

Holographic Constraint on Scale Separation

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We propose a new consistency condition for the compatibility of a gravitational effective field theory in AdS with a dual holographic description in terms of a family of large- N CFTs. Using large- N factorization of correlation functions combined with a properly defined notion of single- and multiparticle operators, we argue that the cubic scalar bulk couplings for fields dual to operators with extremal arrangements of the conformal dimensions should vanish. We apply this criterion to the 4d $\mathcal{N} = 1$ effective supergravity theory describing the simplest DGKT AdS₄ vacua in type IIA string theory and show that it is nontrivially satisfied. In addition, we calculate explicitly all nonvanishing three-point correlation functions of low-lying scalar operators in the putative 3D CFTs dual to these AdS₄ string theory backgrounds.

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Introduction—The AdS/CFT correspondence provides a powerful quantitative framework to address open questions in the physics of strongly coupled QFTs and quantum gravity. Guided by top-down examples arising from string and M theory and the numerous successful tests of the correspondence, it is possible to delineate some of the universal features of holographic CFTs; see [1,2]. Using these general properties one can then proceed to deduce broader lessons for all consistent quantum gravitational theories in AdS without relying on a specific embedding in a UV complete theory. Our goal in this work is to proceed along these lines and argue for a general constraint imposed by holography on the properties of gravitational EFTs in AdS.

Our starting point is simple and well-known. Suppose that a holographic theory in AdS has scalar fields ϕ_i dual to operators with dimensions Δ_i that obey the relation $\Delta_i = \Delta_j + \Delta_k$. If the scalar fields have a cubic interaction in the AdS bulk, this *extremal* arrangement of the operator dimensions leads to a divergence in the three-point Witten diagram and thus to a seemingly ill-defined three-point function in the dual CFT. This problem was appreciated in the early days of AdS/CFT since it arises in the paradigmatic holographic duality between type IIB string theory on AdS₅ × S⁵ and the 4d $\mathcal{N} = 4$ SYM theory [3]. The resolution in this specific example relies on the details of the structure of type IIB supergravity and the $\mathcal{N} = 4$ SYM theory. The effective supergravity action conspicuously ensures that the extremal three-point vertices

in the bulk vanish and thus the problematic Witten diagrams do not arise. As emphasized in [4], this is realized in the dual $\mathcal{N} = 4$ SYM theory after carefully taking into account the mixing between single- and multitrace operators.

Inspired by this we revisit the problem of calculating extremal three point functions in AdS/CFT and arrive at an important conclusion. The general principles of holography combined with large- N factorization of the correlation functions of holographic CFTs imply that in any effective theory of gravity in AdS with finitely many matter fields the cubic couplings of scalar fields dual to operators with extremal arrangements of the conformal dimensions must vanish. In the context of string compactifications, this condition amounts to a new holographic constraint on the UV completeness of scale separated AdS backgrounds of string and M theory; see [5] for a review.

This holographic constraint may at first appear of limited utility since one may expect that for generic EFTs in AdS the spectrum of massive scalars need to be finely tuned for extremal arrangements to be possible. Interestingly, string theory constructions lead to large classes of scale separated AdS vacua with minimal or no supersymmetry for which these seemingly fine-tuned extremal arrangements of scalar operator dimensions are ubiquitous. This fact was recently emphasized in [6–8] for the DGKT AdS₄ vacua in massive type IIA string theory [9], and it is this class of models that will be the main example on which we test our new constraint. As we show below, in the simplest supersymmetric and non-supersymmetric AdS₄ DGKT vacua based on a $T^6/\mathbb{Z}_3^2$ IIA orientifold compactification with fluxes, the cubic extremal couplings vanish after a nontrivial cancellation. Encouraged by this we proceed and calculate explicitly the nonextremal three-point-correlation functions of all light scalar operators in this model.

