Generalized dimensional reduction of $D = 6$ (2,0) chiral supergravity

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We report on the Scherk–Schwarz reduction of $D = 6$ (2,0) supergravity coupled to matter and its interpretation as $D = 5$ gauged $N = 2$ supergravity of no-scale type.

PACS: 04.65.+e, 04.50.+h, 11.10.Kk
Key words: Scherk–Schwarz , 6D supergravity, dimensional reduction

1 Introduction

In the present contribution we report on the Scherk–Schwarz (SS) dimensional reduction [1] on S^1 of $D = 6$ ungauged supergravity theories with eight supercharges [2]. The reduction gives supergravities in five dimensions with a flat gauge group. Such flat gaugings appear in four dimensions in the context of both, SS [3, 4, 5] and flux compactifications [6].

The Scherk–Schwarz mechanism relies on the presence of a global symmetry group of the higher dimensional theory. The class of no-scale supergravities at $D = 5$ that we obtain depend then on the global symmetry of the $D = 6$ theory [7].

In the (2,0) (minimal) theories [7, 8, 9] there are three kinds of matter multiplets: the vector multiplet which has no scalars, the tensor multiplet with 1 scalar and the hypermultiplet with four scalars. If we have n_T tensor multiplets and n_H hypermultiplets, the scalar manifold is a product

$$\frac{SO(1, n_T)}{SO(n_T)} \times \mathcal{M}_Q \,, \tag{1}$$

where $\mathcal{M}_Q$ is a quaternionic manifold of quaternionic dimension n_H [10]. The SS phase is, in general, a combination of isometries of both manifolds.

The graviton multiplet contains a self dual tensor field, while the tensors from the tensor multiplets are anti-self dual. We denote the set of tensor fields as B^r, $r = 0, \ldots, n_T$, with B^0 pertaining to the graviton multiplet.

When vector multiplets are present, the vectors (A^x, $x = 1, \ldots, n_V$) couple to the tensor fields and their interaction term is of the form [11, 12, 13, 14]

$$C_{rxy} B^r \wedge F^x \wedge F^y \,, \quad F^x = \mathrm{d}A^x \,,$$

with C_{rxy} = constant. This term is related by supersymmetry to the kinetic term of the vectors

$$C_{rxy} b^r F^x \wedge {}^* F^y \,. \qquad (2)$$

The fields b^r, $r = 0, \ldots, n_T$ satisfy the constraint

$$\eta_{rs} b^r b^s = 1 \,,$$

which defines the manifold $SO(1, n_T)/SO(n_T)$. The terms (2) explicitly break the $SO(1, n_T)$ symmetry, unless the vector fields A^x transform under some n_V–dimensional representation R_V of $SO(1, n_T)$ with the property that $\mathrm{Sym}(R_V \otimes R_V)$ contains the vector representation. In that case, the constants C_{rxy} can be chosen as invariant couplings. This happens, for instance, if R_V is a spinor representation of $SO(1, n_T)$. Remarkably, this choice leads after dimensional reduction on S^1 to the real special geometries which are homogeneous (in particular, symmetric) spaces [15, 16, 17, 18].

Under this assumption, the SS reduction produces a theory with a flat gauge group of the form $U(1) \ltimes R_V$, where the $U(1)$ generator is in the Cartan subalgebra (CSA) of the maximal compact subgroup $SO(n_T)$ of the global symmetry $SO(1, n_T)$. The $U(1)$ group is gauged by the vector coming from the metric in dimension six. The tensors are in a vector representation of $SO(1, n_T)$, so they are charged under U(1) (except for some singlets as B^0).

We remark that in order to introduce a SS phase in the tensor–vector multiplet sector it is actually sufficient that the constants C_{rxy} preserve a $U(1)$ subgroup of $SO(1, n_T)$, which is a much weaker assumption. The examples that we will consider in this paper have the full $SO(1, n_T)$ symmetry.

