Gauge Charges from Supergravity

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Abstract. Some recent results in the study of four dimensional supergravity flux compactifications are reviewed, discussing in particular the role of torsion on the compactification manifold in generating gauge charges for the effective four dimensional theories.

1. Introduction

Many efforts in theoretical physics have been devoted, in the last decades, to look for a quantum theory describing in a consistent and unified way gauge interactions with gravity. Such a theory, if any, should incorporate the Standard Model of elementary particles but is also expected to give answers to some problems left open after the huge theoretical and phenomenological successes of the Standard Model, like the understanding of the hierarchy problem, of confinement in QCD, of the tiny value of the cosmological constant, or of the black-hole entropy.

In this spirit, superstring theory appears nowadays as a very promising candidate for the quantum theory underlying Nature: it is indeed a consistent and finite theory for quantum gravity, giving unification of gauge interactions with gravity. It is described by a two dimensional conformal theory defined at the Planck scale, spanning a curved ten dimensional target space. When the spectrum is extended to include solitonic configurations (D-branes [1]) it may reveal one extra eleventh dimension (M-theory) [2, 3].

The massless excitations of superstrings are described, in a low energy limit, by effective supersymmetric field theories, named supergravity theories, which contain supersymmetric couplings of gravity with matter and gauge fields. Supergravity theories may be defined in any space-time dimensions up to eleven, however they get predictive power when interpreted as effective theories for superstrings or M-theory. If superstring theory has to provide a natural explanation of interactions in our real world, some trick must be present so that in particular the theory, at low energies, looks four dimensional. The simplest way to do so is to imagine that the vacuum configuration for space-time is
not ten dimensional Minkowski space, but is instead of the form $\mathbb{R}^{(1,3)} \times M_6$, where $M_6$ is some six-dimensional compact manifold, whose size is so small that at the length-scales experimented in our low energy world it cannot be observed. Then, any tensor field of the ten dimensional theory looks, in the four dimensional compactified theory, as a collection of fields of various masses and spins [4] (the procedure of compactification from higher dimensions to four was first studied in five dimensions by Kaluza and Klein [5]). In particular, the ten dimensional fields give rise to a set of massless four dimensional scalars (moduli) describing metric and topological properties of the compactification manifold $M_6$. Any choice of the vacuum background of string theory selects one among a lot of possible inequivalent compactified space-times, and this implies giving a particular expectation value to the scalars. Note that these scalars are just the ones appearing in supergravity theories. Giving them some particular vacuum expectation values means, going down to very low energies, fixing the values of the physical observables in the Standard Model describing strong and electro-weak interactions.

We are going to focus on this intertwining role of supergravity between the stringy (Planck) scale and low energy physics. We will discuss in particular how gauge interactions arise in supergravity theories, with particular attention to the role of the torsion in the compactification manifold as gauge charge.

This contribution is organized as follows: In section 2 the role of scalar fields and the way gauge charges appear in supergravity are surveyed; in section 3 the example of no-scale models as gauged supergravity theories is discussed, with particular attention to the Scherk–Schwarz model, while section 4 focuses on the role of torsion on the compactification manifold in flux compactifications and on its relation to the Scherk–Schwarz model. Conclusions and outlook are left to the concluding section. The bibliography section is not exhaustive; somewhere we preferred to give reference to review papers where more extended references on the subject may be found. For a more complete list of references we refer to [6].

2. Scalars in supergravity and gauging

A peculiar feature of supergravity is that the spectrum includes fermionic and bosonic fields with spin $\leq 2$, among which scalars. Supersymmetry strongly constrains the couplings, and in particular it implies that the scalars belong to non linear $\sigma$-models:

$$L_{\text{kin|scalars}} = g_{\ell m}(\phi) \partial_\mu \phi^\ell \partial^\mu \phi^m.$$ (1)

