{\mu\nu} X^{\mu\nu\rho\sigma} F_{\rho\sigma} \right),
\end{equation}

where the tensor \( X \) is an infinite family of HD terms

\begin{align}
X^{\mu\nu\rho\sigma} &= I_{\mu\nu}^{\rho\sigma} - 8\gamma_{1,1} C_{\mu\nu}^{\rho\sigma} - 4\gamma_{2,1} C^2 I_{\mu\nu}^{\rho\sigma} - 8\gamma_{2,2} C_{\mu\nu}^{\alpha\beta} C_{\alpha\beta}^{\rho\sigma} \\
&\quad - 4\gamma_{3,1} C^3 I_{\mu\nu}^{\rho\sigma} - 8\gamma_{3,2} C^2 C_{\mu\nu}^{\rho\sigma} - 8\gamma_{3,3} C_{\mu\nu}^{\alpha_1\beta_1} C_{\alpha_1\beta_1}^{\alpha_2\beta_2} C_{\alpha_2\beta_2}^{\rho\sigma} + \ldots.
\end{align}

In the equation above, \( I_{\mu\nu}^{\rho\sigma} = \delta_{\mu}^{\rho} \delta_{\nu}^{\sigma} - \delta_{\mu}^{\sigma} \delta_{\nu}^{\rho} \) is an identity matrix and \( C^\alpha = C_{\mu\nu}^{\alpha_1\beta_1} C_{\alpha_1\beta_1}^{\alpha_2\beta_2} \ldots C_{\alpha_{n-1}\beta_{n-1}}^{\mu\nu} \). \( g_F^2 \) is an effective dimensionless gauge coupling and we set

\(^{1}\) We have set the AdS radius \( L = 1 \) without loss of generality.
\(^{2}\) The other models, for instance [22–33], have also been developed to produce the effect of homogeneous disorder.
$g_F = 1$ in the numerical calculation. The action (4) is constructed in terms of double EM field strengths, which is sufficient for linear response, coupled to any number of symmetry-allowed curvature tensors, which go beyond the Weyl action studied in [15, 19, 20, 34–48]. Note that the $X$ tensor possess the following symmetries

$$X_{\mu\nu\rho\sigma} = X_{[\mu\nu][\rho\sigma]} = X_{\rho\sigma\mu\nu}.$$  \hspace{1cm} (6)

When we set $X_{\mu\nu\rho\sigma} = I_{\mu\nu\rho\sigma}$, the theory (4) reduces to the standard Maxwell theory.

Since the action (4) is an infinite family of $n$ powers of the Weyl tensor $C$, we truncate it to the level $n = 2$, which is the 6 derivatives and the focus in this paper. On the other hand, since the effect of the coupling terms $\gamma_1$ and $\gamma_2$ is similar, we mainly focus on the term of $\gamma_1$ through this paper. For convenience, we denote $\gamma_1, 1 = \gamma$ and $\gamma_2, i = \gamma_i (i = 1, 2)$ in what follows.

### III. BOUNDS ON THE COUPLING

In this section, we examine the bounds on the coupling over the EA-AdS background (2). To this end, we turn on the perturbations of gauge field and write down the corresponding linearized perturbative equations in momentum space [15, 20]

$$A_t'' + \left( \frac{f'}{f} - \frac{X_1'}{X_1} + 2 \frac{X_3'}{X_3} \right) A_t' + \left( - \frac{p^2 \hat{\omega}^2 X_1}{f X_3} + \frac{X_1' X_3'}{X_1 X_3} - \frac{f' X_3}{f X_3} - \frac{X_1' X_3'}{X_1 X_3} \right) A_t = 0,$$

$$A_y'' + \left( \frac{f'}{f} + \frac{X_6'}{X_6} \right) A_y' + \frac{p^2}{f^2} \left( \omega^2 \frac{X_2}{X_6} - \hat{q}^2 \frac{X_1}{X_6} \right) A_y = 0,$$  \hspace{1cm} (7)

where $\hat{\mu} = (\hat{\omega}, \hat{q}, 0)$ are the dimensionless quantities defined as

$$\hat{\omega} \equiv \frac{\omega}{4\pi T} = \frac{\omega}{p}, \quad \hat{q} \equiv \frac{q}{4\pi T} = \frac{q}{p}, \quad p \equiv p(1) = 4\pi T.$$  \hspace{1cm} (8)

