HOLOGRAPHIC THERMAL CORRELATORS: ELECTROMAGNETIC AND GRAVITATIONAL PERTURBATIONS OF THE RN-AdS$_4$ GEOMETRY

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We are currently exploring axial perturbations in the linearized Einstein field equations for a Reissner-Nordström black hole in four-dimensional anti-de Sitter space. A framework is constructed to study holographic thermal two-point correlation functions, with particular focus on the one related to conductivity. Preliminary non-perturbative numerical results are presented on their dependence on different parameters, along with physical interpretation and agreement with the expected analytic asymptotics of the correlator.

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INTRODUCTION

Holographic duality, often referred to as the AdS/CFT correspondence, is a framework based on the conjecture that weakly coupled gravitational theory is equivalent to strongly coupled field theory in the large N limit (where N intuitively represents the effective number of degrees of freedom). It was proposed by Maldacena in 1997 [1], and a proper mathematical formulation was developed in 1998 by Gubser, Klebanov, and Polyakov [2] and independently by Witten [3] (GKPW). Using GKPW rules, one can show that we can construct a duality between a QFT in d dimensions with the boundary of $(d+1)$-dimensional bulk gravitational theory. Over the last several decades, numerous theoretical checks and simulations have supported the conjecture.

This field is actively developing due to the possibility of studying strong-coupling effects in a non-perturbative manner. It is actively used in studies of string theory [1], strongly interacting quark-gluon plasma [4–6], and high-temperature superconductors [7, 8].

The interest in the Reissner-Nordström anti-de Sitter (RN-AdS) metric comes from the properties that it endows to dual field theory. Insertion of the black hole into AdS space is equivalent to considering a field theory with finite temperature. If one adds the charge for the black hole, it corresponds to implementing a finite chemical potential. Combining all together, we obtain quite a peculiar, strongly coupled field theory with temperature and chemical potential. It can be useful to study this model itself [9]. In the context of quantum matter, the RN-AdS metric is usually utilized as a background for the description of a holographic superconductor [7, 10].

In this work, we explore the perturbation in the four-dimensional RN-AdS metric and apply holographic duality to the study of dual field theory. In particular, we are interested in two-point correlation functions that describe the dynamics of field theory and usually can be identified with a clear physical interpretation (e.g., electric, thermal, and thermoelectric conductivities).

The point of interest is a holographic analysis of the so-called Gertsenshtein-Zel'dovich (GZ) effect, which is observed in the bulk gravitational theory and closely related to these conductivities. In 1961 [11], Gertsenshtein proposed the possibility of resonance between electromagnetic (EM) waves and gravitational waves (GWs) due to the same speed of propagation of these waves. He used the linearized Einstein field equations to prove that GWs produce EM waves through wave resonance as they move through a strong magnetic field. Similarly, EM waves passing through a strong magnetic field produce GWs [12]. While the analysis is purely classical, in quantum language, the external field can be viewed as 'catalyzing' a resonant conversion between photon and graviton states, much like neutrino flavor mixing. In our case, the role of "catalyst" is played by a charged black hole [13], and we study some properties of the discussed effect. This proceeding presents work in progress and preliminary analysis available only for electric conductivity, and we plan to cover the other two correlators in the near future.

1. SETUP

To consider a system with a charged black hole, the initial action should include gravitational and electromagnetic interactions. The simplest form that one can consider is the Einstein-Maxwell action that has the form:

$$S = \frac{1}{16\,\pi}\int d^d x\,\sqrt{-g}\,\left[R - 2\Lambda - F_{\mu\nu}F^{\mu\nu}\right], \qquad (1)$$

where $g_{\mu\nu}$ is the metric and g is its determinant, R is the Ricci scalar, Λ is the cosmological constant, and $F_{\mu\nu}$ is the electromagnetic tensor.

