Anomaly-Free Brane Worlds in Seven Dimensions

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Abstract

We present an orbifold compactification of the minimal seven-dimensional supergravity. The vacuum is a slice of $AdS_7$ where six-branes of opposite tension are located at the orbifold fixed points. The cancellation of gauge and gravitational anomalies restricts the gauge group and matter content on the boundaries. In addition anomaly cancellation fixes the boundary gauge couplings in terms of the gravitational constant, and the mass parameter of the Chern-Simons term.
1 Introduction

The geometry of extra spacetime dimensions has recently played a prominent role in theories beyond the Standard Model. A generic feature of these models is that one extra dimension is compactified on a line segment (or orbifold) where boundary worlds exist at the end points. In particular, the Standard Model gauge and matter fields are normally assumed to be confined on the boundaries while gravity propagates in the bulk. It is this geometric separation of the two boundary worlds, interacting only gravitationally in the bulk, which has provided new insight into the hierarchy problem [1, 2].

In the generic brane world scenario there is no restriction on the possible gauge group structure on the boundary. Clearly such a restriction on the gauge group of a higher-dimensional theory could help to explain the particle content of the low-energy world. The one notable example is the Horava-Witten (HW) theory in eleven dimensions [3], where the cancellation of gauge and gravitational anomalies restricts the gauge group on the ten-dimensional boundaries to be $E_8$. The bulk eleven-dimensional (11d) theory is then further interpreted to be the strongly coupled limit of the ten-dimensional $E_8 \times E_8$ heterotic string theory [3].

Apart from ten dimensions, gravitational anomalies also exist in six (and two) dimensions [4]. Thus, in this work we shall show that in seven-dimensional brane worlds the gauge group structure and matter content is similarly restricted on the six-dimensional boundaries by gauge and gravitational anomalies. Unlike the HW theory where there is a unique gauge group, we will show that many more possibilities exist in the seven-dimensional theory, which are not necessarily dimensional reductions of the HW theory. It should be stressed that, as in the HW case, supersymmetry is a central element in our construction. This is because supersymmetry dictates the possible fields allowed in the bulk as well as on the boundaries, and furthermore restricts their possible couplings.

The vacuum of the seven-dimensional theory will be a slice of $AdS_7$, where six-branes of opposite tension are located at the orbifold fixed points. This leads to a localized gravity, tensor and hypermultiplet. Moreover, anomaly cancellation will require the addition of extra vector, tensor and hypermultiplets on the boundaries. The boundary theory must then have locally supersymmetric couplings to the seven-dimensional bulk supergravity multiplet. We will find that the gauge couplings are fixed by the anomaly cancellation in terms of the bulk gravitational coupling, and a mass parameter of the bulk Chern-Simons term.

Furthermore, by the AdS/CFT correspondence [5], our seven-dimensional bulk theory is dual to a strongly coupled six-dimensional (6d) conformal field theory (CFT), in much the same way that the HW theory is dual to the strongly coupled $E_8 \times E_8$ heterotic string theory. This dual correspondence provides a way to further understand the properties of strongly coupled six-dimensional conformal field theories.
Let us consider the minimal $\mathcal{N} = 2$ seven-dimensional (7d) gauged supergravity [6, 7, 8]. The gravity multiplet in this theory consists of the graviton $g_{MN}$, an antisymmetric three-form $A_{MNK}$, an $SU(2)$ triplet of vectors $A_{Mi}$, a scalar $\phi$, and the $SU(2)$ pseudo-Majorana gravitino $\psi_M$ and spinor $\chi^i$. A dual version where the three-form is replaced by a two-form is discussed in [9, 10]. The capital Latin indices $M,N = 0,1,\ldots,6$ are 7d spacetime indices, while $i,j = 1,2$ label the $SU(2)$ R-symmetry group. The bosonic part of the 7d action is

