Holography for the Ishibashi-Kawai-Kitazawa-Tsuchiya Matrix Model

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Holographic dualities that relate type II strings on near-horizon Dp brane geometries to super Yang-Mills theories with sixteen supercharges in $p+1$ dimensions provide nonconformal generalizations of the famous AdS/CFT correspondence. For the extremal case $p=-1$, this suggests a holographic duality for the Ishibashi-Kawai-Kitazawa-Tsuchiya (IKKT) matrix model—super Yang-Mills theory reduced to zero dimensions. Despite intriguing and highly nontrivial results in the IKKT model, this duality has remained largely unexplored so far. In this Letter, we consider the lowest supermultiplet of gauge invariant operators of the model and identify its states with the lowest Kaluza-Klein fluctuations of (Euclidean) IIB supergravity on the D(-1) instanton background. We construct its holographic bulk realization as a one-dimensional maximal supergravity with 32 supercharges and local SO(10) invariance, capturing the full nonlinear dynamics. Analyzing the bulk Killing spinor equations, we construct the most general half-supersymmetric solutions, which typically break SO(10). We present their uplifts to IIB supergravity, and furthermore to pp waves in twelve dimensions. These results provide the minimal setup for conducting precision tests of holography involving Einstein gravity.

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A particularly interesting corner of the holographic dualities is the correspondence between the backgrounds of Dp-branes and their nonconformal field theory duals. Unlike the cases of conformal field theory duals, which are captured by the celebrated AdS/CFT correspondence, nonconformal dualities open a window into a less symmetric dynamics characterized by scale-dependent behavior. These dualities preserve sixteen supercharges and have been studied since the early days of holography [1,2]. In this context, the dual field theory is the reduction of ten-dimensional super Yang-Mills theory down to $p+1$ dimensions.

Extrapolating this correspondence to the extremal case of $p=-1$ suggests a holographic duality between strings on the D(-1) instanton background and super Yang-Mills theory reduced to zero dimensions. The latter is known as the Ishibashi-Kawai-Kitazawa-Tsuchiya (IKKT) matrix model, and has by itself been conjectured as a candidate for a nonperturbative formulation of type IIB superstring theory [3]. The dual D(-1) instanton background is a half-supersymmetric solution of Euclidean IIB supergravity with flat spacetime metric and nontrivial dilaton and axion [4–7]. Despite the compelling and highly nontrivial results that have been obtained in the IKKT model [8–11], this holographic duality has remained largely unstudied, not at least due to various subtleties and degeneracies on both sides of the correspondence. However, recent work has started to probe holography in the polarized IKKT matrix model [12–14], a supersymmetric $\mathrm{SO}(3)\times\mathrm{SO}(7)$ invariant mass deformation of the original model [15].

In this Letter, we explore the correspondence between gauge invariant operators of the IKKT model and Kaluza-Klein fluctuations around the D(-1) instanton background. In particular, we consider the lowest BPS multiplet of such operators and identify its states with the lowest Kaluza-Klein fluctuations of IIB supergravity. We argue that their full nonlinear dynamics is described by a one-dimensional maximal supergravity with 32 supercharges and local SO(10) invariance, which we construct explicitly. For a subsector of the theory, we present the explicit nonlinear embedding into the Euclidean IIB theory. This paves the way for the holographic computation of correlation functions along the lines of [16].

We also study the Killing spinor equations of the bulk theory and determine the general family of half-supersymmetric solutions in terms of ten real constants. We provide their uplift to the IIB theory and recover the D(-1) instanton as a special case. All these solutions turn out to have a flat spacetime metric in ten dimensions and can be further uplifted to twelve-dimensional pp waves. We discuss their significance for the holographic description of supersymmetric deformations of the IKKT model.

