\[ \mathcal{N} = 6 \] gauged supergravities
from generalized dimensional reduction

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Abstract

We construct new \[ \mathcal{N} = 6 \] gauged supergravities in four and five dimensions using generalized dimensional reduction. Supersymmetry is spontaneously broken to \[ \mathcal{N} = 4, 2, 0 \] with vanishing cosmological constant. We discuss the gaugings of the broken phases, the scalar geometries and the spectrum. Generalized orbifold reduction is also considered and an \[ \mathcal{N} = 3 \] no-scale model is obtained with three independent mass parameters.

PACS numbers: 04.50.+h, 04.65.+e, 11.25.Mj, 11.30.Qc.

Keywords: Extended Supersymmetry, Supersymmetry Breaking, Field Theories in Higher Dimensions, Supergravity Models.
1 Introduction

It has recently been shown that generalized dimensional reduction [1] from $D$ to $D-1$ dimensions [2, 3, 4, 5], as well as orientifold reduction with fluxes [6, 7, 8], can be used to derive new gauged supergravities with spontaneous supersymmetry breaking, which escaped previous classifications [9]. The reason why this type of gaugings had not generally been considered before is that dimensional reduction does not give the standard effective theory, but a dualized one, where a subgroup of the $R$-symmetry is embedded in a hidden sector [10]. This happens because, in the higher-dimensional theory, part of the supersymmetry automorphism group is contained in the Lorentz group. To obtain these supergravities directly in lower dimensions, it is necessary to move first in the dualized theory, then perform the gauging. This is in fact the procedure followed in [11, 12] to derive these theories for $\mathcal{N} = 3, 4$. Generalized reduction represents a very powerful tool to do this systematically and it was used in [2, 3] to get new $\mathcal{N} = 2, 4, 8$ gauged supergravities in diverse dimensions. In this letter we apply it to derive $\mathcal{N} = 6$ no-scale models in $D = 4$ and $D = 5$.

The Scherk–Schwarz twist is constructed by using the global symmetries of the higher-dimensional theory: as a result, a non-semisimple subgroup of these is gauged. Supersymmetry breaking can be partial or total, depending on the choice of the twist parameters. It is also possible to halve the number of surviving supersymmetries without reducing the number of independent breaking parameters by adding an orbifold projection (see for example [3, 13]). In this way we obtain flat-gauged $\mathcal{N} = 6$, $D = 4, 5$, and $\mathcal{N} = 3$, $D = 4$ no-scale models with three independent mass scales.

The plan of the paper is as follows. In the next section we review the structure of $\mathcal{N} = 6$ supergravities in diverse dimensions, the field contents, the symmetries and earlier results. In section 3, we describe generalized reductions from $D$ to $D-1$ dimensions, discussing the general properties of the method. The details of the $D = 5 \rightarrow 4$ and $D = 6 \rightarrow 5$ cases are presented in sections 4 and 5 respectively, while orbifold is discussed in section 4.1. The last section contains our conclusions.

2 $\mathcal{N} = 6$ in diverse dimensions

$\mathcal{N} = 6$ supergravities did not receive as much attention as their $\mathcal{N} = 8$ and $\mathcal{N} = 4$ sisters, since no string theory exists with such an amount of supercharges in its unbroken phase. However, it has been shown [14] that existing superstring theories can be reduced to $\mathcal{N} = 6$ supergravities in four dimensions via asymmetric orbifolds. Beyond four dimensions, $\mathcal{N} = 6$ supergravities exist only in $D = 5$ and $D = 6$. The $D = 5$ theory was first derived in [15], and its $AdS_5$ gauged version was constructed in [16]. In $D = 6$ two theories exist, the $(6,0)$ and the $(4,2)$. The former has neither graviton nor gravitino and will not be considered. The latter, constructed in [17, 18], is anomalous and has no lagrangian, as the $(2,0)$ and the $(4,0)$. The anomaly can be
cancelled only by adding two gravitino supermultiplets, thus recovering the $N = 8$ supergravity [18]. However, this does not represent a problem since we are interested only in the effective $D = 5$ theory obtained by dimensional reduction.

