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Compactifications on twisted tori with fluxes and free differential algebras

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Abstract

We describe free differential algebras for non-Abelian one and two form gauge potentials in four dimensions deriving the integrability conditions for the corresponding curvatures. We show that a realization of these algebras occurs in M-theory compactifications on twisted tori with constant four-form flux, due to the presence of antisymmetric tensor fields in the reduced theory.

1. Introduction

Flux compactifications on twisted tori provide interesting examples of string and M-theory compactifications where most of the moduli fields are stabilized [2–8]. Particular cases of such compactifications include heterotic string, type II orientifold models and M-theory in the presence of constant p-form fluxes (where p depends on the particular string model and $p = 4$ in M-theory). When fluxes and (or) Scherk–Schwarz geometrical fluxes are turned on, interesting gauge algebraic structures emerge which in most cases have the interpretation of a gauged Lie algebra [5,9–12].

In this case the Maurer–Cartan equations (zero curvature conditions) read

$$dA^\Lambda + \frac{1}{2} f^\Lambda{}_{\Sigma\Gamma} A^\Sigma \wedge A^\Gamma = 0, \qquad (1.1)$$

where integrability implies the Jacobi identities

$$f^\Lambda{}_{[\Sigma\Gamma} f^\Pi{}_{\Delta]\Lambda} = 0. \qquad (1.2)$$

This comes from the vanishing of the cubic term

$$d^2 A^\Pi = -\frac{1}{2} f^\Lambda{}_{\Sigma\Gamma} f^\Pi{}_{\Delta\Lambda} A^\Sigma \wedge A^\Gamma \wedge A^\Delta = 0. \qquad (1.3)$$

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When fundamental tensor fields are present in the theory, in absence of gauge couplings in the supergravity theory, one can transform them into scalars and this is the way the full duality symmetry (sometimes called U-duality) is recovered. However, in presence of non-abelian gauge couplings, an obstruction can arise in the dualization of such antisymmetric tensors, so that the theory only preserves some subalgebra of the full duality group. Moreover, the gauged algebra structure may be more complicated than an ordinary Lie algebra and in fact, as noted in [10] for generic Scherk–Schwarz and form flux couplings it turns out to be a free differential algebra (FDA) [13–17].

In the case of M-theory, we will show that its Maurer–Cartan equations are equivalent to the integrability conditions for the 4-form G_{IJKL} and for the vielbein 1-form in $D = 11$. This also will explain how the Lie algebra part of the free differential algebra is deformed in the presence of generic Scherk–Schwarz and form flux couplings.

2. The free differential algebra and its Maurer–Cartan equations

The generalization of (1.1) to a free differential algebra including 2-form gauge fields B_i consists of the following (zero-curvature) system

$$\mathcal{F}^\Lambda = dA^\Lambda + \frac{1}{2} f^\Lambda{}_{\Sigma\Gamma} A^\Sigma \wedge A^\Gamma + m^{\Lambda i} B_i = 0, \quad (2.1)$$

$$\mathcal{H}_i = dB_i + (T_\Lambda)_i{}^j A^\Lambda \wedge B_j + k_{i\Lambda\Sigma\Gamma} A^\Lambda \wedge A^\Sigma \wedge A^\Gamma = 0, \quad (2.2)$$

where $f^\Lambda{}_{\Sigma\Gamma}$, $(T_\Lambda)_i{}^j$, $m^{\Lambda i}$ and $k_{i\Lambda\Sigma\Gamma}$ are the structure constants of the FDA.

The integrability condition of this system comes from the Bianchi identities

$$d\mathcal{F}^\Lambda = 0, \quad (2.3)$$

$$d\mathcal{H}_i = 0. \quad (2.4)$$

From (2.3), by setting to zero the terms proportional to A^3 and $A \wedge B$ polynomials we get

$$f^\Lambda{}_{\Sigma[\Gamma} f^\Sigma{}_{\Pi\Delta]} + 2m^{\Lambda i} k_{i\Gamma\Pi\Delta} = 0, \quad (2.5)$$

$$f^\Lambda{}_{\Sigma\Gamma} m^{\Sigma j} + m^{\Lambda i} (T_\Gamma)_i{}^j = 0, \quad (2.6)$$

