Large superconformal near-horizons from M-theory

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We report on a classification of supersymmetric solutions to 11D supergravity, with SO(2, 2) × SO(3) isometry, which are AdS/CFT dual to 2D CFTs with \( \mathcal{N} = (0, 4) \) supersymmetry. We recover the Maldacena, Strominger, Witten (MSW) near-horizon with small superconformal symmetry and identify a class of \( \text{AdS}_3 \times S^2 \times S^2 \times CY_2 \) geometries with emergent large superconformal symmetry. This exhausts known compact geometries. Compactification of M-theory on CY_2 results in a vacuum of 7D supergravity with large superconformal symmetry, providing a candidate near-horizon for an extremal black hole and a potential new setting to address microstates.

INTRODUCTION

To obey the second law of thermodynamics, black holes must possess entropy, which Bekenstein & Hawking showed is proportional to the area of the event horizon [1]. This observation paved the way for the holographic principle and AdS/CFT [2]. One of the earliest AdS/CFT calculations (it predates the conjecture) shows that asymptotic symmetries of gravity in AdS_3 correspond to the Virasoro algebra [3], a feature of 2D CFTs. This observation together with the Cardy formula [4] for the asymptotic growth of states for a CFT with central charge c is enough to provide a microscopic derivation for the Bekenstein-Hawking (BH) entropy [5, 6]. For black holes with AdS_3 near-horizons, this methodology has been an incredible success, culminating in recent years in generalisations to extreme Kerr black holes [7], potential astrophysical bodies [8].

However, Einstein’s gravity is at best an effective description [9], and the BH entropy is expected to be corrected in a candidate UV complete theory, such as M-theory. More concretely, compactifying the 6D M5-brane theory on a four-cycle in a Calabi-Yau three-manifold, CY_3, gives rise to the MSW CFT [10], with \( \mathcal{N} = (0, 4) \) supersymmetry at low energies. The corresponding black hole exhibits the near-horizon AdS_3 × S^2 × CY_3, and subleading corrections to the BH entropy have been shown to perfectly match corrections to the central charge [11, 12].

The MSW CFT exhibits small superconformal symmetry [13], with an SU(2) R symmetry that is manifest in the two-sphere in the dual geometry. However, superconformal symmetries exist [13, 14], rich class of AdS/CFT geometries can be expected, e.g. [16]. In this letter, we identify a new class of M-theory vacua AdS_3 × S^2 × S^2 × CY_2, implying the existence of a distinct class of 2D \( \mathcal{N} = (0, 4) \) CFTs with large superconformal symmetry and R symmetry SU(2) × SU(2). We recall that CFTs with large superconformal symmetry remain largely enigmatic. While constructions based on string theory, such as AdS_3 × S^3 × S^3 × T^1 [17, 18], exist, contrary to small superconformal CFTs, interpretation as a symmetric product CFT is problematic [19].

Our results pertain to general warped AdS_3 × S^2 space-times and are not intended to apply to all M-theory geometries dual to 2D \( \mathcal{N} = (0, 4) \) SCFTs. Within our assumptions, we prove that \( M_6 \) is either CY_3, thus recovering the MSW geometry, or it possesses an additional SU(2) R symmetry that emerges from the supersymmetry analysis. Truncating the emergent SU(2) to U(1), we recover a known class [20, 31] of spacetimes with SU(2) × U(1) isometry [50].

The existence of a class of AdS_3 × S^2 × S^2 × CY_2 solutions to 11D supergravity, with 2D \( \mathcal{N} = (0, 4) \) SCFT duals, comes somewhat as a surprise. In the
case where $CY_2 = T^4$, it was shown long ago that there are geometries related through T-duality to well-known $AdS_3 \times S^3 \times T^4$ solutions in 10D \cite{17}. When $CY_2 = K_3$, the class appears new. It did not feature in a study of wrapped M5-brane geometries \cite{29}. More recently, M-theory geometries dual to 2D $\mathcal{N} = (0, 2)$ SCFTs have been discussed, but where supersymmetry is enhanced to $\mathcal{N} = (0, 4)$, the geometry is either MSW \cite{31, 32}, or no good $AdS_3$ vacuum exists \cite{33, 34}. Moreover, it is expected that M-theory on $K_3$ is dual to heterotic string theory on $T^3 \times S^3 \times CY_2$ \cite{51}. It can be expected our simply-stated results will be of interest to anyone studying the holography of 2D $\mathcal{N} = (0, 4)$ CFTs.

