

1/4 BPS solutions and the AdS₃/CFT₂ correspondenceYolanda Lozano^{1,*}, Niall T. Macpherson^{2,3,†}, Carlos Nunez^{4,‡} and Anayeli Ramirez^{1,§}¹*Department of Physics, University of Oviedo, Avenida Federico Garcia Lorca s/n, 33007 Oviedo, Spain*²*SISSA International School for Advanced Studies, Via Bonomea 265,
34136 Trieste and INFN sezione di Trieste*³*International Institute of Physics, Universidade Federal do Rio Grande do Norte,
Campus Universitario—Lagoa Nova, Natal, RN 59078-970, Brazil*⁴*Department of Physics, Swansea University, Swansea SA2 8PP, United Kingdom*

(Received 7 October 2019; published 23 January 2020)

We discuss new solutions in massive type IIA supergravity with AdS₃ × S² factors, preserving $\mathcal{N} = (0, 4)$ supersymmetry. We propose a duality with a precise family of quivers that flow to $\mathcal{N} = (0, 4)$ fixed points at low energies. These quivers consist on two families of linear quivers coupled by matter fields. Physical observables such as the central charges provide stringent checks of the proposed duality. A formal mapping is presented connecting our backgrounds with those dual to six dimensional $\mathcal{N} = (1, 0)$ conformal field theories (CFTs), suggesting the existence of a flow across dimensions between the CFTs.

DOI: 10.1103/PhysRevD.101.026014

I. INTRODUCTION

An important by-product of the Maldacena conjecture [1], has been a thorough study of supersymmetric and conformal field theories (CFTs) in various dimensions. In particular, the last two decades have witnessed a large effort in the classification of Type II or M-theory backgrounds with AdS_{d+1} factors, see e.g., [2,3]. The solutions are conjectured to be dual to CFTs in d dimensions with different amounts of SUSY, that can then be studied holographically.

Major progress has been achieved when the CFT preserves half of the maximum number of allowed supersymmetries. For the case of $\mathcal{N} = 2$ CFTs in four dimensions, the field theories studied in [4,5] have holographic duals discussed in [6], and further elaborated (among other works) in [7–13]. The case of five-dimensional CFTs was analyzed from the field theoretical and holographic viewpoints in [14–20], among many other interesting works. An infinite family of six-dimensional $\mathcal{N} = (1, 0)$ CFTs was discussed from both the field theoretical and holographic points of view in [21–28]. For three-dimensional $\mathcal{N} = 4$ CFTs, the field theories presented in [29] were discussed holographically in [30–33], among other works.

The case of two-dimensional CFTs and their AdS duals is particularly attractive, due to the interest that CFTs in two dimensions and AdS₃ solutions present in other areas of theoretical physics. This applies in particular to the microscopical study of black holes, where major progress has been achieved [34–39]. This motivated various attempts at finding classifications of AdS₃ backgrounds and studying their dual CFTs [40–58]. $\mathcal{N} = (0, 4)$ AdS₃ solutions remained however largely unexplored, with known cases following mostly from orbifoldings, string dualities or F-theory constructions. Two-dimensional CFTs with $\mathcal{N} = (0, 4)$ supersymmetry constructed in the literature [59–63] await as well their holographic description. In this context, an important recent development has been the complete classification of AdS₃ solutions to massive IIA supergravity with small $\mathcal{N} = (0, 4)$ supersymmetry (and SU(2) structure) achieved in [64]. In this letter we add a new entry to the dictionary between CFTs and string backgrounds with an AdS-factor by proposing explicit CFTs dual to these solutions. We define our CFTs as the IR fixed points of $\mathcal{N} = (0, 4)$ UV finite two dimensional QFTs. These QFTs are described by quivers, consisting of two long rows of gauge groups connected by hypermultiplets and Fermi multiplets. We show that the new background solutions to massive IIA supergravity constructed in [64] contain the needed isometries to be dual to our CFTs. We give an example (further elaborated in [65], where additional examples can be found) that shows agreement between the field theory and holographic calculations of the central charge. Finally, we provide a formal mapping to the AdS₇ solutions constructed in [22] that suggests the existence of a flow across dimensions [66,67] between the dual CFTs.

*ylozano@uniovi.es

†ntmacpher@gmail.com

‡c.nunez@swansea.ac.uk

§anayelam@gmail.com

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International license](https://creativecommons.org/licenses/by/4.0/). Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

II. THE GEOMETRY

The backgrounds of massive type IIA supergravity constructed in [64] were proposed to be dual to $\mathcal{N} = (0, 4)$ CFTs in two dimensions. These solutions have $SL(2) \times SU(2)$ isometries and eight (four Poincaré plus four conformal) supercharges. In this paper we will consider the particular case of the geometries in [64] referred therein as class I. In string frame they read,

$$\begin{aligned} ds^2 &= g_1(ds^2(\text{AdS}_3) + g_2 ds^2(S^2)) + g_3 ds^2(\text{CY}_2) + \frac{d\rho^2}{g_1}, \\ e^{-\Phi} &= g_4, \quad B_2 = g_5 \text{vol}(S^2), \quad \hat{F}_0 = g_6, \quad \hat{F}_2 = g_7 \text{vol}(S^2), \\ \hat{F}_4 &= g_8 d\rho \wedge \text{vol}(\text{AdS}_3) + g_9 \text{vol}(\text{CY}_2). \end{aligned} \quad (1)$$

