

Holographic Probes of Anti-de Sitter Spacetimes

Vijay Balasubramanian^{1*}, Per Kraus^{2†},
Albion Lawrence^{1‡} and Sandip P. Trivedi^{3§}

¹*Lyman Laboratory of Physics, Harvard University
Cambridge, MA 02138, USA*

²*California Institute of Technology
Pasadena, CA 91125, USA*

³*Fermi National Accelerator Laboratory
P.O. Box 500
Batavia IL, 60510, USA*

Abstract

We describe probes of anti-de Sitter spacetimes in terms of conformal field theories on the AdS boundary. Our basic tool is a formula that relates bulk and boundary states – classical bulk field configurations are dual to expectation values of operators on the boundary. At the quantum level we relate the operator expansions of bulk and boundary fields. Using our methods, we discuss the CFT description of local bulk probes including normalizable wavepackets, fundamental and D-strings, and D-instantons. Radial motions of probes in the bulk spacetime are related to motions in scale on the boundary, demonstrating a scale-radius duality. We discuss the implications of these results for the holographic description of black hole horizons in the boundary field theory.

1 Introduction

A recurring theme of recent work is that gravitational theories can sometimes be formulated as gauge theories in fewer dimensions. This point of view has had some encouraging successes, but we still do not understand how the famous problems of

*vijayb@curie.harvard.edu

†perklaus@theory.caltech.edu

‡lawrence@string.harvard.edu

§trivedi@fnth23.fnal.gov

quantum gravity – for example, information loss in black hole evaporation – are solved. All of our intuitions about gravity and spacetime physics are based on a classical, geometric picture valid when \hbar is small and the field configurations macroscopic. In this regime the spacetime physics displays at least approximate locality and causality, and a well-defined geometry in which free particles follow geodesics. These properties seem obscure in the gauge theory formulation. Once we understand their origin, we can investigate precisely when and how they break down.

An important avenue for understanding these issues is the most recent manifestation of the gravity-gauge theory connection: the conjecture [1] that string theory on an anti-de Sitter background is dual to a conformal field theory living on the spacetime boundary. So far, the proposal has been checked by comparing spectra and low-order correlation functions of the dual theories. Such checks are based on a remarkably compact and powerful statement of the equality between certain path integrals in the dual theories [2, 3]. We would like to use this equality to learn how the classical geometric description of the bulk emerges and ultimately breaks down in the holographic boundary representation.

This article begins such a study by describing a variety of spacetime probes from the boundary perspective. Our basic tool is a compact formula relating bulk and boundary states; specifically, the asymptotic behavior of fluctuating, classical bulk fields is related to expectation values of the dual boundary operators in excited boundary states. At the operator level we relate the quantized mode expansion of bulk fields to a mode expansion of boundary operators, so bulk quanta are dual to CFT states created by modes of boundary operators acting on the vacuum. We develop this formalism in Sec. 2 and, more carefully, in Appendix A. This development continues the work in [4] which identified fluctuating supergravity modes as dual to boundary states, and non-fluctuating modes implementing boundary conditions as dual to boundary sources. Here, we will be interested mainly in the classical limit of the fluctuating states and in bulk configurations generated by brane sources.

Using the methods of Sec. 2, we discuss the boundary description of three kinds of probes: D-instantons, F- and D-strings, and dilaton wavepackets. In all cases the characteristic radial position of the bulk probe is mapped to the characteristic scale of the boundary configuration, as understood on general grounds by comparing the action of bulk isometries with conformal transformations on the boundary [1]. This scale-radius duality gives rise to pleasantly physical interpretations of bulk dynamics. For instance, strings or particles move in AdS spacetime to reduce gravitational potential energy; this is dual in the boundary theory to the spreading of localized field distributions to reduce gradient energy. In Euclidean $\text{AdS}_5 \times S^5$ it has been conjectured that a D-instanton at radial position z is dual to an instanton in $d = 4$ SYM with scale size z [5, 6, 7, 8]. As an application of our methods we *derive* this correspondence from the fundamental formulation of the AdS/CFT conjecture given in [2, 3]. We conclude by discussing the meaning of the classical limit on both sides

of the AdS/CFT duality, and by discussing the implications of our results for the holographic representation of black hole horizons in the boundary gauge theory.

2 Relating bulk and boundary states

In order to study how bulk geometry is encoded in the boundary description, we wish to introduce probes into the AdS spacetime. In this section we develop methods that identify boundary configurations corresponding to these probes. Our basic technique is to study the response of the bulk probe to a small change in boundary conditions. The formulation of the AdS/CFT correspondence in [2, 3] then provides the expectation values of operators in the corresponding boundary states. This will also allow us to arrive at a quantized formulation of the bulk theory from the boundary perspective. In this section we will ignore some subtleties [9] which are unimportant for most of the considerations of this paper. A more careful treatment is provided in Appendix A.

We work with Poincaré coordinates for anti-de Sitter space. The metric of AdS_{d+1} in these coordinates is:

$$ds^2 = \frac{R^2}{z^2}(-dt^2 + d\vec{x}_{d-1}^2 + dz^2) \quad (1)$$

(We set $R = 1$ in this section.) Poincaré coordinates only cover a patch of the global spacetime, and $z = 0$ is the boundary of AdS while $z = \infty$ is the horizon. (See [4] for more details and Penrose diagrams.) Euclidean AdS is obtained by taking t to it .

2.1 Euclidean signature

The AdS/CFT correspondence is formulated in [2, 3] as:

$$Z(\phi_i) = e^{-S(\phi_i)} = \langle e^{\int_B \phi_{0,i} \mathcal{O}^i} \rangle \quad (2)$$

where $S(\phi_i)$ is the effective action as a function of the bulk field ϕ_i , $\phi_{0,i}$ is the boundary value of ϕ_i (up to a scaling with the radial coordinate), and \mathcal{O}^i is the dual operator in the CFT. The expectation value on the right is evaluated in the CFT vacuum. We can read (2) as saying that boundary conditions for the bulk theory are dual to sources in the boundary theory. In other words, field theory in Euclidean AdS, expanded around a background approaching $\phi_{0,i}$ at the boundary, is described by a CFT deformed by the addition of a source. By functionally differentiating we find that:

$$\frac{\delta}{\delta \phi_{0,i}(\mathbf{x})} (-S(\phi_i)) = \langle \mathcal{O}^i(\mathbf{x}) \rangle_{\phi_{0,i}} \quad (3)$$

where the subscript $\phi_{0,i}$ indicates that the expectation value on the right hand side is computed in the presence of the source term $\int \phi_{0,i} \mathcal{O}^i$. We will use this relation

to learn about the expectation values of boundary operators in the presence of bulk probes.

Massive scalar: As an example, let us study a massive scalar field with a quadratic bulk action:

$$S(\phi) = \frac{1}{2} \int d^{d+1}x \sqrt{g} (|d\phi|^2 + m^2\phi^2), \quad (4)$$

Consider a solution to the bulk equations of motion that approaches $\phi_0(x)$ at the boundary (up to a scaling with the radial coordinate z). The unique solution regular in the bulk behaves as $\phi(z, \mathbf{x}) = z^{2h_-} \phi_0(\mathbf{x})$ as $z \rightarrow 0$ and can be written as:

$$\phi(z, \mathbf{x}) = c \int d^d \mathbf{x}' \frac{z^{2h_+}}{(z^2 + |\mathbf{x} - \mathbf{x}'|^2)^{2h_+}} \phi_0(\mathbf{x}'). \quad (5)$$

where we have used the bulk-boundary propagator [3] and

$$h_{\pm} = \frac{d}{4} \pm \frac{\sqrt{d^2 + 4m^2}}{4} \equiv \frac{d}{4} \pm \frac{\nu}{2} \quad (6)$$

The presence of this classical configuration corresponds to the addition of a source $\int \phi_0 \mathcal{O}$ to the boundary theory.

We now apply (3) by computing the functional derivative on the left hand side. To do so we perturb around $\phi(z, \mathbf{x})$ by a small fluctuation $\delta\phi$ and evaluate the resulting change in the action. After integrating (4) by parts and using the equations of motion, the variation becomes a surface term at the boundary:

$$\delta S(\phi) = \int_B d\Sigma^\mu \partial_\mu \phi \delta\phi. \quad (7)$$

Evaluating this quantity near the boundary at $z = 0$ is delicate [9]. In this section we will follow [3] by considering contributions to the integrand of (5) from the region $|\mathbf{x} - \mathbf{x}'| \neq 0$. This procedure amounts to ignoring certain contact terms and normalization issues as we discuss in Appendix A and gives, as $z \rightarrow 0$:

$$\frac{\partial\phi}{\partial z} = c(2h_+) z^{2h_+-1} \int d^d \mathbf{x}' \frac{\phi_0(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^{2(2h_+)}}. \quad (8)$$

The perturbation is written as $\delta\phi = z^{2h_-} \delta\phi_0$, and we use $d\Sigma^0 = z^{1-d} d^d \mathbf{x}'$. Then

$$\delta S(\phi) = c(2h_+) \int d^d \mathbf{x} d^d \mathbf{x}' \frac{\phi_0(\mathbf{x}') \delta\phi_0(\mathbf{x})}{|\mathbf{x} - \mathbf{x}'|^{2(2h_+)}}. \quad (9)$$

Using the relation (3) we derive

$$\langle \mathcal{O}(\mathbf{x}) \rangle_{\phi_0} = -c(2h_+) \int d^d \mathbf{x}' \frac{\phi_0(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^{2(2h_+)}}. \quad (10)$$

We have learned that in the presence of the source term $\int \phi_0 \mathcal{O}$ the operator \mathcal{O} has acquired an expectation value given by the right hand side of (10). This matches what we expect from the CFT by direct calculation:

$$\langle \mathcal{O}(\mathbf{x}) \rangle_{\phi_0} = \langle \mathcal{O}(\mathbf{x}) e^{\int \phi_0 \mathcal{O}} \rangle \approx \int d^d \mathbf{x}' \phi_0(\mathbf{x}') \langle \mathcal{O}(\mathbf{x}) \mathcal{O}(\mathbf{x}') \rangle \approx - \int d^d \mathbf{x}' \frac{\phi_0(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^{2(2h_+)}} \quad (11)$$

where the form of the two point correlator follows from scale invariance.