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Let us start from the Euclidean IKKT matrix model, obtained as the reduction of ten-dimensional super Yang-Mills theory to zero dimensions, with the action given by

$$S_{\mathrm{IKKT}} = -\mathrm{Tr}\left[\frac{1}{4}[X_a, X_b][X^a, X^b] - \frac{1}{2}\bar{\Psi}\Gamma^a[X_a, \Psi]\right]. \quad (1)$$

Here, X_a and Ψ^α are bosonic and fermionic SU(N) matrices, respectively, transforming in the **10** and the **16** of SO(10), and Γ^a denotes the SO(10) gamma matrices. To maintain clarity in the equations, we suppress spinor indices and we use $\bar{\Psi}_\alpha \coloneqq \Psi^\beta \mathcal{C}_{\alpha\beta}$ to represent the charge-conjugated spinor, where $\mathcal{C}$ is the symmetric charge conjugation matrix [17]. Throughout the Letter, we will use spinors with 32 components. However, in the IKKT model, the spinor Ψ^α and the supersymmetry parameter satisfy the chirality condition $\Gamma_* \Psi = \Psi$. The action (1) is complex, but this is a standard feature of spinorial actions in Euclidean signature [9].

The IKKT action is invariant under SU(N), SO(10), and supersymmetry transformations [18]. The latter act as

$$\delta_\epsilon X^a = \bar{\epsilon}\Gamma^a \Psi, \qquad \delta_\epsilon \Psi = \frac{1}{2}\Gamma^{ab}\epsilon[X_a, X_b]. \quad (2)$$

The supersymmetry algebra closes on shell into

$$[\delta_{\epsilon_1}, \delta_{\epsilon_2}] = \delta_\lambda^{\mathrm{SU}(N)}, \quad (3)$$

with the SU(N) parameter $\lambda = 2\bar{\epsilon}_2 \Gamma^b \epsilon_1 X_b$. As a consequence, supersymmetry transformations commute on the set of SU(N) invariant single-trace operators

$$\mathcal{O} = \mathrm{Tr}[XX\ldots\Psi\ldots X\ldots\Psi\ldots]. \quad (4)$$

The set of all such operators, after dividing out the field equations of (1), organizes into supermultiplets [19]. These are all of the long type

$$\mathcal{R} \otimes (\mathbf{1} \oplus \mathbf{16} \oplus \mathbf{16} \wedge \mathbf{16} \oplus \mathbf{16} \wedge \mathbf{16} \wedge \mathbf{16} \oplus \ldots), \quad (5)$$

with some SO(10) representation $\mathcal{R}$, except for a tower of protected BPS multiplets

$$\sum_{n=2}^{\infty} \mathcal{B}_n, \quad (6)$$

which are built from a lowest weight state

$$\mathcal{O}^{a_1\ldots a_n} = \mathrm{Tr}\left[X^{((a_1}X^{a_2}\ldots X^{a_n))}\right], \quad (7)$$

where the double brackets $((\ldots))$ refer to traceless symmetrization of indices. These multiplets combine the SO(10) representations

$$\begin{aligned}\mathcal{B}_n = {} & [n, 0000]_n \oplus [n-1, 0001]_{n+\frac{1}{2}} \oplus [n-2, 0100]_{n+1} \\ & \oplus [n-3, 1010]_{n+\frac{3}{2}} \oplus [n-3, 0020]_{n+2} \\ & \oplus [n-4, 2000]_{n+2} \oplus [n-4, 1010]_{n+\frac{5}{2}} \\ & \oplus [n-4, 0100]_{n+3} \oplus [n-4, 0001]_{n+\frac{7}{2}} \\ & \oplus [n-4, 0000]_{n+4}, \end{aligned} \quad (8)$$

with the subscript counting the naive dimension Δ, corresponding to $\Delta_X = 1$, $\Delta_\Psi = \frac{3}{2}$. For generic n, the BPS multiplet $\mathcal{B}_n$ carries

$$\mathrm{No.}_{\mathrm{bos}} = 1 + 128 \dim \mathcal{R}_{[n-2,0,0,0]} = 1 + \mathrm{No.}_{\mathrm{ferm}}, \quad (9)$$

degrees of freedom. For the lowest BPS multiplet $\mathcal{B}_2$, the generic structure (8) degenerates into