The functions g_i are defined in terms of three functions, $u(\rho)$, $\hat{h}_4(\rho)$, $h_8(\rho)$, according to,

$$\begin{aligned} g_1 &= \frac{u}{\sqrt{\hat{h}_4 h_8}}, \quad g_2 = \frac{h_8 \hat{h}_4}{4h_8 \hat{h}_4 + (u')^2}, \quad g_3 = \sqrt{\frac{\hat{h}_4}{h_8}}, \\ g_4 &= \frac{h_8^{\frac{3}{4}}}{2\hat{h}_4^{\frac{1}{4}} \sqrt{u}} \sqrt{4h_8 \hat{h}_4 + (u')^2}, \quad g_6 = h'_8, \quad g_9 = -\partial_\rho \hat{h}_4 \\ g_5 &= \frac{1}{2} \left(-\rho + 2\pi k + \frac{uu'}{4\hat{h}_4 h_8 + (u')^2} \right), \\ g_7 &= -\frac{1}{2} (h_8 - h'_8(\rho - 2\pi k)), \quad g_8 = \left(\partial_\rho \left(\frac{uu'}{2\hat{h}_4} \right) + 2h_8 \right). \end{aligned}$$

Notice that we have written the Page fluxes $\hat{F} = e^{-B_2} \wedge F$. We have also allowed for large gauge transformations $B_2 \rightarrow B_2 + \pi k \text{vol}(S^2)$, with $k = 0, 1, \dots, P$. The transformations are performed every time we cross a ρ -interval $[2\pi k, 2\pi(k+1)]$. The preservation of $\mathcal{N} = (0, 4)$ supersymmetry implies $u'' = 0$. Away from localized sources Bianchi identities also impose $h''_8 = 0$ and $\hat{h}''_4 = 0$ [64].

Below, we present *new* solutions, defined piecewise in the intervals $[2\pi k, 2\pi(k+1)]$. For \hat{h}_4 , h_8 we have,

$$\hat{h}_4(\rho) = \Upsilon \begin{cases} \frac{\beta_0}{2\pi} \rho & 0 \leq \rho \leq 2\pi \\ \alpha_k + \frac{\beta_k}{2\pi} (\rho - 2\pi k), & 2\pi k \leq \rho \leq 2\pi(k+1) \\ \alpha_P - \frac{\alpha_P}{2\pi} (\rho - 2\pi P), & 2\pi P \leq \rho \leq 2\pi(P+1), \end{cases} \quad (2)$$

$$h_8(\rho) = \begin{cases} \frac{\nu_0}{2\pi} \rho & 0 \leq \rho \leq 2\pi \\ \mu_k + \frac{\nu_k}{2\pi} (\rho - 2\pi k) & 2\pi k \leq \rho \leq 2\pi(k+1) \\ \mu_P - \frac{\mu_P}{2\pi} (\rho - 2\pi P) & 2\pi P \leq \rho \leq 2\pi(P+1), \end{cases} \quad (3)$$

while $u(\rho) = \frac{b_0}{2\pi} \rho$.

Imposing the continuity of the Neveu-Schwarz (NS)-sector across the various intervals we find,

$$\mu_k = \sum_{j=0}^{k-1} \nu_j, \quad \alpha_k = \sum_{j=0}^{k-1} \beta_j, \quad (4)$$

which also imply the continuity of the functions \hat{h}_4 , h_8 across intervals. The first derivatives present discontinuities at $\rho = 2\pi k$ where D8 and D4 sources are located.

A. Page charges

The Page charges are important observable quantities characterizing a supergravity solution. They are quantized, and are the ones that are related to the ranks of the gauge or global groups of the dual CFT. They are obtained integrating the Page fluxes, according to $(2\pi)^{7-p} g_s \alpha'^{(7-p)/2} Q_{Dp} = \int_{\Sigma_{8-p}} \hat{F}_{8-p}$.¹ This implies the quantization of some of the constants in Eqs. (2)–(3). In the interval $[2\pi k, 2\pi(k+1)]$ we find,

$$\begin{aligned} Q_{D8} &= 2\pi F_0 = \nu_k, \quad Q_{D6} = \frac{1}{2\pi} \int_{S^2} \hat{F}_2 = \mu_k. \\ Q_{D4} &= \frac{1}{8\pi^3} \int_{\text{CY}_2} \hat{F}_4 = \Upsilon \frac{\text{Vol}(\text{CY}_2)}{16\pi^4} \beta_k, \\ Q_{D2} &= \frac{1}{32\pi^5} \int_{\text{CY}_2 \times S^2} \hat{F}_6 = \Upsilon \frac{\text{Vol}(\text{CY}_2)}{16\pi^4} \alpha_k. \end{aligned} \quad (5)$$