Interpretation: The AdS/CFT correspondence in (2) states that turning on a bulk mode which behaves as $z^{2h_-} \phi_0(\mathbf{x})$ near the boundary is dual to including a source term $\int \phi_0 \mathcal{O}$ in the CFT. As discussed in [4], the growth of such modes near the boundary indicates that they are non-fluctuating classical backgrounds. In effect, the presence of the mode redefines the Hamiltonian of the theory, since fluctuations should take place on top of this background. This is mirrored in the CFT by a modification of the action by the addition of a source term.

In the bulk, a mode with leading boundary behaviour $z^{2h_-} \phi_0$ induces a *subleading* component behaving as $z^{2h_+} \tilde{\phi}$. (This is seen by expanding (5) in powers of z .) The corresponding statement in the CFT is that the addition of the source induces an expectation value for \mathcal{O} . In fact, our analysis showed that $\langle \mathcal{O}(\mathbf{x}) \rangle_{\phi_0} \sim \tilde{\phi}(\mathbf{x})$, so that operator expectation values and bulk field components behaving as z^{2h_+} are precisely dual. This duality is the prevailing theme of the present work.

Bulk sources: In the above example, the solution $\phi(z, \mathbf{x})$ was completely determined by the boundary value ϕ_0 and the requirement of regularity in the bulk. However, as we shall see in Sec. 3, this uniqueness fails when we admit singular fields corresponding to sources in the bulk. Such bulk sources contribute subleading pieces to the fields at the boundary which modify (5) and contribute to operator expectation values. So once again we will find that subleading pieces of the bulk fields are dual to boundary expectation values. In this way, we will show in Sec. 3 that (3) implies that a bulk D-instanton in AdS₅ is dual to an instanton in the boundary Yang-Mills theory.

2.2 Lorentzian signature

The crucial new feature of Lorentzian signature is that the bulk wave equation admits propagating, normalizable mode solutions. Such modes describe the physical, low energy excitations of the spacetime; their explicit forms have been worked out in ([10, 11, 12, 4]) and they behave as z^{2h_+} near the boundary. These normalizable modes form the Hilbert space of the bulk theory. The possible boundary conditions for fields in AdS spacetime are encoded by the choice of non-normalizable mode solutions behaving as z^{2h_-} near the boundary. As argued in [4], the normalizable and

non-normalizable solutions are dual to states and sources respectively in the boundary conformal field theory. Here we make explicit the map between bulk and boundary states. So given a bulk field ϕ_i approaching $z^{2h_-}\phi_{0,i}$ at the boundary, we write the Lorentzian bulk-boundary correspondence as:

$$Z(\phi_i) = e^{iS(\phi_i)} = \langle s | e^{i \int_B \phi_{0,i} \mathcal{O}^i} | s \rangle \quad (12)$$

Here $|s\rangle$ represents the CFT state that is dual to the bulk state. Operator expectation values in excited CFT states will differ from their vacuum values; we write:

$$\frac{\delta}{\delta\phi_{0,i}(\mathbf{x})} S(\phi_i) = \langle s | \mathcal{O}^i(\mathbf{x}) | s \rangle_{\phi_{0,i}} \quad (13)$$

where the subscript indicates that the expectation is computed in the presence of a source term. We will also show that this can be turned into a statement relating quantized field operators in the bulk to operators in the boundary.

Massive scalar: As an example, we return again to the free massive scalar with action (4). Now a general nonsingular solution of the bulk equations approaching $z^{2h_-}\phi_0(\mathbf{x})$ at the boundary can be written as:

$$\phi(z, \mathbf{x}) = \phi_n(z, \mathbf{x}) + c \int d^d \mathbf{x}' \frac{z^{2h_+}}{(z^2 + |\mathbf{x} - \mathbf{x}'|^2)^{2h_+}} \phi_0(\mathbf{x}') \quad (14)$$

where ϕ_n is a normalizable mode, and $|\mathbf{x} - \mathbf{x}'|^2 = -(t - t')^2 + \sum_{i=1}^{d-1} (x_i - x'_i)^2$. Here we have used a bulk-boundary propagator obtained by continuation from Euclidean space in [3]. There are ambiguities in this choice whose meaning is discussed in Sec. 2.3 and Appendix A. We now repeat the procedure used in Euclidean signature, dropping contact terms as before, and paying attention to the extra contribution from ϕ_n . (There are subtleties in this procedure - see Appendix A.) Using $\phi_n(z, \mathbf{x}) \rightarrow z^{2h_+}\tilde{\phi}_n(\mathbf{x})$ as $z \rightarrow 0$, we find:

$$\langle \tilde{\phi}_n | \mathcal{O}(\mathbf{x}) | \tilde{\phi}_n \rangle_{\phi_0} = (2h_+)\tilde{\phi}_n(\mathbf{x}) + c(2h_+) \int d^d \mathbf{x}' \frac{\phi_0(\mathbf{x}')}{|\mathbf{x} - \mathbf{x}'|^{2(2h_+)}} \quad (15)$$

where we have indicated that the CFT is in the excited state $|\tilde{\phi}_n\rangle$.¹ So \mathcal{O} gets an expectation value from two distinct contributions: from the excited state and from the source that has been turned on. Note that $|\tilde{\phi}_n\rangle$ is a ‘‘coherent’’ state on the boundary in which operators have non-vanishing expectation values.

¹The normalizations produced by this naive Lorentzian treatment are not correct - see Appendix A.

Interpretation: In Lorentzian AdS, normalizable and non-normalizable modes are dual to states and sources respectively. As we have seen, operator expectation values are affected by both the state and the source. Nevertheless, as in the Euclidean case, the component of the total bulk field behaving as z^{2h-} defines the source while the component behaving as z^{2h+} gives rise to the boundary expectation value. We are free to set ϕ_0 to zero if we wish, so that the sources are turned off – then we are studying states of the original unmodified CFT.

Bulk sources: In the example above, we considered linearized wave equations for a massive scalar. It is possible to consider fully non-linear solutions with possible bulk singularities due to sources. For instance, the field configuration arising from a D-brane in the AdS geometry would be of this type. The treatment based on (13) is equally valid in this case – linearized fluctuations around the fully non-linear solution lead by integration by parts to the same surface integrals. When ϕ_0 vanishes (15) will also continue to hold, but in general nonvanishing ϕ_0 will push the field configuration into the non-linear regime and bulk interactions will become important. Equipped with (15), in the next sections we will determine the dual boundary description of various solitonic objects in the bulk.

2.3 Operator formulation

We would also like a more microscopic mapping at the level of individual quantum states. This is obtained by regarding fields as quantized operators. In particular, we may write ϕ_n in terms of a mode expansion:

$$\hat{\phi}_n = \sum_k \left[a_k \phi_{n,k} + a_k^\dagger \phi_{n,k}^* \right] \quad (16)$$

and similarly for the boundary operator \mathcal{O} ,

$$\hat{\mathcal{O}} = \sum_k \left[b_k \tilde{\phi}_{n,k} + b_k^\dagger \tilde{\phi}_{n,k}^* \right], \quad (17)$$

where as before, $\tilde{\phi}(\mathbf{x})$ is the boundary value of the component of the bulk mode $\phi(z, \mathbf{x})$ scaling as z^{2h+} as $z \rightarrow 0$. Note that the modes appearing in the expansion of $\hat{\phi}_n$ satisfy an on-shell condition $(\square - m^2)\phi_{n,k} = 0$, whereas the modes $\tilde{\phi}_{n,k}$ do not satisfy any wave equation on the boundary, but are instead a fully complete set of functions. Interpreting (15) as an operator statement we conclude that $a_k = b_k$, and that $a_k^\dagger |0\rangle = |k\rangle$ where $|k\rangle$ is a “one particle” state created by a single application of b_k^\dagger . (It is intriguing that creation and annihilation operators of elementary bulk fields are related to composite operators on the boundary, leading us to identify $\sum_k b_k^\dagger b_k$ as a particle number operator.) In other words, bulk states described by quanta occupying normalizable modes are dual to CFT states described by acting on the vacuum with

modes of the appropriate boundary operator. This provides a direct correspondence between bulk and CFT states.²

Choice of vacuum: Quantum field theory in curved space can usually accommodate a variety of inequivalent vacua, corresponding to different definitions of positive frequency. In the present context, the choice of vacuum affects the formalism in two places. First, mode solutions in (16,17) have positive frequency with respect to a particular time coordinate, here taken to be Poincaré time. A different choice of time can lead to an inequivalent vacuum state, related to the original vacuum by a Bogoliubov transformation. Second, we have made a particular choice for the form of the bulk-boundary propagator in (14) which we obtained by continuation from Euclidean space. This is the appropriate propagator to use when perturbing around the Poincaré vacuum. However, alternative vacua can be chosen by changing the mode expansions (16,17) and modifying the bulk-boundary propagator. The latter modification will involve adding a normalizable mode to the original propagator, which leaves unchanged its z^{2h_-} dependence near $z = 0$. These issues are discussed in greater detail in Appendix A

2.4 Radial isometry and boundary scale transformation

In subsequent sections we will use the formalism developed above to study the boundary representation of probes in the bulk. A recurring theme will be a duality between characteristic radial positions in the bulk and characteristic scales in the boundary. Let us review how this arises. The Poincaré metric (1) has a radial isometry:

$$(\vec{x}, t) = \mathbf{x} \rightarrow \lambda \mathbf{x} \quad ; \quad z \rightarrow \lambda z \quad (18)$$

As we have discussed, boundary expectation values for a massive scalar are dual to the component of the bulk field scaling as $z^{2h_+} \tilde{\phi}(\mathbf{x})$ near the boundary. So consider a one parameter family of bulk solutions of the form:

$$\phi^\lambda(z, \mathbf{x}) = \phi(\lambda z, \lambda \mathbf{x}) \quad (19)$$

According to (15), the expectation value of the boundary operator in the corresponding state behaves as:

$$\langle \phi^\lambda | \mathcal{O}(\mathbf{x}) | \phi^\lambda \rangle = \lambda^{2h_+} \langle \mathcal{O}(\lambda \mathbf{x}) \rangle \quad (20)$$

So the radial isometry generates scale transformations on the boundary. In the examples studied below we will explicitly see that the boundary configuration spreads out as the bulk probe falls towards the Poincaré horizon.