$X_i, i = 1, \ldots, 6$, are the matrix elements of $X^{AB}_A = \text{diag}(X_i(u))$ with $A, B \in \{tx, ty, tu, xy, xu, yu\}$, which encode the essential information of tensor $X_{\mu\nu\rho\sigma}$. In particular, $X_1(u) = X_2(u) = X_5(u) = X_6(u)$ and $X_3(u) = X_4(u)$ because of the symmetry of the background geometry (2) and the structure of $X$ tensor (5). Note that here we only consider the equations for $A_t$ and $A_y$ since $A_t' = -\frac{\hat{q} X_2}{X_3} A_x'$ and we have chosen the radial gauge as $A_u = 0$. In addition, the equations of EM dual theory can be obtained by letting $X_i \rightarrow \hat{X}_i = 1/X_i$. 
FIG. 1: The DC conductivity $\sigma_0$ versus $\hat{\alpha}$ for some fixed $\gamma_1$ and other coupling parameters vanishing.

A. Bounds from DC conductivity

In [18], it has been found that for the SS-AdS geometry and the subspace of $\gamma_1 \neq 0$ but other parameters vanishing, $\gamma_1$ is unconstrained from the anomalies of causality and instabilities. But an additional constraint from $Re\sigma(\omega) \geq 0$, especially DC conductivity being positive, gives $\gamma_1 \leq 1/48$. Here we shall further examine this bound over EA-AdS geometry from conductivity.

To this end, we write down the expression of DC conductivity [15, 19, 20]

$$\sigma_0 = \sqrt{-gg^{xx}} \sqrt{-g^{tt} g^{uu}} X_1 X_5 \big|_{u=1}.$$  \hfill (9)

It can be explicitly worked out as up to 6 derivatives

$$\sigma_0 = 1 - \frac{2}{3} \gamma f''(1) - \left(\frac{4}{3} \gamma_1 + \frac{1}{9} \gamma_2\right)f''(1)^2,$$  \hfill (10)

with

$$f''(1) = -2 - \frac{4(-1 - 2\hat{\alpha}^2 + \sqrt{1+6\hat{\alpha}^2})}{\hat{\alpha}^2}. \hfill (11)$$

We can easily find that for arbitrary $\hat{\alpha}$, the upper bound $\gamma_1 = 1/48$ preserves for other parameters vanishing. FIG.1 clearly shows this result.

Before closing this subsection, we would like to present some comments on the DC conductivity $\sigma_0$ from 6 derivative term. First, similar with the case of 4 derivative term [20], there is a specific value $\hat{\alpha} = 2/\sqrt{3}$, for which $\sigma_0$ is independent of the coupling parameter $\gamma_1$. Second, for the case of 4 derivatives, $\sigma_0$ is monotonic function of $\hat{\alpha}$ [20], while for one of 6 derivatives, $\sigma_0$ is non-monotonic and has upper/lower bound setting by $\sigma_0 = 1$. 

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B. Bounds from anomalies of causality and instabilities

To examine the causality and instabilities, we package the equations (7) into a more convenient Schrödinger form,

\[-\partial_z^2 \psi_i(z) + V_i(u) \psi_i(z) = \hat{\omega}^2 \psi_i(z),\]  

(12)

where \(dz/du = p/f\). Note that in order to remove linear terms of \(\partial_z\) and obtain the above equation, we have introduced an auxiliary functions \(G_i(u)\) to write

\[A_i(u) = G_i(u) \psi_i(u)\]

with \(A_{\bar{t}}(u) := A'_{\bar{t}}(u)\) and \(i = \bar{t}, y\). In addition, \(V_i(u)\) are the effective potentials, which are decomposed into the momentum dependent part and the independent one

\[V_i(u) = \hat{q}^2 V_{0i}(u) + V_{1i}(u),\]  

(13)

where

\[V_{0\bar{t}} = f \frac{X_1}{X_3}, \quad V_{0y} = f \frac{X_3}{X_1},\]  

(14)

\[V_{1\bar{t}} = \frac{f}{4p^2 X_1^2} \left[3f(X'_1)^2 - 2X_1(f X'_1)\right],\]  

(15)

\[V_{1y} = \frac{f}{4p^2 X_1^2} \left[-f(X'_1)^2 + 2X_1(f X'_1)\right].\]  

(16)

Since \(V_{\bar{t}} = V_y|_{X_1 \rightarrow \hat{X}_1}\), we only consider the case of \(V_{\bar{t}}\) in what follows.