SO(2, 2) × SO(3)-INARIANT SPACETIMES

We recall that bosonic sector of 11D supergravity consists of a metric, $g$, and a three-form potential, $C$, with four-form field strength, $G = dC$. The equations of motion follow from the action

$$ S = \frac{1}{2\kappa^2} \int *R - \frac{1}{2} G \wedge *G - \frac{1}{6} C \wedge G \wedge G. $$

Supersymmetric solutions satisfy the Killing spinor equation (KSE):

$$ \nabla_M \eta + \frac{1}{288} \left[ \Gamma_M^{NQR} \Gamma_{NQR} - 8\delta^M_N \Gamma_{PQR} \right] G_{NPQR} \eta = 0, $$

where $M, N = 0, \ldots, 10$, $\nabla_M \eta = \partial_M \eta + \frac{1}{2} \omega_M \Gamma_{NP} \eta$, with spin connection $\omega$, and $\eta$ is a Majorana Killing spinor. It is well-known that the Einstein equation is implied by the KSE once the Bianchi identity, $dG = 0$, and equation of motion for $C$ hold \cite{38}.

2D $\mathcal{N} = (0, 4)$ CFTs enjoy both $SO(2, 2)$ conformal symmetry and $SU(2) \simeq SO(3)$ R symmetry, which motivates the general ansatz

$$ ds^2 = e^{2A} \left[ \frac{1}{m^2} d\sigma^2(AdS_3) + e^{2B} d\sigma^2(S^2) + ds^2(M_6) \right], $$

$$ G = \frac{1}{m^3} \text{vol}(AdS_3) \wedge A + \text{vol}(S^2) \wedge \mathcal{H} + \mathcal{G}, $$

where $m$ is the inverse $AdS_3$ radius, $A, B$ denote scalar warp factors and $A, \mathcal{H}, \mathcal{G}$ are respectively closed one, two and four-forms. The curvatures of symmetric spaces are canonically normalised and fields depend only on the coordinates of the internal 6D Riemannian manifold $M_6$.

In order to characterise the internal space and the fields, we decompose the 11D gamma matrices \cite{39}.

$$ \Gamma_\mu = \tau_\mu \otimes \sigma_3 \otimes \gamma_7, \quad \Gamma_\alpha = 1_2 \otimes \sigma_\alpha \otimes \gamma_7, $$

and 11D Killing spinor,

$$ \eta = \psi \otimes e^{A/2} [\chi_+ \otimes \epsilon_+ + \chi_- \otimes \epsilon_-], $$

where $\mu = 0, 1, 2$ label $AdS_3$ directions, $\alpha = 1, 2$ denote those of $S^2$, $m = 1, \ldots, 6$ correspond to $M_6$ and we define $\gamma_7 \equiv i\gamma_{123456}$. $\psi$ is a solution to the $AdS_3$ KSE, $\nabla_\mu \psi = \frac{1}{2} \tau_\mu \psi$, resulting in Poincaré spinors of definite chirality, while $\chi_{\pm}$ denote an $SU(2)$-doublet satisfying the KSE on $S^2$, $\nabla_\alpha \chi_{\pm} = \pm i \sigma_\alpha \chi_{\mp}$, with $\chi_- = \sigma_3 \chi_+$. It is a common feature of Refs. \cite{30, 39, 40} that the Majorana condition is not manifest, however conjugate spinors, $\bar{\eta}^c$, may easily be constructed e. g. \cite{41}. Following the decomposition through, one determines the effective 6D KSE equations in terms of $\epsilon_{\pm}$ \cite{38} and recasts them in terms of conditions on differential forms \cite{12}, which we illustrate later.

We stress that there is a priori no relation between $\epsilon_{\pm}$, even if one is to be expected \cite{52}. In related work, Ref. \cite{40} simplified the problem by omitting a term in the four-form flux, which enabled a simplification of the KSE analysis, before showing that the omitted term could not be reconciled perturbatively. This term was later ruled out in general \cite{43}. In the current setting, this simplification involves fixing $A = \mathcal{G} = 0$. However, since geometries with non-zero $A, \mathcal{H}, \mathcal{G}$ can be generated via T-duality \cite{27}, this simplification is difficult to motivate.