respectively. From (2.4) we get three conditions from the vanishing of the terms proportional to $B \wedge B$, $B \wedge A \wedge A$ and from A^4 terms:

$$(T_\Lambda)_i{}^{(j} m^{\Lambda k)} = 0, \quad (2.7)$$

$$(T_\Lambda)_i{}^j f^\Lambda{}_{\Sigma\Gamma} - 2(T_{[\Sigma})_i{}^k (T_{\Gamma]})_k{}^j + 6m^{\Lambda j} k_{i\Lambda\Sigma\Gamma} = 0, \quad (2.8)$$

$$3 f^\Lambda{}_{[\Sigma\Gamma} k_{i\Pi\Delta]\Lambda} - 2(T_{[\Pi})_i{}^j k_{j\Sigma\Gamma\Delta]} = 0. \quad (2.9)$$

When $m^{\Lambda i} = 0$, the condition (2.5) implies for the A^Λ the ordinary Lie algebra Jacobi identities. Eq. (2.8) tells us that $(T_\Lambda)_i{}^j$ is a representation of the Lie algebra and (2.9) states that $k_{i\Lambda\Sigma\Gamma}$ is a cocycle of the Lie algebra. When $m^{\Lambda i} k_{i\Gamma\Pi\Delta} \neq 0$ (2.5) gives the departure from an ordinary Lie algebra for the f structure constants.

3. FDA from M-theory on twisted tori with fluxes

As an example of a concrete realization of the free differential algebra (2.1) and (2.2), we will now describe the one obtained by compactification of M-theory on twisted tori in the presence of fluxes considered in [10]. The compactification of M-theory to 4 dimensions provides 28 vector fields G^I_μ, $A_{\mu IJ}$ and 7 2-form tensor fields $A_{\mu\nu I}$. This means that we can identify the generic indices Λ, i of our FDA as follows $\Lambda = \{I, IJ\}$, $i = I$. Furthermore, one has to write the single indices I, J in the same position as Λ, i, but the antisymmetric couples $IJ, KL, \ldots$ are written as upper indices if $\Lambda, \Sigma, \ldots$ are lower ones and as lower indices if $\Lambda, \Sigma, \ldots$ are upper ones.

If one considers first the case when only form fluxes are turned on, the Lie algebra is

$$[Z_I, Z_J] = g_{IJKL} W^{KL}, \qquad [Z_I, W^{JK}] = [W^{IJ}, W^{KL}] = 0, \quad (3.1)$$

which is the central extension of an Abelian gauge algebra. In this case the only non-vanishing structure constants are [10]

$$f^\Lambda{}_{\Sigma\Gamma} = f_{[IJ]KL} = g_{IJKL}, \quad (3.2)$$

$$k_{i\Lambda\Sigma\Gamma} = k_{IJKL} = \frac{1}{6} g_{IJKL}, \quad (3.3)$$

while $m^{\Lambda i} = (T_\Lambda)_i{}^j = 0$. It then follows that (2.5) and (2.6) are trivially satisfied and g_{IJKL} is arbitrary. This result is a consequence of the very degenerate structure of the Lie algebra (3.1).

An intermediate richer example comes in the case of Scherk–Schwarz fluxes τ_{IJ}^K and vanishing 4-form flux. This is the case considered in the pioneering papers of Scherk–Schwarz [1,2]. In this case $k_{i\Lambda\Sigma\Gamma} = 0$, but $m^{\Lambda i}$ and $(T_\Lambda)_i{}^j$ do not vanish. In fact, the non-vanishing parts of these structure constants are

$$\begin{aligned} &m^{\Lambda i} \neq 0 \quad \text{for } \Lambda = [IJ], \quad i = K, \\ &m_{IJ}{}^K = \tau_{IJ}^K, \end{aligned} \tag{3.4}$$

$$\begin{aligned} &(T_\Lambda)_i{}^j \neq 0 \quad \text{for } \Lambda = I, \quad i = J, \quad j = K, \\ &(T_I)_J{}^K = -\tau_{IJ}^K. \end{aligned} \tag{3.5}$$

