$\frac{1}{2}$-BPS Domain wall from $N = 10$ three dimensional gauged supergravity

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Abstract: We explicitly construct $N = 10$ Chern-Simons gaged supergravity in three dimensions with non-semisimple gauge group $SO(5) \ltimes T^{10}$. The gauge group is embedded in $E_{6(-14)}$ which is the isometry group of the 32-dimensional scalar manifold $E_{6(-14)}/SO(10) \times U(1)$. The resulting theory is on-shell equivalent to $SO(5)$ Yang-Mills gauged supergravity coming from dimensional reduction on $S^1$ of $SO(5) \ N = 5$ gauged supergravity in four dimensions. We discuss the spectrum of the corresponding reduction. The $SO(5) \ltimes T^{10}$ gauged supergravity, describing the reduced theory, admits a $\frac{1}{2}$-BPS domain wall vacuum solution whose explicit form is also given. This provides an example of a domain wall in non-maximal gauged supergravity.

Keywords: AdS-CFT Correspondence, Gauge-gravity correspondence, Supergravity models.
1. Introduction

Chern-Simons gauged supergravity in three dimensions has a very rich structure due to the duality between scalars and vectors in three dimensions. There are many possible gauge groups since there is no restriction on the number of vector fields that act as gauge fields [1, 2], or equivalently, no restriction on the dimension of the gauge group provided that it can be embedded in the global symmetry group and consistent with supersymmetry. Any number of vector fields can be introduced via Chern-Simons terms which do not give rise to extra degrees of freedom. The theory is also useful in the study of AdS$_3$/CFT$_2$ correspondence, see for example [3] for a nice review.

To understand AdS$_3$/CFT$_2$ correspondence in the context of string/M theory, the embedding of three dimensional gauged supergravity in ten or eleven dimensions is required. The usual procedure to obtain lower dimensional supergravities from higher dimensional theories is the Kaluza-Klein (KK) dimensional reduction. The general U-duality covariant formulation of three dimensional gauged supergravities is in the form of Chern-Simons theory in which the gauge fields enter the Lagrangian through the Chern-Simons terms [4]. On the other hand, dimensional reductions result in Yang-Mills type gauged supergravity in which gauge kinetic terms are in the form of conventional Yang-Mills terms. The known class of Chern-Simons gauge groups that gives equivalent Yang-Mills type theory is of non-semisimple type [5]. Any Yang-Mills type Lagrangian can be rewritten in the Chern-Simons form by introducing two gauge fields and a compensating scalar for each Yang-Mills gauge field. This makes non-semisimple gauge groups more interesting in finding effective theories of string/M theory in three dimensions.

Some embeddings of three dimensional gauged supergravities into higher dimensions have appeared so far. These examples include $N = 2, 4, 8, 16$ gauged supergravities from reductions on spheres and Calabi-Yau manifold in [6, 7, 8, 9, 10, 11, 12] and recently various $N = 2$ theories from wrapped D3-branes of [13]. In this paper, we will give another example of this embedding namely $N = 10$ gauged supergravity with $SO(5) \ltimes \mathbf{T}^{10}$ gauge group. Due to the above mentioned equivalent between Chern-Simons and Yang-Mills type gauged supergravities, this should potentially describe $N = 5$ gauged supergravity in four dimensions with gauged group $SO(5)$ reduced on $S^1$. The latter has been constructed in [14]. It has been shown in [15] that the theory admits two $AdS_4$ critical points, an $N = 5$ supersymmetric point with $SO(5)$ gauge symmetry and a non-supersymmetric point with $SO(3)$ residual gauge symmetry. The theory has also been studied in the context of holographic superconductor in [16]. The non-supersymmetric critical point is perturbatively stable with all mass-squares above the BF-bound.
Unlike the four dimensional analogue which has maximally supersymmetric $AdS_4$ ground state, we will find that the reduced theory in three dimensions admits only a $\frac{1}{2}$-BPS domain wall as a vacuum solution. This is in contrast to compact and non-compact gaugings of the same theory studied in [17] that admits maximally supersymmetric $AdS_3$ critical points. The loss of supersymmetry after $S^1$ reduction has been pointed out in the context of non-semisimple gaugings in three dimensions in [10]. A general result on $S^1$ reduction of AdS spaces has been given in [18]. There are many known $\frac{1}{2}$-BPS domain walls in higher dimensional gauged supergravities, see for example [19, 20, 21, 22, 23, 24] as well as in lower dimensions, see [25] and [26] for three-and two-dimensional solutions. These domain walls are important in the context of the DW/QFT correspondence [27, 28, 29] which is a generalization to non-conformal field theories of the original AdS/CFT correspondence [30]. They are also useful in the study of domain wall/cosmology [31, 32, 33].

The paper is organized as follow. In section 2, we review the general structure of $N$ extended gauged supergravities in three dimensions including all relevant formulae and notations. The $SO(5) \times T^{10}$ gauged supergravity and the associated domain wall solution are discussed in section 3. We then discuss possible higher dimensional origin of the resulting theory from $S^1$ dimensional reduction of $N = 5$ $SO(5)$ gauged supergravity in four dimensions. We finally give some conclusions and comments in section 5. All details and explicit calculations are given in appendix A. In appendix B, we will explore possible non-semisimple gauge groups of $N = 9$ gauged supergravity in three dimensions.