The generator of the group $U(1)$ may also have a component on the isometries of the quaternionic manifold [19]; in particular, it may have a component in the CSA of the $SU(2)$ R–symmetry, then breaking supersymmetry (notice that this can happen even if hypermultiplets are not present, corresponding to a $D = 5$ Fayet–Iliopoulos term). The SS reduction leads to a positive semidefinite potential also in this case. The $D = 5$ interpretation of the theory must correspond to a gauging with the term $V_R = 0$ (see Section 4 and Ref. [20]).

The report is organized as follows. In Sections 2, 3, 4 we discuss the dimensional reduction in presence os SS phases, included the induced scalar potential and its comparison with $D = 5$ gauged supergravity. In Section 5 we discuss the conditions for uplifting (oxidation) of the $D = 5$ theory to a chiral theory in $D = 6$.

2 Generalized dimensional reduction

We give here the qualitative features of the SS reduction of a general (2,0) theory from $D = 6$ to $D = 5$ and show how it produces an $N = 2$ theory in $D = 5$ with tensor, vector and hyper multiplets, and a flat gauge group.

Let us consider a $D = 6$ theory with n_T tensor multiplets, n_V vector multiplets and n_H hypermultiplets. These theories are anomalous unless the condition

$$n_H - n_V + 29\, n_T = 273 \qquad (3)$$

is satisfied [21].

It was shown in Ref. [11] that when performing a standard dimensional reduction to $D = 5$ on an anomaly–free (2,0) theory, we obtain a particular class of $N = 2$, $D = 5$ theories. After the reduction, the geometry of the hypermultiplets ($\mathcal{M}_Q$) remains unchanged. The scalar manifold of the vector and tensor multiplets has a real special geometry [25]. Let $\mathcal{M}_R$ be this manifold in $D = 5$ and $d = \dim \mathcal{M}_R$.

On general grounds, real–special geometry consists essentially on an embedding of $\mathcal{M}$ in a manifold of dimension $d + 1$ through a cubic polynomial constraint

$$\mathcal{V} = d_{IJK} t^I t^J t^K = 1\,, \quad I, J, K = 1, \ldots, d + 1\,.$$

The metric induced by the embedding from the metric in the higher dimensional manifold a_{IJ},

$$a_{IJ} = -\tfrac{1}{2}\,\partial_I \partial_J \ln \mathcal{V}\,, \quad g_{ij} = a_{IJ} \partial_i t^I \partial_j t^J \big|_{\mathcal{V}=1}\,, \quad i, j = 1, \ldots, d\,. \tag{4}$$

In the following, we will denote $G_{IJ} = a_{IJ}|_{\mathcal{V}=1}$. When the $D = 5$ theory comes from a dimensional reduction from $D = 6$, $d = n_T + n_V + 1$ (the extra scalar coming from the metric), and the cubic polynomial takes the particular form

$$\mathcal{V} = 3\,(z\eta_{rs} b^r b^s + C_{rxy}\, b^r a^x a^y)\,; \quad r = 0, 1, \cdots, n_T\,; \quad x = 1, \cdots, n_V\,. \tag{5}$$

η_{rs} is the $(1, n_T)$ Lorentzian metric related to the space $SO(1, n_T)/SO(n_T)$ (parametrized by b^r) in (1), $z = \sqrt{g_{55}} = \mathrm{e}^\sigma$ is the Kaluza–Klein scalar and $a^x = A^x_6$ are the axions.

We now focus on the cases when $SO(1, n_T)$ is a global symmetry. This demands the coupling C_{rxy} to be an invariant coupling in the sense explained in Section 1. One could then introduce a SS phase in the CSA of $SO(n_T)$. Some of the vector and tensor multiplets are charged under this generator, so they acquire mass. In the $D = 5$ interpretation the vectors gauge a non-abelian flat group, but their scalar partners give no contribution to the scalar potential, in agreement with the known results on $D = 5$ gauged supergravity [22, 23]. The gauging of flat groups in the context of $N = 2$ supergravity has not been considered in previous classifications [24]. These gaugings are always of no-scale type due to the particular structure of the critical points [26].

