On Asymmetric Orbifolds
and the $D = 5$ No-modulus Supergravity

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Abstract

We examine whether any type II asymmetric orbifolds have the same massless spectrum as the dimensional reduction of $D = 5$ simple supergravity, which, besides the eleven-dimensional supergravity, is the only known supergravity above four dimensions with no moduli. We attempt to construct such models by further twisting the orbifolds which yield $D = 4$, $N = 4$ pure supergravity to find that, unfortunately, none of the models have that spectrum. We provide supergravity arguments explaining why this is so. As a by-product, we list all possible momentum-winding lattices that give $D = 4$, $N = 4$ pure supergravity.
1 Introduction

Besides the eleven-dimensional supergravity, $D = 5$ simple supergravity is the only known supergravity above four dimensions with no moduli, and in fact they have many similarities [1]-[5]. Because $D = 5$ simple supergravity contains no scalar fields and particularly no dilaton, it cannot be realized as low-energy theories of any perturbative string compactifications. It is rather mysterious why such a supergravity exists in only five dimensions and how it arises in string theories.

In view of the similarity of the Lagrangians, one would naturally speculate that some $D = 4, N = 2$ string compactification might lead to the $D = 5$ theory in its strong coupling, just as type IIA string theory becomes M theory [6, 7] in this limit. This idea of finding isolated points in string/M theory moduli space is an old one and some works based on it can be found in the literature [8]-[11]. The underlying motivation is to find mechanisms of stabilizing the moduli (see e.g. [12]) through the search of any ‘$D = 5$ analogue of M theory’.

We will, therefore, look for $D = 4, N = 2$ string compactifications with the same massless field content as the dimensional reduction of the $D = 5$ supergravity. Note that this dimensionally reduced theory has only one $N = 2$ gravity multiplet and one vector multiplet but no hypermultiplet, while any Calabi-Yau or symmetric orbifold compactifications necessarily contain the universal hypermultiplet, of which the dilaton is a member, hence it cannot be realized by them. Thus, in this paper we will consider asymmetric orbifolds [13].

Our strategy is the following: We first construct asymmetric orbifolds with $D = 4, N = 4$ pure supergravity as their low-energy limits, following ref.[11]. We then examine if any residual symmetries of the invariant lattices fix the shift vector and, at the same time, reduce the number of massless vectors to 1/3. Finally, if there is one, we check if any new massless moduli appear in the twisted sector. We will show that unfortunately none of the models lead to the desired supergravity in this framework. We will also give supergravity arguments explaining why this is so.

Although our attempt is unsuccessful, it provides a no-go statement and will be a step toward the understanding the string-theory origin of the $D=5$ no-modulus supergravity. As a by-product, we list all (in this framework) possible momentum-winding lattices that give $D = 4, N = 4$ pure supergravity by using the classification of conjugacy classes of

\footnote{We concentrate on type II compactifications because $D = 5$ simple supergravity is obtained [14] from a consistent truncation of the $D = 11$ supergravity (and hence the $D = 4$ theory from IIA theory).}
Weyl group elements.

In Section 2, we recall the asymmetric orbifold constructions of $N = 4$ pure supergravity, and give a list of possible momentum-winding lattices. Section 3 is devoted to the details of the constructions. In Section 4, we provide supergravity arguments suggesting that any orbifold does not seem likely to realize the desired supergravity as its low-energy theory. Finally we summarize our conclusions in Section 5.

2 $N = 4$ pure supergravity from asymmetric orbifolds

We start with a compactification of type II theories on a six-dimensional torus whose metric is given by the matrix of some simply-laced semi-simple Lie algebra $g$. Our convention is $\alpha' = 2$ and the radii $R = 1$ so that it is a self-dual torus. Turning on an appropriate $B$ field, the Narain lattice $\Gamma_{6,6}(g) = (p_L, p_R)$ takes the form

$$p_L, p_R \in \Lambda_W(g), \quad p_L - p_R \in \Lambda_R(g),$$

(2.1)