2. $N = 10$ gauged supergravity in three dimensions with non-semisimple gauge groups

Before going to the detail of the construction, we briefly review the general structure of three dimensional gauged supergravities and apply it to the construction of $N = 10$ gauged supergravity with non-semisimple gauge group $SO(10) \times T^{10}$. We will keep the number of supersymmetry to be $N$ for conveniences and later set $N = 10$. In general, the matter coupled supergravity in three dimensions is in the form of a non-linear sigma model coupled to supergravity. For $N > 4$, supersymmetry demands that the scalar target manifold must be a symmetric space of the form $G/H$ in which $G$ and $H$ are the global symmetry group and its maximal compact subgroup, respectively [34]. In particular, for $N > 8$, supersymmetry determines the scalar manifold uniquely. In the present case of $N = 10$, the scalar manifold is given by the coset space $E_{6(-14)}/SO(10) \times U(1)$ which is a 32-dimensional Kahler manifold.
Coupling of the sigma model to N-extended supergravity requires the presence of \( N-1 \) almost complex structures \( f^{\dot{P}} \), \( \dot{P} = 2, \ldots, N \) on the scalar manifold. The tensors \( f^{IJ} = f^{[IJ]} \), \( I, J = 1, \ldots, N \), constructed by the relation

\[
f^{1P} = -f^{P1} = f^P, \quad f^{PQ} = f^{[P} f^{Q]}. \tag{2.1}
\]
generate the \( SO(N) \) R-symmetry in a spinor representation under which scalar fields transform. On symmetric scalar manifolds of the form \( G/H \), the maximal compact subgroup \( H = SO(N) \times H' \) contains the R-symmetry \( SO(N) \) and another compact subgroup \( H' \) commuting with \( SO(N) \). In \( N = 10 \) theory, the group \( H' \) is simply \( U(1) \). The \( G \)-generators \( t^M, M = 1, \ldots, \text{dim}G \), can be split into \( (T^{IJ}, X^\alpha) \) generating, respectively, \( SO(N) \times H' \) and non-compact generators \( Y^A \) corresponding to \( \text{dim}G - \text{dim}H \) scalars. The global symmetry group \( G \) is characterized by the following algebra

\[
\begin{align*}
[T^{IJ}, T^{KL}] &= -4\delta^{[I[K} T^{L]J]}, \\
[T^{IJ}, Y^A] &= -\frac{1}{2} f^{I[AB} Y_B, \\
[X^\alpha, X^\beta] &= f^{\alpha\beta\gamma} X^\gamma, \\
[X^\alpha, Y^A] &= h^{\alpha A} Y^B, \\
[Y^A, Y^B] &= \frac{1}{4} f^{AB} T^{IJ} + \frac{1}{8} C_{\alpha\beta} h^{\beta AB} X^\alpha. \tag{2.2}
\end{align*}
\]

The tensors \( f^{IJ} \) are related to \( SO(N) \) gamma matrices, \( \Gamma^I_{A\dot{A}} \) in which \( A \) and \( \dot{A} \) label spinor and conjugate spinor representations, respectively, by

\[
f^{IJ} = -\frac{1}{2} \Gamma^{IJ} = -\frac{1}{4} \left( \Gamma^I \Gamma^J - \Gamma^J \Gamma^I \right). \tag{2.3}
\]

\( C_{\alpha\beta} \) and \( f^{\alpha\beta\gamma} \) are \( H' \) invariant tensor and \( H' \) structure constants, respectively. The \( H' \) group is generated in the \( SO(N) \) spinor representation by matrices \( h^{\alpha A}_{\phantom{\alpha A}B} \). The coset manifold whose coordinates are given by \( d = \text{dim}(G/H) \) scalar fields \( \phi^i, i = 1, \ldots, d \) can be described by a coset representative \( L \). The usual formulae for a coset space are

\[
\begin{align*}
L^{-1} t^M L &= \frac{1}{2} \mathcal{V}^M_{\phantom{M}IJ} T^{IJ} + \mathcal{V}^M_{\phantom{M}A} X^\alpha + \mathcal{V}^M_{\phantom{M}A} Y^A, \tag{2.4} \\
L^{-1} \partial_i L &= \frac{1}{2} Q_i^{IJ} T^{IJ} + Q_i^A X^\alpha + c_i^A Y^A. \tag{2.5}
\end{align*}
\]

which will be useful later on. \( e_i^A \) is the vielbein on the scalar manifold while \( Q_i^{IJ} \) and \( Q_i^A \) are \( SO(N) \times H' \) composite connections. Scalar matrices \( \mathcal{V} \) will be used to define the moment maps below.

Gaugings of supergravities in various space-time dimensions are efficiently described in a \( G \)-covariant way by the so-called embedding tensor formalism [1]. In essence, the embedding tensor \( \Theta_{\mathcal{M}\mathcal{N}} \) is a symmetric gauge invariant tensor that acts as a projector
from the global symmetry group $G$ to a particular gauge group. Gauge covariant
derivatives describing the minimal coupling of the gauge fields $A^M_\mu$ to other fields also
involve the embedding tensor. For example, the covariant derivative on scalar fields is
given by
\[ D_\mu \phi^i = \partial_\mu \phi^i + g \Theta_{MN} A^M_\mu X^N_i \]  
where $X^N_i$ are Killing vectors generating isometries on the scalar manifold and $g$ is the
gauge coupling constant.

In order to define a viable gauging, the embedding tensor has to satisfy the so-called
quadratic constraint
\[ \Theta_{PL} f^{KL} (M \Theta_{N})^K = 0, \]  
which is the requirement that the gauge generators $\Theta_{MN}^N$ form a closed algebra, or
equivalently the gauge group is a proper subgroup of $G$. Furthermore, for supersym-
metry to be preserved in the gauging process, the embedding tensor needs to satisfy
the projection constraint
\[ P_{R_0} \Theta_{MN} = 0. \]  
This condition comes from supersymmetry, but it should be noted that the constraint
in this form is obtained by regarding the scalar manifold to be a symmetric space.

It is useful to introduce the T-tensor given by the moment map of the embedding
tensor by scalar matrices $V^M_A$, obtained from (2.4),
\[ T_{AB} = V^M_A \Theta_{MN} V^N_B. \]  
The T-tensor transforms under the maximal compact subgroup $H$ and consists of var-
dious components such as $T^{IJKL}$, $T^{IA}$ and $T^{AB}$. Since fermions transform under $H$,
the fermion couplings will be written in term of the T-tensor or linear combinations
of its components as we will see below. For any supersymmetric gauging, supersymmetry
requires only that the T-tensor satisfies the projection
\[ P_{\Box} T^{IJKL} = 0 \]  
where $\Box$ is the Riemann tensor-like representation of $SO(N)$. In the case of symmetric
scalar manifolds which are of interest in this paper, this constraint can be lifted to
the constraint on the embedding tensor given in (2.8) in which the $G$-representation
$R_0$, branched under $SO(N)$, contains $\Box$ representation of $SO(N)$. Any subgroup of
$G$ whose embedding tensor satisfies the above constraints is called admissible gauge

group.