By varying this action with respect to metric, one can obtain the following Einstein's field equations (EFE) in the presence of the electromagnetic field:

$$R_{\mu\nu} - \frac{1}{2}Rg_{\mu\nu} + \Lambda g_{\mu\nu} = 8\,\pi T_{\mu\nu} = = 2\left(g^{\lambda\sigma}F_{\mu\lambda}F_{\nu\sigma} - \frac{1}{4}g_{\mu\nu}F_{\lambda\sigma}F^{\lambda\sigma}\right), \qquad (2)$$

where we define energy-momentum tensor as:

$$T_{\mu\nu} = \frac{1}{4\pi}\left(g^{\lambda\sigma}F_{\mu\lambda}F_{\nu\sigma} - \frac{1}{4}g_{\mu\nu}F_{\lambda\sigma}F^{\lambda\sigma}\right). \qquad (3)$$

Taking into account the definition of electromagnetic tensor via potential $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, and varying with respect to electromagnetic potential, one can obtain other equations of motion:

$$\partial_\mu\left(\sqrt{-g}F^{\mu\nu}\right) = \nabla_\mu F^{\mu\nu} = 0. \quad (4)$$

We are interested in the spherically symmetric vacuum solution of charge Q and mass M in hyperbolic space with constant negative curvature $\Lambda < 0$. As we mentioned previously, this solution is the Reissner-Nordström anti-de Sitter metric. In four dimensions ($d = 4$), RN-AdS$_4$ metric can be written as (system of units: $\hbar = c = G = e^2/4\pi = 1$):

$$ds^2 = -f(r)dt^2 + f(r)^{-1}dr^2 + r^2 d\Omega_2^2, \quad (5)$$

where $f(r) = 1 - \frac{2M}{r} + \frac{Q^2}{r^2} - \frac{1}{3}\Lambda r^2$.

From the symmetry of the solution, we can see that the value of the electromagnetic potential can depend only on r, and only A_t is a non-vanishing component, which has the form:

$$A_t = \frac{Q}{r} - \frac{Q}{r_+}, \quad (6)$$

where r_+ is the outer event horizon, combination $Q/r_+ = \mu$ is a vanishing condition on the event horizon, and also plays the role of the chemical potential in holographic duality.

2. SOLVING LINEARIZED EQUATIONS

Let us derive the linearized gravitational wave equations for the RN black hole in AdS space. Consider small linear perturbation of the background metric $\bar{g}_{\mu\nu}$ and electromagnetic potential $\bar{A}_\mu$, which we can write as

$$g_{\mu\nu} = \bar{g}_{\mu\nu} + h_{\mu\nu}, \quad (7)$$

$$A_\mu = \bar{A}_\mu + \alpha_\mu, \quad (8)$$

where the perturbation $h_{\mu\nu}$ goes to zero at infinity fast enough that the metric $g_{\mu\nu}$ is asymptotically $\bar{g}_{\mu\nu}$. Furthermore, $h_{\mu\nu}$ and its derivatives are assumed to be small enough so that we can ignore higher-order terms.

Additionally, to leave only physical degrees of freedom, we should implement Lorenz and transverse-traceless (TT) gauges for these perturbations. These conditions on the fields take the next form:

$$\nabla_\mu h^{\mu\nu} = 0, \quad h^\mu_\mu = 0, \quad \nabla_\mu \alpha^\mu = 0. \quad (9)$$

We now have all the necessary ingredients to derive the linearized system of equations. Substituting Eqs. (7) and (8) into Eqs. (2) and (4), taking into account Eq. (9), and subtracting the background solution, we obtain the desired expressions. We omit the intermediate steps and proceed directly to the result. In the linear approximation, the general form of the axial perturbations can be expressed as (all other components vanish):

$$\begin{aligned} h_{t\phi} &= h_{\phi t} = e^{-i\omega t}w(r,\theta), \\ h_{r\phi} &= h_{\phi r} = e^{-i\omega t}q_2(r,\theta), \quad (10) \\ h_{\theta\phi} &= h_{\phi\theta} = e^{-i\omega t}q_3(r,\theta), \end{aligned}$$

$$\alpha_\phi = e^{-i\omega t}B(r,\theta), \quad (11)$$

where we explicitly write the time-dependence for the ingoing wave and assume momentum is zero. From the Lorenz gauge equations (9), we can always fix one component, and we fix the value of $w(r,\theta)$.

