

Stable Nonsupersymmetric Anti-de Sitter Vacua of Massive IIA Supergravity

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Stable nonsupersymmetric anti-de Sitter (AdS) vacua of string theory are widely believed not to exist. In this Letter, we analytically compute the full bosonic Kaluza-Klein spectrum around the G_2 -invariant nonsupersymmetric AdS_4 solution of massive IIA supergravity and show that it is perturbatively stable. We also provide evidence that six other nonsupersymmetric AdS_4 solutions of massive IIA supergravity are perturbatively stable. Since previous studies have indicated that these AdS vacua may also be non-perturbatively stable, our findings pose a challenge to the swampland conjecture.

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Constructing stable nonsupersymmetric vacua of string theory is one of its foremost challenges and crucial for applications to phenomenology, cosmology, and holography. Unlike their supersymmetric counterparts, such vacua are not protected by supersymmetry arguments and positive mass theorems [1]. The stability of nonsupersymmetric anti-de Sitter (AdS) vacua is particularly interesting, since they would provide a crucial tool for applying holography to realistic systems in condensed matter or indeed QCD. Moreover, nonsupersymmetric AdS vacua might be a stepping stone to understanding time-dependent de Sitter solutions in string theory.

The stability of nonsupersymmetric AdS vacua has proven particularly contentious in recent years. Various arguments have been put forward based on the weak gravity conjecture [2] that all nonsupersymmetric AdS vacua of string theory are unstable, ultimately leading to the AdS swampland conjecture [3]. Moreover, despite years of trying, no fully fledged case of a stable nonsupersymmetric AdS vacuum is known in string theory. Indeed, a powerful method for constructing nonsupersymmetric AdS vacua involves uplifting nonsupersymmetric solutions of lower-dimensional supergravities via a consistent truncation. However, all but a small number of these AdS solutions are already unstable within the lower-dimensional supergravity, with some of the scalar fields violating the Breitenlohner-Freedman (BF) bound [4]. Notable exceptions to this are the nonsupersymmetric $\text{SO}(3) \times \text{SO}(3)$ -

invariant AdS_4 vacuum of four-dimensional $\mathcal{N} = 8$ $\text{SO}(8)$ gauged supergravity [5] obtained by consistent truncation of 11-dimensional supergravity on S^7 [6], and seven nonsupersymmetric AdS_4 vacua of four-dimensional $\mathcal{N} = 8$ $\text{ISO}(7)$ gauged supergravity (see Appendix A of Ref. [7] for a summary and original references) obtained by consistent truncation of massive IIA supergravity on S^6 [8]. Remarkably, for these nonsupersymmetric AdS_4 vacua, all 70 scalar fields in the $\mathcal{N} = 8$ supergravity multiplet have masses above the BF bound [7,9].

Although these vacua are stable in four dimensions, their higher-dimensional stability is far from guaranteed. For example, the Kaluza-Klein (KK) spectrum around such vacua could contain tachyonic scalars violating the BF bound, or the AdS vacua may exhibit nonperturbative instabilities. Understanding the higher-dimensional stability has long been a challenge for these vacua, since both the computation of Kaluza-Klein spectra and the systematic search for nonperturbative instabilities (see, e.g., Ref. [10]) used to be notoriously difficult. However, within the last year, a new and powerful method based on exceptional field theory (ExFT) [11] was developed in Refs. [12,13] for computing Kaluza-Klein spectra of vacua of maximal gauged supergravities. At the same time, Ref. [14] showed that probe branes can be used to easily search for signals, dubbed brane-jet instabilities, of nonperturbative instabilities [10]. These new methods were used to show that the $\text{SO}(3) \times \text{SO}(3)$ -invariant AdS_4 vacuum has tachyonic Kaluza-Klein modes [15] and suffers from M2-brane-jet instabilities [14], with brane-jet instabilities also arising for other AdS vacua [16].