where $\Lambda_R(g) (\Lambda_W(g))$ is the root (weight) lattice of a simply-laced (in order for the Cartan matrix to be interpreted as a metric) Lie algebra $g$. The necessary and sufficient condition for this (with the metric $G_{ij}$ of the torus being the Cartan matrix) is that $(B_{ij} - G_{ij})/2$ have integer entries, where $i, j$ are the coordinates of the torus. Then $\Gamma_{6,6}(g)$ is invariant under the independent actions on the left- and the right-lattices of the Weyl reflection group $W(g)$ (T-duality). One can use these automorphisms of the lattice to twist the model on this special background. We define the action of $g_L, g_R \in W(g)$ on a state with definite momenta $|p_L, p_R>$ as

$$|p_L, p_R> \rightarrow e^{2\pi i (p_L v_L - p_R v_R)} |g_L p_L, g_R p_R>,$$

(2.2)

where, $(v_L, v_R)$ are called the shift vectors. The oscillators are transformed similarly.

Let us now construct asymmetric orbifolds whose low-energy limits are $N = 4$ pure supergravity [16, 17] following ref.[1]. We consider abelian asymmetric orbifolds twisted by a group generated by a pair of Weyl group elements $(g_L, g_R)$. To obtain an $N = 4$ theory whose supersymmetries come only from the right-moving sector, we choose $g_R = 1$ and $g_L \in SO(6)$ but $\notin SU(3)$. Conjugacy classes of Weyl group elements have been classified in the mathematical literature [18, 19, 20]. The statement is that each simple Lie algebra $g$ has a particular set of Weyl group elements known as the ‘primitive elements’ (see e.g. [19] for the precise definition), and any conjugacy class of Weyl group elements of $g$ corresponds to a primitive element of some regular subalgebra of $g$, that is, the
subalgebra whose Dynkin diagram is obtained by removing one node from the extended Dynkin diagram of $\mathfrak{g}$ (or the one by repeating this procedure on the Dynkin diagrams so obtained).

In the rank = 6 case, eigenvalues of a Weyl group element $w$ of order $N$ are of the form

$$\{\epsilon^{r_1}, \epsilon^{r_2}, \epsilon^{r_3}, \epsilon^{N-r_1}, \epsilon^{N-r_2}, \epsilon^{N-r_3}\},$$  \hspace{1cm} (2.3)

where $r_i$ and $N - r_i$ are positive integers and $\epsilon = e^{2\pi i/N}$. The condition for $w$ not to lie in $SU(3)$ is that any of $r_1 \pm r_2 \pm r_3 \neq 0 \text{ mod } N$. Furthermore, if $N$ is even, modular invariance require that the sum $\sum_i r_i$ must be even [21]. An exhaustive search using the list of the primitive elements shows that these requirements leave only four possible Weyl group elements listed in Table.

| Weyl group elements | Eigenvalues | Order | Momentum-winding lattice ($\mathfrak{g}$) |
|---------------------|-------------|-------|---------------------------------|
| $E_6(a_1)$          | $(\epsilon^1, \epsilon^2, \epsilon^3, \epsilon^4, \epsilon^5, \epsilon^6)$ | 9     | $E_6$                           |
| $A_4 \oplus A_2$    | $(\epsilon^3, \epsilon^6, \epsilon^9, \epsilon^{12}, \epsilon^5, \epsilon^{10})$ | 15    | $A_4 \oplus A_2$                |
| $D_4(a_1) \oplus A_2$ | $(\epsilon^3, \epsilon^6, \epsilon^9, \epsilon^4, \epsilon^8)$ | 12    | $D_4 \oplus A_2$                |
| $A_2 \oplus A_1 \oplus A_1 \oplus A_1$ | $(\epsilon^2, \epsilon^4, \epsilon^3, \epsilon^3)$ | 6     | $A_2 \oplus A_1 \oplus A_1 \oplus A_1 \oplus A_1$ |

Table: List of possible Weyl group elements. $\epsilon = e^{2\pi i/N}$, where $N$ is the order of the element. $a_1$ is a label distinguishing different primitive elements.