Implementing these variables into the linearized perturbation equations, one obtains the system of four equations. We can eliminate the angular dependence and simplify equations significantly by employing the symmetry of the problem and finding the solution with the following ansatz [14]:

$$q_2(r,\theta) = \frac{3\tilde{q}_2(r)}{L^2 + L - 2} C_{L+1}^{(-1/2)}(\cos\theta), \quad (12)$$

$$q_3(r,\theta) = r^2\tilde{q}_3(r)\frac{C_{L+2}^{(-3/2)}(\cos\theta)}{\sin\theta}, \quad (13)$$

$$B(r,\theta) = \frac{-3\tilde{B}(r)}{2\omega i\sqrt{L^2 + L - 2}} C_{L+1}^{(-1/2)}(\cos\theta), \quad (14)$$

where $C_n^\mu(\cos\theta)$ denotes the Gegenbauer polynomials, $L \in 2,3 \dots N$ and plays role of orbital quantum number. It is convenient to rewrite $\tilde{q}_2(r)$ in the form:

$$\tilde{q}_2(r) = -r^2\left(\frac{rA(r)}{\Delta} + \tilde{q}_{3,r}(r)\right), \quad (15)$$

where $\Delta = r^2 f(r)$.

We can rewrite the equations as a combination of $A(r)$ and $\tilde{B}(r)$ (hereafter omitting the tilde and the explicit r-dependence), where the former variable represents the gravitational perturbation and the latter the electromagnetic one. These equations can be diagonalized by introducing variables:

$$\begin{aligned} Z_1 &= (-p_1p_2)^{1/2}A + p_1B, \quad (16) \\ Z_2 &= p_1A - (-p_1p_2)^{1/2}B, \\ p_1 &= 3M + \sqrt{9M^2 + 4Q^2(L^2 + L - 2)}, \\ p_2 &= 3M - \sqrt{9M^2 + 4Q^2(L^2 + L - 2)}. \end{aligned}$$

And the final result is two one-dimensional Schrodinger-type wave equations [15]:

$$\boldsymbol{\Lambda}^2 Z_i = V_i^- Z_i, (i = 1,2), \quad (17)$$

with the potential of the form:

$$V_i = \frac{\Delta}{r^5}\left(L(L+1)r - p_j\left(1 + \frac{p_i}{(L^2 + L - 2)r}\right)\right), \quad (18)$$

$$(i,j = 1,2; i \neq j),$$

and where $\boldsymbol{\Lambda}^2 = \frac{d}{dr^{*2}} + \omega^2$ and $dr^* = r^2\Delta^{-1}dr$ are Eddington–Finkelstein coordinates.

This result coincides with one from [14, 16]. The dependence of Z_1 and Z_2 on both A and B describes the mixing induced by the GZ effect and encodes the interaction between the gravitational and electromagnetic fields. In the special case $Q = 0$, we observe the decoupling of the equations, meaning that the charge of the black hole in this configuration acts as a catalyst for the GZ effect, since without it, this effect disappears. This limit corresponds to the Schwarzschild-AdS$_4$ black hole, for which gravitational and electromagnetic perturbations change separately.

We should impose the physical boundary conditions on the horizon of the black hole at r_+ to solve (17). The potential $V(r) \sim f(r) \sim 0$ at the event horizon. So, near the horizon, one expects the following behavior of the solution:

$$Z(r) \sim e^{-i\omega r_*} \sim (r - r_+)^{\frac{-i\omega}{f'(r_+)}},\ r_* \to -\infty, \quad (19)$$

where the minus sign in the exponent corresponds to ingoing waves, since we expect only incoming waves to be possible.