In this Letter, we will investigate the Kaluza-Klein spectrum of the seven nonsupersymmetric AdS_4 vacua of the maximal $\text{ISO}(7)$ supergravity, which correspond to

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compactifications of massive IIA supergravity on various deformed S^6 spheres and are stable within the four-dimensional supergravity. In particular, we focus on the vacuum with the largest bosonic symmetry, G_2 . Remarkably, and in contrast to the $SO(3) \times SO(3)$ symmetric AdS₄ vacuum of the SO(8) theory, the AdS₄ vacuum with G_2 symmetry has shown to be Dp-brane-jet stable for $p = 2, 4, 6, 8$, while some other nonperturbative channels have also been ruled out [7]. Moreover, a partial search for the other six AdS₄ vacua has also found no signs of D2-brane-jet instabilities [7]. Finally, because the compactification spaces have S^6 topology supported by fluxes, these vacua do not suffer from decays due to “bubbles of nothing” of the type argued in Ref. [17]. This suggests that these vacua may be nonperturbatively stable, making it even more important to study their perturbative Kaluza-Klein stability in light of the weak gravity and swampland conjectures. This is what we set up to do in this Letter. In particular, we compute the full Kaluza-Klein spectrum of the G_2 -invariant vacuum, providing the first analytic spectrum of a nonsupersymmetric vacuum, and use this to prove its perturbative stability. Moreover, we collect numerical evidence for the perturbative stability of the six other AdS₄ vacua.

In Refs. [12,13], ExFT was used to derive mass matrices for the Kaluza-Klein spectrum around any vacuum of $\mathcal{N} = 8$ gauged supergravities in four and five dimensions that arises from a consistent truncation of ten- or 11-dimensional supergravity. Let us briefly review some of the salient features of ExFT and Kaluza-Klein spectroscopy.

ExFT is a reformulation of ten- or 11-dimensional supergravity, which unifies the metric and flux degrees of freedom within a manifestly $E_{7(7)}$ covariant formulation. Its bosonic sector

$$\{g_{\mu\nu}, \mathcal{M}_{MN}, \mathcal{A}_\mu^M\}, \quad \begin{aligned} \mu &= 0, \dots, 3, \\ M &= 1, \dots, 56 \end{aligned} \quad (1)$$

consists of an external and an internal metric $g_{\mu\nu}$, \mathcal{M}_{MN} , respectively, with the latter parametrizing the coset space $E_{7(7)}/\text{SU}(8)$, together with vector fields \mathcal{A}_μ^M transforming in the **56** of the group $E_{7(7)}$.

As shown in Refs. [12,13], a general Kaluza-Klein fluctuation around a vacuum that uplifts from four-dimensional gauged supergravity can be expressed as a product of the modes of the consistent truncation, with a complete basis of functions on the compactification manifold. A powerful feature of this method is that this complete basis of functions can be chosen to be the scalar harmonics \mathcal{Y}^Σ of the compactification with a metric that preserves the largest possible symmetry group G_{max} in the lower-dimensional gauged supergravity. In the case of the maximal ISO(7) supergravity investigated in this Letter, the internal space topology is S^6 , the largest possible

symmetry group is $G_{\text{max}} = \text{SO}(7)$, and we can choose the \mathcal{Y}^Σ to be the scalar harmonics on the round S^6 .

The fluctuation ansatz of the ExFT fields (1) around an AdS₄ vacuum is given by [12,13]

$$\begin{aligned} g_{\mu\nu}(x, y) &= \rho^{-2} \left(\dot{g}_{\mu\nu}(x) + \sum_{\Sigma} \mathcal{Y}^\Sigma h_{\mu\nu, \Sigma}(x) \right), \\ \mathcal{A}_\mu^M(x, y) &= \rho^{-1} (U^{-1})_{\underline{A}}^M \sum_{\Sigma} \mathcal{Y}^\Sigma A_{\mu}^{\underline{A}, \Sigma}(x), \\ \mathcal{M}_{MN}(x, y) &= U_M^{\underline{A}} U_N^{\underline{B}} \left(\delta_{\underline{A}\underline{B}} + \mathcal{P}_{\underline{A}\underline{B}, I} \sum_{\Sigma} \mathcal{Y}^\Sigma j_{I, \Sigma}(x) \right), \end{aligned} \quad (2)$$