²Related issues are discussed in [13].

3 Instanton probes

Our first example of a bulk probe is a D-instanton in $\text{AdS}_5 \times S^5$. The methods developed in the previous section will show that it is dual to a boundary Yang-Mills instanton. A particular consequence is a duality between the radial position of the bulk object and the scale size on the boundary. D-instanton solutions in AdS space have been discussed in [6, 7, 8] where the close similarity between the bulk dilaton profile and the boundary instanton was noted. Our main point in this section is that this fact follows from the general considerations of Sec. 2. This *derives* the duality between the D-instanton and the Yang-Mills instanton from the fundamental formulation of the AdS/CFT correspondence in (2).

According to the methods of Sec. 2, to determine the boundary expectation values corresponding to a D-instanton we require the form of the bulk solution near the boundary. We will use Poincaré coordinates for AdS_5 , so that the metric is given by (1) with $R^4 = 4\pi g_s N \alpha'^2$. D-instanton solutions have been presented in [6, 7, 8]; only the dilaton (ϕ) and the axion (χ) are turned on, while the $\text{AdS}_5 \times S^5$ Einstein metric is unchanged. Dimensional reduction of these fields on S^5 produces a Kaluza-Klein tower of modes on AdS_5 . Here we are only interested in the behaviour as $z \rightarrow 0$ of the massless five dimensional dilaton and axion that couple to $\text{Tr}(F^2)$ and $\text{Tr}(F\tilde{F})$ on the boundary.³ This is given by:⁴

$$e^\phi = g_s + c \frac{z^4 \tilde{z}^4}{[\tilde{z}^2 + |\vec{x} - \vec{x}_a|^2]^4} \dots, \quad (21)$$

$$\chi = \chi_\infty \pm (e^{-\phi} - 1/g_s) \quad (22)$$

Here \tilde{z} is the radial position of the D-instanton. The constant c in (21) can be determined by requiring that the D instanton carries the correct axionic charge. That is:

$$\frac{1}{2\kappa_{10}^2} \int e^{2\phi} \partial_\mu \chi dS^\mu = 2\pi, \quad (23)$$

yielding

$$c = \frac{24\pi}{N^2}. \quad (24)$$

³These couplings were first studied in [14].

⁴There is some disagreement between [6, 7] and [8] concerning whether the the D-instanton should be localized on the S^5 . The difference will lie in the excitation of the Kaluza-Klein harmonics that are massive fields on AdS_5 . Including these excitations which fall off faster at the boundary will give expectation values to dual higher dimension operators that were identified in [3]. The authors of [6, 7, 8] all agree on the asymptotic form of the massless AdS_5 fields which is all that we require here.

Boundary expectations: We can now use the methods of Sec. 2 to derive the expectation values of boundary operators. The five-dimensional dilaton action is:

$$S_\phi = -\frac{1}{4\kappa_5^2} \int d^5x \sqrt{g} g^{\mu\nu} \partial_\mu \phi \partial_\nu \phi + \dots \quad (25)$$

where $1/\kappa_5^2 = V_5 R^5 / \kappa_{10}^2$, κ_{10} is the 10 dimensional Newton constant and $V_5 = \pi^3$ is the volume of the unit 5-sphere. We start with the dilaton background for the D-instanton (21) and add a small perturbation $\delta\phi$. The resulting change in the action is a boundary term:

$$\delta S = -\frac{1}{2\kappa_5^2} \int d^4x \frac{R^3}{z^3} \delta\phi \partial_z \phi \quad (26)$$

It follows that the functional derivative with respect to the boundary configuration ϕ_0 is: $(\delta S / \delta \phi_0) = -(1/2\kappa_5^2) R^3 \partial_z \phi$. Using ϕ in (21) and the relation $R^4 = \kappa_{10} N / 2\pi^{5/2}$ gives:

$$\frac{\delta S}{\delta \phi_0(\vec{x})} = -\frac{48}{4\pi g_s} \frac{\tilde{z}^4}{[\tilde{z}^2 + |\vec{x} - \vec{x}_a|^2]^4} . \quad (27)$$

We learn from [3, 2] that the AdS₅ massless dilaton couples to $Tr(F^2)$ in the boundary CFT. Choosing the normalization $S_{YM} = (1/4g_{YM}^2) \int d^4x Tr(F^2) + \dots$ for the Yang-Mills action gives:

$$\frac{\delta S_{YM}}{\delta \phi_0(\vec{x})} = -\frac{1}{4g_{YM}^2} \langle Tr(F^2(\vec{x})) \rangle . \quad (28)$$

Equating (27) and (28) and using $4\pi g_s = g_{YM}^2$, we find:

$$\frac{1}{4g_{YM}^2} \langle Tr F^2(\vec{x}) \rangle = \frac{48}{g_{YM}^2} \frac{\tilde{z}^4}{[\tilde{z}^2 + |\vec{x} - \vec{x}_a|^2]^4} , \quad (29)$$

which is exactly the Yang-Mills field strength in an instanton background. So, as advertised, a D-instanton in the bulk is precisely dual to a Yang-Mills instanton in the boundary theory.

This discussion of the dilaton can be extended to the axion which yields the expected $\langle F\tilde{F} \rangle$ for an instanton. Since the AdS₅ metric is unchanged, we learn that the expectation value of the stress tensor $\langle T_{\mu\nu} \rangle$ vanishes in an instanton background. This is easily checked; the stress tensor is:

$$T_{\mu\nu} = \frac{1}{4} g_{\mu\nu} Tr F^{\rho\sigma} F_{\rho\sigma} - Tr F_\mu^\rho F_{\rho\nu} . \quad (30)$$

The two terms cancel for the Yang-Mills instanton. Similarly, the NS-NS B-field in an S -wave on S^5 is known to be dual to a dimension 6 operator in the Yang-Mills theory [15]. This operator was derived in [16, 17]:

$$\mathcal{O}_{\mu\nu}^{(6)} \propto \text{tr} \left[\frac{1}{2} F_{[\nu\alpha} F^{\alpha\beta} F_{\beta\mu]} + \frac{1}{8} F_{\alpha\beta} F^{\alpha\beta} F_{\mu\nu} \right] , \quad (31)$$

where we have antisymmetrized the indices μ, ν . Again, it is simple to check that the first term vanishes upon antisymmetrization and the second term vanishes identically.⁵ Finally, the D-instanton action $S_{Dinst} = 2\pi/g_s$ coincides with the YM instanton action $S_{YM} = 8\pi^2/g_{YM}^2$ using $4\pi g_s = g_{YM}^2$.

Scale-radius duality and bulk locality: The above duality between the D-instanton and YM instanton provides the first example of a phenomenon we will call scale-radius duality. The characteristic radial position of the D-instanton is \tilde{z} . On the boundary \tilde{z} is the characteristic scale of the YM instanton. A D-instanton closer to the horizon at $\tilde{z} = \infty$ is mapped into a fatter boundary object [5]. More specifically, the action of the isometries (18) translates the D-instanton both radially and parallel to the boundary. The corresponding conformal transformation of the boundary instanton rescales it and translates it at the same time.

This relation has interesting consequences for the emergence of local physics in the bulk when we apply our techniques to multi-instanton solutions. Consider two D-instantons at very different radial positions z , but at the same coordinate \vec{x} parallel to the boundary. These are dual to two coincident YM instantons with widely different scale sizes. Locality of the bulk objects is expected at large N when classical physics is valid. In this limit the collective coordinates of the boundary configuration should approximately decouple into two separate sets associated with instantons of two different scale sizes (*i.e.* the metric on the moduli space is block-diagonal in this region). Such behaviour typically occurs for instantons at large spatial distances. Here we learn that a large difference in scale size will also cause a separation of collective coordinates. Turning this around, the approximate non-interaction of collective coordinates of coincident instantons at widely different scales translates into approximate locality of the bulk physics.

We can also consider the interaction between instantons and anti-instantons. By evaluating the probe action of a D-instanton in the background of the anti-D-instanton, we find a bulk interaction of the form:

$$\delta S \sim \frac{(z_1 z_2)^4}{(z_1 - z_2)^8}, \quad (32)$$

where the D-instantons are at coincident \vec{x} positions and $z_{1,2}$ are their radial positions. In deriving (32) we have used the asymptotic form of the dilaton, which follows from the required fall-off of the dilaton and $SO(1,5)$ invariance (it is also consistent with [6, 7, 8]). From (32) it follows that in the conformal large N Yang-Mills theory we expect the interaction between coincident instantons and anti-instantons to fall

⁵Note, however, that the symmetrized part of the first term does not vanish. In general we expect that there will be all sorts of combinations of F that will not vanish in this background. This is not a surprise – for example, in the presence of a classical bulk configuration, we expect that many operators corresponding to multi-particle bulk states will have expectation values.

off as the eighth power of difference in scale size. Such behavior is not evident in perturbative gauge theory, and presumably arises from the sum of planar diagrams in the large N limit.

4 String probes

Fundamental and D-strings are particularly interesting probes of Lorentzian AdS_5 . Working in Poincaré coordinates (1), we will find the Yang-Mills description of string solitons stretched parallel to the AdS boundary, and placed at fixed radial positions \tilde{z} . Once again, the characteristic position in the bulk will be mapped to the characteristic scale on the boundary, and motion towards the Poincaré horizon appears as a fattening of the boundary flux tube. A string placed at a fixed radial position is not a solution to the equations of motion since it can reduce its potential energy by falling towards the horizon. Nevertheless, since the analysis of a static string is technically clearer, we imagine that it is stabilized by an external force. We will find that the corresponding boundary flux could reduce its energy by spreading and must be similarly stabilized.⁶ This analysis is readily generalized to slowly moving strings.