We want to eliminate the rapidly oscillating exponent near the horizon that strongly influences the accuracy of the numerical solution. Taking into account (19), one can use the following substitution:

$$Z(r) = e^{-i\omega r_*}\psi(r). \quad (20)$$

Exploring asymptotical behavior in the limit $r \to +\infty$, one finds that the solution of the equation converges to:

$$\psi(r) \to \psi_{(0)} + \frac{\psi_{(1)}}{r} + \cdots,\quad r \to +\infty, \quad (21)$$

where $\psi_{(0)}$ and $\psi_{(1)}$ are non-normalizable and normalizable parts of the solution.

3. CALCULATIONS OF THE HOLOGRAPHIC TWO-POINT FUNCTIONS

The holographic duality is a conjecture that connects classical gravitational bulk theory with strongly-coupled field theory on its boundary. As mentioned in the introduction, the main pillar of this idea is captured in the GKPW formula:

$$\left\langle e^{\int d^d x\, J(x)\mathcal{O}(x)} \right\rangle_{QFT} = \int D\phi\, e^{iS_{bulk}|_{\phi(x,r=\infty)=J(x)}}. \quad (22)$$

Here the field theory lives on the boundary, i.e, in the asymptotic limit $r \to +\infty$.

Consider $A(t,\theta)$ is a leading-order (non-normalizable) boundary value of a field in gravitational theory in the limit of $r \to +\infty$, meaning that $\phi(x, r = \infty) = A(x) = J(x)$, i.e corresponds to the source of the boundary operator $\mathcal{O}(x)$. Field theory has one less dimension than the initial gravitational theory, so we consider $d = 3$ QFT. The meaning of the sub-leading (normalizable) part $B(x)$ can be found in the following relation:

$$\langle \mathcal{O}_i(x) \rangle = i \frac{\delta}{\delta A(x)} S_{bulk}|_{\phi(x,r=\infty)=A(x)} \propto B(x), \quad (23)$$

where the normalizable term corresponds to the vacuum expectation value (VEV) of the boundary operator $\mathcal{O}_i(x)$.

We can write the relation between source $A_{\mathcal{O}_B}(x)$ and response $B_{\mathcal{O}_A}(x)$ using the retarded Green's function, which follows from the theory of linear response and the Kubo formula:

$$\int d^d y\, A_{\mathcal{O}_B}(y) G^R_{\mathcal{O}_A \mathcal{O}_B}(x,y) = B_{\mathcal{O}_A}(x), \quad (24)$$

where you can read off $B_{\mathcal{O}_A}(x)$ and $A_{\mathcal{O}_B}(y)$ from the large r behavior of the solution and $d^d y = dt d\Omega$ is $\mathbb{R} \times S^2$ volume element (our boundary field theory has this dimension). Also, we are implicitly summing over all operators involved and over coordinate indices.

The system can be presented as a matrix relation in momentum space. Using an appropriate form of Green's functions with the Fourier transform and extension of perturbations at the boundaries as:

$$\alpha_{\phi(0/1)}(x) = e^{-i\omega t}\alpha_{\phi(0/1)}(\omega) C^{(-1/2)}_{L+1}(\cos\theta), \quad (25)$$

$$h_{\theta\phi(0/1)}(x) = e^{-i\omega t} h_{\theta\phi(0/1)}(\omega) \frac{C^{(-3/2)}_{L+2}(\cos\theta)}{\sin\theta},$$

one can obtain the following relation

$$\begin{pmatrix} \alpha_{\phi(1)}(\omega) \\ h_{\theta\phi(1)}(\omega) \end{pmatrix} = \\ = \begin{pmatrix} G^R_{\alpha_\phi,\alpha_\phi}(\omega) & G^R_{\alpha_\phi,h_{\theta\phi}}(\omega) \\ G^R_{h_{\theta\phi},\alpha_\phi}(\omega) & G^R_{h_{\theta\phi},h_{\theta\phi}}(\omega) \end{pmatrix} \begin{pmatrix} \alpha_{\phi(0)}(\omega) \\ h_{\theta\phi(0)}(\omega) \end{pmatrix}. \quad (26)$$