Now we consider the subspace of \(\gamma_1 \neq 0\) but other parameters vanishing. Since \(X_1 = X_3\) in this case, one has \(V_{0\bar{t}} = f(u)\), which satisfies the constraint in the limit of large momentum \((\hat{q} \rightarrow \infty)\),

\[0 \leq V_{0\bar{t}}(u) \leq 1,\]  

(17)

for any \(\gamma_1\) and \(\hat{\alpha}\). While for the small momentum region, although the dominative potential \(V_{1\bar{t}}\) develops a negative minimum (see top plots in FIG.2), there are no unstable modes in these regions by analyzing the zero energy bound state in the potential \(V_{1\bar{t}}\) (see bottom plots in FIG.2), which is [49]

\[\tilde{n}_{1\bar{t}} = I/\pi + 1/2, \quad I \equiv \left(n - \frac{1}{2}\right) \pi = \int_{u_0}^{u_1} \frac{p}{f(u)} \sqrt{-V_{1\bar{t}}(u)} du,\]  

(18)

where \(n\) is a positive integer and the interval \([u_0, u_1]\) defines the negative well.

After examining the instabilities for small and large momentum limit, we turn to the case for some finite momentum. For this aim, we show the potentials \(V_{\bar{t}}(u)\) for different \(\gamma_1\) and
FIG. 2: Top plots: The potential $V_{1\bar{f}}$ for sample $\gamma_1$ and $\hat{\alpha}$. Bottom plots: $\bar{n}_{1\bar{f}}$ as the function $\hat{\alpha}$ for given $\gamma_1$. They clearly exhibit that $\bar{n}_{1\bar{f}}$ is always less than unit.

FIG. 3: The potentials $V_{\bar{f}}(u)$ with different $\gamma_1$ and $\hat{\alpha}$ at some finite momentum.
FIG. 4: Left plot: The optical conductivity with $\gamma_1 = -1$ for different $\hat{\alpha}$. Right plot: The optical conductivity with $\gamma_1 = 0.02$ for different $\hat{\alpha}$.

$\hat{\alpha}$ at some finite momentum in FIG. 3. It is obvious that the potential is always positive, which indicates that no unstable modes appear even for the finite momentum.

Therefore, we conclude that the region $\gamma_1 \in (-\infty, +\infty)$ is still physically viable for EA-AdS geometry. Combining with the constraint from conductivity, we have $\gamma_1 \in (-\infty, 1/48]$.

IV. CONDUCTIVITY

In this section, we study the optical conductivity by solving the second equation in (7) for $\hat{q} = 0$ with ingoing boundary condition at horizon. And then it can be read off in terms of

$$\sigma(\omega) = \frac{\partial_u A_y(u, \hat{\omega}, \hat{q} = 0)}{i\omega A_y(u, \hat{\omega}, \hat{q} = 0)}.$$  \hfill (19)

As has been shown in [18], for $\gamma_1 = -1$ and $\hat{\alpha} = 0$, the optical conductivity $\sigma(\omega)$ exhibits a sharp Drude-like peak at small frequency, then a gap and eventually rises and saturates to the value $\sigma = 1$ at large frequency (left plot in FIG. 4). While for $\gamma_1 = 0.02$ and $\hat{\alpha} = 0$, it displays a hard-gap-like behavior at small frequency, then a pronounced peak before saturating (right plot in FIG. 4). The holographic result is similar with that of the $O(N) NL\sigma M$ in the large-$N$ limit [2] as pointed out in [18].

Now we study the effect of the homogeneous disorder. For small $\hat{\alpha} \ll 1$, we find that the transport is well described by the standard Drude formula (see FIG. 5)

$$\sigma(\hat{\omega}) = \frac{K\tau}{1 - i\hat{\omega}\tau}. \hfill (20)$$

TABLE I shows some examples of the relaxation time $\tau$ for different $\hat{\alpha}$. It decreased with the increase of $\hat{\alpha}$, indicating that the dissipation of this system gradually becomes obvious. With
FIG. 5: The low frequency behavior of the optical conductivity. The real line is fitted with the standard Drude formula (20).

FIG. 6: The low frequency behavior of the optical conductivity. The real line is fitted with the standard Drude formula (20).