We have used that the magnetic part $\hat{F}_{6,\text{mag}} = \hat{f}_6$ is

$$\hat{f}_6 = \frac{\Upsilon}{2} (h_4 - h'_4(\rho - 2\pi k)) \text{vol}(S^2) \wedge \text{vol}(\text{CY}_2). \quad (6)$$

Besides, we count one NS-five brane every time we cross the value $\rho = 2\pi k$ (for $k = 1, \dots, P$). The total number of NS-five branes is $Q_{\text{NS}} = \frac{1}{4\pi^2} \int_{\rho \times S^2} H_3 = P + 1$.

The study of the Bianchi identities (see [65] for the details), shows that dissolved in flux, we have “color” D2 and D6 branes. We also find that D4 and D8 branes play the role of “flavor,” appearing explicitly as delta-function corrections of the Bianchi identities. For the interval $[2\pi(k-1), 2\pi k]$, we calculate

$$N_{D8}^{[k-1,k]} = \nu_{k-1} - \nu_k, \quad N_{D4}^{[k-1,k]} = \beta_{k-1} - \beta_k, \quad (7)$$

$$N_{D6}^{[k-1,k]} = \mu_k = \sum_{i=0}^{k-1} \nu_i, \quad N_{D2}^{[k-1,k]} = \alpha_k = \sum_{i=0}^{k-1} \beta_i. \quad (8)$$

We then have a Hanany-Witten brane setup [68], that in the interval $[2\pi(k-1), 2\pi k]$ (bounded by NS-five branes), has

¹In what follows, we set $\alpha' = g_s = 1$.

TABLE I. $\frac{1}{8}$ -BPS brane intersection underlying our geometry. (x^0, x^1) are the directions where the 2d CFT (dual to our AdS₃) is defined. (x^2, \dots, x^5) span the CY₂, on which the D6 and the D8-branes are wrapped. x^6 is the direction associated with ρ . Finally (x^7, x^8, x^9) are the transverse directions realizing the SU(2)-symmetry associated to S².

	0	1	2	3	4	5	6	7	8	9
D2	x	x					x			
D4	x	x						x	x	x
D6	x	x	x	x	x	x	x			
D8	x	x	x	x	x	x		x	x	x
NS5	x	x	x	x	x	x				

$N_{D6}^{[k-1,k]}$, $N_{D2}^{[k-1,k]}$ color branes and $N_{D8}^{[k-1,k]}$, $N_{D4}^{[k-1,k]}$ flavor branes (see Table I and Fig. 1).

To close the geometric part of our study, we use the formalism in [69,70] to calculate the holographic central charge. The result is (see [65] for a derivation),

$$c_{\text{hol}} = \frac{3\pi}{2G_N} \text{Vol}(\text{CY}_2) \int_0^{2\pi(P+1)} \hat{h}_4 h_8 d\rho. \quad (9)$$

Since the backgrounds have localized singularities, associated with the presence of D-branes, observables calculated using the geometry are trustable as long as the numbers ν_k, β_k, b_0, P are large.

III. THE FIELD THEORY

In this section we discuss the two-dimensional CFTs dual to the backgrounds given by Eqs. (1)–(3). They are defined as the strongly coupled IR fixed points of QFTs that in the (weakly coupled) UV are constructed from the “building block” depicted in Fig. 2. We have an SU(N) gauge group with the matter content of a two-dimensional $\mathcal{N} = (4, 4)$ vector multiplet in the adjoint of SU(N). This gauge group is joined with other (gauge or global) symmetry groups SU(\hat{P}), SU(R) and SU(Q). The connection with the SU(\hat{P}) symmetry group is mediated by

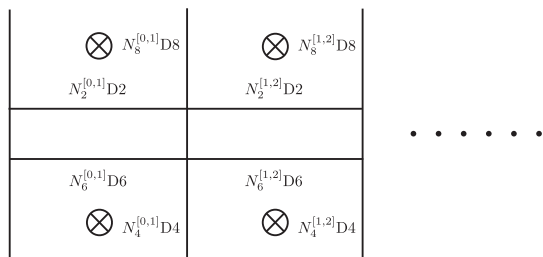


FIG. 1. The generic Hanany-Witten setup associated with our backgrounds. The vertical lines are NS-five branes. The horizontal lines represent D2 and wrapped D6 branes. The crosses indicate D4 and wrapped D8 branes.

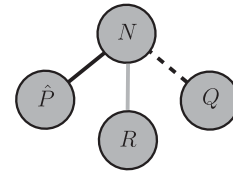


FIG. 2. The building block of our theories.

$\mathcal{N} = (4, 4)$ hypermultiplets running over the black solid line, that with the SU(R) symmetry group via $\mathcal{N} = (0, 4)$ hypermultiplets that propagate over the grey lines, and that with SU(Q) via $\mathcal{N} = (0, 2)$ Fermi multiplets that run over the dashed line.² All these multiplets transform in the bifundamental representation of the gauge groups. A similar (but not the same) field content was used in [61].