Strings as fluxes: First, we establish that fundamental (F) and Dirichlet (D) strings in AdS_5 are described by electric and magnetic fluxes in the boundary gauge theory. This is easily shown by starting with the worldvolume action for D3-branes (with Higgs fields suppressed):

$$S_{D3} = -T_3 \text{Tr} \int d^4 \sigma \left\{ e^{-\phi} \sqrt{-\det(G_{mn} + 2\pi\alpha' F_{mn} + B_{mn})} - \frac{1}{4} \epsilon^{mnpq} C_{mn}^{(2)} F_{pq} \right\} \quad (33)$$

(We have written the nonabelian Born-Infeld action appropriate to commuting background fields [18].) The action includes terms of the form:

$$\int d^4 \sigma B^{mn} \text{Tr} F_{mn} \quad \text{and} \quad \int d^4 \sigma \epsilon^{mnpq} C_{mn}^{(2)} \text{Tr} F_{mn}, \quad (34)$$

where B_{mn} and $C_{mn}^{(2)}$ are the NS-NS and RR 2-forms respectively. An F-string extended in the x^i direction should couple to B_{ti} , and so is described by a non-vanishing value of $E_i = \text{Tr} F_{ti}$, in other words an electric flux.⁷ Similarly, the D-string corresponds to a magnetic flux $B_i = \epsilon_{ijk} \text{Tr} F_{jk}$. This is consistent with S-duality, which

⁶Related discussions appear in [13].

⁷The source for the $U(1)$ electric flux in the Yang-Mills theory is not associated with the dynamical part of B_{mn} , which couples to a dimension 6 operator in the CFT [15, 16], but rather with pure gauge degrees of freedom which contribute to surface integrals for conserved charges at infinity. It may seem surprising that (34) shows that bulk strings are related to the $U(1)$ part of the $U(N)$ boundary theory since this factor has been argued to decouple [19, 3, 20]. However, it is more accurate to consider the $U(1)$ as “frozen” after inclusion of all external probes and VEVs. The remaining $SU(N)$ part of the theory is dynamical and the rest of this section studies the expectation of $\text{Tr}(F^2)$ with a trace in $SU(N)$. We are grateful to O. Aharony for correspondence regarding this issue.

interchanges F and D strings in the bulk, and electric and magnetic fields in the Yang-Mills theory.

Bulk strings: To describe bulk F and D strings we start with their worldsheet actions, which include the terms

$$S_F = -\frac{1}{2\pi\alpha'} \int d^2\sigma \left\{ \sqrt{-\det g_{\mu\nu} \partial_m X^\mu \partial_n X^\nu} - \frac{1}{2} \epsilon^{mn} B_{\mu\nu} \partial_m X^\mu \partial_n X^\nu \right\} \quad (35)$$

$$S_D = -\frac{1}{2\pi\alpha' g_s} \int d^2\sigma \left\{ e^{-\phi} \sqrt{-\det g_{\mu\nu} \partial_m X^\mu \partial_n X^\nu} - \frac{1}{2} \epsilon^{mn} C_{\mu\nu}^{(2)} \partial_m X^\mu \partial_n X^\nu \right\} \quad (36)$$

We have included couplings due to the string frame metric $g_{\mu\nu}$, the dilaton ϕ , and the 2-forms $B_{\mu\nu}$, $C_{\mu\nu}^{(2)}$. Now consider static strings extended in the x^1 direction, parallel to the boundary. In static gauge,

$$t(\sigma^m) = \sigma^0 \quad z(\sigma^m) = \tilde{z} = \text{constant} \quad x^1(\sigma^m) = \sigma^1 \quad \vec{x}_\perp(\sigma^m) = \vec{x}_{a\perp}. \quad (37)$$

Here $\vec{x}_{a\perp}$ are directions orthogonal to the string but parallel to the boundary. Evaluating S_F , S_D in the AdS₅ background (1) gives the potential energy of the static strings (per unit coordinate length):

$$V_F = \frac{g_{YM}}{2\pi} N^{1/2} \frac{1}{\tilde{z}^2} \quad V_D = \frac{2}{g_{YM}} N^{1/2} \frac{1}{\tilde{z}^2}, \quad (38)$$

Here we used the relations $R^4 = 4\pi g_s N \alpha'^2$ and $4\pi g_s = g_{YM}^2$ that are appropriate to AdS₅. The F and D string potentials are related by S-duality: $g_{YM}^2/4\pi \rightarrow (g_{YM}^2/4\pi)^{-1}$.

Boundary expectations: A bulk string is dual to a boundary CFT state in which various operators have expectation values. For example, the actions (35,36) will induce long range fields for B_{tx^1} , $C_{tx^1}^{(2)}$ in the presence of F and D strings respectively. The results of Sec. 2 and the couplings (34) then yield nonvanishing values for $\langle \text{Tr} F_{tx^1} \rangle_F$ and $\langle \text{Tr} F_{x^2 x^3} \rangle_D$, corresponding to electric and magnetic fluxes. Rather than evaluating these explicitly, we focus on the expectation value for $\text{Tr} F^2$, which couples to the dilaton ϕ . Note that this trace is in the $SU(N)$ part of the gauge group. First we work out the long range dilaton field produced by string sources via their linear coupling to the dilaton. Although it may appear from (35) that the F-string does not couple to the dilaton, this is simply because the actions are written in the string frame. Working instead in the Einstein frame, with $g_{\mu\nu}^E = (g_s e^\phi)^{1/2} g_{\mu\nu}$, we find the couplings:

$$S_F = -\frac{g_{YM}}{4\pi} N^{1/2} \int d^2\sigma \frac{\phi}{\tilde{z}^2} \quad S_D = \frac{1}{g_{YM}} N^{1/2} \int d^2\sigma \frac{\phi}{\tilde{z}^2}. \quad (39)$$

In order to obtain the long range dilaton field we will need the asymptotic form of a Green's function, satisfying Dirichlet boundary conditions, for the equation:

$$\frac{1}{2\kappa_5^2} \nabla^2 G_D(\vec{x}_\perp, z) = \frac{z^5}{R^5} \delta(z - \tilde{z}) \delta(\vec{x}_\perp - \vec{x}_{a\perp}), \quad (40)$$

where the right hand side contains a delta function transverse to the direction in which the string extends. This is readily solved to give:

$$G_D(\vec{x}, z) = -\frac{4\pi}{N^2} \frac{z^4 \tilde{z}^4}{[\tilde{z}^2 + |\vec{x}_\perp - \vec{x}_{a\perp}|^2]^3}. \quad (41)$$

(This can also be obtained from our earlier results for the D-instanton by integrating (21) with respect to the positions t_a and x_a^1 and multiplying by the appropriate normalization.) Now it follows from (39) and (41) that

$$\phi_F = \frac{g_{YM}}{N^{3/2}} \frac{z^4 \tilde{z}^2}{[\tilde{z}^2 + |\vec{x}_\perp - \vec{x}_{a\perp}|^2]^3} \quad (42)$$

and

$$\phi_D = -\frac{4\pi}{g_{YM} N^{3/2}} \frac{z^4 \tilde{z}^2}{[\tilde{z}^2 + |\vec{x}_\perp - \vec{x}_{a\perp}|^2]^3}. \quad (43)$$

These asymptotic fields yield the expectation values

$$\langle \text{Tr} F^2 \rangle_F = \frac{2g_{YM}^3 N^{1/2}}{\pi^2} \frac{\tilde{z}^2}{[\tilde{z}^2 + |\vec{x}_\perp - \vec{x}_{a\perp}|^2]^3} \quad (44)$$

$$\langle \text{Tr} F^2 \rangle_D = -\frac{8g_{YM} N^{1/2}}{\pi} \frac{\tilde{z}^2}{[\tilde{z}^2 + |\vec{x}_\perp - \vec{x}_{a\perp}|^2]^3} \quad (45)$$

We see that the boundary configuration corresponding to a static string in the bulk is spread over a region with scale size \tilde{z} . This analysis can be generalized to a slowly moving string by using retarded Green's functions instead of (41); \tilde{z} is then replaced by its value at retarded time.

Scale-radius duality and bulk locality: We have learned that a bulk string placed at \tilde{z} is dual to a flux tube spread over a region with characteristic scale \tilde{z} – another example of scale-radius duality. In the bulk a string will fall towards the horizon (large \tilde{z}) to minimise gravitational potential energy. Correspondingly, the gauge field strength in the boundary theory will spread out to minimize gradient energy, asymptotically going to zero. The AdS/CFT correspondence implies that the equation governing the spreading in the boundary is the geodesic equation for strings in the bulk. At present it is difficult to analyze this directly from the boundary perspective, but we gain some insights from [21],⁸ where a system of $p+2$ and p branes in flat space is studied. The authors found that a D-string can be included in a D3-brane $SU(N)$ gauge theory as a Z_N flux after compactifying two directions transverse to the D-string. Their analysis showed that the minimum energy configuration with fixed Z_N flux is pure gauge, with vanishing field strength. Taking a large compactification radius, this agrees with the picture in (44-45) where $\langle \text{Tr}(F^2) \rangle$

⁸We thank A. Sen for bringing this reference to our notice.

vanishes as $\tilde{z} \rightarrow \infty$. For a classical, static field configuration which is purely electric or magnetic, $(1/g_{YM}^2)\text{Tr}F^2$ is proportional to the energy density.⁹ One expects that the total energy for a field configuration of fixed flux and size \tilde{z} goes like $1/\tilde{z}^2$, in agreement with the bulk potentials (38).