$$\mathcal{B}_2 = \mathbf{54}_{+2} \oplus \mathbf{144}_{+\frac{5}{2}} \oplus \mathbf{120}_{+3} \ominus \mathbf{45}_{+4} \ominus \mathbf{16}_{+\frac{9}{2}}, \quad (10)$$

consistent with the counting (9). While in higher dimensions states of negative multiplicity in the lowest multiplet correspond to the Goldstone modes to be subtracted from a massive state, here they account for the global SO(10) and supersymmetry, respectively. The first three terms in (10) correspond to the operators

$$\mathcal{O}^{ab} = \mathrm{Tr}[X^a X^b] - \frac{1}{10}\delta^{ab}\mathrm{Tr}[X^c X^c], \quad (11)$$

$$\mathcal{O}^a = \mathrm{Tr}[X^a \Psi] - \frac{1}{9}\mathrm{Tr}[X_b \Gamma^{ab}\Psi], \quad (12)$$

$$\mathcal{O}^{abc} = \mathrm{Tr}\left[X^a[X^b, X^c]\right] - \frac{1}{8}\mathrm{Tr}[\bar{\Psi}\Gamma^{abc}\Psi], \quad (13)$$

on which the supersymmetry algebra (3) closes with

$$\delta_\epsilon \mathcal{O}^{ab} = \frac{9}{5}\bar{\epsilon}\Gamma^{(a}\mathcal{O}^{b)}, \quad (14)$$

$$\delta_\epsilon \mathcal{O}^a = \frac{1}{18}\left(7\Gamma_{bc}\epsilon\mathcal{O}^{abc} - \Gamma^{abcd}\epsilon\mathcal{O}_{bcd}\right), \quad (15)$$

and $\delta_\epsilon \mathcal{O}^{abc} = 0$. As a consequence of the nilpotent structure of the supersymmetry transformations, the vev deformations of the IKKT model (1) by components of $\mathcal{O}^{ab}$ preserve all 16 supercharges.

In the higher-dimensional holographic dualities, it can be shown that the BPS part of the SYM spectrum in $p+1$ dimensions precisely coincides with the full spectrum of supergravity fluctuations around the Dp-brane near horizon background [19]. In particular, in the known higher-dimensional cases, the lowest BPS multiplet in the Yang-Mills spectrum is realized as the lowest Kaluza-Klein supergravity multiplet in the bulk. In turn, its dynamics is described by a maximal gauged supergravity in $(p+2)$

dimensions, with gauge group $\mathrm{SO}(9-p)$, obtained by reduction from ten dimensions on S^{8-p} [2]. Accordingly, for IKKT, the associated bulk theory should correspond to a one-dimensional SO(10)-gauged supergravity combining the lowest Kaluza-Klein fluctuations (10) around the D(-1)-instanton background. In the following, we construct this model explicitly.

We begin by presenting the dual half-supersymmetric instanton solution of Euclidean IIB supergravity [4,5]. The metric in the Einstein frame, the axion χ, and the dilaton Φ are given by

$$ds_{10}^2 = dr^2 + r^2 d\Omega_9^2, \qquad e^{\Phi} = H = -\chi^{-1}, \quad (16)$$

while all the other p-forms vanish. The harmonic function H reads

$$H = h + \frac{Q}{r^8}. \quad (17)$$

Here, $d\Omega_9$ is the line element on the unit 9-sphere, such that the metric is flat and that the solution has an SO(10) isometry. The constant Q is proportional to the instanton charge, while the constant h can be shifted by the $\mathrm{SL}(2,\mathbb{R})$ symmetry of the IIB theory [6]. Note that the sphere radius r, which plays the role of Euclidean time, serves as the holographic coordinate. The string frame metric describes a wormhole geometry, but becomes flat in the near-horizon limit [7]. This contrasts with Dp-brane near-horizon geometries for $p \geq 0$ which correspond to domain wall solutions [2].