SUPERSYMMETRY CONDITIONS

We review the salient conditions on bilinears, defined in the appendix, which we construct from spinors $\epsilon_{\pm}$ \cite{39}, which encapsulate the local supersymmetry conditions we must solve. Firstly, supersymmetry demands that the following bilinears vanish \cite{39},

$$ W^- = X^+ = \text{Re}(Y) = \tilde{Z} = 0. $$

Moreover, the remaining bilinears are constrained

$$ 2mV^+ + e^{-B} \text{Im}(Y) $$

$$ = \frac{e^{-3A}}{2} \left[ \frac{1}{24!} \text{Im}(L^3)_{mn} (*6\mathcal{G})^{mn} + K^+ A^m \right]. $$

$$ \hat{Y} = -\frac{i}{2m e^B} W^+, \quad Z = -\frac{i}{2m e^B} X^-. $$

Thus, there are only three real scalars, $V^\pm, W^+$, and one complex scalar, $X^-$, which can be independent.

From the vector spinor bilinears, one can identify four real Killing vectors on $M_6$ \cite{39}, three of which, $\text{Im}(\tilde{K}^3), \text{Re}(K^4)$ and $\text{Im}(K^4)$ extend to symmetries of the overall solution \cite{49}. In contrast, the $S^2$ warp factor (also $\mathcal{H}$) depends on $K^+$, thus hinting at spacetimes with larger supersymmetry groups \cite{52}. Thankfully, $K^+$ may be truncated out consistently provided $V^+ = 0$, i. e. for 6D spinors $\epsilon_{\pm}$ with equal norm. Henceforth, we consider $\epsilon_+ \epsilon_+ = \epsilon_- \epsilon_-$, so that $V^- = K^+ = 0$. 


The scalars satisfy differential constraints \[39\],

\[
\begin{align*}
dV^+ &= 0, \\
d[e^{-B}\text{Im}(Y)] &= 0, \\
e^{-3A}d[e^{3A}X^+] &= -2m\tilde{K}^4, \\
e^{-3A}d[e^{3A}W^+] &= 2m\text{Re}(K^3),
\end{align*}
\]

while the vectors must satisfy

\[
\begin{align*}
d[e^{3A+B}K^-] &= -e^{-B}\text{Im}(Y)\mathcal{H} + e^{3A}\tilde{L}^1, \\
d[e^{6A+B}\text{Re}(\tilde{K}^3)] &= -e^{3A+B}Y\mathcal{G} + e^{6A+B}\text{Re}(L^3) \\
d[e^{6A+B}\text{Im}(\tilde{K}^3)] &= -e^{3A}W^+\mathcal{H} + 2me^{6A+2B}L^1 \\
&\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\quad\qua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\( \beta_1 = \frac{1}{2} \left( \frac{L}{R_2} \cot \frac{\zeta}{2} \alpha_4 - i \frac{L}{R_1} \alpha_1 \right), \quad \beta_4 = \frac{\beta_2}{\beta_1} \beta_3, \)
\( \beta_3 = -\frac{1}{2} \left( \frac{L}{R_2} \cot \frac{\zeta}{2} \alpha_2 + i \frac{L}{R_1} \alpha_3 \right), \quad \beta_2 = -\frac{\alpha_2}{\alpha_1} \beta_1, \)