The scalars play an important role in the field theory, as they generalize the Higgs field of the Standard Model. The low energy observables (gauge couplings, mass parameters, the cosmological constant, etc...) are given in terms of vacuum expectation values (v.e.v.) of scalars. The cosmological constant $\Lambda$, for example, is the v.e.v. of the scalar potential

$$\Lambda = < V(\phi) >,$$ (2)
while the gauge couplings and $\theta$-term appear in the lagrangian as the v.e.v. of scalar dependent functions:

$$L_{\text{kin|vec}} = \gamma_{\Lambda\Sigma}(\Phi) F^\Lambda \wedge F^\Sigma + \theta_{\Lambda\Sigma}(\Phi) F^\Lambda \wedge F^\Sigma.$$  \hfill (3)

However, the v.e.v. of the scalar fields fixing the value of the couplings in the effective supergravity theories are in fact determined by the choice of a vacuum in the corresponding underlying superstring theory. What happened, with the introduction of superstring theory as a unifying theory describing in particular particle physics, is that the Standard Model problem of having many free parameters to be fixed has turned into the problem of an enormous degeneracy of vacua for string theory. As we can see, in this way we do not gain any predictive power on the quantities characterizing our world; however, we obtain the important conceptual achievement of unifying gravity with the other interactions and to give a natural understanding for the origin of all the parameters involved (completely arbitrary in the Standard Model). Moreover, there is the hope that some non-perturbative stringy effect, inducing supersymmetry breaking, could also lift the vacuum degeneracy leaving fewer accessible vacua.

One is then driven to the problem of understanding how gauge charges and interactions arise in supergravity.

Supergravity theories are built from symmetry requirements: local invariance under the supersymmetry algebra. The bosonic sector includes gravity, abelian gauge fields and uncharged scalars; to get a theory invariant under the local supersymmetry algebra there is no need to introduce gauge charges. The gauge charges are introduced via a deformation of the theory named the gauging procedure, which corresponds to promote a global symmetry of the theory (an isometry of the scalar manifold) to gauge symmetry and to turn on interactions with the corresponding gauge fields of the theory (which may become non-abelian). Consistency of the deformation with supersymmetry implies the presence of a scalar potential, where the dynamical information on the model are codified.

However, if we want to study supergravity theories as effective theories of superstrings, it must be possible to interpret the gauge charges of a given supergravity model in terms of data corresponding to some superstring background. For example, ungauged supergravity models (with no gauge charges) appear as the low energy limit of perturbative Type II superstrings compactified on tori without any non-trivial flux for the tensor fields (p-forms) of the theory in the internal directions. On the other hand, starting from a ten dimensional theory, three different mechanisms have been identified to generate gauge charges in the four dimensional effective supergravity models $\mathbb{N}$ $\mathbb{S}$ $\mathbb{U}$ $\mathbb{K}$:

- by turning on charges corresponding to non abelian gauge fields already in the ten dimensional theory (from gauge fields living on D-branes, for Type I or II theories, or from the heterotic sector, for the case of the heterotic string);
- by turning on non trivial p-form fluxes in the internal compact manifold;
• by turning on charges associated to the geometry of the compactification manifold (à la Scherk and Schwarz).

The emerging scenario is that the first class of charges should contain the gauge group of the Standard Model, while the other two classes of charges should be responsible for supersymmetry breaking, via the super-Higgs mechanism.

Note that the dualities among string theories still map one into the other theories in the presence of charges, provided that also the charges are transformed accordingly.

3. An explicit example: the no-scale models

No-scale models are a particular class of supergravity theories emerged in the 80’s for their phenomenological interest [11, 12]. They have seen a revived interest from theoretical physicists after that some recent progress in string theory (D-branes [1], Randall–Sundrum scenarios [13]) allowed to find them a nice interpretation in terms of the classes of ten dimensional charges introduced in the previous section.