To construct the maximally supersymmetric theory that governs the lowest Kaluza-Klein fluctuations around this background, we will first construct a (nonsupersymmetric) consistent truncation of IIB Euclidean supergravity on the sphere S^9 as a special case of a more general consistent S^d truncation of $(d+1)$-dimensional dilaton-axion gravity. We then extend the resulting one-dimensional theory for $d=9$ to include the full field content of the supermultiplet (10) by imposing maximal supersymmetry.

The starting point of our construction is the $(d+1)$-dimensional bosonic Lagrangian

$$\mathcal{L}_{d+1} = |E|\left(R - \frac{1}{2}\partial_\mu \Phi \partial^\mu \Phi + \frac{1}{2} e^{2\Phi} \partial_\mu \chi \partial^\mu \chi\right), \quad (18)$$

in Euclidean spacetime. We keep the dimension arbitrary for now, but later focus on the $d=9$ case, where the Lagrangian (18) describes the axion and dilaton sector of Euclidean IIB supergravity [20]. Note that all fields in this action are real [21], but in Euclidean signature the sign of the axion kinetic term has changed [4]. Consequently, the coset parametrized by Φ and χ is $\mathrm{SL}(2)/\mathrm{SO}(1,1)$. This relative sign is essential for ensuring that the energy momentum tensor of the flat space solution (16) vanishes.

When expanded around the sphere S^d, the $(d+1)$-dimensional theory (18) admits a consistent truncation to a finite set of Kaluza-Klein modes. The nonlinear dynamics of these modes are governed by a one-dimensional theory, which is outlined below. The Kaluza-Klein truncation *Ansatz* for the $(d+1)$-dimensional fields is inspired by the *Ansätze* of [23,24] and reads

$$ds_D^2 = e^{\frac{8d\phi}{d-1}} \Delta \mathrm{e}^2\, dt^2 + g^{-2} e^{\frac{8\phi}{d-1}} T_{ij}^{-1} D\mu^i D\mu^j, \quad (19)$$

$$e^{\Phi} = e^{-4\phi} \Delta^{-1}, \quad (20)$$

$$\chi = -\frac{1}{2\mathrm{e} g}\left(T_{ij}^{-1} D_t T_{kj} \mu^i \mu^k + \frac{8\dot{\phi}}{d-1}\right), \quad (21)$$

where the embedding coordinates of the unit S^d-sphere μ^i, with $i = 1, \ldots, d+1$, satisfy $\mu^i \mu^i = 1$, and where $\Delta = T_{ij}\mu^i\mu^j$, The symmetric, positive definite matrix $T_{ij} \in \mathrm{SL}(d+1)$ contains the lower-dimensional scalar fields that only depend on the Euclidean time coordinate t. The other one-dimensional fields include the Einbein e, a dilaton ϕ, and a set of $\mathrm{SO}(d+1)$ gauge fields $A_{[ij]}$. The latter appear implicitly through the gauge covariant differential D, which acts as follows

$$D\mu^i = d\mu^i - g A_{ji}\, \mu^j dt, \quad (22)$$

$$D_t T_{ij} = \dot{T}_{ij} - 2g A_{k(i} T_{j)k}, \quad (23)$$

with $DT_{ij} =: D_t T_{ij} dt$ and where g is the gauge coupling constant. The dot represents the derivative with respect to the time t. The local $\mathrm{SO}(d+1)$ invariance of the resulting one-dimensional gravity theory is inherited from the isometry group of the sphere S^d. Its Lagrangian reads

$$\mathcal{L}_1 = \frac{8(d+1)}{(d-1)} \mathrm{e}^{-1} \dot{\phi}^2 + \frac{1}{4} \mathrm{e}^{-1} D_t T_{ij}^{-1} D_t T_{ij} - \frac{1}{2} \mathrm{e} g^2 e^{8\phi} (2 T_{ij} T_{ij} - (T_{ii})^2). \quad (24)$$

The scalar potential appearing in the second line is characteristic of sphere truncations of gravity theories. To explicitly prove the consistency of this Kaluza-Klein truncation, it is necessary to show that, after substitution of the truncation *Ansatz*, the $(d+1)$-dimensional field equations reduce to those derived from the one-dimensional Lagrangian (24). We have checked that this is indeed the case and the details of the computation will be presented elsewhere [25].