where the Kaluza-Klein fluctuations for the metric, vector fields, and scalars are labeled by $h_{\mu\nu, \Sigma}(x)$, $A_{\mu}^{\underline{A}, \Sigma}$, and $j_{I, \Sigma} \in \mathfrak{e}_{7(7)} \ominus \mathfrak{su}(8)$, respectively. The latter appear under projection $\mathcal{P}_{\underline{A}\underline{B}, I}$, with $I = 1, \dots, 70$ resulting from the expansion of the group element \mathcal{M}_{MN} . On the other hand, $\rho(y) \in \mathbb{R}^+$ and $U_M^{\underline{A}}(y) \in E_{7(7)}$ denote the scaling function and the twist matrix, respectively, encoding the consistent truncation to $\mathcal{N} = 8$ gauged supergravity [18]. In the fluctuation ansatz (2), the twist matrix appears dressed with the scalar matrix of the four-dimensional supergravity $\mathcal{V}_{\underline{M}}^{\underline{A}} \in E_{7(7)}/\text{SU}(8)$ evaluated at the scalar configuration specifying the vacuum of the maximal $D = 4$ supergravity, i.e.,

$$U_M^{\underline{A}}(y) = U_M^{\underline{M}}(y) \mathcal{V}_{\underline{M}}^{\underline{A}}. \quad (3)$$

The fluctuation ansatz (2) induces mass matrices for the Kaluza-Klein spectrum which are entirely expressed through the embedding tensor of the $\mathcal{N} = 8$ gauged supergravity $X_{\underline{MN}}^{\underline{P}}$ and the linear action $\mathcal{T}_{\underline{M}}^{\underline{\Sigma}, \Omega}$ of the G_{max} Killing vector fields on the scalar harmonics, both dressed by the scalar matrix $\mathcal{V}_{\underline{M}}^{\underline{A}}$. The matrices $\mathcal{T}_{\underline{M}}^{\underline{\Sigma}, \Omega}$ are explicitly defined as

$$L_{\mathcal{K}_{\underline{M}}} \mathcal{Y}^\Sigma = -\mathcal{T}_{\underline{M}}^{\underline{\Sigma}, \Omega} \mathcal{Y}^\Omega, \quad (4)$$

where $\mathcal{K}_{\underline{M}}$ are the Killing vectors generating G_{max} , which can be extracted from the twist matrix $U_M^{\underline{M}}$. In turn, the dressed objects are defined as

$$\begin{aligned} X_{\underline{AB}}^{\underline{C}} &= (\mathcal{V}^{-1})_{\underline{A}}^{\underline{M}} (\mathcal{V}^{-1})_{\underline{B}}^{\underline{N}} X_{\underline{MN}}^{\underline{P}} \mathcal{V}_{\underline{P}}^{\underline{C}}, \\ \mathcal{T}_{\underline{A}}^{\underline{\Sigma}, \Omega} &= (\mathcal{V}^{-1})_{\underline{A}}^{\underline{M}} \mathcal{T}_{\underline{M}}^{\underline{\Sigma}, \Omega}. \end{aligned} \quad (5)$$

The mass matrices are obtained by linearizing the ExFT field equations with the fluctuation ansatz (2) [12,13]. For the purpose of this Letter, we give the scalar mass matrix in a yet more compact form:

$$\begin{aligned} \mathcal{M}_{I\Sigma, J\Omega}^{(\text{scalar})} &= \mathcal{M}_{IJ}^{(0)} \delta_{\Sigma\Omega} + \delta_{IJ} \mathbb{M}_{\Sigma\Omega}^{(2)} + \mathcal{N}_{IJ}{}^C \mathcal{T}_{\underline{C}\Sigma\Omega} \\ &\quad - \frac{1}{6} (\Pi^T \Pi)_{I\Sigma, J\Omega}. \end{aligned} \quad (6)$$