We can summarize the field content of the three theories [$D = 4$, $D = 5$ and the (4,2) of $D = 6$] as

$$D = 4 : \left[1, 6, 15 + 1, 20 + 6, 15 + \overline{15}\right]_0$$

$$[e^\alpha_\mu, \psi^I_\mu, V^{IJ}_\mu + V^I_\mu, \chi^{IJK} + \chi^I, \phi^{IJ} + \tilde{\phi}_{IJ}]_0$$

$$D = 5 : \left[1, 6, 14 + 1, 14 + 6, 14\right]_0$$

$$[e^A_M, \psi^i_M, V^{ij}_M + V_M^{ij} + V^{i}_{MM} + V^{i}_{MN} + V^{i}_{MN} + V^{i}_{MM}, \chi^{ab} + \chi^a, \phi^{ab}]_0$$

$$D = 6 : \left[1, (4, 1) + (1, 2), (5, 1)_2 + (1, 1)_2 + (4, 2), (5, 2) + (4, 1), (5, 1)\right]_0$$

$$[e^A_M, \psi^a_M + \psi^a_M, V^{-ab}_{MN}, V^+_{MN}, V^+_{MN} + V^{a,a}_M + V^{a,a}_M + \chi^a + \phi^{ab}]_0 ;$$

we used the following short-hand notation for supermultiplets

$$[(J = 2), (J = 3/2), (J = 1), (J = 1/2), (J = 0)]_{0/m},$$

where the subscript $(0/m)$ indicates a massless or massive multiplet, $(J = n)$ indicates, for each value of the spin $J = n$, the field representations of the corresponding automorphism groups

1. $U(6), USp(6)$ and $USp(4) \times USp(2)$ and finally the subscript $(2)$ stands for a two-form. Space-time indices $\mu, \nu, ..., M, N, ..., \hat{M}, \hat{N}, ...$ refer to $D = 4, 5, 6$, while the indices $I, J, ..., i, j, ..., a, b, ..., a, b, ...$ refer to $SU(6), USp(6), USp(4)$ and $USp(2)$ respectively. Repeated indices belonging to the same group are antisymmetrized and, in the case of $USp(2N)$, they are also $\Omega$-traceless ($\Omega$ being the $USp(2N)$ symplectic metric). The 2-forms $V^{(-/+)}_{MN}$ are (anti) self-dual antisymmetric tensors in six dimensions. These theories, besides supersymmetry and the invariance under general coordinate transformations, have a non-compact global symmetry, called $U$-duality ($G$), whose maximal compact subgroup is represented by the $R$-symmetry group ($H$). The manifolds $G/H$ define the following non-linear $\sigma$-models

$$\mathcal{M}_{D=4} = \frac{SO^*(12)}{U(6)}, \quad \mathcal{M}_{D=5} = \frac{SU^*(6)}{USp(6)},$$

$$\mathcal{M}_{D=6} = \frac{SU^*(4) \times SU^*(2)}{USp(4) \times USp(2)} \simeq \frac{SO(5, 1)}{SO(5)},$$

for the scalar fields in each theory $^2$. 

$^1$For simplicity, in the rest of the paper we will report only the number of states for each spin $J = n$ without making field representations explicit.

$^2$More details on groups $SU^*(2N), SO^*(2N)$ and their algebras can be found, for example, in [19].
3 Generalized reduction

Generalized reduction from $D$ to $D-1$ dimensions corresponds to imposing twisted periodicity conditions along the compact dimension, namely

$$
\Phi(x^\mu, y + 2\pi r) = U \Phi(x^\mu, y),
$$

where $U = \exp[T]$ is a symmetry of the theory. To switch back to periodic fields, it is sufficient to rescale them with the factor $U(y) = \exp[y/(2\pi r)T]$. If the twist is chosen to be in the $U$-duality compact subgroup, supersymmetry can be broken spontaneously without producing a cosmological constant (non-compact twists, in general, lead to theories with unstable scalar potentials and therefore will not be considered here). The gauged group generated in the effective theories is a non-semisimple subgroup of the $U$-duality group, in particular the semidirect product of $U(1)$ and $n$ shift symmetries, corresponding to the higher-dimensional gauge invariance along the compact direction. This can be easily deduced by looking at the reduction of field strengths and kinetic terms. It is easy to verify, for instance, that field strengths with flat indices