In general, gaugings need some modifications to the original ungauged Lagrangian
by fermionic mass-like terms and a scalar potential, at order $g$ and $g^2$, respectively.
Also, the supersymmetry transformation rules need to be modified at order \( g \). In what follow, we will need the scalar potential and fermionic supersymmetry transformations. They are written in terms of the \( A_{1}^{I\bar{J}} \) and \( A_{2i}^{I\bar{J}} \) tensors which are in turn constructed from various components of the T-tensor

\[
A_{1}^{I\bar{J}} = -\frac{4}{N-2} T^{IM,JM} + \frac{2}{(N-1)(N-2)} \delta^{I\bar{J}} T^{MN,MN},
\]

\[
A_{2j}^{I\bar{J}} = \frac{2}{N} T^{I\bar{J},j} + \frac{4}{N(N-2)} f^{M(Im} T^{J)m}_{j} + \frac{2}{N(N-1)(N-2)} \delta^{I\bar{J}} f^{KL} m T^{KL,m}. \tag{2.12}
\]

The scalar potential is simply given by

\[
V = -\frac{4}{N} g^{2} \left( A_{1}^{I\bar{J}} A_{1}^{I\bar{J}} - \frac{1}{2} N g^{ij} A_{2i}^{I\bar{J}} A_{2j}^{I\bar{J}} \right). \tag{2.13}
\]

The metric \( g_{ij} \) on the target manifold is related to the vielbein by \( g_{ij} = e_{i}^{A} e_{j}^{A} \). We also note here that the quadratic constraint (2.7) can be written in terms of \( A_{1}^{I\bar{J}} \) and \( A_{2i}^{I\bar{J}} \) as

\[
2 A_{1}^{IK} A_{1}^{KJ} - N A_{2i}^{IK} A_{2j}^{K} = \frac{1}{N} \delta^{I\bar{J}} \left( 2 A_{1}^{KL} A_{1}^{KL} - N A_{2i}^{KL} A_{2j}^{KL} \right). \tag{2.14}
\]

The fermionic field content of the \( N \) extended supergravity in three dimensions consists of \( N \) gravitini \( \psi_{\mu}^{I} \) and \( d \) spin-\( \frac{1}{2} \) fields \( \chi^{iI} \). The latter is written in an overcomplete basis and subject to the projection constraint

\[
\chi^{iI} = \frac{1}{N} \left( \delta^{I\bar{J}} \delta_{j}^{i} - f^{I\bar{J}j} \right) \chi^{jJ} \tag{2.15}
\]

giving rise to \( d \) independent \( \chi^{iI} \) fields. The fermions \( \chi^{iI} \) can be redefined such that they transform in a conjugate spinor representation of \( SO(N) \) via

\[
\chi^{\dot{I}} = \frac{1}{N} e_{i}^{A} e_{\dot{I}}^{A} \chi^{iI}. \tag{2.16}
\]

The corresponding supersymmetry transformations are as follow:

\[
\delta \psi_{\mu}^{I} = \mathcal{D}_{\mu} \epsilon^{I} + g A_{1}^{I\bar{J}} \gamma_{\mu}^{\bar{J}} \epsilon^{J}, \tag{2.17}
\]

\[
\delta \chi^{iI} = \frac{1}{2} (\delta^{I\bar{J}} \delta_{j}^{i} - f^{I\bar{J}j}) \mathcal{D}^{j} \epsilon^{J} - g N A_{2i}^{I\bar{J}} \epsilon^{J}, \tag{2.18}
\]

where only relevant terms are given and

\[
\mathcal{D}_{\mu} \epsilon^{I} = \partial_{\mu} \epsilon^{I} + \frac{1}{4} \omega_{\mu}^{ab} \epsilon^{ab} + \partial_{\mu} \phi^{I} Q_{1}^{J} \epsilon^{J} + g \Theta_{\mathcal{M}N} A_{\mu}^{M} \gamma^{NI} \epsilon^{J}. \tag{2.19}
\]
Gauge groups of interest to us are non-semisimple groups of the form $G_0 \ltimes T^{\dim G}$. The translational symmetry $T^{\dim G}$ consists of $\dim G$ commuting generators which transform as an adjoint representation under $G_0$. This type of gauge groups gives rise to the on-shell equivalent Yang-Mills gauged supergravity coming from dimensional reductions of some higher dimensional theory. The $G_0 \ltimes T^{\dim G}$ gauge group whose generators are respectively $J^m$ and $T^m$, $m = 1, \ldots, \dim G$ is characterized by the following algebra

$$[J^m, J^n] = f_{mn}^k J^k, \quad [J^m, T^n] = f_{mn}^k T^k, \quad [T^m, T^n] = 0 \quad (2.20)$$

where $f_{mn}^k$ are $G_0$ structure constants. We will denote the $G_0$ and $T^{\dim G}$ parts of the gauge group by $a$ and $b$, respectively. As shown in [3], the corresponding embedding tensor consists of two parts, one with the coupling between $a$ and $b$ types $\Theta_{ab}$ and the other with the coupling between $b$ and $b$ types $\Theta_{bb}$. The full embedding tensor can be written as

$$\Theta = g_1 \Theta_{ab} + g_2 \Theta_{bb} \quad (2.21)$$

with $g_1$ and $g_2$ being the coupling constants. Supersymmetry constraint (2.8) may impose some relation on $g_1$ and $g_2$ such that eventually there is only one coupling. Both $\Theta_{ab}$ and $\Theta_{bb}$ are given by the Cartan-Killing form of $G_0$, $\eta_{mn}^{G_0}$, which is non-degenerate since $G_0$ is semisimple. The above information is sufficient for our discussion in this paper. The interested readers are invited to consult [4] and [5] for more a detailed discussion about three dimensional gauged supergravity with non-semisimple gauge groups.

3. $SO(5) \ltimes T^{10}$ gauged supergravity and $\frac{1}{2}$-BPS domain wall solution

In this section, we explicitly construct $N = 10$ gauged supergravity with $SO(5) \ltimes T^{10}$ gauge group. We begin with the scalar manifold $E_6(-14)/SO(10) \times U(1)$ and use $E_6$ generators given in [33] and [30]. The non-compact form $E_6(-14)$ is constructed by using the “Weyl unitarity trick”. We follow the same construction and notation as in [17] to which we refer the readers for more details.