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Holography—Our starting point is a two-derivative gravitational effective theory in $d + 1$ dimensions coupled to a *finite* number of matter fields of spin up to $3/2$. For our discussion only the scalar fields in this gravitational EFT play a role and we assume the following effective action:

$$S = \eta \int d^{d+1}x \sqrt{g} [-R + F^{jk} \partial_\mu \phi_j \partial^\mu \phi_k + V], \quad (1)$$

written in Euclidean signature. Here $F^{jk}(\phi)$ specifies the kinetic terms for the real scalar fields ϕ_i and we assume that the potential $V(\phi)$ has at least one critical point corresponding to an AdS_{d+1} solution. For concreteness, we will assume that the UV cutoff scale in the EFT is given by the $(d + 1)$ -dimensional Newton constant G and we adopt the normalization $\eta = 1/(16\pi G_N)$. We can expand the Lagrangian around the AdS vacuum to cubic order to find

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & -R - \frac{d(d-1)}{L^2} + \frac{1}{2} (\partial_\mu \phi_i)^2 + \frac{1}{2} \frac{m_i^2}{L^2} \phi_i^2 \\ & + d_{ijk} \phi_i \partial_\mu \phi_j \partial^\mu \phi_k + \frac{c_{ijk}}{L^2} \phi_i \phi_j \phi_k + \dots, \end{aligned} \quad (2)$$

and we can choose to normalize the AdS scale as $L = 1$. We will assume that this gravitational EFT can be UV completed in a consistent theory of quantum gravity and the AdS vacuum has a holographic dual description in terms of a sequence of CFTs. The central charge of these CFTs, defined in terms of the two-point function of the stress tensor, is proportional to the gravitational coupling $c \sim \eta$, and we assume that it is large, i.e., $c \gg 1$. In addition, to ensure that we have a weakly coupled effective theory in AdS we assume that the masses m_i and the cubic couplings c_{ijk} and d_{ijk} in (2) do not scale with η .

Several clarifications are in order. The type of gravitational EFTs described above includes, but is not limited to, scale separated AdS vacua of string or M theory; see [5]. In these setups the cutoff scale of the EFT is set by the KK scale of the internal manifold. Indeed, the DGKT AdS_4 vacua that we discuss below are specific examples of such scale separated vacua. The putative holographic dual CFT should contain operators of spin more than 2. By assumption there are no such states in the gravitational EFT. This means that there is a higher-spin gap in the CFT spectrum. There should certainly be a large parameter in the sequence of CFTs that determines this gap but we remain agnostic as to whether this parameter is related to c . We note that the type of gravitational EFTs in AdS we discuss here suffer from a cosmological constant problem, as nicely summarized in [10], which makes them ideal targets to study with the tools of AdS/CFT. Finally, we stress that the gravitational EFTs we focus on are qualitatively different from the consistent truncations of supergravity discussed in the AdS/CFT literature. A consistent truncation typically results from a compactification of string/M theory in which

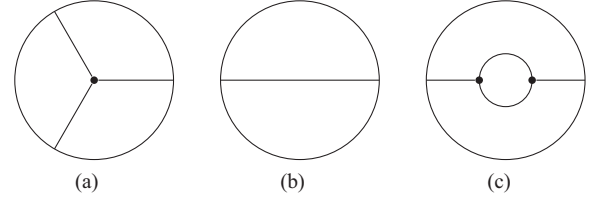


FIG. 1. Two-point (b),(c) and three-point (a) Witten diagrams discussed in the text.

the KK and AdS scales are of the same order and thus the effective theory in AdS_{d+1} has *infinitely* many fields of spin up to 2.