The cancellation of gauge anomalies constrains the ranks of the different symmetry groups. Using the contribution to the gauge anomaly coming from each multiplet, see [72], we find that for the SU(N) gauge group the cancellation of the anomaly imposes,

$$2R = Q. \quad (10)$$

Our quiver gauge theories are then obtained by “assembling” the building blocks of Fig. 2 such that there is anomaly cancellation for all gauge groups.

In turn, the central charge of the IR CFT is calculated by associating it with the correlation function of U(1)-R-symmetry currents (computed in the UV-description above). At the conformal point, the (right-moving) central charge is related to the, two point, U(1)_R current correlation function, such that (see [73]),

$$c = 6(n_{\text{hyp}} - n_{\text{vec}}). \quad (11)$$

The central charge is then obtained by counting the number of $\mathcal{N} = (0, 4)$ hypermultiplets and subtracting the number of $\mathcal{N} = (0, 4)$ vector multiplets in the UV description. Note that the SU(2) R-symmetry does not mix with the Abelian flavor symmetries, and it is not necessary to go through a c-extremization procedure [74].

A. The proposed duality

Our proposal relates the backgrounds in Eqs. (1)–(3) with quiver field theories obtained by assembling the building block depicted in Fig. 2. For generic functions \hat{h}_4, h_8 this results in the quiver shown in Fig. 3, associated to the Hanany-Witten setup in Fig. 4. The reader can check that the cancellation of gauge anomalies implies for a generic SU(α_k) color group, in the interval $[2\pi(k - 1), 2\pi k]$

²Notice that $\mathcal{N} = (0, 2)$ Fermi multiplets are allowed in $\mathcal{N} = (0, 4)$ theories as long as they do not transform under the SO(4) R-symmetry (see [71]).

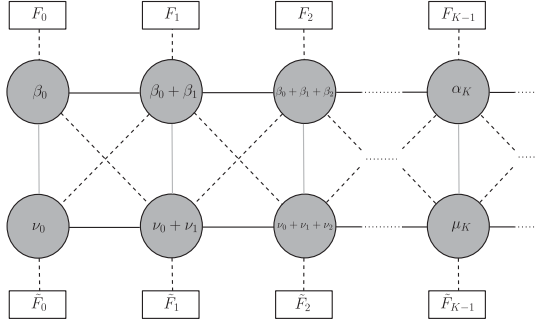


FIG. 3. A generic quiver field theory whose IR is dual to the holographic background defined by Eqs. (1)–(3).

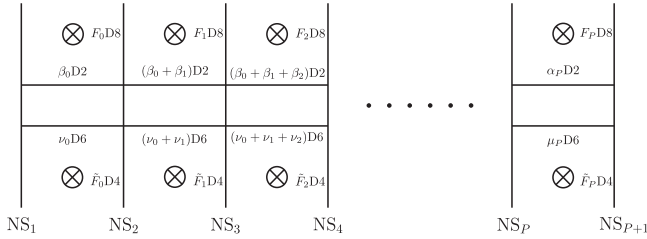


FIG. 4. Hanany-Witten setup associated with our generic quiver in Fig. 3.

$$F_{k-1} + \mu_{k+1} + \mu_{k-1} = 2\mu_k \rightarrow F_{k-1} = \nu_{k-1} - \nu_k, \quad (12)$$

which, according to (7), is precisely the number of flavor D8 branes in the $[2\pi(k-1), 2\pi k]$ interval of the brane setup. Things work analogously if D6 are replaced by D2 (or $\mu_k \leftrightarrow \alpha_k$) and D8 by D4 ($\nu_k \leftrightarrow \beta_k$), and we work with a generic lower-row gauge group $SU(\mu_k)$.

The central charge of the quiver is calculated using expression (11). We find,

$$\frac{c}{6} = \sum_{j=1}^P (\alpha_j \mu_j - \alpha_j^2 - \mu_j^2 + 2) + \sum_{j=1}^{P-1} (\alpha_j \alpha_{j+1} + \mu_j \mu_{j+1}). \quad (13)$$

In [65] we present various examples in which this expression agrees with the holographic central charge computed according to (9). This should hold in the limit in which the number of nodes P and the ranks of each gauge group α_i, μ_i are large, which is when the supergravity backgrounds are trustable. We present one such example below. Notice that both the global symmetries and isometries (space-time, SUSY, and flavor), as well as the ranks of the gauge (colour) groups do match in both descriptions, the latter being given by the numbers α_k, μ_k in (4).