In either case, no larger semi-simple simply-laced Lie algebra has the algebra in the last column as its regular subalgebra. Therefore, the momentum lattice is uniquely determined for each Weyl group element in this table.

Since these twists yield the fields of $D = 4$, $N = 4$ pure supergravity already in the untwisted sector, we will take the shift vectors which are not in the dual of the invariant lattices [11] so as to avoid extra massless fields from the twisted sector. Let us examine each case in some detail.

2.1 $E_6(a_1)$

The first $E_6$ case is a known example [22, 11]. In the $A_2 \oplus A_2 \oplus A_2$ basis the metric and the $B$ field are

\[
G_{ij} = \begin{bmatrix}
2 & 1 & 1 & 2 \\
1 & 2 & 1 & 2 \\
2 & 1 & 1 & 2 \\
1 & 2 & 1 & 2
\end{bmatrix}, \quad B_{ij} = \begin{bmatrix}
3 & 3 & 3 \\
-3 & 3 & 3 \\
-3 & 3 & 3 \\
3 & 3 & 3
\end{bmatrix}.
\hspace{1cm} (2.4)
\]
We can take
\[
g_L = (\epsilon^1, \epsilon^2, \epsilon^4, \epsilon^5, \epsilon^7, \epsilon^8), \quad v_L = 0,
g_R = 1, \quad v_R = \left(\frac{1}{9}, \frac{1}{9}, -\frac{2}{9}, \frac{1}{9}, \frac{1}{9}, -\frac{2}{9} \right)
\] (2.5)
with \(\epsilon = e^{2\pi i/9}\). The shift vectors are expressed in terms of the weights of its regular subalgebra \(A_2 \oplus A_2 \oplus A_2\), where the two simple roots of \(A_2 = SU(3)\) are denoted by \((1, -1, 0)\) and \((0, 1, -1)\) in this notation. This twist gives a \(D=4, N=4\) multiplet as the massless field already in the untwisted sector. For the twisted sector to have no massless fields \(v_R\) has been so chosen that any of \(jv_R\) for \(j = 1, \ldots, 8\) does not lie in \(\Lambda_W(E_6)\). Since \(\frac{1}{2}v_R^2 = \frac{1}{9}\), the level-matching condition is satisfied.

2.2 \(A_4 \oplus A_2\)

The next example is the case \(A_4 \oplus A_2\). We found the following two shift vectors:
\[
g_L = (\epsilon^3, \epsilon^6, \epsilon^9, \epsilon^{12}, \epsilon^5, \epsilon^{10}), \quad v_L = 0,
g_R = 1, \quad v_R = \left(\frac{9}{15}, -\frac{4}{15}, -\frac{1}{15}, 0, 0, \frac{1}{5}, -\frac{1}{5}, 0\right)
\] (2.6)
and
\[
g_L = (\epsilon^3, \epsilon^6, \epsilon^9, \epsilon^{12}, \epsilon^5, \epsilon^{10}), \quad v_L = 0,
g_R = 1, \quad v_R = \left(\frac{1}{15}, -\frac{4}{15}, -\frac{3}{15}, 0, 0, 0, 0\right)
\] (2.7)
with \(\epsilon = e^{2\pi i/15}\). The shift vectors are again written as vectors in the weight spaces. The notation for the \(A_2\) piece is the same as above, and \((1, -1, 0, 0, 0), (0, 1, -1, 0, 0), (0, 0, 1, -1, 0)\) and \((0, 0, 0, 1, -1)\) correspond to the simple roots for the \(A_4\) piece. Clearly in both cases \(15v_R \in \Lambda_R(A_4 \oplus A_2) \subset \Lambda_W(A_4 \oplus A_2)\), and \(jv_R\) for any \(j = 1, \ldots, 14\) does not belong to the weight lattice. The level-matching condition is also satisfied.