The two-point functions of the operators dual to ϕ_i are fixed by their conformal dimensions Δ_i which are related to the masses as $m_i^2 = \Delta_i(\Delta_i - d)$. The three-point functions of these scalar operators are fully fixed by the conformal symmetry in terms of Δ_i and an overall constant which can be computed holographically using the effective action (2); see [11] for a review. As discussed in [12], the CFT three-point-function coefficients C_{ijk} are determined by the following linear combination of cubic bulk couplings:

$$c'_{ijk} = c_{ijk} + \frac{m_i^2 - m_j^2 - m_k^2}{2} d_{ijk}, \quad (3)$$

and take the form $C_{ijk} = \eta^{-2} c'_{ijk} A_{\Delta_i \Delta_j \Delta_k}$. The factor $A_{\Delta_i \Delta_j \Delta_k}$ arises from evaluating the bulk three-point Witten diagram [see Fig. 1(a)], and is given by [13]

$$A_{\Delta_i \Delta_j \Delta_k} = \frac{\Gamma(\Theta_i) \Gamma(\Theta_j) \Gamma(\Theta_k) \Gamma(\Theta - \frac{d}{2})}{2\pi^d \Gamma(\Delta_i - \frac{d}{2}) \Gamma(\Delta_j - \frac{d}{2}) \Gamma(\Delta_k - \frac{d}{2})}, \quad (4)$$

where $\Theta = (\Delta_i + \Delta_j + \Delta_k)/2$ and $\Theta_n = \Theta - \Delta_n$. Importantly, $A_{\Delta_i \Delta_j \Delta_k}$ diverges when one of the $\Theta_i = -n$ with $n \in \mathbb{Z}_{\geq 0}$. For $n = 0$, this corresponds to an *extremal* arrangement of dimensions, i.e., $\Delta_i = \Delta_j + \Delta_k$, while the cases with $n > 0$ are called *superextremal*. We now proceed to quantify this divergence and propose a mechanism for obtaining finite three-point functions consistent with the structure of a holographic CFT.

Bulk: A basic tenet of holography is that for each single-particle bulk field ϕ_i there is a corresponding dual operator, which we will refer to as a “single-particle operator” (SPO) denoted by $\mathcal{O}_{\phi_i}^{(s)}$ [14], whose classical conformal dimension $\Delta_i^{(0)}$ is determined by the mass of ϕ_i via $m_i^2 = \Delta_i^{(0)}(\Delta_i^{(0)} - d)$. In the semiclassical regime, the scaling of the holographic two-point and three-point functions can be deduced from the effective Lagrangian (2) and the tree-level Witten diagrams depicted in Figs. 1(a) and 1(b). One finds that every propagator in AdS scales as $1/c$ and every three-point interaction vertex scales as c . This leads to the following scaling of the two-point and

three-point functions,

$$\langle \mathcal{O}_{\phi_i}^{(s)} \mathcal{O}_{\phi_i}^{(s)} \rangle \sim \frac{1}{c}, \quad \langle \mathcal{O}_{\phi_i}^{(s)} \mathcal{O}_{\phi_j}^{(s)} \mathcal{O}_{\phi_k}^{(s)} \rangle \sim \frac{1}{c^2}. \quad (5)$$

A further consequence of the EFT (1) is that the anomalous dimension of SPOs is $1/c$ suppressed; i.e., at large c the dimensions Δ_i have the expansion

$$\Delta_i = \Delta_i^{(0)} + \frac{\gamma_i}{c} + \dots \quad (6)$$

Here γ_i is a real number independent of c and the result follows from considering the bulk one-loop correction to the two-point function of $\mathcal{O}_i$ [see Fig. 1(c)], with only a *finite* number of fields running in the loop.

Field theory: The scalar operator spectrum of the dual CFT consists of “single-trace” operators $\mathcal{O}_i^{\text{ST}}$ as well as multitrace operators, which arise from the coincident point limit of single-trace ones. Of particular importance for our discussion are double-trace operators, which we denote by $\mathcal{O}_{i,j}^{\text{DT}} \equiv [\mathcal{O}_i^{\text{ST}} \mathcal{O}_j^{\text{ST}}]$. One can distinguish between single- and double-trace operators by considering the scaling of their correlators with c . Due to large- N factorization, see [2] for a detailed discussion, we have that the two-point- and three-point-functions scale as