Finally, we want to note that to uplift (oxidate [27, 28]) to $D = 6$ a five dimensional $N = 2$ supergravity a necessary condition is that the cubic polynomial defining the real special geometry has the form (5). All the homogeneous spaces with real special geometry fall in this category. These spaces have been classified in Refs. [15, 16, 17, 18]; they were denoted as $L(q, P, \dot{P})$ in Ref. [18]. We explain here this notation. Let $q = n_T - 1$ and let $\mathcal{D}_{n_T}$ be the real dimension of an irreducible representation of $\mathrm{Spin}(1, n_T)$. For $n_T = 1, 5 \mod 8$ there are two inequivalent real or pseudoreal (quaternionic) representations. Let P and $\dot{P}$ denote the number of copies of such representations ($\dot{P} = 0$ for $n_T \neq 1, 5 \mod 8$). Then, $n_V = (P + \dot{P})\mathcal{D}_{n_T}$. The R–symmetry group of $\mathrm{Spin}(1, n_T)$ in the representation $(P, \dot{P})$ is denoted by $\mathcal{S}_q(P, \dot{P})$ (see Table 3 of Ref. [18]).

When $\dot{P}=0$, the notation $L(q,P)=L(q,P,0)$ is used. The symmetric spaces [25] correspond to the particular cases $L(1,1)$, $L(2,1)$, $L(4,1)$, $L(8,1)$, $L(-1,P)$ and $L(0,P)$. We also have $L(q,0)=L(0,q)$. They are reported in Table 2. of Ref. [18]. The examples of SS reductions reported in this paper will actually fall in this class.

3 $D=5$ massive tensor multiplets

The $D=5$ theory obtained through an ordinary Kaluza–Klein dimensional reduction contains n_T+n_V+1 vector multiplets. This is because the (anti) self-duality condition in $D=6$

$$\partial_{[\mu}B^r_{\nu\rho]}=\pm\frac{1}{3!}\,\epsilon_{\mu\nu\rho\lambda\tau\sigma}\partial^\lambda B^{r|\tau\sigma}\,,\quad \mu,\nu=1,\ldots,6 \tag{6}$$

tells us that in $D=5$ the two form $B^r_{\mu\nu}$ is dual to the vector $B^r_{\mu 6}$ $(\mu,\nu=1,\ldots 5)$.

We want now to perform a SS generalized dimensional reduction instead. Let ${M^r}_s=-{M_s}^r$ be the SS phase in the CSA of the global symmetry $SO(n_T)\subset SO(1,n_T)$. The form B^0 (of the gravitymultiplet) is inert under $SO(n_T)$, so in the rest of this subsection the value $r=0$ is excluded and $r=1,\ldots n_T$. The $D=6$ anti self-duality condition gives now

$$\begin{aligned}\partial_{[\mu}B^r_{\nu\rho]}&=\frac{1}{3!}\,\epsilon_{\mu\nu\rho\lambda\tau 6}\left(\partial^6 B^{r|\lambda\tau}+2\partial^\lambda B^{r|\tau 6}\right)=\\&=\frac{1}{3!}\,\epsilon_{\mu\nu\rho\lambda\tau}\left({M^r}_s B^{s|\lambda\tau}+F^{r|\lambda\tau}\right),\quad \mu,\nu=1,\ldots,5\,,\end{aligned} \tag{7}$$

where $F^r_{\lambda\tau}=2\partial_{[\lambda}B^r_{\tau]6}$.