$$
V_{\alpha \beta} \rightarrow \rho^{-2\gamma} e^\mu_{\alpha} e^\nu_{\beta} U(y) \left( B_{\mu \nu} + MB[\mu A_{\nu}] + V_y A_{\mu \nu} \right) = \rho^{-2\gamma} e^\mu_{\alpha} e^\nu_{\beta} U(y) \left( \hat{B}_{\mu \nu} + V_y A_{\mu \nu} \right),
$$

$$
V_{\alpha \hat{y}} \rightarrow \rho^{-\gamma-1} e^\mu_{\alpha} U(y) \left[ \partial_\mu V_{\hat{y}} - M (B_{\mu} + A_{\mu} V_{\hat{y}}) \right] = \rho^{-\gamma-1} e^\mu_{\alpha} U(y) D_\mu V_{\hat{y}},
$$

where we have not reported internal indices explicitly,

$$
M \equiv U(y)^{-1} \partial_y U(y),
$$

$$
B_{\mu} \equiv V_{\mu} - V_y A_{\mu},
$$

$$
\hat{B}_{\mu \nu} \equiv B_{\mu \nu} + MB[\mu A_{\nu}], \quad B_{\mu \nu} \equiv \partial_\mu B_{\nu}, \quad A_{\mu \nu} \equiv \partial_\mu A_{\nu},
$$

$\rho = e^\hat{y}_{\mu}$ is the radion and $A_\mu = \rho^{-1} e^\hat{y}_{\mu}$ is the graviphoton. The $V_{\alpha \beta}$ term in eq. (4) produces field strengths that are covariant with respect to the group $U(1) \otimes T^n$, with structure constants given by the Scherk–Schwarz mass matrix $M$. The compact factor is gauged by the graviphoton coming from the higher-dimensional metric, while the translation symmetries were gauged by the corresponding vectors. The term $V_{\alpha \hat{y}}$, on the other hand, produces the covariant kinetic term for the axions $V_{\hat{y}}$, which can be shifted away in the unitary gauge, giving mass to the vectors through the Higgs mechanism. As discussed in [20, 2, 3] Chern–Simons terms, besides getting covariant, produce extra contributions which are necessary for the invariance of the lower-dimensional theory.

\[\text{3In deriving eq. (4) we adopted the vielbein parametrization of [1], i.e. the triangular gauge}\]

$$
\epsilon^A_M = \left( e^\hat{y}_{\mu}, e^\mu_{\alpha}, e^\mu_{\hat{y}}, 0, e^\hat{y}_{\mu} \right).
$$
It is straightforward to verify that the generalized reduction works in the same way for the fermionic sector. In fact, the \((\alpha\beta)\) term of the flat derivative \(\epsilon^A_M e^B_N \partial_M \psi_N\) makes the lower-dimensional gravitino kinetic term, covariant with respect to the gauged group, while the \((\alpha\tilde{y})\) one gauges the \(\psi_y\) component. Analogously to \(V_y\) with the vector fields \(V_\mu\), \(\psi_y\) is reabsorbed by gravitinos, which get the mass \(M\) through the super-Higgs effect. Finally the scalar manifold vielbein \(P_\mu\) becomes covariant as well, while the extra component produces a potential \(\text{tr}(P_y^2)\) that vanishes at its extremum.

The gauged supergravity obtained in this way is then a spontaneously broken one, with vanishing vacuum energy. We now discuss more in detail the \(N = 6\) case.

4 The \(D = 4\) effective theory

Although the four-dimensional \(N = 6\) \(U\)-duality group is \(SO^*(12)\), when we consider the effective theory obtained by dimensional reduction from the \(D = 5\) one, only part of this group will be a symmetry of the action. It is then useful to decompose the corresponding algebra with respect to the five-dimensional one \([su^*(6)]\):

\[
so^*(12) = usp(6) \oplus [su^*(6) \text{ mod } usp(6)] \oplus so(1,1) + (14 + 1)_r + (14 + 1)',
\]