Two strings at large radial separations do not interact very much. This feature can be seen by examining the collective coordinates of the bulk solitons – each string has an approximately independent set. On the boundary the corresponding statement is that flux tubes of very different scale sizes have approximately independent collective fluctuations, even when they have coincident centers.¹⁰ The interactions of bulk strings are also causal in the classical limit. For instance, a fluctuation on one string will only affect the other after a time lag. This translates into a typical interval required for the spread of boundary fluctuations from one scale to another. As in the case of D-instantons, the separation of boundary collective fluctuations and the time lag for interactions between scales are only expected to emerge in an approximate sense. At a more fundamental level, the exact dynamics dictated by the CFT description will imply deviations from bulk locality and causality.

5 Dilaton Wave Packet

Finally we study massless dilaton wavepackets in the bulk of AdS_5 . In previous sections we studied pointlike sources and found that the bulk position translated into a boundary scale. The situation is more complicated for wavepackets because the bulk object already has a characteristic scale which will also get reflected on the boundary. The correct approach is to study a family of objects related in the bulk by the AdS isometry (18). For the D-instanton and string probes, this isometry simply translates the bulk objects. The component of the translation parallel to the boundary becomes a translation in the CFT, while the radial translation becomes a *spatial* rescaling. We will see that the isometry (18) both translates and changes the size of dilaton wavepackets. This is reflected in the boundary theory in the *spatio-temporal* width of $\langle \text{Tr}(F^2) \rangle$.

⁹Since we do not expect classical Yang-Mills theory to be accurate in this context, the following discussion is meant only to indicate the qualitative behavior of the field configuration.

¹⁰Of course, an object like a flux tube, not being a stable soliton, does not have collective coordinates in the usual sense. By “collective coordinates” we mean fluctuations of the boundary fields that preserve the overall scale and shape of the tube. For example, we can imagine endowing the tube with a ripple in its shape or a transverse velocity. Such motions would occur on time scales distinct from the rate of spreading of the tube.

Bulk wavepackets: A normalizable dilaton wavepacket can be constructed by superimposing mode solutions [10, 11, 12, 4]:

$$\phi_\lambda(z, \vec{x}, t) = \int d^3k d\rho C_\lambda(\vec{k}, \rho) (\rho z)^2 J_2(\rho z) e^{i(\vec{k} \cdot \vec{x} - \omega t)}. \quad (46)$$

Here J_2 is a Bessel function and $\omega^2 - \rho^2 - \vec{k}^2 = 0$. The profile C_λ is a Gaussian centered at $(\lambda \vec{k}_0, \lambda \rho_0)$ with a width $\sigma \lambda^2$:

$$C_\lambda(\vec{k}, \rho) = \frac{1}{\lambda^4} a e^{[-\frac{(\vec{k} - \lambda \vec{k}_0)^2 + (\rho - \lambda \rho_0)^2}{2\lambda^2 \sigma} - i\frac{\rho \alpha}{\lambda}]}, \quad (47)$$

With this definition, wavepackets with different values of λ are related by the radial isometry (18):

$$\phi_\lambda(z, \vec{x}, t) = \phi_{(\lambda=1)}(\lambda z, \lambda \vec{x}, \lambda t). \quad (48)$$

So we expect to see a manifestation of the scale-radius duality in the dual boundary dynamics.

For sufficiently early or late times¹¹ we can use the stationary phase approximation and the asymptotic form of the Bessel function for large arguments to study (46). We find, self-consistently, that at large $|t|$, the packet ϕ_λ is centered at:

$$\vec{x} = \frac{\vec{k}_0}{\omega_0} t \quad ; \quad z = \frac{\rho_0}{\omega_0} |t| + \frac{\alpha}{\lambda}, \quad (49)$$

As λ increases, the center of the wave packet moves radially away from the horizon. At large times the detailed form of the wavepacket is still somewhat complicated but the key features can be understood by setting $\alpha = 0$. This gives:

$$\phi_\lambda(z, x_i, t) = a \sqrt{\sigma} (\rho_0 z)^{3/2} \left(\frac{\omega_0}{t}\right)^{3/2} e^{-\frac{\sigma \lambda^2}{2} (\vec{x} \cdot \hat{\zeta}_0 - t)^2} e^{-\frac{1}{2} \frac{x_T^2 \omega_0^2}{t^2 \sigma}} e^{\frac{i}{2} \frac{x_T^2 \lambda \omega_0}{t}} e^{i(\lambda \vec{\zeta}_0 \cdot \vec{x} - \lambda \omega_0 t)}. \quad (50)$$

Here \vec{x} is the 4-vector (z, x_i) , $\vec{\zeta}_0$ is the 4-vector, (ρ_0, \vec{k}_0) and x_T stands for the spatial distance in 4 dimensions transverse to the wave packet's momentum along $\vec{\zeta}_0$. So, as $t \rightarrow -\infty$ the wave packet is a shock wave that emerges from the horizon travelling along $\hat{\zeta}$. Initially its energy is large and it travels like a massless particle along the light cone. But with time the pull of gravity gets stronger and begins to reflect the wavepacket back. As this happens the shock wave contracts into a localized lump in the direction perpendicular to its motion. At this stage (50) is no longer valid. Eventually the state turns around completely, gathering itself into a shock wave again, this time hurtling towards the horizon in the far future and spreading out transverse to its direction of motion. This spreading transverse to the direction of motion will be familiar to reader as the standard behaviour of relativistic wavepackets in flat space.

In (50) we had set $\alpha = 0$; reinstating it does not change the qualitative features of the wave packet. In particular the widths $\sigma \lambda^2$ and $\frac{\omega_0^2}{t^2 \sigma}$, which govern the spreading

¹¹We need $|\lambda t| \gg \omega_0 / \rho_0^2$ and $|\lambda t| \gg \omega_0 / \sigma$.

parallel and perpendicular to the direction of motion stay the same. Thus, shock waves closer to the horizon (smaller λ from (49)) are also more spread out along their direction of motion.

Boundary Description: Using the results of Sec. 2 we relate the asymptotic behaviour of the dilaton wave packet to the expectation value of $\langle Tr F^2 \rangle$:

$$F^2(\vec{x}) \equiv \frac{1}{4g_{YM}^2} \langle \lambda | Tr(F^2(\vec{x})) | \lambda \rangle = \int d^3k d\rho 4C_\lambda(\vec{k}, \rho) \rho^4 e^{i(\vec{k} \cdot \vec{x} - \omega t)}, \quad (51)$$

Here the states $|\lambda\rangle$ are related to each other by a scale transformation dual to the isometry (18). Two limiting cases are instructive: $\rho_0 \gg |\vec{k}_0|$ and $\rho_0 \ll |\vec{k}_0|$.

When $\rho_0 \gg |\vec{k}_0|$ the wavepacket starts at early times with most of its momentum in the radial direction. Then (51) gives:

$$|\lambda t| \gg \frac{\omega_0}{\sigma} : \quad F^2(\vec{x}) \propto \left(\frac{\omega_0}{t}\right)^{\frac{3}{2}} e^{-\frac{\sigma\lambda^2}{2}(t+\frac{\alpha}{\lambda})^2} e^{-\frac{1}{2}\frac{\vec{x}^2\omega_0^2}{t^2\sigma}} e^{\frac{i}{2}\frac{\vec{x}^2\omega_0\lambda}{t}} e^{-i\omega_0(\lambda t + \alpha)} \quad (52)$$

$$|\lambda t| \ll \frac{\omega_0}{\sigma} : \quad F^2(\vec{x}) \propto e^{-\frac{\sigma\lambda^2}{2}(t+\frac{\alpha}{\lambda})^2} e^{-\frac{1}{2}(\vec{x}^2\sigma\lambda^2)} e^{-i\omega_0(\lambda t + \alpha)} \quad (53)$$

At early and late times ($|\lambda t| \gg \frac{\omega_0}{\sigma}$) the bulk state looks like a shock wave moving in the radial direction and at small intermediate times ($|\lambda t| \ll \frac{\omega_0}{\sigma}$) the bulk state is being reflected by the AdS geometry and turned around.

When $\rho_0 \ll |\vec{k}_0|$ the wavepacket starts at early times with most of its momentum parallel to the boundary. Then we find:

$$|\lambda t| \gg \frac{\omega_0}{\sigma} : \quad F^2(\vec{x}) \propto \left(\frac{\omega_0}{t}\right)^{\frac{3}{2}} e^{-\frac{\sigma\lambda^2}{2}(\vec{x} \cdot \hat{k}_0 - t)^2} e^{-\frac{1}{2}\frac{x_T^2\omega_0^2}{t^2\sigma}} e^{\frac{i}{2}\frac{x_T^2\omega_0\lambda}{t}} e^{i\lambda(\vec{k}_0 \cdot \vec{x} - \omega_0 t)} \quad (54)$$

$$|\lambda t| \ll \frac{\omega_0}{\sigma} : \quad F^2(\vec{x}) \propto e^{-\frac{\sigma\lambda^2}{2}(\vec{x} \cdot \hat{k}_0 - t)^2} e^{-\frac{1}{2}(x_T^2\lambda^2\sigma)} e^{i\lambda(\vec{k}_0 \cdot \vec{x} - \omega_0 t)}, \quad (55)$$

Here $x_T^2 = \vec{x}^2 - (\vec{x} \cdot \hat{k}_0)^2 + (\frac{\alpha}{\lambda})^2$. At early and late times ($|\lambda t| \gg \frac{\omega_0}{\sigma}$) the bulk state looks like a shock wave moving parallel to the boundary.

Scale-radius duality: We have just derived the boundary description of a class of wavepackets ϕ_λ that are related by the radial isometry (18). We learned from (49) that the characteristic spatial center of the packet depends linearly on $1/\lambda$. Packets that are closer to the horizon also have bigger bulk widths. The basic lesson we learn from (52) - (55) is that these characteristic bulk features map on the boundary to a characteristic *spatio-temporal* scale. Packets which are characteristically closer to the horizon (and spatially wider) map to boundary fields with a greater spread in space and time.