B. An example

Let us discuss an example that illustrates the duality proposed above. We consider the quiver with two rows of

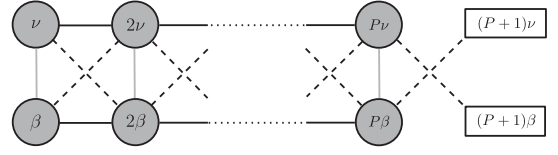


FIG. 5. Quiver consisting of two rows of linearly increasing color groups, terminated with the addition of flavour groups.

linearly increasing color groups depicted in Fig. 5. For an intermediate gauge node $SU(k\nu)$ we have $Q = 2k\beta$, $R = k\beta$. This implies that (10) is satisfied and any generic intermediate gauge group is not anomalous. If we refer to the last gauge group in the upper-row $SU(P\nu)$ we have that $Q = (P+1)\beta + (P-1)\beta = 2P\beta$ and $R = P\beta$. As a consequence (10) is satisfied and the gauge group $SU(P\nu)$ is also not anomalous. The same occurs for the lower-row gauge groups.

The counting of (0,4) hypermultiplets and vector multiplets gives

$$n_{\text{vec}} = \sum_{j=1}^P (j^2(\nu^2 + \beta^2) - 2),$$

$$n_{\text{hyp}} = \sum_{j=1}^{P-1} j(j+1)(\nu^2 + \beta^2) + \sum_{j=1}^P j^2 \nu \beta, \quad (14)$$

from which the central charge of the IR CFT can be computed,

$$c = 6\nu\beta \left(\frac{P^3}{3} + \frac{P^2}{2} + \frac{P}{6} \right) - 3(\nu^2 + \beta^2)(P^2 + P) + 12P$$

$$\sim 2\nu\beta P^3. \quad (15)$$

In turn, the holographic description of the system is in terms of the functions,

$$h_8(\rho) = \begin{cases} \frac{\nu}{2\pi} \rho & 0 \leq \rho \leq 2\pi P \\ \frac{\nu P}{2\pi} (2\pi(P+1) - \rho) & 2\pi P \leq \rho \leq 2\pi(P+1), \end{cases}$$

$$\hat{h}_4(\rho) = \Upsilon \begin{cases} \frac{\beta}{2\pi} \rho & 0 \leq \rho \leq 2\pi P \\ \frac{\beta P}{2\pi} (2\pi(P+1) - \rho) & 2\pi P \leq \rho \leq 2\pi(P+1). \end{cases}$$

Using (9) and a convenient choice for the constant Υ , gives rise to the holographic central charge,

$$c_{\text{hol}} = 2\nu\beta P^3 \left(1 + \frac{1}{P} \right) \sim 2\nu\beta P^3. \quad (16)$$

We thus find perfect agreement between the field theory and holographic calculations. In [65] other examples of dual holographic pairs are discussed that provide stringent support to our proposed duality.

IV. MAPPING TO AdS₇ BACKGROUNDS

A subclass of the solutions discussed in [64] can be related to the AdS₇ solutions to massive IIA supergravity constructed in [22]. As opposed to the mappings in [21], this mapping is not one-to-one, due to the presence of additional D2-D4 branes in the AdS₃ solutions, whose backreaction introduces extra 4-form and 6-form fluxes, and reduces the supersymmetries by a half. Using this map it is possible to give an interpretation to the 2d CFTs dual to the AdS₃ solutions as associated to D2-D4 defects in the D6-NS5-D8 brane setups dual to the AdS₇ solutions in [22], wrapped on the CY₂. Thus, the word defect is here used to indicate the presence of extra branes in Hanany-Witten brane set-ups that would otherwise arise from compactifying higher dimensional branes.

The explicit map, discussed in detail in [75], reads,

$$\rho \leftrightarrow 2\pi z, \quad u \leftrightarrow \alpha, \quad h_8 \leftrightarrow -\frac{\ddot{\alpha}}{81\pi^2}, \quad \hat{h}_4 \leftrightarrow \frac{81}{8}\alpha,$$

$$ds^2(\text{AdS}_3) + \frac{3^4}{2^3} ds^2(\text{CY}_2) \leftrightarrow 4ds^2(\text{AdS}_7). \quad (17)$$

This transforms the original backgrounds in (1) into the AdS₇ backgrounds constructed in [22]

$$\frac{ds_{10}^2}{\pi\sqrt{2}} = 8\sqrt{-\frac{\alpha}{\ddot{\alpha}}} ds^2(\text{AdS}_7) + \sqrt{-\frac{\ddot{\alpha}}{\alpha}} dz^2 + \frac{\alpha^{3/2}(-\ddot{\alpha})^{1/2}}{\Delta} ds^2(\text{S}^2),$$

$$e^{2\Phi} = 2^{5/2}\pi^5 3^8 \frac{(-\alpha/\ddot{\alpha})^{3/2}}{\dot{\alpha}^2 - 2\alpha\ddot{\alpha}}, \quad F_0 = -\frac{\ddot{\alpha}}{162\pi^3}$$

$$B_2 = \pi \left(-z + k + \frac{\alpha\dot{\alpha}}{\Delta} \right) \text{vol}(\text{S}^2),$$

$$\hat{F}_2 = \frac{1}{162\pi^2} (\dot{\alpha} - \ddot{\alpha}(z-k)) \text{vol}(\text{S}^2), \quad \Delta = \dot{\alpha}^2 - 2\alpha\ddot{\alpha}, \quad (18)$$