$$\begin{aligned} \langle \mathcal{O}_i^{\text{ST}} \mathcal{O}_i^{\text{ST}} \rangle &\sim 1, \quad \langle \mathcal{O}_{i,j}^{\text{DT}} \mathcal{O}_{i,j}^{\text{DT}} \rangle \sim 1, \\ \langle \mathcal{O}_i^{\text{ST}} \mathcal{O}_j^{\text{ST}} \mathcal{O}_k^{\text{ST}} \rangle &\sim \frac{1}{\sqrt{c}}, \quad \langle \mathcal{O}_i^{\text{ST}} \mathcal{O}_j^{\text{ST}} \mathcal{O}_{i,j}^{\text{DT}} \rangle \sim 1, \end{aligned} \quad (7)$$

where we have normalized the operators such that their two-point functions do not scale with c . Note that the last scaling is due to an allowed “GFF contraction” between $\mathcal{O}_{i,j}^{\text{DT}}$ with $\mathcal{O}_i^{\text{ST}}$ and $\mathcal{O}_j^{\text{ST}}$ which is represented in the bulk dual by a disconnected Witten diagram.

The holographic dictionary: Holography relates the two notions of operators introduced above. In a generic situation we simply have that single-particle states in the bulk are dual to single-trace operators, i.e. $\mathcal{O}_{\phi_i}^{(s)} = 1/\sqrt{c} \mathcal{O}_i^{\text{ST}}$, where we included an overall factor of $1/\sqrt{c}$ to match the scaling of two-point functions in (5) and (7).

However, in a situation where the theory at hand allows for an *extremal* arrangement; i.e., when three bulk fields $\phi_{i,j,k}$ have masses such that the condition

$$\Delta_i^{(0)} + \Delta_j^{(0)} = \Delta_k^{(0)} \quad (8)$$

is satisfied, there is an inherent *ambiguity* in the identification of the operator $\mathcal{O}_{\phi_k}^{(s)}$ due to a degeneracy in the CFT spectrum. Since both the single-trace operator $\mathcal{O}_k^{\text{ST}}$ and the double-trace operator $\mathcal{O}_{i,j}^{\text{DT}}$ have classical dimension $\Delta_k^{(0)}$, we need to consider the possibility of mixing,

$$\mathcal{O}_{\phi_k}^{(s)} = \frac{1}{\sqrt{c}} \left(\mathcal{O}_k^{\text{ST}} + \frac{B}{c^\beta} \mathcal{O}_{i,j}^{\text{DT}} + \dots \right), \quad (9)$$

for real β and an order-one number B which depends on the specific CFT. The dots stand for possible higher-trace admixtures. As proposed in [4,15,16] in the context of $\mathcal{N} = 4$ SYM theory, and more generally in [12], the precise linear combination in (9) is determined by demanding that the SPO $\mathcal{O}_{\phi_k}^{(s)}$ is *orthogonal to all multitrace operators*:

$$\langle \mathcal{O}_{\phi_k}^{(s)} [\mathcal{O}_{i_1}^{\text{ST}} \dots \mathcal{O}_{i_n}^{\text{ST}}] \rangle \stackrel{!}{=} 0, \quad n \geq 2. \quad (10)$$

While this provides a nonperturbative definition of SPOs, it may be difficult to implement practically in an abstract holographic CFT. Nevertheless, one can derive a lower bound on the exponent β in (9) by demanding consistency with the scaling of two-point functions in (5) and (7). To this end, we can use (9) and find the following scaling of the two-point function of SPOs,

$$\begin{aligned} \langle \mathcal{O}_{\phi_k}^{(s)} \mathcal{O}_{\phi_k}^{(s)} \rangle &= \frac{1}{c} \left[\langle \mathcal{O}_k^{\text{ST}} \mathcal{O}_k^{\text{ST}} \rangle + \frac{B^2}{c^{2\beta}} \langle \mathcal{O}_{i,j}^{\text{DT}} \mathcal{O}_{i,j}^{\text{DT}} \rangle + \dots \right] \\ &\sim \frac{1}{c} \left[1 + \frac{1}{c^{2\beta}} + \dots \right], \end{aligned} \quad (11)$$

where we used the large- c scalings in (7) and omitted further subleading terms. The first term on the last line of (11) yields the expected $1/c$ scaling as in (5). For this not to be spoiled by the second term, we need that β satisfies the inequality

$$\beta \geq 0. \quad (12)$$

One can argue along similar lines that the scalings with c of potential higher-trace terms in (9) obey an analogous bound.