$$S_{\text{bosonic}} = \frac{1}{\kappa^2} \int d^7x \sqrt{-g} \left[ \frac{1}{2} R - \frac{\sigma^4}{48} F_{MNPQ}^2 - \frac{\sigma^2}{4} F_{MN}^j F_{MN}^i \right] - \frac{1}{2} (\partial_M \phi)^2$$

$$+ \frac{i}{48\sqrt{2}} F_{MNPQ} \left( F_{KLi}^j R^i_j - \frac{2ig}{3} \text{tr}(A_K A_L A_R) \right) \epsilon^{MNPQKLR}$$

$$+ 60(m - \frac{2}{5} h\sigma^4)^2 - 10(m + \frac{8}{5} h\sigma^4)^2 + \frac{h}{36} \epsilon^{KLMNPQR} F_{KLMN} A_{PQR} \right], \tag{1}$$

where $m = -g\sigma^{-1}/(5\sqrt{2})$, $\sigma = \exp(-\phi/\sqrt{5})$, and $g$ is the $SU(2)$ gauge coupling. It can be shown that the 7d supergravity Lagrangian is invariant under $x_7 \to -x_7$ provided that [11]

$$A_{MNP} \to -A_{MNP}, \quad A_{Mi}^j \to -A_{Mi}^j, \quad h \to -h, \quad m \to -m. \tag{2}$$

Since the $\mathbb{Z}_2$ transformation, $x_7 \to -x_7$ is a symmetry of the theory, we will consider the compactification down to six dimensions on the orbifold $S^1/\mathbb{Z}_2$. Then the only fields which survive at the orbifold fixed points are the $\mathbb{Z}_2$ singlets, and form the following 6d multiplets

$$(g_{\mu\nu}, A_{\mu\nu}^\pm, \psi_\mu), \quad \text{gravity}$$

$$(A_{\mu\nu}, \phi, \chi_i^l), \quad \text{tensor}$$

$$(A_i^j, \xi, \psi^i), \quad \text{hypermultiplet} \tag{3}$$

where $A_{\mu\nu}^\pm = A_{\mu\nu}^\pm, A_i^j = A_i^j, \xi = g_{77}, \psi_i = \psi_i^j$, and $\psi_\mu^i = \psi_{-\mu}^i$ are left-handed symplectic Majorana-Weyl fermions while $\chi^i = \chi_i^l$, and $\psi^i = \psi_{-i}^l$ are right-handed.

The compactification of the 7d supergravity theory on an orbifold results in a chiral $\mathcal{N} = (0,1)$ 6d theory with the massless spectrum [3] provided that under $x_7 \to -x_7$ Eq. (2) is satisfied. However, the relations (2) do not necessarily respect the supersymmetry transformations at the boundaries. For example, since the parameters $h$ and $m$ are odd, the variation of the kinetic energy terms will produce $\delta$-function terms. In order to make the truncated theory on the orbifold supersymmetric we must introduce six-branes at the orbifold fixed points with specific boundary potentials. This is very similar to the five-dimensional supersymmetric
Randall-Sundrum model\footnote{12, 13}, where supersymmetry requires the introduction of brane tensions. If we introduce the boundary potential term\footnote{11}

\[ S_0 = \int d^6x \int dy \sqrt{-g} 20(m - \frac{2}{5} h\sigma^4) \left[ \delta(y) - \delta(y - \pi R) \right], \quad (4) \]

then the complete action $S_7 + S_0$ will be supersymmetric, where $S_7$ is the full 7d action including fermionic terms, and $y$ denotes the seventh coordinate $x_7$.