The 78 generators of $E_6$ constructed in [30] are labeled by $c_i, i = 1, \ldots, 78$. The $SO(10)$ R-symmetry is generated by $c_i, i = 1, \ldots, 21, 30, \ldots, 36, 45, \ldots, 52, 71, \ldots, 78$ and $\tilde{c}_{33}$. We need to relabel these generators to the form of $T^{IJ}$ in our $SO(N)$ covariant formalism. This has already been done in [17], but we will repeat it in appendix A for convenience. The group $H' = U(1)$ is generated by $\tilde{c}_{70}$ whose definition and that of $\tilde{c}_{33}$
can be found in appendix A.

The non-compact generators can be identified as

\[
Y^A = \begin{cases} 
  ic_{A+21} & \text{for } A = 1, \ldots, 8 \\
  ic_{A+28} & \text{for } A = 9, \ldots, 16 \\
  ic_{A+37} & \text{for } A = 17, \ldots, 32 
\end{cases}
\]  

(3.1)

We can then use (2.2) to extract the tensors \( f^{IJ} \) whose components are computed by

\[
f^{IJ}_{AB} = -\frac{1}{3} \text{Tr} \left( [T^{IJ}, Y^A] Y^B \right).
\]  

(3.2)

Notice that the generators have normalizations \( \text{Tr}(T^{IJ}T^{IJ}) = -6 \) and \( \text{Tr}(Y^AY^A) = 6 \), no sum on \( IJ \) and \( A \).

We now construct generators of the gauge group \( SO(5) \ltimes T^{10} \). This group is embedded in \( USp(4,4) \subset E_6(-14) \). The maximal compact subgroup \( USp(4) \times USp(4) \subset USp(4,4) \) is identified as the \( SO(5) \times SO(5) \) subgroup of the R-symmetry \( SO(10) \). Recall that the 32 scalars transform as \( 16^+ + 16^- \) under \( SO(10) \times U(1) \). Under \( SO(5) \times SO(5) \), the scalars transform as

\[
16^+ + 16^- = (4,4)^+ + (4,4)^-.
\]  

(3.3)

We then identify \( SO(5) \) part of the gauge group as the diagonal subgroup \( SO(5)_{\text{diag}} \subset SO(5) \times SO(5) \) under which scalars transform as

\[
16^+ + 16^- = (4 \times 4)^+ + (4 \times 4)^- \\
= (1 + 10 + 5)^+ + (1 + 10 + 5)^-.
\]  

(3.4)

In this decomposition, we see that there are two singlets under \( SO(5)_{\text{diag}} \). The adjoint representation \( 10^+ \) and \( 10^- \) will be used to construct the translational generators of \( T^{10} \).

The explicit form of the corresponding gauge generators are as follow. The \( SO(5)_{\text{diag}} \) generators are given by

\[
J^{ij} = T^{ij} + T^{i+5,j+9}, \quad i, j = 1, \ldots, 5
\]  

(3.5)

while the \( T^{10} \) generators are found to be

\[
t^{ij} = T^{ij} - T^{i+5,j+5} + \tilde{Y}^{ij}, \quad i, j = 1, \ldots, 5
\]  

(3.6)

where \( \tilde{Y}^{ij} \) are given in appendix A.

The embedding tensor is of the form

\[
\Theta = g_1 \Theta_{ab} + g_2 \Theta_{bb}
\]  

(3.7)
where $\Theta_{ab}$ and $\Theta_{bb}$ are given by the Cartan-Killing form of $SO(5)$. The supersymmetry constraint requires $g_2 = 0$ meaning that there is no coupling among $T^{10}$ generators. This is similar to $N = 16$ and $N = 8$ theories with $SO(8) \ltimes T^{28}$ gauge group studied in \cite{11, 24}.

We are now in a position to study the scalar potential of the resulting gauged supergravity. Following the technique of \cite{37}, we begin with scalar fields which are singlets under the semisimple part of the gauge group, $SO(5)$. They are given by $1^\pm$ in (3.4) and correspond to the non-compact generators

\begin{align*}
Y_{s1} &= Y_3 - Y_5 - Y_{12} + Y_{16} + Y_{17} - Y_{18} + Y_{27} + Y_{29}, \\
Y_{s2} &= Y_4 + Y_8 + Y_{11} + Y_{13} + Y_{22} - Y_{23} + Y_{28} - Y_{32}.
\end{align*}

(3.8)

Accordingly, the coset representative is parametrized by

\[ L = e^{aY_{s1}} e^{bY_{s2}}. \]

(3.9)

Using the formulae (A.4) and (A.5), we can compute $A_{1IJ}$ and $A_{2I}^J$ by using a computer program Mathematica. The scalar potential is computed to be

\[ V = -6 e^{4(a-b)} (1 + e^{8b}) g^2 \]

(3.10)

where we have denoted $g_1$ simply by $g$. The presence of the $e^a$ factor implies that the potential has no critical point. We then expect the vacuum solution to be a domain wall.

To find a domain wall solution, we adopt the usual domain wall ansatz for the metric

\[ ds^2 = e^{2A} dx_{1,1}^2 + dr^2. \]

(3.11)