The field equations deriving from (24) can be conveniently expressed in an $\mathrm{SO}(d+1)$ gauge where $T_{ij} = \delta_{ij} e^{\varphi_j}$, with $\sum_i \varphi_i = 0$. Additionally, we exploit the time reparametrization invariance of the theory to set $\mathrm{e} = 1$. Varying with respect to the Einbein, the dilaton, and the scalars φ_i,

respectively, leads to

$$\frac{16(d+1)}{d-1}\dot{\phi}^2 = \frac{1}{2}\sum_k \dot{\varphi}_k^2 - g^2 e^{8\phi} V_0, \tag{25}$$

$$\frac{4(d+1)}{d-1}\ddot{\phi} = -g^2 e^{8\phi}\, V_0, \tag{26}$$

$$\frac{1}{2}\ddot{\varphi}_i - \frac{4}{d-1}\ddot{\phi} = g^2\, e^{8\phi}\left(2e^{2\varphi_i} - \sum_k e^{\varphi_i+\varphi_k}\right), \tag{27}$$

with the potential

$$V_0 = 2\sum_k e^{2\varphi_k} - \left(\sum_k e^{\varphi_k}\right)^2. \tag{28}$$

The gauge fields A_{ij} are set to zero by their own field equations, or decouple from the dynamics. For vanishing φ_i, the system admits an $\mathrm{SO}(d+1)$-invariant solution, which for $d=9$ uplifts to the D-instanton (16) via the truncation *Ansatz* (19)–(21). This solution actually belongs to a broader class of half-supersymmetric solutions, which we will derive below.

As discussed above, we expect the holographic duality to relate the IKKT operators (11)–(13) to the lowest Kaluza-Klein fluctuations of the full IIB supergravity, corresponding to the multiplet (10). In principle, one should thus extend the S^9 truncation *Ansatz* (19)–(21) to include all the ten-dimensional p-forms. This would likely require elaborate techniques based on exceptional geometry that are beyond the scope of this Letter. We instead opt to directly construct the maximally supersymmetric extension of the one-dimensional Lagrangian (24) for $d=9$.

Because of the triviality of the tangent space in one dimension, the fermionic fields of the resulting supergravity theory simply consist of Grassman variables transforming in different SO(10) representations. They include the gravitino ψ^α, as well as the fermions λ^α and χ_a^α. We employ here the SO(10) spinor notations introduced in (1). Note, however, that unlike the chiral spinor Ψ in the IKKT model, a supergravity spinor such as the gravitino ψ^α (or the supersymmetry parameter ϵ^α) carries 32 independent components. Likewise, the vector-spinor χ_a^α satisfies the trace condition $\Gamma^a\chi_a = 0$ and carries 2×144 components. This indeed is the proper off-shell bulk content to realize the on-shell multiplet (10).

In addition to the fermionic fields, maximal supersymmetry also necessitates enlarging the bosonic sector by adding 120 axion fields $a_{[ijk]}$, according to (10). Up to second order in the axions and in the fermion fields, the resulting supergravity Lagrangian takes the form [26]