where the function $\alpha(z)$ satisfies the equation $\ddot{\alpha} = -162\pi^3 F_0$. As analyzed in [22], $\alpha(z)$ encodes the information about the 6d (1,0) dual CFT, which is realized in a D6-NS5-D8 Hanany-Witten setup. Using the mapping defined by (17) it is possible to obtain an AdS₇ solution in the class of [22] from an AdS₃ solution. In turn, the D6-NS5-D8 *sector* of the AdS₃ solution is obtained by compactifying the D6-NS5-D8 branes that underlie the AdS₇ solution on the CY₂, while it is necessary to add the D2-D4 *sector*, encoded by the functions u and h_4 (see [75] for the details) to achieve conformality and fully determine the AdS₃ solution.

The holographic central charge of the 6d CFTs dual to the AdS₇ solutions was computed in [25],

$$c_{\text{AdS}_7} = \frac{1}{G_N} \frac{2^4}{3^8} \int dz(-\alpha\ddot{\alpha}). \quad (19)$$

In turn, the holographic central charge of the 2d CFTs is given in (9). Using the mapping given by (17) this becomes,

$$c_{\text{AdS}_3} \leftrightarrow \frac{3}{2^3 G_N} \text{Vol}(\text{CY}_2) \int dz(-\alpha\ddot{\alpha}) = \frac{3^9}{2^7} \text{Vol}(\text{CY}_2) c_{\text{AdS}_7}.$$

This kind of relation is ubiquitous when calculating the holographic central charges for “flows across dimensions.” Our result strongly suggests that we can obtain our CFTs by compactifying the D6-NS5-D8 system underlying the 6d (1,0) CFT on a CY₂. Conformality in the lower dimensional theory however requires the presence of “defect” D2 and D4 branes, represented by the fluxes F_4, F_6 in (1). Flows of this type were studied in [76,77], but these do not reach an AdS₃ fixed point. It would be interesting to find the explicit RG flows that deform the six-dimensional $\mathcal{N} = (1, 0)$ CFT to reach a two-dimensional $\mathcal{N} = (0, 4)$ conformal fixed point in the IR.

V. CONCLUSIONS

This paper presents a new entry in the mapping between CFTs and AdS-supergravity backgrounds, for the case of two-dimensional (small) $\mathcal{N} = (0, 4)$ CFTs and backgrounds with AdS₃ × S² factors. We have reported new solutions of the type AdS₃ × S² × CY₂, belonging to class I in the classification in [64], with compact CY₂ and piecewise continuous defining functions. We have proposed explicit 2-d dual CFTs based on the Hanany-Witten setups implied by the Page charges of the solutions. We matched the background isometries and the global symmetries (both space-time and flavor) of the CFTs, and checked the agreement between the holographic and field theory central charges. The CFTs are defined as the IR limit of UV finite long quivers with (0,4) SUSY, that generalize 2-d (0,4) quivers previously discussed in the literature [61,63]. We presented a map between a subclass of the solutions in [64] and the AdS₇ backgrounds dual to six dimensional $\mathcal{N} = (1, 0)$ CFTs [22–24]. This mapping suggests the possibility of finding an RG flow across dimensions between the dual CFTs.

This paper just scratches the surface of a rich line of work. In the forthcoming papers [65,75] we will present various checks of our proposed duality. As a by-product we will obtain explicit completions of the background obtained via non-Abelian T-duality on AdS₃ × S³ × CY₂, along the lines of [9,18,33,78,79].

ACKNOWLEDGMENTS

We would like to thank Giuseppe Dibitetto, Gaston Giribet, S. Prem Kumar, Daniel Thompson, Alessandro Tomasiello and Stefan Vandoren for very useful discussions. Y.L. and A.R. are partially supported by the Spanish government Grant No. PGC2018-096894-B-100 and by the Principado de Asturias Grant No. FC-GRUPIN-IDI/2018/000174. N. T. M. is funded by the Italian Ministry

of Education, Universities and Research under the Prin project “Non Perturbative Aspects of Gauge Theories and Strings” (2015MP2CX4) and INFN. C.N. is a Wolfson Fellow of the Royal Society. A.R. is supported by CONACyT-Mexico. We would like to acknowledge

the Mainz Institute for Theoretical Physics (MITP) of the DFG Cluster of Excellence PRISMA⁺ (Project ID 39083149) and the Theory Unit at CERN for their hospitality and partial support during the development of this work.