No-scale supergravities are models with an automatically vanishing cosmological constant (at tree level), and where supersymmetry is spontaneously broken. The gravitino mass $m_{3/2}$ is a sliding scale at tree level, dynamically fixed by radiative corrections, with possibility of a hierarchical suppression with respect to the Planck scale. The characteristic features of the various no-scale models are codified in the choice of the gauge group and in the structure of the scalar potential. However, they all share the feature of having a semi-positive definite scalar potential with flat directions. From a purely supergravity point of view, no-scale models are typically related to gaugings of suitable non-semisimple global symmetry groups of the Lagrangian (flat gaugings) [14, 15]. A peculiar feature of this kind of models is that the spectrum appears separated in an observable sector, (where the standard Model is defined), and an hidden sector (which is decoupled for energy scales $E << M_P$) interacting with the observable sector only through gravity.

Phenomenologically viable no-scale models have been constructed in string compactifications [16] where the SM/GUT lives on space-time filling D-branes (observable sector) while the hidden sector is given by the bulk of the ten dimensional ambient space. A hierarchical supersymmetry breaking [17] with stabilization of (some of the) moduli may be achieved in the bulk (hidden sector) by turning on background fluxes in the compactification manifold. Interesting models of flux compactifications have been found mostly in Type IIB theory by switching on appropriate RR and NS three-form fluxes in the internal directions [18] (other examples were obtained in heterotic M-theory, due to a non-vanishing G-flux sourced by the boundaries [9], and on the Type IIA front [10, 19]).

Alternatively, no-scale models may be found as well from generalized dimensional reduction à la Scherk–Schwarz (SS) [20] of eleven dimensional supergravity (or any truncation thereof). This is a generalized type of dimensional reduction on tori,
corresponding to a twist in the boundary conditions of the fields in the internal directions, with a phase dependent on a global symmetry of the theory (SS phase). The twist originates a no-scale model with non abelian couplings and a positive definite scalar potential, which, in general, yields spontaneous supersymmetry breaking. The charges obtained this way appear to be associated to the internal manifold geometry. However until recently they could not find a direct interpretation in the context of flux compactifications of superstring or M theory. We are going to see in the following section that such an interpretation is in fact possible in the presence of torsion in the internal manifold. Let us first briefly review what the SS mechanism is.

3.1. The Scherk-Schwarz mechanism

To describe how the mechanism works, let us consider the simple example of a free complex scalar field $\phi(x, y)$ on $\mathbb{R}^{1,3} \times S^1$, with lagrangian

$$\mathcal{L} = \frac{1}{2} \partial_{\hat{\mu}} \phi \partial^{\hat{\mu}} \phi^*, \quad (4)$$

where $\hat{\mu} = (\mu, 4); \mu = 0, \ldots, 3; x^4 = y$. We want to interpret this model from a four dimensional point of view, supposing that the $S^1$ has a tiny radius. However, instead of the Kaluza–Klein (KK) ansatz, let us impose the following generalized ansatz of dimensional reduction:

$$\phi(x, y) = e^{imy} \sum_{n=-\infty}^{+\infty} \phi_n(x)e^{iny/R}, \quad (y \sim y + 2\pi R). \quad (5)$$

We observe that this ansatz is not well defined, because $\phi(x, y)$ is multivalued on $S^1$. However, since the lagrangian (4) is invariant under global $U(1)$ transformations

$$\phi \rightarrow e^{i\alpha} \phi \quad (6)$$

then the phase $e^{imyR}$ has no physical effects. The five dimensional kinetic lagrangian (4) generates, in the dimensional reduction to four dimensions, a “scalar potential” (a mass term, in fact, in this case) $V = m^2 > 0$. All the masses are shifted by $m$, and the zero-mode has a mass $m$.

The key ingredient, which allows to implement the SS mechanism and to introduce in the dimensionally reduced theory the “charge” $m$, is the global invariance of the lagrangian (4) under the symmetry group $U(1)$. The charge is then introduced as a $y$-dependent phase in the dimensional reduction ansatz (5). The same mechanism, applied to theories invariant under more general global symmetries, allows to turn on more general charges related to the internal geometry.