To see this, first consider the case $\rho_0 \gg |\vec{k}_0|$. The bulk packet starts as a radial shockwave coming from the past horizon, reflects at intermediate times and returns

as a shockwave to the future horizon. From (52) the boundary F^2 starts with a very small amplitude and a huge spread in the spatial directions. At intermediate times (53) tells us that F^2 is a Gaussian in space and at late times the amplitude decreases again while F^2 spreads in space. The temporal profile, like the intermediate time spatial profile, is Gaussian with a scale set by λ . Since λ also indexes the radial isometries relating bulk packets, we are once again seeing a scale-radius duality.

As another example, consider the case $\rho_0 \ll |\vec{k}_0|$. Now the bulk packet starts as a shockwave parallel to the boundary, is reflected at intermediate times and returns as a shockwave to the future horizon. From (54) and (55) we learn that the boundary F^2 is a shockwave spreading out spatially in the directions perpendicular to the motion at early and late times. Other than this, the behaviour is exactly parallel to the case $\rho_0 \gg |\vec{k}_0|$. The profile in time, like the intermediate time spatial profile, is a Gaussian with a scale set by λ .

The analysis of this section is only a first step in a more complete study. For example, it would be interesting to understand how the bulk scattering of two shock waves is mirrored in the boundary theory. Understanding this would help uncover how bulk locality emerges from the boundary description.

6 Discussion and conclusions

6.1 Bulk versus boundary dynamics

The classical limit: Much of this article has addressed the CFT description of classical bulk probes. These states are “classical” because they contain a very large number of particles, and propagate in backgrounds with small curvature. There is a subtlety in defining such a limit in view of the “stringy exclusion principle” advocated in [23] for AdS₃. Given the AdS/CFT correspondence, this “exclusion principle” imposes a bound on the occupancy of certain bulk states. Nevertheless, as N increases, the maximum occupancy increases also. So, in the large N limit our considerations are valid.

Hairy holography: We have provided an explicit prescription for relating states in the bulk and boundary theories. Roughly speaking, the prescription works because normalizable modes in the bulk extend to the boundary and serve as a kind of “hair”, determining the boundary state. Each mode falls off exponentially (in physical distance), but the volume of the boundary grows exponentially as well allowing for a significant effect. We have found that the asymptotic value of the “hair” appears directly in the expectation values of operators in the boundary state, giving a precise realization of the holographic proposal of ‘t Hooft and Susskind [22]. Our analysis has been mainly at a linearised level and back reaction on the metric was neglected.

We hope to extend our work in the future to situations like the formation of a black hole where the nonlinearities of gravity are more important.

S-matrix, bulk commutators and local physics: Using the operator relation between bulk and boundary fields, we can reiterate some points made in [5, 4]. It is apparent that transition amplitudes between physical states in the bulk can be computed from the boundary theory. Specifically, prepare a bulk state $|\Psi\rangle = a_{k_1}^\dagger \cdots a_{k_n}^\dagger |0\rangle$ and evolve it forward in time as $|\Psi\rangle \rightarrow e^{-iHt}|\Psi\rangle$. Precisely the same operation can be performed in the boundary theory: the expansion (17) and the equivalence $a_k = b_k$ allow one to prepare the initial state, and time evolution with respect to the CFT Hamiltonian is identified with time evolution in the bulk. Thus transition amplitudes in the boundary theory can be reinterpreted as amplitudes in the bulk.

An unusual feature of this map is that the radial coordinate in the bulk does not appear in the boundary mode expansion (17). This makes it difficult to check aspects of bulk locality such as the commutation of operators at spacelike separation. However, as we have explicitly shown, the boundary theory has access to data on the radial position of localized bulk probes in the characteristic scale of their boundary images. As we have discussed, classical locality of the bulk physics is mapped, at large N , into the approximate independence of collective fluctuations of objects at different boundary scales. So the breakdown of locality in quantum gravity should be understood in terms of incomplete decoupling of scales in the boundary theory at finite N . We are investigating this issue and hope to report on it elsewhere.

6.2 Holographic description of horizons

We have accumulated enough tools to suggest how spacetime causal structure will be reflected in the boundary theory. Here we present a qualitative discussion - further details will appear elsewhere.

Black holes and thermal states: AdS-Schwarzschild black holes [24] and the BTZ black hole [25] (see [26] for a review) both have maximally extended solutions with two asymptotic regions, each with a timelike boundary at spatial infinity.¹² In such spacetimes, the bulk Hilbert space is a product of two identical copies, each accessed by a single asymptotic region [28]. For example, thermal states such as the Hartle-Hawking vacuum are written as correlated tensor products of states:

$$|\text{HH}\rangle = \sum_n e^{-\beta\omega_n} |n, \omega_n\rangle_I \otimes |n, \omega_n\rangle_{II} , \quad (56)$$

¹²The disconnected topology of the boundary in the BTZ case can be seen directly from its orbifold construction [27].

where Hilbert spaces I and II are formally identical. Tracing over one copy in the product produces the thermal ensemble accessible to the other asymptotic observer. In fact, this construction is a standard method of describing real-time, finite-temperature field theories; one studies operators that access one Hilbert space that is correlated appropriately with another. The auxiliary Hilbert space then functions as an external bath which thermalizes the system. The situation is entirely parallel from the boundary perspective. The boundary of spacetime has two disconnected components and so the CFT Hilbert space factorizes into a product of two identical pieces. The choice of bulk vacuum is reflected in the CFT vacuum as discussed in Sec. 2.3 and Appendix A. Tracing over one boundary component leaves a thermal state accessible to one asymptotic region.

Scale-radius duality: We would like to use the scale-radius duality discussed in this paper to study motions towards the horizon from the boundary perspective. For pure AdS, scale-radius duality originates in the dual action of the bulk radial isometry and the boundary scale transformations. In fact, these are not symmetries of the black hole. The curvature of Schwarzschild and the discrete identifications of BTZ break the isometry group, and the corresponding thermal boundary state breaks conformal invariance. Nevertheless, there are arguments that the bulk-boundary duality continues to hold since the spacetime is *asymptotically* AdS [1, 3, 29, 30]. What is more, BTZ black holes *locally* enjoy the same symmetries as AdS and so motions of local probes continue to map to motions in boundary scale. (These points can be made quantitatively using the methods of this paper, as we hope to discuss elsewhere.)

Horizons from the boundary perspective: Armed with the scale-radius duality, we introduce a bulk probe that starts near the boundary and falls towards the horizon. From the boundary perspective, the position of the horizon is represented by the thermal scale. The boundary probe starts life at a very high scale much above the thermal bath. This separation of scales allows it to spread unimpeded as though the bath was absent. As the bulk object falls through the horizon, its boundary dual reaches the thermal scale. Falling through the horizon in the bulk is reflected on the boundary as thermalization due to interaction with the thermal bath. Note that this does not mean that the probe state has “mixed” with the density matrix describing the black hole. Rather, interactions with the thermal bath make the probe state look like a typical state in the ensemble. As the bulk object penetrates to the singularity increasing the black hole mass, thermalization of its boundary dual raises the boundary temperature. The horizon as a causal construct preventing extraction of information is only “real” to the degree that thermalization obscures the history of a state.

Black holes from collapse: We can also consider black holes formed from collapsing shells of matter. Again, from the boundary perspective, the state will spread out until it reaches the scale characteristic of the temperature of the black hole it has created. The degree to which the resulting horizon is sharp will be the degree to which the final state is difficult to measure due to the complicated way that the information about the configuration is spread out over modes at low spatiotemporal scales. The causal structure of the black hole appears as a statistical phenomenon and is precise in the thermodynamic limit.

6.3 Conclusions

In summary, we have developed techniques for describing bulk probes from the boundary perspective. Using our methods, properties of solutions to the bulk equations of motion can be translated into properties of states and expectation values in the boundary theory. We have argued that there is a map between the quantized mode expansions of bulk fields and boundary operators and used our approach to demonstrate a scale-radius duality for several probes. Finally, we outlined the application of our methods to the study of black hole causal structure. This work constitutes a preliminary attempt to address the emergence and eventual breakdown of local spacetime from the gauge theory perspective towards quantum gravity.

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A Correlators and propagator ambiguities

In this appendix we will give a more careful discussion of the methods illustrated in Sec. 2. In particular we will discuss the meaning of the ambiguities in defining the bulk-boundary propagator for Lorentzian spacetimes. We will continue to work in Poincaré coordinates (1).

A.1 Euclidean signature

As in Sec. 2 we study the bulk-boundary correspondence (2) for a classical massive scalar with action (4) that couples to an operator with dimension $2h_+$ (6). In Sec. 2 we worked in position space and followed the procedure of [3] relating the boundary contribution to the bulk action to CFT correlators. In fact, as discussed in [9] this procedure violates Ward identities and must be modified. The correct algorithm is to evaluate a suitably normalized bulk action at $z = \epsilon$ prior to taking $\epsilon \rightarrow 0$.