-
- [1] J. M. Maldacena, *Int. J. Theor. Phys.* **38**, 1113 (1999); *Adv. Theor. Math. Phys.* **2**, 231 (1998).
- [2] J. P. Gauntlett, D. Martelli, J. Sparks, and D. Waldram, *Classical Quantum Gravity* **21**, 4335 (2004).
- [3] J. B. Gutowski and G. Papadopoulos, *J. High Energy Phys.* **02** (2015) 145.
- [4] E. Witten, *Nucl. Phys.* **B500**, 3 (1997).
- [5] D. Gaiotto, *J. High Energy Phys.* **08** (2012) 034.
- [6] D. Gaiotto and J. Maldacena, *J. High Energy Phys.* **10** (2012) 189.
- [7] R. A. Reid-Edwards and B. Stefanski, Jr., *Nucl. Phys.* **B849**, 549 (2011).
- [8] O. Aharony, L. Berdichevsky, and M. Berkooz, *J. High Energy Phys.* **08** (2012) 131.
- [9] Y. Lozano and C. Nunez, *J. High Energy Phys.* **05** (2016) 107.
- [10] I. Bah, *J. High Energy Phys.* **09** (2015) 163.
- [11] C. Nunez, D. Roychowdhury, and D. C. Thompson, *J. High Energy Phys.* **07** (2018) 044.
- [12] C. Nunez, D. Roychowdhury, S. Speziali, and S. Zacarias, *Nucl. Phys.* **B943**, 114617 (2019).
- [13] I. Bah, F. Bonetti, R. Minasian, and E. Nardoni, *J. High Energy Phys.* **06** (2019) 123.
- [14] E. D’Hoker, M. Gutperle, A. Karch, and C. F. Uhlemann, *J. High Energy Phys.* **08** (2016) 046.
- [15] E. D’Hoker, M. Gutperle, and C. F. Uhlemann, *Phys. Rev. Lett.* **118**, 101601 (2017).
- [16] E. D’Hoker, M. Gutperle, and C. F. Uhlemann, *J. High Energy Phys.* **05** (2017) 131.
- [17] M. Gutperle, A. Trivella, and C. F. Uhlemann, *J. High Energy Phys.* **04** (2018) 135.
- [18] Y. Lozano, N. T. Macpherson, and J. Montero, *J. High Energy Phys.* **01** (2019) 116.
- [19] M. Fluder and C. F. Uhlemann, *Phys. Rev. Lett.* **121**, 171603 (2018).
- [20] O. Bergman, D. Rodriguez-Gomez, and C. F. Uhlemann, *J. High Energy Phys.* **08** (2018) 127.
- [21] F. Apruzzi, M. Fazzi, A. Passias, A. Rota, and A. Tomasiello, *Phys. Rev. Lett.* **115**, 061601 (2015).
- [22] F. Apruzzi, M. Fazzi, D. Rosa, and A. Tomasiello, *J. High Energy Phys.* **04** (2014) 064.
- [23] D. Gaiotto and A. Tomasiello, *J. High Energy Phys.* **12** (2014) 003.
- [24] S. Cremonesi and A. Tomasiello, *J. High Energy Phys.* **05** (2016) 031.
- [25] C. Nunez, J. M. Penin, D. Roychowdhury, and J. Van Gersel, *J. High Energy Phys.* **06** (2018) 078.
- [26] K. Filippas, C. Nunez, and J. Van Gersel, *J. High Energy Phys.* **06** (2019) 069.
- [27] I. Brunner and A. Karch, *J. High Energy Phys.* **03** (1998) 003.
- [28] A. Hanany and A. Zaffaroni, *Nucl. Phys.* **B529**, 180 (1998).
- [29] D. Gaiotto and E. Witten, *Adv. Theor. Math. Phys.* **13**, 721 (2009).
- [30] E. D’Hoker, J. Estes, and M. Gutperle, *J. High Energy Phys.* **06** (2007) 022.
- [31] E. D’Hoker, J. Estes, M. Gutperle, and D. Krym, *J. High Energy Phys.* **08** (2008) 028.
- [32] B. Assel, C. Bachas, J. Estes, and J. Gomis, *J. High Energy Phys.* **08** (2011) 087.
- [33] Y. Lozano, N. T. Macpherson, J. Montero, and C. Nunez, *J. High Energy Phys.* **11** (2016) 133.
- [34] J. M. Maldacena, A. Strominger, and E. Witten, *J. High Energy Phys.* **12** (1997) 002.
- [35] C. Vafa, *Adv. Theor. Math. Phys.* **2**, 207 (1998).
- [36] R. Minasian, G. W. Moore, and D. Tsimpis, *Commun. Math. Phys.* **209**, 325 (2000).
- [37] A. Castro, J. L. Davis, P. Kraus, and F. Larsen, *Int. J. Mod. Phys. A* **23**, 613 (2008).
- [38] B. Haghighat, S. Murthy, C. Vafa, and S. Vandoren, *J. High Energy Phys.* **01** (2016) 009.
- [39] C. Couzens, H. H. Lam, K. Mayer, and S. Vandoren, *J. High Energy Phys.* **08** (2019) 043.
- [40] E. Witten, *J. High Energy Phys.* **07** (1997) 003.
- [41] D. Kutasov, F. Larsen, and R. G. Leigh, *Nucl. Phys.* **B550**, 183 (1999).
- [42] N. Seiberg and E. Witten, *J. High Energy Phys.* **04** (1999) 017.
- [43] F. Larsen and E. J. Martinec, *J. High Energy Phys.* **11** (1999) 002.
- [44] N. Kim, *J. High Energy Phys.* **01** (2006) 094.
- [45] J. P. Gauntlett, O. A. P. Mac Conamhna, T. Mateos, and D. Waldram, *Phys. Rev. Lett.* **97**, 171601 (2006).
- [46] J. P. Gauntlett, N. Kim, and D. Waldram, *J. High Energy Phys.* **04** (2007) 005.
- [47] N. Kim, *Phys. Rev. D* **86**, 067901 (2012).
- [48] H. Kim, K. K. Kim, and N. Kim, *J. High Energy Phys.* **08** (2007) 050.
- [49] E. O. Colgain, J. B. Wu, and H. Yavartanoo, *J. High Energy Phys.* **08** (2010) 114.
- [50] D. Tong, *J. High Energy Phys.* **04** (2014) 193.
- [51] Y. Lozano, N. T. Macpherson, J. Montero, and E. O. Colgain, *J. High Energy Phys.* **08** (2015) 121.