In particular, let us now study the case of generalized dimensional reduction of $4+n$ dimensional gravity on $\mathcal{M}_{4+n} = \mathbb{R}^{1,3} \times S^1 \times T^{n-1}$. The result will be compared, in the next section, with an ordinary KK dimensional reduction of $4+n$ dimensional gravity on the same manifold $\mathcal{M}_{4+n}$ but in the presence of a torsion background.

Let us suppose to formally perform the reduction in two steps:
• first consider the KK compactification to five dimensions on \( T^{n-1} \). One gets a five dimensional theory on \( \mathbb{R}^{1,3} \times S^1 \) with global invariance \( SL(n-1) \), which is the isometry group of the torus \( T^{n-1} \).

• then we may use this global invariance for building a SS phase-matrix for a generalized dimensional reduction from five to four dimensions.

From the \( 4 + n \)-dimensional vielbein one gets the following fields in the theory reduced to five dimensions: the five dimensional vielbein \( V_\mu^\hat{a} \), \( n-1 \) vectors \( B^m_\mu \) and \( \frac{n(n-1)}{2} \) scalars \( \phi^i, \sigma \) parametrizing the coset \( \frac{SL(n-1)}{O(n-1)} \times O(1,1) \), whose coset-representatives are the internal components of the \( 4 + n \) dimensional vielbein, \( V^i_m(\phi^i, \sigma) \).

The SS ansatz for the dimensional reduction (restricted to the zero-modes) from five to four dimensions is (we take \( x^4 \equiv y^1 \) and no dependence of the fields on the extra internal coordinates \( x^5 = y^2, \ldots, x^{n-1} = y^{n-4} \))

\[
B^m_\mu(x, y^1) = U^m_n(y^1)B^m_\mu(x) \\
V^i_m(x, y^1) = (U^{-1})^n_m(y^1)V^i_n(x)
\]  

(7)

where \( U = \exp[My^1], M \in SO(n-1) \). The theory reduced to four dimensions has non abelian gauge charges which come from the dimensional reduction of the Levi-Civita connection on \( \mathcal{M}_{4+n} \), using the generalized ansatz (7). In particular one finds \( [20, 14] \)

\( m = (4, m); i = (4, i) \)

\[
F^m_{\mu \nu} = \partial_{[\mu} B^m_{\nu]} + f^m_{n \rho} B^n_{[\mu} B^\rho_{\nu]} \\
D_\mu V^i_m = \partial_\mu V^i_m + f^i_{n \rho} B^n_\mu V^i_\rho
\]  

(8) \hspace{1cm} (9)

with structure constants given by the SS phases:

\[
f^m_{4n} = M^m_n.
\]  

(10)

Let us make an observation which will be important in the next section. If the theory in \( 4 + n \) dimensions contains also p-forms besides the metric, then the SS prescription requires that the scalars coming from the dimensional reduction of the p-forms have to be rotated as well with the SS phases.

For the case, e.g., of a one-form \( A_m \) one has

\[
A_m(x, y^1) = (U^{-1})^n_m(y^1)A_n(x),
\]  

(11)

so that

\[
F_{4m}(x, y^1) = M^m_n A_n(x).
\]  

(12)

We may then conclude that the SS mechanism is surprisingly efficient in introducing geometrical charges in the theory. However it appears as a bit artificial prescription, which is not completely satisfying from the microscopical point of view.

We are going to see in the next section that the SS mechanism may be understood as an ordinary flux compactification if we assume the presence of a constant background torsion in the internal manifold.
4. The SS mechanism as a flux compactification with internal torsion

The purpose of the present section is to review how theories originating from SS dimensional reduction may alternatively be found as the result of a KK dimensional reduction in the presence of an internal torsion. In particular we shall consider a $D+n$ dimensional pure gravity theory compactified to $D$ dimensions on an $n$-torus $T^n$ showing that, by switching on an appropriate constant background torsion in the torus (a torsion flux), we may generate the same $D$-dimensional theory as the one described in section 3.1, originating from a SS reduction from $D+1$ dimensions with global symmetry generator chosen within the global symmetry algebra $\mathfrak{sl}(n-1, \mathbb{R})$. This is based on [6].