on the right-hand side of the above equation, we have respectively: the \(R\)-symmetry inherited from the five-dimensional one, the algebra of the \(D = 5\) scalar manifold, the dilatation \(so(1,1)\) associated with the radion \(e^5\), and the \((14 + 1)_r\) shift symmetries coming from the corresponding \(D = 5\) vector fields. The remaining \((14 + 1)'\) generators are the compact elements that would extend \(USp(6)\) to \(U(6)\), but now fall into the magnetic subgroup of the \(Sp(32,R)\) duality group acting on the vector field strengths and their duals [21]. These generators can still be rotated into the electric subgroup, recovering the \(U(6)\) symmetry of the action, by means of a duality transformation. However, when the gauging will be turned on via the Scherk–Schwarz mechanism, the embedding showed above will be frozen, and the \(U(6)\) and \(USp(6)\) formulations will no longer be equivalent.

The scalar geometry can be decomposed as well:

\[
\frac{SO^*(12)}{U(6)} = \left[ \frac{SU^*(6)}{USp(6)} \times SO(1,1) \right] \otimes (T^{14+1}),
\]

showing how \(D = 5\) scalar fields fit into the four dimensional coset manifold.

The rank-3 group \(USp(6)\) can thus be used for the twist, producing a mass matrix \(M = U^{-1}(y) \partial_y U(y)\) with three independent mass parameters \(m_{1,2,3}\).

The spectrum is twice degenerate, as this mechanism cannot produce chirality. The mass for each field can be read directly from \(M\) and depends only on the way the field transforms under \(USp(6)\). The mass spectrum, already deduced in [22] \(^4\), is summarized in Table 1 and satisfies the mass formula \(\text{str}M^4 = \text{str}M^2 = 0\) [23]. Depending on the number of mass parameters

\(^4\)We take this opportunity to correct a typo in table 7 of [22].
turned on, supersymmetry can be broken totally or partially to $\mathcal{N} = 4, 2, 0$. The gauged group can be $U(1)^{\mathcal{O}} T^{12,10,8,6}$, depending on the number of massive vectors.

The broken phases then appear to be:

$$\mathcal{N} = 6 \rightarrow 4$$
$$[1, 4, 6, 4, 2]_0 + \{2 \times [0, 1, 4, 6, 4]_m\} + 2 \times [0, 0, 1, 4, 6]_0$$

$$\mathcal{N} = 6 \rightarrow 2$$
$$[1, 2, 1, 0, 0]_0 + 3 \times [0, 0, 1, 2, 2]_0 + 2 \times [2 \times [0, 1, 2, 1, 0]_m] \nonumber$$
$$+ 2 \times [2 \times [0, 0, 1, 2, 1]_m] + 2 \times [2 \times [0, 0, 0, 1, 2]_m]$$

$$\mathcal{N} = 6 \rightarrow 0$$
$$[1_0, 6_m, 12_m + 4_0, 20_m, 12_m + 6_0]$$

where, as expected [4, 22], massive multiplets in curly brackets are $\frac{1}{2}$ BPS short multiplets.

These theories thus represent the $\mathcal{N} = 6$ no-scale models in four dimensions.

### 4.1 Orbifold projection

We can also obtain chiral theories by compactifying on the orbifold $S^1/\mathbb{Z}_2$. The orbifold reduction can be implemented by assigning $\mathbb{Z}_2$ parities to fields

$$\Phi(-y) = \mathbb{Z}_2 \Phi(y), \quad (8)$$

in such a way that the theory remains invariant under the transformation $y \rightarrow -y$, and eliminating odd-parity fields that have no zero modes. This projection produces chirality by halving the

| Field | Mass | No. of states |
|-------|------|--------------|
| $e^{\alpha}_\mu + e^{5}_\mu + \tilde{e}^{5}_\mu$ | 0 | $1 + 1 + 1$ $1 + 0 + 1$ |
| $\psi^{i}_\mu + \tilde{\psi}^{i}_5$ | $m_{1,2,3}$ | $6 + 6$ $3 + 3$ |
| $V^{ij}_\mu + V^{ij}_5$ | $m_{\ell} \pm m_{\ell'}$ $\ell > \ell'$ | $12 + 12$ $6 + 6$ |
| | 0 | $2 + 2$ $0 + 2$ |
| $V_\mu + V_5$ | 0 | $1 + 1$ $0 + 1$ |
| $\chi^{ijk}$ | $|m_1 \pm m_2 \pm m_3|$ | 8 $4$ |
| | $|m_{1,2,3}|$ | 6 $3$ |
| $\phi^{ij}$ | $|m_{1,2,3}|$ | 6 $3$ |
| | $m_{\ell} \pm m_{\ell'}$ $\ell > \ell'$ | 12 $6$ |
| | 0 | 2 $2$ |