2.3 \(D_4(a_1) \oplus A_2\)

We can take
\[
g_L = (\epsilon^3, \epsilon^3, \epsilon^9, \epsilon^9, \epsilon^4, \epsilon^8), \quad v_L = 0,
g_R = 1, \quad v_R = \left(\frac{4}{12}, \frac{1}{12}, \frac{1}{12}, 0, \frac{1}{12}, \frac{1}{12}, -\frac{2}{12}\right)
\] (2.8)
with \(\epsilon = e^{2\pi i/12}\). The first four entries of \(v_R\) is a vector in the \(D_4\) weight space, where we take \((1, -1, 0, 0), (0, 1, -1, 0), (0, 0, 1, -1)\) and \((0, 0, 1, 1)\) as its simple roots.
2.4 \( A_2 \oplus A_1 \oplus A_1 \oplus A_1 \oplus A_1 \)

Similarly,
\[
g_L = (\epsilon^2, \epsilon^4, \epsilon^3, \epsilon^3, \epsilon^3), \quad v_L = 0,
\]
\[
g_R = 1, \quad v_R = \left(1, \frac{1}{6}, \frac{1}{6}, \frac{1}{6} ; \frac{1}{6}, \frac{1}{6} \right)
\]
with \( \epsilon = e^{2\pi i/6} \) are a solution. The fundamental weight of \( A_1 \) is denoted by \( \sqrt[3]{2} \).

3 \( N = 2 \) models by a further twist

We will now attempt to construct \( N = 2 \) models which have a graviton, two vectors and two scalars as their only massless bosonic fields. As already mentioned in Introduction, we try to construct such models by accompanying a further twist in the right sector to reduce the number of vectors (and at the same time the number of supersymmetries), while keeping the shift vector to avoid extra moduli from the twisted sector. We only consider the case where the order of the new right twist is a divisor of the order of the left twist since otherwise the order of the whole group would change and the level-matching property of the original shift vector would be lost.

3.1 \( E_6(a_1) \)

The extended Dynkin diagram of \( E_6 \) has a \( \mathbb{Z}_3 \) outer automorphism, and the shift vector is invariant under the \( \mathbb{Z}_3 \) transformation generated by
\[
\begin{bmatrix}
    0 & 1 & 0 \\
    0 & 0 & -1 \\
    -1 & 0 & 0
\end{bmatrix},
\]
where each block represents the action on one of the three \( A_2 \) weight spaces. This belongs to \( SU(2) \), and hence breaks one half of supersymmetries. Moreover, six vectors coming from the NS-NS massless states \( \psi^\mu_{-1/2} \bar{\psi}^i_{-1/2} | \text{vac} > \) (\( \mu = 2,3 \) and \( i = 4, \ldots, 9 \) label the noncompact and compact coordinates,) are reduced to two by the permutation. Thus the twist
\[
g_L = (\epsilon^1, \epsilon^2, \epsilon^5, \epsilon^7, \epsilon^8), \quad v_L = 0,
\]
\[
g_R = \text{eq.(3.10)}, \quad v_R = \left(\frac{1}{9}, \frac{1}{9}, -\frac{2}{9} ; \frac{1}{9}, \frac{1}{9}, -\frac{2}{9} \right)
\]
precisely yields the desired massless fields in the untwisted sector. However, although the shift vector is preserved, it turns out that the twisted sector also have extra massless
fields. This can been seen as follows: Because of the new $Z_3$ right twist, the invariant lattice of (3.11) changed from the original $\Lambda_R(E_6)$ to

$$\{(v, v, -v) \mid v \in \Lambda_R(A_2)\}.$$  

(3.12)

Then its dual lattice becomes

$$\left\{ \frac{1}{3}(w, w, -w) \mid w \in \Lambda_W(A_2) \right\},$$

(3.13)

to which the shift vector $v_R$ belongs. Thus extra massless fields arise from the twisted sector.

3.2 $A_4 \oplus A_2$

The situation is worse. The normal lattice of the shift vector has neither $Z_3$ nor $Z_5$ symmetry (Recall that the order of the left twist is 15 in this case.), but has only $Z_2$ symmetry, for both cases found in the last section. Thus the untwisted sector has more than two vector fields and the model is not the one we wanted to have.