To obtain frequency dependence $\alpha_{\phi(0/1)}(\omega)$ and $h_{\theta\phi(0/1)}(\omega)$, one should expand the large r behavior at the boundary:

$$h_{\theta\phi}(r) \sim r^2 h_{\theta\phi(0)}(\omega) + c_1(\omega) + \frac{h_{\theta\phi(1)}(\omega)}{r}, \quad (27)$$

$$\alpha_\phi(r) \sim \alpha_{\phi(0)}(\omega) + \frac{\alpha_{\phi(1)}(\omega)}{r}, \quad (28)$$

$$h_{r\phi}(r) \sim \frac{c_2(\omega)}{r} + \frac{c_3(\omega)}{r^3}, \quad (29)$$

$$h_{t\phi}(r) \sim r^2 h_{t\phi(0)}(\omega) + c_4(\omega) + \frac{h_{t\phi(1)}(\omega)}{r}. \quad (30)$$

We indeed have all the constants and dependence on $\psi_{1,2(0)}$ and $\psi_{1,2(1)}$, although their explicit forms are rather cumbersome and offer little practical insight. Accordingly, we omit them here. But we should say several comments, the constants c_1 and c_4 depend on the sources of the master's equations $\psi_{1(0)}$ and $\psi_{2(0)}$ in the same way as $h_{\theta\phi(0)}$ and $h_{t\phi(0)}$. Therefore, contain no additional information and cannot be interpreted as independent response coefficients; to obtain the physical response, one must consider the next-order terms in the expansion. In the case of $h_{r\phi}$, we find no non-renormalizable contribution corresponding to a source term, implying that this component does not contribute at the boundary. The constants c_2 and c_3 need not be specified.

It is also important to note that although we have three non-zero components, only two of them are independent. We choose $h_{t\phi}$ to be dependent on $h_{\theta\phi}$ and α_ϕ, so the Green's function matrix is consistently represented as a 2×2 matrix, which can also be used to determine the Green's functions associated with $h_{t\phi}$.

4. RESULTS

Now turn to the numerical evaluation of the two-point function for the RN-AdS$_4$ black hole. We begin with the uncharged case, $Q = 0$, which reduces to the Schwarzschild-AdS$_4$ geometry, and then examine $Q = 1.8$ (preliminary) to illustrate how the charge modifies the dynamics. Throughout our analysis, we fix the horizon radius and cosmological constant at $r_+ = 1$ and $\Lambda = -3$. For these parameter values, the extremal charge is $Q_{ext} = 2$.

The plots exhibit only minor variations as the charge is tuned. The most pronounced deviations appear in the intermediate-frequency regime, where both the hydrodynamic approximation and the operator product expansion (OPE) break down.

The dashed lines on the real part of the Green's function correspond to the result that we obtain for the photon in the AdS space. This approximation is only applicable up to the first pole. From a gravitational point of view, to interact with a black hole, the perturbation must have enough energy to go through the centrifugal barrier; for lower energies (the well region), one expects similar behavior to pure AdS space.

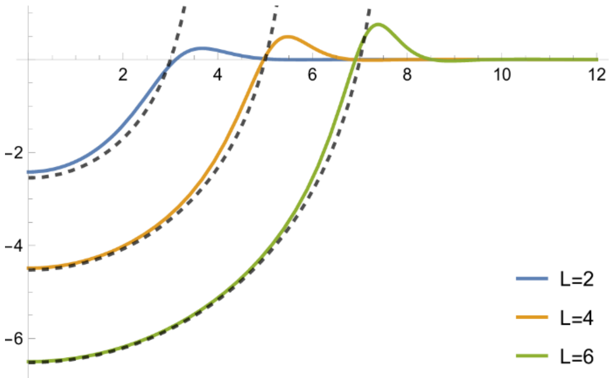

Fig. 1. The real part of Green's function $G^R_{\alpha_\phi,\alpha_\phi}$ for $Q=0$. The dashed lines correspond to the case of the photon in pure AdS space