The other non-vanishing structure constants occur for $f^\Lambda{}_{\Sigma\Gamma}$ when

$$\begin{aligned} &\Lambda = I, \qquad \Sigma = J, \qquad \Gamma = K, \qquad f^I{}_{JK} = \tau_{JK}^I, \\ &\Lambda = [IJ], \qquad \Sigma = K, \qquad \Gamma = [LM], \\ &f_{[IJ]K}{}^{[LM]} = -2\tau_{K[I}^{[L}\delta_{J]}^{M]}. \end{aligned} \tag{3.6}$$

In this case (2.7) is identically satisfied and (2.6), (2.8) are identical to (2.5), which reads as $\tau_{[IJ}^L \tau_{K]L}^M = 0$. Note that $m^{\Lambda i}$ corresponds to a "magnetic" mass term for the B_i field.

The $f^\Lambda{}_{\Sigma\Gamma}$ structure constants in (3.6) define the Scherk–Schwarz algebra for M-theory:

$$\begin{aligned} &\left[W^{IJ}, W^{KL}\right] = 0, \qquad [Z_I, Z_J] = \tau_{IJ}^K Z_K, \\ &\left[Z_I, W^{JK}\right] = 2\tau_{IL}^{[J} W^{K]L}. \end{aligned} \tag{3.7}$$

Let us now consider the general case when both τ_{IJ}^K and g_{IJKL} are non-vanishing. In this case the last term in (2.5) is non-vanishing for $\Lambda = [IJ]$, $\Sigma = K$, $\Gamma = L$ and $\Pi = M$. It reads

$$\tau_{IJ}^N g_{KLMN}. \tag{3.8}$$

If this term does not vanish the f structure constants do not define a Lie algebra. In this case (2.5) (as also (2.9)) becomes

$$\tau_{[IJ}^N g_{KLM]N} = 0. \tag{3.9}$$

This condition has the 11-dimensional interpretation of the integrability condition of the 4-form field strength [10].

All other equations are satisfied as a consequence of the τ Jacobi identities $\tau_{[IJ}^L \tau_{K]L}^M = 0$. These follow from (2.5) by taking $\Lambda, \Sigma, \Gamma, \Pi = IJKL$. It is obvious that if the stronger condition (3.8) holds then the $f^\Lambda{}_{\Sigma\Gamma}$ define an ordinary Lie algebra. This happens if the Scherk–Schwarz fluxes τ_{IJ}^K have the K index complementary to the flux coupling g_{IJKL}. This can actually be realized in certain type II orientifold models.

To summarize, we have shown that for generic Scherk–Schwarz couplings τ_{IJ}^K and 4-form flux g_{IJKL}, the M-theory gauge algebra is a free differential algebra rather than an ordinary Lie algebra. The equations

$$\tau_{[IJ}^M \tau_{K]M}^L = 0, \tag{3.10}$$

$$\tau_{[IJ}^N g_{KLM]N} = 0, \tag{3.11}$$

are the integrability conditions for the FDA. When the stronger condition $\tau_{IJ}^N g_{KLMN} = 0$ holds then the $f^\Lambda{}_{\Sigma\Gamma}$ define an ordinary Lie algebra whose commutators read [10]

$$\begin{aligned} &[Z_I, Z_J] = g_{IJKL} W^{KL} + \tau_{IJ}^K Z_K, \\ &\left[Z_I, W^{JK}\right] = 2\tau_{IL}^{[J} W^{K]L}, \\ &\left[W^{IJ}, W^{KL}\right] = 0. \end{aligned} \tag{3.12}$$

It is interesting to note that in M-theory compactified on a twisted torus with 4-form flux turned on $m^{\Lambda i}$ and g_{PQRS} have the physical interpretation of magnetic and electric masses for the antisymmetric tensors B_I. This is clear looking at the covariant field strength

$$\mathcal{F}^\Lambda = dA^\Lambda + \frac{1}{2} f^\Lambda{}_{\Sigma\Gamma} A^\Sigma \wedge A^\Gamma + m^{\Lambda I} B_I. \tag{3.13}$$

This expression appears quadratically in the (kinetic part of the) Lagrangian together with the coupling

$$g_{IJKL} B_M \wedge dA_{NP} \epsilon^{IJKLMNP}, \tag{3.14}$$

which comes from the 11-dimensional Chern–Simons term $F \wedge F \wedge A$. It is amusing to note that the consistency condition [18,19] for electric and magnetic contributions to the mass is in this case a consequence of (3.9).