$$\begin{aligned}
\mathcal{L}_1 = {}& 10\mathrm{e}^{-1}\dot{\phi}^2 + \frac{1}{4}\mathrm{e}^{-1}D_t T_{ij}^{-1} D_t T_{ij} - \frac{1}{12}\mathrm{e}^{-1}e^{-2\phi}T_{ij}^{-1}T_{kl}^{-1}T_{mn}^{-1}D_t a_{ikm}D_t a_{jln} + 20\bar{\lambda}\mathcal{D}_t\lambda + 2\bar{\chi}^a\mathcal{D}_t\chi_a \\
& - \frac{1}{2}\mathrm{e}g^2 e^{8\phi}\left(2T_{ij}T_{ij} - (T_{ii})^2 + \frac{1}{2}e^{30\phi}\left(a_{ijk}a_{lmn}T_{il}T_{jm}^{-1}T_{kn}^{-1} - 2a_{ijk}a_{ijl}T_{kl}^{-1}\right)\right) \\
& + 4g\mathcal{X}\bar{\psi}\lambda - 2g\mathcal{X}^{ab}\bar{\psi}\Gamma_b\Gamma_*\chi_a - 21\mathrm{e}g\mathcal{X}\bar{\lambda}\Gamma_*\lambda - 16\mathrm{e}g\mathcal{X}^{ab}\bar{\lambda}\Gamma_b\chi_a - \mathrm{e}g\left(4\mathcal{X}^{ab} - \frac{1}{10}\delta^{ab}\mathcal{X}\right)\bar{\chi}_a\Gamma_*\chi_b \\
& - 20\mathrm{e}^{-1}\bar{\psi}\Gamma_*\lambda\dot{\phi} + 2\mathrm{e}^{-1}\bar{\psi}\Gamma^b\chi^a\mathcal{P}_{ab} - \frac{1}{2}\mathrm{e}^{-1}e^{-\phi}\bar{\chi}^a\Gamma^{bc}\psi\, p_{abc} - \frac{1}{12}e^{-\phi}\bar{\chi}^a\Gamma^{bcd}\chi^a\, p_{bcd} - e^{-\phi}\bar{\chi}^a\Gamma^b\chi^c\, p_{abc} \\
& - \frac{1}{6}\mathrm{e}^{-1}e^{-\phi}\bar{\lambda}\Gamma^{abc}\Gamma_*\psi\, p_{abc} - e^{-\phi}\bar{\lambda}\Gamma^{abc}\lambda\, p_{abc} - e^{-\phi}\bar{\chi}^c\Gamma^{ab}\Gamma_*\lambda\, p_{abc} + \mathcal{L}_{\mathrm{Yuk}}[a_{ijk}],
\end{aligned} \tag{29}$$

where $D_t a_{ijk} = \dot{a}_{ijk} - 3gA_{l[i}a_{jk]l}$. The first two terms of the first two lines coincide with the nonsupersymmetric Lagrangian (24), while the Yukawa couplings in the third line involve the scalar combinations

$$\mathcal{X}^{ab} \coloneqq e^{4\phi}\left(\mathcal{T}^{ab} - \frac{1}{10}\delta^{ab}\mathcal{T}^{cc}\right), \qquad \mathcal{X} \coloneqq e^{4\phi}\mathcal{T}^{aa}, \tag{30}$$

that transform in the **54** and **1** of SO(10), respectively. Here, we have introduced the SO(10) counterpart $\mathcal{T}^{ab}$ of the scalar matrix T_{ij}. These unimodular symmetric matrices are both expressed as

$$T_{ij} = V_i{}^a V_j{}^a, \qquad \mathcal{T}^{ab} = V_i{}^a V_i{}^b, \tag{31}$$

in terms of a representative $V_i{}^a \in \mathrm{SL}(10)$ for the coset space $\mathrm{SL}(10)/\mathrm{SO}(10)$. The associated (gauge covariant) Maurer-Cartan current decomposes into

$$\begin{aligned}
(V^{-1})_a^i D_t V_{ib} &= (V^{-1})_a^i(\dot{V}_{ib} - gA_{ji}V_{jb}) \\
&\eqqcolon \mathcal{P}_{(ab)} + \mathcal{Q}_{[ab]},
\end{aligned} \tag{32}$$

where the component $\mathcal{P}_{ab}$ naturally couples to the fermions in the Lagrangian and in the supersymmetry transformations. The composite connection $\mathcal{Q}_{ab}$ enters the definition of the SO(10) covariant derivative $\mathcal{D}_t$ that acts on fermions.