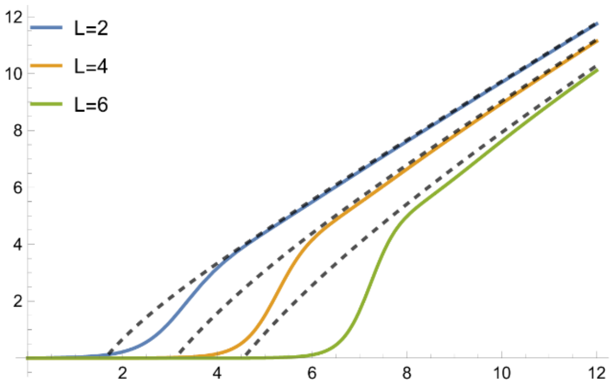

Fig. 2. The imaginary part of Green's function $G^R_{\alpha_\phi,\alpha_\phi}$ for $Q=0$. The dashed lines correspond to the (31)

When frequencies become significant enough to cross the barrier, the approximation breaks, which makes sense, as a black hole absorbs the energy, making it impossible to form resonant normal modes. Instead, one can observe a smooth finite peak that corresponds to the quasinormal modes that decay with time.

The plots show that larger L and Q are better described by the approximation, which is yet to be understood. The smoothness and height of peaks also depend on L and Q.

The negative value of the Green's function at low frequencies implies that the system is a perfect conductor. Moreover, the increasing depth of the Green's function with a higher angular momentum quantum number L reveals that modes with larger L exhibit even stronger conductive behavior. This enhancement can be understood in terms of the centrifugal barrier, which dominates the low-frequency gravitational potential and becomes more pronounced as L increases, effectively deepening the potential well for higher-L modes. This also explains the increase in energy required for higher L perturbations to cross the black hole's barrier. At very high frequencies, one expects that all the perturbations can penetrate the black hole's barrier and are eventually dissipated by the imaginary part.

Using a photon in the AdS space approximation, we find the high-frequency behavior of the imaginary part:

$$G^R_{\alpha_\phi,\alpha_\phi} \sim \frac{2}{\pi^2}\omega + \frac{\pi\big(2-3L(1+L)\big)}{6\omega}, \omega \to +\infty. \quad (31)$$

These asymptotic results are presented as a dashed line on the imaginary part of Green's function. The formula does not depend on the charge Q, meaning that in the high-energy limit, the charge of the black hole is irrelevant in this approximation. As high-frequency perturbations can propagate freely inside the system and eventually end up in the black hole, leading to dissipation [17].

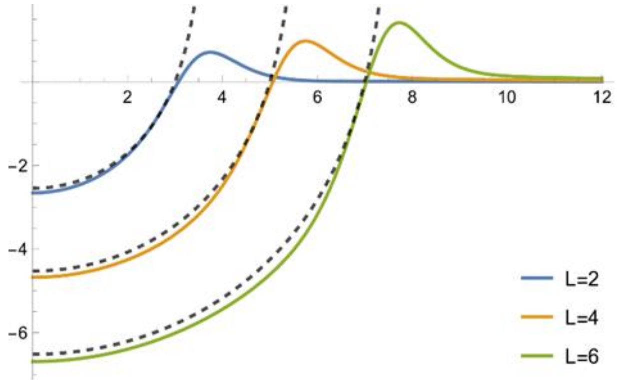

Fig. 3. The real part of Green's function $G^R_{\alpha_\phi,\alpha_\phi}$ for $Q=1.8$ (preliminary). The dashed lines correspond to the case of the photon in pure AdS space