- [52] O. Kelecki, Y. Lozano, J. Montero, E. O. Colgain, and M. Park, *Phys. Rev. D* **93**, 086010 (2016).
- [53] L. Eberhardt, *J. High Energy Phys.* **02** (2018) 087.
- [54] C. Couzens, C. Lawrie, D. Martelli, S. Schafer-Nameki, and J. M. Wong, *J. High Energy Phys.* **08** (2017) 043.
- [55] C. Couzens, D. Martelli, and S. Schafer-Nameki, *J. High Energy Phys.* **06** (2018) 008.
- [56] G. Dibitetto, G. Lo Monaco, A. Passias, N. Petri, and A. Tomasiello, *Fortschr. Phys.* **66**, 1800060 (2018).
- [57] N. T. Macpherson, *J. High Energy Phys.* **05** (2019) 089.
- [58] N. S. Deger, C. Eloy, and H. Samtleben, *J. High Energy Phys.* **10** (2019) 145.
- [59] B. Haghighat, A. Iqbal, C. Kozcaz, G. Lockhart, and C. Vafa, *Commun. Math. Phys.* **334**, 779 (2015).
- [60] B. Haghighat, C. Kozcaz, G. Lockhart, and C. Vafa, *Phys. Rev. D* **89**, 046003 (2014).
- [61] A. Gadde, B. Haghighat, J. Kim, S. Kim, G. Lockhart, and C. Vafa, *J. High Energy Phys.* **02** (2018) 143.
- [62] C. Lawrie, S. Schafer-Nameki, and T. Weigand, *J. High Energy Phys.* **04** (2017) 111.
- [63] A. Hanany and T. Okazaki, *J. High Energy Phys.* **03** (2019) 027.
- [64] Y. Lozano, N. T. Macpherson, C. Nunez, and A. Ramirez, [arXiv:1908.09851](https://arxiv.org/abs/1908.09851).
- [65] Y. Lozano, N. T. Macpherson, C. Nunez, and A. Ramirez, [arXiv:1909.10510](https://arxiv.org/abs/1909.10510).
- [66] A. Donos and J. P. Gauntlett, *J. High Energy Phys.* **08** (2014) 006.
- [67] F. Benini, N. Bobev, and P. M. Cricigno, *J. High Energy Phys.* **07** (2016) 020.
- [68] A. Hanany and E. Witten, *Nucl. Phys.* **B492**, 152 (1997).
- [69] N. T. Macpherson, C. Nunez, L. A. P. Zayas, V. G. J. Rodgers, and C. A. Whiting, *J. High Energy Phys.* **02** (2015) 040.
- [70] Y. Bea, J. D. Edelstein, G. Itsios, K. S. Kooner, C. Nunez, D. Schofield, and J. A. Sierra-Garcia, *J. High Energy Phys.* **05** (2015) 062.
- [71] E. Witten, *J. Geom. Phys.* **15**, 215 (1995).
- [72] S. Franco, D. Ghim, S. Lee, R. K. Seong, and D. Yokoyama, *J. High Energy Phys.* **09** (2015) 072.
- [73] P. Putrov, J. Song, and W. Yan, *J. High Energy Phys.* **03** (2016) 185.
- [74] F. Benini and N. Bobev, *J. High Energy Phys.* **06** (2013) 005.
- [75] Y. Lozano, N. T. Macpherson, C. Nunez, and A. Ramirez, *J. High Energy Phys.* **12** (2019) 013.
- [76] G. Dibitetto and N. Petri, *J. High Energy Phys.* **12** (2017) 041.
- [77] G. Dibitetto and N. Petri, *J. High Energy Phys.* **01** (2018) 039.
- [78] Y. Lozano, C. Nunez, and S. Zacarias, *J. High Energy Phys.* **09** (2017) 008.
- [79] G. Itsios, Y. Lozano, J. Montero, and C. Nunez, *J. High Energy Phys.* **09** (2017) 038.