Equation (7) can be rewritten as a self-duality condition for a massive two–form in five dimensions [29]. Assume that the Cartan element M is invertible; then we can define

$$\hat{B}^r_{\mu\nu}=B^r_{\mu\nu}+{(M^{-1})^r}_s F^s_{\mu\nu}\,,$$

so

$$\partial_{[\mu}\hat{B}^r_{\nu\rho]}={M^r}_s\frac{1}{3!}\,\epsilon_{\mu\nu\rho\lambda\tau}\hat{B}^{s|\lambda\tau}\,,$$

that is,

$$\mathrm{d}\hat{B}^r={M^r}_s{}^*\hat{B}^s\,.$$

For n_T even, an element M with non zero eigenvalues $\pm\mathrm{i}m_\ell\neq 0$ $(\ell=1,\ldots,n_T/2)$ is invertible. Then we have $n_T/2$ complex massive two–forms. For n_T odd, the matrix M has at least one zero-eigenvalue. The corresponding antisymmetric tensor B^{r_0} is a gauge potential which can be dualized to a vector. If some other eigenvalue m_ℓ is zero, the same argument applies and there will be a couple of tensors (or one complex tensor) which can be dualized to vectors.

Summarizing, in the five dimensional theory there are $2n\leq n_T$, massive tensor multiplets (or n complex ones) and n_T-2n+1 abelian vector multiplets, one of them formed with the vector which is dual (after reduction to $D=5$) to the self dual tensor present in the $D=6$ graviton multiplet. This vector is a singlet of the global symmetry group.

4 The induced scalar potential and its extrema

In this section we compute the scalar potential of the SS reduced theory.

The scalar potential comes from the kinetic term of the scalar fields [1]. The only scalars at $D=6$ are in the tensor and hyper multiplets, which parametrize the manifold in (1). We denote by φ^i, $i=1,\ldots,n_T$ the coordinates on $SO(1,n_T)/SO(n_T)$, and let

$$v^a = v^a_i \partial_\mu \varphi^i \, \mathrm{d}x^\mu = v^a_\mu \, \mathrm{d}x^\mu \,, \quad a=1,\ldots,n_T$$

be the pull back to space time of the vielbein one form. Similarly, the quaternionic manifold $\mathcal{M}_Q$ [10, 30], with holonomy $SU(2)\times USp(2n_H)$, has coordinates q^u, $u=1,\ldots 4n_H$ and vielbein

$$\mathcal{U}^{\alpha A} = \mathcal{U}^{\alpha A}_u \partial_\mu q^u \, \mathrm{d}x^\mu = \mathcal{U}^{\alpha A}_\mu \, \mathrm{d}x^\mu \,, \quad \alpha = 1,\cdots,2n_H\,, \quad A=1,2\,.$$

There is still a scalar mode coming from the metric $\mathrm{e}^\sigma = \sqrt{g_{55}}$.

For the scalar potential we obtain

$$V(\sigma,\varphi,q) = V^{SS}_T + V^{SS}_H = \mathrm{e}^{-8\sigma/3}\left[v^a_6(\varphi)v_{6a}(\varphi) + \mathcal{U}^{\alpha A}_6(q)\,\mathcal{U}^{\beta B}_6(q)\mathbb{C}_{\alpha\beta}\epsilon_{AB}\right], \quad (8)$$

where $\mathbb{C}_{\alpha\beta}$ and ϵ_{AB} are the antisymmetric metrics.

We see that this potential is semipositive definite. The critical points occur at

$$v^a_6(\varphi)=0 \quad \text{and} \quad \mathcal{U}^{\alpha A}_6(q) = 0\,, \quad (9)$$

so $V=0$ at the critical points, which are then Minkowski vacua. The scalar σ is not fixed, so the theory is of no-scale type. Notice that (9) implies

$$v^a_6(\varphi) = v^a_i M_i^j \varphi^j = 0,\,, \quad \mathcal{U}^{\alpha A}_6(q) = \mathcal{U}^{\alpha A}_u M^u_v q^v = 0\,.$$

If the mass matrices have some vanishing eigenvalues, then this results in some moduli of the theory, other than σ. For n_T odd, since the tensor multiplet mass matrix has always one vanishing eigenvalue, there are at least two massless scalars. There are three massless vectors in this case.