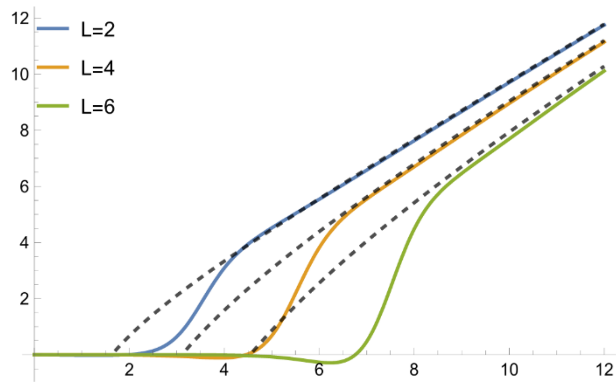

Fig. 4. The imaginary part of Green's function $G^R_{\alpha_\phi,\alpha_\phi}$ for $Q=1.8$ (preliminary). The dashed lines correspond to the (31)

One can also consider the hydrodynamic limit to study the low-frequency behavior of the imaginary part. From the hydrodynamic analysis [7, 18], one can find the following relation:

$$\mathrm{Im}\left[G^R_{\alpha_\phi,\alpha_\phi}\right] \sim \omega\sigma_Q, \quad (32)$$

where σ_Q is a constant calculated from theory, and low-frequency behavior for our system indeed agrees with this limit.

In this holographic model, the boundary theory behaves as a perfect conductor closely analogous to graphene [18], albeit defined on a spherical geometry. One can also find the Drude peak at $\omega = 0$ for the real part of the conductivity, which is not apparent in the plot, but can be calculated via Kramers-Kronig relations.

CONCLUSIONS

In this work, we obtained analytical solutions for axial perturbations of the metric and electromagnetic potential in the RN-AdS$_4$ background within the framework of linearized gravity. We further investigate the structure of the unperturbed metric and highlight the emergence of the Gertsenshtein-Zel'dovich effect.

After imposing ingoing boundary conditions at the horizon, we turned our attention to the principles of holographic duality. We demonstrated how the leading and subleading terms in the near-boundary expansion encode the source and the vacuum expectation value, and connected them, using linear response theory and the Kubo formula, with the two-point correlation function of the dual operator. Finally, we extracted the asymptotic behavior of both the metric and electromagnetic perturbations at the AdS boundary.

Concentrating on the Green's function $G^R_{\alpha_\phi,\alpha_\phi}$, which directly governs the conductivity of the boundary theory, we numerically computed its real and imaginary parts for different values of the charge. These preliminary results were then benchmarked against theoretical predictions in two asymptotic regimes: the hydrodynamic and the OPE limits. As expected, the low-frequency behavior of the real part of the correlator closely reproduces that of an electromagnetic perturbation in pure AdS. Meanwhile, the imaginary part exhibits excellent agreement with both the hydrodynamic and OPE predictions. In the intermediate-frequency window, we uncover a nontrivial behavior of the dual, strongly coupled system.

Looking ahead, we plan to cross-check and extend our analysis to the remaining two-point correlation functions introduced in this work. In particular, we are interested in the non-diagonal Green's function that represents the interaction between different operators and, for our system, can be interpreted as a consequence of the GZ effect.

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ГОЛОГРАФІЧНІ КОРЕЛЯЦІЙНІ ФУНКЦІЇ ЗА СКІНЧЕННОЇ ТЕМПЕРАТУРИ: ЕЛЕКТРОМАГНІТНІ ТА ГРАВІТАЦІЙНІ ЗБУРЕННЯ 4-ВИМІРНОЇ ЗАРЯДЖЕНОЇ ЧОРНОЇ ДІРИ В АНТИДЕСІТТЕРІВСЬКОМУ ПРОСТОРІ

А.В. Анатайчук, О.В. Жибоєдов

Ми досліджуємо аксіальне збурення в лінеаризованих польових рівняннях Ейнштейна для зарядженої чорної діри в чотиривимірному антідесіттерівському просторі. Ми будуємо систему для вивчення голографічних двоточкових кореляційних функцій, де ми зосереджуємося на тій, що пов'язана з провідністю. Представлено попередні непертурбативні числові результати, їх залежність від різних параметрів, їх фізична інтерпретація та їх узгодження з очікуваними аналітичними асимптотиками корелятора.