Table 1: Spectrum of $D = 4$ supergravity obtained by generalized reduction of $\mathcal{N} = 6$, $D = 5$ supergravity. The last two columns give the number of states for the reduced $\mathcal{N} = 6$ and $\mathcal{N} = 3$ theories respectively.
number of gravitinos (and thus the number of supersymmetries). The naive orbifold reduction leads to an $\mathcal{N} = 3$ supergravity coupled to three vector multiplets, whose scalar geometry is

$$\frac{SU(3,3)}{U(3) \times SU(3)}.$$  

A Scherk–Schwarz twist can be turned on by using a Cartan subalgebra of $USp(6)$, which anticommutes with $\mathbb{Z}_2$, along the lines of [3] for the $\mathcal{N} = 4$ case. The resulting theory is a matter-coupled flat $\mathcal{N} = 3$ supergravity, spontaneously broken to $\mathcal{N} = 2,1,0$, with three independent mass parameters. The spectrum is reported in Table 1 and is equal to the $\mathcal{N} = 6$ one without the four massless vectors and with the degeneracy of the massive states halved.

Because the graviphoton $A_\mu$ is always odd under the $\mathbb{Z}_2$ parity, the gauged group is an abelian $T^m$ ($m$ being the number of massive vectors). The field content for each theory is then:

$$\mathcal{N} = 3 \rightarrow 2$$

$$[1,2,1,0,0]_0 + [0,1,4,6,4]_m + [0,0,1,2,2]_0 + [0,0,0,2,4]_0$$

$$\mathcal{N} = 3 \rightarrow 1$$

$$[1,1,0,0,0]_0 + 2 \times [0,1,2,1,0]_m + 2 \times [0,0,1,2,1]_m$$

$$+ 2 \times [0,0,0,1,2]_m + 3 \times [0,0,0,1,2]_0$$

$$\mathcal{N} = 3 \rightarrow 0$$

$$[1_0,3m,6_m,10_m,6_m + 6_0]$$

where the massive fields are now arranged into long non-BPS multiplets [4, 22].

We can recognize in this theory the one constructed by Tsokur and Zinovev by dualization [11], the derivation given here shows in a simple way its five-dimensional origin, and explains why its spectrum still satisfies the $\mathcal{N} = 6$ supersymmetry constraint $\text{str} \mathcal{M}^4 = \text{str} \mathcal{M}^2 = 0$.

This theory can also be obtained from a $T^6/\mathbb{Z}_2$ orientifold reduction of Type IIB supergravity with fluxes turned on [24, 6]. Then we found another low-energy duality between flux and Scherk–Schwarz compactifications, analogous to the one found in [8] for the $\mathcal{N} = 4$ case.

5 The $D = 5$ effective theory

Finally we discuss the $D = 6 \rightarrow 5$ case. As before, it is useful to look at the decomposition of the $D = 5$ $U$-duality algebra $\mathfrak{su}^*(6)$ in terms of the $D = 6$ one $[\mathfrak{su}^*(4) \oplus \mathfrak{su}^*(2)]$:

$$\mathfrak{su}^*(6) = \mathfrak{usp}(4) \oplus \mathfrak{usp}(2) \oplus [\mathfrak{su}^*(4) \mod \mathfrak{usp}(4)] \oplus \mathfrak{so}(1,1) + \mathfrak{so}(4,2)_r + \mathfrak{so}(4,2)^t,$$  

where it is simple to identify respectively the $D = 6$ $R$-symmetry, the coset algebra of the higher-dimensional scalar manifold, the dilatation associated to the modulus $e_6^\delta$, the eight shift
Table 2: Spectrum of the $\mathcal{N} = 6$, $D = 5$ reduced supergravity