The M-theory FDA also includes a 3-form gauge field C which is a singlet. The zero-curvature condition for this 3-form is

$$\begin{aligned} &dC + m^{ij} B_i \wedge B_j + m^i_{\Lambda\Sigma} A^\Lambda \wedge A^\Sigma \wedge B_i \\ &\quad + t_\Lambda A^\Lambda \wedge C + k_{\Lambda\Sigma\Gamma\Delta} A^\Lambda \wedge A^\Sigma \wedge A^\Gamma \wedge A^\Delta = 0. \end{aligned} \tag{3.15}$$

In the M-theory FDA, the only non-vanishing terms are $k_{IJKL} \sim g_{IJKL}$ and $m^I_{JK} \sim \tau^I_{JK}$, with all the other components and t_Λ and m^{ij} vanishing. In this case the Bianchi identity is trivially satisfied because a 5-form in $D=4$ identically vanishes. However, the curvature of C can be determined by demanding its full invariance under all gauge transformations.

4. Non-zero curvature case

The previous Maurer–Cartan equations (2.1) and (2.2), which entail the "structure constants" relations (2.5)–(2.9) can be lifted to non-zero curvature, so obtaining covariant Bianchi identities for the curvatures. In the case of M-theory with Scherk–Schwarz fluxes turned on this procedure essentially reproduces the covariant curvatures G of Section 3.4 of [2]. When also the constant 4-form fluxes $F_{IJKL}=g_{IJKL}$ are turned on, then one gets generalized curvatures which are covariant under the combined 1-form and 2-form gauge transformations considered in Section 2 of [10].

An interesting new feature of the curvatures is the presence in $\mathcal{H}_I$ of a "contractible generator" [13], i.e., in physical language, of a curvature itself (which also exists in the ungauged theory)

$$\mathcal{H}_I = dB_I + \mathcal{F}^J \wedge A_{IJ}, \tag{4.1}$$

where $\mathcal{F}^J = dA^J$. This is a kind of Green–Schwarz (mixed) Chern–Simons term which modifies the gauge transformations of B_I so that $\mathcal{H}_I$ is invariant under the gauge transformations

$$\begin{aligned}
&\delta B_I = d\Lambda_I - \epsilon_{IJ}\mathcal{F}^J,\\
&\delta A^I = d\omega^I,\\
&\delta A_{IJ} = d\epsilon_{IJ}.
\end{aligned} \tag{4.2}$$

The (ungauged) Bianchi identity is now

$$d\mathcal{H}_I = \mathcal{F}^J \wedge \mathcal{F}_{IJ}, \tag{4.3}$$

which satisfies $d^2\mathcal{H}_I = 0$ and is also invariant under the gauge transformations (4.2).

Let us now consider the case when $g_{IJKL}\neq 0$ (but $\tau_{IJ}^K = 0$), so that

$$\begin{aligned}
&\mathcal{F}_{IJ} = dA_{IJ} + \frac{1}{2} g_{IJKL} A^K \wedge A^L,\\
&\mathcal{F}^I = dA^I.
\end{aligned} \tag{4.4}$$

Then the $\mathcal{H}_I$ curvature reads

$$\begin{aligned}
\mathcal{H}_I = {} & dB_I + \mathcal{F}^J \wedge A_{IJ}\\
& + \frac{1}{6} g_{IJKL} A^J \wedge A^K \wedge A^L,
\end{aligned} \tag{4.5}$$

and the coefficient of the $\mathcal{F}\wedge A$ term is fixed, relative to the A^3 term in such a way that $d\mathcal{H}_I = \mathcal{F}^J \wedge \mathcal{F}_{IJ}$. Now $\mathcal{H}_I$ and its Bianchi identity are invariant under the gauge transformations