3.3 The other cases

A similar analysis shows that neither of the choices (2.8), (2.9) leads to a model with one gravity and one vector multiplet. The resulting models have extra massless fields in the twisted sector and/or the untwisted sector. This is because, again, either the invariant lattices of the $N = 4$ models do not have enough residual symmetries or the new less dense invariant lattices give rise to extra massless states in the twisted sector. Although we have explicitly confirmed this only for the shift vectors given in Section 2, the situation appears to be the same for the other cases since, after all, choosing a shift vector not lying in the weight lattice $\Lambda_W(g)$ is somewhat contradictory to the demands of preserving enough symmetries.

4 Supergravity arguments

Finally, let us consider from the point of view of supergravity actions why we did not get any asymmetric orbifold model which realizes the dimensional reduction of $D = 5$ simple supergravity. In general, any type II orbifold compactification has a graviton, a dilaton and an anti-symmetric two-form field $B$ as massless fields coming from the NS-NS sector.
On the other hand, the dimensionally reduced bosonic action takes the form [2]

\[
2\kappa_4 S_4 = \int d^4x E^{(4)} \left( R^{(4)} - \frac{3}{2} \partial_\mu \rho \partial^\mu \rho - \frac{3}{2} \rho^{-2} \partial_\mu A \partial^\mu A - \frac{1}{4} \rho^3 B^{(KK)}_{\mu\nu} B^{(KK)\mu\nu} \right.
\]

\[\left. - \frac{3}{4} \rho F^{(4)}_{\mu\nu} F^{(4)\mu\nu} - \frac{3}{4} E^{(4)-1} \epsilon^{\mu\nu\tau\sigma} F_{\mu\nu} F_{\tau\sigma} A, \right) \tag{4.14}
\]

where the five-dimensional fields have been parameterized as

\[
E^{(5)\hat{\rho}} = \begin{bmatrix} \rho^{-1/2} E^{(4)\alpha} & \rho B^{(KK)}_{\hat{\rho}} \\ 0 & \rho \end{bmatrix},
\]

\[
A^{(5)}_{\hat{\rho}} = \begin{cases} A_{\mu} & (\mu = 0, \ldots, 3), \\ A_{4} = A, \end{cases}
\tag{4.15}
\]

and \( F^{(4)}_{\mu\nu} = F_{\mu\nu} + B^{(KK)}_{\mu\nu} A, F'_{\mu\nu} = 2\partial_{[\mu} A'_{\nu]}, A'_{\mu} = A_{\mu} - B^{(KK)}_{\mu} A \). \( \rho \) is naturally identified as the dilaton, and therefore the other scalar field \( A \) must be the dual to the \( B \) field if (4.14) is a low-energy action of an orbifold. However, if the dual transformation is applied on \( A \) (This can be done if we take the vector \( A_{\mu} \) as fundamental rather than the Kaluza-Klein invariant one \( A'_{\mu} \)), the dualized action becomes non-polynomial in \( A_{\mu} A_{\mu} \) and, unlike type IIA supergravity, is hard to be regarded as a string effective action. This is in contrast with \( D = 4, N = 4 \) pure supergravity, which can be successfully dualized [23] to give a part of the heterotic string action. Also, the minimally coupled \( D = 4, N = 2 \) supergravity [24] of the type in ref. [25], obtained [16] by a consistent truncation of \( N = 4 \) pure supergravity, can be easily dualized, in which the two vector fields enter in the action symmetrically. In our case, however, the two vectors are asymmetric, carrying different \( SO(2) \) charges [3]. Therefore the scalar field \( A \) cannot be interpreted as an axion.