The supersymmetry transformation of $\chi^{II}$, $\delta \chi^{II} = 0$ from equation (2.18), gives the following equations

\begin{align*}
b' \gamma_r e^I + \frac{1}{2} g (1 - e^{4b}) e^{2(a-b)} e^I &= 0, \quad I = 1, \ldots, 5, \quad (3.12) \\
b' \gamma_r e^I - \frac{1}{2} g (1 - e^{4b}) e^{2(a-b)} e^I &= 0, \quad I = 6, \ldots, 10, \quad (3.13) \\
\phi' e^I - g \frac{e^{2(a+b)(1+e^{4b})}}{1 + e^{8b}} e^I &= 0, \quad I = 1, \ldots, 5, \quad (3.14) \\
\phi' e^I + g \frac{e^{2(a+b)(1+e^{4b})}}{1 + e^{8b}} e^I &= 0, \quad I = 6, \ldots, 10, \quad (3.15)
\end{align*}

where we have used $'$ to denote the derivative $\frac{d}{dr}$ and $\phi' A^I = \frac{1}{6} \text{Tr} (L^{-1} L' Y^A)$. We will now impose the projection conditions $\gamma_r e^I = -\epsilon^I$ for $I = 1, \ldots, 5$ and $\gamma_r e^I = \epsilon^I$ for $I =$
$\epsilon^I$ has two real components. The projectors then reduce the supersymmetry by a fraction of $\frac{1}{2}$. With these two projectors, we end up with two independent equations

$$b' = \frac{1}{2} g(1 - e^{4b})e^{2(a-b)}, \quad (3.16)$$

$$a' = -g e^{2(a+b)(1+e^{4b})} \left(1 + e^{8b}\right). \quad (3.17)$$

The supersymmetry variation of the gravitini $\psi^{I}_{\mu}$, $\delta \psi^{I}_{\mu} = 0$ from equation (2.17) after using the above projectors, gives rise to

$$e^{4b} = 1, \quad (3.18)$$

$$A' = 2g \left(1 + e^{4b}\right) e^{2(a-b)} \quad (3.19)$$

where we have used the spin connection $\omega_{\hat{\mu}\hat{\nu}} = A' \delta_{\hat{\mu}}^\hat{\nu}$ with $\hat{\mu}, \hat{\nu} = 0, 1$.

We see from (3.18) that supersymmetry demands $b = 0$. Equation (3.16) is now trivially satisfied, and equation (3.17) becomes

$$a' + e^{2a}g = 0. \quad (3.20)$$

The solution is easily obtained to be

$$a = -\frac{1}{2} \ln \left(2gr + C_1\right) \quad (3.21)$$

where $C_1$ is an integration constant. Substituting into equation (3.19) gives

$$A' = 4ge^{2a} = \frac{4g}{C_1 + 2gr} \quad (3.22)$$

whose solution is, with another integration constant $C_2$,

$$A = C_2 + 2 \ln \left(2gr + C_1\right). \quad (3.23)$$

As in other solutions of this type, the residual supersymmetry is generated by the Killing spinors given by $\epsilon^i = e^{\hat{i}} \epsilon^{i}_{0\pm}$, $i = 1, \ldots, 5$ with the constant spinors $\epsilon^i_{0\pm}$ satisfying $\gamma_\nu \epsilon^i_{0\pm} = \pm \epsilon^i_{0\pm}$. The full symmetry of this solution is $ISO(1,1) \times SO(5)$ with the unbroken $N = (5,5)$ Poincare supersymmetry in notation of the dual two-dimensional field theory.

The two integration constants $C_1$ and $C_2$ can be set to zero by shifting the coordinate $r$ and rescaling the coordinates $x^\mu$. We can also write down the solution in the
form of warped $AdS_3$ by introducing the new coordinate $\rho = -\frac{1}{4g^2 r}$ in term of which the metric becomes

$$ds^2 = \frac{1}{(4g^2 \rho)^2} \left( dx_{1,1}^2 + d\rho^2 \right). \quad (3.24)$$

We end this section by considering subgroups of $SO(5) \ltimes T^{10}$ namely $SO(4) \ltimes T^6$ and $(SO(3) \ltimes T^3) \times (SO(2) \ltimes T^1) \sim U(2) \ltimes T^4$. It can be checked that both of them are not admissible.

4. Higher dimensional origin

In this section, we discuss higher dimensional origin of the $SO(5) \ltimes T^{10} N = 10$ gauged supergravity constructed in the previous section. By the general result of [3], this theory is on-shell equivalent to the $SO(5)$ Yang-Mills gauged supergravity which can be obtained from $S^1$ reduction of $N = 5$ gauged supergravity in four dimensions with $SO(5)$ gauge group. The four dimensional theory has been constructed in [14] and can be obtained as a truncation of the maximal $N = 8$ gauged supergravity. In the notation of [14], the field content of this theory contains one graviton $e^i_M$ or $g_{MN}$, five gravitini $\psi^i_M$, eleven spin-$\frac{1}{2}$ fields $\chi^{ijk}$ and $\chi^{678}$, ten scalars $\phi^i$ and $\phi_i$ living in the coset space $SU(5,1)/U(5)$ and ten vector fields $A^i_M$ being $SO(5)$ gauge fields. Here, $M, N = 0, 1, 2, 3$ and $a, b = 0, 1, 2, 3$ are four dimensional space-time and tangent space indices respectively while $i, j = 1, \ldots, 5$ are $SU(5)$ indices except for $A^i_M$ which transform in the adjoint representation of $SO(5)$.

If we reduce this theory on $S^1$ along the $x^3$ direction, we find the following fields in three dimensions. The metric $g_{MN}$ gives the non-dynamical three dimensional metric $g_{\mu\nu}$, the graviphoton $g_{\mu 3}$ and a scalar $g_{33}$. The $SO(5)$ gauge fields result in the three dimensional gauge fields of the same gauge group $A^i_\mu$ and ten scalars $A^i_M$ transforming in the adjoint representation of $SO(5)$. Finally, the ten scalars $(\phi^i, \phi_i)$ obviously become the three dimensional scalars.

A spinor in four dimensions give rise to two spinors in three dimensions. We then obtain ten gravitini $\psi^i_\mu$ from $\psi^i_M$ and ten spin-$\frac{1}{2}$ fields $\psi^i_3$. There are additional $20 + 2$ spin-$\frac{1}{2}$ fields from the reduction of $\chi^{ijk}$ and $\chi^{678}$, respectively. In three dimensions, the metric and gravitino do not have any dynamics. We then find 32 fermionic on-shell degrees of freedom from $(\psi^i_3, \chi^{678}, \chi^{ijk})$. We can also dualize $A^i_\mu$ and $g_{\mu 3}$ to $10 + 1$ scalars. All together, we end up with 32 scalars from $(\phi^i, \phi_i, g_{33}, g_{\mu 3}, A^i_\mu, A^i_3)$. This is the same as in $N = 10$ gauged supergravity.