Here, $\mathcal{M}_{IJ}^{(0)}$ and $\mathbb{M}_{\Sigma\Omega}^{(2)} = -(\mathcal{T}_{\underline{A}} \mathcal{T}_{\underline{A}})_{\Sigma\Omega}$ are the mass matrices of the four-dimensional supergravity scalars and of the spin-2 fluctuations, respectively. The matrices in the last two terms are given by

$$\begin{aligned} \Pi_{\underline{A}\Sigma, I\Omega} &= \delta_{\Sigma\Omega} X_{\underline{A}C}{}^D \mathcal{P}_{\underline{C}D, I} - 12 \mathcal{P}_{\underline{A}D, I} \mathcal{T}_{\underline{D}\Omega}{}^\Sigma, \\ \mathcal{N}_{IJ}{}^C &= -4(X_{\underline{C}A}{}^B + 12X_{\underline{A}B}{}^C) \mathcal{P}_{\underline{A}D}{}^I \mathcal{P}_{\underline{B}D}{}^J. \end{aligned} \quad (7)$$

In particular, the matrix $\Pi_{\underline{A}\Sigma, I\Omega}$ features in the linearized covariant scalar derivatives, $\partial_\mu j_{I, \Sigma} - \mathcal{A}_\mu{}^{C\Omega} \Pi_{\underline{C}\Omega, I\Sigma}$ and thus encodes the projection onto the Goldstone scalars. Accordingly, it is orthogonal to the mass matrix and encodes the vector mass matrix as $\mathcal{M}^{(\text{vector})} = \Pi \Pi^T$.

Let us now specialize to the specific ISO(7) gauged supergravity whose vacua we are interested in. In terms of $\text{SL}(7) \subset E_{7(7)}$ defined by the branching of the **56** as

$$A^{\underline{M}} \rightarrow \{A^{[ab]}, A^{a8}, A_{[ab]}, A_{a8}\}, \quad (8)$$

with $a = 1, \dots, 7$, the only nonzero components of the embedding tensor are

$$X_{\underline{M}\underline{N}}{}^P = \begin{cases} X_{abcd}{}^{ef} = -X_{ab}{}^{ef}{}_{cd} = -8\delta_{[a}^e \delta_{b][c} \delta_{d]}^f, \\ X_{abc8}{}^{d8} = -X_{ab}{}^{d8}{}_{c8} = -2\delta_{[a}^d \delta_{b]c}, \\ X_{c8ab}{}^{d8} = -X_{c8}{}^{d8}{}_{ab} = -2\delta_{c[a} \delta_{b]}^d, \\ X^{c8}{}_{ab}{}^{d8} = -X^{c8}{}_{d8}{}^{ab} = 2c\delta_{[a}^c \delta_{b]}^d. \end{cases} \quad (9)$$

Here, c is a magnetic parameter that renders the gauging of dyonic type [19,20] and which is identified with the Romans mass parameter [21] in ten dimensions.

For the fluctuation ansatz, we will be using the scalar harmonics on the round S^6 , which can be expressed as symmetric traceless polynomials in the elementary harmonics \mathcal{Y}^a satisfying $\mathcal{Y}^a \mathcal{Y}^b \delta_{ab} = 1$. These are given by

$$\mathcal{Y}^\Sigma = \{1, \mathcal{Y}^a, \mathcal{Y}^{a_1 a_2}, \dots, \mathcal{Y}^{a_1 \dots a_\ell}, \dots\}, \quad (10)$$

with $\mathcal{Y}^{a_1 \dots a_\ell} \equiv \mathcal{Y}^{((a_1 \dots a_\ell))}$ and $((\))$ denotes traceless symmetrization. Therefore, the $\mathcal{T}_{\underline{M}}{}^\Sigma{}_\Omega$ now simply become the SO(7) generators in the $[\ell, 0, 0]$ representation, with the only nonzero components given by

$$\mathcal{T}_{ab}{}^{c_1 \dots c_\ell}{}_{d_1 \dots d_\ell} = 4\ell \delta_{[a}^{(c_1} \delta_{b]((d_1} \delta_{d_2}{}^{c_2} \dots \delta_{d_\ell)}{}^{c_\ell))}. \quad (11)$$