The relevance of torsion in flux compactifications was recently understood in the study of T-duality between Type IIB and Type IIA superstrings [21, 22]. Indeed it was found that the T-dual of the charge associated to the flux of $H^{(3)}_{NS}$ is not a p-form flux but internal torsion. Our geometrical interpretation allows to study also the no-scale SS models in the context of flux compactifications of superstring (or M) theory.

A delicate point in our analysis is the notion of gauge invariance in the presence of torsion. Indeed, we found that to have a precise identification of the SS models in terms of torsion a generalization of the notion of gauge transformation is needed, corresponding to the coupling of the torsion to the gauge fields.

4.1. Coupling torsion - gauge fields: the general idea

The torsion tensor may be defined as an antisymmetric contribution

$$T^P_{MN} \equiv \tilde{\Gamma}^P_{[MN]} = \frac{1}{2} (\tilde{\Gamma}^P_{MN} - \tilde{\Gamma}^P_{NM}),$$

(13)

to the affine connection $\tilde{\Gamma}$. In this case, the request of a metric connection ($\nabla_P g_{MN} = 0$) determines the connection coefficients as

$$\tilde{\Gamma}^P_{MN} = \Gamma^P_{MN} + K^P_{MN},$$

(14)

where $\Gamma^P_{MN} = \Gamma^P_{NM}$ are the coefficients of the Levi-Civita connection while

$$K^P_{MN} = \frac{1}{2} (T^P_{MN} - T^P_{MN} - T^P_{NM}).$$

(15)

contains the torsion contribution. For further definitions and conventions, we refer to [6].

To discuss the origin of the coupling of the torsion to gauge fields let us consider, as an example, the case of a vector field $A_M$. The field strength feels the effect of the torsion background via the principle of general covariance:

$$F^P_{MN} \equiv \nabla_{[M}A_{N]} = \partial_{[M}A_{N]} + T^P_{MN}A_P.$$  

(16)

When $T \neq 0$, $F$ is not invariant under the usual gauge transformation

$$\delta A_M(X) = \partial_N \Lambda(X).$$

(17)
However, it is possible \cite{23} to make the torsion compatible with the presence of gauge fields by introducing the generalized definition of gauge transformation

\[
\delta_C A_M(X) = C^N_M(X) \partial_N \Lambda(X),
\]  
(18)

where the point-dependent matrix \( C^N_M(X) \) has to be constrained by the request of (generalized) gauge invariance of the field strength

\[
\delta_C F_{MN} = 0.
\]  
(19)

The general solution, found in \cite{23} for a gauge field in \( d \) dimensions coupled to torsion in \( d \) dimensions is

\[
C^N_M = \delta^N_M e^\phi, \quad T^P_{MN} = \delta^P_M \partial_N \phi.
\]  
(20)

There is only one torsion degree of freedom allowed, the scalar field \( \phi \). Let us now see that this restriction may be largely relaxed if the idea of \cite{23} is combined with the Kaluza–Klein ansatz in a dimensional reduction context, and that it is then possible to turn on an internal torsion with many independent components. The different components of the torsion generate geometrical charges that reproduce the SS phases while giving them a geometrical interpretation.

4.2. Coupling torsion / gauge fields and dimensional reduction

Let us consider a gauge field \( A_M \) in the \( 4+n \) dimensional manifold \( \mathcal{M}_{4+n} = \mathbb{R}^{1,3} \times T^n = \mathbb{R}^{1,3} \times S^1 \times T^{n-1} \), with coordinates \( X^M = (x^\mu, y, y^m) \), \( m = 1, \cdots n-1 \in T^{n-1} \) and \( \mu = 0, 1, 2, 3 \in \mathbb{R}^{1,3} \); \( y \) parametrizes \( S^1 \) (it is the same manifold as in section 3.1). We ask that the KK ansatz holds true:

\[
A_M = A_M(x); \quad \Lambda = \Lambda(x).
\]  
(21)