\( \hat{\alpha} \) further increasing before reaching \( \hat{\alpha} = 2/\sqrt{3} \), the Drude-like peak disappears, indicating a crossover from a coherent to an incoherent metallic phase. After \( \hat{\alpha} \) is beyond \( \hat{\alpha} = 2/\sqrt{3} \), with it increasing further, a peak exhibits again (left plot in FIG.4). It is different from that of 4 derivatives showing in [20], in which the low frequency conductivity of the holographic system with \( \gamma > 0 \) but other coupling parameter vanishing exhibits a dip as \( \hat{\alpha} \) becomes larger. But the peak is not well fitted by the standard Drude formula (see FIG.6). It indicates that we need new formula to fit it and we shall further explore it in future. While for \( \gamma_1 = 0.02 \) and \( \hat{\alpha} \neq 0 \), a soft gap exhibits at low frequency replacing the hard gap because of the effect of homogeneous disorder. And the sharp resonant peak at medium frequency for \( \hat{\alpha} = 0 \) gradually decreases with the increase of \( \hat{\alpha} \). We expect that the \( O(N) \) \( NL\sigma M \)
in the large-$N$ limit [2] exhibits some similar phenomena when some effects of disorder are introduced. We shall explore it in future.

| $\hat{\alpha}$ | 0   | 0.01 | 0.1 | 0.5 |
|----------------|-----|------|-----|-----|
| $\tau$        | 87.62 | 86.64 | 83.67 | 35.27 |

**TABLE I**: The relaxation time $\tau$ of the holographic system with $\gamma_1 = -1$ fitted by the standard Drude formula (20) for different $\hat{\alpha}$.

V. CONCLUSION AND DISCUSSION

In this letter, we mainly study the charge response from higher derivatives, in particular the 6 derivatives, over the homogeneous disorder background. In [18], it has been shown that the holographic result from higher derivatives is similar with that of the $O(N) \ N L \sigma M$ in the large-$N$ limit [2]. When the homogeneous disorder effect by axions is introduced into this holographic system with $\gamma_1 < 0$ (small enough, for example $\gamma_1 = -1$), we find that before $\hat{\alpha}$ reaching $\hat{\alpha} = 2/\sqrt{3}$, there is a crossover from a coherent to incoherent metallic phase with $\hat{\alpha}$ increasing. As $\hat{\alpha}$ is beyond $\hat{\alpha} = 2/\sqrt{3}$ and further amplified, a peak exhibits again at low frequency. But it cannot be well fitted by the standard Drude formula and call for new formula. While the holographic system with the limit of $\gamma_1 \rightarrow 1/48$ is an insulator with a hard-gap-like at low frequency and a pronounced peak emerging at medium frequency. But the homogeneous disorder effect drives the hard-gap-like at low frequency into the soft gap and suppresses the pronounced peak at medium frequency. We expect that when some effects of disorder are introduced into the $O(N) \ N L \sigma M$ in the large-$N$ limit, some similar phenomena can be observed, which deserves further studying.

It is also interesting and illuminating to compare the main properties of the conductivity from 6 derivatives ($\gamma = 0$ and $\gamma_1 \neq 0$) here with that from 4 derivatives ($\gamma \neq 0$ and $\gamma_1 = 0$) in our previous work [20]. First, for the case of 4 derivatives, since the causality results to a bound on $\gamma$, the transport is incoherent for any $\hat{\alpha}$. While for the case of 6 derivatives, there is coherent charge transport with long-lived excitations for $\gamma_1 < 0$ and small $\hat{\alpha}$. Second, there is a dip in optical conductivity with $\gamma < 0$ and $\gamma_1 = 0$ for small $\hat{\alpha}$, which is similar with the excitation of vortices. But for the case with $\gamma_1 > 0$ and $\gamma = 0$, the optical conductivity exhibits a hard-gap-like or soft gap at low frequency and a resonant
peak at medium frequency. Third, for the case of 4 derivatives, the homogeneous disorder effect drives the peak (dip) in the low frequency optical conductivity into its contrary. But for the case of 6 derivatives, the homogeneous disorder effect cannot drive the peak (gap) into its contrary. It is an interesting phenomenon and deserves further exploring for other higher derivative terms.

There are lots of open questions deserving further exploration. For example, we can study the charge responses from higher derivatives with homogeneous disorder at full momentum and energy spaces. Also we can test the robustness of the phenomena observed here when the disorder is introduced by the other way, for instance [22–33].

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