We stress once more that the harmonics (10) constitute an appropriate basis for the analysis of all vacua of this theory, regardless of the specific S^6 metric occurring at the vacua. With the data in Eqs. (9) and (11) specific to the maximal

ISO(7) supergravity, we can now proceed and compute the KK spectrum of the metric, vector, and scalar fluctuations around a given AdS₄ vacuum specified by the scalar matrix $\mathcal{V}_{\underline{M}}{}^{\underline{A}}$ via the dressed tensors in Eq. (5).

The maximal ISO(7) supergravity possesses a rich structure of AdS₄ vacua. To date, 60 such vacua have been identified up to discrete degeneracies (see Appendix A of Ref. [7]), all of which are contained within a \mathbb{Z}_2^3 -invariant sector of the theory [22]. This sector describes an $\mathcal{N} = 1$ supergravity coupled to seven chiral multiplets with complex scalar components z_i ($i = 1, \dots, 7$). Thus, the 56-bein entering the fluctuation ansatz (2) via the dressed twist matrix in Eq. (3) parametrizes a coset subspace

$$\mathcal{V}_{\underline{M}}{}^{\underline{A}}(z_i) \in [\text{SL}(2)/\text{SO}(2)]^7 \subset E_{7(7)}/\text{SU}(8). \quad (12)$$

The Kähler potential $K(z_i, \bar{z}_i)$ and holomorphic superpotential $W(z_i)$ for the complex z_i are given by

$$\begin{aligned} K(z_i, \bar{z}_i) &= -\sum_{i=1}^7 \log[-i(z_i - \bar{z}_i)], \\ W(z_i) &= 2g[z_1 z_2 z_3 + z_1 z_6 z_7 + z_2 z_5 z_7 + z_3 z_5 z_6 \\ &\quad + (z_1 z_5 + z_2 z_6 + z_3 z_7) z_4 + c], \end{aligned} \quad (13)$$

where g is the coupling constant in the four-dimensional supergravity that relates to the (inverse) radius of S^6 .

This truncated sector contains seven nonsupersymmetric AdS₄ vacua, yet stable within $D = 4$ supergravity. I.e., all 70 scalars from the $\mathcal{N} = 8$ supergravity multiplet have masses above the BF bound. At these vacua, the ISO(7) symmetry of the maximal theory is broken to G_{res} [23–25], and the mass spectra organize into representations of G_{res} .

Employing the techniques reviewed above, we now compute the full bosonic Kaluza-Klein spectra around these nonsupersymmetric AdS₄ vacua. Let us first focus on the G_2 -invariant vacuum located at

$$z_{1,2,3,4,5,6,7} = 2^{-\frac{4}{3}}(-1 + \sqrt{3}i)c^{\frac{1}{3}}, \quad (14)$$

with AdS radius $L^2 = 2^{-\frac{10}{3}} 3^{\frac{3}{2}} g^{-2} c^{\frac{1}{3}}$, and which was first constructed directly in ten dimensions in Ref. [26]. The mass spectrum of bosonic fields in the maximal supergravity was presented in Ref. [23], and a first rudimentary study of its higher-dimensional stability was performed in Ref. [27]. Here we will compute its full bosonic Kaluza-Klein spectrum by exploiting the large residual symmetry group G_2 to analytically evaluate the mass matrices (6).