Where does this scalar come from? We can gain insight into this question by examining how \( D = 5 \) simple supergravity is obtained by a consistent truncation from \( D = 11 \). It is known that the eleven-dimensional supergravity (bosonic) action [28]

\[
2\kappa_{11} S_{11} = \int d^{11}x E^{(11)} \left( R^{(11)} - \frac{1}{48} F^{(11)}_{MNPQ} F^{(11)MNPQ} \right.
\]

\[\left. - \frac{1}{124} E^{(11)-1} \epsilon_{M_1 \cdots M_{11}} A^{(11)}_{M_1 M_2 M_3} F^{(11)}_{M_4 \cdots M_7} F^{(11)}_{M_8 \cdots M_{11}} \right) \tag{4.16}
\]

is consistently truncated to \( D = 5 \) simple supergravity

\[
2\kappa_5 S_5 = \int d^5x E^{(5)} \left( R^{(5)} - \frac{3}{4} F^{(5)}_{\mu\nu} F^{(5)\mu\nu} \right.
\]

\[\left. - \frac{1}{4} E^{(5)-1} \epsilon_{\mu_1 \cdots \mu_5} A^{(5)}_{\mu_1} F^{(5)}_{\mu_2 \mu_3} F^{(5)}_{\mu_4 \mu_5} \right) \tag{4.17}
\]
by setting \[ E^{(11)A}_M = \begin{bmatrix} E^{(5)\hat{\mu}}_{\hat{\nu}} & 0 \\ 0 & \delta_{\hat{m}\hat{\alpha}} \end{bmatrix}, \]
\[ A^{(11)}_{MNP} = \begin{cases} A_{\hat{\mu}56} = A_{\hat{\mu}78} = A_{\hat{\mu}910} = A_{\hat{\nu}}, \\ 0 & \text{otherwise}, \end{cases} \tag{4.18} \]
where \( M, N, \ldots \) and \( A \) are the eleven-dimensional curved and local-Lorentz indices, and \( \hat{\mu}, \hat{\nu} \ldots \) and \( \hat{\alpha} \) are the corresponding five-dimensional indices. \( \hat{m} \) and \( \hat{a} \) are those of the flat six-dimensional internal space \( T^6 \). The reduced \( D = 4 \) action is obtained by further parameterizing the five-dimensional fields as (4.15).

Now if the fifth direction \( M = 4 \) is thought of as tangent to the \( S^1 \) direction which is compactified to yield type IIA theory, then \( \rho \) is certainly the dilaton, while \( A = A_{456} = A_{478} = A_{4910} \) are the components of the NS-NS two-form \( B \) field with both indices tangent to \( T^6 \). The \( B_{\mu\nu} \) components are truncated and do not appear in the spectrum. This is the reason why the dualization of \( A \) does not give a neat result. Another difference from the massless spectrum of the orbifold is that both of two vector fields are in the R-R sector, while what we have trying to restore are the vectors in the NS-NS sector. Of course, in general one can think of U dualities which interchange NS-NS and R-R, but an explicit investigation shows that no \( E_7 \) element maps the dimensionally reduced supergravity to the minimally coupled one. Thus we conclude that there is no reason to expect that the reduction of \( D = 5 \) supergravity to \( D = 4 \) is realized by the orbifolds that we have considered. Also, even though there exists a totally different asymmetric orbifold which is not related to \( N=4 \) pure supergravity but has one \( N = 2 \) gravity and one vector multiplet as its massless fields, it would not be related to the five-dimensional theory in any limit.

## 5 Conclusions

We have tried to construct, by further twisting the \( N = 4 \) pure supergravity models, asymmetric orbifolds whose massless fields are the same as the dimensional reduction of \( D = 5 \) simple supergravity. We have found no example of models, and argued that this is in some sense natural because the second scalar is not the axion and the two vectors should come from the R-R sector.

We have seen that it is hard to fix all moduli but one \( N = 2 \) vector multiplet; it cannot be achieved by Calabi-Yau compactifications, nor by orbifolds. The string-theory origin of \( D = 5 \) simple supergravity still remains obscure. However, the arguments in the
last section indicate another possibility of finding \( D = 5 \) simple supergravity in string theories: We have seen that the \( B_{\mu\nu} \) components with four-dimensional spacetime indices are truncated, and we know models in which this truncation occurs: The orientifolds. It would be interesting to investigate whether any orientifold model realizes the \( D = 4 \) spectrum, and in case there is such a model, whether it has a decompactifying limit to \( D = 5 \).

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