The supersymmetric vacuum is now the one which satisfies the Killing equations $\delta \psi_{Mi} = \delta \chi_i = 0$. Assuming that all bulk fields are zero except for the scalar field $\phi$, we find that

\[ \langle \sigma \rangle = \left( \frac{g}{8\sqrt{2h}} \right)^{\frac{1}{5}}. \quad (5) \]

In this vacuum the bulk action becomes\footnote{11}

\begin{align*}
S_{\text{bulk}} &= S_7 + S_{(0)} + S_{(\pi R)}, \quad (6) \\
S_7 &= \int d^6x \int dy \sqrt{-g} \left[ \frac{1}{2} M^5 R - \Lambda_7 \right], \quad (7) \\
S_{(y^*)} &= \int d^6x \sqrt{-g_6} \left[ \mathcal{L}_{(y^*)} - \Lambda_{(y^*)} \right], \quad (8)
\end{align*}

where $g_6$ is the induced metric on the six-brane located at $y^*$. The cosmological constants are given by

\[ \Lambda_7 = -15M^5k^2; \quad \Lambda_{(0)} = -\Lambda_{(\pi R)} = 10M^5k, \quad (9) \]

where

\[ k = \left( \frac{h\sigma^4}{16} \right)^{\frac{1}{5}}. \quad (10) \]

The Einstein equations for the combined bulk and boundary action\footnote{6} can be solved to obtain a seven-dimensional Randall-Sundrum solution

\[ ds^2 = e^{-2k|y|} dx_6^2 + dy^2, \quad (11) \]

where $0 \leq y \leq \pi R$ and $k$ is the AdS curvature scale which is given by\footnote{10}. Note that supersymmetry automatically guarantees the fine-tuning conditions\footnote{9} required to obtain the Randall-Sundrum solution. This leads to a slice of $AdS_7$, where the 6d gravity multiplet is localized on the UV brane at $y^* = 0$, while the tensor and hypermultiplet are localized on the IR brane at $y^* = \pi R$. 

3
3 Anomaly cancellation with a boundary

The orbifold compactification of the 7d supergravity theory has resulted in a theory
with a localized gravity multiplet as well as a tensor, and a hypermultiplet. However,
unlike the five-dimensional case where arbitrary matter can be added to the
boundaries \[12\], in the slice of \(AdS_7\) the 6d fermions of the vector, tensor and gravity
multiplets will in general lead to gravitational and gauge anomalies. The cancellation
of these anomalies will restrict the possible gauge groups and matter content on the
boundary.

In six dimensions the anomalies are formally described by an 8-form, \(I_8\). For \(n_V\)
vector multiplets, \(n_T\) tensor multiplets and \(n_H\) hypermultiplets the requirement for
the cancellation of the irreducible \(\tr R^4\) term in \(I_8\), where \(R\) is a curvature 2-form,
leads to the condition \[14\]
\[ n_V - n_H - 29n_T + 273 = 0 . \]
(12)
From the dimensional reduction of the bulk gravity multiplet there will be one tensor
multiplet and one hypermultiplet in the 6d theory. However, this theory by itself is
anomalous and we are forced to introduce extra boundary fields. In particular we
will be interested in introducing vector multiplets on the boundary which will also
produce gauge anomalies. The gauge group should be \(G_1 \times G_2\), where each \(G_i\) factor is
localized at the two fixed points (for simplicity we will only consider semisimple \(G\)).
Notice that the anomaly need not be equally distributed between the fixed points as
in the HW case. In addition the anomaly eight-form \(I_8\) should satisfy
\[ \frac{\partial^2 I_8}{\partial \tr F_1^2 \partial \tr F_2^2} = 0 , \]
(13)
where \(F_1, F_2\) are the gauge field strengths of the \(G_1 \times G_2\) gauge group. This condition
simply means that there is no charged matter in the bulk and that the only interaction
between the two six-branes is purely gravitational. In this way we will be able to
cancel the anomaly locally on each boundary.