Improved procedure in momentum space The improved procedure of [9] is easiest to implement in momentum space. The unique solution to the wave equation $(\square - m^2)\phi = 0$ with momentum \vec{k} parallel to the boundary is the Bessel function [2, 9, 4]:

$$\phi \propto e^{i\vec{k}\cdot\vec{x}} z^{\frac{d}{2}} K_\nu(|k|z) \phi_0(\vec{k}) . \quad (57)$$

At the boundary $z \rightarrow 0$, $\phi \rightarrow A(z) + B(z)$ where $A = z^{2h_-}(1 + \dots)$ and $B = z^{2h_+}(1 + \dots)$ and the ellipses indicate series in even powers of z .¹³ For $m^2 > 0$ this is divergent as $z \rightarrow 0$ and requires regulation. According to [9] we work at $z = \epsilon$ with the normalization $\phi = C(\epsilon)\phi_0(\vec{k})e^{i\vec{k}\cdot\vec{x}}$. The authors of [9] set $C(\epsilon) = 1$, but since the scaling of ϕ as it approaches the boundary is important, we choose $C = \epsilon^{2h_-}$:

$$\phi(z, \vec{k}) = \epsilon^{2h_-} \frac{z^{\frac{d}{2}} K_\nu(|k|z)}{\epsilon^{\frac{d}{2}} K_\nu(|k|\epsilon)} \phi_0(\vec{k}) e^{i\vec{k}\cdot\vec{x}} \quad (58)$$

The bulk action (4) then reduces to a boundary term [2, 9]:

$$S = \lim_{z \rightarrow \epsilon} z^{1-d} \delta(\vec{k} + \vec{k}') \phi_0(\vec{k}') \epsilon^{2h_-} \frac{(|k'|z)^{\frac{d}{2}} K_\nu(|k'|z)}{(|k'|\epsilon)^{\frac{d}{2}} K_\nu(|k'|\epsilon)} \partial_z \epsilon^{2h_-} \frac{(|k|z)^{\frac{d}{2}} K_\nu(|k|z)}{(|k|\epsilon)^{\frac{d}{2}} K_\nu(|k|\epsilon)} . \quad (59)$$

As we discussed, $z^{d/2} K_\nu(z) \rightarrow A(z) + B(z) = z^{2h_-}(1 + \dots) + z^{2h_+}(1 + \dots)$. Putting this in (59) as $\epsilon \rightarrow 0$ yields some singular contact terms that we drop and a finite term arising from the mixture of the A and B terms in the numerator and denominator. These give the expected behaviour of the 2-point function:

$$\langle \mathcal{O}(\vec{k}') \mathcal{O}(\vec{k}) \rangle \propto \delta(\vec{k}' + \vec{k}) |k|^{2\nu} . \quad (60)$$

The position space procedure used in Sec. 2 gives different normalizations because the bulk-boundary propagator used there approaches a delta function at the $z = 0$ boundary rather than at $z = \epsilon$.

¹³For integral ν there are some log terms also, but they do not change the basic argument.

Operator expectation values From Sec. 2 we know that turning on source in the boundary theory will lead to nontrivial operator expectation values. These are given by the first variation of $S(\phi)$:

$$\delta S(\phi) = \int_{z \rightarrow 0} d\Sigma^\mu \partial_\mu \phi \delta \phi. \quad (61)$$

where, as in Sec. 2 we have added a small perturbation of the form $\delta\phi = z^{2h-} \delta\phi_0$. In momentum space, with the cutoff procedure prescribed above, $\delta\phi(z = \epsilon, \vec{k}) = \epsilon^{2h-} \delta\phi_0(\vec{k})$. Combining this with $\partial_z \phi$ at $z = \epsilon$, using (61) and dropping contact terms as before, we get a finite one-point function:

$$\langle \mathcal{O}(\vec{k}) \rangle_{\phi_0(\vec{k})} = \langle \mathcal{O}(\vec{k}) \mathcal{O}(-\vec{k}') \rangle_{\phi_0(\vec{k}')} \quad (62)$$

where the two-point function was given in (60) and is evaluated here in the absence of a source. So the interesting part of the one-point function – the part which does not come from coincidence of the source and the operator insertion – arises from the subleading part of the source term which scales as z^{2h+} at the boundary. The bulk of this paper rests on the independent specification of this subleading part near $z = 0$ via the addition of bulk probes and the consequent modifications of operator expectations on the boundary.

A.2 Lorentzian signature

We again begin by discussing free, massive scalar fields. In Lorentzian signature, normalizable solutions to the wave equation exist; so specifying the fields at $z = 0$ does not uniquely specify the field configuration in the bulk. As discussed in [4] the normalizable solutions form the bulk Hilbert space which is dual to the space of boundary states. There is also a spectrum of non-normalizable modes that implement boundary conditions and are dual to boundary sources.

In the supergravity effective action, the normalizable solutions can appear in two places. First, classical field theory in the bulk involves expanding the bulk action around classical, normalizable backgrounds. We will see that this corresponds to turning on expectation values for CFT operators. Secondly, the bulk-boundary propagator is not uniquely specified by the asymptotic behaviour $z^{2h-} \delta^d(\vec{x} - \vec{x}')$ as $z \rightarrow 0$, since a normalizable solution vanishing at the boundary can always be added to it. This ambiguity is related to the choice of vacuum for the theory and will affect the correlation functions.

Working in momentum space and in Poincaré coordinates (1), we write $\vec{k} = (\omega, \vec{q})$ for the momentum parallel to boundary ($\vec{y} = (t, \vec{x})$). For $k^2 > 0$ (spacelike momenta) the solutions are identical to the Euclidean case and there is no normalizable solutions. For $k^2 < 0$, there are two solutions which are smooth in the interior [4]. One solution is

$$\phi^-(\vec{k}, z) \propto z^{\frac{d}{2}} J_{-\nu}(|k|z) e^{i\vec{k} \cdot \vec{y}} \quad (63)$$

when ν is nonintegral.¹⁴ This solution is not normalizable (for $\nu > 1$ – see [4] for a discussion of the case $\nu < 1$); it behaves as z^{2h-} when $z \rightarrow 0$. So the mode is non-fluctuating and is dual to a source term at the boundary. An independent solution is:

$$\phi^+(\vec{k}, z) = z^{\frac{d}{2}} J_\nu(|k|z) e^{i\vec{k}\cdot\vec{y}}, \quad (64)$$

which is normalizable and behaves as z^{2h+} when $z \rightarrow 0$.

Operator expectation values: The most general solution $\phi(\vec{k}, z)$ which is asymptotic to $\epsilon^{2h-} \phi_0(\vec{k}) e^{i\vec{k}\cdot\vec{y}}$ when $z = \epsilon$ is:

$$\phi(\vec{k}, z) = \Phi(\vec{k}) e^{i\vec{k}\cdot\vec{y}} \phi^+(\vec{k}, z) + \epsilon^{2h-} \frac{\phi^-(\vec{k}, z) + A(\vec{k}) \phi^+(\vec{k}, z)}{\phi^-(\vec{k}, \epsilon) + A(\vec{k}) \phi^+(\vec{k}, \epsilon)} e^{i\vec{k}\cdot\vec{y}} \phi_0(\vec{k}). \quad (65)$$

Here $\Phi(\vec{k})$ is the Fourier component of the classical field configuration we wish to study and

$$K_A(\vec{k}, \epsilon, z) = \epsilon^{2h-} \frac{\phi^-(\vec{k}, z) + A(\vec{k}) \phi^+(\vec{k}, z)}{\phi^-(\vec{k}, \epsilon) + A(\vec{k}) \phi^+(\vec{k}, \epsilon)} \quad (66)$$

is the bulk-boundary Green's function. As we can see, K_A has an ambiguity which is encoded by $A(\vec{k})$. Once again, we insert this into the (quadratic) action and integrate by parts.¹⁵ The result is:

$$S_A(\Phi, \phi_0) = \frac{1}{2} \int_{z=\epsilon} d^d k d^d k' \delta^d(\vec{k} + \vec{k}') z^{1-d} \partial_z \left[\Phi(\vec{k}) \phi^+(\vec{k}, z) K_A(\vec{k}', \epsilon, z) \phi_0(\vec{k}') \right. \\ \left. + \phi_0(\vec{k}) K_A(\vec{k}', \epsilon, z) \phi_0(\vec{k}') \right] \quad (67)$$

From this we can read off the one-point function:

$$\langle \mathcal{O}(\vec{k}) \rangle_{\phi_0} = \lim_{z \rightarrow \epsilon} \partial_z \left[\left(\Phi(-\vec{k}) \phi^+(-\vec{k}, z) + \phi_0(-\vec{k}) \right) K(\vec{k}, \epsilon, z) \right] \quad (68)$$

When $\phi_0 = 0$, it is easy to see that

$$\langle \mathcal{O}(\vec{k}) \rangle = \frac{d}{2} \Phi(-\vec{k}) \quad (69)$$

The two-point function is:

$$\langle \mathcal{O}(\vec{k}) \mathcal{O}(\vec{k}') \rangle_{\phi_0} = \langle \mathcal{O}(\vec{k}) \rangle_{\phi_0} \langle \mathcal{O}(\vec{k}') \rangle_{\phi_0} + \lim_{z \rightarrow \epsilon} \delta^d(\vec{k} + \vec{k}') \partial_z K(\vec{k}, \epsilon, z) \quad (70)$$

In this case, the parts of K which behave as z^{2h+} and z^{2h-} are explicitly labeled by ϕ^+ and ϕ^- . As in the Euclidean case, after dropping contact terms the finite piece as $\epsilon \rightarrow 0$ comes combinations of ϕ^- terms in the numerator and ϕ^+ terms in the denominator (and vice-versa) of (66). Thus the connected part of the two-point function will depend quite strongly on $A(\vec{k})$.

¹⁴When ν is integral, $\phi^-(\vec{k}, z) \propto z^{\frac{d}{2}} Y_\nu(|k|z) e^{i\vec{k}\cdot\vec{y}}$

¹⁵In fact the integration by parts gives an oscillating surface term at $z \rightarrow \infty$. This should be understood in terms of an infrared cutoff in the gauge theory.

Interpretation of propagator ambiguities: The dependence of these correlators on Φ is easy to understand. Turning on such a classical background implies that the bulk theory is in a “coherent” state in which many modes have an expectation value. The map between bulk and boundary states discussed in Sec. 2 implies that the CFT should also be in a “coherent” state built from the modes of the dual boundary operator. Not surprisingly, the CFT operator has an expectation value in this state. Roughly speaking, the expectation value is dual the part of ϕ that behaves as z^{2h_+} at the boundary.