We give $SO(5)_{gauge}$ representations of the reduced fields in table [3] from which we have omitted the non-dynamical fields $g_{\mu\nu}$ and $\psi^i_\mu$. We have kept $\phi^i$ and $\phi_i$ separately.
Table 1: Representations of three dimensional fields resulted from $S^1$ reduction of $N = 5$ gauged supergravity in four dimensions.

to emphasize their four dimensional origin. We now consider the representation of the 32 scalars in $E_6(-14)/SO(10) \times U(1)$ coset space under the $SO(5)$ part of the gauge group. Recall that under $SO(10) \times U(1)$, the scalars transform as $16^+ + 16^-$. Under $SO(10) \times U(1) \supset SU(5) \times U(1) \supset SO(5)$ in which the $U(1)$ is the $U(1)$ subgroup of $U(5) \subset SO(10)$, we find

$$16^+ + 16^- \rightarrow (1^-_5 + 5^-_3 + 10^-_1)^+ + (1^-_5 + 5^-_3 + 10^-_1)^-$$
$$\rightarrow (1 + 5 + 10) + (1 + 5 + 10)$$

(4.1)

We find perfect agreement with table 1. Reference [38] is very useful in this decomposition. In the formalism of [4], the fermions $\chi^A$ transform as $\overline{10}^+ + 10^-$ under $SO(10) \times U(1)$. Similar decomposition gives $2 \times (1 + 5 + 10)$ under $SO(5)$ gauge group. This is again the representations obtained from $S^1$ reduction shown in table 1.

The result of [39] suggests that three dimensional supergravity with $E_6$ coset manifold can be obtained from dimensional reduction on a torus, $S^1$ in the present case, of a supergravity theory with $A_5$ coset manifold in four dimensions. Reference [39] consider only maximally non-compact $E_6$ and other types Lie groups. The result here should provide an example of a non-maximally non-compact $E_6$ ($E_6(-14)$) coset obtained from a non-maximally non-compact $A_5$ $SU(5,1)$ coset in four dimensions. Furthermore, the general formulae for toroidal reductions given in the appendix of [39] should also be applicable in this case.

5. Conclusions and discussions

In this paper, we have constructed $N = 10$ $SO(5) \ltimes T^{10}$ gauged supergravity in three
dimensions. We have found that the resulting theory admits a 1/2-BPS domain wall as a vacuum solution. The solutions preserves \( N = (5, 5) \) Poincare supersymmetry in two dimensions with ten supercharges. The solution is similar to the domain wall from the \( S^7 \) compactification of type II string theory discussed in [1]. This solution is the vacuum solution of the maximal \( N = 16 \) \( SO(8) \times T^{28} \) gauged supergravity. The solution given here provides an example of a domain wall in non-maximal gauged supergravity and might be useful in the DW/QFT correspondence as well as its applications.

We have also discussed possible higher dimensional origin of this theory. This is given by \( S^1 \) reduction of \( N = 5 \) \( SO(5) \) gauged supergravity in four dimensions. We have found that the spectrum of the reduction matches with the constructed three dimensional gauged supergravity. If the \( N = 5 \) four dimensional theory is reduced on \( S^1/\mathbb{Z}_2 \), it could give rise to \( N = 5 \) gauged supergravity in three dimensions. Indeed, the latter in general has scalar manifold \( USp(4, k)/USp(4) \times USp(k) \) [32]. We have seen that the \( SO(5) \times T^{10} \) gauge group is embedded in \( USp(4, 4) \subset E_{6(-14)} \). We then expect that \( N = 5 \) \( SO(5) \) gauged supergravity in four dimensions reduced on \( S^1/\mathbb{Z}_2 \) should give \( N = 5 \) \( SO(5) \times T^{10} \) gauged supergravity in three dimensions with scalar manifold \( USp(4, 4)/USp(4) \times USp(4) \) containing 16 scalars. It turns out that the latter theory admits \( SO(5) \times T^{10} \) gauge group. The details will be reported in subsequent work [1]. Unlike the \( N = 10 \) theory, the \( N = 5 \) truncation admits maximally supersymmetric \( AdS_3 \) vacuum solution. This truncation should be similar to the case of \( N = 8 \) \( SO(8) \times T^{28} \) gauged supergravity with \( SO(8, 8)/SO(8) \times SO(8) \) scalar manifold studied in [25]. This theory is a truncation of \( N = 16 \) \( SO(8) \times T^{28} \) gauged supergravity with scalar manifold \( E_{8(8)}/SO(16) \).

Due to the similar structure as in the above examples, we would like to briefly discuss the case of \( N = 12 \) gauged supergravity. The scalar manifold is the 64-dimensional quaternionic manifold \( E_{7(-5)}/SO(12) \times SU(2) \). The gauge group should be \( SO(6) \times T^{15} \) embedded in \( SU(4, 4) \subset E_{7(-5)} \). The \( SO(6) \) is again identified as \( SO(6)_{\text{diag}} \subset SO(6) \times SO(6) \subset SO(12) \). The 64 scalars transform under \( SO(12) \times SU(2) \) as \( (32, 2) \) and under \( SO(6) \times SO(6) \times SU(2) \) as \( (4, 4) + (4, 4), 2) \). Then, under the \( SO(6) \) part of the gauge group, we find the representation for scalars \( (4 \times 4 + 4 \times 4), 2) = (1 + 15 + 1 + 15, 2) \). The non-compact generators in the \( 15 \) should combine with \( SO(6) \times SO(6) \) generators to form the \( T^{15} \) part of the gauge group. The fermions transform as \( (32, 2) \) under \( SO(12) \times SU(2) \) and \( ((4, 4) + (4, 4), 2) \) under \( SO(6) \times SO(6) \times SU(2) \). Under \( SO(6) \), they transform as \( (10 + 6 + 10 + 6, 2) \).