The normalized spin-2 mass matrix reduces to the G_2 Casimir operator

$$L^2 \mathbb{M}_{\Sigma\Omega}^{(2)} = -L^2 (\mathcal{T}_{\underline{A}} \mathcal{T}_{\underline{A}})_{\Sigma\Omega} = \frac{3}{2} (\text{Cas}_{G_2})_{\Sigma\Omega}, \quad (15)$$

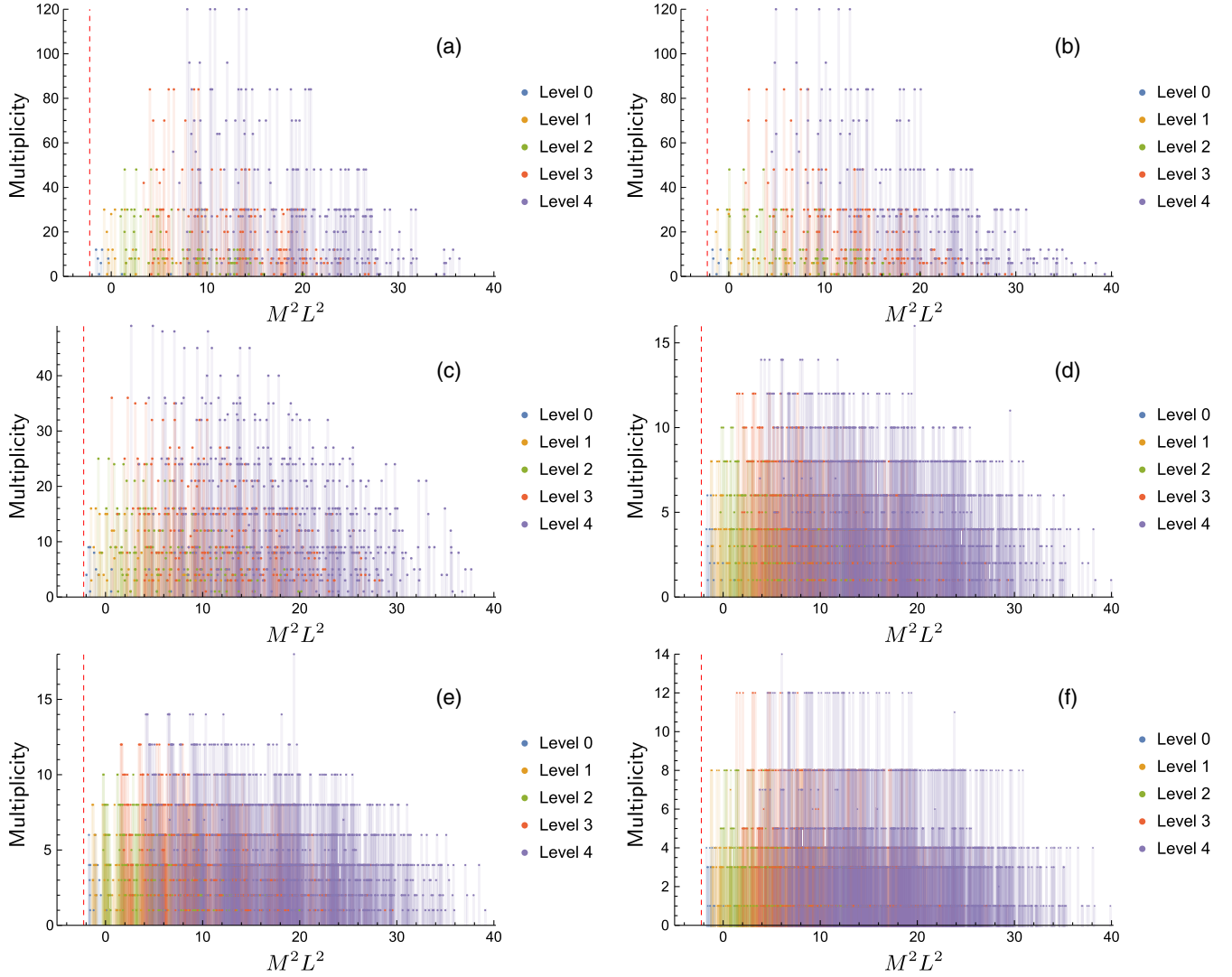


FIG. 1. Normalized masses $M^2 L^2$ and multiplicities for the first four KK levels of spin-0 fluctuations around the six non-supersymmetric and BF stable $\text{AdS}_4 \times S^6$ solutions of massive IIA supergravity given in Table I. The subfigure labels (a)–(f) correspond to the ID's of the different vacua in Table I. The dashed red line marks the BF bound $M^2 L^2 = -\frac{9}{4}$.

whose eigenvalues are given by [28]

$$C_{[n_1, n_2]} = \frac{1}{3}(5n_1 + n_1^2 + 9n_2 + 3n_1 n_2 + 3n_2^2) \quad (16)$$

on the $[n_1, n_2]$ representation. The spin-2 fluctuations at level ℓ transform in the $[\ell, 0]$ representation, which gives the mass eigenvalues $\frac{1}{2}\ell(\ell + 5)$, in accordance with the results in Ref. [29].

For the scalar Kaluza-Klein modes, the first three terms of the (normalized) mass matrix (6) can be shown to combine into

$$6\delta_{IJ}\delta_{\Sigma\Omega} + 3(\text{Cas}_{G_2})_{\Sigma\Omega}\delta_{IJ} - \frac{3}{2}(\text{Cas}_{G_2})_{I\Sigma, J\Omega}. \quad (17)$$

where the two Casimir operators act on the representations of the spin-2 harmonics and on the scalar fluctuations $j_{I\Sigma}$,