| Field | Mass | No. of states |
|-------|------|---------------|
| $e_M^A + e_M^6 + e_6^6$ | $0$ | $1 + 1 + 1$ |
| $\psi_M^a + \psi_6^a$ | $|m_{1,2}|$ | $4 + 4$ |
| $\psi_M^6 + \psi_6^6$ | $|m_3|$ | $2 + 2$ |
| $V_{MN}^{ab}$ | $|m_1 \pm m_2|$ | $2c_2$ |
| $V_{MN}^{a}$ | $0$ | $1$ |
| $V_{MN}^{a,a} + V_{6}^{a,a}$ | $|m_{1,2} \pm m_3|$ | $8 + 8$ |
| $\chi_{ab,a}^{a}$ | $|m_1 \pm m_2 \pm m_3|$ | $8$ |
| | $|m_3|$ | $2$ |
| $\chi^{a}$ | $|m_{1,2}|$ | $4$ |
| $\phi^{ab}$ | $|m_1 \pm m_2|$ | $4$ |
| | $0$ | $1$ |

generators in the $(4, 2)$ of $USp(4) \times USp(2)$ associated to the corresponding $D = 6$ vector fields, and the compact generators of $USp(6)/[USp(4) \times USp(2)]$. The scalars of the $D = 6$ coset manifold, together with those coming from the compactification procedure, are thus embedded into the $D = 5$ one as follows:

$$SU^*(6) = \left[ SU^*(4) \times SO(1, 1) \right] \otimes \mathcal{T}^{(4,2)}.$$ (11)

Three independent parameters are still available for the twist $|m_{1,2}$ from $USp(4)$ and $m_3$ from $USp(2)$]. The spectrum can easily be derived by looking at how fields transform under the $R$-symmetry group, and is reported in Table 2. Note that four anti self-dual forms complexify into two massive complex antisymmetric tensors fields in $D = 5$ ($2c_2$), while the remaining two can be dualized to vectors. As in $D = 4$, we can have partial or total supersymmetry breaking to $\mathcal{N} = 4, 2, 0$. Note that $\mathcal{N} = 4$ can be obtained in two different ways: by turning on $m_3$ we get a $USp(4)$ invariant action, while turning on either $m_1$ or $m_2$ we have only $USp(2) \times USp(2)$. The gauged groups are $U(1)\otimes \mathcal{T}^{8,6,4}$, depending on the number of vectors acquiring a mass. The broken phases for $\mathcal{N} = 6 \rightarrow 4, 2, 0$ in $D = 5$ then are

$$\mathcal{N} = 6 \rightarrow 4 \quad USp(4)$$
$$[1, 4, 6, 4, 1]_0 + \{2 \times [0, 1, 4, 5, 0]_m \} + [0, 0, 1, 4, 5]_0$$

$$\mathcal{N} = 6 \rightarrow 4 \quad USp(2) \times USp(2)$$
$$[1, 4, 6, 4, 1]_0 + \{2 \times [0, 1, 2 + 1_2^5, 5, 2]_m \} + [0, 0, 1, 4, 5]_0$$

$$\mathcal{N} = 6 \rightarrow 2$$
$$[1, 2, 1, 0, 0]_0 + 2 \times [0, 0, 1, 2, 1]_0$$
In this case anomalies will not allow us to upgrade the result to six dimensions. We simply used the $D = 6$ theory together with generalized dimensional reduction as a tool to derive the $D = 5, \mathcal{N} = 6$ no-scale model, analogously to what had been done in [3] for the $(4,0) D = 6$ theory.

6 Conclusions

We constructed $\mathcal{N} = 6$ no-scale models in $D = 4, 5$ through generalized dimensional reduction, and discussed the gauging, the spectrum and the scalar geometries. The theories obtained in this way belong to a class of gauged supergravities not considered by earlier classifications. Moreover the theories presented in this paper, together with the one constructed in [16], are the only known $\mathcal{N} = 6$ gauged supergravities. We also rederived the $\mathcal{N} = 3, D = 4$ theory of [11], giving it a five-dimensional interpretation.

Acknowledgements

We are grateful to S. Ferrara and F. Zwirner for helpful discussions and comments on the manuscript.

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