We now consider \( S^1 \) reduction of \( N = 6 \) \( SO(6) \) gauged supergravity in four dimensions which is also a truncation of \( N = 8 \) \( SO(8) \) gauged supergravity [42]. The bosonic fields are \( (g_{MN}, \phi_{AB}, \phi_{AB}, A_{M}^{AB}, A_{M}) \) where the 30 scalars \( (\phi_{AB}, \phi_{AB}) \) live in the coset space \( SO^*(12)/U(6) \) and \( A, B = 1, \ldots, 6 \), see [42] for more detail. The fermionic fields

\[ -12 - \]
are given by \((\psi_A^M, \chi^A, \chi^{ABC})\). After \(S^1\) reduction, the dynamical bosonic fields are given by \((g_{\mu 3}, g_{33}, \phi^{AB}, \phi_{AB}, A_\mu, A_3, A^{AB}_\mu, A^{AB}_3)\) transforming as \((1 + 1 + 15 + 15 + 1 + 1 + 15 + 15)\) under \(SO(6)\) gauge group. After dualizing the vector fields, we end up with 64 scalars with correct \(SO(6)\) representations as in \(N = 12\) gauged supergravity. The reduced dynamical fermionic fields are \((\psi_3^A, \chi^{ABC}, \chi^A)\) transforming under \(SO(6)\) as \(2 \times (6 + 10 + 10 + 6)\) which are indeed the same as those in \(N = 12\) theory. The factor of 2 comes from the fact that a four dimensional spinor gives two three dimensional spinors.

Finally, similar to the discussion in the \(N = 5\) case, we expect that the \(S^1/\mathbb{Z}_2\) reduction should give \(N = 6\) \(SO(6) \times T^{15}\) gauged supergravity on three dimensions with scalar manifold \(SU(4,4)/S(U(4) \times U(4))\) whose compact and non-compact gauge groups have been explored in \([13]\). The possibility of non-semisimple gauge groups is under investigation \([41]\).

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A. Useful formulae and details

In this appendix, we give some details of \(N = 10\) gauged supergravity with \(SO(5) \times T^{10}\) gauge group constructed in the main text. First of all, the \(SO(10)\) R-symmetry
generators $T^{ij}$ are explicitly given by

\[
T^{12} = c_1, \quad T^{13} = -c_2, \quad T^{23} = c_3, \quad T^{34} = c_6, \quad T^{14} = c_4, \quad T^{24} = -c_5,
\]
\[
T^{15} = c_7, \quad T^{25} = -c_8, \quad T^{35} = c_9, \quad T^{45} = -c_{10}, \quad T^{56} = -c_{15}, \quad T^{16} = c_{11},
\]
\[
T^{26} = -c_{12}, \quad T^{46} = -c_{14}, \quad T^{36} = c_{13}, \quad T^{17} = c_{16}, \quad T^{27} = -c_{17}, \quad T^{47} = -c_{19},
\]
\[
T^{37} = c_{18}, \quad T^{67} = -c_{21}, \quad T^{57} = -c_{20}, \quad T^{78} = -c_{36}, \quad T^{18} = c_{30}, \quad T^{28} = -c_{31},
\]
\[
T^{48} = -c_{33}, \quad T^{38} = c_{32}, \quad T^{68} = -c_{35}, \quad T^{58} = -c_{34}, \quad T^{29} = -c_{46}, \quad T^{19} = c_{45},
\]
\[
T^{49} = -c_{48}, \quad T^{39} = c_{47}, \quad T^{69} = -c_{50}, \quad T^{59} = -c_{49}, \quad T^{89} = -c_{52}, \quad T^{79} = -c_{51},
\]
\[
T^{1,10} = -c_{71}, \quad T^{2,10} = c_{72}, \quad T^{3,10} = -c_{73}, \quad T^{4,10} = c_{74}, \quad T^{5,10} = c_{75},
\]
\[
T^{6,10} = c_{76}, \quad T^{7,10} = c_{77}, \quad T^{8,10} = c_{78}, \quad T^{9,10} = -\tilde{c}_{53}
\]

where $\tilde{c}_{53}$ and $\tilde{c}_{70}$ are defined by

\[
\tilde{c}_{53} = \frac{1}{2} c_{53} + \frac{\sqrt{3}}{2} c_{70} \quad \text{and} \quad \tilde{c}_{70} = -\frac{\sqrt{3}}{2} c_{53} + \frac{1}{2} c_{70}.
\]

Also, notice a typo in the sign of $T^{9,10}$ in \[17\].

The $\tilde{Y}^{ij}$ part of the translational generators $T^{10}$ is constructed from the following non-compact generators

\[
\tilde{Y}^{12} = \frac{1}{2} (Y_3 - Y_{12} + Y_{17} + Y_{29} + Y_5 - Y_{16} + Y_{18} - Y_{27}),
\]
\[
\tilde{Y}^{13} = \frac{1}{2} (Y_2 + Y_{14} + Y_{21} - Y_{26} - Y_1 + Y_{15} - Y_{19} - Y_{25}),
\]
\[
\tilde{Y}^{14} = \frac{1}{2} (Y_{31} - Y_7 - Y_6 - Y_{30} - Y_9 + Y_{10} + Y_{20} - Y_{24}),
\]
\[
\tilde{Y}^{15} = \frac{1}{2} (Y_{15} - Y_{14} + Y_{25} - Y_{26} - Y_1 - Y_2 + Y_{19} + Y_{21}),
\]
\[
\tilde{Y}^{23} = \frac{1}{2} (Y_1 + Y_2 + Y_{15} - Y_{14} + Y_{19} + Y_{21} - Y_{25} + Y_{26}),
\]
\[
\tilde{Y}^{24} = \frac{1}{2} (Y_{10} + Y_9 - Y_{30} - Y_{31} + Y_6 - Y_7 - Y_{20} - Y_{24}),
\]
\[
\tilde{Y}^{25} = \frac{1}{2} (Y_2 - Y_1 - Y_{25} - Y_{26} - Y_{14} - Y_{15} + Y_{19} - Y_{21}),
\]
\[
\tilde{Y}^{34} = \frac{1}{2} (Y_8 - Y_4 - Y_{11} - Y_{28} + Y_{13} - Y_{32} + Y_{22} + Y_{23}),
\]
\[
\tilde{Y}^{35} = \frac{1}{2} (Y_{18} + Y_{17} - Y_{12} + Y_{27} - Y_{29} - Y_{16} - Y_5 - Y_3),
\]
\[
\tilde{Y}^{45} = \frac{1}{2} (Y_8 + Y_4 - Y_{11} - Y_{28} - Y_{13} + Y_{32} - Y_{23} + Y_{22}).
\]
This choice is of course not unique.