We are now ready to discuss the situation in which the effective Lagrangian (2) contains nonvanishing extremal couplings $c'_{ijk} \neq 0$. In this case the holographic three-point functions of the SPOs have the following large- c behavior:

$$\langle \mathcal{O}_{\phi_i}^{(s)} \mathcal{O}_{\phi_j}^{(s)} \mathcal{O}_{\phi_k}^{(s)} \rangle = \frac{1}{c} \cdot c'_{ijk} \frac{\hat{a}}{(\gamma_i + \gamma_j - \gamma_k)} + \dots, \quad (13)$$

which follows from plugging (6) into (4) and expanding to leading order in the large- c limit. Here $\hat{a}$ is a numerical factor depending on the conformal dimensions $\hat{a} = \Gamma(\Delta_k^{(0)} - d/2) / [2^{3/2} \pi^{d/4} \Gamma(\Delta_k^{(0)})] \prod_{n=\{i,j,k\}} \sqrt{\Gamma(\Delta_n^{(0)}) / \Gamma(\Delta_n^{(0)} + 1 - d/2)}$. The enhancement of the scaling in (13) to $1/c$ instead of the expected $1/c^2$ in (5) arises due to a pole in the Γ functions in (4).

On the other hand, we can use the definition of the SPO in (9) for $\mathcal{O}_{\phi_k}^{(s)}$ and use (7) to find the following scaling of

the same correlator with c :

$$\langle \mathcal{O}_{\phi_i}^{(s)} \mathcal{O}_{\phi_j}^{(s)} \mathcal{O}_{\phi_k}^{(s)} \rangle = \frac{1}{c^{3/2}} \left[\frac{A}{\sqrt{c}} + \frac{B}{c^\beta} \cdot D + \dots \right], \quad (14)$$

where A and D are order-one numbers that depend on the particular CFT and can be computed from the three-points functions $\langle \mathcal{O}_i^{\text{ST}} \mathcal{O}_j^{\text{ST}} \mathcal{O}_k^{\text{ST}} \rangle$ and $\langle \mathcal{O}_i^{\text{ST}} \mathcal{O}_j^{\text{ST}} \mathcal{O}_{i,j}^{\text{DT}} \rangle$.

The expressions in (13) and (14) need to have the same large- c scaling. This is only possible if the second term in (14) becomes the leading contribution at large c , which in turn fixes β to be

$$\beta = -\frac{1}{2}. \quad (15)$$

This however is in contradiction with the bound on β in (12). We thus find that a nonvanishing bulk extremal cubic coupling implies that the notion of SPO for ϕ_k breaks down due to $\beta < 0$ in (9). This means that the gravitational EFT described by the bulk effective Lagrangian (2) has a dual CFT which fails to obey large- N factorization and to describe the dynamics of the single-particle weakly coupled bulk field ϕ_k , thereby violating the standard premise of holography.

The only way to resolve this issue is to either give up on the standard notion of a holographic dual for the particular gravitational EFT in AdS with extremal operator arrangements in the spectrum, or have all extremal cubic couplings c'_{ijk} vanish. This amounts to no cubic interactions in the bulk and thus the extremal three-point correlators vanish. Note that a vanishing of c'_{ijk} can arise in two ways. One can have the trivial situation where both the potential and derivative cubic terms in the Lagrangian (2) vanish, $c_{ijk} = d_{ijk} = 0$. More interestingly, c'_{ijk} can vanish due to a cancellation between nonzero c_{ijk} and d_{ijk} couplings. The second option has recently been emphasized in the case of a 4d $\mathcal{N} = 1$ holographic SCFT [12], and as we describe next this is also realized in a simple example of the DGKT construction. We close this discussion by noting that our argument can be repeated also for superextremal arrangements of the conformal dimensions [see below (4)] to conclude that in the bulk dual of a consistent holographic CFT all cubic superextremal couplings c'_{ijk} should vanish. Finally, we note that in holographic CFTs that are not described by a $(d+1)$ -dimensional EFT with finitely many fields, or for which c_{ijk} or d_{ijk} scale with c , there could be nonvanishing extremal cubic couplings, see [17] for a recent discussion and [18] for a top-down example in type IIB string theory.