In addition to the freedom to specify a classical background/coherent state via Φ , there is an ambiguity arising from the freedom to specify $A(\vec{k})$.¹⁶ The origin of the ambiguity is the freedom to specify a vacuum state. In standard field theory, Wick rotation from a Euclidean signature specifies a certain propagator. But if we were simply to ask that the two-point function satisfy the appropriate inhomogenous differential equation, we would have a much wider range of choices; we could pick some combination of the advanced, retarded, and Feynman propagators, and we could add propagators that solved the associated homogenous differential equation. Our choice would depend on boundary conditions and the choice of vacuum. Here we choose A via analytic continuation from the Euclidean case; the combination of $J_\nu(|\vec{k}|z)$ and $J_{-\nu}(|\vec{k}|z)$ will be that equal to $K_\nu(i|\vec{k}|z)$:

$$K_\nu(i|\vec{k}|z) = \frac{\pi e^{i\pi\nu/2} J_{-\nu}(-|\vec{k}|z) - e^{-i\pi\nu/2} J_\nu(-|\vec{k}|z)}{2 \sin(\nu\pi)}. \quad (71)$$

In some spacetimes (e.g. those containing black holes) there are inequivalent vacua defined with respect to different times. Presumably this would be reflected in the the bulk-boundary propagator.

B The Bulk Response: Another method

The relation between the dilaton field and the vev of $Tr F^2$ was obtained in Sec. 3 by considering the response of the bulk and boundary theories to small perturbations. An alternate method for computing the bulk response in the presence of branes is to do things in the opposite order. Instead of starting with the bulk soliton and evaluating the effect of a perturbation, we could consider a perturbation of AdS and ask about the change in the action on including the soliton. Here we study the D-instanton using the latter technique.

Consider a perturbation of the dilaton in AdS_5 which reduces to a delta function at $\vec{x} = \vec{x}_a$ on the boundary. The bulk value of this perturbation is:

$$\delta\phi = \frac{6}{\pi^2} \frac{z^4}{[z^2 + |\vec{x} - \vec{x}_a|^2]^4}. \quad (72)$$

¹⁶We are grateful to Tom Banks and Emil Martinec for a discussion of this issue.

The D-instanton is a source for the dilaton and axion and so the linearized effects of a D-instanton at $z = \tilde{z}$ and $\vec{x} = \vec{x}_1$ can be incorporated by adding a term in the action of the form:

$$S_{Disnt} = 2\pi \int J(e^{-\phi} + i\chi) \quad (73)$$

where $J = \delta(x_0 - \tilde{x}_0)\delta(\vec{x} - \vec{x}_a)$ stands for the D instanton source. The coefficient in front of the integral in eq. (73) is obtained by requiring that the D-instanton action is given by $S = 2\pi/g_s$. Writing $e^{-\phi} = \frac{1}{g_s}(1 - \delta\phi)$ and substituting for $\delta\phi$ from (72) gives:

$$S_{Disnt} = -\frac{48}{4\pi g_s} \frac{x_0^4}{[x_0^2 + |\vec{x} - \vec{x}_a|^2]^4}. \quad (74)$$

This is exactly equal to the bulk response in (27).

This method gives the same bulk response as in Sec. 3 because we are working in the linearized limit. An analogy with electrodynamics is useful. Introduce a small perturbation of the electrostatic potential surrounding a system of charges. The energy of the resulting system can be computed from the action for the electromagnetic field and the source coupling to the field. Then the bulk contribution vanishes by the equations of motion and we are left with a surface integral like (26). This method is followed in most of this paper. Alternatively, the energy is $E = \sum_i q_i V_i$, where q_i and V_i are the charge and the potential at the position of each charge respectively. So the change in energy is determined by the perturbation of the potential at the location of each charge. This is exactly analogous to (73).

The same reasoning works for a general linear system. The supergravity theory under consideration here is certainly not linear, but the D-branes act as small sources. Since they carry a charge of order $\sim 1/g_s$ the corresponding changes in the supergravity fields are of order $\mathcal{O}(g_s^2 \times 1/g_s)$. In the large N limit that we are working in, these changes are of order $\mathcal{O}(1/N)$ and are therefore small¹⁷. Thus to leading order in $1/N$ (and for purposes of calculating the one-point functions) one can work with a supergravity action expanded to quadratic order about the AdS background. The resulting system is in effect linear and the two methods for evaluating the bulk response must then agree.

References

- [1] J. Maldacena, “The large N limit of superconformal field theories and supergravity”, Harvard preprint HUTP-98-A097, hep-th/9712200.

¹⁷For example, from eq. (21), eq.(23), we see that the dilaton is changed from its asymptotic value by an amount $\delta\phi = \frac{1}{N} \frac{24\pi}{g_s N} \frac{x_0^4 \tilde{x}_0^4}{[\tilde{x}_0^2 + |\vec{x} - \vec{x}_a|^2]^4}$. In the large N limit where $g_s N$ is kept fixed this is of $\mathcal{O}(1/N)$.

- [2] S.S. Gubser, I.R. Klebanov and A.M. Polyakov, “Gauge theory correlators from noncritical string theory”, Princeton preprint PUPT-1767, hep-th/9802109.
- [3] E. Witten. “Anti-de Sitter space and holography”, IAS preprint IASSNS-HEP-98-15, hep-th/9802150.
- [4] V. Balasubramanian, P. Kraus and A. Lawrence, “Bulk vs. boundary dynamics in anti-de Sitter spacetime”, Harvard preprint HUTP-98/A028, Caltech preprint CALT68-2171, hep-th/9805171
- [5] T. Banks and M.B. Green, “Non-perturbative effects in $AdS_5 \times S^5$ string theory and $d = 4$ SUSY Yang-Mills”, hep-th/9804170.
- [6] C-S. Chu, P-M. Ho and Y-Y. Wu, “D-instanton in $AdS(5)$ and instanton in $SYM(4)$ ”, hep-th/9806103.
- [7] I.I. Kogan and G. Luzón, “D-instantons on the boundary”, hep-th/9806197.
- [8] M. Bianchi, M.B. Green, S. Kovacs and G. Rossi, “Instantons in supersymmetric Yang-Mills and D-instantons in IIB superstring theory”, hep-th/9807033.
- [9] D.Z. Freedman, S.D. Mathur, A. Matusis and L. Rastelli, “Correlation functions in the CFT_d/AdS_{d+1} correspondence”, MIT preprint MIT-CTP-2727, hep-th/9804058.
- [10] S.J. Avis, C.J. Isham and D. Storey, “Quantum field theory in anti-de Sitter space-time”, *Phys. Rev.***D18** (1978) 3565.
- [11] P. Breitenlohner and D.Z. Freedman, “Positive energy in anti-de Sitter backgrounds and gauged extended supergravity”, *Phys. Lett.***B115** (1982) 197; “Stability in gauged extended supergravity”, *Ann. Phys, NY***144** (1982) 249.
- [12] L. Mezincescu and P.K. Townsend, “Stability at a local maximum in higher-dimensional anti-de Sitter space and applications to supergravity”, *Ann. Phys, NY* **160** (1985) 406.
- [13] T. Banks, M.R. Douglas, G. Horowitz and E.J. Martinec, Talk delivered by T. Banks at Strings '98 and work in progress, <http://www.itp.ucsb.edu/>
- [14] I.R. Klebanov, “World volume approach to absorption by nondilatonic branes”, *Nucl. Phys.***B496** (1997) 231, hep-th/9702076; S.S. Gubser, I.R. Klebanov and A.A. Tseytlin, String theory and classical absorption by three-branes, *Nucl. Phys.***B499** (1997) 217, hep-th/9703040.
- [15] A. Rajaraman, “Two-form fields and the gauge theory description of black holes”, hep-th/9803082.

- [16] S.R. Das and S.P. Trivedi, “Three-brane action and the correspondence between $N=4$ Yang-Mills theory and anti-de Sitter space”, hep-th/9804149.
- [17] S. Ferrara, M.A. Lledo and A. Zaffaroni, ”Born-Infeld Corrections to D3 brane action in $AdS_5 \times S^5$ and $N = 4, d = 4$ primary superfields”, hep-th/9805082.
- [18] A.A. Tseytlin, “On nonabelian generalization of Born-Infeld action in string theory”, *Nucl. Phys.***B501** (1997) 41, hep-th/9701125.
- [19] H.J. Kim, L.J. Romans and P. van Nieuwenhuizen, “Mass spectrum of chiral ten-dimensional $N = 2$ supergravity on S^5 ”, *Phys. Rev.* —**D32** (1985) 389.
- [20] O. Aharony and E. Witten, “Anti-de Sitter space and the center of the gauge group”, hep-th/9807205.
- [21] E. Gava, K.S.Narain and M.H. Sarmadi, ”On the bound states of p - and $(p+2)$ -Branes”, hep-th/9704006.
- [22] G. 't Hooft, “Dimensional reduction in quantum gravity”, Salamfest 1993:0284-296, gr-qc/9310026; L. Susskind, “The world as a hologram”, *J.Math.Phys.*36:6377-6396,1995, hep-th/9409089.
- [23] J. Maldacena and A. Strominger, “ AdS_3 black holes and a stringy exclusion principle”, Harvard preprint HUTP-98/A016, hep-th/9804085.
- [24] S.W. Hawking and D. Page, “Thermodynamics of black holes in anti-de Sitter space”, *Comm. Math. Phys.* **87** (1983) 577.
- [25] M. Banados, C. Teitelboim, and J. Zanelli, “The Black hole in three-dimensional space-time”, *Phys. Rev. Lett.* **69** 1849, hep-th/9204099 ; M. Banados, M. Henneaux, C. Teitelboim, and J. Zanelli, “Geometry of the (2+1) black hole”, *Phys. Rev.* **D48** 1506, gr-qc/9302012.
- [26] S. Carlip, “The (2+1)-dimensional black hole”, *Class. Quant. Grav.* **12** (1995) 2853, gr-qc/9506079.
- [27] G.T. Horowitz and D. Marolf, “A new approach to string cosmology”, hep-th/9805207.
- [28] N.D. Birrell and P.C.W. Davies, *Quantum fields in curved space*, Cambridge University Press (1982).
- [29] N. Itzhaki, J. Maldacena, J. Sonnenschein and S. Yankielowicz, “Supergravity and the large- N limit of theories with 16 supercharges”, *Phys. Rev.* **D58** (1998), hep-th/9802042.

- [30] E. Witten, “Anti-de Sitter space, thermal phase transition, and confinement in gauge theories”, hep-th/9803131.