The scalar matrices for the moment maps are given by

\[
\begin{align*}
\mathcal{V}_{ij,IJ}^a &= -\frac{1}{6} \text{Tr}(L^{-1}j^{ij}LT^{IJ}), \\
\mathcal{V}_{ij,IJ}^b &= -\frac{1}{6} \text{Tr}(L^{-1}t^{ij}LT^{IJ}), \\
\mathcal{V}_{ij,IA}^a &= \frac{1}{6} \text{Tr}(L^{-1}j^{ij}LY^A), \\
\mathcal{V}_{ij,IA}^b &= \frac{1}{6} \text{Tr}(L^{-1}t^{ij}LY^A),
\end{align*}
\]

from which the T-tensor follows

\[
\begin{align*}
T^{IJ,KL} &= g \left( \mathcal{V}_{ij,IJ}^a \mathcal{V}_{ij,KL}^b + \mathcal{V}_{ij,IJ}^b \mathcal{V}_{ij,KL}^a \right) \\
T^{IJ,A} &= g \left( \mathcal{V}_{ij,IJ}^a \mathcal{V}_{ij,A}^b + \mathcal{V}_{ij,IJ}^b \mathcal{V}_{ij,A}^a \right)
\end{align*}
\]

Using these together with (2.11), (2.12) and (2.13), we can find the tensors \( A_{1IJ} \) and \( A_{2IJ} \) as well as the scalar potential.

**B. Non-semisimple gauging of \( N = 9 \) gauged supergravity in three dimensions**

We will consider \( N = 9 \) gauged supergravity in three dimensions. The corresponding scalar manifold is given by the 16-dimensional \( F_4(-20)/SO(9) \) coset space. Some vacua of the compact and non-compact gaugings of this theory have been studied in [44]. In this appendix, we will explore the possibilities of non-semisimple gauge groups which are crucial for embedding the theory in higher dimensions. Notice that the construction of \( E_6 \) given in [36] is based on the \( F_4 \) group given in [35]. We can simply remove the last 26 matrices \( c_i, i = 53, \ldots, 78 \) from \( E_6 \) to get the group \( F_4 \) generated by \( c_i, i = 1, \ldots, 52 \) as has been used in [44]. All 52 matrices are effectively \( 26 \times 26 \) matrices since all elements in the last row and last column are zero.

The \( SO(9) \) R-symmetry generators are \( T^{IJ} \) in (A.1) with \( I, J = 1, \ldots, 9 \), and non-compact generators are the first 16 generators of (3.1), \( Y^A, A = 1, \ldots, 16 \). In the case of \( F_4(4)/USp(6) \times SU(2) \) which is a scalar manifold of \( N = 4 \) theory studied in [45], \( SO(4) \ltimes T^6 \) can be gauged consistently with supersymmetry by the embedding of \( SO(4) \ltimes T^6 \) in \( SO(5, 4) \subset F_4(4) \). In the present case, the embedding of \( SO(3) \ltimes T^3 \) in \( USp(2, 2) \subset USp(4, 2) \times SU(2) \subset F_4(-20) \) should be possible.

To identify generators of this group, we first consider the \( SO(4) \ltimes T^6 \) subgroup of the
$SO(5) \times T^{10}$ in section 3. Obviously, the $SO(4)$ part is generated by $J^{ij}$, $i, j = 1, \ldots, 4$. We then consider $\tilde{Y}^{ij}$ with $i, j = 1, \ldots, 4$. It can be verified that by removing $Y_{17}$ to $Y_{32}$ from $\tilde{Y}^{ij}$, the resulting generators, see appendix A,

$$
\begin{align*}
\tilde{Y}^{12} &= \frac{1}{2} (Y_3 - Y_{12} + Y_5 - Y_{16}), \\
\tilde{Y}^{13} &= \frac{1}{2} (Y_2 + Y_{14} - Y_1 + Y_{15}), \\
\tilde{Y}^{14} &= \frac{1}{2} (Y_{10} - Y_7 - Y_6 - Y_{30} - Y_9), \\
\tilde{Y}^{23} &= \frac{1}{2} (Y_1 + Y_2 + Y_{15} - Y_{14}), \\
\tilde{Y}^{24} &= \frac{1}{2} (Y_{10} + Y_9 + Y_6 - Y_7), \\
\tilde{Y}^{34} &= \frac{1}{2} (Y_8 - Y_4 - Y_{11} + Y_{13})
\end{align*}
$$

(B.1)

still transform in the adjoint representation of $SO(4)$. It turns out that when combined into $t^{ij}$, the resulting generators do not commute. Therefore, it is not possible to find $SO(4) \ltimes T^6$ subgroup of $F_4(-20)$. On the other hand, we can form two $SU(2)_{\pm}$ subgroups from these generators by introducing the self-dual and anti-self-dual $SO(4)$ generators

$$
\begin{align*}
J^1_+ &= J^{12} + J^{34}, & J^2_+ &= J^{13} - J^{24}, & J^3_+ &= J^{14} + J^{23}, \\
t^1_+ &= t^{12} + t^{34}, & t^2_+ &= t^{13} - t^{24}, & t^3_+ &= t^{14} + t^{23}
\end{align*}
$$

(B.2)

and

$$
\begin{align*}
J^1_- &= J^{12} - J^{34}, & J^2_- &= J^{13} + J^{24}, & J^3_- &= J^{14} - J^{23}, \\
t^1_- &= t^{12} - t^{34}, & t^2_- &= t^{13} + t^{24}, & t^3_- &= t^{14} - t^{23}.
\end{align*}
$$

(B.3)

It can be readily verified that each set of generators forms $SO(3) \ltimes T^3 \sim SU(2) \ltimes T^3$ algebra but generators $t^a_\pm$ from the two sets do not commute with each other. Although this subgroup can be embedded in $F_4(-20)$, it is not admissible namely it cannot be gauged in a way that is consistent with supersymmetry. Embedding in higher dimensions aside, it seems to be difficult (if possible) to find non-semisimple gaugings of the $N = 9$ theory.

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