$$\delta B_I = d\Lambda_I - \epsilon_{IJ}\mathcal{F}^J + \frac{1}{2}\omega^M g_{MIJK} A^J \wedge A^K, \tag{4.6}$$

$$\delta A^I = d\omega^I, \tag{4.7}$$

$$\delta A_{IJ} = d\epsilon_{IJ} - g_{IJKL}\omega^K A^L. \tag{4.8}$$

Analogously, the threefold antisymmetric tensor C curvature is

$$dC - \mathcal{F}^I \wedge B_I + \frac{1}{4!} g_{IJKL} A^I \wedge A^J \wedge A^K \wedge A^L, \tag{4.9}$$

which is invariant under the gauge transformations

$$\begin{aligned}
\delta C = {} & d\Sigma + \mathcal{F}^I \wedge \Lambda_I\\
& - \frac{1}{6} g_{IJKL}\omega^I \wedge A^J \wedge A^K \wedge A^L,
\end{aligned} \tag{4.10}$$

$$\delta B_I = d\Lambda_I - \epsilon_{IJ}\mathcal{F}^J + \frac{1}{2}\omega^M g_{MIJK} A^J \wedge A^K, \tag{4.11}$$

$$\delta A^I = d\omega^I. \tag{4.12}$$

Note that the dC field strength is a Lagrange multiplier and can be algebraically eliminated from the Lagrangian giving a contribution to the scalar potential.

5. Concluding remarks

In the present Letter we have considered the free differential algebra which comes from M-theory compactified on a twisted torus with constant 4-form fluxes. This is just a special case of the Maurer–Cartan equations described by (2.1) and (2.2). A similar situation arises in type IIA theories since in this case charged antisymmetric tensor fields are also present. However, in this case one can find a particular set of geometrical fluxes which can be consistently set to vanish and then the Lie algebra structure is recovered because the condition

$$m^{\Lambda i} k_{i\Sigma\Gamma\Pi} = 0, \tag{5.1}$$

is satisfied. Such examples were described in [10].

The FDA given by the system of curvatures $\mathcal{F}^\Lambda$, $\mathcal{H}_i$ can be recast in the form of an ordinary Lie algebra if (some of the) B_i are redefined so that the quadratic term in $A^\Sigma \wedge A^\Gamma$ is absorbed in the new $\tilde{B}_i$ [13]. This can be done at most for rank(m) tensors fields, which can be the same as the range of the i indices provided that this is smaller than that of the vector fields Λ, as in the M-theory case. Explicitly, for those B_α for which the subblock $m^{\alpha\beta}$ is invertible, one can introduce the definition ($\Lambda = \{\alpha, A\}$)

$$\tilde{B}_\alpha \equiv B_\alpha + \frac{1}{2} m^{-1}_{\alpha\beta} f^\beta{}_{\Lambda\Sigma} A^\Lambda \wedge A^\Sigma, \tag{5.2}$$

so that the new zero curvature conditions read

$$d\tilde{B}_\alpha = 0, \tag{5.3}$$

$$\mathcal{F}^\alpha = dA^\alpha + m^{\alpha\beta}\tilde{B}_\beta = 0, \tag{5.4}$$

$$\mathcal{F}^A = dA^A + \frac{1}{2} f^A{}_{BC} A^B \wedge A^C = 0. \tag{5.5}$$

The new Lie algebra is defined by the structure constants $f^A{}_{BC}$ and this is obtained by deleting the A^α generators from the original algebra. This is the quotient of the original algebra with the subalgebra related to the A^α vectors. It is an obvious consequence of the Jacobi identities for $\mathcal{F}^A$ that $f^A{}_{\Lambda\Sigma} = 0$ whenever Λ or Σ take values in the α range.