Consider the case of \(n_T = 1\). Assuming that the irreducible part of the anomaly \(\tr R^4\) is cancelled, then the remaining reducible part (normalized as in \[15\]) is given by
\[ I_8 = (\tr R^2)^2 + \frac{1}{6} \tr R^2 \left( X_1^{(2)} + X_2^{(2)} \right) - \frac{2}{3} \left( X_1^{(4)} + X_2^{(4)} \right) , \]
(14)
where \(X_i^{(n)} = \Tr F_i^n + \sum_i n_i \tr_i F_i^n\), and \(\Tr, \tr_i\) are traces in the adjoint and the
\(R_i\) representation, respectively, whereas \(n_i\) is the number of hypermultiplets in the
representation \(R_i\). The anomaly \[14\] can be cancelled by the Green-Schwarz mechanism \[16\] provided that it can be factorized into the form
\[ I_8 = \left( \tr R^2 + u_i \sum_i \tr F_i^2 \right) \left( \tr R^2 + v_i \sum_i \tr F_i^2 \right) , \]
(15)
where \( u_i, v_i \) are constants. This ensures that at the massless level the theory is anomaly free.

However, from the 7d orbifold perspective the 6d anomaly must be distributed between the two fixed points. The bulk topological Chern-Simons term plays a crucial role in cancelling the anomaly by a local Green-Schwarz mechanism as occurs in the 11d HW theory. Thus, by writing \( I_8 = I_8^{(1)} + I_8^{(2)} \) and demanding the local factorization

\[
I_8^{(i)} = (c_i \text{tr} R^2 + a_i \text{tr} F_i^2) \left( \text{tr} R^2 + b_i \text{tr} F_i^2 \right)
\]

(16)

where \( a_i, b_i, c_i \) are constants and \( c_1 + c_2 = 1 \), the two terms in the sum \( I_8 \) vanish by a local Green-Schwarz mechanism at each orbifold fixed point [17]. It can be shown [16] that the factorization (16) is indeed possible as long as \( \alpha_i \text{tr} F_i^4 = 0 \). For \( \alpha_i \neq 0 \), this occurs for all the irreps of \( E_8, E_7, E_6, F_4, G_2, SU(3), SU(2), U(1) \), for the 28 of \( \text{Sp}(4) \) and \( SU(8) \), and all the irreps of \( \text{SO}(2n) \) with highest weight \((f_1, f_2, f_1, -f_2, 0, ..., 0)\) in the Gel’fand-Zetlin basis [18].

Let us now present some examples which illustrate the possible matter content that is allowed on the six-branes from anomaly cancellation. In the case where there is only one tensor multiplet in the 6d theory \( (n_T = 1) \), arising from the dimensional reduction of the bulk theory, we are lead to the constraint

\[
n_H = n_V + 244.
\]

(17)

As discussed earlier we will assume that on each boundary there is a gauge group \( G_i \). In addition under \( G_1 \times G_2 \) let us suppose that the total number of hypermultiplets consists of the following representations

\[
n_1(d_{F_1}, 1) + n_2(1, d_{F_2}) + (n_S + 1)(1, 1),
\]

(18)

where \( d_{F_i} \) is the dimension of the fundamental representation of the group \( G_i \), and \( n_{1,2}, n_S \) are the numbers of each representation. Note that we have automatically included the extra singlet hypermultiplet (or radion multiplet) arising from the dimensionally reduced bulk theory. Thus, assuming the constraint (17) is satisfied together with (15) and (16), we find the following solutions for \( G_1 = G_2 = G \):

| \( G \times G \) | \( n_1 + n_2 \) | \( n_S \) |
|----------------|----------------|------|
| \( G_2 \times G_2 \) | 20 | 131 |
| \( F_4 \times F_4 \) | 10 | 87 |
| \( E_6 \times E_6 \) | 12 | 75 |
| \( E_7 \times E_7 \) | 8 | 61 |

In particular we see that not only the gauge groups, but also the number of hypermultiplet generations is restricted on the boundaries. For example in the \( E_6 \times E_6 \)
case, if one boundary contains 3 generations of the fundamental 27 then the other boundary must have 9 generations. There is also an \((n_1, n_2, n_5)\) solution \((2, 7, 156)\), and similar exceptions exist for the other gauge groups. It is also possible to have two different gauge groups distributed between the fixed points. These include \(E_8 \times E_7, E_8 \times G_2, E_8 \times F_4, E_7 \times G_2, E_7 \times F_4, E_6 \times G_2, \) and \(F_4 \times G_2\). These exceptions and other possibilities will be presented in Ref. [11]. Our solutions differ from the usual compactifications of the HW theory because in HW compactifications there is matter charged under both local gauge groups [19, 20]. This is true even in compactifications of weakly coupled string theory [21].