where \( \varphi_1 + \varphi_2 = \varphi_3 + \varphi_4 \) and we have redefined \( m = L^{-1}, R_1 = e^{\beta_1}, R_2 = e^{\beta_2}/\sqrt{4m^2 e^{\beta_2} - 1}. \) With these expressions, we determine \( W^+ = V^+ \cos \zeta, X^- = V^+ \sin \zeta \sin \varphi_1 + \varphi_2 \) and solve \([11]\) and \([12]\) to show the warp factor \( e^A \) is a constant and
\[
 e^5 = -R_2 d\zeta, \quad e^6 = -R_2 \sin \zeta d\chi, \quad (26)
\]
where we have defined \( d\chi = d(\varphi_1 + \varphi_2). \) This allows us to identify the one-forms dual to the Killing vectors,
\[
 \text{Im}(K^3) = -\frac{LV^+}{2} \sin^2 \zeta d\chi, \quad K^4 = -\frac{LV^+}{2} e^x (d\zeta + i\cos \zeta d\chi), \quad (27)
\]
which correspond to an emergent \( SU(2). \) We can ensure the Killing vectors are canonically normalised through the choice \( V^+ = 2R_2^2/L. \) Solving the remaining supersymmetry conditions, one arrives at the conclusion that \( \chi \) aside, the other angular parameters are constant, with \( \mathcal{M}_6 \) being a direct-product of \( S^2 \) and \( CY_2, \) more concretely \( T^4 \) or \( K_3. \) The final expression for the four-form flux reads
\[
 G = \frac{2e^{3A}}{LV^+} \left[ -R_1^2 \tilde{L}^1 \wedge \text{vol}(S^2) + *_6 \text{Im}(L^3) \right]. \quad (28)
\]
It is easy to check that the equations of motion are satisfied, in line with expectations \([38]\). We also see that both \( \xi_1, \xi_2 \) and conjugates need to appear in the spinor. This may be contrasted with the spinor considered in Ref. \([12]\), which is not the most general, and would appear to preclude this outcome. For this reason, setting \( \beta_1 = \alpha_4 = 0 \) in \([24]\), one recovers the results of existing classifications \([29, 30]\). Setting \( A = 0, \) since the overall warp-factor is constant, we can confirm the radii satisfy
\[
 \frac{4}{L^2} = \frac{1}{R_1^2} + \frac{1}{R_2^2}. \quad (29)
\]
The ratio between \( S^2 \) radii, \( \alpha, \) corresponds to the supergroup \( D(2, 1; \alpha), \) with bosonic subgroup \( SL(2, \mathbb{R}) \times SU(2) \times SU(2). \)

To establish the connection to minimal ungauged supergravity in 7D \([40]\), we exploit the following consistent Kaluza-Klein reduction ansatz:
\[
ds^2_{11} = e^{-2B} ds^2_5 + e^{2B} ds^2(CY_2),
\]
\[
 G = F + \sum_{a=1}^{3} F^a \wedge J^a, \quad (30)
\]
where \( J^a \) denote the three self-dual harmonic two-forms of \( CY_2, \) \( B \) is a scalar and \( F \) and \( F^a \) are respectively field strengths corresponding to a three-form and one-form potentials, \( F = dC, F^a = dA^a. \) The resulting action in Einstein frame in 7D is
\[
 \mathcal{L}_7 = R \text{vol}_7 - \frac{36}{5} dB \wedge *_7 dB - \frac{1}{2} e^{2B} F \wedge *_7 F
 - e^{-2B} F^a \wedge *_7 F^a - F \wedge F^a \wedge A^a. \quad (31)
\]
To cast the action in the original notation of ref. \([46],\) one should employ the following redefinitions:
\[
 B = \sqrt{\frac{5}{6}} \phi, \quad F_{us} = \sqrt{2} F_{\text{them}}, \quad F_{us}^a = \sqrt{2} F_{\text{them}}^a. \quad (32)
\]

**DISCUSSION**

We have initiated a classification of all solutions to 11D supergravity with \( SO(2, 2) \times SO(3) \) isometry. This is the simplest geometric signature of a supergravity solution dual to a 2D CFT with \( \mathcal{N} = (0, 4) \) supersymmetry, including the MSW CFT. In the process, we have identified a novel class of near-horizon geometries in M-theory with large superconformal symmetry. Compactifying M-theory on \( CY_2, \) we identify a resulting \( AdS_3 \times S^2 \times S^2 \) vacuum to 7D supergravity, thus providing a candidate near-horizon for an extremal black hole and a potential new controlled setting to count black hole microstates.

The M-theory geometry provides a unifying description of well-known \( AdS_3 \times S^2 \times S^2 \times S^1 \) geometries of Type II string theory through T-duality \([27]\), and heterotic vacua via M-theory/heterotic duality \([45]\). A careful treatment of the central charge reveals the expected form of a large superconformal algebra \([47]\)
\[
c \sim \frac{k^+ k^-}{k^+ + k^-}. \quad (33)
\]
with affine \( SU(2)_L \) current algebras at levels \( k^\pm \) related to the quantised charges, yet where \( c \sim N^2, \) for large charge \( N, \) and not the more usual \( c \sim N^3 \) of geometries corresponding to M5-branes.