In the M-theory case, the rank of $m^{\Lambda i}$ is encoded in the Scherk–Schwarz fluxes τ^K_{IJ} regarded as a 7×21 triangular matrix. A quadratic submatrix can have at most rank 7 so the Lie algebra spanned by the A^A is at least 21-dimensional. When describing the algebra in terms of its generators, one must delete the generators $W^{\tilde{L}\tilde{K}}$ whose gauge fields are absorbed by the antisymmetric tensors. The resulting Lie algebra is obtained by all Z_K, W^{LK} generators but the $W^{\tilde{L}\tilde{K}}$, which is an Abelian subalgebra. A simple example is the case when τ^K_{IJ} correspond to a "flat group". In this case $B_I = \{B_0, B_\alpha\}$ and $A^\Lambda = \{A^\alpha, A^A\}$, with $A^\alpha = A_{0\alpha}$ and $A^A = \{A^0, A^\alpha, A_{\alpha\beta}\}$. The original structure constants follow from $\tau^K_{IJ} = \tau^\alpha_{0\beta} = t^\alpha{}_\beta$, where $\alpha, \beta = 1, \ldots, 6$, and t is an invertible antisymmetric matrix, (this means that the 3 skew eigenvalues are non-zero). The redefined tensor fields are

$$\tilde{B}_\alpha = t^\delta{}_\gamma t^{-1\beta}_\alpha A_{\beta\delta} \wedge A^\gamma + A_{0\alpha} \wedge A^0 + B_\alpha, \tag{5.6}$$

$$\tilde{B}_0 = B_0 - A^\alpha \wedge A_{0\alpha}, \tag{5.7}$$

and the zero curvatures conditions read

$$dA^0 = 0, \tag{5.8}$$

$$dA^\alpha + t^\alpha{}_\beta A^0 \wedge A^\beta = 0, \tag{5.9}$$

$$dA_{0\alpha} + t^\beta{}_\alpha \tilde{B}_\beta = 0, \tag{5.10}$$

$$dA_{\alpha\beta} + 2t^\gamma{}_{[\alpha} A^0 \wedge A_{\beta]\gamma} = 0, \tag{5.11}$$

$$d\tilde{B}_\alpha = 0, \tag{5.12}$$

$$d\tilde{B}_0 - t^\beta{}_\alpha A^\alpha \wedge A^\gamma \wedge A_{\beta\gamma} = 0. \tag{5.13}$$

Note that the Jacobi identities of the τ do not set any constraint on the t matrices. If we split the generators into Z_0, Z_α, $W^{0\alpha}$, $W^{\alpha\beta}$, it is immediate to see that the index α goes over six values and the gauge fields $A_{\mu 0\alpha}$ disappear from the gauge algebra. The generator algebra becomes then

$$\begin{aligned} &\left[Z_0, W^{\alpha\beta}\right] = 2\tau^{[\alpha}_{0\gamma} W^{\beta]\gamma}, \qquad [Z_0, Z_\alpha] = \tau^\gamma_{0\alpha} Z_\gamma, \\ &\left[Z_\alpha, W^{\beta\gamma}\right] = [Z_\alpha, Z_\beta] = \left[W^{\alpha\beta}, W^{\gamma\delta}\right] = 0, \end{aligned} \tag{5.14}$$

which is the usual (22-dimensional) flat Scherk–Schwarz algebra. This algebra becomes 24- or 26-dimensional if one or two eigenvalues of the t matrix vanish. The same reasoning applies when form fluxes are present. In this case the commutators of the Z_α are not vanishing and the gauge algebra get modified.

Note that the physical interpretation of this reduction of the FDA to a minimal part and a contractible one [13] corresponds to the anti-Higgs mechanism where antisymmetric tensors absorb vector fields to become (dual to) massive vectors. The quotient Lie algebra is the unbroken gauge algebra. It is interesting to see that, due to the cubic terms of the 1-forms A^Λ in the $\mathcal{H}_i$ curvature (this only happens when the 4-form flux is present), the quadratic part of the $\mathcal{F}^\Lambda$ curvature does not correspond to an ordinary Lie algebra before the quotient has been taken.

Another interesting generalization is to extend such FDA to the fermionic sector of the theory, since the $D = 4$ theory has $N = 8$ local supersymmetry. Such program was originally carried out in $D = 11$ in [14] and its extension to the present compactification should be possible.

We finally remark that the different structures of the 4-dimensional effective theories obtained when the gauge algebra is a FDA or an ordinary Lie algebra are reflected in different scalar potentials. This fact may have important consequences when looking for complete moduli stabilization in such compactifications.

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