DGKT—We now focus on a well-known class of scale separated AdS₄ vacua in massive type IIA string theory arising from a compactification on CY manifolds in the presence of background fluxes and O6 planes. The simplest example of such a construction is based on the orbifold $T^6/\mathbb{Z}_3$ [9]. In the absence of any background flux the 4D

effective theory has 13 complex scalars—three corresponding to the Kähler moduli of the two-cycles in the manifold, one arising from the 10d dilaton combined with a C_3 axion, and nine associated to the collapsed two-cycles at the nine fixed points of the orbifold action. Notably, all complex structure moduli are frozen by the orientifold involution. As shown in [9], all these massless scalars acquire a mass when one turns on the Romans mass, m_0 , the H_3 flux p , the F_6 flux, e_0 , and the three F_4 fluxes, e_i , through the independent four-cycles in the internal space. In the simplest model on which we focus here, all other fluxes are turned off. In the absence of D6 branes, C_7 tadpole cancellation implies $m_0 p < 0$ and without loss of generality we take $m_0 > 0$, $p < 0$. Importantly, the e_i fluxes are not constrained and can be scaled to be parametrically large. This ensures a weakly coupled IIA regime in which both α' and g_s corrections are small and the volume of $T^6/\mathbb{Z}_3$, which determines the KK scale is much smaller than the size of the noncompact 4D space-time. As emphasized in [9], the 9 moduli associated with the orbifold singularities can be made parametrically heavier than the KK scale by appropriately tuning the fluxes through the singular two-cycles. With all these ingredients we arrive at a 4d $\mathcal{N} = 1$ effective supergravity theory which is valid below the KK scale set by the large fluxes e_i and contains four chiral multiplets and the gravity multiplet. We now proceed to study this model through a holographic lens.

The 4D $\mathcal{N} = 1$ supergravity theory is specified by a Kähler potential and a superpotential that depend on the Kähler moduli scalars t_i and the axiodilaton S [19,20],

$$K = -\log(8\kappa) - 4 \log [\text{Im}(S)] - \sum_{i=1}^3 \log [\text{Im}(t_i)],$$

$$W = e_0 + e_i t_i - \kappa m_0 t_1 t_2 t_3 - 2pS, \quad \kappa = \frac{1}{8\sqrt{3}}. \quad (16)$$

The kinetic terms are $\mathcal{L}_K = -2K_{I\bar{J}} \partial_\mu \Phi^I \partial^\mu \bar{\Phi}^{\bar{J}}$, where $\Phi_I = \{t_i, S\}$ and $K_{I\bar{J}} = \partial_{\Phi^I} \partial_{\bar{\Phi}^{\bar{J}}} K$. The scalar potential is given by the standard $\mathcal{N} = 1$ supergravity expression

$$V = 2e^K \left(\sum_I K^{I\bar{J}} D_I W D_{\bar{J}} \bar{W} - 3|W|^2 \right), \quad (17)$$

where $D_I W = \partial_I W + W \partial_I K$. One can show that the only critical point of the potential is at the following values of the scalar fields

$$\text{Re}(t_i) = 0, \quad \text{Im}(t_i) = \frac{1}{|e_i|} \sqrt{\frac{5}{3} \left| \frac{e_1 e_2 e_3}{\kappa m_0} \right|},$$

$$\text{Re}(S) = \frac{e_0}{2p}, \quad \text{Im}(S) = \frac{1}{|p|} \sqrt{\frac{80}{27} \left| \frac{e_1 e_2 e_3}{\kappa m_0} \right|}. \quad (18)$$