We suppose to have a constant torsion background on the torus \( T^n \), and that the gauge field is coupled to the torsion as in \cite{13}. Then, \( F = \nabla A \) is not gauge invariant, but it is invariant under the generalized gauge transformation à la HRRS:

\[
\delta_C A_M = C^N_M \partial_N \Lambda = C^\nu_M \partial_\nu \Lambda.
\]  
(22)

If we further assume:

\[
C^N_M = C^N_M(y),
\]  
(23)

then the generalized gauge invariance condition \( \delta_C F = 0 \) has now more general solutions (see \cite{6} for details), allowing the following non-zero entries for the matrices \( C \) and \( T \):

\[
C^N_M : (C^N_M, C^n_M); \quad T^M_{NP} : (T^\mu_{\nu 4}, T^m_{NP})
\]  
(24)

with the only constraint

\[
C^\nu_M(y) = \delta^\nu_\mu e^\phi(y); \quad T^\mu_{\nu 4} = \frac{1}{2} \delta^\mu_\nu \partial_4 \phi(y).
\]  
(25)

By choosing

\[
T^m_{4n} = M^m_n \in SO(n-1), \quad (\phi = 0),
\]  
(26)
we have a consistent theory, where
\[ \delta_C A_\mu = \delta A_\mu = \partial_\mu \Lambda, \quad \delta_C A_4 = 0, \quad \delta_C A_m = 0. \]  
(27)
Moreover it gives
\[ F_{4n} = M^m_n A_m \]  
(28)
which precisely reproduces the effect of the geometrical SS phase for a vector, equation \([12]\).

Let us now consider the KK dimensional reduction of the metric in the presence of the torsion background \([20]\), and let us compare the result with the one of Section 3.1. As in section 3.1, the theory reduced to four dimensions has non-abelian gauge charges which come from the dimensional reduction of the affine connection on \(\mathcal{M}_{4+n}\), and we find the same result as before
\[ F_{\mu \nu}^{m} = \partial_{[\mu} B_{\nu]}^{m} + f_{\nu}[\mu B_{\nu]}^{p} \]  
(29)
\[ D_{\mu} V_{i}^m = \partial_{\mu} V_{i}^{m} + f^{m}_{\nu \rho} B_{\nu \rho} V_{i}^{m} \]  
(30)
However, the interpretation is now different. In this case the charges do not come from the dimensional reduction of the metric (which is in fact an ordinary KK reduction) but from the torsion contribution to the affine connection \([12]\), which, differently from section 3.1 in this case is not any more Levi–Civita. The structure constants of the four dimensional gauge theory are now given by the components of the torsion
\[ f_{4n}^m = T_{4n}^m. \]  
(31)
The comparison between the two interpretations is summarized in the table below.

| \(f_{4n}^m\) | S-S | Torsion |
|---|---|---|
| \(M_{n}^m\) = SS phases | \(T_{4n}^m\) = torsion |
| \(F_{4m} : A_m(x, y) = (e^{-M y})^m_n A_n(x)\) | \(F_{4m} = T_{4n}^m A_n\) |
| \(V_{m}^i : (e^{-M y})^m_n V_{n}^i(x)\) | \(V_{n}^i(x)\) |
| \(\Gamma_{NP}^M : \Gamma_{[NP]}^M = 0\) | \(\Gamma_{[NP]}^M \neq 0\) |

5. Conclusions

We discussed how gauge charges and interactions do emerge in supergravity. We focused in particular on the role of the charges related to the geometry, and shown that the SS mechanism is in fact an ordinary compactification in the presence of an internal torsion background. The scenario that seems to emerge is that all the bulk charges, responsible for supersymmetry breaking, may be generated by compactifications in the presence of fluxes and internal torsion.

There are various directions for future investigations, which include a similar analysis of theories corresponding to compactifications on more general manifolds, the construction of a \(D = 4\) gauged supergravity deriving from a general torsion background, and the effects of a dynamical torsion on this setting.
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