Likewise, the axions couple to the fermions via the dressed current

$$p_{abc} \coloneqq (V^{-1})^i_a (V^{-1})^j_b (V^{-1})^k_c D_t a_{ijk}. \tag{33}$$

Note that the absence of a kinetic term for the gravitino in the Lagrangian is expected, as the standard Rarita-Schwinger term vanishes in one dimension. Note also that in the last line of (29), to keep the equations manageable, we have refrained from presenting explicitly the Yukawa couplings involving the axion fields. Comprehensive results, including higher-order terms in the axions, will be provided in a follow-up publication [25].

The supergravity Lagrangian (29) is invariant under local SO(10) transformations and supersymmetry. For the aims of this Letter, it is sufficient to present the supersymmetry variations of the fields in the sector where $a_{ijk} = 0$. The bosons in this truncation transform as

$$\delta_\epsilon \mathrm{e} = \bar{\epsilon}\psi, \qquad \delta_\epsilon \phi = \bar{\epsilon}\Gamma_* \lambda, \tag{34}$$

$$\delta_\epsilon V_i{}^a = \frac{1}{2}\bar{\epsilon}\Gamma^{(a}\chi^{b)} V^b_i, \tag{35}$$

$$\begin{aligned}\delta_\epsilon A_{ij} = {} & e^{4\phi}\mathrm{e}(4\bar{\epsilon}\Gamma^{ab}\lambda - 2\bar{\epsilon}\Gamma^a\Gamma_*\chi^b)V^{ab}_{ij} \\ & + e^{4\phi}\bar{\epsilon}\Gamma^{ab}\Gamma_*\psi V^{ab}_{ij},\end{aligned} \tag{36}$$

with $V^{ab}_{ij} \coloneqq V^a_{[i}V^b_{j]}$, while the fermions transformations read

$$\delta_\epsilon \psi = \mathcal{D}_t \epsilon - \frac{1}{4} g\mathrm{e}\,\mathcal{X}\Gamma_*\epsilon, \tag{37}$$

$$\delta_\epsilon \lambda = \frac{1}{2}\mathrm{e}^{-1}\dot{\phi}\Gamma_*\epsilon + \frac{1}{10} g\mathcal{X}\epsilon, \tag{38}$$

$$\delta_\epsilon \chi_a = \frac{1}{2}\mathrm{e}^{-1}\Gamma^b \epsilon \mathcal{P}_{ab} - \frac{1}{2} g\mathcal{X}_{ab}\Gamma^b\Gamma_*\epsilon. \tag{39}$$

In this sector, it becomes straightforward to study the Killing spinor equations. The gamma matrix structure implies that supersymmetry can either be broken completely, or by half, where in the latter case ϵ^α is restricted by $\Gamma_*\epsilon = \pm\epsilon$. We focus on the positive chirality solutions to align with the chirality convention of the IKKT model spinor. In the diffeomorphisms and SO(10) gauges adopted previously, where $\mathrm{e} = 1$ and $V^a_i = \delta^a_i e^{\varphi_i/2}$, and upon redefining

$$\phi_i \coloneqq \varphi_i - \phi, \qquad e^\phi = \prod_i e^{-\phi_i/10}, \tag{40}$$

the first order Killing equations reduce to

$$\partial_\tau \phi_i = -e^{\phi_i} \Rightarrow e^{\phi_i} = \frac{1}{\tau + c_i}, \tag{41}$$

with real constants c_i. It is straightforward to verify that this yields a solution to the equations of motion (25)–(27) for any choice of the c_i. Here, $\tau(t)$ is a redefined time coordinate that satisfies $\dot{\tau} = -2ge^{5\phi}$. We will show in [25] that these are the most general solutions preserving 16 supercharges, as turning on the axions necessarily leads to further breaking of supersymmetry.