Our work has two interesting implications. Firstly, it is striking that the \( AdS_3 \times S^2 \times S^2 \times CY_2 \) geometries are not identifiable as \( AdS_3 \) limits of wrapped M5-branes \([29]\). This suggests the M5-brane picture is novel and motivates further study to understand anomaly in-flow \([11]\). Secondly, as we have shown, since 11D supergravity compactifies on \( CY_2 \) to 7D minimal supergravity, the \( AdS_3 \times S^2 \times S^2 \) solution hints at being the near-horizon of an extremal black hole. While such solutions have in principle been classified \([45]\), we are not aware of a near-horizon uniqueness theorem in 7D, cf. \([49]\). Assuming a black hole exists, strong parallels to the MSW case, with M-theory compactified on Calabi-Yau, are expected to facilitate a microscopic derivation of the entropy. Since the small superconformal algebra is recovered from the large one through a decompactification of a two-sphere,
it is tempting to speculate that contact with the MSW results may be made in the same limit.

Lastly, we remark that we have assumed $SU(2)$-structure, and more general solutions with identity structure are known to exist [27]. We hope to extend the classification to consider more general internal manifolds in future work [47].

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Spinor bilinears

In our conventions, the 6D gamma matrices are Hermitian $\gamma^\dagger_m = \gamma_m$ and anti-symmetric $\gamma^T_m = -\gamma_m$. Consistent with the symmetries of the gamma matrices, given $\epsilon_\pm$, we can define an exhaustive set of scalar

$$V^\pm = \frac{1}{2}(\epsilon^\dagger_+ \epsilon_+ \pm \epsilon^\dagger_- \epsilon_-),$$

$$W^\pm = \frac{1}{2}(\epsilon^\dagger_+ \gamma \gamma \epsilon_+ \pm \epsilon^\dagger_- \gamma \gamma \epsilon_-),$$

$$X^\pm = \frac{1}{2}(\epsilon^{T\dagger}_+ \epsilon_+ \pm \epsilon^{T\dagger}_- \epsilon_-),$$

$$Y = \epsilon^\dagger_+ \epsilon_-, \quad \tilde{Y} = \epsilon^\dagger_+ \gamma \gamma \epsilon_-, \quad Z = \epsilon^{T\dagger}_+ \epsilon_-, \quad \tilde{Z} = \epsilon^{T\dagger}_+ \gamma \gamma \epsilon_-,$$

(34)

and vector spinor bilinears:

$$K^\pm_m = \frac{1}{2}(\epsilon^\dagger_+ \gamma_m \epsilon_+ \pm \epsilon^\dagger_- \gamma_m \epsilon_-),$$

$$K^0_m = \frac{i}{2}(\epsilon^\dagger_+ \gamma_m \gamma \gamma \epsilon_+ \pm \epsilon^\dagger_- \gamma_m \gamma \gamma \epsilon_-),$$

$$K^3_m = \epsilon^\dagger_+ \gamma_m \epsilon_-, \quad \tilde{K}^3_m = \epsilon^\dagger_+ \gamma_m \gamma \gamma \epsilon_-, \quad K^4_m = \epsilon^{T\dagger}_+ \gamma_m \epsilon_-,$$

$$\tilde{K}^4_m = \epsilon^{T\dagger}_+ \gamma_m \gamma \gamma \epsilon_-,$$

(35)

where factors of $i$ ensure vectors are real. We define the following two-forms:

$$L^1_{mn} = \frac{i}{2}(\epsilon^\dagger_+ \gamma_m \epsilon_+ \pm \epsilon^\dagger_- \gamma_m \epsilon_-),$$

$$L^3_{mn} = \epsilon^\dagger_+ \gamma_m \epsilon_-, \quad L^4_{mn} = \epsilon^{T\dagger}_+ \gamma_m \epsilon_-,$$

where notation follows Ref. [33].

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Indeed, the existence of the $SU(2) \times U(1)$ isometry is a puzzle, since typically emergent isometries are $R$ symmetries, but $SU(2) \times U(1)$ does not fit with a known $N = 4$ superconformal algebra. Recently, the first explicit example of a geometry in this class was identified and the emergent $U(1)$ symmetry shown not to be an $R$ symmetry, but rather the M-theory circle.

It has been suggested that the dilaton is singular in the near-horizon of intersecting NS5-branes and this violates the partial integration argument of ref. We thank G. Papadopoulos for correspondence.

Naive supersymmetry counting works here: four $AdS_3$ spinors and two $S^2$ spinors already give eight supersymmetries, so we just expect one from $M_6$.

We have checked that $\tilde{K}^+$ is non-zero for maximally supersymmetric $AdS_7 \times S^4$.

We have corrected a sign typo in [14].