The SS potential given in (8) should be compared to the most general gauging of $N=2$, $D=5$ supergravity [22, 23, 32]

$$V_{D=5} = V_T + V_H + V_R\,; \quad V_T \geq 0\,; \quad V_H \geq 0\,,$$

where V_T and V_H are the contributions of tensor and hypermultiplets (separately positive) and V_R is the contribution from vector and gravity multiplets due to the quaternionic Killing prepotential P^X_I, $X=1,2,3$ [30].

For a $D=5$ gauging corresponding to a SS reduction, we then need $V_R=0$.

The explicit form of V_R is [23, 22, 32]

$$V_R = -4t_{IJ}G^{IK}G^{JL}P^X_K P^X_L = -\frac{4}{3}\left(\frac{1}{3}\,(t^{-1})^{IJ} + t^I t^J\right) P^X_I P^X_J\,, \qquad (10)$$

where $t_{IJ} = d_{IJK}t^K$ and $G^{IJ} = a_{IJ}|_{\mathcal{V}=1} = -\frac{1}{3}\,(t^{-1})^{IJ} + t^I t^J$ (see (4)).

Even when there are no hypermultiplets, this term is not necessarily zero, because one can take a constant prepotential, $P^{X_0}_I = g_I =$ constant (the rest zero.). g_I is the $N=2$ Fayet–Iliopoulos parameter, and we retrieve the particular form of V_R found in Ref. [20].

Equation (10) can also be written as

$$V_R = -4d^{IJK}t_I P^X_J P^X_K\,, \qquad (11)$$

where indices are lowered and raised with the metric G_{IJ}. For symmetric spaces one has $d_{IJK} = d^{IJK}$ [20].

From the point of view of the SS reduction, the constant prepotential corresponds to an $SU(2)$ phase, which in absence of hypermultiplets only gives masses to the fermions. Therefore we must have $V_R = 0$ for any value of t^I in the reduced theory. Moreover, since this depends only on the real special geometry (see (11)), this conclusion also holds in presence of hypermultiplets.

In the SS reduction the vector gauging the $U(1) \subset SU(2)$ is the partner of the scalar $z = \mathrm{e}^\sigma$, so $P^X_z \neq 0$ and the rest are zero. $V_R = 0$ then requires

$$(t^{-1})^{zz} = -3(t^z)^2 \iff G^{zz} = 2(t^z)^2\,. \qquad (12)$$

Let us consider some particular examples of theories with $V_R = 0$. Setting $n_V = 0$, equation (5) becomes

$$\mathcal{V} = 3z\eta_{rs}b^r b^s$$

and one can check that (12) holds [20]. It also holds for the spaces $L(0,P)$. More generally, it holds for all symmetric spaces with real special geometry because of the relations $d_{IJK} = d^{IJK}$ and $d_{zzI} = 0$. They readily imply $V_R = 0$.

We have checked that there are in fact counterexamples to the condition $V_R = 0$ among the theories classified in [17, 18] which are all of the form (5), so $V_R = 0$ is a further restriction satisfied by the $D=5$ real geometries that can be uplifted (oxidated) to $D=6$. It would be interesting to know, in the general case, what the conditions on the coefficients $C^r_{xy} = \eta^{rs}C_{sxy}$ in (5) are to have $V_R = 0$.

We will see in the next subsection that the possible resolution of this puzzle lies in the cancellation of anomalies of the six dimensional theory.