respectively. In contrast, the last term in Eq. (6) only affects the Goldstone scalars and ensures that these appear with zero eigenvalues. Putting this together with Eq. (16), we find that the normalized scalar masses at level ℓ allow for a surprisingly compact expression in terms of the quadratic G_2 Casimir eigenvalues as

$$M_{[n_1, n_2]_\ell}^2 L^2 = (\ell + 2)(\ell + 3) - \frac{3}{2}C_{[n_1, n_2]}, \quad (18)$$

where $[n_1, n_2]_\ell$ denotes the G_2 representation that the Kaluza-Klein modes at level ℓ appear in. The form of the scalar mass matrix (6) implies that this mass formula also extends to the spin-1 sector. In particular, when evaluated for the first two levels, it reproduces the vector masses computed in Ref. [30].

TABLE I. Identifier, residual symmetry group, location, and AdS radius of the six nonsupersymmetric and BF-stable AdS₄ vacua of the maximal ISO(7) supergravity with SU(3), SO(4), U(1) × SO(3), and SO(3) symmetry.

ID	G_{res}	$c^{-\frac{1}{3}}z_i$	$g^2 c^{-\frac{1}{3}}L^2$
(a)	SU(3)	$z_{1,2,3} = -0.2698 + 0.7329i, z_{4,5,6,7} = 0.4915 + 0.6618i$	0.512522
(b)	SU(3)	$z_{1,2,3} = -0.4547 + 0.8376i, z_{4,5,6,7} = 0.3348 + 0.6012i$	0.511579
(c)	SO(4)	$z_{1,2,3,5,6,7} = -0.4123 + 0.6512i, z_4 = -0.0678 + 1.1466i$	0.510363
(d)	U(1) × SO(3)	$z_{1,7} = -0.5061 + 0.7328i, z_{2,3,5,6} = -0.3520 + 0.6240i, z_4 = -0.2732 + 0.9609i$	0.511595
(e)	U(1) × SO(3)	$z_{1,7} = -0.4869 + 0.7937i, z_{2,3,5,6} = -0.3382 + 0.6059i, z_4 = -0.3736 + 0.9082i$	0.511536
(f)	SO(3)	$z_{1,7} = -0.5065 + 0.7285i, z_{2,6} = -0.3663 + 0.6256i, z_{3,5} = -0.3403 + 0.6259i, z_4 = -0.2671 + 0.9623i$	0.511594