The value of the cosmological constant at this critical point determines the scale of the AdS₄ vacuum and reads,

$$V_{\text{AdS}} \equiv -\frac{6}{M_P^2 L^2} = -\frac{243\sqrt{3}\kappa^{3/2}p^4|m_0|^{5/2}}{6400(e_1^2 e_2^2 e_3^2)^{3/4}}. \quad (19)$$

There are in fact four distinct AdS_4 vacua with the same value of the cosmological constant, which are distinguished by the signs of the fluxes $s_i \equiv \text{sign}(m_0 e_i)$. When all $s_i = -1$ the critical point of V also obeys $D_I W = 0$ and thus we preserve $4d \mathcal{N} = 1$ supersymmetry. For all other choices of signs supersymmetry is broken.

If we take the large e_i flux parameters to be all of the same order, $e_i \sim \bar{e}$, the central charge of the dual CFT scales as $c \sim \bar{e}^{9/2}$; see [21]. It is in this regime of large $\bar{e} \gg 1$ that we have a scale separated AdS_4 vacuum of type IIA supergravity. Indeed, setting $L = 1$ in (19) and comparing to (1) and (2) we see that $\eta = (M_P^2/2) \gg 1$ as needed for the applicability of our holographic discussion.

To compute the holographic two-point and three-point functions for the eight scalar fields in these AdS_4 vacua, we need to expand the supergravity Lagrangian to cubic order. At the quadratic level, we choose a basis for the scalar fluctuations in which the kinetic terms are canonically normalized and the mass matrix is diagonal. We provide more details on this choice of basis in Supplemental Material [22]. The effective quadratic Lagrangian is then given by

$$\begin{aligned} \mathcal{L}^{(2)} = & \frac{1}{2} \sum_{\alpha=0}^3 (\partial \varphi_\alpha)^2 + 35\varphi_0^2 + 9 \sum_{i=1}^3 \varphi_i^2 \\ & + \frac{1}{2} \sum_{\alpha=0}^3 (\partial a_\alpha)^2 + \frac{43-45s}{2} a_0^2 + \frac{25+15s}{2} \sum_{i=1}^3 a_i^2, \end{aligned} \quad (20)$$

where $\{\varphi_0, \varphi_i\}$ denote the fluctuations of the four scalars, while $\{a_0, a_i\}$ are those of the pseudoscalars. Importantly, the masses depend only on the product of signs $s \equiv s_1 s_2 s_3$ and therefore we find only two distinct mass spectra for the four AdS vacua. Here we focus on the supersymmetric choice $s = -1$ and relegate the analysis for $s = 1$ to Supplemental Material [22]. The SPOs dual to the four scalar fields have conformal dimensions given by

$$\Delta_\varphi = \{10, 6, 6, 6\}, \quad (21)$$

whereas the conformal dimensions of the operators dual to the pseudoscalars are

$$\Delta_a = \{11, 5, 5, 5\}. \quad (22)$$

Interestingly, this spectrum of conformal dimensions is not specific to the $T^6/\mathbb{Z}_3^2$ choice of internal CY orientifold and arises universally for all other DGKT vacua constructed from CY orientifolds [7]. The fact that the conformal dimensions of the scalars and pseudoscalars differs by 1 is compatible with the $3D \mathcal{N} = 1$ supersymmetry preserved by the putative dual SCFT; see [23].

At cubic order the effective Lagrangian for the eight scalars reads,

$$\begin{aligned} \mathcal{L}^{(3)} = & \sum_{\alpha=0}^3 \sum_{0 \leq \beta \leq \gamma} d_{\varphi_\alpha a_\beta a_\gamma} \varphi_\alpha \partial_\mu a_\beta \partial^\mu a_\gamma \\ & + \sum_{0 \leq \alpha \leq \beta \leq \gamma} c_{\varphi_\alpha \varphi_\beta \varphi_\gamma} \varphi_\alpha \varphi_\beta \varphi_\gamma + \sum_{\alpha=0}^3 \sum_{0 \leq \beta \leq \gamma} c_{\varphi_\alpha a_\beta a_\gamma} \varphi_\alpha a_\beta a_\gamma. \end{aligned}$$

The precise values of the cubic couplings depend only on the sign s and thus we find only two distinct sets of three-point functions for the four AdS_4 vacua. Here we focus on the supersymmetric choice $s = -1$ and discuss the other sign in Supplemental Material [22].