5 Uplifting $N = 2$ at $D = 5$ to (2,0) $D = 6$ theory

In $D=6$, $(2,0)$ chiral theories it was found that there is, in general, a clash between the gauge invariance of the two–forms and the gauge invariance of the 1 forms (vector fields). For generic couplings C_{rxy}, in the abelian case, the n_V $U(1)$ currents J_x are not conserved but satisfy the equation [14, 11]

$$\mathrm{d}^* J_x = \eta_{rs} C^r_{xy} C^s_{zw} F^y \wedge F^z \wedge F^w \,. \tag{13}$$

This violation of the gauge invariance implies also a violation of supersymmetry because the theory is formulated in the Wess–Zumino gauge and the supersymmetry algebra closes only up to gauge transformations [11]. The current is conserved if the constants C^r_{xy} satisfy the condition

$$\eta_{rs} C^r_{x(y} C^s_{zw)} = 0 \,. \tag{14}$$

This condition is equivalent to the seemingly stronger condition

$$\eta_{rs} C^r_{(xy} C^s_{zw)} = 0 \tag{15}$$

because $C^r_{xy} = C^r_{yx}$. This can also be seen from the fact that the anomaly polynomial [11]

$$A \sim \eta_{rs} C^r_{xy} C^s_{zw} F^x \wedge F^y \wedge F^z \wedge F^w$$

vanishes if (15) holds. It is interesting to observe in this respect that among the homogeneous spaces in Ref. [17, 18] only the symmetric spaces, with the exception of the family $L(-1, P)$, $P > 0$, satisfy this condition [17, 18].

Also, we must note that the symmetric spaces satisfying (15) do have in fact $V_R = 0$, while for the homogeneous, non symmetric cases there are counterexamples.

Condition (14) is only required for a $D = 6$ ungauged supergravity. If the theory in $D = 6$ is already gauged, the terms in the right hand side of (13) may be compensated by (one loop) quantum anomalies through a Green–Schwarz mechanism, namely, the Lagrangian becomes a Wess–Zumino term [9, 12, 14]. The $D = 6$ potential is semipositive definite and simply given by [8]

$$V_{D=6} \simeq \sqrt{g} P^X_x P^X_y (C^{-1})^{xy} \,, \quad \text{where } C_{xy} = C_{rxy} b^r \,.$$

The $D = 6$ supersymmetric vacua occur at $P^X_x = 0$. An hypermultiplet can be "eaten" by a vector multiplet, making it massive. Note that there are not BPS particle multiplets in $D = 6$. The additional contribution to the potential in $D = 5$ is

$$\sqrt{g_5}\, \mathrm{e}^{-2\sigma/3} P^X_x P^X_y (C^{-1})^{xy} \,.$$

Since in this case V_R needs not to vanish, one may find new vacua in the SS reduction.

As an illustration of spaces satisfying (15), we give the spectrum of tensor, vector and hypermultiplets for the exceptional symmetric spaces in Table 1. Note that the values of n_T and n_V are given by the uplifting (oxidation) procedure of Ref. [31, 28]. These spaces are contained in the classification of Refs. [17, 18] and consequently have a cubic polynomial of the form (5). The number n_H instead is fixed by the gravitational anomaly cancellation (3). For generic SS phases, the

Table 1. Exceptional symmetric spaces

$L(q,P)$	$L(1,1)$	$L(2,1)$	$L(4,1)$	$L(8,1)$
(n_T,n_V,n_H)	(2,2,217)	(3,4,190)	(5,8,136)	(9,16,28)

$L(1,1)$ model has one massless scalar and two massless vectors. All the other exceptional models have two massless scalars and three massless vectors.

There are no other solutions in the series $L(q,P,\dot{P})$. It is obvious that for non homogeneous spaces the constants C^r_{xy} are rather arbitrary and there may be much more solutions to the uplifting condition.

However, in order to have a SS phase in the tensor and vector multiplet sector, non homogeneous spaces should have at least a residual $U(1)$ isometry.

The work of S.F. has been supported in part by the D.O.E. grant DE-FG03-91ER40662, Task C, and in part by the European Community's Human Potential Program under contract HPRN-CT-2000-00131 Quantum Space-Time, in association with INFN Frascati National Laboratories.

The work of M. A. Ll. has been supported by the research grant BFM 2002-03681 from the Ministerio de Ciencia y Tecnología (Spain) and from EU FEDER funds and by D.O.E. grant DE-FG03-91ER40662, Task C.

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