Note also that in the \(n_T = 1\) case, there are no solutions with \(SU(n), SO(n)\) or \(Sp(n)\) gauge groups because the cancellation of the quartic Casimir does not allow one to simultaneously satisfy (15) and (16). Nevertheless, as will be shown in [11] such solutions can exist for \(n_T > 1\). However, in this case one must employ the generalized Green-Schwarz mechanism of Ref. [22].

Finally note that the six-dimensional theory may still be ill defined due to non-perturbative anomalies [23, 24, 18]. Global anomalies exist as long as \(\pi_6(G)\) is non-trivial. In our case only the gauge group \(G_2\) may be plagued by global anomalies since \(\pi_6(G_2) = \mathbb{Z}_3\). In particular, with \(n_F\) fundamentals of \(G_2\), the condition for the absence of global anomalies is \(n_F = 1 \mod 3\) [25]. Thus, for the \(G_2 \times G_2\) case, the absence of non-perturbative anomalies further restricts the values of \((n_1, n_2)\) in the above table.

4 Bulk-boundary action

The addition of vector, tensor and hypermultiplets propagating on the boundaries implies that there must exist locally supersymmetric couplings of the six-dimensional multiplets to the seven-dimensional supergravity multiplet propagating in the bulk. In general, the boundary action can be written in the form \(S_{\text{boundary}} = S_0 + S_{YM} + S_H + S_T\), where \(S_0\) is given in [4] and \(S_{YM}, S_H, S_T\) are the boundary actions for the vector, hyper and tensor multiplets, respectively.

Let us consider first the addition of vector multiplets on the boundary. One can show [11] that the combined action \(S_{\text{bulk}} + S_{\text{boundary}}\) is locally supersymmetric up to fermionic bilinear terms where

\[
S_{YM} = -\frac{1}{\lambda^2} \int d^6x \sqrt{-g} \left\{ \frac{\sigma^{-2}}{4} F_{\mu\nu}^\alpha F^{\alpha\mu\nu} + \frac{1}{2} \lambda^a \Gamma^\mu D_\mu \lambda^a + \frac{\sigma^{-1}}{4} \bar{\psi}_\mu \Gamma^{\nu\rho} \Gamma^\mu \lambda^a F_{\nu\rho}^\alpha \\
+ \frac{\sigma^{-1}}{2\sqrt{3}} \lambda^a \Gamma^{\mu\nu} \chi - \frac{\sigma^{-2}}{24\sqrt{2}} \lambda^a \Gamma^{\mu\nu\rho} \lambda^a F_{\mu\nu\rho\gamma} + \frac{i\sigma}{2\sqrt{2}} \lambda^a \Gamma^\mu F_{\mu\gamma} \lambda^a \right\},
\]  

(19)
and the supersymmetry transformations are
\begin{align}
\delta A^a_\mu &= \frac{1}{2} \sigma \bar{\epsilon} \Gamma_\mu \lambda^a , \\
\delta \lambda^a &= -\frac{1}{4} \Gamma^{\mu\nu} F^a_{\mu\nu} \epsilon .
\end{align}

As in the HW theory one requires the modification of the Bianchi identity for $F_{\mu\nu\rho\tau}$. The modified Bianchi identity has the effect of changing the Chern-Simons term in $S_{\text{bulk}}$ into a Green-Schwarz term. The Green-Schwarz term precisely cancels the anomalous variation of the effective action for six-dimensional Weyl fermions provided that the boundary gauge coupling, $\lambda$, satisfies
\begin{equation}
\lambda^2 = 8\kappa \sqrt{\frac{3\pi^3 h}{\gamma} ,}
\end{equation}
where $\gamma$ is a constant defined by $X_i^{(4)} = \gamma (\text{tr} F^2)^2$. This relation is similar to that found in the HW theory, except that now there is an extra dependence on the topological mass parameter, $h$. In the HW theory the coefficient of the Chern-Simons term is fixed by supersymmetry whereas in seven dimensions the theory is supersymmetric up to the arbitrary parameter, $h$.