For the spin-0 Kaluza-Klein modes, the relevant representations at level ℓ are the product of those of level 0, $2([2, 0] \oplus [1, 0] \oplus [0, 0])$, with $[\ell, 0]$, taking care to remove the Goldstone bosons. For generic ℓ , this gives

$$\begin{aligned}
 &2[\ell - 4, 2] \oplus [\ell - 3, 1] \oplus [\ell - 3, 2] \oplus 2[\ell - 2, 0] \\
 &\oplus 3[\ell - 2, 1] \oplus 2[\ell - 2, 2] \oplus [\ell - 1, 0] \\
 &\oplus 3[\ell - 1, 1] \oplus 5[\ell, 0] \oplus [\ell, 1] \\
 &\oplus [\ell + 1, 0] \oplus 2[\ell + 2, 0], \quad (19)
 \end{aligned}$$

with care needed for the degeneracies at low ℓ . Similarly, the spin-1 Kaluza-Klein modes at generic level ℓ transform in the representations

$$\begin{aligned}
 &[\ell - 3, 1] \oplus [\ell - 3, 2] \oplus 3[\ell - 2, 1] \oplus 3[\ell - 1, 0] \\
 &\oplus 3[\ell - 1, 1] \oplus 2[\ell, 0] \oplus [\ell, 1] \oplus 3[\ell + 1, 0], \quad (20)
 \end{aligned}$$

again with care needed for the degeneracies at low ℓ . Evaluation of the mass formula (18) on the list of G_2 representations in Eq. (19) for the tower of spin-0 Kaluza-Klein modes shows that the masses increase with increasing Kaluza-Klein level ℓ and are positive. This proves the perturbative stability of the nonsupersymmetric G_2 vacuum.

We finally turn to the other six nonsupersymmetric AdS₄ solutions of massive IIA supergravity with smaller G_{res} , which correspond to compactifications of massive IIA supergravity on various deformed S^6 spheres and are summarized in Table I. The small residual symmetry groups G_{res} do not allow us (for the moment) to analytically resolve these mass spectra. Instead, we revert to a numerical evaluation of the mass matrices (6) up to and including KK level $\ell = 4$, with results displayed in Fig. 1. Remarkably, the analysis shows that all scalars on these levels are perturbatively stable with their normalized masses lying above the BF bound. Moreover, as is pictured in Fig. 1, the lowest-lying normalized masses at each Kaluza-Klein level ℓ increase monotonically with the level, thus suggesting that the full spectra will also be stable.

To summarize, we have analytically computed the full Kaluza-Klein spectrum of the nonsupersymmetric and

G_2 -invariant AdS₄ × S^6 background of massive IIA supergravity, resulting in the closed mass formula (18) encoding the entire scalar spectrum. This spectrum does not contain masses below the BF bound at any level in the KK tower of spin-0 fluctuations. Therefore, it presents an example of a nonsupersymmetric, yet perturbatively stable, solution of ten-dimensional massive IIA supergravity. We recall that this solution is a well-defined background of massive IIA string theory [8,31], and that the analysis of Ref. [7] has moreover excluded Dp -brane-jet instabilities for $p = 2, 4, 6, 8$. We have also presented evidence that six other nonsupersymmetric AdS₄ × S^6 backgrounds of massive IIA supergravity with smaller residual symmetry are also perturbatively stable in ten dimensions.

Since these AdS₄ vacua have already passed a number of nontrivial tests regarding their nonperturbative stability, they seem to provide counterexamples to the belief that there are no stable nonsupersymmetric AdS vacua of string theory. Therefore, our findings here challenge the AdS swampland conjecture [3]. If the swampland conjecture is to hold, there must be an alternative decay channel for these nonsupersymmetric AdS₄ vacua, which it would be imperative to unearth.

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