The uplifts of these solutions to half-supersymmetric Euclidean IIB supergravity solutions are obtained using the truncation *Ansatz* (19)–(21), and read

$$ds^2_{10} = \frac{1}{4g^2}(\mu^i\mu^i e^{\phi_i})d\tau^2 + d\mu^i d\mu^i \frac{e^{-\phi_i}}{g^2}, \tag{42}$$

$$e^\Phi = e^{-5\phi}(\mu^i\mu^i e^{\phi_i})^{-1} = -\chi^{-1}, \tag{43}$$

with the explicit form of the functions ϕ_i in (41), and where all the other IIB p-forms vanish. For any values of the constants c_i, the IIB metric (42) still turns out to be flat in the Einstein frame. When all the constants c_i are equal, the solution reduces to the SO(10) invariant D(−1) background (16) (in the near-horizon limit [7]) that is dual to the IKKT model. More generally, for I sets of $n_I \geq 2$ equal constants, we expect these solutions to be the analogs of the smeared D3-brane solutions studied in [27], describing $\prod_I \mathrm{SO}(n_I)$-invariant distributions of N instantons spread across the nine transverse directions. They should be dual to the supersymmetric vev deformations of the IKKT model, associated with the corresponding components of the operator (11). Finally, following [28,29], the flat solutions (43) can be further uplifted on a torus to the following pp-wave solutions of pure Lorentzian gravity in twelve dimensions,

$$ds^2_{12} = dx_+ dx_- + e^\Phi dx_- dx_- + ds^2_{10}, \tag{44}$$

with the light-cone coordinates $x_\pm = x_{12} \pm x_{11}$, and where x_{11} is the Lorentzian time.

In this Letter, we have constructed maximal supergravity in one dimension, finally completing the list of maximal supergravities across all spacetime dimensions. The model, by itself, stands out as a new and distinguished example of supersymmetric quantum mechanics, featuring 32 supercharges, and an intricate highly nonlinear scalar potential which deserves further exploration. In particular, its quantization could be approached using techniques that bypass the technical difficulties of quantum field theory. It would also be interesting to explore whether the nonperturbative regime of the model can be studied using supersymmetric localization.

As a maximal gauged supergravity, this model is distinct in that it specifically builds on a sphere compactification without an obvious higher-dimensional origin for the limit $g \to 0$ of vanishing coupling constant. In maximal supergravities in higher dimensions $D > 1$, the limit $g \to 0$

reproduces the maximal supergravity obtained by toroidal reduction with a symmetry enhancement to the rank $(11-D)$ exceptional group E_{11-D}. For the present model this would suggest a highly nontrivial realization of the hyperbolic exceptional group E_{10} on the $g \to 0$ limit of (29), in line with the longstanding conjecture of [30]. One may expect further intriguing connections to the constructions of [31,32].

From the holographic perspective, the model yields a gravitational (bulk) realization of the lowest BPS multiplet of gauge invariant operators in the IKKT matrix model. As such, it sets the stage for a holographic computation of correlation functions among the operators (11)–(13) along the lines of [16]. This offers a new angle on the large N limit of IKKT which has many subtleties on its own. It would be most interesting to compare this to the existing results in this model [8–11]. Another ambitious question is the existence of a Lorentzian analog of our gravitational dual whose solutions seem to require Euclidean signature. In particular, this would have to capture the qualitative differences of the Lorentzian matrix model, revealed in recent numerical simulations [33,34].

We have also in this Letter constructed the general 1/2-BPS solutions of the model, which all uplift to IIB solutions with flat spacetime metric, and preserve all the supersymmetries of the $D(-1)$ background. We expect them to be dual to states in the "Coulomb branch" of the IKKT model, and it would be interesting to study the corresponding distributions of instantons in more detail. It will also be interesting to explore potential relations to known operators deformations of the IKKT model, in particular to the so-called polarized IKKT matrix model [15] for which recent progress has been achieved in [12–14], as well as to the deformations that are dual to the spherical branes solutions of [35,36]. Another recently explored deformation is based on holography for the $D(-1)/D7$ system [37] whose near-horizon background shares many similarities with our solutions [38].

Let us finally note, that the holographic duality discussed here is rather special in that the IIB spacetime metric of the $D(-1)$ instanton background is actually flat. Yet, the identification of states dual to the matrix model operators builds on the expansion of supergravity fields into S^9 sphere harmonics, just as in higher dimensions, and close to the standard AdS/CFT correspondence. It is tempting to speculate whether this model may in fact hold lessons for flat space holography.

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