Similarly we can introduce hypermultiplets on the boundary. Under the supersymmetry transformations
\begin{align}
\delta \varphi^\alpha &= \frac{1}{2} \sigma^{-1/2} V^\alpha_{\alpha_1} \epsilon^i \xi^X , \\
\delta \zeta^X &= \frac{1}{2} \sigma^{1/2} V_{\alpha_1} \Gamma^\mu \partial_\mu \varphi^\alpha \epsilon^i ,
\end{align}
the locally supersymmetric boundary action for neutral hypermultiplets is
\begin{equation}
S_H = \int d^6 x \sqrt{-g} \left[ -\frac{1}{2} g_{\alpha\beta} (\varphi^i \partial_\mu \varphi^\alpha \partial_\mu \varphi^\beta - \frac{1}{2} \bar{\xi} Y \Gamma^\mu D_\mu \xi Y^i \bar{\xi} Y \Gamma^\mu \partial_\mu \varphi^\alpha \chi^i \\
+ \frac{\sigma^{1/2}}{2} \bar{\psi}_\mu \Gamma^\nu \Gamma^\mu \partial_\mu \varphi^\alpha V_{\alpha_1} \xi^X \zeta^Y + \frac{\sigma^{1/2}}{24 \sqrt{2}} \bar{\xi} Y^i \bar{\xi} \Gamma^{\mu\nu\rho} \xi^Y F^{(i)}_{\mu\nu\rho} \right] ,
\end{equation}
where $\varphi^\alpha (\alpha, \beta = 1, ..., 4n_H)$, and $\zeta^X (X, Y = 1, ..., 2n_H)$ are the scalars and fermions of the $n_H$ hypermultiplets, respectively, and $g_{\alpha\beta}$ is the metric of the scalar manifolds. Similarly, as in the case of the vector multiplets the Bianchi identity for $F_{\mu\nu\rho\tau}^i$ must be modified which results in a correction to the supersymmetry transformation of $F_{\mu\nu\rho\tau}^i$. This is crucial in showing that the scalar manifold is quaternionic, as required by 6d $\mathcal{N} = (0, 1)$ local supersymmetry. Details will be presented elsewhere.
5 Conclusion

We have presented a new class of models with a boundary, where the gauge group structure on the boundary is determined by the cancellation of gauge and gravitational anomalies. The vacuum of the bulk theory is a slice of $AdS_7$ with localized gravity. Anomaly cancellation also places constraints on the possible boundary matter, and determines the boundary gauge coupling in terms of the bulk gravitational constant, and the mass parameter of the Chern-Simons term. By the AdS/CFT correspondence our seven-dimensional brane world is dual to a six-dimensional conformal field theory. Much like the five-dimensional counterpart [27, 28, 29], this conformal field theory is defined with a cutoff, and couples to gravity. Boundary fields on the UV (IR) brane are identified as fundamental (composite) states in the CFT, and the strong coupling regime of this six-dimensional theory is described by our seven-dimensional solution.

The six-dimensional boundary theories also have phenomenological interest. For example, the hierarchy problem is naturally solved and the gauge group structure on the boundaries can contain the Standard Model gauge group. Moreover, there are also possible monopole compactifications on $S^2$, like those considered in [30], which give rise to chiral four-dimensional $\mathcal{N} = 1$ theories. The low energy particle content would then be fixed by anomaly cancellation (see also [31]), and the hierarchy problem is explained by the warped bulk. These issues as well as cosmological implications remain to be investigated.

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