The spectrum in (21) and (22) exhibits two possible extremal arrangements, namely (10,5,5) and (11,6,5). In the first case, the corresponding cubic couplings are determined by the matrices

$$d_{\varphi_0 a_i a_j} = \frac{4}{3\sqrt{13}} M_{ij}, \quad c_{\varphi_0 a_i a_j} = -\frac{100}{3\sqrt{13}} M_{ij}, \quad (23)$$

while in the second case we find

$$d_{\varphi_j a_i a_0} = -\frac{4}{3\sqrt{13}} M_{ij}, \quad c_{\varphi_j a_i a_0} = -\frac{160}{3\sqrt{13}} M_{ij}, \quad (24)$$

with $i, j = 1, 2, 3$ and the matrix M_{ij} given by

$$M_{ij} = \begin{pmatrix} 0 & 0 & 3\sqrt{\frac{2}{13}} \\ 0 & -3 & 0 \\ 3\sqrt{\frac{2}{13}} & 0 & -\frac{33}{13} \end{pmatrix}. \quad (25)$$

Remarkably, for both sets of extremal cubic couplings (23) and (24) we find a nontrivial cancellation in c'_{ijk} after using (3)

$$c'_{\varphi_0 a_i a_j} = c_{\varphi_0 a_i a_j} + \frac{70-10-10}{2} d_{\varphi_0 a_i a_j} = 0, \quad (26)$$

and similarly for the couplings in (24). This in turn implies that the corresponding extremal 3pt-function coefficients C_{ijk} are zero.

We therefore conclude that the holographic consistency condition described in the previous section is satisfied in the 4d gravitational EFT arising from the DGKT construction. As we show in Supplemental Material [22] the same type of nontrivial cancellations happen for the non-supersymmetric AdS_4 vacua with $s = 1$ for both extremal and super-extremal three-point functions. Most nonextremal scalar three-point couplings allowed by parity are nonzero and we present their explicit values in [22].

Discussion—There are several generalizations of our results that would be worth pursuing. It is important to understand whether our new holographic criterion is obeyed by other DGKT AdS_4 vacua based on more complicated CY orientifolds since the operator spectrum

in these setups always allows for extremal arrangements [7]. It may however be technically challenging to analyze explicitly the cubic interactions in these examples in full generality. It would also be interesting to apply our criterion to proposed AdS_3 scale separated vacua in string theory with integer conformal dimensions and extremal arrangements in the operator spectrum, like the ones in [24,25] and [26]. We note in passing that our constraint is not readily applicable to the KKLT and LVS setups; see [5] for a review, since in the corresponding AdS_4 vacua of type IIB string theory the holographic central charge is not tunable and the conformal dimensions of low-lying scalar operators do not generally admit extremal arrangements. Given our explicit results for the 3pt-functions for the $T^6/\mathbb{Z}_3$ DGKT model it will be interesting to calculate also the scalar 4pt-functions as well as the $1/c$ corrections to the conformal dimensions. To this end one would need to study exchange and one-loop Witten diagrams in the 4d $\mathcal{N} = 1$ effective supergravity.

The putative family of 3D CFTs holographically dual to DGKT-type AdS_4 vacua remains as puzzling as ever due to its several exotic features—a large spectral gap for operators of spin < 2 as well as the absence of relevant or marginal operators. Moreover, there are other constraints on scale-separated AdS_4 vacua that the DGKT construction does not pass; see [27,28]. We hope that our results will help shed new light on this mysterious system on the interface of string theory and holography.

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